

Layer Intercalibration of the ATLAS Electromagnetic Calorimeter and CP-odd Higgs Boson Couplings Measurements in the Four-Lepton Decay Channel with Run 2 Data of the LHC

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Personal contributions

This manuscript includes both details and analyses that are of general interest or are necessary to understand the rest of its content.

Chapter 1 gives an introduction, both mathematical and phenomenological, to the Standard Model of particle physics.

Chapter 2 introduces the LHC. Section 2.7 on the LHC filling scheme is necessary to understand the studies developed in Chapter 6 on the LAr electromagnetic calorimeter intercalibration, and Section 2.9 on parton distribution functions is necessary to understand Chapter 9 on the Higgs CP analysis.

Chapter 3 and Chapter 4 detail how collisions are recorded in the ATLAS detector and how the electronic signal from each subdetector is used to reconstruct electrons, muons, *etc.* that are used in the analyses. In particular Section 3.4.1 and 3.4.2 give general information on the calorimetry and are useful to understand the details of the calibration detailed in Chapter 6. Section 3.4.3 is required to understand Chapter 6

Chapter 5 gives an overview of the electron calibration procedure in ATLAS (see note below). This Chapter details analyses to which I have not directly contributed to, but most of them depend on my contribution to the layer intercalibration.

One of the two main topics of this Thesis is detailed in Chapter 6: the intercalibration of the first and second layer of the electromagnetic calorimeter. My contributions to this analysis are the following:

- I showed that the Run 1 method used for this calibration was not performing as expected in Run 2.
- I developed the extrapolation method described and checked its validity.
- I provided the estimate of the systematics, including the new ones inherent to the change of method.

The cross-checks described in Section 6.5 and Section 6.7 were performed by other people in the group.

Chapter 7 gives more precise information on the Higgs phenomenology at the LHC. In particular, the EFT framework used in Chapter 9 is introduced in Section 7.3.

Chapter 8 describes the general $H \rightarrow 4\ell$ analysis: event selection, background estimation, systematics, categorisation, which are common to all $H \rightarrow 4\ell$ analyses. My contribution to this is

- the maintaining of the data and MC central processing code for the $H \rightarrow 4\ell$ group.

I have been responsible for providing analysis n-tuples to the group until early 2017, and the analysis presented in this Chapter is the most up-to-date one and so includes many changes that were developed since then.

The other main topic of this Thesis is the search for a CP -odd contribution to the Higgs boson couplings (in the $H \rightarrow 4\ell$ decay channel). This analysis is detailed in Chapter 9 and my contributions are:

- the investigation and validation of a CP -sensitive observable (Section 9.1),
- in particular the development of an analysis code tool interfacing the MADGRAPH matrix element computation to the analysis (Section 9.1.6),
- the validation of the morphing method in the case of this analysis (Section 9.2.2),
- the estimation of the sensitivity of this analysis in various EFT bases (Section 9.3),
- including the validation of the simulated samples,
- the preliminary evaluation of systematic uncertainties impacting the measurements (Section 9.4).

Contents

Introduction	11
1 The Standard Model of particle physics	15
1.1 Mathematical framework: Quantum Field Theory	15
1.1.1 Notions of group theory	16
1.1.2 Least action principle and Lagrangian formalism	18
1.1.3 Gauge theories	18
1.1.4 Symmetries and Noether theorem	21
1.1.5 Non abelian symmetry groups: the example of $SU(3)$	22
1.1.6 Electroweak theory and missing mass terms	23
1.1.7 Spontaneous symmetry breaking, BEH mechanism, and gauge boson masses	25
1.1.8 Fermion masses	28
1.1.9 Mathematical formulation of the Standard Model	29
1.2 Discovery of the Higgs boson and first results	31
1.2.1 Theoretical constraints before the discovery	31
1.2.2 Direct searches at LEP, Tevatron, and the LHC	32
1.2.3 Successes and limitations of the Standard Model	36
1.2.4 Status of the Higgs boson after Run 1 at the LHC	38
2 The Large Hadron Collider	43
2.1 Why colliding particles	43
2.2 Motivations for a multi-TeV proton collider at CERN	44
2.3 A hadron collider in the LEP tunnel	45
2.4 The CERN accelerator complex	46
2.5 The Large Hadron Collider	46
2.6 The LHC experiments	49
2.7 LHC filling scheme	50
2.8 Luminosity and number of events	51
2.9 Parton distribution functions at hadron colliders	52
2.10 LHC performance	54
3 ATLAS: A Toroidal LHC ApparatuS	57
3.1 Requirements for the ATLAS detector	58
3.2 ATLAS coordinate systems	60
3.3 The inner detector and the solenoid magnet	61
3.3.1 The pixel detector	63

3.3.2	The Strip Semiconductor Tracker (SCT)	63
3.3.3	The Transition Radiation Tracker (TRT)	63
3.3.4	Central solenoid magnet	64
3.4	The electromagnetic and hadronic calorimeters	65
3.4.1	Principles of calorimetry	65
3.4.2	Application to the ATLAS electromagnetic calorimeter	67
3.4.3	Central and endcap electromagnetic calorimeter	68
3.4.4	Central hadronic calorimeter	76
3.4.5	The hadronic endcap and forward calorimeters	77
3.5	The muon spectrometer and the toroid magnets	80
3.6	Luminosity monitoring	82
3.7	Trigger system	85
4	Reconstruction and identification of physics objects	89
4.1	Reconstruction of tracks in the ID	89
4.2	Reconstruction of tracks in the MS	90
4.3	Reconstruction of energy in calorimeter cells	90
4.4	Electrons and photons	92
4.5	Muons	95
4.6	Jets and b -tagging	98
4.7	Missing transverse energy	99
5	Calibration of the electromagnetic calorimeter	101
5.1	Electron and photon calibration overview	102
5.2	MVA-based calibration	103
5.3	E_1/E_2 intercalibration	106
5.4	Pre-Sampler scales	106
5.4.1	Overview of the PS scale determination	106
5.4.2	Determination of the correlation factor $A(\eta)$	107
5.4.3	Determination of the correction factor $b_{1/2}(\eta)$	107
5.4.4	Results	109
5.4.5	PS scales with muons	110
5.5	Third layer energy calibration	111
5.6	Material estimation in front of the calorimeter	111
5.6.1	Reasons for a precise material estimate and methodology	111
5.6.2	Results and validation	112
5.6.3	Nominal geometry for Run 2	113
5.7	$Z \rightarrow ee$ energy scales	115
5.7.1	Overview	115
5.7.2	Estimation of α_{ij} and c'_{ij} using the template method	118
5.7.3	Estimation of α_i and c'_i using the Z lineshape method	119
5.7.4	$Z \rightarrow ee$ scale results and systematics	120
5.8	Energy scale and resolution extrapolation over the energy range	121
5.9	Cross-checks of energy scales and resolution extrapolation	125
5.9.1	Methodology	125
5.9.2	Determination of the J/ψ prompt fraction	125
5.9.3	Extraction of residual scales using the lineshape method	126
5.9.4	Results and systematics	126

6	ECal first and second layers intercalibration	129
6.1	Overview of the layer intercalibration	130
6.2	Geometry considerations: η definition and binning	133
6.2.1	Relation between cell η and track η	133
6.2.2	Calorimeter geometry impact in the barrel for $1.30 < \eta_{\text{corr}} < 1.48$	134
6.2.3	Calorimeter geometry impact in the endcap for $1.375 < \eta_{\text{corr}} < 1.50$	135
6.2.4	Final binning	136
6.3	Number of cells summed	137
6.4	Extraction of the muon energy deposit using an analytical fit	138
6.4.1	Energy distribution fit with Landau * Gaussian convolution	138
6.4.2	Systematic investigation	139
6.4.2.1	Number of cells in S2	139
6.4.2.2	Bunch train effects	139
6.4.2.3	Tighter isolation requirements	141
6.4.2.4	Ambient noise subtraction	143
6.4.2.5	First hint: pileup dependency	145
6.4.3	Noise modelling with templates	148
6.4.4	Extrapolating the layer intercalibration to $\langle\mu\rangle = 0$	148
6.4.5	Results	153
6.4.6	Systematic uncertainties of the MPV method	153
6.4.6.1	Validation of the extrapolated results	155
6.4.6.2	Choice of the second cell in S2	155
6.4.6.3	Leakage energy in S1	157
6.4.6.4	Leakage energy in S2	159
6.4.6.5	Combined plot of MPV systematic uncertainties	161
6.5	Estimation of the muon energy deposit using the TM method	161
6.5.1	Truncated mean definition	161
6.5.2	Truncated mean results	164
6.5.3	Systematics of the truncated mean method	164
6.6	Comparison of muon MPV and TM method, and to Run 1 result	165
6.6.1	Comparison of MPV and TM uncertainties in Run 2	166
6.6.2	Final result: Run 2 combination of the MPV and TM methods	168
6.6.3	Run-1 / Run-2 comparison	169
6.7	Cross-check: electron method	169
6.7.1	E_1/E_2 follow up with $Z \rightarrow ee$ electrons	170
6.7.2	Results and compatibility with the muon method	172
	Conclusion of the layer intercalibration analysis	173
7	Higgs boson phenomenology and EFT	175
7.1	Theoretical aspects of the SM Higgs boson	176
7.1.1	Higgs boson production modes at pp colliders	176
7.1.1.1	Gluon fusion: ggF	176
7.1.1.2	Vector-boson fusion: VBF	177
7.1.1.3	Associated production with a vector boson: VH	178
7.1.1.4	Associated production with a $t\bar{t}$ or $b\bar{b}$ pair: $t\bar{t}H/b\bar{b}H$	179
7.1.2	Decay channels	179
7.1.3	Spin and parity	183
7.2	Beyond the Standard Model	184
7.2.1	The κ -framework	185

7.2.2	Simplified Template Cross-Sections (STXS)	186
7.3	Effective Field Theories: EFT	186
7.3.1	Energy power counting	186
7.3.2	A simple example: the Fermi theory of β decay	188
7.3.3	Standard Model EFT and the Warsaw basis	189
7.3.4	The Higgs basis	191
7.3.5	Basis translation and the SMEFTSIM package	192
7.3.6	Higgs Characterisation framework	192
7.4	Higgs couplings and CP -mixing results from LHC Run 1 and Run 2	194
7.4.1	STXS results	194
7.4.2	κ -framework	194
7.4.3	EFT interpretation	196
7.4.4	Direct CP measurements and couplings structure functions	198
7.4.5	Direct CP measurement with single parameter parametrisation of the effective Lagrangian	200
8	Experimental aspects of the Higgs to Four Leptons decay	203
8.1	Overview of the analysis	204
8.2	Data and simulated samples	205
8.2.1	Data samples	205
8.2.2	Steps of an event generation	206
8.2.2.1	Hard-scatter	206
8.2.2.2	Additional radiations	207
8.2.2.3	Hadronisation	207
8.2.3	Simulated signal samples	207
8.2.4	Simulated background samples	208
8.3	Event selection	208
8.3.1	Trigger requirements	209
8.3.2	Physics objects selection	209
8.3.3	Lepton quadruplet building	210
8.3.4	Lepton quadruplet selection	211
8.4	Backgrounds estimation	212
8.4.1	Non-resonant background	212
8.4.2	Reducible $\ell\ell + \mu\mu$ background	213
8.4.3	$\ell\ell + \mu\mu$ transfer factor systematics uncertainties	218
8.4.4	Reducible $\ell\ell + ee$ background	220
8.4.5	Extrapolation to the SR using the $Z+X$ CR	222
8.4.6	Categorised background yield and determination of the background shapes	224
8.5	Systematic uncertainties common to all $H \rightarrow 4\ell$ analyses	225
8.5.1	Experimental uncertainties	226
8.5.2	Theory uncertainties on the signal modelling	227
8.5.3	Background uncertainties	228
8.6	Event categorisation for the couplings analysis	228
8.6.1	STXS categorisation	229
8.6.2	Improved category separation using neural networks	230
8.7	Expected yields with the full Run-2 dataset	232
9	Measurement of CP-odd Higgs-boson couplings	239
9.1	Optimal observables	240

9.1.1	<i>CP</i> -sensitive observables in the $H \rightarrow 4\ell$ decay channel and in the VBF production mode	240
9.1.2	Definition of Optimal Observables	242
9.1.3	Technical generation of matrix elements with MADGRAPH	243
9.1.4	Dependence on the PDF choice	245
9.1.5	Impact of changing the hypothesis provided to the matrix element computing	246
9.1.6	Development of a MADGRAPH-based simplified matrix element computation package	247
9.2	Morphing technique	250
9.2.1	Mathematical definition of the morphing method	250
9.2.2	Validation of the $\mathcal{OO}_{1,jj}$ morphed shapes	252
9.3	Preliminary studies	252
9.3.1	Additional BSM simulated samples	254
9.3.2	Sensitivity using VBF signal samples only	254
9.3.3	Limits from cross-section results	255
9.3.4	Validation of the morphing technique	257
9.3.5	Sensitivity including $ZZ^{(*)}$ and other Higgs boson production modes	257
9.4	Preliminary results	259
9.4.1	Expected sensitivity with fully-simulated samples	259
9.4.2	Systematic uncertainties assessment	261
9.4.3	Study of a potential bias from usage of neural networks	262
	Summary and conclusion of the <i>CP</i> -odd couplings analysis	264
	Conclusion	267
	Appendices	269
	A Layout of distorted geometries	271
	B STXS Stage-0 categorisation yields	275
	C Synthèse en français	277
C.1	Contexte expérimental et théorique	277
C.1.1	Le Modèle Standard et le boson de Higgs	277
C.1.2	Phénoménologie du boson de Higgs	278
C.1.3	Spin-parité (<i>CP</i>) du boson de Higgs : motivation et résultats récents	279
C.1.4	Au-delà du Modèle Standard : théories effectives	280
C.2	Le LHC et le détecteur ATLAS	281
C.2.1	Le LHC (Grand Collisionneur de Hadrons)	281
C.2.2	Le détecteur ATLAS	282
C.2.3	Aperçu du calorimètre électromagnétique	283
C.3	Étalonnage du calorimètre électromagnétique d'ATLAS	285
C.3.1	Étalonnage relatif L1/L2 : objectifs et méthode	285
C.3.2	Étalonnage relatif L1/L2 : résultats initiaux et limites	287
C.3.3	Étalonnage relatif L1/L2 : méthode d'extrapolation et patrons de bruit	288
C.3.4	Étalonnage relatif L1/L2 : résultats finals et systématiques	289
C.4	Étude de couplages <i>CP</i> -impairs du boson de Higgs	291
C.4.1	L'analyse $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$	291
C.4.2	Observables optimales	292
C.4.3	Étude des couplages impairs par <i>CP</i> du boson de Higgs	294

List of Abbreviations	297
List of Figures	299
List of Tables	309
Bibliography	313
References for Chapter 1 — The Standard Model of particle physics	313
References for Chapter 2 — The Large Hadron Collider	318
References for Chapter 3 — ATLAS: A Toroidal LHC ApparatuS	321
References for Chapter 4 — Reconstruction and identification of physics objects	322
References for Chapter 5 — Calibration of the electromagnetic calorimeter	323
References for Chapter 6 — ECal first and second layers intercalibration	326
References for Chapter 7 — Higgs boson phenomenology and EFT	327
References for Chapter 8 — Experimental aspects of the Higgs to Four Leptons decay	332
References for Chapter 9 — Measurement of CP -odd Higgs-boson couplings	336
References for Chapter A — Layout of distorted geometries	337
References for Chapter B — STXS Stage-0 categorisation yields	337
References for Chapter C — Synthèse en français	337
Acknowledgements / Remerciements	343

Introduction

The Standard Model of particle physics was established during the second half of the twentieth century and provides the current best description of fundamental particles and their interactions. One key element of the model is the spontaneous symmetry breaking mechanism, necessary to give mass to the particles. This mechanism is associated with a particle, the Higgs boson, whose discovery in 2012 by the ATLAS and CMS experiments at the Large Hadron Collider (LHC) completed the Standard Model.

However, despite being tested up to a formidable accuracy, the Standard Model fails to explain some observed phenomena, such as the matter–antimatter asymmetry, and is therefore considered to be the low energy realisation of a higher-energy theory. As no new particle emerged after the discovery of the Higgs boson, searches for new physics turned to precision measurements to try to spot hints of physics beyond the Standard Model. Such hints are therefore actively investigated in many areas of particle physics. Given the Higgs boson discovery is still recent, large room for possible beyond the Standard Model physics effects are left in this sector. Interest thus rapidly moved to measurements of its properties to try to spot a deviation from the Standard Model prediction: cross-sections, branching ratios, spin and parity (CP).

The CP symmetry, and more particularly its violation, is one of the three Sakharov conditions necessary for the baryogenesis. CP violation is predicted and has already been seen in the quark sector and hints are present in the neutrino sector, but would be pure beyond the Standard Model physics if found in the Higgs sector. The spin-parity of the Higgs boson is predicted to be 0^+ by the Standard Model. Run 1 analyses have excluded the pure spin-parity states 0^- , 1^+ , 1^- , 2^+ , 2^- at more than 99% confidence level, leaving 0^+ as the only pure state compatible with the observations. However, these analyses relied on the assumption that the observed particle is a pure spin-parity state, and CP -mixed states are therefore not yet completely excluded. Such cases reveals especially interesting since a mixing between 0^+ and 0^- states would induce CP violation in the Higgs sector.

The CP analyses are embedded in the Higgs boson anomalous couplings measurements. Such studies were started during Run 1 and were interpreted in terms of couplings modifiers in the so-called κ -framework. However, with an integrated luminosity five times higher than in Run 1 and a Higgs boson production cross-section twice as much at 13 TeV than at 8 TeV, the Run 2 dataset offers ten time more data to study more precisely the Higgs boson properties. As a consequence, a more general framework is introduced to interpret more finely the Higgs boson couplings results: effective field theories.

The Higgs boson decaying into a pair of Z bosons, each of them subsequently decaying into a pair of leptons (electrons or muons) was one of the discovery channels in 2012, and remains one of the major channels for Higgs properties measurements. Despite a low yield, the $H \rightarrow 4\ell$ has low backgrounds and offers a high sensitivity with a signal-to-background ratio slightly over

two. This channel therefore is a natural choice for the studies of the Higgs couplings to vector bosons. Focusing on the vector-boson fusion (VBF) and Higgs-strahlung (VH) production modes, anomalous Higgs boson couplings to vector bosons appear both on the production and decay side, enhancing potential beyond the Standard Model effects. Thanks to a high momentum transfer, the VBF production mode provides the best sensitivity to such effects.

While cross-section measurements, now known at the 25 % level for the VBF mode, provide constraints on a possible CP -odd contribution to the couplings between the Higgs boson and the vector bosons, they do not allow to distinguish CP -even from CP -odd effects and to determine the sign of an eventual CP -odd contribution. Separation of these effects can nevertheless be achieved through angular information between the final-state particles. Such angular information is found in the numerous variables built from the high multiplicity $H \rightarrow 4\ell$ final state, or from jet angular separation in VBF collisions. However, such variables only partially use the kinematic information available from the final state. To improve the sensitivity to CP effects, an observable based on matrix element computation is built, making use of all the kinematic information and therefore called *optimal observable*. This observable is also able to unambiguously discriminate CP -even from CP -odd effects.

This document is organised as follows. In Chapter 1, the mathematical construction of the Standard Model is presented. The constraints put on the Higgs boson before its discovery are summarised. These include the theoretical constraints as well as the limits put by the Large Electron–Positron collider and Tevatron experiments until their respective shutdown, in 2000 and 2011. Finally, the results leading to the Higgs boson discovery at the LHC in 2012 and the first studies performed on its properties during Run 1 are recalled.

Chapter 2, describes the motivations for building the LHC at CERN. The chain of injectors providing protons to the LHC is detailed and the parameters of the LHC and its injectors are summarised. The performance of the LHC during Run 2, achieving a record centre-of-mass energy of 13 TeV and a record luminosity up to $2 \times 10^{34} \text{ cm}^{-2} \text{ s}^{-1}$ is discussed.

The ATLAS detector is described in Chapter 3. The various subsystems, their specifications and their role in the event reconstruction and particle identification are detailed. Strong emphasis is put on calorimetry and the liquid argon electromagnetic calorimeter as its calibration is one of the main topic developed in this Thesis.

Chapter 4 gives a short description of how physics objects (electrons, muons, *etc.*) are reconstructed from the electric signals recorded in the detector. The energy reconstruction in the liquid argon calorimeter cells and the impact of the additional collision in bunch crossing (pileup) in the calorimeter is discussed.

In order to achieve precise physics results, objects such as electrons, muons, photons, *etc.* need to be precisely measured and calibrated. The liquid argon electromagnetic calorimeter plays a central role in the measurement of electron and photon energy, and its energy calibration is detailed in Chapter 5. The calibration procedure includes a succession of steps improving the uniformity and linearity of the energy response in the calorimeter. The final energy scales are derived from $Z \rightarrow ee$ events at the end of the calibration procedure, providing an electron energy resolution at the permil level.

One specific step of the calibration is the calorimeter layers intercalibration, described in Chapter 6. This step corrects for residual electronics miscalibration and cross-talk effects. This task is performed using muons, allowing to probe fine regions of the detector, but whose energy deposit in the calorimeter of the same order as the noise. As a consequence, this measurement is sensitive to the noise, and in particular the one induced by pileup. With harsher pileup conditions in Run 2, a new method is developed to achieve a similar precision than in Run 1 despite the higher pileup level.

In Chapter 7, the current state-of-the-art results on the Higgs boson are recalled. Theoretical measurement and interpretation frameworks such as simplified template cross-section, the κ -framework or effective field theories are presented. Recent experimental results on Higgs boson couplings are shown as reference.

The Chapter 8 introduces the general $H \rightarrow 4\ell$ analysis within the ATLAS experiment. The background and systematics estimates common to all $H \rightarrow 4\ell$ analyses are presented. The event categorisation in production modes, and its refinement using neural networks is detailed

The measurements of the CP -odd Higgs boson to vector boson couplings is performed in Chapter 9. The construction and validation of the aforementioned optimal observables is detailed. The morphing technique used to interpolate between the simulated points is described. The expected limits on various effective field theory coefficients are derived using such optimal observables and compared to limits from previous cross-section measurements. Finally, preliminary systematic uncertainties are assessed.

The Standard Model of particle physics

Contents

1.1 Mathematical framework: Quantum Field Theory	15
1.1.1 Notions of group theory	16
1.1.2 Least action principle and Lagrangian formalism	18
1.1.3 Gauge theories	18
1.1.4 Symmetries and Noether theorem	21
1.1.5 Non abelian symmetry groups: the example of $SU(3)$	22
1.1.6 Electroweak theory and missing mass terms	23
1.1.7 Spontaneous symmetry breaking, BEH mechanism, and gauge boson masses	25
1.1.8 Fermion masses	28
1.1.9 Mathematical formulation of the Standard Model	29
1.2 Discovery of the Higgs boson and first results	31
1.2.1 Theoretical constraints before the discovery	31
1.2.2 Direct searches at LEP, Tevatron, and the LHC	32
1.2.3 Successes and limitations of the Standard Model	36
1.2.4 Status of the Higgs boson after Run 1 at the LHC	38

1.1 Mathematical framework: Quantum Field Theory

Particle physics describes the world of tiny objects, subject to the rules of quantum mechanics, and which may be travelling at high speed (close to the speed of light), like in particle accelerators. We therefore need a framework that handles both quantisation and special relativity. Such framework is founded by the Quantum Field Theory (QFT).

Notation: to simplify and compact the writing of the mathematical framework, the following standard notations are used throughout the rest of this Chapter.

- The covariant notation X_μ ($\mu = 0, 1, 2, 3$, similarly for other greek indices) denotes a quantity depending on space-time coordinates (position, momentum, *etc.*). By convention, X_0 corresponds to the time coordinate, and $\mathbf{X} = (X_1, X_2, X_3)$ to the space coordinates.

- The derivative with respect to the μ^{th} space-time coordinate $\partial/\partial x_\mu$ is abbreviated ∂_μ .
- Repeated greek indices denote an implicit summation as in the Einstein convention: $X_\mu Y^\mu = X^0 Y_0 - \mathbf{X} \cdot \mathbf{Y} = X_0 \cdot Y_0 - X_1 \cdot Y_1 - X_2 \cdot Y_2 - X_3 \cdot Y_3$. For repeated latin indices, the convention is a simple summation: $X_a Y_a = X^a Y^a = X_1 \cdot Y_1 + X_2 \cdot Y_2 + \dots$.

The commonly admitted natural units $\hbar = c = 1$ are used.

This Section describes the steps that lead to the derivation of the Standard Model Lagrangian in Section 1.1.9. Renormalisation and quantization methods are not discussed (see for example Ref. [1]).

1.1.1 Notions of group theory

A group is a set G with the internal operator $*$ fulfilling the following axioms:

- Closure** $\forall a, b \in G, a * b \in G,$
Associativity $\forall a, b, c \in G, (a * b) * c = a * (b * c) = a * b * c,$
Identity $\exists e \in G, \forall a \in G, a * e = e * a = a,$
Inverse $\forall a \in G, \exists b \in G, a * b = b * a = e.$ The inverse b of a is commonly noted a^{-1} or $-a$.

The number of elements of such group is called the *order* of the group, and can be finite or infinite. Note that no assumption is made on the commutativity of the operator (groups with a commutative internal law are called *abelian groups*).

Simple examples of groups include the (finite) group of symmetries of an equilateral triangle (identity plus 3 axial plus 2 rotations symmetries), or the (infinite) group of integers with additive operator $(\mathbb{Z}, +)$. Groups can however contain completely abstract elements that can only hardly be mentally grasped. Fortunately, given two groups $(G, \cdot)$ and $(G', \times)$, if there exists an homomorphism $\phi : G \rightarrow G'$ such that

$$\forall g_1, g_2 \in G, \phi(g_1 \cdot g_2) = \phi(g_1) \times \phi(g_2), \quad (1.1)$$

then G' is called a *representation* of G .

In particular, if there exists an homomorphism $\phi : G \mapsto \text{GL}_n(\mathbb{K})$ ¹ for a finite dimension n , then ϕ is said to be a matrix representation of degree n of the group G . This will be verified for all groups useful in quantum physics, which will therefore be conveniently represented by a group of square matrices. A matrix representation of degree n is said to be irreducible if no set of $m \times m$ matrices with $m < n$ can be found which forms a valid matrix representation of G .

For infinite groups, two types are distinguished: the discrete groups (such as $(\mathbb{Z}, +)$), and the continuous (or parametrised) groups whose elements depend on one or several continuous parameters. An example of a continuous group is the group of rotations in $\mathbb{R}^2$: an irreducible matrix representation of this group is

$$\left\{ R(\theta) = \begin{pmatrix} \cos \theta & -\sin \theta \\ \sin \theta & \cos \theta \end{pmatrix} \right\}_{\theta \in \mathbb{R}}. \quad (1.2)$$

An infinite number of matrices (of dimension 2) are necessary to describe all such rotations, but they can be parametrised by a continuous parameter θ (the angle of rotation). A continuous

¹ $\uparrow \text{GL}_n(\mathbb{K})$ is the *general linear group* over the field $\mathbb{K}$, i.e. the group of invertible matrices of dimension n . Here, $\mathbb{K}$ is the field of real or complex numbers, $\mathbb{R}$ or $\mathbb{C}$.

group with finite degree matrix representation is called a *Lie group*, and such a kind of structure is of prime importance for the development of the Standard Model.

It can be shown that the matrices U representing a Lie group can be written in the form

$$U = \exp(\boldsymbol{\alpha} \cdot \mathbf{T}) = \exp(\alpha_a T_a), \quad (1.3)$$

with repeated a indices standing for an implicit sum over the a index, where the T_a matrices composing $\mathbf{T}$ are called *generators* of the Lie group and the α_a composing $\boldsymbol{\alpha}$ are the continuous parameters. It can further be shown that

$$T_a = \left. \frac{\partial U}{\partial \alpha_a} \right|_{\alpha_a=0}. \quad (1.4)$$

In the above case of 2D rotations, $\boldsymbol{\alpha}$ is limited to one parameter θ and there is only one generator:

$$J = \begin{pmatrix} 0 & -1 \\ 1 & 0 \end{pmatrix}. \quad (1.5)$$

One more element of group theory is needed before returning to physics: the Lie algebra structure constants. Let us consider a slightly more complex example: the group of rotations in $\mathbb{R}^3$. This group depends on 3 parameters $\theta_x, \theta_y, \theta_z$ for the 3 rotations axes x, y, z respectively. The matrix representation of such group is:

$$\{R_x(\theta_x) + R_y(\theta_y) + R_z(\theta_z)\}_{\theta_x, \theta_y, \theta_z \in \mathbb{R}}, \quad (1.6)$$

with

$$R_x(\theta_x) = \begin{pmatrix} 1 & 0 & 0 \\ 0 & \cos \theta_x & -\sin \theta_x \\ 0 & \sin \theta_x & \cos \theta_x \end{pmatrix}, R_y(\theta_y) = \begin{pmatrix} \cos \theta_y & 0 & \sin \theta_y \\ 0 & 1 & 0 \\ -\sin \theta_y & 0 & \cos \theta_y \end{pmatrix}, R_z(\theta_z) = \begin{pmatrix} \cos \theta_z & -\sin \theta_z & 0 \\ \sin \theta_z & \cos \theta_z & 0 \\ 0 & 0 & 1 \end{pmatrix}. \quad (1.7)$$

The generators of this group are

$$J_x = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & -1 \\ 0 & 1 & 0 \end{pmatrix}, \quad J_y = \begin{pmatrix} 0 & 0 & 1 \\ 0 & 0 & 0 \\ -1 & 0 & 0 \end{pmatrix}, \quad J_z = \begin{pmatrix} 0 & -1 & 0 \\ 1 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}. \quad (1.8)$$

Defining $\boldsymbol{\theta} = (\theta_x, \theta_y, \theta_z)$ and $\mathbf{J} = (J_x, J_y, J_z)$, we have $R(\boldsymbol{\theta}) = \exp(\boldsymbol{\theta} \cdot \mathbf{J})$. However, since the matrices J_i, J_j, J_k do not commute, this is not equal to $\exp(\theta_x J_x) \exp(\theta_y J_y) \exp(\theta_z J_z)$. The commutation relations between the J matrices read

$$[J_x, J_y] = J_z, \quad [J_y, J_z] = J_x, \quad [J_z, J_x] = J_y, \quad (1.9)$$

where $[X, Y] \doteq XY - YX$ is called the *commutator* of X and Y . This can be rewritten as

$$[J_i, J_j] = \epsilon_{ijk} J_k \quad (i, j, k \in \{x, y, z\}), \quad (1.10)$$

where ϵ_{abc} is the fully anti-symmetric 3-index tensor, so-called Levi-Civita tensor.

Leaving the specific case of the 3D rotations and returning to the general case, the commutation relations between the generators of a group are noted

$$[T_a, T_b] = f_{abc} T_c, \quad (1.11)$$

and the f_{abc} are called the *structure constants* of the Lie algebra associated to the group. These structure constants are sufficient to derive all the structure of the Lie algebra related to the group. In the above example, the Levi-Civita tensor is the structure constant of the group of rotations in $\mathbb{R}^3$.

1.1.2 Least action principle and Lagrangian formalism

Quantum physics, just as classical mechanics, rely on a key assumption: the least action principle. However, contrary to classical mechanics, quantum physics cannot predict with an arbitrary precision both the position and the momentum of a particle at the same time. As a consequence, the content of a system is not described by a point-like particle but rather with a complex field depending on space-time coordinates commonly denoted $\phi(x_\mu)$ whose squared amplitude ($|\phi(x_\mu)|^2$) can be interpreted as a presence probability density at a given point of space and time. The following development is made with the consideration of fields, nevertheless the formalism is based and still close to its original classical counterpart.

Expressed in terms a field ϕ , the action S is defined as

$$S = \int d^4x \mathcal{L}(\phi, \partial_\mu \phi, x_\mu), \quad (1.12)$$

where $\mathcal{L}$ is the *Lagrangian density*, a function that depends on the field $\phi(x^\mu)$, its first derivative $\partial_\mu \phi(x^\mu)$ and on the space-time coordinate itself x^μ . The least action principle states that the time evolution of the system follows a path in phase-space that yields an extremal action (usually minimal): $\delta S = 0$. From that assumption, the fundamental Euler-Lagrange equations can be derived:

$$\frac{\partial \mathcal{L}}{\partial \phi} - \partial_\mu \frac{\partial \mathcal{L}}{\partial (\partial_\mu \phi)} = 0. \quad (1.13)$$

This can be generalised for an arbitrary number of fields, yielding as many Euler-Lagrange equations as fields present. Developing Eq. 1.13 leads to the equations of motion for the fields.

1.1.3 Gauge theories

Moving group theory to physics, a symmetry is defined as a transformation that leaves the equations of motion, or equivalently the Lagrangian, unchanged under its action. A symmetry is called either global if it is independent of the space-time coordinates, or local if the transformation is a function of the space-time location. A simple illustration is given in Figure 1.1. Symmetries are further qualified as a space-time symmetry if the space-time coordinates change under the given symmetry, or as an internal symmetry if the symmetry only acts on internal degrees of freedom (that can be more or less abstract: charge, spin, isospin, *etc.*) of the Lagrangian.

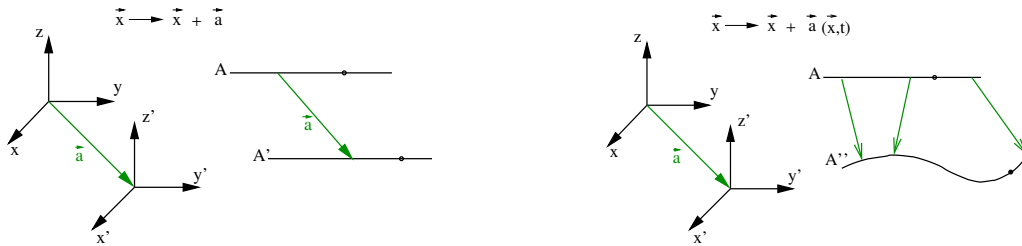


Figure 1.1 – Example trajectories before and after a global (left) and local (right) space translation (space-time symmetry). [2]

In physics, gauge theories are defined as field theories whose Lagrangian remains invariant under the action of a local symmetry. Gauge theories are the cornerstone of the mathematical formulation of the Standard Model, as their properties can be derived from symmetries and a few first principles. The issue of renormalisation is not be discussed here, but it should be noted that all gauge theories were proven to be renormalisable thanks to a result by 't Hooft and Veltman in 1972 [3], making them of fundamental interest for particle physics.

As an example, let us consider the free (*i.e.* without interaction) Dirac Lagrangian, describing the propagation of a free, massive fermionic field (*i.e.* a Dirac spinor) $\psi(x)$ of mass m :

$$\mathcal{L}_{\text{Dirac}}^{\text{free}} = \bar{\psi}(i\partial - m)\psi, \quad (1.14)$$

where

$$\bar{\psi} = \psi^\dagger \gamma^0, \quad (1.15)$$

with $\phi^\dagger$ the adjoint (transpose conjugate) of ψ , and

$$\partial = \gamma^\mu \partial_\mu. \quad (1.16)$$

The Dirac matrices γ^μ are 4×4 matrices that can be expressed from the Pauli matrices σ^k :

$$\gamma^0 = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}, \quad \gamma^k = \begin{pmatrix} 0 & \sigma^k \\ -\sigma^k & 0 \end{pmatrix}, \quad (1.17)$$

with

$$\sigma^1 = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}, \quad \sigma^2 = \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix}, \quad \sigma^3 = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}. \quad (1.18)$$

If we define a global continuous transformation of parameter α under the $U(1)$ group²:

$$\psi(x) \rightarrow e^{i\alpha} \psi(x), \quad (1.19)$$

then $\bar{\psi}(x) \rightarrow e^{-i\alpha} \bar{\psi}(x)$ and $\mathcal{L}_{\text{Dirac}}^{\text{free}} \rightarrow e^{-i\alpha} \bar{\psi}(i\partial - m)e^{i\alpha} \psi = \mathcal{L}_{\text{Dirac}}^{\text{free}}$ since the derivative do not act on $e^{i\alpha}$. The transformation of Eq. 1.19 is thus a symmetry for the free Dirac Lagrangian.

If we now consider a similar but *local* transformation of $U(1)$, called a *gauge transformation*:

$$\psi(x) \rightarrow U(x)\psi(x) = e^{i\alpha(x)}\psi(x), \quad (1.20)$$

then $\bar{\psi}(x) \rightarrow e^{-i\alpha(x)}\bar{\psi}(x)$ and

$$\begin{aligned} \mathcal{L}_{\text{Dirac}}^{\text{free}} &\rightarrow e^{-i\alpha(x)} \bar{\psi} i \partial (e^{i\alpha(x)} \psi) - e^{-i\alpha(x)} \bar{\psi} m e^{i\alpha(x)} \psi \\ &= e^{-i\alpha(x)} \bar{\psi} i \left(i \partial(\alpha(x)) e^{i\alpha(x)} \psi + e^{i\alpha(x)} \partial \psi \right) - \bar{\psi} m \psi \\ &= e^{-i\alpha(x)} \bar{\psi} i e^{i\alpha(x)} \partial \psi - e^{-i\alpha(x)} \bar{\psi} \partial(\alpha(x)) e^{i\alpha(x)} \psi - \bar{\psi} m \psi \\ &= \bar{\psi} i \partial \psi - \bar{\psi} \partial(\alpha(x)) \psi - \bar{\psi} m \psi \\ &= \mathcal{L}_{\text{Dirac}}^{\text{free}} - \bar{\psi} \partial(\alpha(x)) \psi, \end{aligned} \quad (1.21)$$

which is different from $\mathcal{L}_{\text{Dirac}}^{\text{free}}$. The free Dirac Lagrangian is therefore not invariant under a local $U(1)$ symmetry, in other words it is not gauge invariant under $U(1)$.

This can be intuitively understood as a problem in the definition of the derivative: since the transform is local (depends on the space-time), the definition of the derivative is not the same in each point of space-time (see for example Chapter 15.1 of Ref. [1] for a detailed explanation). It can be shown that the derivative definition can be fixed by changing it to a covariant derivative D_μ , whose form

$$D_\mu = \partial_\mu + iqA_\mu(x), \quad (1.22)$$

²↑ The unitary group of n dimension $U(n)$ is the group of $n \times n$ complex matrices M such that $M^\dagger M = 1$. In one dimension, such objects are assimilated to a scalar field and form the unit complex circle $\mathbb{U} = \{z \in \mathbb{C}, |z| = 1\}$. In our case they represent a shift of phase of the field.

(q an arbitrary constant) can be derived from minimal assumptions. It can further be shown that the newly introduced A_μ field should transform under the same gauge transformation by

$$A_\mu(x) \rightarrow A_\mu(x) - \frac{1}{q} \partial_\mu \alpha(x). \quad (1.23)$$

It can be shown that the number of additional fields required to recover the gauge invariance is equal to the number of generators of the symmetry group.

With the new covariant derivative definition in Eq. 1.22, Eq. 1.14 is therefore changed into:

$$\mathcal{L}_{\text{Dirac}}^{\text{gauge inv.}} = \bar{\psi}(i\cancel{D} - m)\psi, \quad (1.24)$$

and it can be shown that this Lagrangian is gauge-invariant under $U(1)$. The field $A_\mu(x)$ needed to recover the gauge invariance is called a *gauge field*.

If we now expand a bit the $\mathcal{L}_{\text{Dirac}}^{\text{gauge inv.}}$:

$$\begin{aligned} \mathcal{L}_{\text{Dirac}}^{\text{gauge inv.}} &= \bar{\psi} i \gamma^\mu (\partial_\mu + i q A_\mu) \psi - \bar{\psi} m \psi \\ &= \bar{\psi} (i \gamma^\mu \partial_\mu - m) \psi - q \bar{\psi} \gamma^\mu A_\mu \psi \\ &= \mathcal{L}_{\text{Dirac}}^{\text{free}} - q \bar{\psi} A \psi. \end{aligned} \quad (1.25)$$

The second term relates an anti-fermion $\bar{\psi}$ a fermion ψ and the field A_μ , in other words is an interaction term! To complete this Lagrangian, one should still allow the vector field A_μ to propagate, *i.e.* add a free-field term for A_μ . Again using the covariant derivative definition and only requiring gauge invariance, it can be shown that

$$F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu \quad (1.26)$$

is the only possibility. It can further be checked that $F_{\mu\nu}$ is also gauge invariant under $A_\mu(x) \rightarrow A_\mu(x) - \frac{1}{q} \partial_\mu \alpha(x)$:

$$\begin{aligned} F_{\mu\nu} &\rightarrow \partial_\mu \left(A_\nu - \frac{1}{q} \partial_\nu \alpha \right) - \partial_\nu \left(A_\mu - \frac{1}{q} \partial_\mu \alpha \right) \\ &= \partial_\mu A_\nu - \partial_\nu A_\mu - \frac{1}{q} (\partial_\mu \partial_\nu \alpha - \partial_\nu \partial_\mu \alpha) = F_{\mu\nu}, \end{aligned} \quad (1.27)$$

since $\partial_\mu \partial_\nu = \partial_\nu \partial_\mu$.

It is interesting to note that massive gauge fields are forbidden by imposing gauge invariance: if we add a term $\frac{1}{2} m^2 A_\mu A^\mu$ into the Lagrangian, then under $U(1)$ gauge transformation:

$$\frac{1}{2} m^2 A_\mu A^\mu \rightarrow \frac{1}{2} m^2 \left(A_\mu - \frac{1}{q} \partial_\mu \alpha \right) \left(A^\mu - \frac{1}{q} \partial^\mu \alpha \right) = \frac{1}{2} m^2 A_\mu A^\mu + \dots, \quad (1.28)$$

additional terms appear that cannot be compensated for. The only possibility to maintain gauge invariance is to have $m = 0$.

To summarise, starting from the Lagrangian for a free fermion field in Eq. 1.14 and only imposing local gauge invariance, we manage to make a gauge field appear, and to create an interaction term as in Eq. 1.25. Casting this result to a concrete case, we can interpret A_μ as the photon field (q will be the electric charge) as entering the quantum electrodynamics (QED) Lagrangian:

$$\mathcal{L}_{\text{QED}} = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} + \bar{\psi}(i\cancel{D} - m)\psi \quad (1.29)$$

$$= -\frac{1}{4} (\partial_\mu A_\nu - \partial_\nu A_\mu) (\partial^\mu A^\nu - \partial^\nu A^\mu) + \bar{\psi} i \gamma^\mu \partial_\mu \psi - m \bar{\psi} \psi + q \bar{\psi} \gamma^\mu A_\mu(x) \psi. \quad (1.30)$$

An easy way of representing the propagation and interactions of fields was found by Richard Feynman, in the so-called Feynman diagrams. The so-called associated Feynman rules for the mathematical formulation of the diagrams can be deduced from the Lagrangian but are not detailed here. The Feynman diagrams appearing in QED are described in Figure 1.2. By convention, time is represented on the horizontal axis (flowing left to right) and space on the vertical axis. Fermions are represented by a solid line with an arrow following time (anti-fermion with a reverse arrow), and gauge bosons by a wavy line.

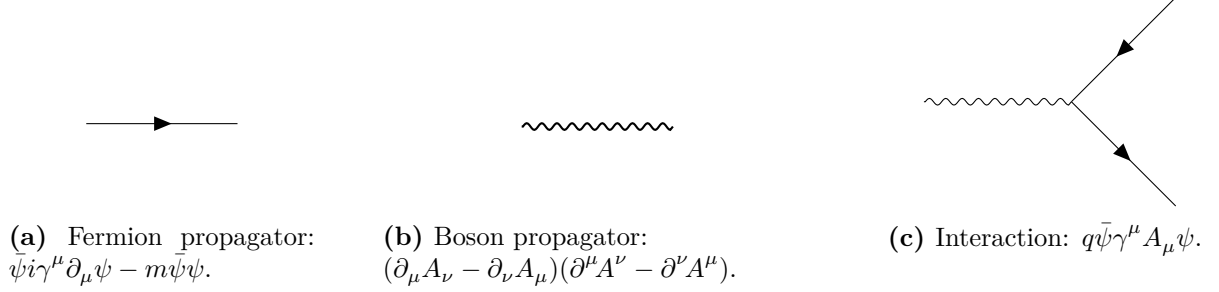


Figure 1.2 – Feynman diagrams associated with the Dirac Lagrangian under $U(1)$ gauge symmetry.

1.1.4 Symmetries and Noether theorem

To help simplify a problem, one can search for symmetries of the system. For example, the movement of a test particle in a central potential is much easier to describe in spherical coordinates than in cartesian ones. This is due to the rotational symmetry of the system: some quantities do not change after a rotation of the reference frame for example.

Starting from this consideration, another major building block for the mathematical formulation of the Standard Model is Noether's theorem. This theorem was first demonstrated by Emmy Noether in 1915 [4] and states that to each continuous symmetry leaving the action of a system unchanged, there is a corresponding conserved quantity, and conversely. In terms of gauge theory, this translates into a conserved current for each gauge symmetry, generically denoted J^μ and such that $\partial_\mu J^\mu = 0$. This four-current arises from the Euler-Lagrange equations under the symmetry (see for example Chapter 2.2 of Ref. [1]). It can further be shown that one conserved current arises for each generator of the group of symmetry.

In the $U(1)$ example above for QED (one generator), it can be shown that the conserved current is

$$j^\mu = \bar{\psi}\gamma^\mu\psi. \quad (1.31)$$

Examples of continuous space-time symmetries and their associated conserved quantities are listed in Table 1.1. The Standard Model was finally built by looking at the observed symmetries of the systems, and writing the most general Lagrangian which follows such symmetries. This is quickly shown in Section 1.1.9.

Table 1.1 – Examples of symmetries and their conserved quantities.

Symmetry	Conserved quantity
Time translation	Energy
Space translation	Momentum
Space rotation	Angular momentum

Among other symmetries, it is worth mentioning the C , P and T symmetries. C is called *charge conjugation* and transforms a particle into its anti-particle. P and T are space-time

symmetries called *parity* and *time reversal* transformation respectively. They act on a four-vector $x^\mu = (t, \mathbf{x})$ such that $T(x^\mu) = (-t, \mathbf{x})$, $P(x^\mu) = (t, -\mathbf{x})$. Note that these are discrete, not continuous symmetries, and are therefore not the object of Noether's theorem. By construction, any quantum field theory is Lorentz-invariant and therefore invariant under the *CPT* symmetry. However it will be shown that the *C*, *P*, and *T* are not exact symmetries themselves and are individually violated.

1.1.5 Non abelian symmetry groups: the example of $SU(3)$

In the above example of Section 1.1.3, the symmetry, although abstract, is relatively simple: only one generator so there is no issue of commutation between operators. $U(1)$ is then called an *abelian group* as discussed in Section 1.1.1. Thanks to a result from Yang and Mills in 1954 [5], the same reasoning holds for non-abelian groups. Looking at the $SU(3)$ group³, its group structure is the following

$$[\lambda^a, \lambda^b] = if^{abc}\lambda^c, \quad (1.32)$$

where the λ^a matrices are the 8 Gell-Mann matrices, *i.e.* the generators of $SU(3)$. If we apply this group of symmetries to the Dirac Lagrangian seen in Eq. 1.14 and require it to be gauge-invariant under

$$\psi \rightarrow e^{i\alpha^a(x)\lambda^a} \psi, \quad (1.33)$$

with ψ now having 3 components

$$\psi = \begin{pmatrix} \psi_1 \\ \psi_2 \\ \psi_3 \end{pmatrix}, \quad \bar{\psi} = (\bar{\psi}_1 \quad \bar{\psi}_2 \quad \bar{\psi}_3), \quad (1.34)$$

each ψ_i being a Dirac spinor, then we should define the covariant derivative

$$D_\mu = \partial_\mu - ig \frac{\lambda^a}{2} A_\mu^a(x), \quad (1.35)$$

with the eight new gauge fields $A_\mu^a(x)$ transforming as

$$A_\mu^a(x) \rightarrow A_\mu^a(x) + \frac{1}{g} \partial_\mu \alpha^a(x) + f^{abc} A_\mu^b(x) \alpha^c(x). \quad (1.36)$$

Note that this is similar to Eq. 1.23, with an additional term $f^{abc} A_\mu^b(x) \alpha^c(x)$ due to the non-abelian part of the group. One can then show that the free field structure for $A_\mu^a(x)$ is

$$F_{\mu\nu}^a = \partial_\mu A_\nu^a - \partial_\nu A_\mu^a + gf^{abc} A_\mu^b A_\nu^c. \quad (1.37)$$

This is again close to Eq. 1.26, but with the additional term $gf^{abc} A_\mu^b A_\nu^c$ due to the non-abelian structure of $SU(3)$. There is an interesting feature of the gauge field propagation in non-abelian theories. The development of $F_{\mu\nu}^a F^{\mu\nu,a}$ exhibits terms such as $\partial_\mu A_\mu^a - \partial_\nu A_\nu^a$ as seen in the abelian case, but there are also terms such as $ig f^{abc} A_\mu^b A_\nu^c (\partial^\mu A^{\nu,b} - \partial^\nu A^{\mu,c})$ and $g^2 f^{abc} f^{ade} A_\mu^b A_\nu^c A^{\mu,d} A^{\nu,e}$ which correspond to gauge field self interactions with 3 or 4 field propagators! In conclusion, non-abelian gauge theories contain gauge fields which are self interacting, and the additional Feynman diagrams from Figure 1.3 must be considered.

As no assumption on the abelian character of the group is made in Noether's theorem, it extends straightforwardly to non-abelian gauge theories and it can be shown that there is a

³↑ The *S* in $SU(n)$ mean “special”: $SU(n)$ is the subgroup of $U(n)$ such that $SU(n) = \{M \in U(n), \det(M) = 1\}$.

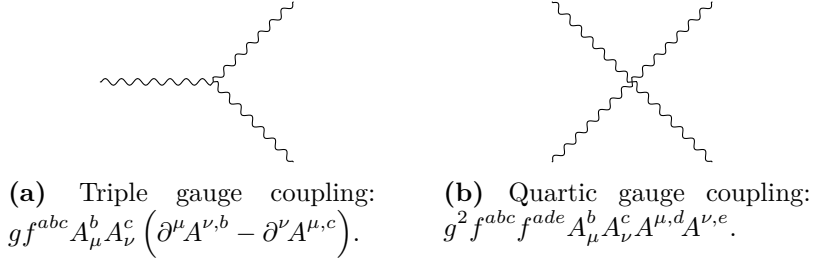


Figure 1.3 – Additional Feynman diagrams associated with the Dirac Lagrangian under non-abelian gauge symmetry $SU(n)$.

conserved current for each of the eight generators. Similarly to Eq. 1.31, one can therefore derive the conserved currents generically written as

$$j_a^\mu = \bar{\psi}\gamma^\mu\lambda_a\psi. \quad (1.38)$$

From the Standard Model viewpoint, $SU(3)$ will describe quantum chromodynamics (QCD), associated to quarks and gluons. Similarly to QED where the conserved current can be shown to ensure electric charge conservation, the QCD conserved currents ensure that the *colour charge* is conserved. We then denote the QCD Lagrangian

$$\mathcal{L}_{\text{QCD}} = -\frac{1}{4}F_{\mu\nu}^aF^{\mu\nu,a} + \bar{\psi}(i\not{D} - m)\psi \quad (1.39)$$

$$\begin{aligned} &= -\frac{1}{4}(\partial_\mu A_\nu^a - \partial_\nu A_\mu^a)(\partial^\mu A^{\nu,a} - \partial^\nu A^{\mu,a}) \\ &\quad - g_s f^{abc} A_\mu^b A_\nu^c (\partial^\mu A^{\nu,b} - \partial^\nu A^{\mu,b}) \\ &\quad - \frac{g_s^2}{4} (f^{abc} A_\mu^b A_\nu^c) (f^{ade} A^{\mu,d} A^{\nu,e}) \\ &\quad + \bar{\psi}i\gamma^\mu\partial_\mu\psi - m\bar{\psi}\psi \\ &\quad + g_s \bar{\psi}\gamma^\mu A_\mu^a \lambda^a \psi. \end{aligned} \quad (1.40)$$

The reasoning presented here actually applies to $SU(n)$ for any n . In particular, another interesting group for the description of the Standard Model is $SU(2)$ which is used to model the weak nuclear interaction. The generators of this group are the already-mentioned Pauli matrices, therefore 3 additional bosons will be added.

1.1.6 Electroweak theory and missing mass terms

The unification of the weak and electromagnetic theories was proposed by Glashow [6], Salam [7] and Weinberg [8] in the 1960's to form the electroweak (EW) theory. The gauge symmetry of this group is described by $SU(2)_I \times U(1)_Y$, where I stands for the weak isospin and Y for the weak hypercharge. These are related to the electric charge by the Gell-Mann-Nishijima formula, independently introduced in Refs. [9] and [10]:

$$Q = \frac{Y}{2} + I_3. \quad (1.41)$$

The EW theory yields four gauge fields: three from $SU(2)_I$ which are denoted W_μ^a , and one from $U(1)_Y$ which is denoted B_μ .

The EW theory is a chiral theory, meaning it does not act the same way on right- and left-handed spinors. This is due to the parity violation of the weak interaction observed by Wu [11]

and later confirmed by Garwin, Lederman and Weinrich [12]. The left and right spinors ψ_L and ψ_R are defined from a spinor ψ by

$$\psi_L = \frac{1 - \gamma_5}{2} \psi \quad \text{and} \quad \psi_R = \frac{1 + \gamma_5}{2} \psi \quad (1.42)$$

so that $\psi = \psi_L + \psi_R$, where $\gamma_5 = i\gamma_0\gamma_1\gamma_2\gamma_3 = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}$ with γ_i for $i = 0, 1, 2, 3$ are the Dirac matrices already mentioned (Eq. 1.17). The gauge transformation of $SU(2)_I$ transforms left doublets but not right singlets, so that the singlets do not interact through the weak interaction. The left doublets are composed of

$$\psi_L \in \left\{ \begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_L, \begin{pmatrix} \nu_\mu \\ \mu^- \end{pmatrix}_L, \begin{pmatrix} \nu_\tau \\ \tau^- \end{pmatrix}_L, \begin{pmatrix} u_\alpha \\ d_\alpha \end{pmatrix}_L, \begin{pmatrix} c_\alpha \\ s_\alpha \end{pmatrix}_L, \begin{pmatrix} t_\alpha \\ b_\alpha \end{pmatrix}_L \right\}, \quad (1.43)$$

while the right singlets are composed of

$$\psi_R \in \left\{ e_R^-, \mu_R^-, \tau_R^-, u_{R,\alpha}, c_{R,\alpha}, t_{R,\alpha}, d_{R,\alpha}, s_{R,\alpha}, b_{R,\alpha} \right\}, \quad (1.44)$$

with α running over the colour indices. An important point is that neutrinos are not included in the right spinors: indeed right-handed neutrinos have not yet been observed. This makes an important difference with respect to the spontaneous symmetry breaking explained below, and neutrinos are considered to not acquire their mass through this mechanism (the neutrino mass mechanism is not detailed here). The various fermions and their properties under $SU(2)_I \times U(1)_Y$ are summarised in Table 1.2.

Table 1.2 – Summary table of the Standard Model fermion content and their properties in the electroweak theory. The ν_R would have 0 charge for all columns, meaning it does not couple to the electroweak interaction. Q is the electric charge, related to the third component I_3 of the weak isospin I and to the weak hypercharge Y by Eq. 1.41. α runs over the colour indices.

Index notation	First family	Second family	Third family	Q	I	I_3	Y
$L_L^i = \begin{pmatrix} \nu_L^i \\ \ell_L^i \end{pmatrix}$	$\begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_L$	$\begin{pmatrix} \nu_\mu \\ \mu^- \end{pmatrix}_L$	$\begin{pmatrix} \nu_\tau \\ \tau^- \end{pmatrix}_L$	0 -1	1/2 1/2	1/2 -1/2	-1 -1
$Q_L^i = \begin{pmatrix} u_L^i \\ d_L^i \end{pmatrix}$	$\begin{pmatrix} u_\alpha \\ d_\alpha \end{pmatrix}_L$	$\begin{pmatrix} c_\alpha \\ s_\alpha \end{pmatrix}_L$	$\begin{pmatrix} t_\alpha \\ b_\alpha \end{pmatrix}_L$	2/3 -1/3	1/2 1/2	1/2 -1/2	1/3 1/3
ℓ_R^i	e_R^-	μ_R^-	τ_R^-	-1	0	0	-2
u_R^i	$u_{R,\alpha}$	$c_{R,\alpha}$	$t_{R,\alpha}$	2/3	0	0	4/3
d_R^i	$d_{R,\alpha}$	$s_{R,\alpha}$	$b_{R,\alpha}$	-1/3	0	0	-2/3

The Lagrangian for the EW theory therefore reads

$$\mathcal{L}_{\text{EW}}^{\#} = -\frac{1}{4} W_{\mu\nu}^a W^{\mu\nu,a} - \frac{1}{4} B_{\mu\nu} B^{\mu\nu} + \sum_{\text{L fermions}} i\bar{\psi}_L \not{D} \psi_L + \sum_{\text{R fermions}} i\bar{\psi}_R \not{D} \psi_R, \quad (1.45)$$

where the $\#$ will be explained later. The covariant derivative D_μ is defined by

$$D_\mu = \partial_\mu - ig \frac{\sigma_a}{2} W_\mu^a - ig' \frac{Y}{2} B_\mu, \quad (1.46)$$

where g and g' are arbitrary constants called *coupling constants* in the following. From that, we have the usual

$$B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu \quad (1.47)$$

and

$$W_{\mu\nu}^a = \partial_\mu W_\nu^a - \partial_\nu W_\mu^a - g\epsilon^{abc}W_\mu^b W_\nu^c, \quad (1.48)$$

where ϵ^{abc} is the fully antisymmetric Levi-Civita tensor.

One might have remarked that contrary to the examples developed in Sections 1.1.3 and 1.1.5, no mass term is present in this Lagrangian, even for the fermion fields. Indeed, if we try to add a mass to the fermions, then the Lagrangian includes a mass term

$$m\bar{\psi}\psi = m(\bar{\psi}_L + \bar{\psi}_R)(\psi_L + \psi_R) = m(\bar{\psi}_L\psi_L + \bar{\psi}_R\psi_R + \bar{\psi}_L\psi_R + \bar{\psi}_R\psi_L), \quad (1.49)$$

and we have

$$\bar{\psi}_L\psi_L = \psi^\dagger \frac{1-\gamma^5}{2} \gamma^0 \frac{1-\gamma^5}{2} \psi = \psi^\dagger \gamma^0 \frac{1+\gamma^5}{2} \frac{1-\gamma^5}{2} \psi = 0 \quad (1.50)$$

and similarly for ψ_R . The mass term must therefore come from $\bar{\psi}_L\psi_R + \bar{\psi}_R\psi_L$. However, this expression is not gauge invariant under $SU(2)_I$, thus one must have $m = 0$. Moreover, as previously discussed, mass terms for the gauge fields are also forbidden. This is an issue since both quarks and charged leptons on the one hand, and the weak bosons W^+ , W^- and Z^0 are known to be massive.

1.1.7 Spontaneous symmetry breaking, BEH mechanism, and gauge boson masses

The solution to both of these problems comes from the spontaneous symmetry breaking mechanism. This idea originally came from solid-state physics and magnetism and was later applied to particle physics by Nambu and Jona Lasinio [13–16]. Guralnik, Hagen and Kibble [17], Brout and Englert [18], and Higgs [19–21] all independently applied the idea to gauge theories, leading to what is now called the *Higgs* or *BEH mechanism*.

A symmetry is said to be spontaneously broken when it is a symmetry of the Lagrangian but not of the ground state of the system. An example of such system is a pencil standing perfectly vertically in equilibrium on a table, which presents a rotational symmetry around the vertical axis. However, standing vertically is not a stable position for the pencil, *i.e.* this is not its ground state. Once the pencil has fallen, due to a small perturbation, it lies on the table pointing in a specific direction and so the rotational symmetry is now broken.

The idea developed by Brout, Englert and Higgs consists of introducing an $SU(2)$ doublet of a complex scalar field ϕ with weak hypercharge $Y = 1$ defined by

$$\phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix} = \frac{1}{\sqrt{2}} \begin{pmatrix} \phi_1 + i\phi_2 \\ \phi_3 + i\phi_4 \end{pmatrix} \quad (1.51)$$

and with potential

$$V(\phi^\dagger\phi) = -\mu^2\phi^\dagger\phi + \lambda|\phi^\dagger\phi|^2, \quad (1.52)$$

where μ and λ are positive, real constants. The minimas of $V(\phi^\dagger\phi)$ are solutions of

$$\phi^\dagger\phi = \frac{1}{2}(\phi_1^2 + \phi_2^2 + \phi_3^2 + \phi_4^2) = \frac{\mu^2}{2\lambda}, \quad (1.53)$$

which is a sphere in the 4-dimensional space. At this point, the Goldstone theorem [22] states that one massless state, a *Goldstone boson*, appears for each broken symmetry.

Bearing in mind that we eventually want to break $SU(2)_I \times U(1)_Y$ to allow the W^+ , W^- and Z^0 masses to appear but keep $U(1)_{\text{EM}}$ unbroken and thus the photon massless, one particular but common choice of solution is to impose $\phi_1 = \phi_2 = \phi_4 = 0$. This is pictured in Figure 1.4.

Then, $\phi_3^2 = \frac{\mu^2}{\lambda}$ which we shall denote v^2 and will be later called the *vacuum expectation value* of the field ϕ . By choosing a particular direction ϕ_3 for the ground state, we have broken the symmetry from Eq. 1.53. With such a choice, we can rewrite $\phi = \begin{pmatrix} 0 \\ v \end{pmatrix}$, and then expand ϕ at first order around its minimum:

$$\phi = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + h(x) \end{pmatrix}, \quad (1.54)$$

$h(x)$ being a real function.

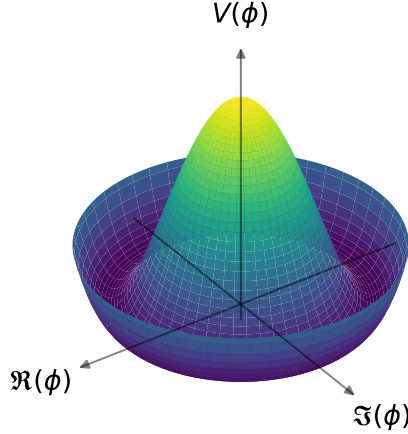


Figure 1.4 – Illustration of the Higgs potential $V(\phi^\dagger\phi) = -\mu^2\phi^\dagger\phi + \lambda|\phi^\dagger\phi|^2$. The potential has a rotational symmetry around the point $|\phi| = 0$, but its minimum (ground state) is not at this value.

The Lagrangian describing such a complex scalar field doublet, and left invariant by a $SU(2)_I \times U(1)_Y$ gauge symmetry reads

$$\mathcal{L}_{\text{EW}}^{\text{bosonic}} = -\frac{1}{4}W_{\mu\nu}^a W^{\mu\nu,a} - \frac{1}{4}B_{\mu\nu}B^{\mu\nu} + (D_\mu\phi)^\dagger(D^\mu\phi) - V(\phi^\dagger\phi), \quad (1.55)$$

with D_μ defined as in Eq. 1.46: $D_\mu = \partial_\mu - ig\frac{\sigma_a}{2}W_\mu^a - ig'\frac{Y}{2}B_\mu$, with $Y = 1$ by hypothesis. If we put Eq. 1.54 into Eq. 1.55, several interesting terms appear. First we get terms like

$$\mathcal{L}_{\text{EW}}^{\text{bosonic}} \supset \frac{1}{2}\partial_\mu h \partial^\mu h - \frac{v^2}{2}\lambda h^2 - \lambda v h^3 - \frac{\lambda}{4}h^4 + \dots, \quad (1.56)$$

which describes the propagation of a scalar field ϕ with mass

$$m_H = \sqrt{2\lambda v^2}. \quad (1.57)$$

We additionally get triple and quartic self couplings for the Higgs field. These corresponds to the Feynman diagrams of Figure 1.5.

Second, if we focus on

$$\left| \left(ig\frac{\sigma_a}{2}W_\mu^a + ig'\frac{Y}{2}B_\mu \right) \phi \right|^2 \quad (1.58)$$

which emerges from the covariant derivative, this term can be written after explicitly computing the contraction $\sigma^a W_\mu^a$ and expanding around the vacuum expectation value:

$$\left| \frac{-i}{2} \begin{pmatrix} gW_\mu^3 + g'B_\mu & g(W_\mu^1 - iW_\mu^2) \\ g(W_\mu^1 + iW_\mu^2) & -gW_\mu^3 + g'B_\mu \end{pmatrix} \begin{pmatrix} 0 \\ v \end{pmatrix} \right|^2. \quad (1.59)$$

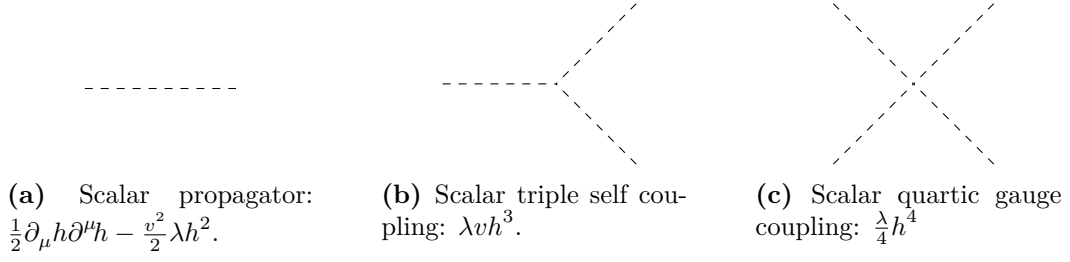


Figure 1.5 – Feynman diagrams associated with complex scalar field propagation with potential as in Eq. 1.52: $V(\phi^\dagger\phi) = -\mu^2\phi^\dagger\phi + \lambda|\phi^\dagger\phi|^2$.

A similar term with h instead of v will appear, leading to a direct interaction between the scalar and the gauge bosons as pictured in Figure 1.6.

Then if we define

$$W_\mu^\pm = \frac{W_\mu^1 \mp iW_\mu^2}{\sqrt{2}}, \quad (1.60)$$

we can simplify Eq. 1.58 into

$$\mathcal{L}_{\text{EW}}^{\text{bosonic}} \supset \frac{1}{4}v^2 g^2 W_\mu^+ W^{\mu-} + \frac{v^2}{8} \begin{pmatrix} W_\mu^3 & B_\mu \end{pmatrix} \begin{pmatrix} g^2 & -gg' \\ -gg' & g'^2 \end{pmatrix} \begin{pmatrix} W^{\mu,3} \\ B^\mu \end{pmatrix}. \quad (1.61)$$

From this, we can see that the W_μ^3 and B_μ states are mixed by the matrix. However, this matrix can be diagonalised, with eigenvalues 0 and $g^2 + g'^2$. After defining

$$\sin \theta_W = \frac{g'}{\sqrt{g^2 + g'^2}} \quad \text{and} \quad \cos \theta_W = \frac{g}{\sqrt{g^2 + g'^2}}, \quad (1.62)$$

we obtain

$$\begin{pmatrix} Z_\mu \\ A_\mu \end{pmatrix} = \begin{pmatrix} \cos \theta_W & -\sin \theta_W \\ \sin \theta_W & \cos \theta_W \end{pmatrix} \begin{pmatrix} W_\mu^3 \\ B_\mu \end{pmatrix} \quad (1.63)$$

and θ_W is called the *Weinberg angle*. Replacing in Eq. 1.55, we have

$$\mathcal{L}_{\text{EW}}^{\text{bosonic}} \supset -\frac{1}{2}W_{\mu\nu}^+ W^{-\mu\nu} - \frac{1}{4}Z_{\mu\nu}Z^{\mu\nu} - \frac{1}{4}A_{\mu\nu}A^{\mu\nu} + m_W^2 W_\mu^+ W^{-\mu} + \frac{1}{2}m_Z^2 Z_\mu Z^\mu, \quad (1.64)$$

with

$$m_W^2 = \frac{1}{4}g^2 v^2 \quad \text{and} \quad m_Z^2 = \frac{1}{4}v^2(g^2 + g'^2), \quad (1.65)$$

and $m_A = 0$: we have recovered the physical W, Z and γ bosons. We can finally deduce a relation between the W and the Z mass:

$$\cos \theta_W = \frac{m_W}{m_Z}. \quad (1.66)$$

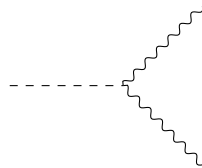


Figure 1.6 – Feynman diagram associated with the interaction between the symmetry breaking complex scalar field and the gauge bosons.

1.1.8 Fermion masses

In Section 1.1.7, we showed how gauge bosons acquire masses through the spontaneous symmetry breaking mechanism. The fermion masses arise from additional terms that we can add to the Lagrangian which will now be gauge invariant thanks to the complex scalar field doublet (Higgs field). These terms are called fermion Yukawa terms.

For the leptons, using the notations from Table 1.2, the Yukawa Lagrangian reads

$$\mathcal{L}_Y^{\text{leptons}} = -y_i^\ell (\bar{L}_L^i \phi \ell_R^i + h.c.). \quad (1.67)$$

Expanding the Higgs field near its vacuum expectation value,

$$\mathcal{L}_Y^{\text{leptons}} = -\frac{y_i^\ell}{\sqrt{2}} v (\bar{\ell}_L^i \ell_R^i + h.c.) - \frac{y_i^\ell}{\sqrt{2}} h (\bar{\ell}_L^i \ell_R^i + h.c.), \quad (1.68)$$

where the y_i^ℓ are the Higgs Yukawa couplings of leptons, and from which we identify the mass term

$$m_\ell^i = \frac{y_i^\ell}{\sqrt{2}} v \quad (1.69)$$

and the Higgs couplings to leptons term. By *h.c.* we denote the *hermitian conjugate*, *i.e.* the † operation of the preceding term. The procedure can be repeated for each lepton family $i = 1, 2, 3$. The associated Feynman diagram for the interaction between the Higgs field and the leptons is shown in Figure 1.7.

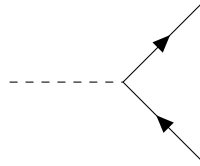


Figure 1.7 – Feynman diagram associated with the interaction between the symmetry breaking complex scalar field and the fermions: $\frac{y}{\sqrt{2}} h (\bar{\psi}_L \psi_R + h.c.)$.

For the quarks, the Yukawa Lagrangian reads

$$\mathcal{L}_Y^{\text{quarks}} = -(y_{ij}^u \bar{Q}_L^i \tilde{\phi} u_R^j + y_{ij}^d \bar{Q}_L^i \phi d_R^j + h.c.), \quad (1.70)$$

where $\tilde{\phi} = i \frac{\sigma^2}{2} \phi^*$, i, j runs over the quark families. The y_{ij}^u and y_{ij}^d matrices are the Higgs Yukawa couplings to up-type and down-type quarks respectively. After expanding near the vacuum expectation value, the quark Yukawa Lagrangian can be written (not focusing on the interaction terms which will yield the same interaction diagram as the one depicted for the leptons in Figure 1.7)

$$\mathcal{L}_Y^{\text{quarks}} = -(\bar{u}_L^i m_{ij}^u u_R^j + \bar{d}_L^i m_{ij}^d d_R^j + h.c.) + \text{interaction}, \quad (1.71)$$

where the m_{ij}^u and m_{ij}^d matrices are the up-type and down-type quark mass matrices. Due to these m_{ij} matrices, the physical quark states are mixed. However, these matrices can be diagonalised to recover the physical mass states. Denoting

$$u_L^i = V_u^{ij} u_L'^j \quad \text{and} \quad d_L^i = V_d^{ij} d_L'^j, \quad (1.72)$$

with u_L^i, d_L^i the quark interaction eigenstates and $u_L'^i, d_L'^i$ the quark mass eigenstates, we get:

$$\mathcal{L}_Y^{\text{quarks}} = - \left(\bar{u}_L'^i \begin{pmatrix} m_u & 0 & 0 \\ 0 & m_c & 0 \\ 0 & 0 & m_t \end{pmatrix}_{ij} u_R'^j + \bar{d}_L'^i \begin{pmatrix} m_d & 0 & 0 \\ 0 & m_s & 0 \\ 0 & 0 & m_b \end{pmatrix}_{ij} d_R'^j + h.c. \right). \quad (1.73)$$

If we now turn to the $W^\pm$ interaction with quarks coming from Lagrangian terms as $\bar{Q}_L^i i \not{D} Q_L^i$, we have

$$\bar{u}_L^i \gamma^\mu d_L^i W_\mu^+ + \bar{d}_L^i \gamma^\mu u_L^i W_\mu^- = \bar{u}_L^i \gamma^\mu (V_u^\dagger V_d)_{ij} d_L^j W_\mu^+ + \bar{d}_L^i \gamma^\mu (V_d^\dagger V_u)_{ij} u_L^j W_\mu^-. \quad (1.74)$$

The $(V_u^\dagger V_d)$ defines the so-called CKM matrix in honour of Cabbibo [23], Kobayashi and Maskawa [24] who introduced the quark mixing and CP -violation mechanism, mixing the mass eigenstates to the interaction eigenstates:

$$\begin{pmatrix} d' \\ s' \\ b' \end{pmatrix} = V_{\text{CKM}} \begin{pmatrix} d \\ s \\ b \end{pmatrix} = \begin{pmatrix} V_{ud} & V_{us} & V_{ub} \\ V_{cd} & V_{cs} & V_{cb} \\ V_{td} & V_{ts} & V_{tb} \end{pmatrix} \begin{pmatrix} d \\ s \\ b \end{pmatrix}. \quad (1.75)$$

Choosing it to act on the down-type quarks is just a matter of convention.

1.1.9 Mathematical formulation of the Standard Model

We now have all the building blocks to construct the Standard Model. The gauge theory is build by

- taking the particle content of the theory, *i.e.* the fermions,
- choosing the group invariances (the gauge symmetries), from which the gauge fields and their interactions will appear,
- including one or several symmetry breaking mechanisms and their corresponding fields.

The fermion content can be chosen arbitrarily. However since the goal is to describe the observations, common sense requires to start from the observed quarks and leptons, referring to the usual 6 quarks (u, d, s, c, t, b) , 3 charged leptons (e, μ, τ) and 3 neutral leptons $(\nu_e, \nu_\mu, \nu_\tau)$. One then impose the gauge symmetry $SU(3)_C \times SU(2)_I \times U(1)_Y$. This is a group product, meaning the Dirac spinors will be composed in a similar way as in Eq. 1.34, each component interacting either with $SU(3)_C$ or $SU(2)_I$ or $U(1)_Y$. This symmetry was chosen as the most simple group reflecting the experimental observations: the $SU(3)_C$ symmetry gives birth to the 8 gluons as seen in Section 1.1.5, and the $SU(2)_I \times U(1)$ leads to the Z, W^+, W^- and γ bosons as described in Section 1.1.7. However, the choice of symmetries is not limited to those: for example, it is possible to extend the Standard Model and impose $SU(5)$ or $SO(10)$ as a gauge group, so that new properties and therefore particles emerge. Finally, the scalar field (Higgs field) is introduced to include mass terms for both fermions and gauge bosons, breaking the $SU(2)_I \times U(1)_Y$ and recovering the physical states for the electromagnetic and weak bosons.

This results in the following (contracted) Lagrangian:

$$\mathcal{L}_{\text{SM}} = -\frac{1}{4} \sum_{a=1}^8 G_{\mu\nu}^a G^{\mu\nu,a} - \frac{1}{4} \sum_{b=1}^3 W_{\mu\nu}^b W^{\mu\nu,b} - \frac{1}{4} B_{\mu\nu} B^{\mu\nu} \quad (1.76)$$

$$+ (D_\mu \phi)^\dagger (D^\mu \phi) - V(\phi^\dagger \phi) \quad (1.77)$$

$$+ \sum_{i=1}^3 -y_i^\ell (\bar{L}_L^i \phi \ell_R^i + h.c.) \quad (1.78)$$

$$+ \sum_{a=1}^3 \sum_{i=1}^3 -(y_{ij}^u \bar{u}_L^{a,i} u_R^{a,j} \tilde{\phi} + y_{ij}^d \bar{d}_L^{a,i} d_R^{a,j} \phi + h.c.) \quad (1.79)$$

$$+ \sum_{i=1}^3 i \bar{L}_L^i \not{D} L_L^i + \bar{\ell}_R^i \not{D} \ell_R^i \quad (1.80)$$

$$+ \sum_{a=1}^3 \sum_{i=1}^3 i \bar{Q}_L^{a,i} \not{D} Q_L^{a,i} + \bar{u}_R^{a,i} \not{D} u_R^{a,i} + \bar{d}_R^{a,i} \not{D} d_R^{a,i}, \quad (1.81)$$

using the notations from Table 1.2 where the a index runs over the colour and i over the lepton families, and with the covariant derivative

$$D_\mu = \partial_\mu - ig_s \frac{\lambda^a}{2} G_\mu^a - ig \frac{\sigma^a}{2} W_\mu^a - ig' \frac{Y}{2} B_\mu. \quad (1.82)$$

To summarise this Lagrangian,

- line 1.76 gives the free propagating terms of the gauge fields, and their self interaction,
- line 1.77 includes the Higgs free propagation, mass term and self couplings, as well as the gauge boson interaction with the Higgs field, *i.e.* the gauge boson mass terms,
- lines 1.78 and 1.79 describe the leptons and quarks Yukawa couplings, that is their masses and interaction with the Higgs field,
- and finally, lines 1.80 and 1.81 represent the propagation of leptons and quarks, as well as their interaction with the gauge bosons.

Note that the Higgs boson does not carry any colour charge and therefore does not directly interact with the gluons, which consequently are massless. Similarly, photons do not directly couple to the Higgs since they are massless. It is also worth noting that the Higgs coupling to fermions is proportional to their masses (Eqs. 1.68 and 1.71), while it is proportional to the squared masses in the case of massive gauge bosons (Eq. 1.64).

The Standard Model (SM) as it stands includes 19 *free* parameters, meaning they are not predicted by the above theoretical formulation and can only be determined by experimental measurements. Many relations link the SM parameters to the others, so the actual choice of free parameter set is a matter of convention (for example $\tan \theta_W = g'/g$ or $m_H = \sqrt{2\lambda v^2}$). A common list is the following:

- the 6 quark masses ($m_u, m_d, m_c, m_s, m_t, m_b$),
- the 3 charged leptons masses (m_e, m_μ, m_τ),
- 3 angles from the CKM matrix ($\theta_{12}, \theta_{23}, \theta_{13}$, plus 1 CP violation phase (δ_{CP}),
- the QCD vacuum angle, also called strong CP violation parameter (θ_{QCD}),
- the 3 gauge couplings for $U(1), SU(2), SU(3)$ (g', g, g_s respectively),
- the Higgs vacuum expectation value (v), and the Higgs boson mass (m_H).

Finally it can be shown that on top of the space-time symmetries mentioned in Table 1.1 and to the gauge symmetries imposing the conservation of colour, weak isospin, hypercharge and electric charge, the Standard Model includes 4 additional *accidental* symmetries (symmetries that are not imposed by the gauge structure) leading to the conservation of the electron, muon, tau numbers (commonly referred to as *lepton number* L) and baryon number (B).

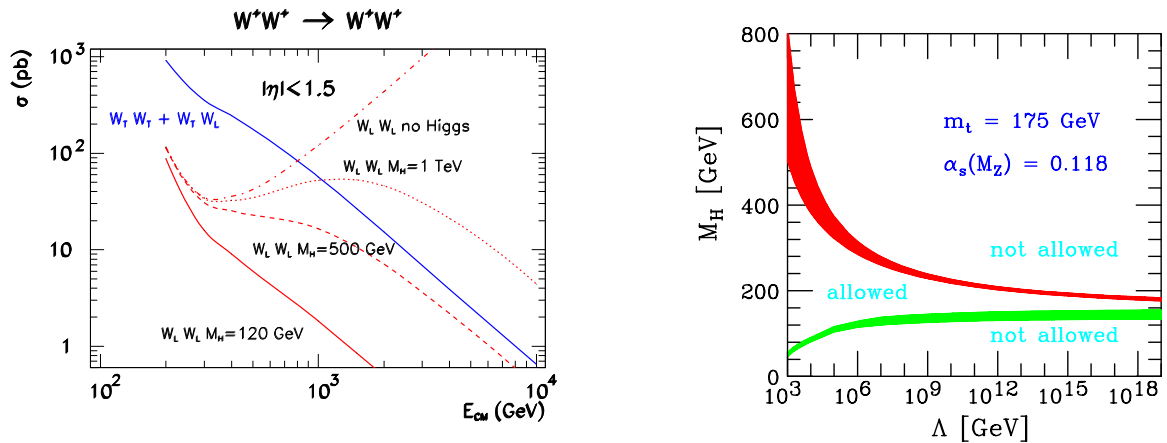
1.2 Discovery of the Higgs boson and first results

1.2.1 Theoretical constraints before the discovery

While the Higgs boson mass is a free parameter of the Standard Model, indirect theoretical considerations could still constrain the range of acceptable mass. These considerations are not detailed here, but in short they are mainly three-fold:

- if the Higgs boson mass is too large, the same-sign $W_L^\pm W_L^\pm \rightarrow W_L^\pm W_L^\pm$ longitudinal scattering diverges as the centre-of-mass energy increases [25]. This is called a violation of *unitarity*, where the probability of an interaction increases beyond one, and is depicted in Figure 1.8a.
- As the energy scale increases (which can be considered as the energy scale where new physics appears), the Higgs self-coupling increases up to a point of divergence [26]. This is called *triviality* and places an upper bound on the Higgs boson mass as shown by the red curve of Figure 1.8b.
- On the contrary, if the self-coupling decreases too much, this leads to a so-called *vacuum instability* [27, 28]. This constraint, depicted by the green curve of Figure 1.8b, places a lower bound on Higgs boson mass.

These theoretical constraints mainly led to an upper bound of ~ 200 GeV to 1 TeV on the Higgs boson mass at the 95% confidence level (CL), depending on the energy scale (see for example Ref. [29] for a more extensive discussion).

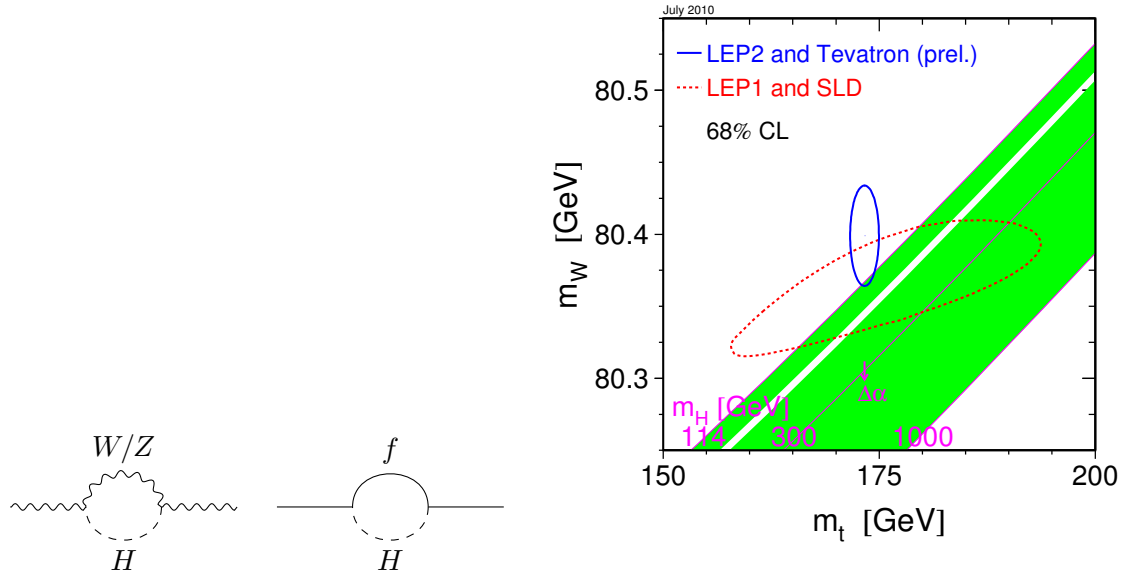


(a) Same sign $W_L^\pm W_L^\pm$ scattering cross-section as a function of the centre-of-mass energy for various values of the Higgs boson mass (W_T and W_L refer to polarisations states).

(b) Higgs boson mass constraints from triviality (red) and vacuum stability arguments (green) as a function of some cut-off scale Λ .

Figure 1.8 – Theoretical Higgs mass constraints from (a) same-sign $W_L^\pm W_L^\pm$ scattering [30], (b) triviality and vacuum stability [29].

Moreover, the Higgs boson also appears in loops, *i.e.* higher order electroweak radiative corrections to the propagators and interactions. The amplitude of these corrections depends on the Higgs boson mass as well as the top quark and W boson masses. With a precise knowledge of the top quark and W boson masses one can then put other indirect constraints on the Higgs boson mass as shown in Figure 1.9. These indirect measurements put a upper limit of the order of ≈ 260 GeV on the Higgs mass at the 95 % CL [31].



(a) Example propagator corrections due to the Higgs field.

(b) Theoretical constraints on the Higgs boson mass from the top quark and W boson mass measurements.

Figure 1.9 – Interplay between the top quark, W and Higgs bosons masses in the Standard Model. [32]

Finally, even if its mass was only loosely constrained, some of the Higgs properties could already be deduced from the theory developed above in Section 1.1. For example, its couplings to the fermions (proportional to their masses) and to the gauge bosons (proportional to their masses squared) are predicted under the Standard Model hypothesis. Other Higgs boson observables as well depend on its mass, and these dependencies are known from the theory. For example its *production cross-section* (*i.e.* probability of production) at proton–proton colliders, its *total width* (linked to its mean lifetime τ by $\Gamma = \hbar/\tau$), and *branching ratios*⁴ as a function of its mass are shown in Figure 1.10.

This prior knowledge of the Higgs boson properties and probable mass domains helped to design experiments and analyses to target these regions. This led to the first studies at LEP and the Tevatron which further constrained the Higgs boson mass, and eventually to the discovery at the LHC in 2012, described in Section 1.2.2.

1.2.2 Direct searches at LEP, Tevatron, and the LHC

As the large majority of particles, the Higgs boson is massive and unstable. As a consequence, excitations of the Higgs field need a large amount of energy to give birth to its physical particle, and the Higgs boson will not appear spontaneously. In order to observe it, we need to artificially excite the Higgs field. This can be achieved at particle colliders: by annihilating particles that

⁴↑ The branching ratio can be understood as the probability for an unstable particle to decay into a specific state. It is defined as $BR_f = \Gamma_f/\Gamma_{\text{tot}}$ where Γ_i is the *partial width* of the final state f such that $\sum_i \Gamma_i = \Gamma_{\text{tot}}$.

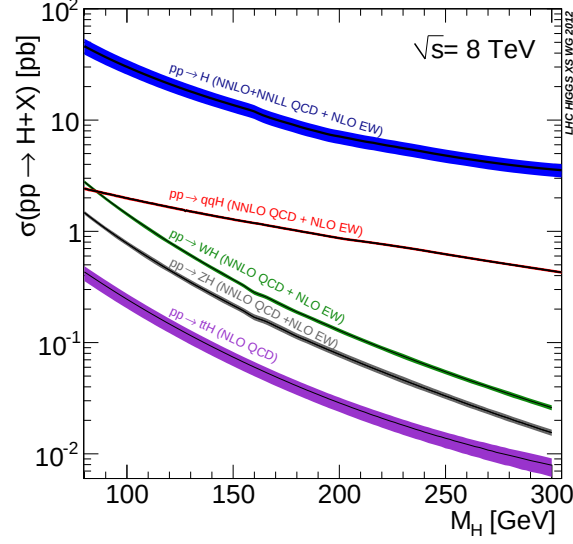
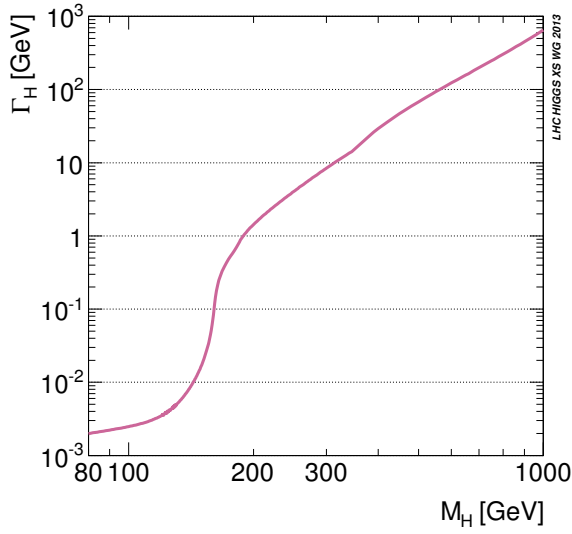
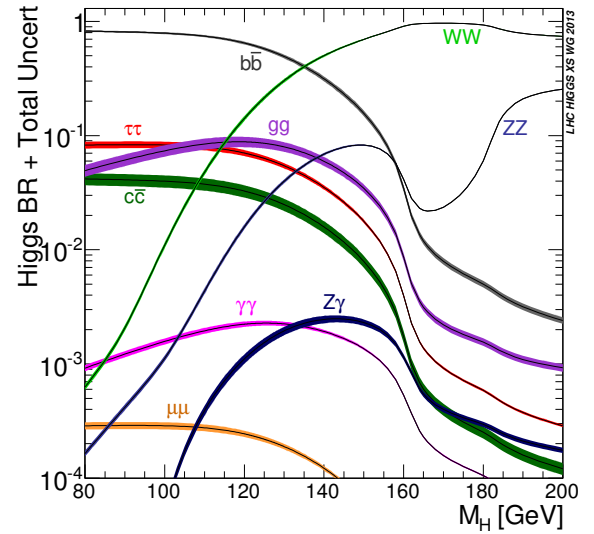
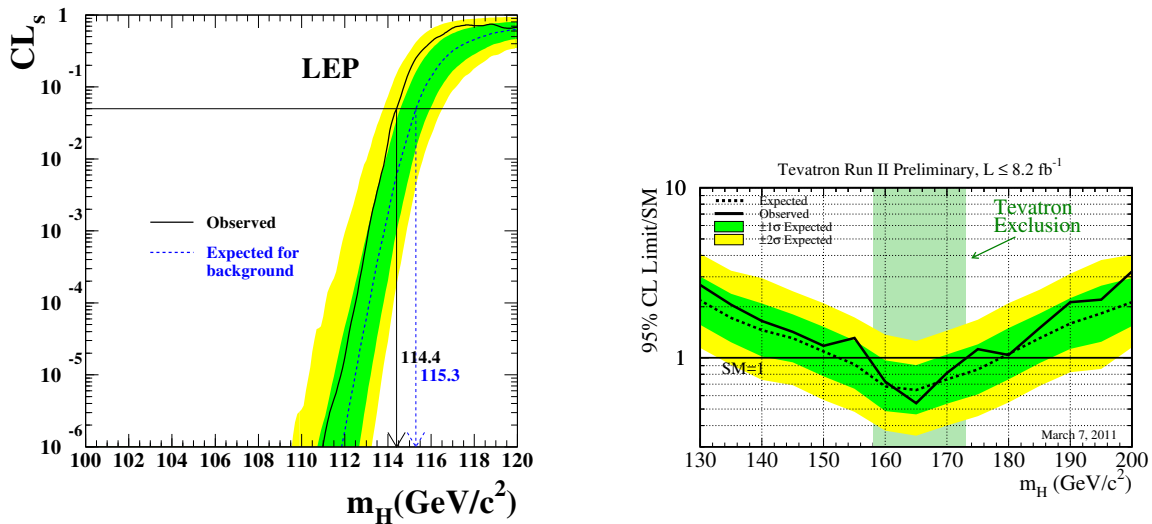

 (a) $\sigma(pp \rightarrow H)$ at $\sqrt{s} = 8 \text{ TeV}$.

 (b) Γ_H .

 (c) $\text{BR}(H \rightarrow X)$.

Figure 1.10 – Evolution of the Higgs boson production cross-section at $\sqrt{s} = 8 \text{ TeV}$ proton–proton colliders (a), total width (b) and branching ratios (c) as a function of the Higgs boson mass. More details will be given in Chapter 7.1. [33, 34]

interact (directly or indirectly) with the Higgs boson, sufficient energy might be released to excite the Higgs field up to a materialisation of a Higgs boson.

First experiments occurred at Fermilab’s Tevatron and CERN’s LEP (Large Electron Positron) colliders, but the direct searches were unsuccessful. The precision measurements conducted by their experiments however helped the theory putting the constraints listed above in Section 1.2.1. LEP eventually stopped in 2000 to prepare for the Large Hadron Collider (LHC, see Chapter 2) and the Tevatron in 2011, shortly after the LHC started. In spite of not finding the Higgs boson, the experiments could exclude mass domains and set limits on possible Higgs boson mass, preparing the ground for the future discovery: by the time the LHC started, the ranges below ~ 115 GeV and between 158 and 185 GeV were excluded at the 95 % CL by the LEP [32] and Tevatron [35] experiments, respectively, as shown in Figure 1.11. Moreover, a fit combining all electroweak precision observables related to the Higgs mass and measured by the LEP and Tevatron experiments could derive an indirect prediction of the Higgs boson mass. The result of this combined fit is shown in Figure 1.12 an exhibits a best fit value around 95 GeV with however large uncertainties. The summary of these hints remarkably pointed toward a low mass Higgs boson, narrowing the likely mass search domain for the LHC (Chapter 2). This high luminosity proton–proton collider, installed in the former LEP tunnel at CERN, started providing high-energy physics collisions to its experiments in 2010.



(a) Higgs boson mass limits from LEP.

(b) 95 % CL Higgs boson mass exclusion limits in the range [100, 200] GeV.

Figure 1.11 – 95 % CL Higgs boson mass exclusion limits from (a) LEP [36] and (b) Tevatron [35].

As the Higgs boson is unstable and with a short lifetime (about 10^{-22} s), experiments do not directly detect it. Instead, the Higgs boson decays into other particles which can in turn be detected (eventually after a chain of decays). Examples of decays include the Higgs boson to a photon pair ($H \rightarrow \gamma\gamma$), the Higgs boson to a pair of Z bosons subsequently decaying into pairs of leptons ($H \rightarrow ZZ^{(*)} \rightarrow 4\ell$), or the Higgs boson decaying to a pair of bottom quarks ($H \rightarrow b\bar{b}$). The Higgs boson is then “reconstructed” from these final states. These decay channels are indicated in Figure 1.10c and more details are given in Chapter 7.

The ATLAS (described in Chapter 3) and CMS experiments were finally able to announce the discovery of a new particle in 2012 [37, 38]. This new particle appeared in the $H \rightarrow \gamma\gamma$, $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ and $H \rightarrow WW$ decay channels at a mass around 125 GeV as seen in Figure 1.13, with properties consistent with the Standard Model Higgs boson.

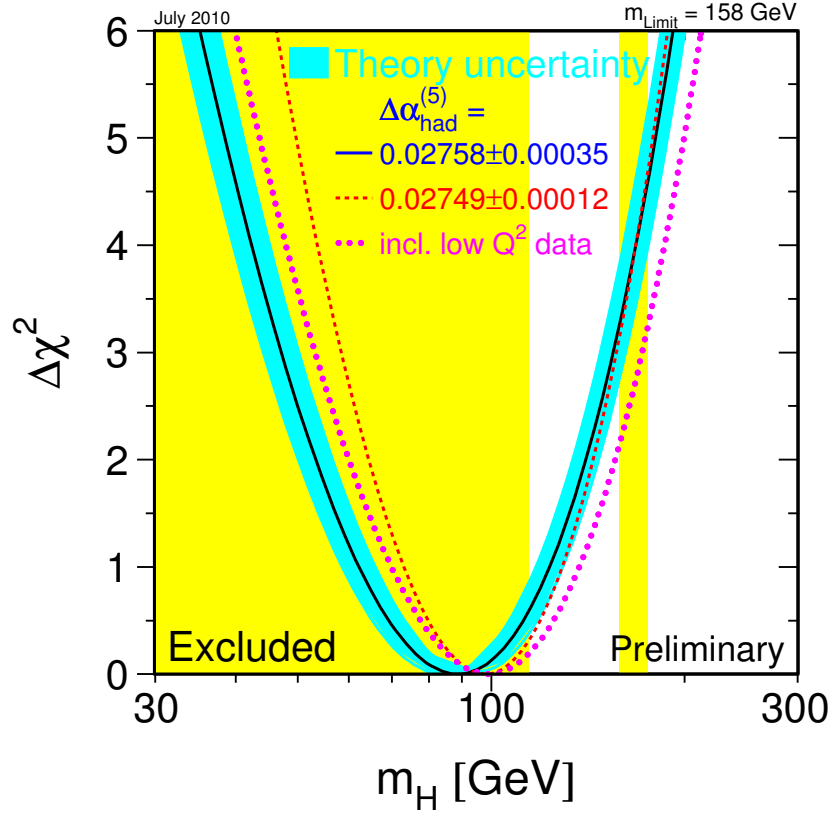
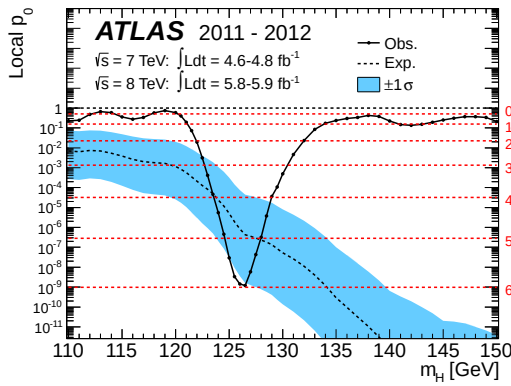
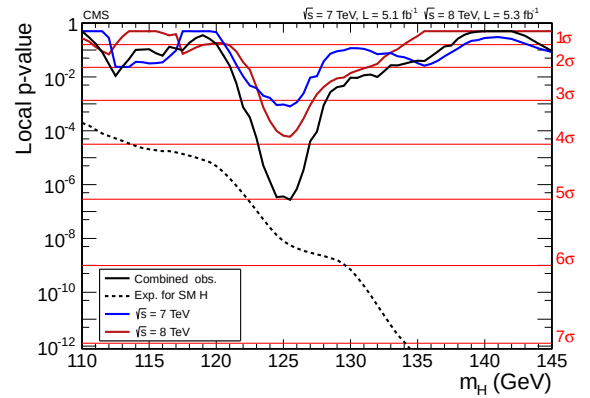


Figure 1.12 – χ^2 of the fit to all electroweak observables as a function of the Higgs mass. The yellow bands represent the indirect 95 % CL limits established by the LEP and Tevatron experiments. [32]



(a) ATLAS result.



(b) CMS result.

Figure 1.13 – Observed local p -value (probability of an excess of events to be consistent with the background) as a function of the hypothesised Higgs boson mass using the 2011 $\sqrt{s} = 7$ TeV and the early 2012 $\sqrt{s} = 8$ TeV data collected at the LHC for (a) ATLAS and (b) CMS. [37, 38]

1.2.3 Successes and limitations of the Standard Model

The Standard Model as described in Section 1.1 has proven extremely accurate in predicting the properties of particles and processes discovered so far, not only in the Higgs sector. As an example, one of its most famous realisation is the computation of the electromagnetic coupling constant (so called *fine structure constant*) at better than one part per billion [39]. Concerning the particle content, the Standard Model has successfully predicted the existence of the Higgs boson, and three (active) lepton families. Such precise predictions can be experimentally tested, and significant experimental deviations of the predictions can provide hints of physics beyond the Standard Model.

However, the tests performed at collider and non-collider experiments have until now proven to match the Standard Model predictions. For example the electromagnetic coupling constant is measured from the electron Landé factor with a precision and value matching the one from Standard Model prediction⁵. The Standard Model was also heavily tested at colliders, especially at LEP whose experiments measured the number of active neutrino families to 2.984 ± 0.008 [40, 41]. Other precision analyses conducted at LEP and Tevatron experiments for example measured the weak mixing angle θ_W and the W boson mass at the level of 0.2‰, and the Z boson mass at the level of a few dozen parts per million level [41]. Combining these many measurements provides an overconstrained system on Standard Model parameters, allowing for consistency tests of the theory.

Precision measurements are also being carried out at the LHC and many parameters will also be investigated, but its focus is at the moment on production cross-sections and decay branching ratios of processes. Figure 1.14 shows a few examples of the diversity of measurements achieved at the LHC, successfully testing the Standard Model over 14 orders of magnitude.

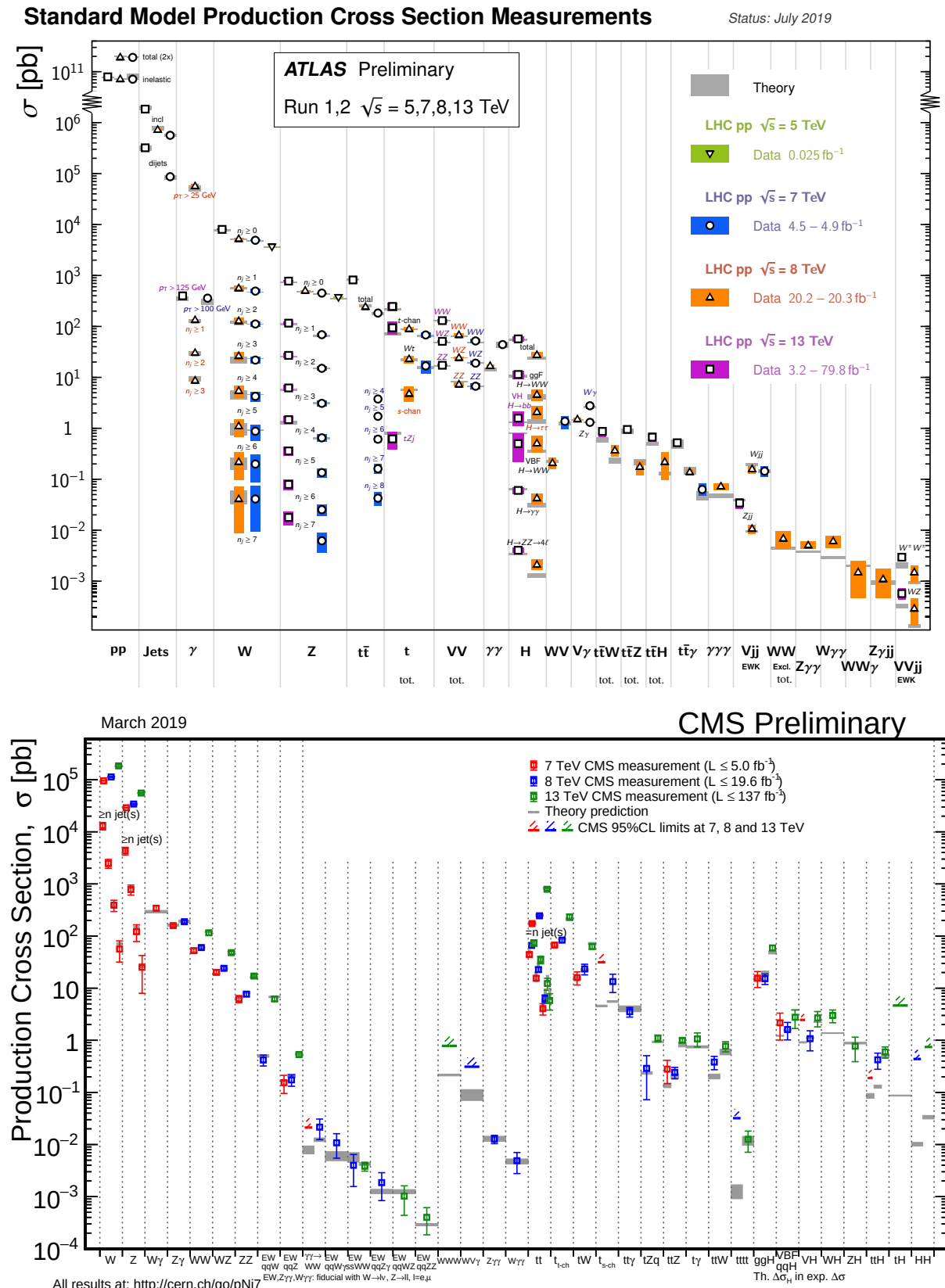
However it is known that the Standard Model cannot be the ultimate theory for particle physics. Indeed, several phenomena cannot be explained by the Standard Model alone, implying that it is consequently an incomplete theory. Examples of the Standard Model limitations follow.

While it provides a unified framework for the electromagnetism and the weak and strong nuclear interactions, gravity is excluded from this description. General relativity (GR), the mathematical theory that describes gravity, has successfully passed all its tests and provided prodigious prediction (such as gravitational waves [44, 45]) just as the Standard Model (SM), but marrying SM to GR has failed so far. Furthermore, the strength of the gravitational interaction is about 40 orders of magnitude lower than the strong nuclear interaction, posing the so-called *hierarchy* problem.

The Standard Model also predicts vanishing masses for the three neutrinos. However, a phenomenon called neutrino oscillation predicted by Pontecorvo [46, 47], observed a few years later [48, 49] and confirmed by several experiments, cannot occur with massless neutrinos. A minimal extension to the Standard Model Lagrangian can account for this phenomenon, but hypotheses are also formulated to describe the neutrino as a Majorana fermion (*i.e.*, the neutrino would be its own anti-particle) instead of a Dirac fermion.

Similarly to the Higgs boson electroweak corrections to the boson and fermion propagators introduced in Figure 1.9a, fermions and bosons also induce higher order corrections to the Higgs boson propagator. In the Higgs boson case however, the corrections induce a large divergence of the Higgs boson mass which has to be compensated with the introduction of a cut-off scale Λ . The drawback is that the theory is not valid above this cut-off scale. To keep the Higgs mass at reasonable values, the loop corrections must cancel with precisions of the order of 10^{-10} , which is known as the *naturalness*, or *fine-tuning*, problem.

⁵↑ Ref. [39] quotes $\alpha^{-1} = 137.035\,999\,173(35)$ (quantum electrodynamics theory), while CODATA quotes $\alpha^{-1} = 137.035\,999\,084(21)$ (experimental measurement).



The Standard Model predicts that matter and antimatter should have a symmetric behaviour, except for the CP symmetry violating terms. CP violation implies that slightly more matter (conventionally) is created than antimatter. However, the so-called *baryon asymmetry* observed at cosmological scales requires the CP violation sources to have an amplitude orders of magnitude higher than the one predicted by the Standard Model.

Another related issue is the so-called *strong CP* problem: while the electro-weak interaction does violate the CP symmetry, the analogue CP -violating parameter θ_{QCD} for the strong interaction is stringently constrained to about $|\theta_{\text{QCD}}| \lesssim 10^{-10}$ experimentally, while nothing suggests such cancellation in the Standard Model.

Finally, on the astrophysical side, baryonic matter represents less than 5% of the energy content of the Universe [50]. The rest would be made of *dark matter* ($\sim 20\%$), introduced by Zwicky [51] to explain the anomaly in the rotation of the galaxies, and of *dark energy*, introduced to explain the expansion of the Universe observed by Hubble [52].

The goal since the Higgs boson discovery is then to search for possible explanations of these unexplained behaviours, looking for deviations from the Standard Model predictions that could lead to new theoretical and experimental discoveries. A powerful framework to generically characterise such deviations is found in effective field theories (EFT), that are introduced in Chapter 7 and used in the main physics analysis of this Thesis described in Chapter 9.

1.2.4 Status of the Higgs boson after Run 1 at the LHC

After completion of Run 1 of the LHC and the Higgs boson discovery mentioned in Section 1.2.2, experiments turned to properties measurements of this new particle. As physics analysis topic of this Thesis (Chapters 8 and 9) use data from the Run 2 of the LHC, this section provides a quick overview of the Higgs results after Run 1 completion.

One of the most crucial measurements is of course the Higgs boson mass, as it is the last free parameter in the Standard Model. The ATLAS and CMS combination in the $H \rightarrow \gamma\gamma$ and $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ channels [53–55] summarised in Figure 1.15 provided a measurement with a precision of about 2%:

$$m_H = 125.09 \pm 0.21 \text{ (stat)} \pm 0.11 \text{ (syst)} \text{ GeV.} \quad (1.83)$$

As discussed in Section 1.2.3, other strong tests of the Standard Model reside in the measurements of production cross-sections and branching ratios. Several analyses mainly using the bosonic decays channels were performed by the ATLAS [56–59] and CMS [60–62], and eventually combined [63–65]. Example results are shown in Figure 1.16 and exhibit a good agreement for both the various signal strengths and the couplings.

Finally, as the Standard Model predicts that the Higgs boson should be a scalar, *i.e.* have a spin-parity $J^{PC} = 0^{++}$ (abridged $J^P = 0^+$), this hypothesis was also tested. First analyses conducted by ATLAS [66–68] and CMS [61, 62, 69, 70] excluded pure 0^- , 1 and 2^+ states better than the 99.9% CL as shown in Figure 1.17. Experiments then turned to exploring exotic states as in Refs. [67, 68, 71] and [64, 70], looking for CP -admixture of the Higgs boson. Some further results are shown in Chapter 7.4.

As Run 1 ended, the Higgs boson discovery was confirmed well above the 5σ threshold, and some of its key parameters were already measured with good accuracy. The LHC had provided an integrated luminosity of around 25 fb^{-1} per experiment combining centre-of-mass energies of 7 and 8 TeV and was turning to the preparation of Run 2. The upgrades brought to the machine during the shutdown allowed an increase in the luminosity and an increase in the centre-of-mass energy up to 13 TeV. As shown in Figure 1.18, increasing the energy from 8 TeV to 13 TeV

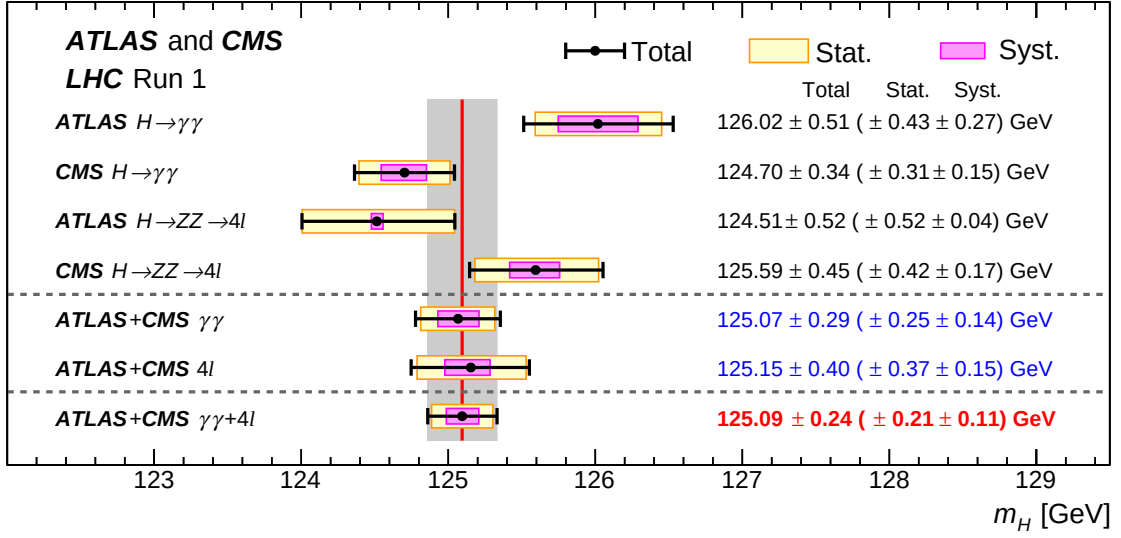
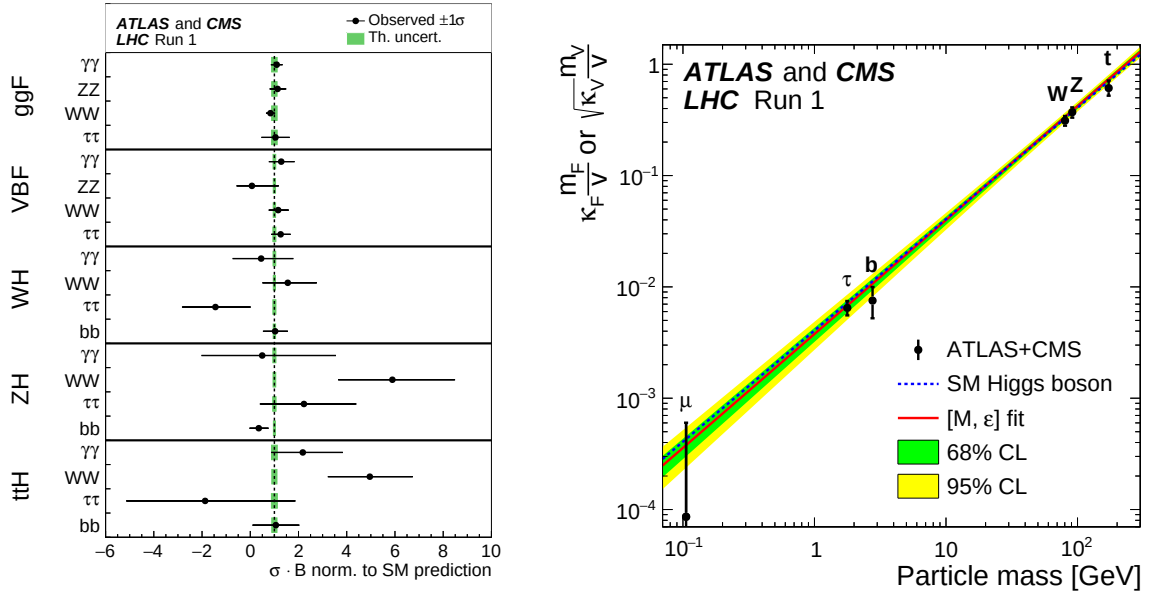


Figure 1.15 – Higgs mass measurements from the ATLAS and CMS experiments in the $H \rightarrow \gamma\gamma$ and $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ decay channels with Run 1 data, and combination of these measurements. [55]



(a) $\sigma_i \cdot \text{BR}^f$ for each processes $i \rightarrow H \rightarrow f$, normalised to their SM value.

(b) Higgs couplings to heavy fermions and weak bosons.

Figure 1.16 – Best fit values for (a) signal strengths and (b) couplings combining ATLAS and CMS measurements in the $H \rightarrow \gamma\gamma$, $H \rightarrow ZZ$, $H \rightarrow WW$, $H \rightarrow b\bar{b}$, $H \rightarrow \tau^+\tau^-$ decay channels using the Run 1 dataset. The various production modes ggF, VBF, WH, ZH, and ttH will be detailed in Chapter 7. [65]

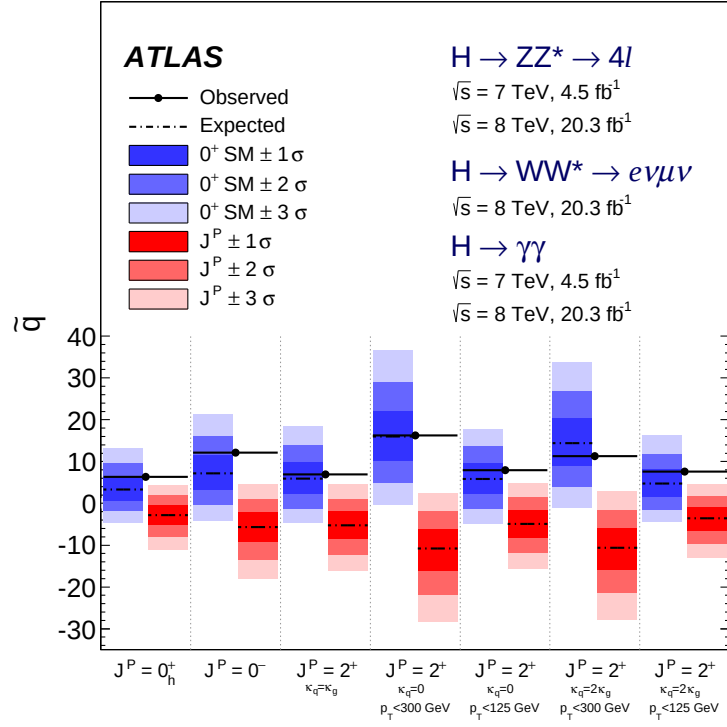
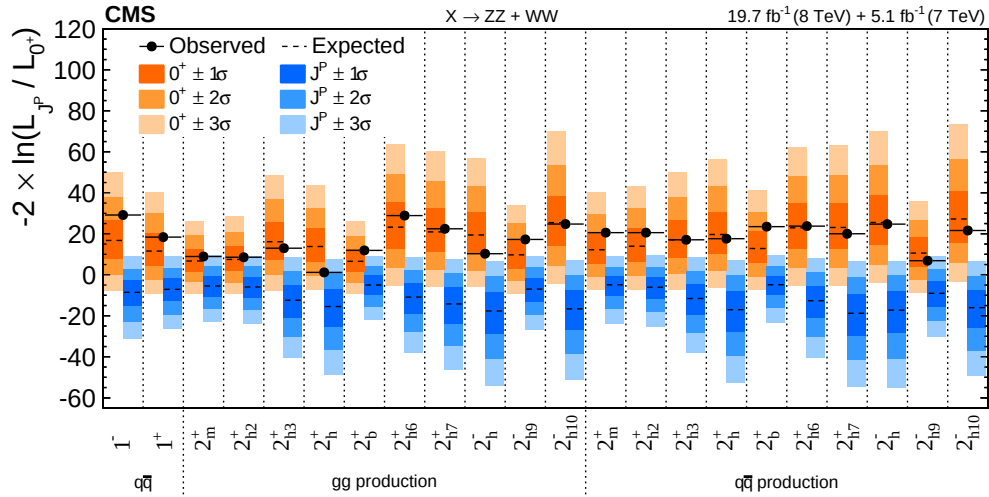

 (a) ATLAS result ($H \rightarrow \gamma\gamma$, $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$, $H \rightarrow WW$).

 (b) CMS result ($H \rightarrow ZZ^{(*)} \rightarrow 4\ell$, $H \rightarrow WW$).

Figure 1.17 – Alternative pure J^P states tested against the SM 0^+ hypothesis with Run 1 data in (a) ATLAS [68] and (b) CMS [70]. q is the test-statistics defined as $\ln(L_{J^P}/L_{0^+})$.

enhances the Higgs production cross-section by a factor of 2.3. Combined with five times more integrated luminosity, the Higgs statistics increased by a factor of ten between Run 1 and Run 2, allowing for even more precise measurements. The full Run 2 dataset is used for the analysis presented in Chapters 8 and 9.

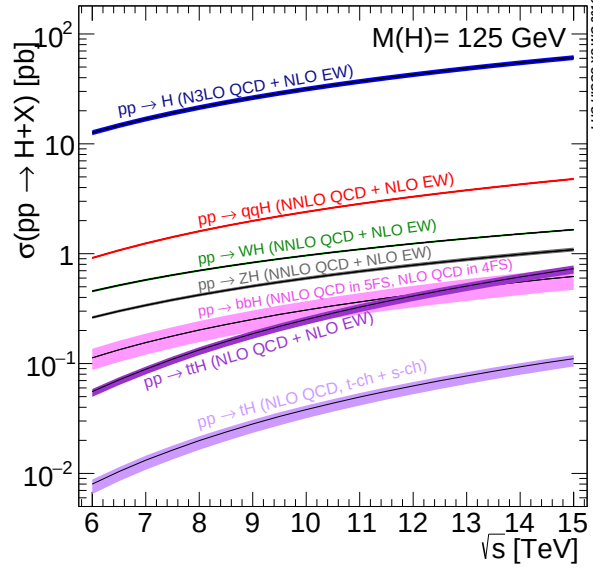


Figure 1.18 – Higgs production cross-section through different modes (see Chapter 7.1.1) as a function of the centre-of-mass energy. [72]

During this Thesis, numerous results using a partial Run 2 dataset were released by the ATLAS [73] and CMS [74] collaborations targeting various Higgs measurements such as cross-sections, branching ratios, couplings or mass measurements. Given the large number of results, an extensive summary is far out of scope of this Thesis and only couplings results will be discussed in Chapter 7.4. It is nevertheless worth mentioning that the mass measurement was also updated with the dataset collected at $\sqrt{s} = 13$ TeV until 2016 by both the ATLAS [75] and CMS [76] experiments:

$$m_H = 124.97 \pm 0.24 \text{ GeV} \quad (\text{ATLAS, Run 1 and 2015+2016, } H \rightarrow \gamma\gamma \text{ and } H \rightarrow 4\ell) \quad (1.84)$$

$$m_H = 125.26 \pm 0.21 \text{ GeV} \quad (\text{CMS, 2016, } H \rightarrow 4\ell). \quad (1.85)$$

The Large Hadron Collider

Contents

2.1	Why colliding particles	43
2.2	Motivations for a multi-TeV proton collider at CERN	44
2.3	A hadron collider in the LEP tunnel	45
2.4	The CERN accelerator complex	46
2.5	The Large Hadron Collider	46
2.6	The LHC experiments	49
2.7	LHC filling scheme	50
2.8	Luminosity and number of events	51
2.9	Parton distribution functions at hadron colliders	52
2.10	LHC performance	54

2.1 Why colliding particles

Most of the particles described by the Standard Model are unstable: the only known stable free particles are photons, electrons, neutrinos, and protons. Due to the confinement property of the QCD, free quarks cannot be observed and instead form bound states called mesons and baryons. Mesons (such as the pions, kaons, *etc.*) are composed of a quark-antiquark pair and are not stable. Baryons (protons, neutrons, *etc.*), composed of three quarks, exhibit only one stable compound: the proton.

The fact that no other particle is stable can be understood as follows: provided a long enough amount of time, a given particle will undergo all physically possible and kinematically allowed interactions (*i.e.* interactions authorised by the Standard Model described in Chapter 1.1, and which conserves the energy, the electric charge, baryon number, *etc.*). Due to energy conservation, the accessible phase-space only include states of lower energy (or equivalently of lower mass). Consequently, at some point, if a decay interaction is allowed, it will occur and lighter mass states will be formed.

The neutrinos, electrons and protons are therefore stable since no interaction allows them to reach a lower energy state. On the contrary, neutrons, for instance, decay through the allowed interaction $n \rightarrow p^+ + e^- + \bar{\nu}$ since $m_n > m_p + m_e + m_\nu$ (939.6 MeV, 938.3 MeV, 0.5 MeV and ~ 0 respectively), and muons decay to electrons through $\mu \rightarrow \nu_\mu + e + \bar{\nu}_e$ since $m_\mu > m_e + m_\nu + m_\nu$. The lightest mesons ($\pi^\pm$ and π^0) are allowed to decay into $\mu + \nu_\mu$ or $\gamma + \gamma$ since they are made of

a colour neutral $q\bar{q}$ pairs (with a null baryon number), and since the muon mass ($m_\mu \approx 106$ MeV) is lower than the pion mass ($m_\pi \approx 135$ GeV). The proton is also colour neutral but made of three quarks so has a baryon number of one, preventing it to decay to lighter mesons. Similarly for bosons, only the photon is stable: as it is massless its decay phase-space vanishes (the also massless gluons cannot be observed for the same reason as quarks: QCD confinement), while the massive Z , W and H bosons decay.

In order to detect an unstable particle, it must therefore firstly be produced, by putting enough energy in a system to create it. This is achieved at particle colliders where two beams of (stable) particles are accelerated and collided head-on, liberating enough energy to produce such unstable particles. The schematic rationale in particle physics is therefore the following: if a particle has not yet been discovered, then either the energy required to produce it is higher than what is currently available (the particle is heavier than the available energy in the centre of mass), or the interactions of this particle with other particles are too rare to have happened a significant amount of time in the current experiments (or both effects at the same time). The first problem is solved by increasing the beam energy in accelerator, thus increasing the energy available in the centre-of-mass of the collision, and the second problem is solved by increasing the *luminosity* of the collisions, *i.e.* the number of collisions produced (Section 2.8). Consequently, the history of experimental particle physics discoveries is closely linked to the progress and advances in particle accelerators.

2.2 Motivations for a multi-TeV proton collider at CERN

A commonly admitted birth date for the LHC project dates back to 1984, at a ECFA-CERN workshop [77]. At that time, the UA1 and UA2 experiments using collisions from the $Spp\bar{p}S$ collider at CERN had just discovered the Z and W boson [78–81], further completing the knowledge of the Standard Model. However, two pieces remained missing: the top quark and the Higgs boson. As these were not discovered at the $Spp\bar{p}S$, achieving a maximum centre-of-mass energy $\sqrt{s} = 900$ GeV for an instantaneous luminosity $\mathcal{L} = 6 \times 10^{30} \text{ cm}^{-2} \text{ s}^{-1}$, a more powerful and luminous colliders was required.

The year 1984 also marks two other milestones for particle accelerators. First, the start of the Tevatron [82] at Fermilab, a high energy proton-antiproton collider providing collisions at $\sqrt{s} = 1.96$ TeV with an instantaneous luminosity of $210 \times 10^{30} \text{ cm}^{-2} \text{ s}^{-1}$ [41], which eventually led to the discovery of the top quark in 1995 by the CDF and $D\bar{D}$ experiments [83, 84]. Second, the start of the civil engineering work for the Large Electron-Positron collider (LEP) at CERN. The LEP collider [85] eventually provided e^+e^- collisions at a centre-of-mass energy up to 209 GeV for an instantaneous luminosity up to $100 \times 10^{30} \text{ cm}^{-2} \text{ s}^{-1}$. However, electron-positron colliders, despite providing a well-controlled collision environment, have a lower discovery potential due to their lower centre-of-mass energy (see Section 2.3).

As a consequence, the possibility was already being studied to build a high-energy proton-proton collider in the same LEP tunnel [86], which would finally become the LHC. The energy and luminosity design of the machine were driven by the will to find the Higgs boson which could hide up to the TeV scale (see Chapter 1.2.1) or new physics going well beyond the TeV scale. Its design [87] reached a centre-of-mass energy of 14 TeV and a luminosity of $2 \times 10^{34} \text{ cm}^{-2} \text{ s}^{-1}$, surpassing by two orders of magnitude the design of the running accelerators at the time.

Exceeding by far the capabilities of the $Spp\bar{p}S$, both the Tevatron and LEP achieved major results in particle physics but fell short of proving the existence of the Higgs boson, which required the LHC to be confirmed. The discovery potential of this machine was realised in 2012 with the discovery of the long-sought Higgs boson, the accelerator achieving that year a record centre-of-mass energy of 8 TeV and instantaneous luminosity of $7.6 \times 10^{33} \text{ cm}^{-2} \text{ s}^{-1}$.

2.3 A hadron collider in the LEP tunnel

Given that $E = \sqrt{p^2 + m^2}$, producing high energy collisions can be achieved either by colliding heavy particles, or by colliding light particles but with a high momentum. Since the heaviest stable particle is the proton (~ 1 GeV), producing beam energies of 7 TeV can only be achieved using the second method, *i.e.* by the acceleration of “light” particles. Particle acceleration is primarily performed by applying an electric field in which the charged particle drifts. While the first accelerators used electrodes to create static electric fields, colliders rather rely on radio-frequency cavities (RF) which also improve beam dynamics: the alternative electric fields used allow for an easy bunching of particles for instance.

Two different concepts of accelerators coexists and are complementary: linear and circular accelerators. Linear accelerators, in their simplest design, consist in a succession of accelerating cavities, with one interaction point in the middle. The drawback of linear colliders however is that, for a given acceleration power of the cavities, reaching a higher beam energy requires more accelerating sections, increasing the cost. In a circular collider on the contrary, as the beam passes multiple times through the same point, far fewer cavities are needed. The drawback is obviously that the beams must be bent in circle, requiring powerful magnets to deflect the particle beams.

Due to this bending, particles are constantly accelerated which induces synchrotron radiation. In a ring of curvature radius ρ , the energy loss by synchrotron radiation per turn is given by [41]:

$$\Delta E = \frac{1}{3} \frac{e^2 \beta^3 \gamma^4}{\rho}, \quad (2.1)$$

where e is the charge of the particles in the beam, β their velocity, and γ the Lorentz boost factor which is proportional to the beam energy. In the final years of the LEP accelerator, such synchrotron radiation losses amounted to 2% per turn for a ~ 100 GeV electron beam. It is therefore hardly conceivable to prepare a higher energy electron beam using the LEP tunnel.

In order to achieve higher energies, two options can be considered: the first solution, assuming an e^+e^- machine, is to increase the curvature radius of the beam, *i.e.* build a larger tunnel. However this implies costly engineering work, and there would be considerable cost reduction if the LHC could instead fit inside the LEP tunnel. The other solution is to use more massive particles: since the Lorentz boost $\gamma = E/m$, the synchrotron emission also scales as m^{-4} . As a consequence, the LHC uses protons instead of electrons, which reduces the energy loss by a factor $(m_p/m_e)^4 \approx 10^{13}$.

Moreover, the constraint of keeping the same tunnel also fixes the maximal energy reachable at the LHC: for a particle of momentum p and charge q , the curvature radius r of the particle in a magnetic field B is given by $r = \frac{p}{qB}$, which translates to:

$$p = qBr. \quad (2.2)$$

Applying this to the LHC case, q is fixed (accelerating protons), r is fixed (LEP tunnel), so the highest possible momentum for the protons is given by the highest magnetic field achievable. At the time of the LHC design, electromagnets such as the one used for the LEP could typically produce fields of the order of 0.2 T [41]. This would result in approximately 200 GeV proton beams, much lower than what is required to reach multi-TeV collisions. However given the state-of-the-art of superconducting magnet development at that time¹, it was forecast that fields as high as 8 T could be achieved. Such magnets were indeed ready for the construction of the LHC, leading to a maximal energy of around 7 TeV per beam.

¹↑ The Tevatron and the HERA colliders were already using superconducting magnets with magnetic fields up to 4 T. However, the LHC magnets need to be cooled down at a lower temperature (1.9 K) to reach the nominal 8.3 T magnetic field, and use super-fluid helium over longer distances, adding to the technological endeavour.

Circular colliders are also usually able to deliver a higher luminosity: in a linear collider, bunches are accelerated, collide, and the remaining non-interacting particles of the bunches are lost. The accelerator must then be refilled, repeating the process as frequently as possible. In a circular collider however, the machine is filled once and the bunches can circulate for hours, providing collisions at each turn.

Finally, linear colliders only provide one interaction point. Circular colliders on the contrary can have several, allowing for multiple experiments to take data at the same time. The LHC for example can provide up to eight interaction points, four of them currently hosting an experiment (see Section 2.6).

2.4 The CERN accelerator complex

CERN since its foundation in 1954 has a long tradition of building long-lasting particle accelerators for fundamental subatomic physics: the first accelerator built at CERN, the Synchro-Cyclotron, started operating in 1957 and stopped in 1990, after 33 years providing 600 MeV beams to experiments. CERN's first collider was the 300 m-long Intersecting Storage Rings (ISR), colliding protons at a centre-of-mass energy up to 62 GeV between 1971 and 1984.

Most of the other CERN accelerators are still in use, and form the CERN accelerator complex as sketched in Figure 2.1. A fraction of their time is used to provide beams to the LHC, serving as an accelerator chain increasing the proton beam energy step-by-step [88]. Yet most of their time is spent providing hadron beams to the various CERN facilities and fixed target experiments.

The journey of a proton colliding in the LHC starts from a gaseous neutral di-hydrogen bottle. After applying an intense electric field, the protons are stripped from their electron and the acceleration begins. The first acceleration stage is performed by the LINAC2 (taking over from LINAC1 in 1978): protons are pushed up to 50 MeV before entering the Proton Synchrotron Booster (PSB, first beam in 1972). This 157 m-long ring accelerates the protons until they reach an energy of 1.4 GeV. They are then sent to the 628 m Proton Synchrotron (PS, started in 1959), accelerating the beam to 25 GeV. The beams are then kicked into the 6.9 km-long Super Proton Synchrotron (SPS, operating since 1976), which prepares 450 GeV beams ready to be injected to the LHC.

Protons finally enter the 26.7 km long LHC rings (Section 2.5), in which they are boosted from 450 GeV up to 7 TeV (nominally, 3.5 to 4 TeV during Run 1, 6.5 TeV during Run 2) in about 20 minutes. In practice though, an LHC beam is not a continuous stream of protons: the beam is structured in regularly separated bunches of $\sim 10^{11}$ protons, and depends on the capabilities of the injectors, which is detailed in Section 2.7. Given the lower particle storage capacity of the previous injector at each stage, it takes several fills of the PS to fill the SPS, and several fills of the SPS to fill the LHC. The full setup of the beams, from empty beams to 6.5 TeV stable beams, take about 2h.

2.5 The Large Hadron Collider

The main characteristics of the LHC are recalled below. More detailed explanations can be found in Refs. [88, 90, 91] and [92].

The layout of the LHC [92] is shown in Figure 2.2. It is not a perfect circle but is instead composed of eight curved sections (called *arcs*) and eight 528 m long straight sections (called *insertions*). The insertions host the experiments (see Section 2.6), the beam dump, or various beam cleaning systems. One insertion also hosts the RF system: each beam is accelerated by

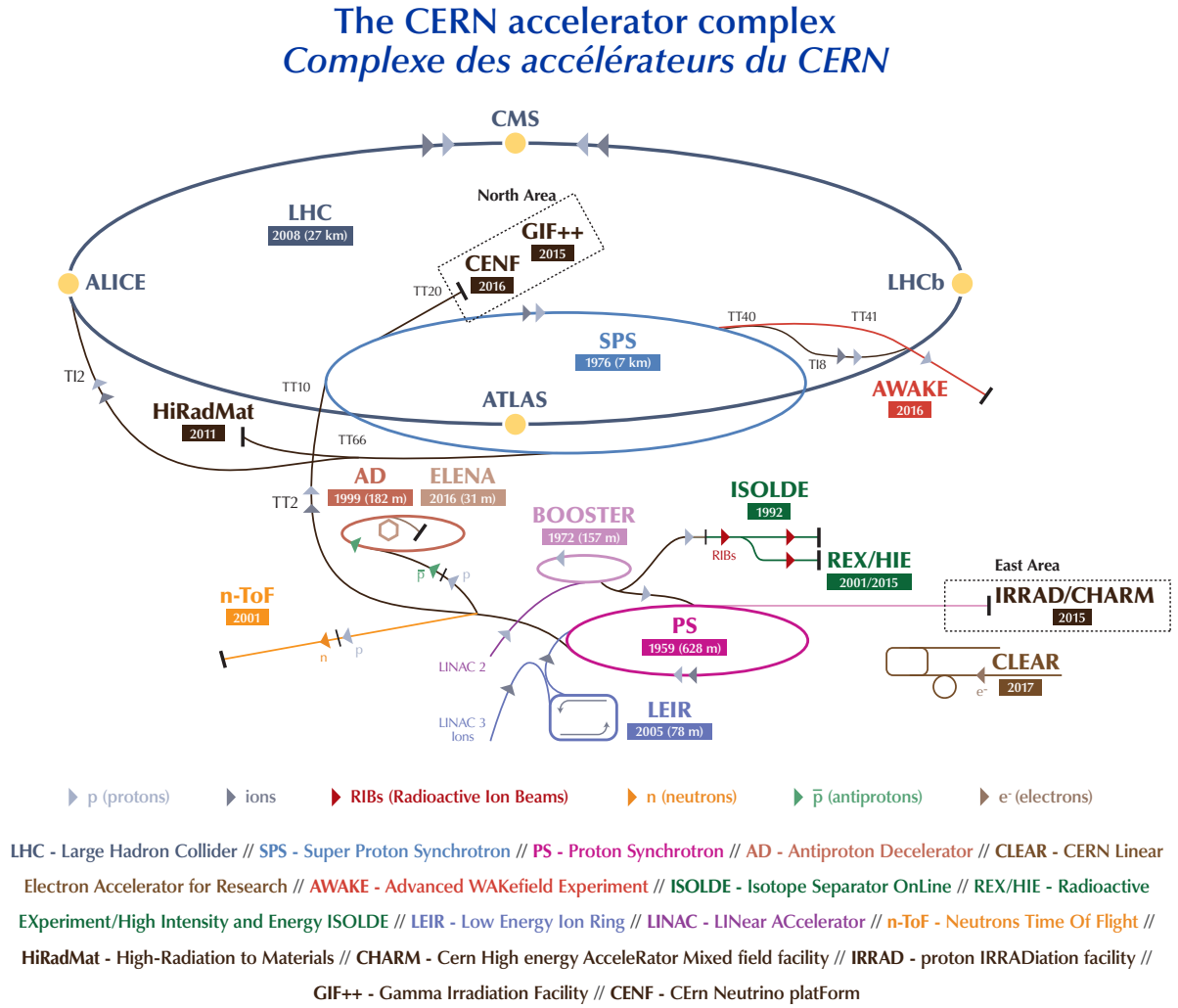


Figure 2.1 – The CERN accelerator complex and the LHC injection chain as of summer 2018. [89]

eight RF cavities operated at 400.8 MHz and 2 MV. The integrated acceleration field amount to 5 MV/m.

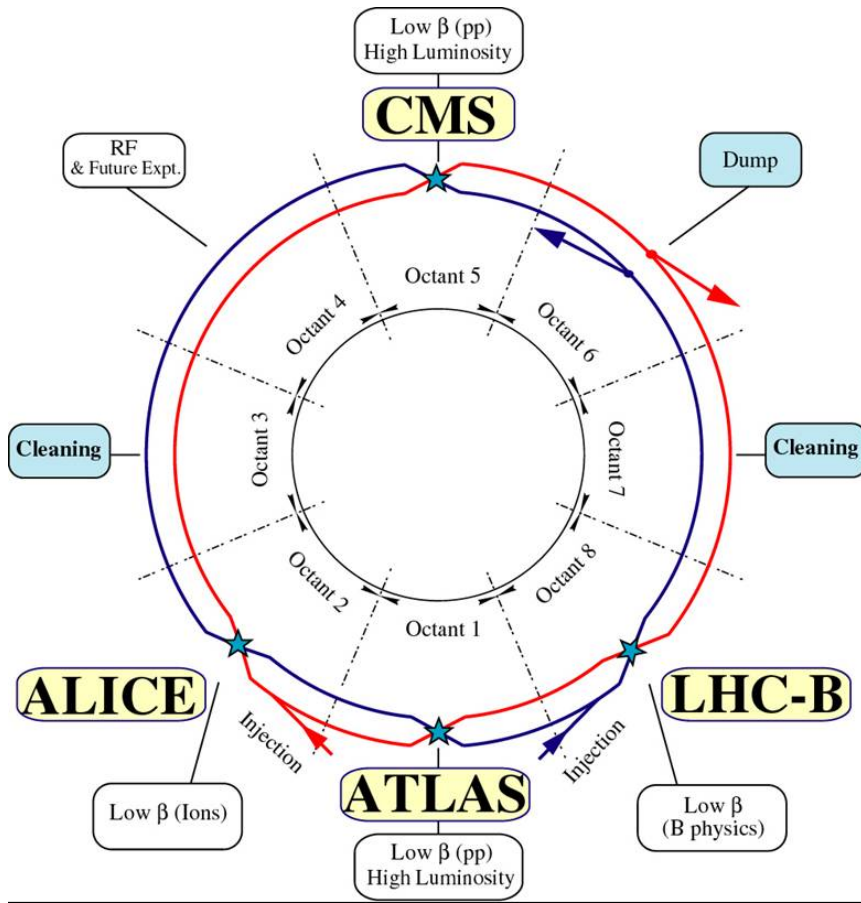


Figure 2.2 – Schematic layout of the main LHC ring. [93]

The arcs host most of the 10 000 magnets composing the LHC rings. The main magnets are the 1232 dipoles that are used to bend the beam. The other are higher order magnets (quadrupoles, sextupoles, *etc.* [94]) whose goal is to focus and correct various defects of the beams. On the contrary to Tevatron, the LHC collides two beams of protons² and not protons against antiprotons: while protons are abundant, antiprotons must be produced by a secondary accelerator and stored for an extended period of time to prepare the beam. Since from the Tevatron experience this is known not to be a trivial task, and given the number of particles needed in the beams to achieve the nominal LHC luminosity, a sufficient number of antiprotons could not be reliably produced so it was decided to collide two beams of protons. Furthermore, the LHC was also designed as a heavy ion collider, with the possibility to collide different types of ions (for example lead–lead or lead–proton collision). The species in the two beams are therefore both positively charged, but their respective charge (and mass) can differ. As a consequence, the magnetic field deflecting each of the beams is opposite, and two beam pipes are needed. Hence, LHC magnets actually combine two magnets, with two apertures for the beams. A schematic of a main LHC dipole is presented in Figure 2.3.

As previously mentioned, to reach the 8.3 T required magnetic field, the magnets need to be in a superconducting state. The design uses a Niobium-Titanium (NbTi) alloy that must be cooled to 1.9 K using superfluid helium to achieve the required current of 11 850 A. The beam pipes also need an advanced vacuum (about 10^{-10} mbar) to prevent beam interaction outside the desired

²↑ The LHC is mainly a proton–proton collider but also, also provides proton–lead and lead–lead collisions as its standard physics programme. These are the primary focus of the heavy-ion programme but are not used in the analyses described in this Thesis and will therefore not be detailed here.

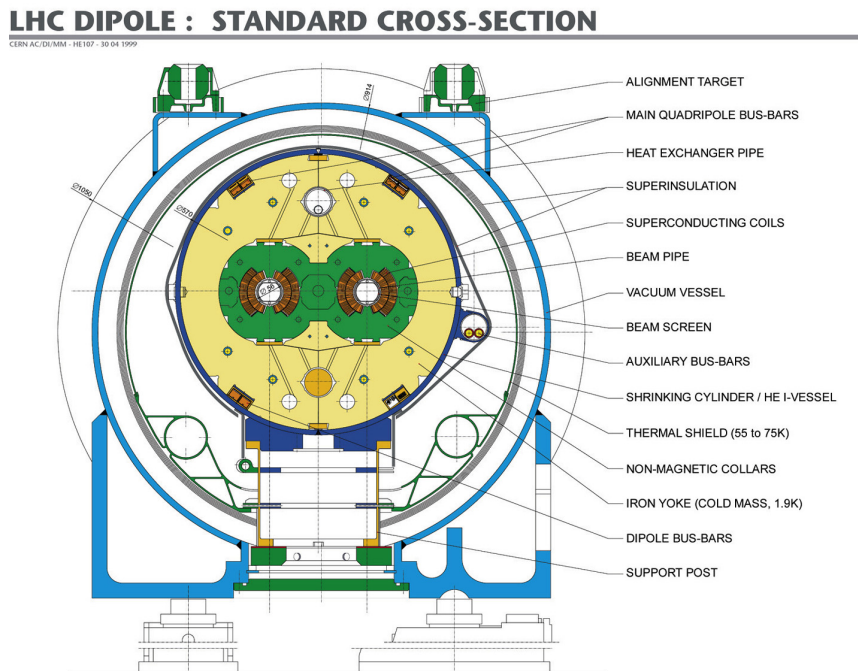


Figure 2.3 – Cross-section of a main LHC dipole showing the beam pipes, the coils, and all the cooling apparatuses. Each dipole is about 15 m long and weights approximately 35 t. [95]

interaction regions. Combining such extreme vacuum and cryogenics technologies into a 27 km long accelerator dissipating ~ 0.4 W/m of synchrotron radiation has been one of the technological endeavours in the construction of the LHC [96].

2.6 The LHC experiments

The LHC is host to four main experiments: ATLAS (A Toroidal LHC ApparatuS [97]) and CMS (Compact Muon Solenoid [98]) are located at the high-luminosity interaction-points of the LHC and benefit from large statistics to unveil rare processes. They nearly cover the full solid angle and can therefore cover a large panel of physics analyses ranging from Standard Model, Higgs boson or top quark precision measurements to searches for Super-Symmetric particles or exotic states.

The LHCb (LHC beauty [99]) experiment focuses on “*b*-physics”, *i.e.* the study of mesons and baryons containing at least one *b*-quark, to test the fundamental symmetries of the Standard Model or their violation (*CP*, lepton number, *etc.*). Given the relatively low mass of *b*-hadrons (5 GeV to 10 GeV) compared to the energy of LHC collisions (13 TeV), their production will mostly be boosted to the forward region. The LHCb detector therefore is optimised for this region, giving its characteristic shape of a rotated pyramid extending only on one side of the interaction point.

The ALICE (A Large Ion Collider Experiment [100]) detector is dedicated to the study of the quark-gluon plasma, a state of matter that is likely to have occurred in the most early instants of the Universe. LHCb and ALICE cannot withstand the flux delivered in the ATLAS and CMS interaction points and have a reduced delivered luminosity, also meaning less pileup.

Three smaller experiments were installed nearby the main ones, with specific physics targets. Both TOTEM (TOTal, Elastic, and diffractive cross-sections Measurements [101]), installed in the forward region of the CMS detector, and LHCf (LHC forward experiment [102]), installed 140 m from the interaction point on both sides of ATLAS, study protons which only scatter by a small deviation during the bunch crossing. These experiments study the cross-section of processes involving small deviations during the collisions, or emulate cosmic rays to interpret and calibrate

large-scale cosmic-ray experiments, respectively. The MoEDAL (Monopole and Exotics Detector At the LHC [103]) experiment was installed more recently, in 2015, near LHCb and is dedicated to the search of magnetic monopoles that could be created during collisions at the LHC.

2.7 LHC filling scheme

The filling scheme of an accelerator ring is determined by the capabilities of its RF system and its injector chain. At the LHC, the RF is operated at ~ 400.8 MHz [93], so the smallest interval that can contain particles is 2.495 ns (approximated to 2.5 ns in the following for sake of simplicity) or ~ 75 cm long. Given the length of the LHC (≈ 26.7 km), this amounts to 35640 possible slots (called *buckets*). Due to limitations of the injectors, the nominal bunch spacing is fixed to 25 ns, so only one every ten bucket can be filled, which defines a bunch: 3564 *bunch positions* are therefore available in the ring. The time measurement at the LHC thus often takes a bunch as basic unit, *i.e.* 25 ns. Collisions happen when two bunches cross each other, at the four experimental interaction points, and each event is identified by a Bunch Crossing Identifier (BCID). Out of these 3564, only 2808 are nominally filled due to dead-times from various systems, as summarised in Figure 2.4: some time must be reserved for the beam ejection kicker ramping (“dump”: ~ 3 μ s or 120 empty bunch positions), and for the kicker rise time of the various accelerators in the chain (PS $\rightarrow$ SPS: ~ 225 ns or 9 empty bunches, and SPS $\rightarrow$ LHC: ~ 950 ns or 39 empty bunches).

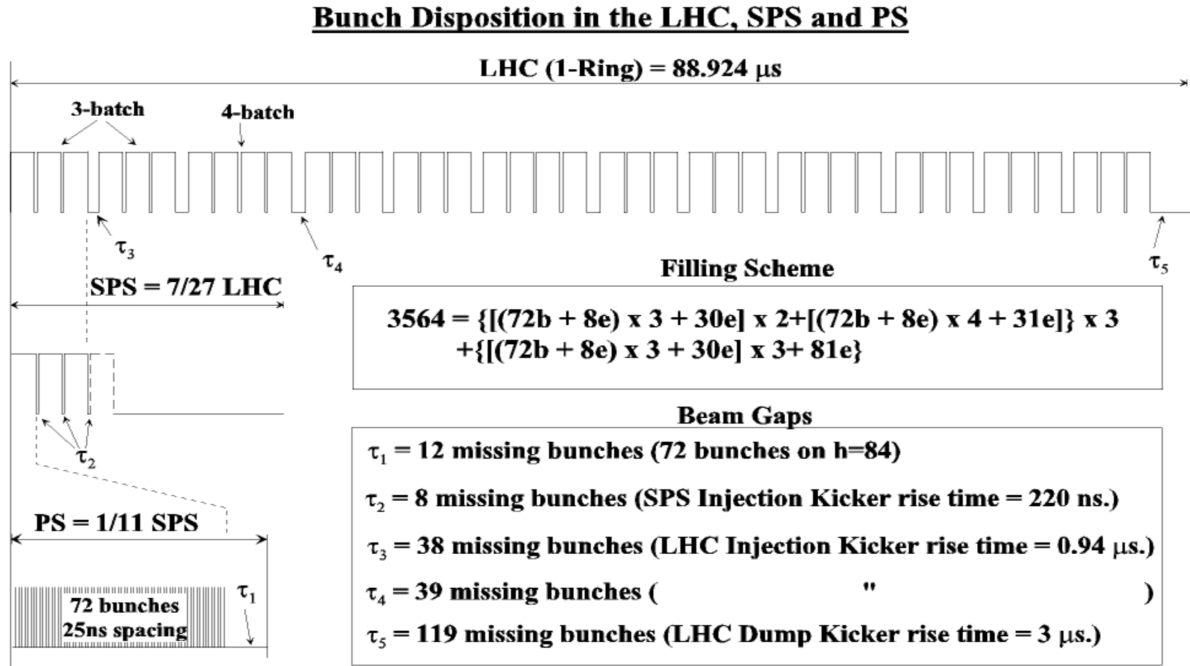


Figure 2.4 – Nominal beam structure at the LHC. The RF cavity allows for 3564 bunch slots of 25 ns, but due to the kicker rise time of the SPS and LHC not all of them can be filled. [93]

In practice, many filling schemes are used corresponding to various optimisations and tests for given LHC conditions. For example, at the start of a year, the LHC is not fully filled, slowly ramping up its total number of bunches in the beams. As an other example, in 2017 a dipole magnet suffered from higher heat load than expected, causing regular beam dumps with the nominal filling scheme. In order to maintain the LHC performances, the total number of bunches had to be reduced and another filling scheme was put in place. After several optimisations the luminosity was able to be kept at the same level as the earlier runs, at the price of harsher environment for the data-taking (higher pileup). A few actual filling schemes used during the

Run 2 are presented in Table 2.1

Table 2.1 – A few examples of the main accelerators filling scheme in Run 2. The batch separation corresponds to the PS $\rightarrow$ SPS kicker rise time, the train separation corresponds to the SPS $\rightarrow$ LHC kicker rise time. In 2017, an issue with a dipole magnet led to the development of an “unusual” filling scheme in order to preserve LHC performances (called “8b4e”).

Period	Colliding bunches	PS filling = 1 batch	SPS filling = 1 train		1 LHC ring	
		# bunch/batch in PS	# batch/train in SPS	batch spacing [$\times 25$ ns]	# train in LHC	train spacing [$\times 25$ ns]
2016	2208	48	2	9	23	≥ 36
2017a	2544	48	2 / 3	8	19	≥ 36
2017b	1866	$\underbrace{8b+4e+8b}_{\times 3}$	1/3/4	8	16	≥ 32
2018		Same as first 2017a periods				

2.8 Luminosity and number of events

As seen in Section 2.2, in order to discover rare processes, a large number of collisions is necessary. The number of events N for a process of cross-section σ relates to the collider luminosity through the following formula:

$$N = \sigma \cdot \epsilon \cdot L. \quad (2.3)$$

This formula exhibits three terms, each one having different dependencies:

- σ is the cross-section³ of the physical process under consideration. It also depend on the kinematics of the initial-state particles. For example the total cross-section for the production of a 125.09 GeV Higgs boson for proton–proton collisions at $\sqrt{s} = 13$ TeV is approximately $\sigma(pp \rightarrow H) \approx 55.6$ pb. Since most often one is interested in a specific decay channel, a more relevant quantity is the cross-section times branching ratio ($\sigma \cdot \text{BR}$). For example, the branching ratio of a 125.09 GeV Higgs decaying to four electrons is around $\text{BR}(H \rightarrow 4e) = 3.28 \times 10^{-5}$, and therefore the cross-section times branching ratio for producing a 125.09 GeV Higgs boson decaying to 4 electrons with proton–proton collisions at $\sqrt{s} = 13$ TeV is approximately $\sigma \cdot \text{BR}(H \rightarrow 4e) \approx 1.82$ pb.
- ϵ relates to detector: it includes effects from incomplete coverage of the phase-space (called acceptance and often denoted $\mathcal{A}$), and various intrinsic inefficiencies of the detector such as the trigger, the particle reconstruction and identification, *etc.*
- Finally the integrated luminosity $L = \int \mathcal{L} dt$, with $\mathcal{L}$ the instantaneous luminosity (defined below), only depends on the accelerator performance and availability.

The instantaneous luminosity $\mathcal{L}$ depends on the beam dynamics, focusing and energy through

$$\mathcal{L} = \frac{N_1 N_2 n_b f_{\text{rev}}}{4\pi \sigma_x \sigma_y} F = \frac{N_1 N_2 n_b f_{\text{rev}} \gamma}{4\pi \epsilon_n \beta^*} F, \quad (2.4)$$

where

- N_1 and N_2 are the numbers of particles per bunch (in this case, protons),

³↑ The cross-section unit corresponds to an area (it intuitively corresponds to the shared area of two cylinders colliding head-on), but values are often too large to be expressed in centimetres. The dedicated unit is the *barn* (and its subdivisions, in particular the pb and fb), which is defined as $1 \text{ b} = 10^{-24} \text{ cm}^2$.

- n_b is the number of bunch per beam,
- f_{rev} is the revolution frequency of the bunches (11.245 kHz at the LHC),
- σ_x and σ_y are the horizontal and vertical *effective* beam sizes at the interaction point (of the order of 15 μm), computed from the individual beam sizes by $\sigma_x^2 = \sigma_{x,1}^2 + \sigma_{x,2}^2$ and $\sigma_y^2 = \sigma_{y,1}^2 + \sigma_{y,2}^2$.
- γ is the Lorentz boost factor,
- ϵ_n is the normalised emittance, related to the beam emittance ϵ by⁴ $\epsilon_n = \gamma\epsilon$,
- β^* is the value of the β function (not the relativistic factor) at the interaction point, where the β function is defined as $\beta = \sigma_x\sigma_y/\epsilon$,
- and F is a geometrical factor due to the small but non-vanishing crossing angle at the interaction point:

$$F = \left(1 + \left(\frac{\theta_c}{2} \frac{\sigma_z}{\sqrt{\sigma_x^2 + \sigma_y^2}} \right)^2 \right)^{-1/2}, \quad (2.5)$$

with σ_z the longitudinal RMS of the bunches (of the order of 10 cm). Such non-zero collision angle is needed to avoid parasitic interactions occurring at other places than the desired interaction point.

Typical values for each year are shown below in Table 2.2.

Typically, an analysis measures a number of events N and wants to deduce the cross-section (times branching ratio) of its process of interest. A precise measurement therefore relies on a precise knowledge of both detector effects and integrated luminosity. For that purpose, the actual instantaneous luminosity is regularly measured in each experiment by dedicated instruments (see Chapter 3.6) and the detectors effects are part of dedicated analysis in the experiments (see Chapter 4 for some examples).

2.9 Parton distribution functions at hadron colliders

Contrary to electrons, protons are not elementary particles: they are composed of three *valence* quarks, and a *sea* of quarks, antiquarks and gluons. These are commonly referred to as *partons*. Consequently at the LHC (or any hadron collider), it is not the protons that actually collide, but one of their partons. Each of the partons carries a fraction of the total proton momentum, which is usually represented by the so-called *Parton Distribution Functions* (PDF). These PDF cannot be predicted by the QCD and must be fitted from experimental data. Inputs to the PDF determination mostly come from results of dedicated deep inelastic electron–proton scattering (DIS), such as the ones provided by the H1 and ZEUS experiments at the HERA collider in DESY (shut down in 2007). However, LHC data is now sufficiently precise to be included in the determination of the PDF. It should also be noted that the DIS results and the derived PDF depend on both momentum fraction x and the momentum transfer Q^2 between the electron and the parton. Several collaborations (for instance: HERA, CTEQ, MSTW, NNPDF, *etc.*) then provide PDFs for each constituent, extrapolated for any Q^2 . An example from the MSTW collaboration [104] is shown in Figure 2.5.

The compositeness of the colliding protons also has an impact on the cross-sections: the total production cross-section for a given process ($pp \rightarrow X$) is in fact the sum over all partonic

⁴↑ The emittance represents the volume of phase-space occupied by the beam (*i.e.* beam size and energy or momentum spread). The normalised emittance relates to the emittance by $\epsilon_n = \beta\gamma\epsilon$, where β and γ are the relativistic factors (here β is not the β function). At the LHC, $\beta \approx 1$, so one can consider $\epsilon_n = \gamma\epsilon$.

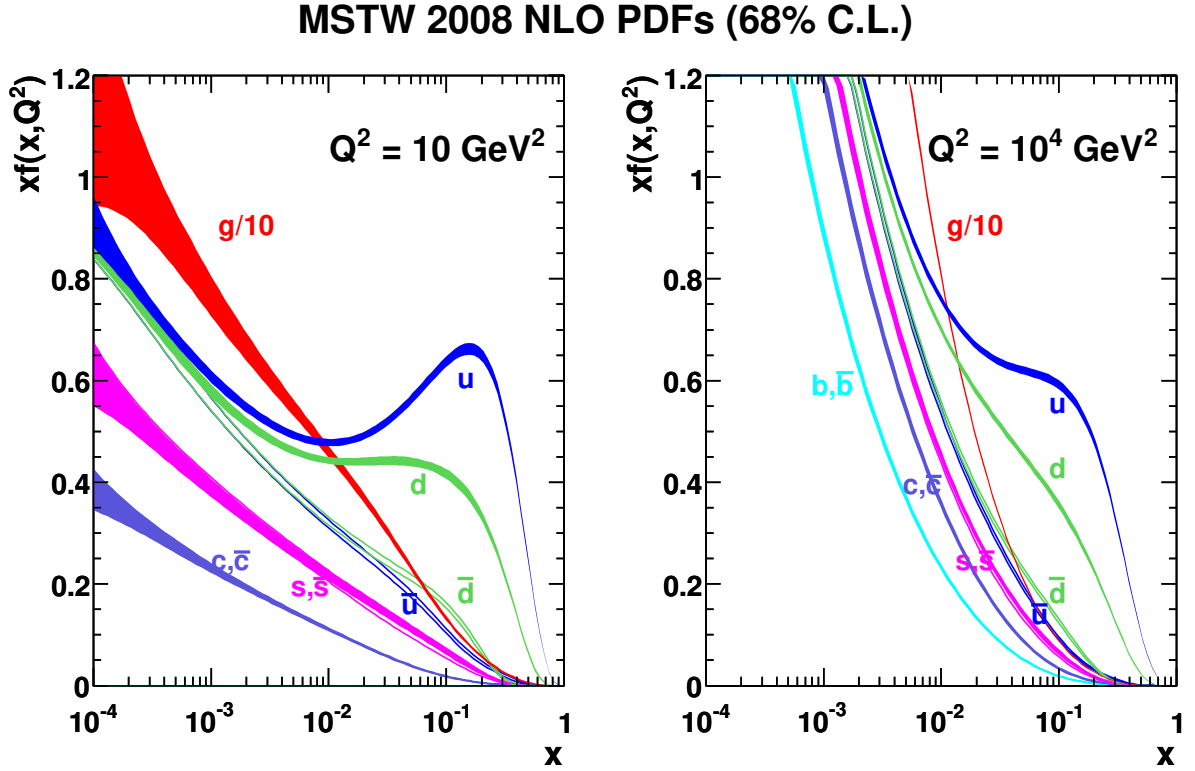


Figure 2.5 – Proton Parton Distribution Functions (PDF) as computed by the MSTW collaboration [104], on the left for a momentum transfer $Q^2 = 10 \text{ GeV}^2$, on the right for $Q^2 = (100 \text{ GeV})^2$. x is the Bjorken variable, representing the fraction of momentum carried by the parton. Since the valence quarks of the proton are 2 ups and one down, the PDF differs between u and $\bar{u}$, and between d and $\bar{d}$. Since the strange ($s, \bar{s}$), charm ($c, \bar{c}$) and bottom ($b, \bar{b}$) contributions come from sea quarks, the contribution is the same for each quark and antiquark. The gluon contribution is scaled by a factor $1/10$.

processes, for all possible momentum of the partons. This mathematically reads

$$\sigma_{pp \rightarrow X} = \sum_{a,b} \int_0^1 dx_1 \int_0^1 dx_2 f_a(x_1) f_b(x_2) \times \hat{\sigma}_{ab \rightarrow X}(x_1, x_2; \hat{s}, \hat{t}, \hat{u}), \quad (2.6)$$

where a and b run over all possible partons (gluon, u , d , s , c and b quarks and their antiquarks), x_1 and x_2 are the momentum fractions carried by the incoming partons, f_i is the PDF for the parton i , $\hat{s}$, $\hat{t}$ and $\hat{u}$ are the *Mandelstam variables* representing the momentum of the incoming and outgoing particles, and $\hat{\sigma}$ is the *partonic cross-section* for the process.

Consequently, while the proton–proton centre-of-mass energy $\sqrt{s}$ is constant for all collisions, the actual centre-of-mass energy, denoted $\sqrt{\hat{s}}$ and defined as the parton centre-of-mass energy, varies. Since it is impossible to know which parton eventually collides for a given event, the exact centre-of-mass energy for this event is unknown. However, this allows to probe multiple centre-of-mass energies with one single beam setup, whereas a lepton collider would have to change its beam settings to “scan” multiple centre-of-mass energies when performing a $\sqrt{s}$ resonance search.

2.10 LHC performance

During Run 2, the LHC has been running with conditions close to or above its nominal design, and with an availability close to 50 % (Figure 2.6). As a single example of its performance, the peak luminosity doubled the design value in 2017 and 2018, leading to more than 150 fb^{-1} delivered to the ATLAS and CMS experiments during the Run 2 (Figure 2.7). A key quantity for the experiment’s data taking is the number of additional interactions per bunch crossing (*pileup*), which increases with the instantaneous luminosity. While this value averaged ~ 20 during Run 1, peaking at ~ 35 , the good LHC performance achieved during Run 2 leads to an average pileup of ~ 34 , peaking above 65, as seen in Figure 2.8. Such conditions complicate the data taking, but experiments also managed to maintain a good data taking quality. The ATLAS and CMS experiments thus each provide a physics dataset of $\sim 140 \text{ fb}^{-1}$ for the full Run 2 analyses. A summary of the key parameters during the past Runs is shown in Table 2.2.

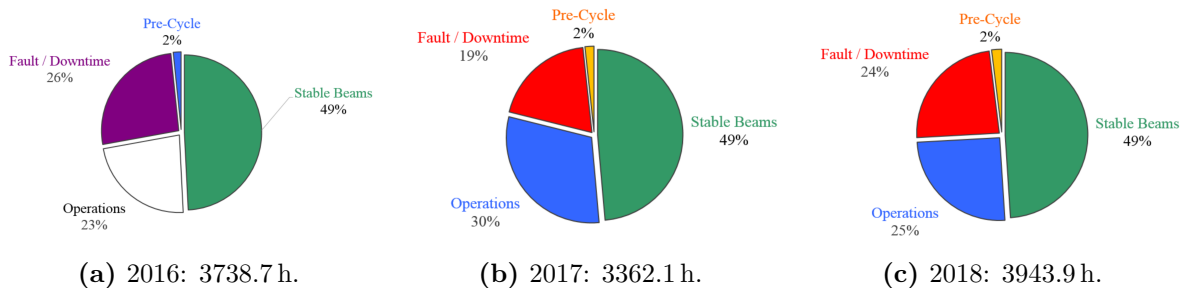


Figure 2.6 – Fraction of the LHC uptime spent in the various mode in (a) 2016 (total 3738.7 h) [105], (b) 2017 (total 3362.1 h) [106], and (c) 2018 (total 3943.9 h) [107]. The average *stable beams* (*i.e.* delivering physics collisions) fraction for the full Run 2 is almost 50 %.

After 2018, the LHC stopped for more than two years to prepare for the Run 3, that should occur between 2021 and 2023 with an increased centre-of-mass energy (likely 13.5 or 14 TeV), and conditions similar to the end of Run 2. A luminosity of about 300 fb^{-1} for ATLAS and CMS is expected at the end of 2023, doubling the dataset recorded during Run 2.

In 2025, the LHC will enter its ultimate phase, called High-Luminosity LHC (*HL-LHC*) which should last from 2025 to 2035. The total integrated dataset should reach 3000 fb^{-1} , thanks to an instantaneous luminosity reaching up to $5 \times 10^{34} \text{ cm}^{-2} \text{ s}^{-1}$. The price of such a high luminosity will however be the high pileup, with a value potentially as high as 200.

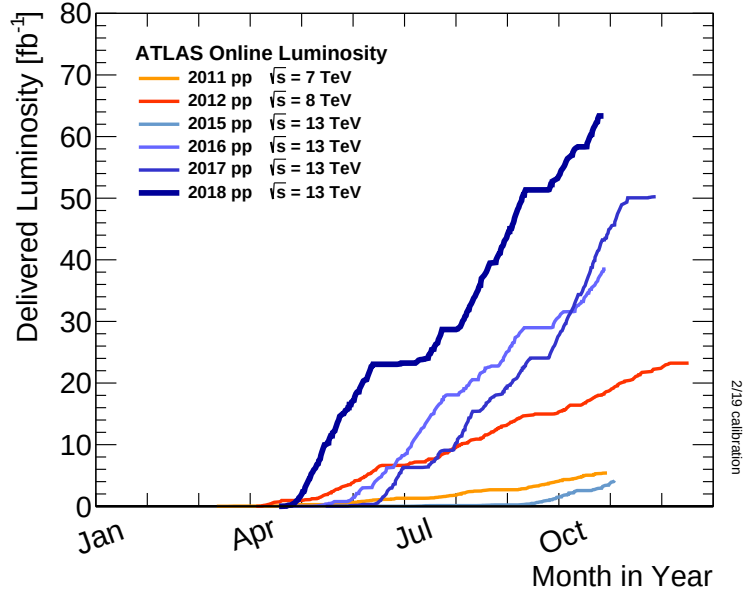
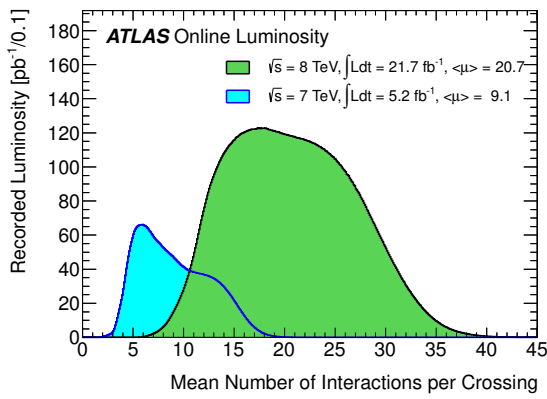
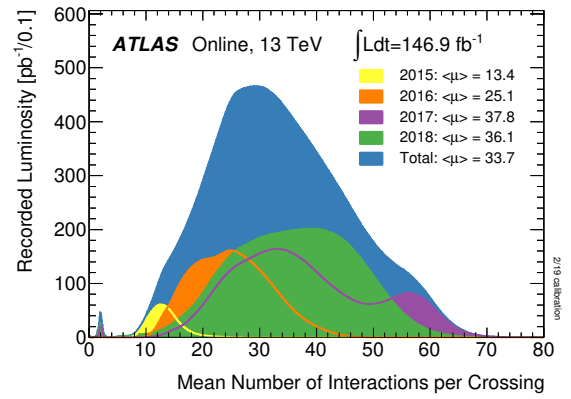


Figure 2.7 – Cumulative delivered luminosity measurement by the ATLAS experiment as a function of time for high-energy pp collision for each LHC running year. [108]



(a) 2011 – 2012 ($\sqrt{s} = 7$ and 8 TeV).



(b) 2015 – 2018 ($\sqrt{s} = 13$ TeV).

Figure 2.8 – Recorded luminosity-weighted distribution of the mean number of interactions per crossing in the ATLAS experiment for each LHC running year of (a) Run 1 and (b) Run 2. [108]

Table 2.2 – LHC configuration parameters during Run 1 and Run 2 (pp collision only) [109, 110]. The 2017 run suffered from a radical change of condition during the summer due to magnet issue. In order to preserve the delivered luminosity, the LHC was filled using a different scheme. The luminosity figures use online ATLAS measurements (initial calibration) [108, 111].

	Design	2011	2012	2015	2016	2017	2018
Beam energy	7	3.5	4	6.5			
Nominal bunch spacing	25	50		25			
Initial # proton per bunch	1.15×10^{11}	1.3	1.5	1.0 - 1.25			
Max. number of bunches	2808	1380	60	2244	2200	2556 / 1868	2556
β^* in ATLAS/CMS	55	100-150	60	80	40	40-30	30-25
Emitance in IP	3.75	~ 2.4	~ 2.4	~ 3.5	~ 2.2	~ 2.2	~ 1.9
Half crossing-angle	143	120	146	145	185 / 140	150-120	160-130
Number of days of high intensity pp operation				88	146	140	145
Machine availability	[%]			69	76	83	79
Stable beam efficiency	[%]			35	49	49	49
Peak instantaneous lumi	$[\times 10^{34} \text{ cm}^{-2} \text{ s}^{-1}]$	0.38	0.76	0.50	1.39	2.06	2.15
Average number of interactions per crossing (pileup)		9.1	20.7	13.5	24.9	37.8	36.1
ATLAS delivered lumi	$[\text{fb}^{-1}]$	5.46	22.8	4.2	38.5	50.2	63.3
ATLAS recorded lumi	$[\text{fb}^{-1}]$	5.08	21.3	3.9	35.6	46.9	60.6
ATLAS physics lumi	$[\text{fb}^{-1}]$	4.57	20.3	3.2	32.9	44.4	58.5

ATLAS: A Toroidal LHC ApparatuS

Contents

3.1	Requirements for the ATLAS detector	58
3.2	ATLAS coordinate systems	60
3.3	The inner detector and the solenoid magnet	61
3.3.1	The pixel detector	63
3.3.2	The Strip Semiconductor Tracker (SCT)	63
3.3.3	The Transition Radiation Tracker (TRT)	63
3.3.4	Central solenoid magnet	64
3.4	The electromagnetic and hadronic calorimeters	65
3.4.1	Principles of calorimetry	65
3.4.2	Application to the ATLAS electromagnetic calorimeter	67
3.4.3	Central and endcap electromagnetic calorimeter	68
3.4.4	Central hadronic calorimeter	76
3.4.5	The hadronic endcap and forward calorimeters	77
3.5	The muon spectrometer and the toroid magnets	80
3.6	Luminosity monitoring	82
3.7	Trigger system	85

In a collision at the LHC, hundreds to thousands of particles can be created at each bunch crossing. However, only particles that live long enough to reach the detector can be measured. Assuming particles are travelling close to the speed of light, and given the tracker size (~ 1 m), particles with a lifetime τ of the order the nanosecond or higher fly long enough to be detected. This includes electrons¹, protons, neutrons ($\tau \approx 900$ s), muons ($\tau \approx 2.2$ μ s), charged pions ($\tau \approx 25$ ns), and kaons ($\tau \sim 10$ ns to 50 ns). In the following, “stable particle” designates any of these *long-lived* particles. Neutrinos, despite being stable, interact too weakly to be detected, and will therefore manifest in the detector as missing energy. Other particles decay too quickly and therefore cannot be directly detected: they are instead deduced from their disintegration products, eventually after several such disintegration steps leading to the above mentioned particles. Moreover, due to QCD, stable and unstable hadrons are often not resolved individually, and mostly show up as groups of particles called *jets*.

¹↑ Since there are only negligible difference of particle and antiparticle interaction with matter, in this Chapter and most of this Thesis no distinction is made between particles and antiparticles so that an electron will denote an electron (negative electron) or a positron (positive electron), muon will denote both the negatively charged muon and the positively charged antimuon, *etc.*

An efficient detector must therefore be able to detect, identify and distinguish between these particles. The layout of a detector is thus designed to fulfil these objectives: first a tracker, to identify charged particles and measure their momenta, which can be complemented by a specialised particle identification system, then a calorimetry system to measure both charged and neutral particle energy by stopping, and finally a muon detector, outside the calorimeter since muons are weakly interacting in the calorimeter and can escape to the outer layers of the detector. To identify the charge and momentum of particles, the tracker must also be enclosed in a magnetic field and a magnet system, while not directly participating to the detection (it is not an “active” area), is needed to complete the characterisation of the detected particles.

3.1 Requirements for the ATLAS detector

The required performances of the detector are linked to the physics case it wants to achieve. As a general purpose detector, ATLAS physics goals are broad and include, for example:

- studies of the W and Z bosons, in particular in their leptonic decay channels, requiring a precise measurement of electrons and muons;
- studies of the top quark, which relies on the measurements of b -hadrons, needing a good vertex resolution for an efficient b -jet identification (b -tagging), and on the measurement of the associated W bosons decaying to leptons or jets;
- searches for unknown (*e.g.* supersymmetric particles) and exotic states, whose tagging often relies on displaced vertices or missing energy in the detector, and a possible high jet multiplicity;
- studies of the Higgs boson in several decay modes, such as in a photon pair, requiring a good energy and spatial resolution of photons, or in a τ -lepton pair subsequently decaying in jets or lighter leptons and neutrinos, demanding precise energy measurements of electrons, muons, and jets, or in a b -quark pair, requiring b -tagging capabilities with good jet energy resolution.

Focusing on the analysis developed in Chapters 8 and 9, namely the Higgs boson decaying to four leptons (electrons or muons), the detector needs to provide a performing energy resolution and identification of electrons, a good momentum resolution of muon, and additionally a good energy resolution of jets for the studies of production modes with associated jets (VBF, ttH , *etc.*). Moreover, as the detector operates in a busy hadronic environment due high luminosity proton collisions, it must also be able to distinguish the pileup from interesting signals. Achieving the desired performance in all these studies requires

- an efficient discrimination between electrons, photons and light hadrons (*e.g.* pions and kaons), so an efficient particle identification system;
- a precise energy resolution for electrons, photons and jets, which translates into a precise calorimetry system with a low upstream material budget;
- a high granularity calorimeter to achieve a precise photon spatial resolution and the discrimination between pileup jets and energy deposits from signal processes;
- a precise measurement of displaced vertices for b -tagging, so a tracker located as close as possible of the interaction point;
- a good muon and electron momentum resolution, so a precise tracker embedded in a high magnetic field;
- and an hermetic 4π coverage to infer the presence of undetected particles from the transverse momentum balance of in the collision.

The ATLAS detector [97] is designed combining these requirements at their best compromise (for example the tracking needs to be as close as possible to the beam with a strong magnetic field and provide as many points as possible, but additional material in front of the calorimeter degrades the calorimeter resolution). Its characteristics and main sub-detectors are summarised throughout this Chapter. An illustrative layout of the detector is shown in Figure 3.1. The full system results in a 44 m long, 25 m diameter, 7000 t detector, which is located at the Point 1 (P1) of the LHC.

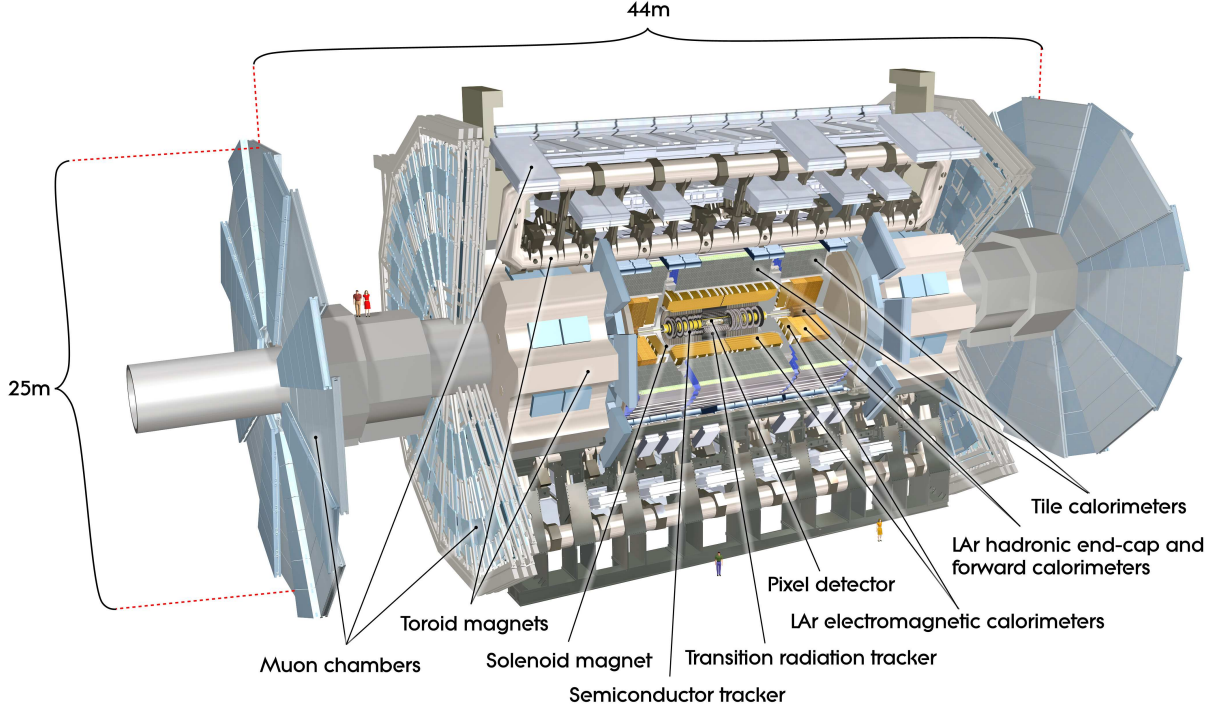


Figure 3.1 – General cut-away view of the ATLAS detector. [97]

The simultaneous use of all the sub-detectors (tracker, calorimeters, muon system) allows for a complete identification of all stable particles. Electrons and photons, while having a similar energy deposit in the calorimeter, can be distinguished by the presence (charged, so electron) or absence (neutral, so photon) of a track in the tracker (ID). The calorimeter system comprises two differently designed elements (ECAL and HCAL), allowing for an efficient distinction between electromagnetic showers from electrons and photons, and hadronic showers from hadrons. Since muons are not stopped by the calorimeters, only depositing a small amount of energy, they travel further and are uniquely identified by the presence of hits in the outer layers of the detector constituted by the muon system (MS).

Finally, neutrinos totally escape the detection, interacting too weakly with matter, but can be deduced from the other elements. Due to the partonic interaction (see Chapter 2.9), the momentum of the colliding partons system cannot be known. However, since the beams collide head-on² it can be assumed that the colliding system has vanishing transverse momentum. By conservation of energy and momentum, the vectorial sum of all outgoing particle energies must also be zero. Thus, the *missing transverse energy* (E_T^{miss}) required to balance the measured transverse energy is interpreted as undetected neutrino(s) flying in that direction.

²↑ The beams actually collide with a small crossing angle (see Chapter 2.8), but this angle ($\theta_c < 150 \mu\text{rad}$) is negligible.

3.2 ATLAS coordinate systems

The ATLAS reference frame is defined by a cartesian system of coordinates defined as follows:

- the origin of the frame is at the centre of the detector,
- the x axis is in the horizontal plane, pointing toward the centre of the LHC,
- the y axis is in the vertical plane, pointing upward,
- the z axis is defined along the beam pipe such that (x, y, z) forms a right-handed system of coordinates.

Due to the cylindrical geometry of the detector, a more convenient cylindrical system of axes (θ, ϕ, z) is defined³: $\phi \in [-\pi, \pi]$ is the azimuthal angle around the beam axis (*i.e.* in the x - y plane), with positive values corresponding to the upper half of the detector, and $\theta \in [0, \pi]$ is the polar angle around the vertical axis (*i.e.* in the x - z plane), with $\theta = 0$ when pointing along $+z$ and $\theta = \pi$ when pointing along $-z$. A more practical coordinate is to use the pseudo-rapidity η ($\in [-\infty, +\infty]$) instead of the polar angle θ which relate by:

$$\eta = -\ln \left(\tan \frac{\theta}{2} \right). \quad (3.1)$$

The pseudo-rapidity is equivalent to the rapidity⁴ for massless particles. A few useful examples of translation between η and θ are given in Table 3.1. The interest of using the pseudo-rapidity lies on the fact that rapidity intervals are invariant under a Lorentz boost. The positive η side (corresponding to positive z coordinate values) is often denoted *A-side*, the other half of the detector is denoted *C-side*.

Table 3.1 – A few examples of η - θ correspondence.

η	0	0.8	1.37	1.52	1.8	2.0	2.4	2.5	3.2	4.9
θ [°]	0	48.4	28.5	24.7	18.8	15.4	10.4	9.4	4.7	0.85
θ [°]	0	60	45	30	15	10	5	1	0.1	0.01
η	0	0.55	0.88	1.32	2.03	2.44	3.13	4.74	7.04	9.34

As mentioned earlier at the end of Section 3.1, since the momentum fraction of the incoming partons in the collision cannot be exactly known, the reliable measurements are based on transverse quantities: for an observable generically denoted $\mathcal{O}$, its transverse component $\mathcal{O}_T$ is defined as

$$\mathcal{O}_T = \sqrt{\mathcal{O}_x^2 + \mathcal{O}_y^2} = \mathcal{O} \sin \theta = \frac{\mathcal{O}}{\cosh \eta}. \quad (3.2)$$

Its longitudinal component can be recovered by $\mathcal{O}_z = \mathcal{O} \tanh \eta$.

Another useful quantity is the definition of the angular separation between two objects in the detector. This is denoted ΔR and defined as

$$\Delta R = \sqrt{(\Delta \eta)^2 + (\Delta \phi)^2}. \quad (3.3)$$

For massless particles, this quantity is also invariant under a Lorentz boost along the z -axis.

³↑ The usual (r, ϕ, z) cylindrical coordinates are recovered with $r = z \tan \theta$.

⁴↑ The rapidity of a particle is defined as $y = \frac{1}{2} \ln \left(\frac{E+p_z}{E-p_z} \right)$. Rapidity and pseudo-rapidity are equivalent for ultra-relativistic particles (for which $p \approx E$): $y = \frac{1}{2} \ln \left(\frac{E+p_z}{E-p_z} \right) \approx \frac{1}{2} \ln \left(\frac{p+p_z}{p-p_z} \right) = \frac{1}{2} \ln \left(\frac{1+\cos \theta}{1-\cos \theta} \right) = -\ln \left(\tan \frac{\theta}{2} \right) = \eta$.

3.3 The inner detector and the solenoid magnet

The purpose of the tracking system is to provide the means of reconstructing each track created during each collision at the LHC. Given the high number of tracks, a finely grained detector is needed, especially in the innermost layers of the tracker. The sensors are also subject to a high particle fluency during their lifetime and must therefore be radiation-hard by designed.

Tracker elements usually provide *hits*, *i.e.* binary information if a particle has passed through this particular element or not: the higher number of points along the trajectory, the better the resolution on the track parameters (position and curvature, leading to vertices and momentum). Tracking systems are therefore arranged in multiple layers, providing a discrete information on the trajectory. The trajectory is then extrapolated toward the beam pipe, providing information on the production vertices: the *primary vertex* that emerges from the actual collision location, and potential *secondary vertices*, displaced by a few millimetres that are the signature of medium-lived particle decays (such as *b*- and *c*- baryons or mesons). Each individual hit information must also be precisely localised, hence requiring small dimension sensor elements (high *pitch*).

The main ATLAS tracking tool is the inner detector (ID, Figures 3.2 and 3.3) which constitutes the innermost part of the ATLAS detector by occupying the radii below 1150 mm and extending over 3512 mm on each side along the *z*-axis. It is composed of three sub-detectors consisting of silicon sensors made of pixels (Section 3.3.1) or strips (Section 3.3.2), and of a gaseous detector allowing for advanced particle identification (Section 3.3.3).

The tracker is also used to provide a precise measurement of the particle momentum, which can be deduced from the curvature of a charged particle track in a magnetic field. The requirements for the ATLAS tracking system is to be able to reconstruct all tracks with $p_T \geq 0.5$ GeV and within $|\eta| < 2.5$. The whole system is therefore enclosed in an intense magnetic field produced by a solenoid magnet (Section 3.3.4).

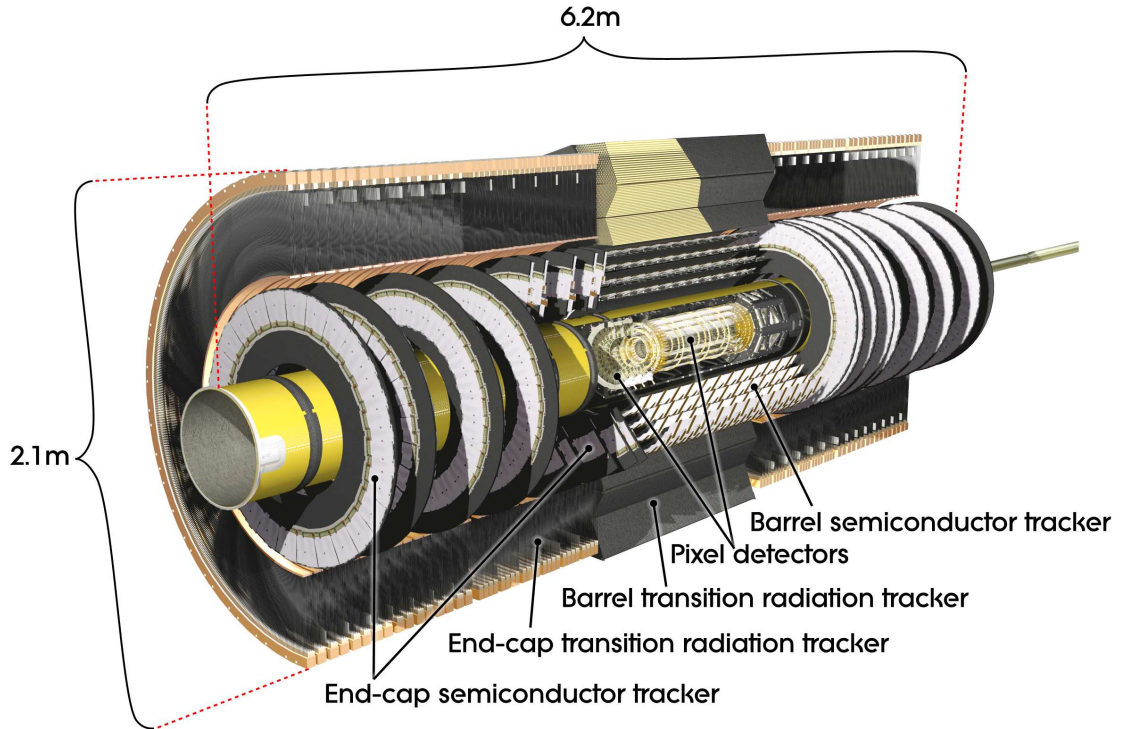


Figure 3.2 – Schematic of the ATLAS Inner Detector. [97]

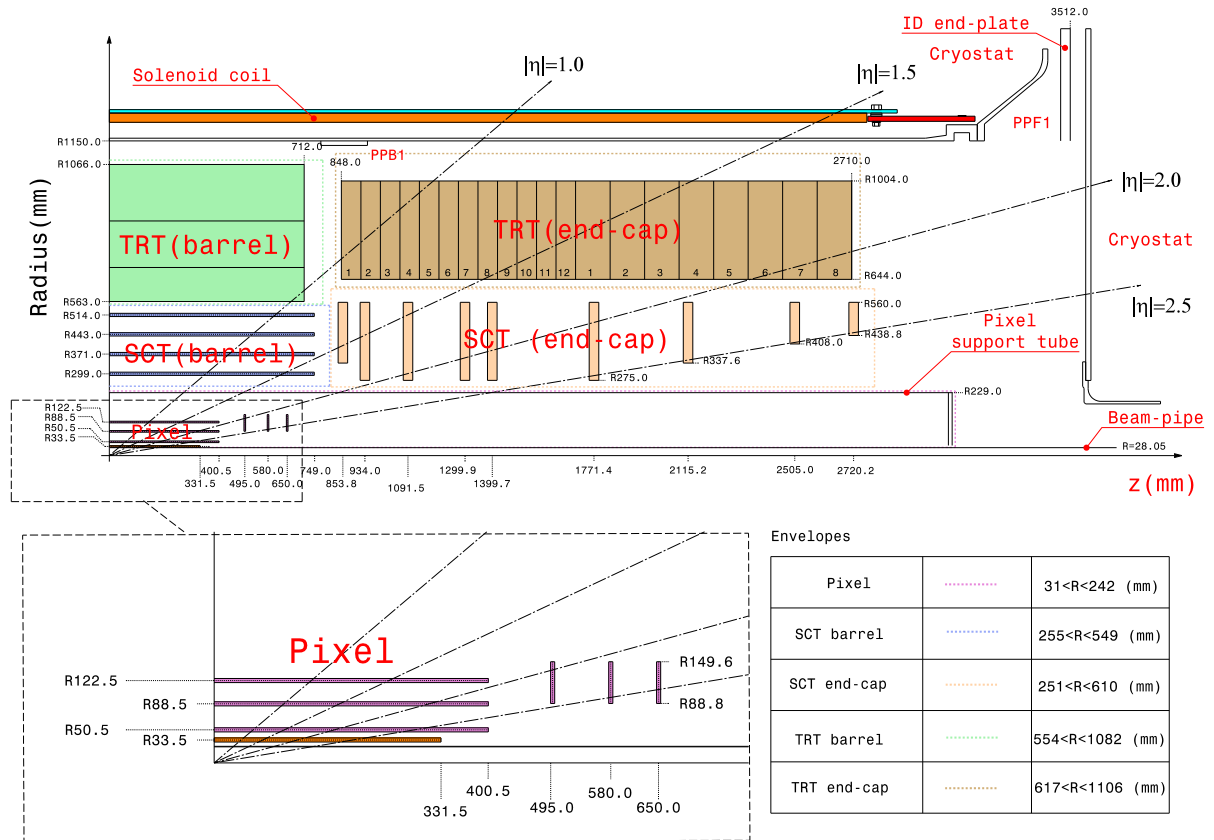


Figure 3.3 — r - z cross-section view of an ATLAS Inner Detector quadrant with the localisation of the active layers and their structural envelope. The bottom panel shows a zoomed-in view of the pixel layers (including the IBL). [112]

3.3.1 The pixel detector

The innermost part of the ID is composed of three layers of pixels (plus the IBL, see below), nominally providing three high precision hits. Three approximately 800 mm long concentric cylindrical layers at radii 50.5 mm, 88.5 mm, and 122.5 mm provide coverage until $|\eta| = 2.0$, and are complemented by three disks perpendicular to the beam axis on each side, ensuring three hits until $|\eta| = 2.5$. The pixel system is designed to be operated at 150 V in its early time, but voltage up to 600 V will be needed to maintain the hit efficiency at the end of their expected lifetime (Run 3), due to radiation damages. The whole pixel structures span the radii from 31 mm to 242 mm.

The pixels are distributed over 1744 sensors of $19 \times 63 \text{ mm}^2$, with 47232 pixels on each sensor. This amounts to approximately 80 million readout channels and cover an area of almost 2 m^2 . Each pixel is typically $50 \times 400 \text{ }\mu\text{m}^2$ in the $R\text{-}\phi \times z$ plane ($R\text{-}\phi \times R$ plane for the endcap), bringing a resolution of about $10 \times 115 \text{ }\mu\text{m}^2$.

The three barrel layers have been complemented during the first long shutdown (2013 – 2014) by the insertion of fourth layer, the Insertable B-Layer (IBL). The beam pipe was specifically radially reduced to allow the positioning of this layer at a radius of 33.5 mm, *i.e.* between the beam pipe and the first pixel layer. The IBL must provide enhanced vertex measurements and sustained tracking performance until the end of Run 3 despite radiation damages to the earlier layers.

3.3.2 The Strip Semiconductor Tracker (SCT)

For cost and readout bandwidth reasons, the next silicon layers use micro-strips instead of pixels. In the barrel each of these strips is 126 mm long in z and $80 \text{ }\mu\text{m}$ wide in $R\text{-}\phi$, and their sizes vary in the endcap (see Table 4.7 of Ref. [97]). Due to their higher dimension compared to the pixels, their resolution is lower but still reaches about $17 \times 580 \text{ }\mu\text{m}^2$ in the $R\text{-}\phi \times z$ plane ($R\text{-}\phi \times R$ plane for the endcap).

The strips are arranged in four concentric layers spanning 1500 mm along the z -axis between $R = 299 \text{ mm}$ and 514 mm in the barrel, and in nine disks perpendicular to the beam axis in each endcap between $z = 853.8 \text{ mm}$ and $z = 2720.2 \text{ mm}$. The strips are oriented along the beam axis (resp. radially) in the barrel (resp. endcaps), and each layer includes two sensors tilted by 40 mrad improving the resolution in the z (resp. R) direction. This layout ensures at least four more additional precision points are added to the tracking. The strips structural devices occupy the volume $R = 255 \text{ mm}$ to 549 mm in the barrel and $R = 251 \text{ mm}$ to 610 mm in the endcap.

The strips cover a total area of 63 m^2 , amounting to ~ 6.3 millions additional channels. Similarly to the pixels, their design operating high voltage is 150 V, but will need to reach 300 V to 350 V at the end of Run 3 to ensure a nominal hit efficiency.

3.3.3 The Transition Radiation Tracker (TRT)

The third sub-detector composing the tracking system is the transition radiation tracker (TRT). It does not rely on semiconductor material but is rather a gaseous detector with a continuous active area in the radii 563 mm to 1066 mm in the barrel and 644 mm to 1004 mm in the endcap. Despite a much lower resolution of $130 \text{ }\mu\text{m}$ in the $R\text{-}\phi$ plane, it provides a high number of points along the track and adds particle identification capabilities to the tracking.

The basic element of the TRT is a tube 4 mm in diameter (called a *straw*) with an anode wire in the middle. It is filled with a xenon or an argon gas mixture which is ionised by the charged

particles going through. The ionisation electrons drift toward the anode, producing a measurable electron avalanche signal. The position along the tube can then be deduced by measuring the time of arrival of the current to the readout electrode.

The second purpose of the TRT is provide discrimination between electrons and charged light hadrons, thanks to a process called *transition radiation*: when a relativistic particle crosses the interface between two media, it can emit a photon with typical energy 5 keV to 30 keV, *i.e.* an X-ray. The probability for such radiation to occur is proportional to the particle boost, *i.e.* to E/m . For a given energy (or equivalently in this case, momentum), it is therefore 270 times more likely to occur for an electron ($m_e \approx 0.5$ MeV) than for a charged pion ($m_{\pi^\pm} \approx 135$ GeV). In the TRT, the straws are separated by polypropylene fibres or foils serving as a radiator medium. The emitted X-rays are absorbed by the xenon in the gas mixture, which increases the signal. Each straw is thus read with two thresholds, one corresponding to the ionisation (normal threshold), and one corresponding to the absorption of the X-ray (*high-threshold*).

In the barrel, each straw is approximately 1440 mm long, and oriented along the beam pipe. The straws are separated by ~ 7 mm, the interstice being filled with polypropylene fibres. They are stacked to provide 73 planes in the R direction, grouped in three rings. In the endcaps, the straws are about 370 mm long and oriented radially. There is a total 160 planes along the z direction on each side, each plane being composed of 768 straws. This amounts to $\sim 350\,000$ readout channels in total. Each of the two endcaps is subdivided in two sets of wheels: the 12 inner wheels have 8 straws each, separated by 8 mm, while the 8 outer wheels have 8 straws spaced by 15 mm (see Table 4.8 in Ref. [97] for more details). The interstice is filled with polypropylene foils.

The barrel part is located in the volume $563\text{ mm} < R < 1066\text{ mm}$ and $|z| < 712\text{ mm}$, and the endcaps in the volume $644\text{ mm} < R < 1004\text{ mm}$ and $848\text{ mm} < |z| < 2710\text{ mm}$. As a consequence, the $|\eta|$ coverage of the TRT stops at $|\eta| = 2.0$. In the covered region, this layout nevertheless allows for an average 36 hits (22 hits in the endcap/barrel transition region around $|\eta| = 1$) to be recorded. For electrons, the number of high-threshold hits is typically 8.

The nominal mixture filling the straws is made of xenon, carbon dioxide and dioxygen in the proportion $\text{Xe}:\text{CO}_2:\text{O}_2 = 70:27:3$. Due to a gas leak in the TRT during Run 2, the first two rings in the barrel and nine endcap wheels are switched to an argon mixture instead (in proportion $\text{Ar}:\text{CO}_2:\text{O}_2 = 70:28.5:1.5$). Due to a lower absorption cross-section of the transition radiation in argon than in xenon, the identification performance is degraded.

3.3.4 Central solenoid magnet

In addition to providing the direction of the tracks, and therefore of their associated particles, the tracker also provides a measurement of their momentum. This is achieved by measuring the bending radius r of the trajectory within a known magnetic field B , using the formula already seen in Eq. 2.2: $p = qBr$. The rotation direction also gives the charge of the particle associated to the track. The magnetic field in the ATLAS inner detector is produced by a 5.8 m long solenoid magnet operated at 7.7 kA, creating a strong 2 T axial field.

The main constraint on the solenoid design is that it is located between the ID and the calorimeter: as additional material in front of the calorimeter impacts the energy resolution, the material budget of the whole solenoid system is reduced to its minimum. This is achieved with two optimisations: first, the conducting material uses a specially designed superconducting Al-stabilised NbTi wire 12 mm thick, allowing for one single layer of conductor winding. Second, the solenoid has no dedicated cryostat but is rather inserted in the same cryostat as the LAr calorimeter, saving two additional insulating walls. Thanks to these optimisations, the solenoid weights only 5.4 t and is 50 mm thick, located between $R = 1.23\text{ m}$ and $R = 1.28\text{ m}$, adding

$0.66 X_0$ at normal incidence.

3.4 The electromagnetic and hadronic calorimeters

The ATLAS calorimeter system provides a measurement of the energy of all detectable particles that are created during a collision. The calorimeters hermetically cover the whole solid angle up to $|\eta| = 4.9$ with various components: the first one is the electromagnetic calorimeter (ECal) using liquid argon as active material and lead as absorbers. It covers the region $|\eta| < 1.48$ in the barrel and $1.375 < |\eta| < 3.2$ in the endcaps. Further on the particle trajectory lies the hadronic calorimeter (HCal), which covers the range $|\eta| < 1.7$ in the barrel. This component uses a different technology, with plastic tiles as a scintillating medium and absorbers made of steel. The endcap hadronic calorimeter (HEC) covers the region $1.5 < |\eta| < 3.2$ and uses a technology similar to the ECal, with argon as active material and a copper absorber. Finally a forward calorimeter (FCal) covers the region $3.1 < |\eta| < 4.9$, using again liquid argon as active medium, but copper or tungsten as absorber. This last calorimeter provides both electromagnetic and hadronic measurements capabilities. Figure 3.4 shows the layout of the different components and are further described below, after a brief reminder of general calorimetry principles. Complementary extensive information can be found in Refs. [113–116] and [97].

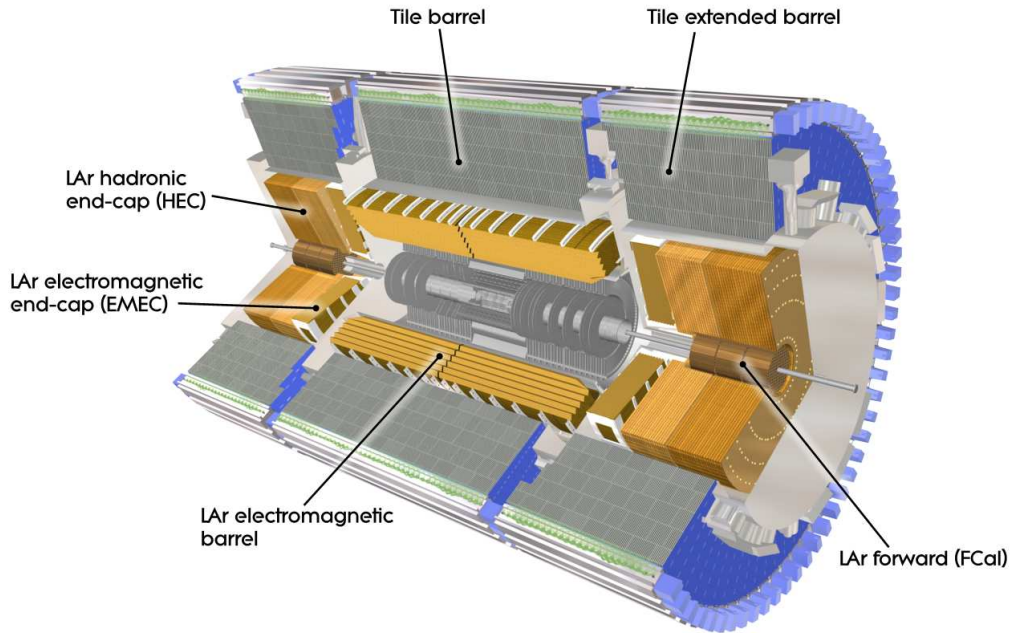


Figure 3.4 — General view of the various calorimeter systems of ATLAS. The various liquid argon components are shown in shades of orange, while the tile calorimeter is shown in shades of green. In the centre, the dark grey represents the ID and the central solenoid. The lighter grey parts on the outside and between the tile calorimeter and the LAr calorimeters show the three cryostat vessels (barrel and two endcaps). [97]

3.4.1 Principles of calorimetry

Energy measurements in calorimeters rely on the measurement of a fraction (ideally, the totality) of the particle energy in an *active* medium by stopping it. Two types of calorimeters can be used: homogeneous or sampling calorimeters. Homogeneous calorimeters consist in a single active medium which at the same time stops the particle and conducts the produced signal (often

optically). They usually have a better resolution than the sampling calorimeters since all the material is active, but are less easy to segment (granular readout). This is the case of the CMS electromagnetic calorimeter made of PbWO_4 crystals which stop high energy electromagnetic particles while producing photons in the visible range, and are therefore transparent in order to readout the signal. Sampling calorimeters, such as those of ATLAS, consist in several layers of active material that produce the signal, separated by layers of *absorbers i.e. passive* material. Due to this presence of passive material, the energy resolution is usually worse than for homogeneous calorimeters. Only details relevant for sampling calorimeters are described in the following.

Depending on their energy, particles interact through various processes with matter. A general example over a broad range of energy is given in Figure 3.5. More specialised figures for electrons and photons interacting with lead is pictured in Figure 3.6. A complete description can be found in Chapter 33 of Ref. [41], the relevant ones for calorimetry at particle accelerators are recalled below.

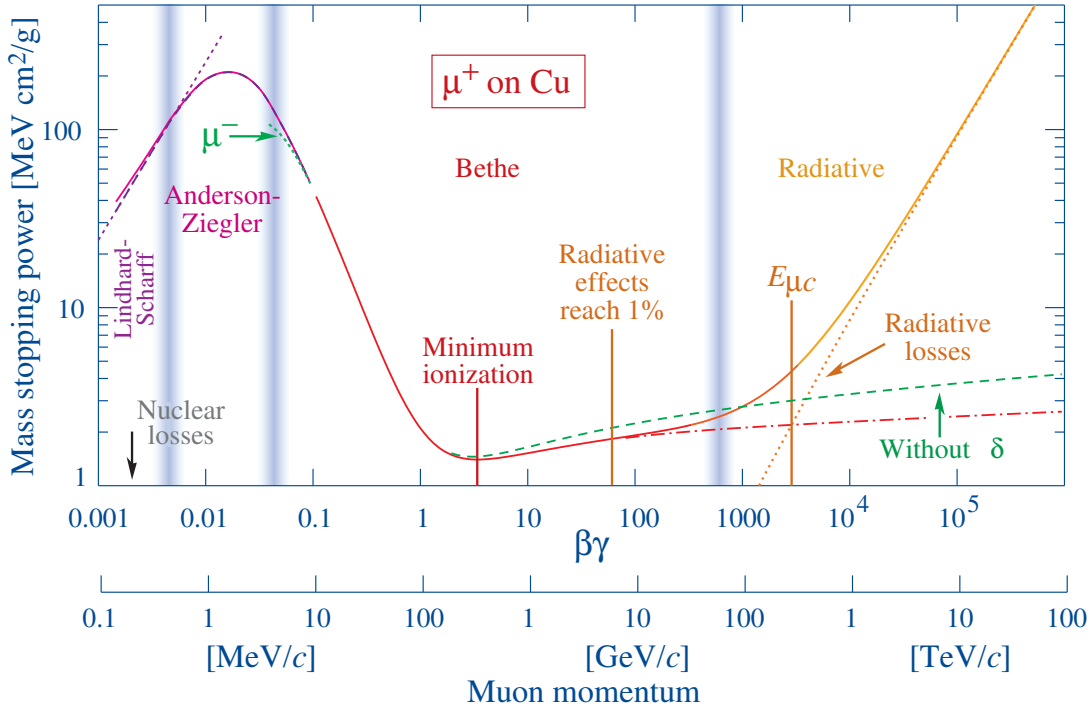


Figure 3.5 – Mass stopping power of muons in copper as a function of $\beta\gamma = p/Mc$. [41]

Ionisation (charged particle with $\mathcal{O}(0.1) < \beta\gamma < \mathcal{O}(10)$): in the regime $\beta\gamma \approx \mathcal{O}(0.1)$ to $\mathcal{O}(10)$, the main process of interaction of a charged particle in matter is the scattering off the electrons of the atoms composing the medium. The transferred energy is enough to unbind an electron from this atom, producing ionisation. This process is described by the Bethe-Bloch formula, giving the mean energy loss by unit length:

$$-\left\langle \frac{dE}{dx} \right\rangle = K z^2 \frac{Z}{A} \frac{1}{\beta^2} \left[\frac{1}{2} \ln \left(\frac{2m_e c^2 \beta^2 \gamma^2 W_{\max}}{I^2} \right) - \beta^2 \right], \quad (3.4)$$

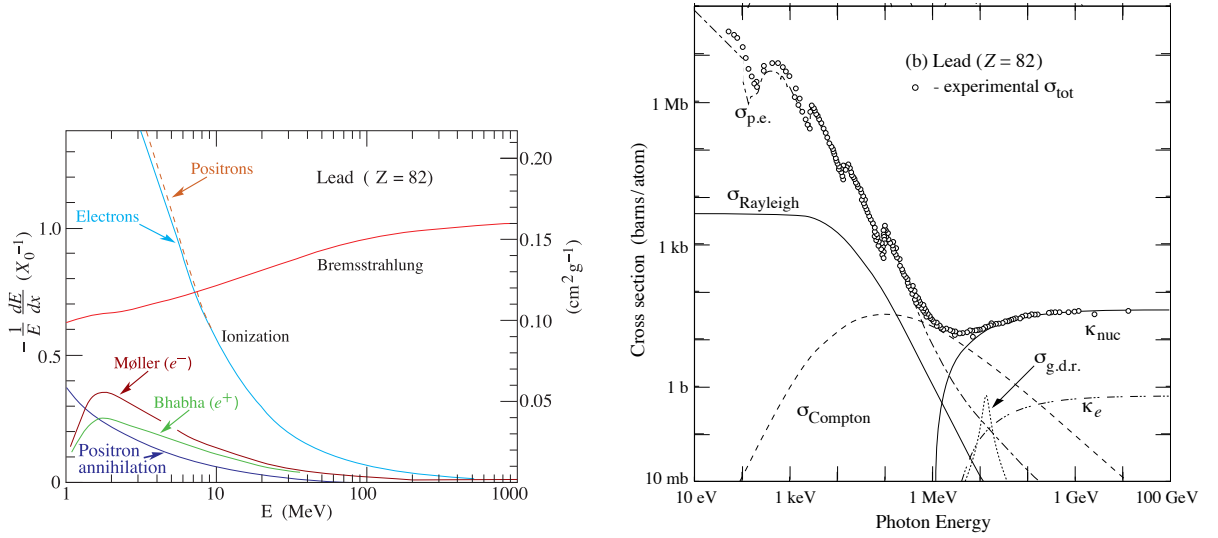
with K a constant, z the charge of the incoming particle, A and Z the atomic mass and nucleus charge of the medium, β and γ the relativistic parameters, I is the ionisation energy of the atoms in the medium, and $T_{\max} = \frac{2m_e c^2 \beta^2 \gamma^2}{1 + 2\gamma m_e/M + (m_e/M)^2}$ where M is the mass of the incoming particle. $\langle dE/dx \rangle$ is often expressed in $\text{MeV cm}^2 \text{g}^{-1}$. The *linear stopping power* of the material is given by $\langle dE/dx \rangle \times \rho$ where ρ is the density of the material, and is expressed in MeV/g . As seen in

Figure 3.5, the energy loss reaches a minimum for a given $\beta\gamma$: particles with this momentum are called *Minimum Ionising Particles* (MIP).

Bremsstrahlung (charged particles with $\beta\gamma > \mathcal{O}(1000)$): at higher energies, the dominant process of charged particle interaction with matter is the emission of a high energy photon, due to the deviation of the incoming particle in the electric field of an atom (similar to Chapter 2.3).

Rayleigh and Compton scattering (photons with $E \lesssim \mathcal{O}(10 \text{ MeV})$): at relatively low energies, photons mainly scatter on the electrons of the atoms constituting the medium.

Photon conversion (photons with $E \gtrsim \mathcal{O}(10 \text{ MeV})$): at high energies, photons traversing a medium mainly undergo electron-positron pair creation in the electromagnetic field of the nuclei (and of the electrons in a lesser measure).



(a) Electron energy loss in lead as a function of its energy.

(b) Cross-section of photon interaction with lead as a function of its energy.

Figure 3.6 – Interaction of electrons and photons with lead: (a) fractional energy loss per radiation length in lead as a function of electron or positron energy, (b) total cross-section of photon interaction in lead as a function of its energy. p.e.: photo-electric effect; Rayleigh: coherent scattering (atom not ionised or excited); Compton: incoherent scattering off of an electron; κ_{nuc} , κ_e : pair production from the nuclear or electron field; g.d.r.: giant dipole resonance and other photo-nuclear interaction, breaking up the nucleus. [41]

3.4.2 Application to the ATLAS electromagnetic calorimeter

Taking the example of a high energy electron (tens to hundreds of GeV) produced in a collision at the LHC and hitting the ATLAS electromagnetic calorimeter, the electron will undergo several bremsstrahlung interactions, losing a fraction of its energy and creating an additional high energy photon each time. Each of these photons then converts into a pair creating two high energy electrons, each of them producing the same effect as previously. The same applies for incoming high energy photons, with just one step shifted. These two processes interplay, creating more and more particles to form a structure called an *electromagnetic cascade* (or *shower*) in the calorimeter, losing energy until the subsequent particles reach a threshold called *critical energy* and denoted E_c .

This energy corresponds to the energy at which an electron loses as much energy by bremsstrahlung as by ionisation⁵. Below this threshold, fewer and fewer particles are created thus stopping the longitudinal development of the shower. The number of particles created in the shower is then proportional⁶ to E/E_c . The energy measured is proportional to the number of final electrons, which have a low enough energy to produce the ionisation signal detected in the calorimeter. The calorimeter resolution therefore depends on the number of particles, so a small E_c is needed. In sampling calorimeters, the shower is mainly produced by the dense material composing the absorbers, with particles spreading out into the active layer providing the measurement. Some of the particles however never reach the active material, so a sampling calorimeter only measure a fraction of the total energy. The number of particles created differs between two showers with similar energy due to Poissonian fluctuations of the processes. This explain the sampling term ($\sigma_E/E \sim \sqrt{E}/E \sim 1/\sqrt{E}$) entering the energy resolution of a calorimeter presented in Eq. 5.12. For example, the critical energy in lead, making the absorber of the ATLAS electromagnetic calorimeter, is $E_c = 7.4 \text{ MeV}$.

A convenient variable describing the stopping power of a medium and therefore the longitudinal extension of a shower is the *radiation length*, denoted X_0 . It is defined as the average path length in the material after which an electron has only $1/e$ of its initial energy left. It can be shown [41] that this also corresponds to $9/7$ of a photon mean free path in the medium (before undergoing a pair creation). To good approximation, the longitudinal extent of a shower is $X_0 \times [\ln(E/E_c) + C]$ where $C = 0.5$ for a photon and -0.5 for an electron: the first interaction of a photon is a conversion, which is less likely than a bremsstrahlung emission for an electron. In the ATLAS electromagnetic calorimeter for example, the radiation length is $X_0 \approx 1.9 \text{ cm}$ (at $|\eta| \approx 0$).

Another useful quantity is the lateral extent of the shower, which is measured by the Molière radius $R_M = X_0 E_s / E_c$. $E_s = 21.2 \text{ MeV}$ [41] is the scale energy of the multiple scattering process ($E_s = m_e c^2 \sqrt{4\pi/\alpha}$), and E_c uses the Rossi definition. The Molière radius represents the radius of a cylinder containing 90% of the shower energy, which in the case of the ATLAS calorimeter is $R_M \approx 3.8 \text{ cm}$ (for $|\eta| \approx 0$).

A summary of atomic and nuclear properties of elements and compounds composing the ATLAS electromagnetic calorimeter are listed in Table 3.2. For an heterogeneous material, its various compounds are combined using (Ref. [41]):

$$\frac{1}{X_0} = \sum_j \frac{w_j}{X_{0,j}} \quad \text{and} \quad \frac{1}{R_M} = \frac{1}{E_s} \sum_j \frac{w_j E_{c,j}}{X_{0,j}}, \quad (3.5)$$

where $w_j = \rho_j \times L_i$ is the weight fraction of compound j , used over a length L_i .

3.4.3 Central and endcap electromagnetic calorimeter

The purpose of the ATLAS electromagnetic calorimeter is to provide a precise measurement of the energy of electrons and photons produced during a collision. It is segmented into three layers, which allows for some pointing capabilities, especially for unconverted photons that do not leave a track in the ID. The first layer is highly segmented in η ⁷ to offer a discrimination between prompt

⁵↑ An alternative definition of the critical energy is the energy for which the loss by ionisation equals the particle energy (Rossi definition).

⁶↑ Assuming a toy model for the showers where after each X_0 (see below), a particle splits into two particles (pair creation $\gamma \rightarrow e^- e^+$ or bremsstrahlung $e^- \rightarrow e^- + \gamma$) of equal energy, and denoting $t = x/X_0$ the longitudinal variable for the shower, we have the number of particles $N(t) = 2^t$ and $E(t) = E_0/N(t) = E_0/2^t$, so $t = \ln(E_0/E)/\ln(2)$. Assuming that the shower stops when $E(t) \approx E_c$, $E(t_{\max}) = E_c = E_0/2^{t_{\max}}$ and therefore $N(t_{\max}) = 2^{t_{\max}} = E_0/E_c$.

⁷↑ A fine separation in ϕ in the first layer is not possible due to bremsstrahlung in upstream material.

Table 3.2 – Atomic and nuclear parameters of some elements composing the ATLAS electromagnetic calorimeter (numbers from Ref. [41]). Structural components are made of stainless steel, approximated here with iron. The kapton is approximated by polyimide film. The prepreg and glue (see Section 3.4.3 and Ref. [114]) are not listed here.

		Liquid Argon LAr	Lead Pb	Iron Fe	Copper Cu	Polyimide film (C ₂₂ H ₁₀ N ₂ O ₅) _n
Atomic number (Z)		18	82	26	29	
Atomic mass number (A)	[g/mol]	39.948	207.2	55.845	63.546	
Density (ρ)	[g/cm ³]	1.396	11.35	7.874	8.960	1.42
Radiation length (X_0)	[g/cm ²]	19.55	6.37	13.84	12.86	40.58
X_0/ρ	[cm]	14.00	0.5612	1.757	1.436	28.57
e^- critical energy (E_c)	[MeV]	32.84	7.43	21.68	19.42	81.01
Molière radius (R_M)	[g/cm ²]	12.62	18.18	13.53	14.05	10.62
R_M/ρ	[cm]	9.043	1.602	1.719	1.568	7.479

photons and pairs of boosted photons from pion decays for example. Its segmentation also allows for a precise electromagnetic shower shape width measurement, helping the electron and photon identification. It is made of two half barrels ($|\eta| < 1.48$) and two endcaps ($1.375 < |\eta| < 2.5$) with full coverage up to $|\eta| = 2.5$. In the transition region between the barrel and the endcap ($1.4 < |\eta| < 1.5$), the measurement combines the information from the two elements but is largely degraded due to a large amount of passive material (readout cables and services for the ID, *etc.*) in front of the calorimeter.

The electromagnetic calorimeter is a sampling calorimeter, using lead as an absorber and liquid argon (LAr) as an active medium for ionisation. Its structure is highly recognisable with its accordion shape. Indeed, if the absorbers were laid in concentric layers with active material in between, the ϕ coverage would be symmetrical, but the readout of all cells would have to be made at the edges of the barrel, requiring longer cable and leading to larger impedances, a slower signal and large cross-talks. On the other hand, if absorbers were made of planar sheets in the R - z plane, the readout would be easier but outgoing particles at nearby ϕ would traverse either only an absorber along its path, or only active material, inducing large differences in the measured energies. The idea solving both the above problems is to have a radial layout but using non-planar absorber sheets: by folding the absorbers in the radial direction, the complete ϕ symmetry is restored. This is demonstrated in the top panel of Figure 3.7.

The barrel is constituted of two 3.2 m long half barrels, occupying the volume between $R \approx 1.4$ m and $R \approx 2$ m. Each half barrel is formed of 16 modules covering $\Delta\phi = \pi/8 \approx 0.4$. Each module is made of 64 stacked absorbers, for a total 1024 absorber in each half barrel.

The total radiation length at $|\eta| = 0$ is about $22 X_0$ [97]. However, since the barrel is uniform as a function of z , particles see an increasing radiation length, both per absorber and total, as $|\eta|$ increases. To have a constant number of X_0 as a function of $|\eta|$, the lead thickness would therefore have to decrease continuously with increasing $|\eta|$. Since this is not practically feasible, the lead thickness instead changes only once, at $|\eta| = 0.8$: the lead absorbers are 1.53 mm thick for $|\eta| < 0.8$, and 1.13 mm thick for $0.8 < |\eta| < 1.48$. This change reduces the sampling fraction seen by particles after $|\eta| > 0.8$. The total radiation length with this change amounts to $22 X_0$ at $|\eta| = 0$, up to $30 X_0$ at $|\eta| = 0.8$, and changes from $24 X_0$ to $33 X_0$ between $|\eta| = 0.8$ and $|\eta| = 1.3$, where is located the “corner” of the barrel. It then decreases until $|\eta| = 1.48$. A summary of the values is described in Table 3.3.

To ensure the rigidity of the absorbers, a 0.2 mm thick stainless steel sheet is glued on each of their side with a resin-impregnated glass-fibre fabric (“prepreg”). While the thickness of the steel is constant across all the barrel, the prepreg thickness also changes at $|\eta| = 0.8$, from 0.13 mm

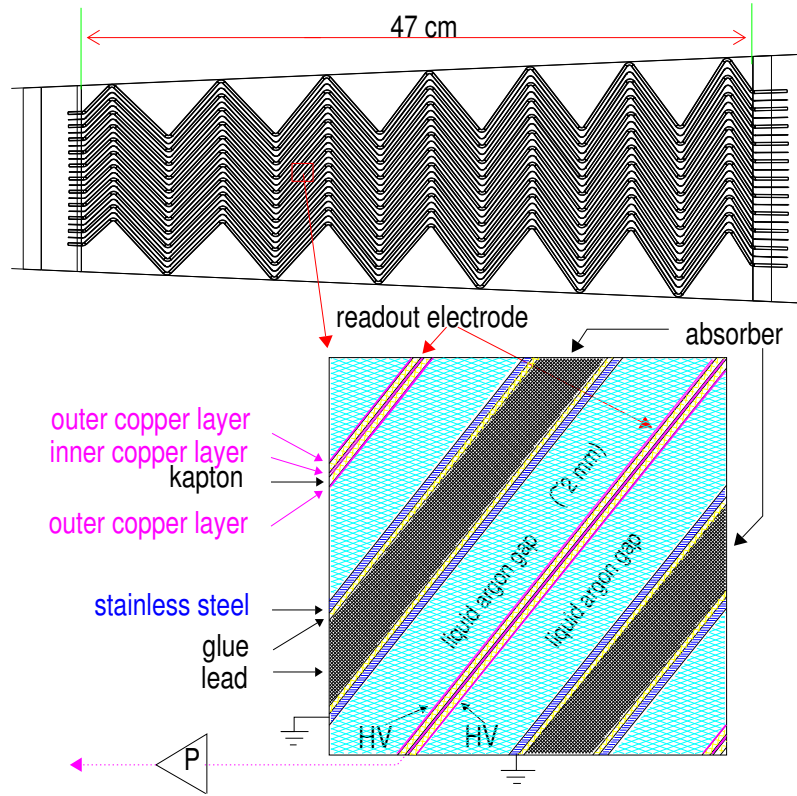


Figure 3.7 – Structure of the active/passive material sandwich for the LAr electromagnetic calorimeter. [114]

to 0.33 mm. The space between two absorbers is filled with two liquid argon gaps, each ~ 2 mm, and separated by a kapton electrode. Each electrode includes three copper layers: the two outer ones are independently fed by two high-voltage sources (to ensure redundancy in case of shorts or other HV failure) providing the 2000 V needed for the ionisation electrons to drift. The stainless steel is grounded to ensure a nominal electric field in the gap. The third kapton copper layer, in between the two others, reads out the ionisation signal (see below). To maintain a constant LAr gap between two absorbers in the barrel, the folding angle changes from 91.9° at inner radius to 68.5° at outer radius. This is further ensured by inserting a honeycomb spacer in the liquid argon gap, which also helps to avoid short-circuits.

A particle going through the gap ionises the liquid argon along its path, creating free electrons that drift toward the electrode. The current created by the electron drift is read out by the electrode in the middle of the kapton by capacitive coupling, creating a signal that rises sharply and then decreases linearly. Indeed, the argon ionisation can be considered instantaneous given the velocity of the ionising passing particle compared to the electron drift velocity in the liquid argon. At the beginning, all the created electrons drift toward the HV electrode, but they are absorbed by the electrode conductor as they reach it and stop contributing to the LAr gap current, decreasing the signal until all electrons have reached the electrode. The total drift time is about 450 ns at $T \approx 89$ K for a ~ 2 mm gap under 2000 V. The electric signal read by the electrodes therefore has a triangle shape, as seen in Figure 3.8. It is then passed through a bipolar shaper, giving the signal the characteristic shape seen in the same Figure. Due to the fast bipolar shaping, the measured response is proportional to the current rather than the total charge, if the signal was fully integrated. The electron drift velocity in the liquid argon depends on the temperature, so the energy response of the calorimeter, proportional to the current, is also sensitive to temperature changes (about 2%/K [117]). In order to maintain the LAr temperature constant, the LAr calorimeters are embedded in a cryostat that keeps the temperature stable at the level of ~ 0.01 K.

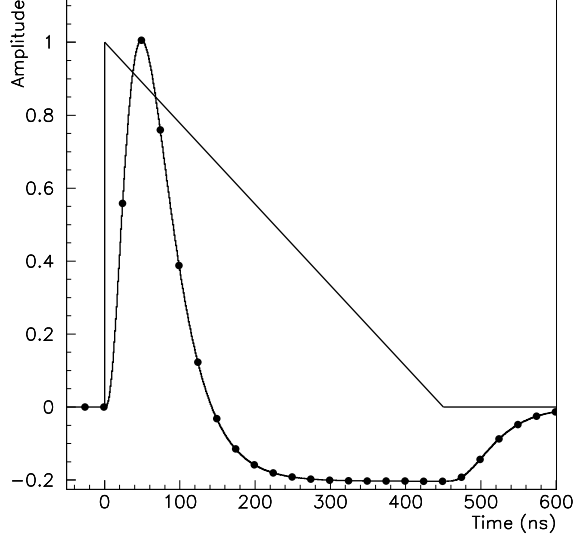


Figure 3.8 – Amplitude versus time for the current in the liquid argon gap (triangular shape), and output signal from the electronics after bipolar shaping. Each dot represents one sampling of 25 ns. [97]

In front of the accordion, an additional component made of a single layer of only 11 mm of liquid argon is inserted (Figure 3.9). This element is called the *Pre-Sampler* (PS) and is used to measure the energy lost in the material in front of the calorimeter (see Chapter 5). The PS is assembled in seven modules covering 0.2×0.2 in $\Delta\eta \times \Delta\phi$, plus one covering $1.4 < |\eta| < 1.52$ (and still 0.2 in ϕ). Electrodes with the depth of the liquid argon layer and covering $\Delta\phi = 0.2$ are inserted perpendicularly to the beam axis. Each anode or cathode is separated from its neighbour cathode or anode along the z direction by 1.9 mm to 2.0 mm, creating an ~ 2 mm argon gap nominally operated at 2 kV (1.2 kV in Run 2). The cathodes are made of two independently fed electrodes providing the high voltage, while the anodes are made of three electrodes: the two outer independently provide the high voltage while the third is used to read out the ionisation signal by capacitive coupling, as for the accordion electrodes. The readout is made by ganging together a sufficient number of electrodes in η to have a granularity of $\Delta\eta = 0.025$ (since they are placed at regular z intervals, the number of electrodes for a constant $\Delta\eta$ interval increases with $|\eta|$), and by etching in two halves the electrodes in ϕ to obtain a granularity of $\Delta\phi = 0.1$. One additional module is placed in the endcaps to cover the $|\eta|$ range $[1.5, 1.8]$ with the same granularity.

The endcap electromagnetic calorimeters are built using the same principles as the barrel part. The outer wheels, covering $1.375 < |\eta| < 2.5$, are made of 768 absorbers each and the inner wheels, covering $2.5 < |\eta| < 3.2$, are made of 256 absorbers each. The absorbers are 1.7 mm thick in the outer wheel and 2.2 mm thick in the inner wheel, and are interleaved with electrodes in a similar way as for the barrel. The folding is made along the z axis, but the angle varies along R .

The endcap argon gaps, however, cannot be maintained constant as a function R : the argon drift gap varies from 2.8 mm to 0.9 mm in the outer wheel and from 3.1 mm to 1.8 mm in the inner wheel (with decreasing radius). In order to compensate for a decreasing gap size from the outer to the inner radii, the high voltage is varied as a function of $|\eta|$ (Figure 3.10): while in the barrel, each HV line feeds a 0.2×0.2 region in $\Delta\eta \times \Delta\phi$ with 2000 V (nominally), in the endcap, the HV is as high as ~ 2500 V around $|\eta| \approx 1.5$ and goes down to ~ 1000 V around $|\eta| \approx 2.5$ and has a finer segmentation, covering smaller areas in η . The detail of nominal HV values is displayed in Table 3.3.

In the accordion for $|\eta| < 2.5$, the electrodes are etched to provide three readout layers, denoted

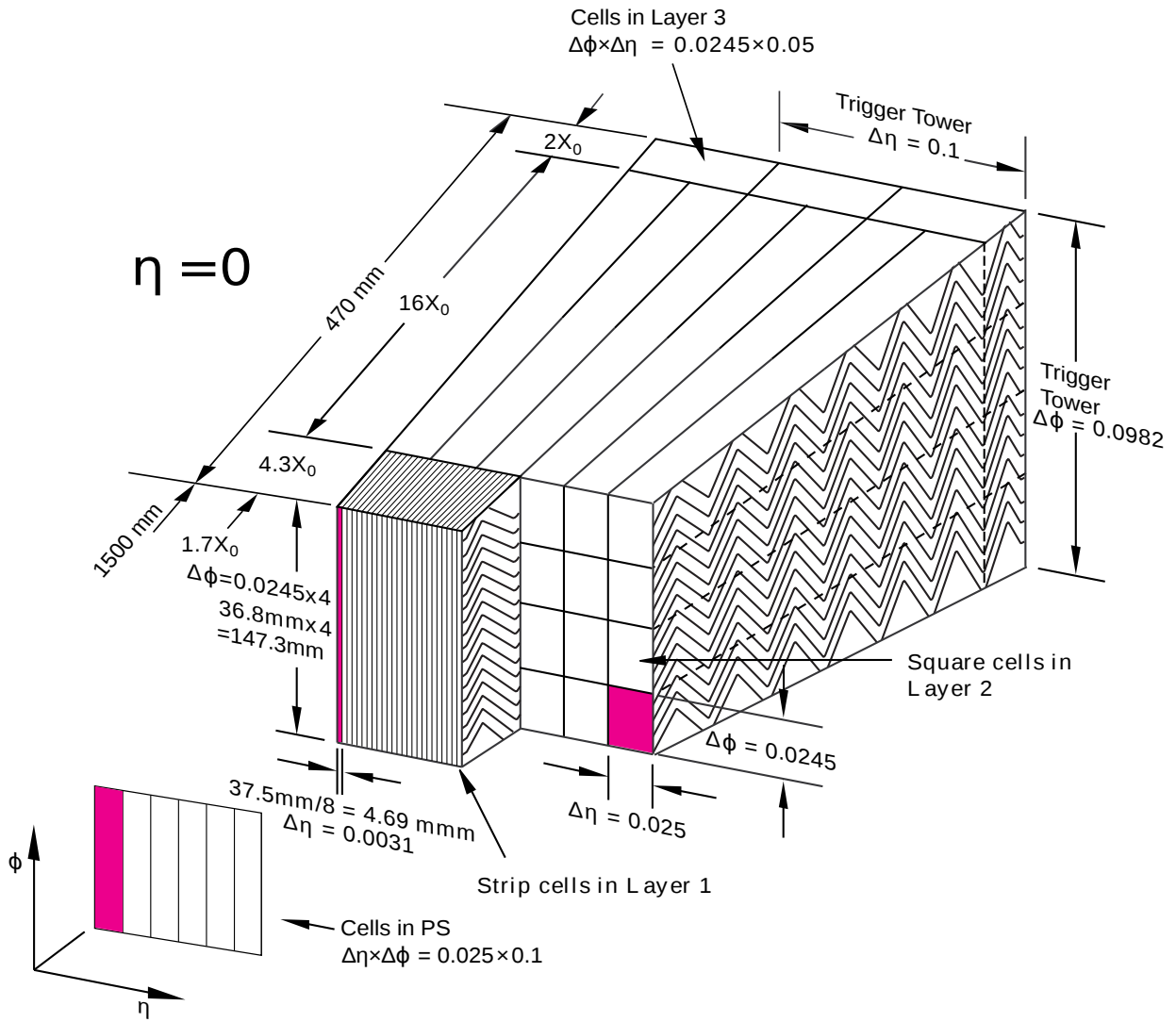


Figure 3.9 – Scheme of the “accordion” electromagnetic calorimeter in the region $0 < \eta < 0.15$. The presampler in the front as well as the three layers are visible. [118]

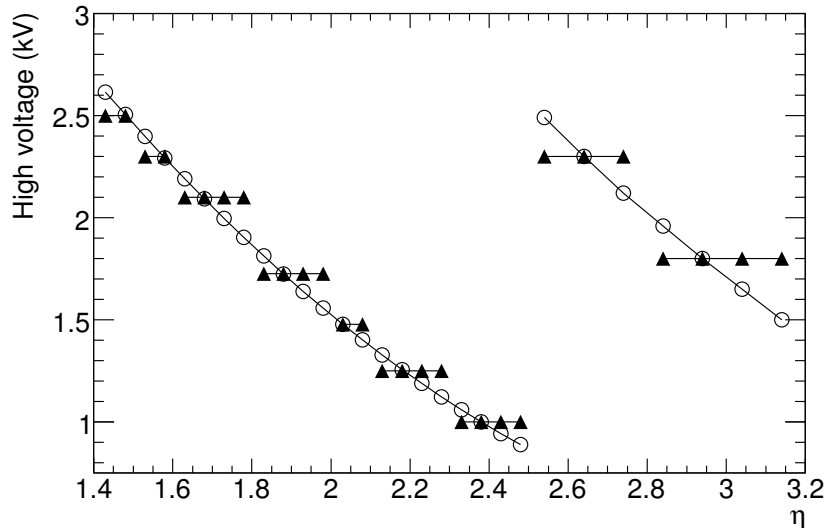


Figure 3.10 – Nominal high-voltage settings in the outer and inner endcap wheels as a function of η . Since the LAr gaps are wider at outer radii than at inner radii, a uniform response can be obtained by continuously varying the HV as a function of the radius (or equivalently, as a function of η), which is shown with the open circle. Given that a single HV line feeds a region covering $\Delta\eta = 0.1$ or 0.2 , this setting can only be approximated, which is shown by the filled triangles. [97]

front (L1), *middle* (L2) and *back* (L3). The front layer corresponds to a constant radiation length of $6X_0$ including the material in front of the calorimeter, the middle layer designed to end after and integrated $22X_0$ of material [114]. The back layer forms the remaining of the calorimeter and is at least $2X_0$. However, as previously mentioned, since the lead thickness is constant as a function of z (with only one change at $|\eta| = 0.8$), the total radiation length seen by the outgoing particles increases as a function of $|\eta|$. In order to maintain a constant number of X_0 in the first and second layer, the length of the electrode strips therefore decreases as a function of $|\eta|$. The third layer is made of the remaining length and its depth and effective radiation length varies by a large factor as a function of $|\eta|$. This can be seen in Figure 3.11 for the barrel and Figure 3.12 for the endcaps.

The electrode geometry ensures the projectivity of the measurements with respect to the interaction point. In the second layer, making up the bulk of the energy measurement, the electrodes are etched in patches of width $\Delta\eta = 0.025$, and four consecutive electrodes are ganged together in ϕ to form a 0.025×0.025 in $\Delta\eta \times \Delta\phi$, which constitutes the standard unit area in the calorimeter. In the third layer, mainly measuring the eventual energy overflowing the ECal toward the HCal, a coarser granularity of $\Delta\eta = 0.05$ is performed, still with $\Delta\phi = 0.025$. In the first sampling, each cell is etched to have $1/8$ of the L2 width in η , but the ϕ granularity is reduced to $\Delta\phi = 0.1$ by ganging 16 electrodes together. Due to their long and thin shape, the L1 cells are often called *strips*. This high η granularity allows to separate two collimated photons from pion decays from prompt photons. In the endcap, the η granularity of the strips is not constant, increasing from $0.025/8$ until $|\eta| = 1.8$ to $0.025/6$ for $1.8 < |\eta| < 2.0$, to $0.025/4$ for $2.0 < |\eta| < 2.4$. Moreover in the edges and in the transition region, the complicated geometry makes the granularity irregular. Detailed information about the granularity may be found in Table 3.3. The total number of readout channels for the ECal amounts to more than 170 000.

While it seems natural to read the first layer in the front of the calorimeter and the third layer in the back, there is no obvious choice for where to put the front-end board of the second layer. Given that the signal must be physically driven from the L2 cell to the board with a conductive area on the electrode, making room among the already dense layout of the front layer strips is not

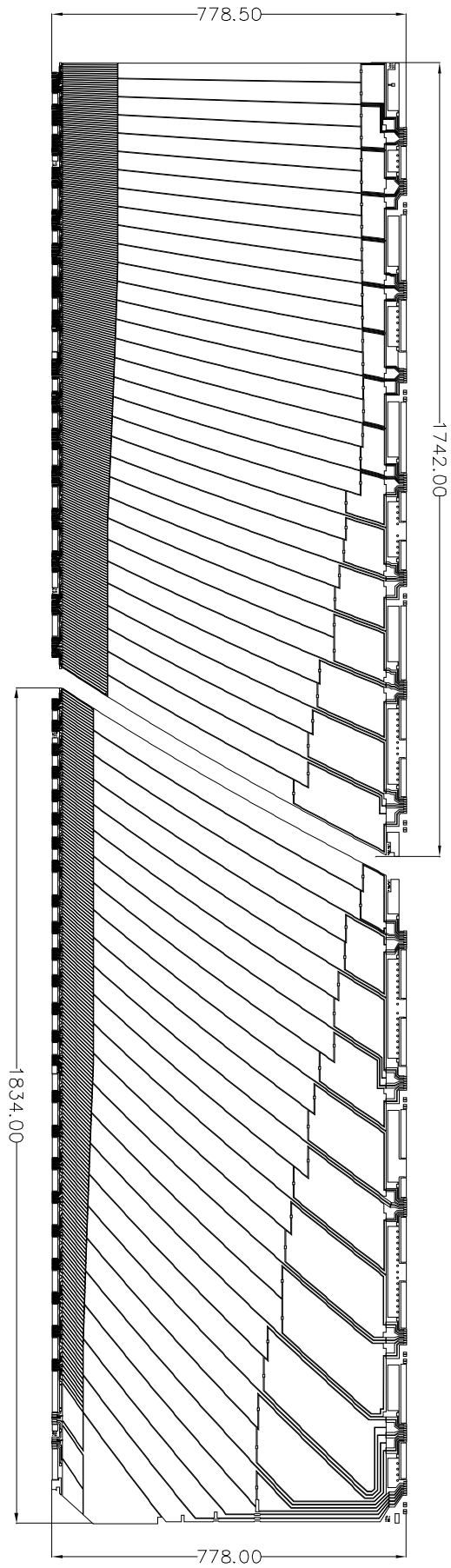


Figure 3.11 — Scheme of the barrel electrodes for the half barrel $0 < \eta < 1.48$, before folding. The beam axis is on the bottom of the plot, with increasing η values toward the right. The change in electrode length at $\eta = 0.8$ is due to a change in the lead absorber sheet thickness at this point (1.53 mm to 1.13 mm). The first and second layer electrode depth slightly decreases as a function of η to keep a constant X_0 . The third layer electrode depth increases as a function of η as the number of X_0 is kept constant as a function of η for the first and second layers, but the calorimeter has a fixed depth as a function of z . [97]

Table 3.3 – Summary of the readout granularity and number of readout channels, the nominal high-voltage settings, and building component parameters in the LAr electromagnetic calorimeter (barrel and endcap). The barrel presampler extends up to $\eta = 1.52$ instead of $\eta = 1.475$ for the accordion. The readout granularity is expressed in $\Delta\eta \times \Delta\phi$ units. The number of channel is for one side only. The total number of channel in the ECal is therefore twice the values given here. The number of calibration channels is not included. Despite having 448 cells for L1 in the region $0 < \eta < 1.4$, the cell closest to $\eta = 0$ is not read leading to only 447 readout channels. The reported high-voltage values are the nominal ones, not the actual settings which are tuned depending on in-situ conditions. The reported high-voltage for the presampler (2000 V) is the Run 1 nominal value. In Run 2 the nominal value was lowered to (1200 V). The high-voltage always supplies zones extending 0.2 in $\Delta\phi$. The LAr gap width and absorber (lead) thickness are reported for the accordion only: the presampler has no absorber and the LAr gap width is constant at about 1.9 mm to 2.0 mm. A layer not present in some η range is indicated with “-”. Data is compiled from Ref. [97]. Empty cells mean the information was not found in that reference.

	Barrel					Endcap outer wheel										EC inner wheel				
η coverage	0	0.8	1.35	1.40	1.475	1.375	1.425	1.5	1.6	1.8	2.0	2.1	2.3	2.4	2.5	2.5	2.8	3.2		
Presampler (PS)			0.025 $\times$ 0.1				-		0.025 $\times$ 0.1			-					-			
Front (L1)		0.025/8 $\times$ 0.1		0.025 $\times$ 0.025	0.025 $\times$ 0.025	0.050 $\times$ 0.1	0.025 $\times$ 0.1	0.025/8 $\times$ 0.1	0.025/6 $\times$ 0.1	0.025/4 $\times$ 0.025	0.025 $\times$ 0.1	0.025 $\times$ 0.1	0.025 $\times$ 0.1	0.025 $\times$ 0.1	0.1 $\times$ 0.1	0.1 $\times$ 0.1	0.1 $\times$ 0.1	0.1 $\times$ 0.1		
Middle (L2)		0.025 $\times$ 0.025		0.075 $\times$ 0.025	0.075 $\times$ 0.025	0.050 $\times$ 0.025		0.025 $\times$ 0.025		0.025 $\times$ 0.025		0.025 $\times$ 0.025		0.025 $\times$ 0.025		0.1 $\times$ 0.1	0.1 $\times$ 0.1			
Back (L3)		0.050 $\times$ 0.025		-	-	-	-	-	0.050 $\times$ 0.025		0.050 $\times$ 0.025		0.050 $\times$ 0.025		0.050 $\times$ 0.025		-	-		
Total PS channels		61 $\times$ 64 = 3904					-	12 $\times$ 64 = 768		-		-		-		-		-		
Total L1 channels		(448 - 1) $\times$ 64 = 28 608		3 $\times$ 256 = 768		216 $\times$ 64 = 13 824													7 $\times$ 64 = 448	
Total L2 channels		57 $\times$ 256 = 14 592				44 $\times$ 256 = 11 264													7 $\times$ 64 = 448	
Total L3 channels		27 $\times$ 256 = 6912		-		20 $\times$ 256 = 5120													-	
η coverage	0	0.8	1.35	1.40	1.475	1.375	1.425	1.5	1.6	1.8	2.0	2.1	2.3	2.4	2.5	2.5	2.8	3.2		
PS HV [V]		+2000 (+1200)					-		-2000 (-1200)		-		-		-		-			
L1+L2+L3 HV [V]		+2000				+2500			+2300	+2100	+1700	+1500	+1250	+1000		+2300	+1800			
LAr gap width [mm]		2.1				2.8 - 0.9													3.1 - 1.8	
Lead thickness [mm]	1.53	1.13				1.7													2.2	
Total radiation length [X ₀]	22 - 30	24 - 33		24 - 38															26 - 36	

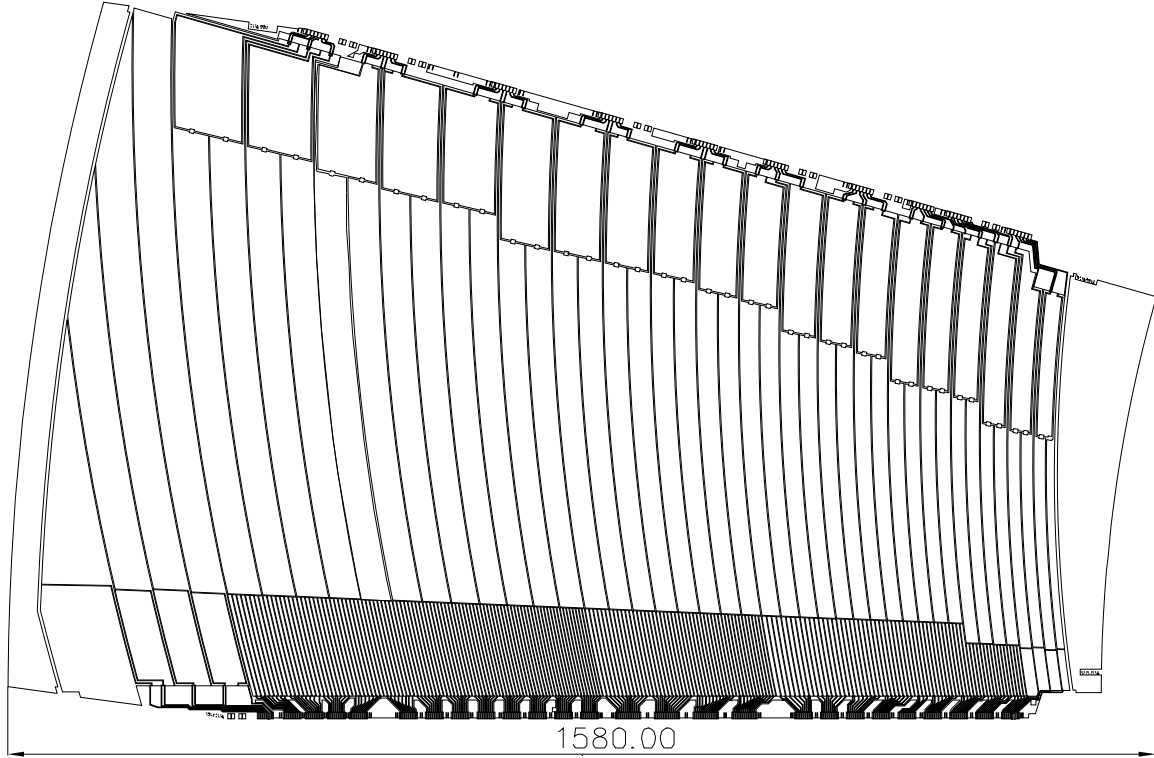


Figure 3.12 – Scheme of the ECal electrodes in the endcap outer wheel ($1.375 < \eta < 2.5$), before folding. The beam pipe is on the right with increasing values of $|z|$ going upward on the figure. The coordinate $\eta = 1.375$ is on the left, and $\eta = 2.5$ is on the right of the plot. [97]

feasible and would moreover induce large cross-talks in the L1 signal. As a consequence, the L2 readout is made in the back of the calorimeter. Such small electrode copper patches driving the signal from the middle cells to the back end can be seen in Figures 3.11 and 3.12 getting around the L3 electrodes.

Since the electrode strips are close to each other, cross-talk effects from one cell to its neighbour cannot be neglected. The largest effect is expectedly seen in the strips, where the first neighbour cross-talk amounts to $\sim 4.3\%$. Subdominant effects come from the S2 neighbour cross-talk below 1%, and from the S2 to S3 cross-talk at about 0.7%.

3.4.4 Central hadronic calorimeter

The barrel hadronic calorimeter (HCal) is also a sampling calorimeter but uses steel as an absorber, and scintillating plastic tiles as active medium (polystyrene). It is made of three parts: one central barrel (5.8 m long) covering the region $|\eta| < 1.0$ and one *extended* barrel part on each side (each 2.6 m long) covering the region $0.8 < |\eta| < 1.7$ (Figure 3.14a). It is located behind the ECal, at radii comprised between 2.28 m to 4.25 m. This extension amounts to a radial length of 7.4 interaction length⁸ (λ).

The HCal is composed of a succession of tiles aligned perpendicularly to the beam axis and following a rotating symmetry around ϕ (Figure 3.14b). An incoming particle in a plastic tile ionises the polystyrene which produces an ultraviolet scintillation light. This light is then guided by a wavelength shifting optic fibre to a photo-multiplier subsequently read out. The two sides of a tile (in ϕ) are read independently by two fibres leading to two independent photo-multipliers,

⁸↑ The interaction length is the equivalent of the radiation length for hadronic interactions. It is denoted λ .

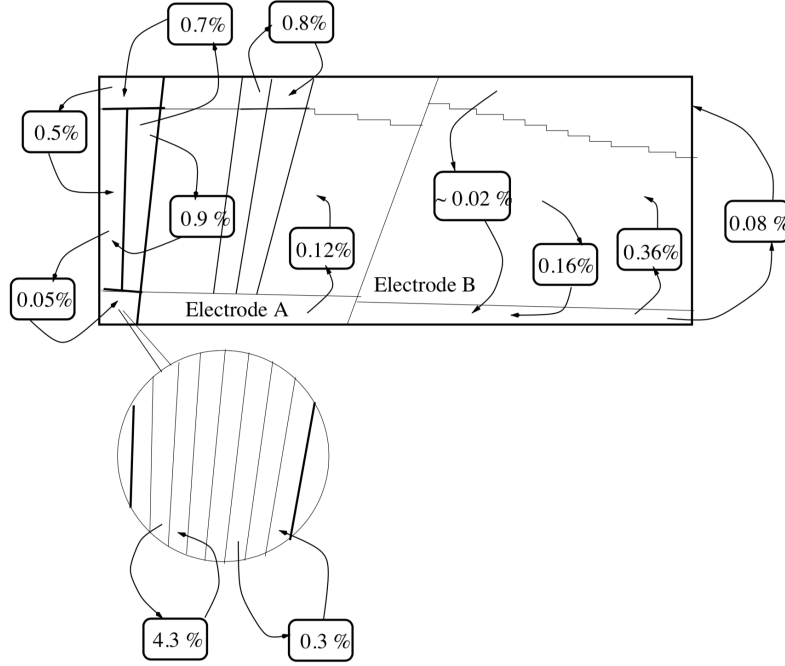


Figure 3.13 – Cross-talk results from the test beam measurements of the LAr barrel module 0 (using the $t_{\max}$ method). [119]

ensuring redundancy in case one fibre is out-of-operation. Each fibre collects the light from several tiles at one or two consecutive radii. The fibres are then grouped together to form a three layers in the r direction of thickness 1.5 , 4.1 and 1.8λ (at $\eta = 0$). The grouping is also made to produce cells of 0.1×0.1 in $\Delta\eta \times \Delta\phi$ in the first and second layer, and 0.1×0.2 in the third layer. This leads to approximately 10 000 channels for the tile hadronic calorimeter.

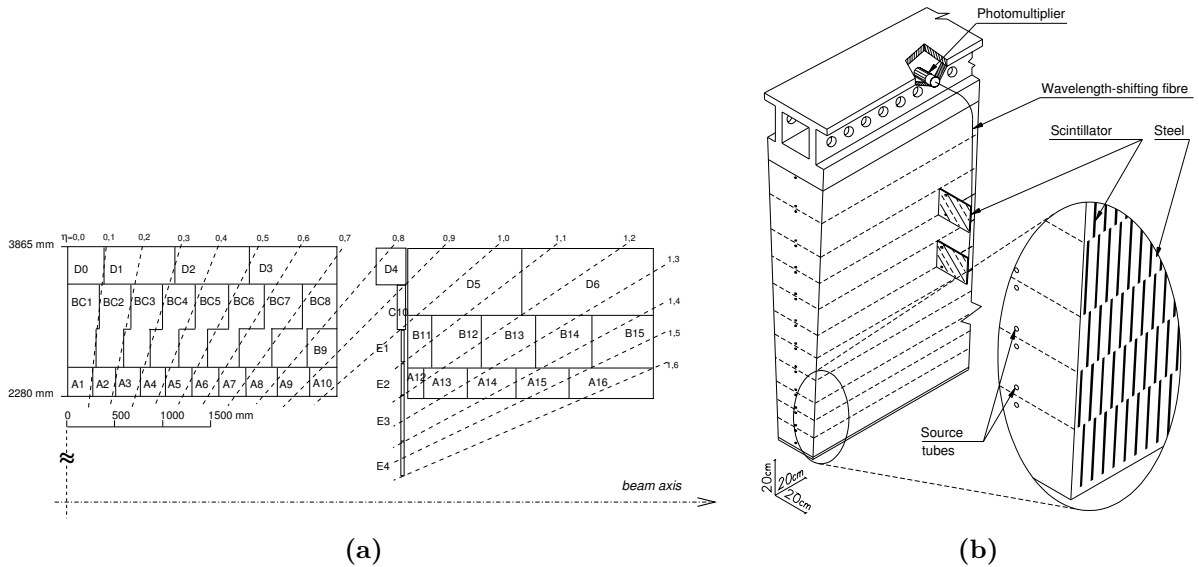


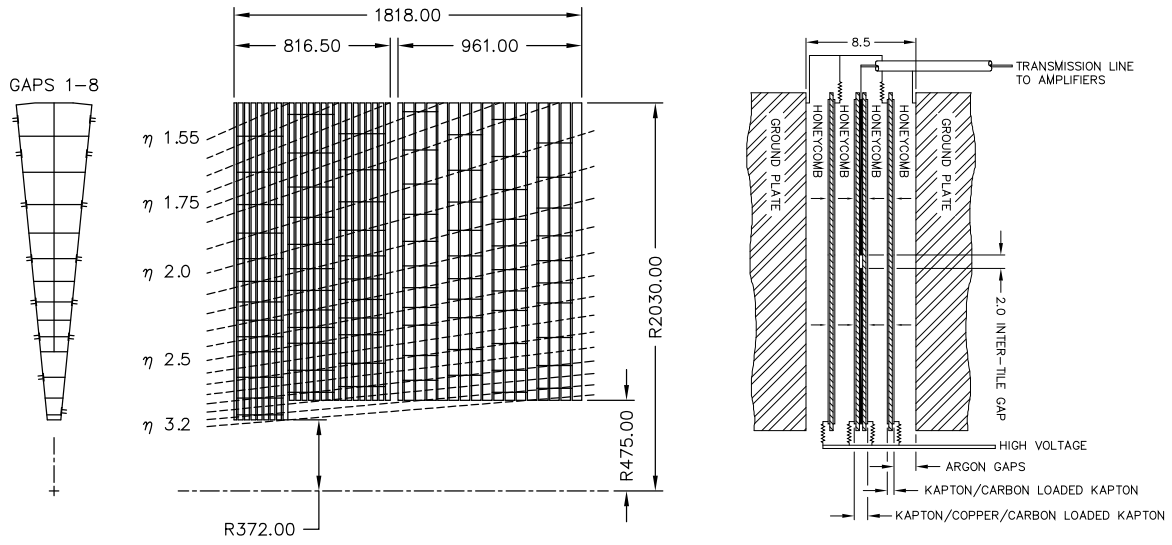
Figure 3.14 – Layout of the tile calorimeter (a), and (b) scheme of HCal module. [97]

3.4.5 The hadronic endcap and forward calorimeters

In the endcap, the expected luminosity does not allow the hadronic calorimeter to use scintillating tiles as it would not withstand high flux of particles. Instead the hadronic endcap calorimeter (HEC) relies on the radiation-hard, granular liquid argon calorimeter, with a more conventional

parallel-plate copper–LAr sampling calorimeter. As represented in Figure 3.15a, it is composed of two wheels (front and rear) with a diameter of 4 m on each side, covering the region $1.5 < |\eta| < 3.2$.

The front wheel includes 24 copper disks each 25 mm thick (plus one 12.5 mm in front), while the rear wheel includes 16 copper disks each 50 mm thick (plus one 25 mm in front). A 8.5 mm LAr gap between all plates is ensured by adding a spacer between the copper plates. The electrode layout splits the 8.5 mm space in four 1.8 mm drift gaps, as displayed in Figure 3.15b. The three electrodes forming the four gaps are each fed with a HV of nominal value of 1800 V, and are etched to provide a granularity of 0.1×0.1 in $\Delta\eta \times \Delta\phi$ for $1.5 < |\eta| < 2.5$, and 0.2×0.2 for $2.5 < |\eta| < 3.2$. Each wheel is further read in two layers in z , offering ~ 2800 readout channel on each side. These data are summarised in Table 3.4. More technical details can be found in Ref. [115].



(a) Schematic R - ϕ (left) and R - z (right) views of the HEC. (b) Schematic of the LAr gap in the HEC.

Figure 3.15 — Schematic of (a) the hadronic endcap calorimeter general layout as well as semi-pointing layout of electrodes (dashed lines), and (b) of its electrode and LAr gap layout. All units are in mm. [97]

Finally, the forward calorimeter (FCal) covers the region $3.1 < |\eta| < 4.9$, ensuring an (almost) 4π hermeticity of the calorimeter systems. It has three 45 cm deep modules, the first one providing with electromagnetic calorimetry and the two other with hadronic calorimetry. It also uses liquid argon active material, with copper absorbers in the electromagnetic part and tungsten absorbers in the hadronic part.

Each FCal module is built with two absorber plates on each side, placed perpendicularly to the beam pipe, with cylindrical rods along the z axis which are partially filled with anode rods (Figure 3.16). Each rod is surrounded by a LAr gap and a concentric absorber layer. The LAr gaps are narrow to avoid ion build up in the high intensity forward region. The signal are read in the front of the calorimeter for the electromagnetic layer and in the back for the hadronic layers. The main parameters of the FCal are summarised in Table 3.4. More technical details can be found in Ref. [116].

Table 3.4 – Summary of the readout granularity and number of readout channels, the nominal high-voltage settings, and building component parameters for the tile hadronic central calorimeter, the LAr hadronic endcap calorimeter and the LAr forward calorimeters. The readout granularity is expressed in $\Delta\eta \times \Delta\phi$ units. The number of channel is for one side only. The number of calibration channels is not included. The reported high-voltage values are the nominal ones, not the actual settings which are tuned depending on in-situ conditions. A layer not present in some η range is indicated with “–”. Data is compiled from Refs. [115], [116] and [97]. Empty cells mean the information was not found in these references.

Module	Tile barrel	Tile extended	HEC front			HEC rear		FCal1 (EM)	FCal2 (Had)	FCal3 (Had)
Absorber	Steel	Steel	Copper			Copper		Copper	Tungsten	Tungsten
η coverage	0–1.0	0.8–1.7	1.5–2.5	2.5–3.2	1.65–2.5	2.5–3.2			3.1–4.9	
Layer 1	0.1×0.1	0.1×0.1	0.1×0.1	0.2×0.2	–	–	$\sim 0.1 \times 0.1$	–	–	–
Layer 2	0.1×0.1	0.1×0.1	0.1×0.1	0.2×0.2	–	–	–	$\sim 0.2 \times 0.2$	–	–
Layer 3	0.2×0.1	0.2×0.1	–	–	0.1×0.1	0.2×0.2	–	–	–	$\sim 0.2 \times 0.2$
Layer 4	–	–	–	–	0.1×0.1	0.2×0.2	–	–	–	–
Total channels per side	~ 5000		$768 + 736 = 1504$			$672 + 640 = 1312$		1008	500	254
LAr HV [V]	not applicable		+1800					+250	+375	+500
LAr gap width [mm]	not applicable		4×1.8					0.269	0.376	0.508
Absorber thickness [mm]	not applicable		25			50				
Total absorption length [λ]	> 7.4							2.66	3.68	3.60

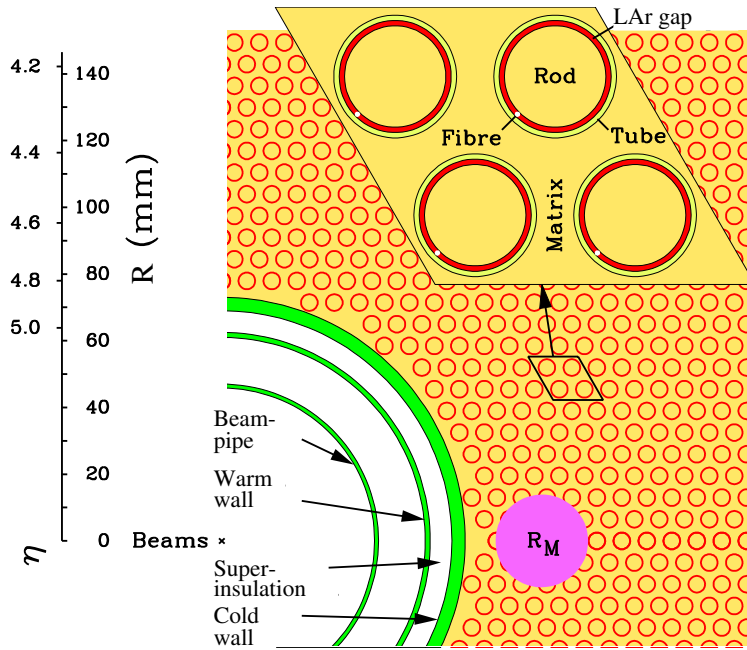


Figure 3.16 – Schematic of the electrode structure in the first layer of the forward calorimeter (FCal1). The Moliere radius R_M is represented by a solid disk. [97]

3.5 The muon spectrometer and the toroid magnets

The outermost layers of the ATLAS detector complement the particle identification and trigger capabilities by a system dedicated to the detection of muons. Indeed, the combined tracker and calorimeter information is not sufficient to identify these particles: the energy deposit of muons in the calorimeters is barely above the noise level (a few hundreds of MeV) and muons escape beyond the depth of the calorimeters. Since muons and charged pions have similar masses ($m_\mu \approx 105 \text{ MeV}$ while $m_{\pi^\pm} \approx 135 \text{ MeV}$), the tracker cannot distinguish between the two. The ATLAS muon spectrometer (MS) therefore uses a tracker-like system providing *hits* and embedded in a toroidal magnetic field to measure the position and momentum of muons.

Given the area needed to be covered, using silicon sensors cannot be considered, and the muon detectors instead rely on largely scalable gaseous detector. They are arranged in *stations* such that at least three stations are crossed by the muon track. Multiple technologies are used, responding to the different requirements as a function of $|\eta|$ in the detector: Monitored Drift Tubes (MDT) make up for most of the precision measurement chambers, covering about 5500 m^2 , and are supplemented by smaller Cathode Strip Chambers (CSC) in the forward region of the innermost station. The faster trigger chambers are composed of Resistive Plate Chambers (RPC) in the barrel, and of Thin Gap Chambers (TGC) in the endcaps. The characteristics of the various technologies used are briefly described below and summarised in Table 3.5.

Table 3.5 – Summary of the parameters for each of the four types of muon chambers. [97]

Type	Station resolution (RMS)			Hits / track		# chambers	# channels
	z/R [mm]	ϕ [mm]	Time [ns]	Barrel	Endcap		
MDT	0.035	–	–	20	20	1150	354k
CSC	0.040 (R)	5	7	–	4	32	30.7k
RPC	10 (z)	10	1.5	6	–	606	373k
TGC	2 – 6 (R)	3 – 7	4	–	9	3588	318k

Monitored Drift Tubes (MDT): the basic element constituting a MDT is a 1 m to 2 m long, 30 mm in diameter tube surrounding an anode wire at ~ 3080 V in its centre, and filled with a gas mixture (Ar:CO₂ in proportion 93:7) that gets ionised when traversed by a charged particle. The measurement of the drift time offers a resolution of 80 μ m per tube. However this drift time is too slow (up to ~ 700 ns) to provide an efficient trigger capability, given the 25 ns collision period. The tubes are gathered into groups of three (middle and outer stations) or four (inner stations) layers to form a chamber and each station is made of a pair of such chambers (see Tables 6.3 and 6.4 of Ref. [97]), achieving a resolution of ~ 35 μ m in the bending plane (η). However the MDT do not provide information on the track ϕ coordinate.

Cathode Strip Chambers (CSC): due to higher radiation level (beam background) in the region $2.0 < |\eta| < 2.7$, the CSC replace the MDT in the innermost station for this region. Each CSC is a multi-wire proportional chamber (MWPC), with anode wire sandwiched between two cathode planes. The ~ 2.54 mm gap in between the two planes is filled with a Ar:CO₂ mixture (in proportion 80:20) and operated under ~ 1900 V: similarly to the MDT, a passing muon ionises the gas mixture, creating an avalanche effect due to the node HV, and ionisation signal is recorded by the strips. The cathodes are segmented in ~ 1.5 mm strips (one with strips perpendicular to the anode wires, the other with strips parallel to the wires), offering a 60 μ m spatial resolution in the bending direction (η , 5 mm in the non-bending direction ϕ). Timing wise, the electron drift time is about 40 ns, yielding a ~ 7 ns resolution per CSC plane. Similarly to the MDT, the CSC stations are made of 4 CSC planes, but are tilted by 11.59° , slightly leaning toward the interaction point, to ensure that tracks impact the planes perpendicularly.

Since the main MDT chambers are too slow to provide trigger capabilities, they must be supplemented by additional detector, less precise but with a faster response.

Resistive Plate Chambers (RPC): the RPC provide triggering capabilities in the barrel ($|\eta| < 1.05$), and complement the missing ϕ information from the MDT in this region. Contrary to the previous detectors, it has no wire and is built from two resistive plates separated by an gaseous insulator medium (C₂H₂F₄:Iso-C₄H₁₀:SF₆ in proportions 94.7:5:0.3) of 2 mm and operated at 9.8 kV, allowing for an avalanche to form at a muon passing. One RPC unit is made of two such gaseous gaps, surrounded by four electrode planes. The signal is read by inductive coupling on conducting strips glued onto the surface of the electrodes. The strips on each of the two plates are oriented perpendicularly to allow for both η and ϕ measurements. The timing resolution is about 5 ns, allowing for a fast trigger decision.

Thin Gap Chambers (TGC): the purpose of the TGC is the same as for the RPC, *i.e.* providing trigger capability and a ϕ coordinate measurement, but for the endcap region $1.1 < |\eta| < 2.4$. The TGC are also multi-wire proportional chambers, with anode wires separated by 1.8 mm and sandwiched between two graphite cathodes at 1.4 mm from the wires. The gap is filled with a highly quenching gas mixture (CO₂:n-C₅H₁₂ in proportions 55:45), allowing for a nominal 2900 V potential to be applied on the wires. The signal is read from the wire and additional strips placed at 1.6 mm from the cathode, orthogonal to the wires. The TGC are assembled in doublets or triplets, which in both case have two strips layers in total (see Figure 6.32 from Ref. [97]).

The layout of the muon detection system with all these elements is displayed in Figure 3.17 for the barrel and in Figure 3.18 in the endcaps. In the barrel, the three stations are installed concentrically and are located at approximate radii of 5 m, 7.5 m and 10 m, *i.e.* one before, one in the middle and one after the toroid magnets in the R direction. The first station is composed of a MDT quadruplet, the second of a MDT triplet surrounded by one RPC unit on each side, and the third by a MDT triplet with a RPC unit in front. This repartition of RPC allows for a

“low- p_T ” (6 GeV to 9 GeV) trigger comparing hits in the first two layers, and a “high- p_T ” (9 GeV to 36 GeV) trigger comparing the two first hits to the expected third hit.

In the endcap, three main MDT stations are installed perpendicularly to the beam pipe and are located at approximate $|z|$ of 7.4 m, 14 m and 21.5 m (one before and two after the endcap toroid vessel). An extra layer at $|z| \approx 10.8$ m is installed to ensure three stations are crossed in the barrel/endcap transition region. The first station is complemented with a TGC doublet in front, while the second station has one TGC triplet in front (reducing false coincidences from background hits) and two TGC doublet behind. In this second station, the trigger information is built from coincidences from two-out-of-three matches in the triplet or three-out-of-four matches in the two doublets.

The modules are arranged in 16 sectors around ϕ , 8 “large” ones placed in between two toroid coils, and 8 “small” ones placed inside the toroids at slightly different radii in the barrel, and slightly different z in the endcap. An overlap is maintained between the large and small modules to ensure a hermetic coverage in the ϕ direction. However, the region $|\eta| \lesssim 0.1$ is occupied by service cables for the calorimeter and the solenoid and is not instrumented with a muon detector. The muon measurement in this region therefore only relies on the tracker and on the calorimeter information.

The muon spectrometer is designed to provide muon momentum measurement from 3 GeV to several TeV (below this range, the energy loss in the calorimeter prevents the muon from going beyond), with a standalone resolution of 3 % for 100 GeV muons and 10 % for 1 TeV muon tracks. This translates into a requirement on the spatial resolution of 50 μm , requiring the knowledge of the precision station location to better than 30 μm , and a need for a magnetic field to bend the muon trajectories.

Eight toroids, symmetrically positioned around ϕ , produce a 0.5 T magnetic field in the barrel ($|\eta| < 1.4$), while eight additional smaller toroids on each side produce a 1 T magnetic field in the endcaps ($1.6 < |\eta| < 2.7$). The eight barrel and endcap toroids are relatively rotated by 22.5° , as seen in Figure 3.19a. In the transition region, the magnetic field map gets more complex as it is the superposition of the barrel and endcap fields. In these areas, the field integral is moreover close to 0, meaning that muon do not get deflected, thus degrading the momentum resolution (Figure 3.19b).

To produce such high fields, a current of 20.5 kA is needed so the magnets are also superconducting, using a NbTi winding that needs to be cooled to ~ 4.5 K. Each of the barrel toroids is 25.3 m long, spans the radii between 4.7 m to 10 m and is enclosed in its own cryostat. On the other hand, the eight endcap cryostats are all embedded in the same vessel, preventing constraining the muon stations to be placed before or after the vessel.

3.6 Luminosity monitoring

While the luminosity is theoretically known from the beam parameters (see Eq. 2.4), its precise measurement is paramount for most analyses as it directly impacts the number of expected events and thus all cross-section measurements (Eq. 2.3). The main detector dedicated to the luminosity measurement in ATLAS is the LUCID (LUMinosity measurement using Cherenkov Integrating Detector) apparatus [120], located at $z = \pm 17$ m from the interaction point. It performs a relative luminosity measurement, detecting inelastic p - p scattering in the forward region. The detector is designed as 16 1.5 m long tubes surrounding the beam line at ~ 10 cm (*i.e.* $|\eta| \approx 5.8$), and filled with a C_4F_{10} gas. This provides a Cherenkov medium with a threshold at 10 MeV for electrons and 2.8 GeV for pions. The Cherenkov light is detected by PMTs, whose signal amplitude can be linked to the number of particles crossing the tube. It is then assumed that the particle count is

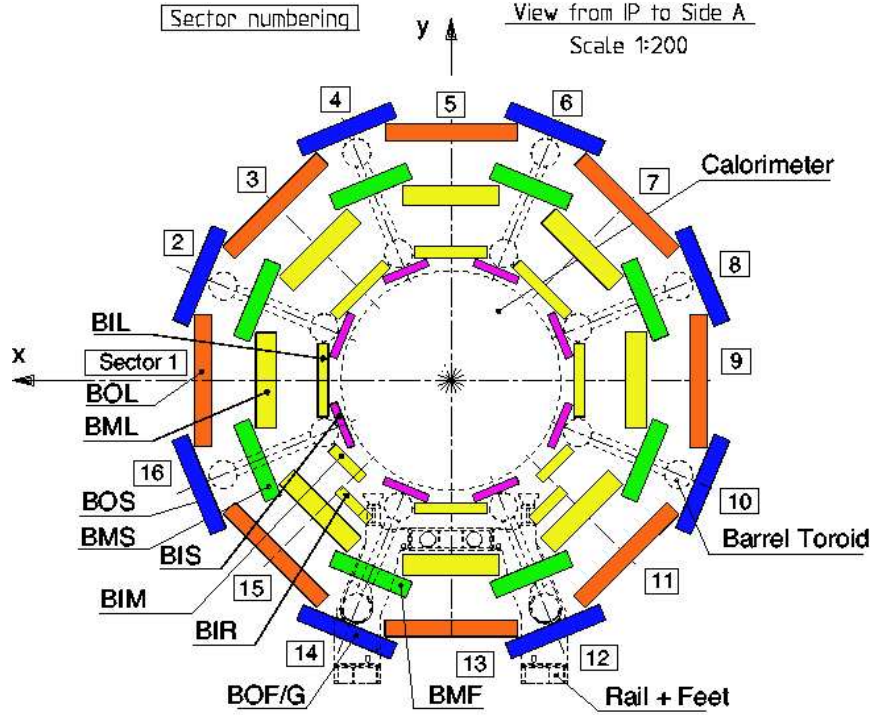


Figure 3.17 – Cross section of the muon system in the barrel seen in the x - y plane. The three layers (Inner, Middle and Outer) of large (yellow and orange) and small (purple, green and blue) stations are visible. The inner stations are composed of three stacked MDT, the middle stations are composed of four stacked MDT, sandwiched between two double gap RPC, the outer stations are composed of four stacked MDT and a double gap RPC. [97]

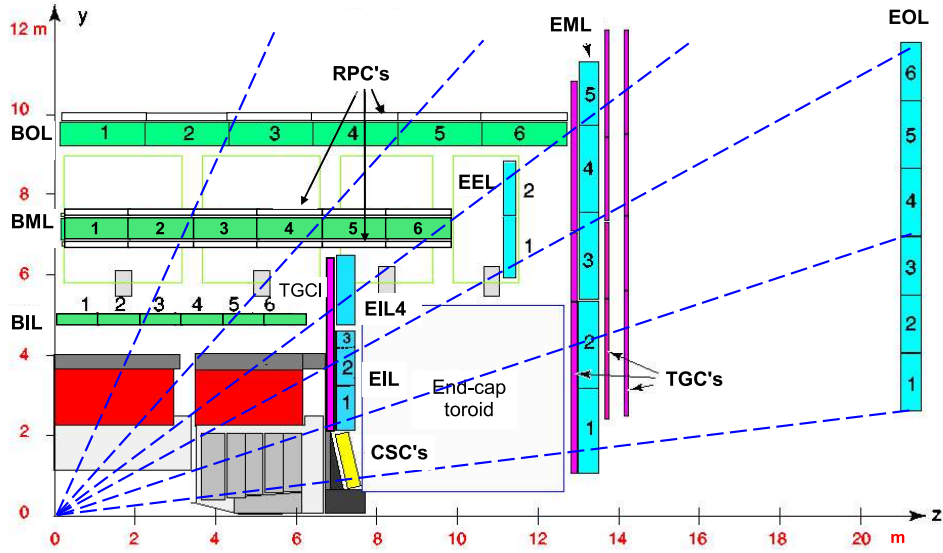


Figure 3.18 – Cross section of a quadrant of the muon system in the endcap ($\eta > 0$) seen in the z - y plane. The MDT stations are represented in green in the barrel and in cyan in the endcap, the RPC with open white boxes, the TGC in purple, and the CSC in yellow. Only the three large stations are drawn. An additional MDT station is installed between the inner and middle ones to ensure 3 stations are crossed in the barrel/endcap transition region. In the centre of the picture, the electromagnetic (grey) and hadronic (red) calorimeters are shown. The six straight line represent various values of η . [97]

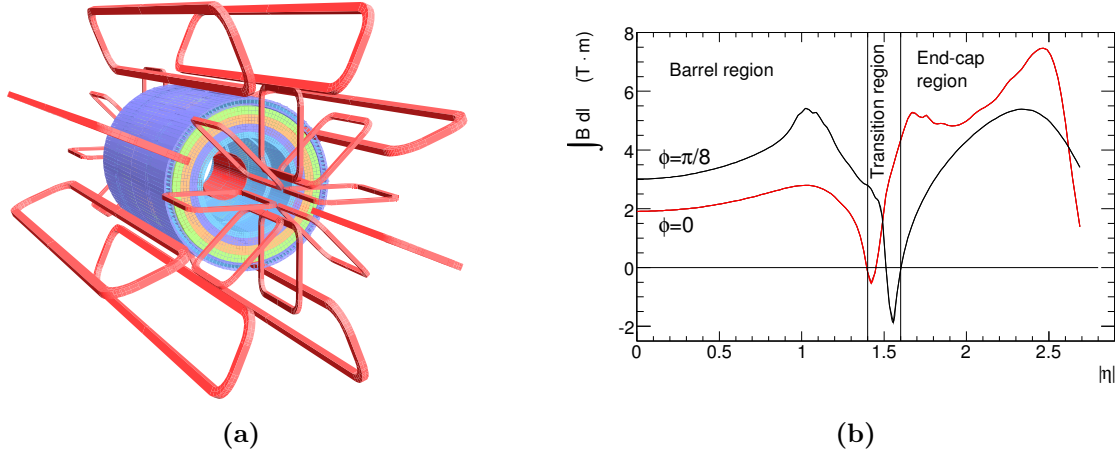


Figure 3.19 – Illustration of the eight barrel and twice eight endcap toroids creating the 0.5 T to 2 T magnetic field in the muon spectrometer. Figure (a) shows the localisation of each toroid with respect to the calorimeter. The integrated field seen by the crossing muons as a function of η is displayed in (b). [97]

proportional to the number of particles created in the collision, *i.e.* proportional to the number of interaction during the bunch crossing, hence the instantaneous luminosity.

The instantaneous luminosity $\mathcal{L}$ relates to the *visible number of interaction per crossing* μ^{vis} by:

$$\mathcal{L} = \frac{\mu^{\text{vis}}}{\sigma^{\text{vis}}} n_b f, \quad (3.6)$$

with n_b and f the number of bunches in the beams and f the revolution frequency of the bunches (see Section 2.8). σ^{vis} is the *visible cross-section* for the measured process, and is calibrated during the so-called *van der Meer (vdM)* [121] scans, with a few scheduled each year. The scans are run with special beam optics, which includes a high β^* and a reduced beam emittance. During these scans, the beams are progressively shifted with respect to each other in the x and y directions, allowing for a measurement of their width: σ_x^{eff} and σ_y^{eff} . For one configuration of the beams, the measured number of interactions per crossing reaches its maximum $\mu_{\text{max}}^{\text{vis}}$, which related to the visible cross-section by:

$$\sigma^{\text{vis}} = \mu_{\text{max}}^{\text{vis}} \frac{2\pi\sigma_x^{\text{eff}}\sigma_y^{\text{eff}}}{N_1 N_2}, \quad (3.7)$$

with N_1 and N_2 the number of protons per bunch (see Section 2.8).

To determine the visible cross-section, the above mentioned LUCID detector uses an algorithm counting the average number of hits per bunch crossing [122–124], integrated over one *lumi block* (~ 60 s): if the combined information from N_{PMT} is used, that N_{HIT} have been recorded during this lumi block, and that this lumi block has seen N_{BC} bunch crossings, the visible cross-section is given by

$$\mu_{\text{vis}} = -\ln\left(1 - \frac{N_{\text{HIT}}}{N_{\text{BC}} N_{\text{PMT}}}\right). \quad (3.8)$$

The integration over a given period of time (lumi block) is required to have enough statistics for each measurement.

Using this method, the total uncertainty on the full Run 2 integrated luminosity reaches 1.7% [124].

3.7 Trigger system

Given the high collision rate at the LHC (40 MHz) and the size of an ATLAS event (~ 1 MiB), recording all the events occurring in the detector would result in writing to disk or tape about 40 TiB/s. Since this is clearly unmanageable, events must be selected in real time (called *online* selection) to achieve a final output rate of 1 kHz to 2 kHz (bandwidth of ~ 2 GiB/s). Moreover, the standard model processes range over fourteen orders of magnitude in terms of cross-section (see Figure 1.14), while the processes of interest at the LHC are rather the rare ones. The selection must therefore be able to recognise event topologies and to reach a fast decision on whether to keep or discard an event.

This selection is performed by a two-stage trigger, which was upgraded for Run 2 [125], presented in Figure 3.20. The first stage, called *L1* (Level-1), is made of custom hardware logic circuits and buffers reaching a decision within $2.5\ \mu\text{s}$, and reduces the event rate from 40 MHz to ~ 100 kHz. It only reads out information from the calorimeter and the muon trigger chambers (RPC and TGC) with a coarse granularity: the calorimeter is read out in region of 0.1×0.1 in η combining all layers, called a calorimeter *trigger tower*. The L1Calo decision is based on a threshold on the total energy from a cluster of 2×2 trigger towers. An isolation criterion can be added by adding a higher bound threshold on the energy summed over the 4×4 surrounding towers. The L1Calo jet triggers also include information from the HCal with same granularity. The L1Muon decision is based on track segments from a coarse reconstruction algorithm. An example L1 rate per group of streams as a function of time in the run is shown in Figure 3.21a.

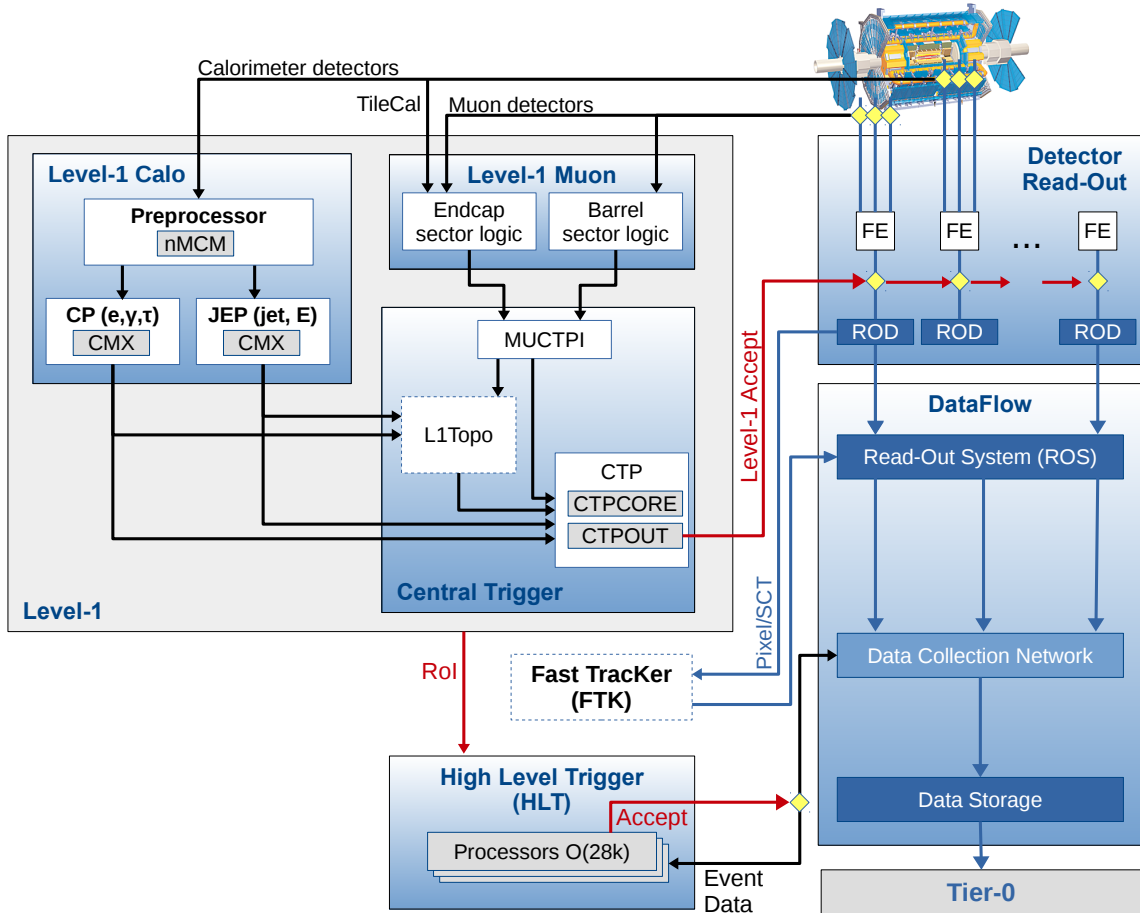


Figure 3.20 – Simplified view of the ATLAS trigger system in Run 2. [125]

The second trigger stage is the High-Level Trigger, which decreases the event rate from 100 kHz to less than 2 kHz. Using a computing farm of $\sim 30\,000$ processors (in 2015, increased during Run 2), the HLT is allowed to retrieve information from the ID tracks as well. It consists in two sub-steps: the first one uses still imprecise but fast reconstruction algorithms, and removes most of the pre-selected events. The second step runs reconstruction algorithms close to the *offline* (*i.e.* analysis) ones, and therefore benefits from a much better energy and track resolution. The event rate out of the HLT is shown in Figure 3.21b. The HLT normally reaches a decision within 1 s, but as seen in Figure 3.22, the time needed to reach a decision can extend largely beyond this limit. Such outliers events are stored in a special *debug* stream which is processed at the end of the run.

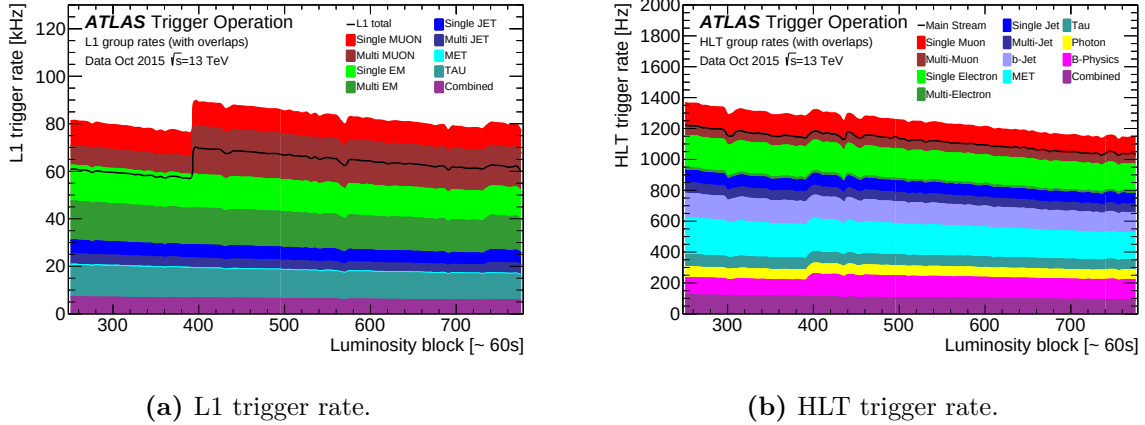


Figure 3.21 – Example of (a) L1 and (b) HLT trigger rates in one of the 2015 runs with a peak luminosity of $4.5 \times 10^{33} \text{ cm}^{-2} \text{ s}^{-1}$. The trigger decisions overlap between the various streams, making the sum higher than the total rate (shown in black line). The step at lumi block ~ 400 is due to the unpre-scaling of the b -physics triggers [125]

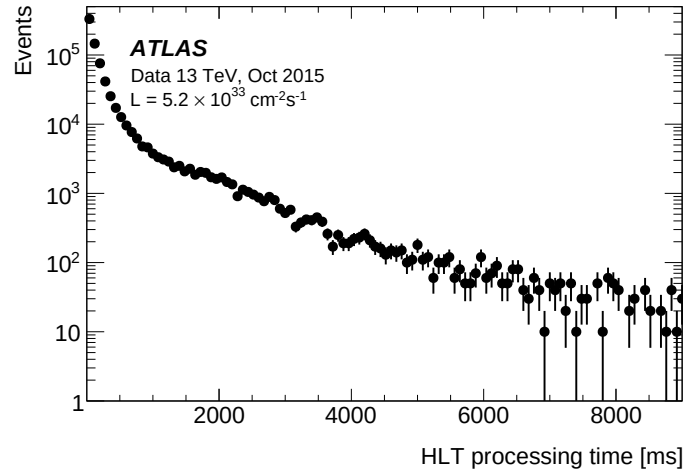


Figure 3.22 – Distribution of HLT processing time per event in one 2015 run with peak instantaneous luminosity of $5.2 \times 10^{33} \text{ cm}^{-2} \text{ s}^{-1}$. [125]

As the run progresses and collisions occur, fewer protons are present in the bunches and the number of collision per bunch crossing decreases, *i.e.* the instantaneous luminosity. The triggers are therefore organised into *menus* for each stream (electron, photon, muon, jet, τ , E_T^{miss} , *etc.*) with different energy thresholds. These menus are included or removed from the trigger system in the course of the runs depending on the instantaneous luminosity in order to use the trigger capability to its maximum bandwidth.

Moreover, some of the low threshold triggers with a too high bandwidth requirements compared to their physics interest can be prescaled. A trigger prescaled by a factor N only randomly accept $1/N$ events that pass its requirements. This prescale value can also be tuned during the run, so that low energy triggers are mainly accepted in the end of the runs.

Reconstruction and identification of physics objects

Contents

4.1	Reconstruction of tracks in the ID	89
4.2	Reconstruction of tracks in the MS	90
4.3	Reconstruction of energy in calorimeter cells	90
4.4	Electrons and photons	92
4.5	Muons	95
4.6	Jets and b -tagging	98
4.7	Missing transverse energy	99

Once the trigger has reached a positive decision for the current event, electric signals are sent out of each sub-detector described in Chapter 3 to be permanently written out, saved on both disk and tape storage. From these raw data corresponding to the binary information of the tracker elements and ADC counts for each calorimeter cell, algorithms are run to reconstruct complete tracks and calorimeter cell clusters that allow to identify and measure physics object properties. The final step is the calibration of the energy, track parameter, momentum, and other primary quantities of such objects. The electron calibration is generically detailed in Chapter 5, and one of the steps which has been studied during this Thesis is described in Chapter 6.

4.1 Reconstruction of tracks in the ID

Tracks are reconstructed from a list of hits in the ID and the MS. Given the high multiplicity of tracks in each collision and consequently the high number of individual hits, building an efficient track finding algorithm is one of the biggest challenges at the LHC, requiring a large amount of computing resources. Once the tracks are reconstructed, vertices are found by extrapolating tracks to the beam location and by associating tracks to each vertex. The two relevant quantities for this association are the transverse and longitudinal impact parameters with respect to the vertex, denoted d_0 and z_0 , respectively.

4.2 Reconstruction of tracks in the MS

In the muon spectrometer, track segments are reconstructed in each station with a linear fit on hits aligned in the bending plane of the toroid field. The MDT are used to determine the coordinate in the bending plane, while the RPC and TGC are used for the coordinate in the orthogonal direction. The full MS segments are seeded in the middle station and extrapolated to the inner and outer stations, looking for matching segments. If at least one of the inner or outer stations has a matching segment (except in the crack region where only one segment is required), the track is considered valid. In case of tracks sharing the same segments, an overlap removal procedure is applied.

4.3 Reconstruction of energy in calorimeter cells

Contrary to tracker hit signals which consist of binary information, the calorimeter cell readout sends an analogue signal formed by an electrical pulse, as shown in Figure 3.8. The signal are then digitised with three analogue-to-digital converters (ADC) using three different gain values: 1, 9.9, and 93 [126] hence named low, medium, and high gain, respectively. Using three gains allows to cover the whole energy range for each cell while keeping a reasonable energy resolution for low energy values. The three gain ranges slightly overlap to avoid non-linear effects in the first ADC counts. Only one gain is sent further, roughly based on the higher gain not saturating the ADC. For ~ 40 GeV electrons, cells in the electromagnetic cluster are mostly read in high gain, while a significant fraction of the cells in a 60 GeV photon cluster are read in medium gain.

The signal shape is sampled every 25 ns, providing approximately 32 points for a complete ionisation cycle. However, due to bandwidth limitations, only a limited number such samples are sent further to the data acquisition system. During Run 1, five of these samples around the peak position of the pulse were read out. Due to the expected higher occupancy in the calorimeter and in order to increase the trigger rate with a constant bandwidth for Run 2, this number is decreased to four samples since 2015.

To recover the full ionisation signal from these four points only, signal shape templates are parametrised for each gain as a function of the amplitude value at each of these points, and the peak amplitude is determined using an optimal filtering method [127]:

$$A = \sum_{j=1}^{N_{\text{samples}}} a_j(s_j - p), \quad (4.1)$$

where s_j represent the amplitude height of each sample in ADC counts, p is the pedestal value determined in regular calibration runs, and the a_j are the optimal filtering coefficients (OFC). Once the shape is sampled, the energy in the cell is reconstructed as [126]

$$E = F_{\mu\text{A} \rightarrow \text{MeV}} \times F_{\text{DAC} \rightarrow \mu\text{A}} \times \frac{1}{\frac{M_{\text{phys}}}{M_{\text{calib}}}} \times G_{\text{ADC} \rightarrow \text{DAC}} \times A, \quad (4.2)$$

where

- A is the amplitude value from the optimal filter method described above, in ADC counts;
- $G_{\text{ADC} \rightarrow \text{DAC}}$ is the gain factor, converting the ADC counts with a digital-to-analogue converter (DAC), in DAC/ADC;
- $M_{\text{phys}}/M_{\text{calib}}$ corrects for the different pulse shapes used for the calibration signal and the physics signals;

- $F_{\text{DAC} \rightarrow \mu\text{A}}$ relates the DAC setting of the electronic board to the current and is determined in regular (typically daily) calibration runs;
- and $F_{\mu\text{A} \rightarrow \text{MeV}}$ relates the incoming particle energy to the ionisation current produced and has been measured during initial test beams [128].

Each of these constants is independently measured for each cell and each gain, allowing for an online correction in case of non nominal HV value for example. The full data acquisition chain from the ionisation signal to reconstructed energy is displayed in Figure 4.1.

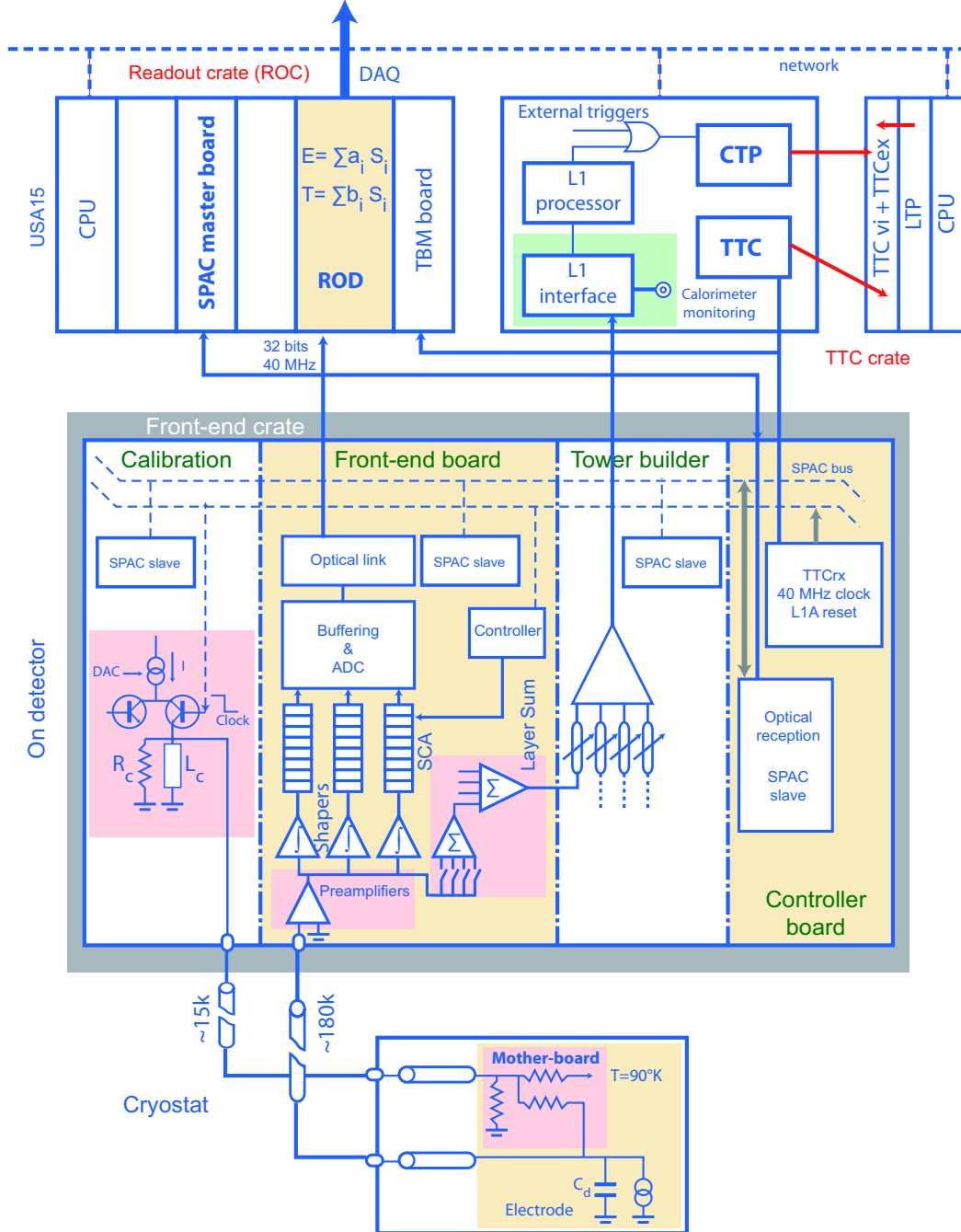


Figure 4.1 – Data acquisition chain for the ATLAS LAr calorimeters. [97]

As seen in Chapter 3.4.3, the signal returned by the calorimeter electronics chain has a bipolar shape. As a consequence, its integral over the pulse duration (about 450 ns) should vanish. Due to that, the contribution of pileup events to the cluster energy should average out to 0.

However, this reasoning is only valid in case of regularly spaced bunches, and if the bunches have all the exact same number of interactions. In practice, the number of interactions per bunch crossing follows a Poisson distribution. Moreover as described in Chapter 2.7, all LHC bunch slots are not filled with protons. For example, the typical LHC filling scheme for 2016 used trains composed of 2 batches separated by 9 empty bunches, with batches containing 48 bunches each. Trains were separated by at least 36 empty bunches. The gap between the two batches therefore is 225 ns which is much smaller than the 450 ns needed by the liquid argon electronic to complete its pulse, leading to non-vanishing energy of pileup events at the beginning of the second batch in the train.

In order to avoid large electronic bias induced by pileup contributions, a correction is therefore derived to account for these energy shifts. Such correction depends on the filling scheme in use at the LHC and of the bunch luminosity. Examples of correction for the standard and the 8b4e (many batches of eight bunches separated by four empty bunches used in the second half of 2017, see Section 2.7) filling schemes are shown in Figure 4.2.

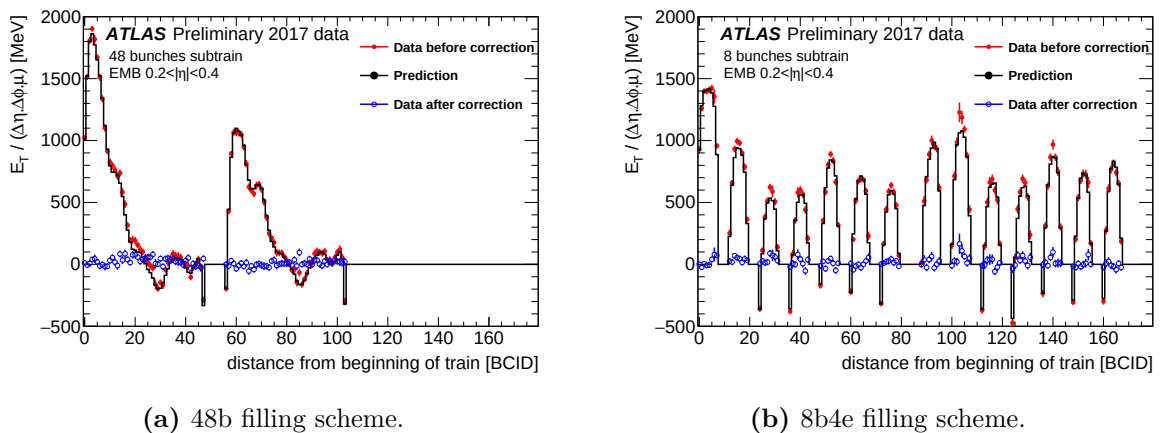


Figure 4.2 – Average transverse energy in zero-bias events per unit $\eta \times \phi$ area and unit number of interactions per bunch crossing as a function of the distance from the beginning of the bunch train (measured in BCID so in 25 ns units), using (a) the standard 48b-9e-48b (or similar, used in most of the Run 2), and (b) the 8b4e filling scheme as used in the second half of the 2017 run. See Chapter 2.7 for details on the LHC filling scheme. [129]

4.4 Electrons and photons

The procedure to reconstruct electrons and photons is similar and starts by building a cluster from cells in the ECal. The main difference between photons and electrons is the absence or presence of a track associated to the cluster. The exact procedure for Run 2 is detailed in Refs. [130] and [131], and has recently been updated in Ref. [132]. It should be noted that the clustering algorithm has changed from using fixed-size clusters (looking for the maximum energy set of cells by sliding a fixed-size window across the calorimeter) in Run 1 to dynamical cluster (briefly described below) for the full Run 2 results¹.

The first step in the electron and photon reconstruction is to build clusters of cells, called *topo-clusters*. The clustering algorithm [133] starts from any calorimeter cell, excluding the PS and first ECal layers to avoid the formation of noise clusters, satisfying a cell energy significance

¹↑ Early partial Run 2 results typically using the 2015 and 2016 dataset only, such as most of the calibration results presented in Chapter 5, used the same “sliding-window” algorithm as in Run 1. This however has a negligible impact on the layer intercalibration using muons described in Chapter 6. The $H \rightarrow 4\ell$ analysis presented in Chapters 8 and 9 uses the full Run 2 dataset and the dynamical clustering.

$|E_{\text{cell}}/\sigma_{\text{cell}}^{\text{noise}}| > 4$, where $\sigma_{\text{cell}}^{\text{noise}}$ is the expected noise (see Figure 4.3). These cells, called *seeds*, are then enriched with all neighbouring cells with a significance greater than two. At this step, two clusters sharing at least one cell are merged together. As a final step, the additional neighbouring cells with significance greater than 0 are added to the cluster. Only clusters with raw energy greater than 400 MeV and with $E_{\text{ECal}}/E_{\text{tot}} > 0.5$ (to suppress pileup clusters, mainly hadronic) are retained.

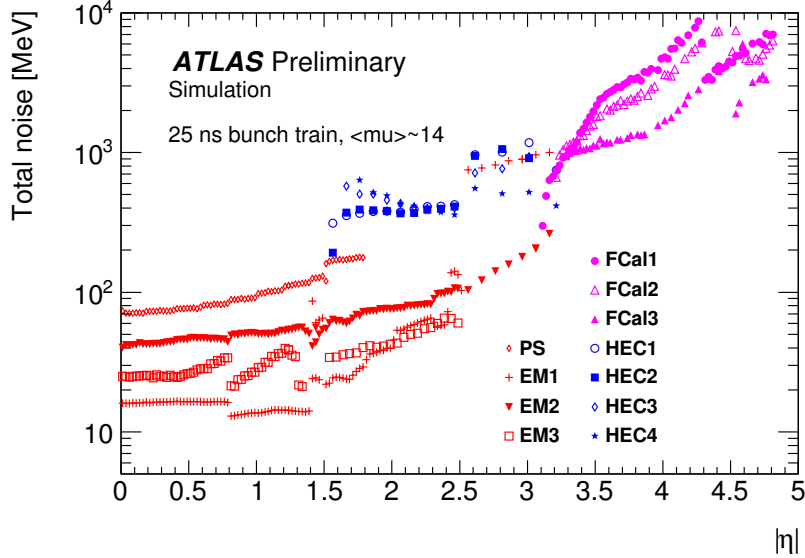


Figure 4.3 – Expected combined electronic and pileup noise (assuming $\langle\mu\rangle = 14$) per cell as a function of $|\eta|$ for each layer of the LAr calorimeter. [118]

An algorithm matching ID tracks to the clusters is then run, and tracks are considered as matching if they verify $|\eta_{\text{track}} - \eta_{\text{cluster}}| < 0.05$ and $-0.10 < q \cdot (\phi_{\text{track}} - \phi_{\text{cluster}}) < 0.05$, where q is the charge associated to the track and the cluster coordinates are taken as the second layer cell barycentre (weighted by their raw energy). The asymmetric $\Delta\phi$ requirement is due to possible radiated photons merged into the cluster. In the case of several matched tracks, the one with best ΔR is kept. Depending on whether the cluster was matched to a track, to a conversion vertex or not matched, the cluster is classified as electron, converted photon, or unconverted photon, respectively. Details of the method can be found in Ref. [132].

In addition to using topo-clusters instead of fixed-size cluster, another improvement in the clustering algorithm [132] has been included in the electron and photon reconstruction: this method consists of including additional *satellite* clusters to the seeding cluster. Such enlarged clusters are called *superclusters*. The underlying motivation for including additional clusters is that an electron passing through the material upstream the calorimeter is likely to emit bremsstrahlung photons. The above defined topo-clusters are considered as an electron supercluster seed if their energy is greater than 1 GeV and if it is associated to a track with four silicon hits, and as a photon supercluster seed if their energy is greater than 1.5 GeV. For electrons, additional satellite clusters in a window 0.125×0.3 are considered if they share the same track as the seed cluster. For converted photons, satellite clusters with a track matching the conversion vertex of the seed cluster are added. Only cells from the ECal are considered for the final supercluster. Furthermore, the η extension of each topo-cluster entering the supercluster is restricted to 0.075 in the barrel, or 0.125 in the endcap. The procedure is illustrated in Figure 4.4. The supercluster building algorithm is run independently for electrons and photons, so the final superclusters follow an ambiguity resolution algorithm sketched in Figure 4.5 to provide the final electrons and photons provided to the analysis. The electron reconstruction efficiency is about 90 % for $E_T = 7$ GeV and is above 97 % after 25 GeV.

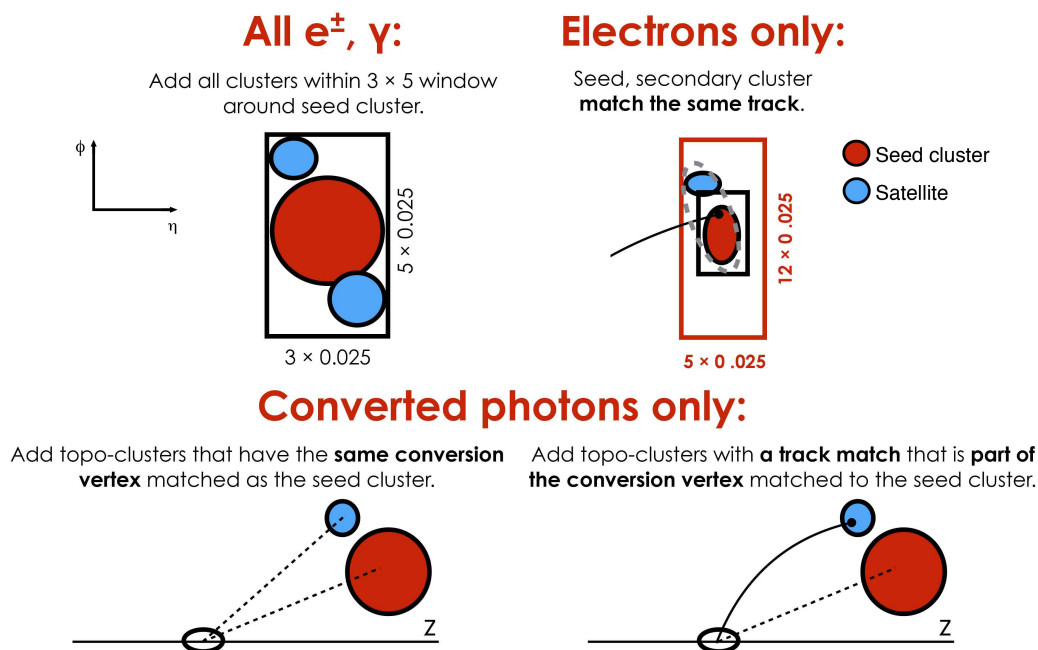


Figure 4.4 – Schematic of the supercluster building algorithm. [132]

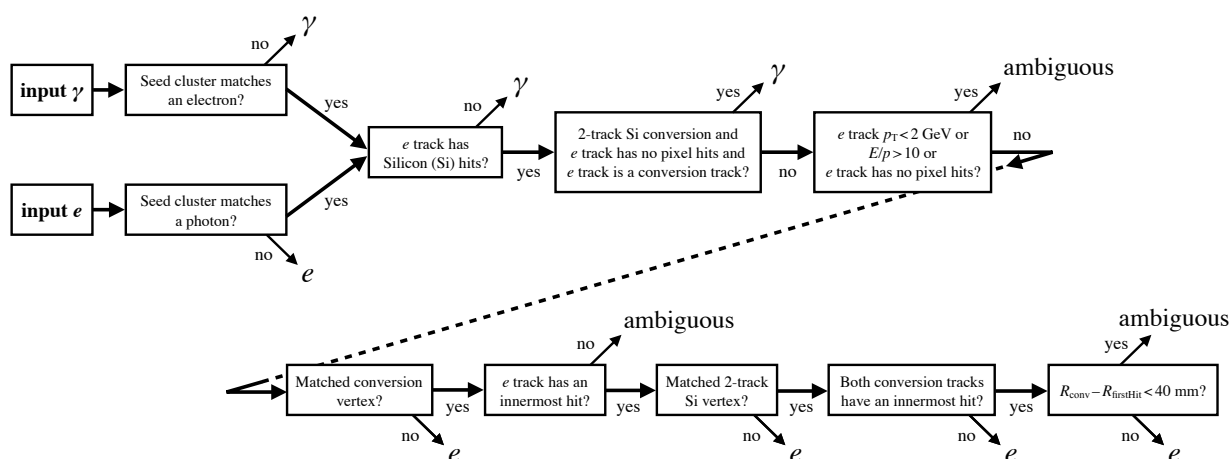


Figure 4.5 – Schematic of the electron/photon supercluster disambiguation algorithm. [132]

After electrons and photon are reconstructed, an additional step of identification is required to distinguish between real electron or photon clusters from clusters originating from light-hadron decays for example. The identification is based on a likelihood using shower shape information from the ECal, but also additional variables defined from the tracker to help distinguishing from hadrons producing displaced vertices and from the HCal to help distinguishing from all types of hadrons (contrary to electrons and photons, hadrons have a large energy fraction deposited in the HCal). The complete list of variables and their definition is given in Table 4.1. Three working points are defined offering various electron efficiency against background rejection. The $H \rightarrow 4\ell$ analysis described in Chapters 8 and 9 uses the loose working point, providing the highest electron identification efficiency (93 % on average) but the lowest background rejection. The breakdown of electron efficiency as a function of pileup and electron energy is shown in Figure 4.6.

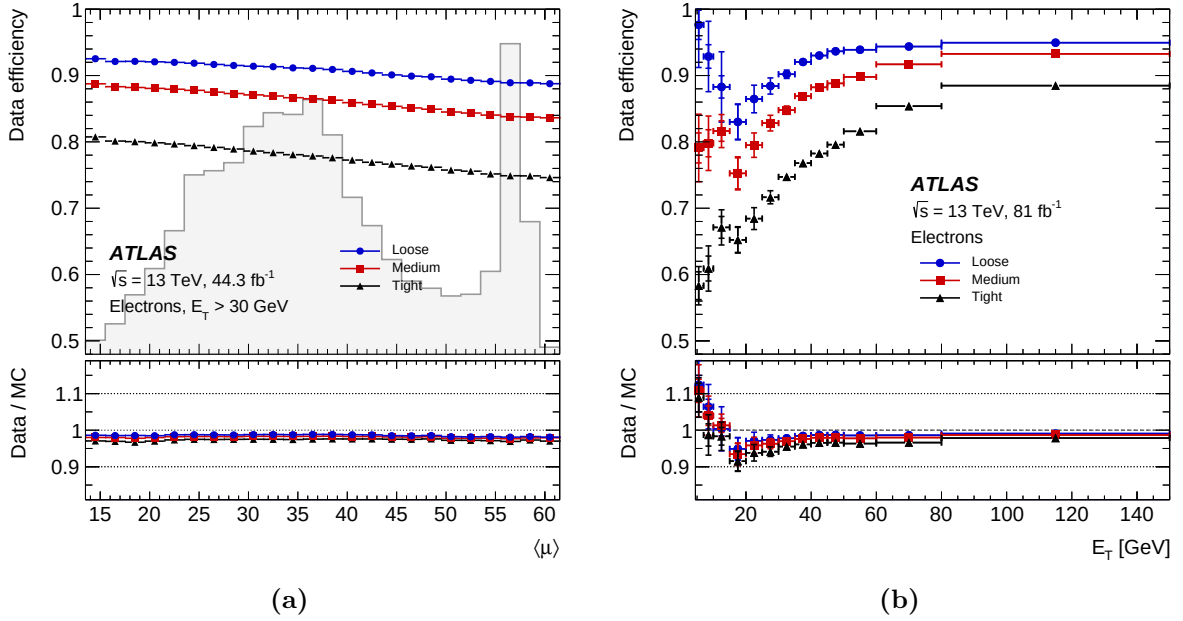


Figure 4.6 – Electron identification efficiency in the ATLAS detector as a function of (a) the pileup (2017 dataset only) and (b) electron p_T (2015 – 2017 dataset). In (a) the shaded histogram recall the $\langle\mu\rangle$ distribution in the 2017 dataset. The $H \rightarrow 4\ell$ analysis uses the *loose* identification criterion ensuring a 90 % electron identification efficiency. A small decrease of 3 % in efficiency is seen across the pileup range. As a function of p_T , the efficiency is around 90 % to 95 % in most of the range, with a dip at 80 % around $E_T = 20$ GeV. The efficiency at low E_T is dominated by the tracker measurement whose precision decreases with p_T while the efficiency at high E_T is dominated by the calorimeter measurement whose precision increases with E_T . The region of the dip corresponds to the transition region between the two regimes. [132]

4.5 Muons

The muon reconstruction procedure is fully described in Ref. [134], and its main elements are recalled here. Muons are primarily reconstructed using the MS information and most of them also include ID information, but some algorithms can provide muon reconstruction using the ID and calorimeter information only.

Combined muons (CB) are seeded from both ID and MS tracks, and the track is combined by using a global fit including all hits in both tracks. During the fit, some of the MS hits can be added or removed to improve the fit quality. Most CB muons are seeded from MS tracks extrapolated to the ID, but another algorithm extrapolating ID tracks to the MS is also used.

Table 4.1 – Listing of the discrimination variables used in the electron and photon identification. [132]

Category	Description	Name	Usage
Hadronic leakage	Ratio of E_T in the first layer of the hadronic calorimeter to E_T of the EM cluster (used over the ranges $ \eta < 0.8$ and $ \eta > 1.37$)	R_{had_1}	e/γ
	Ratio of E_T in the hadronic calorimeter to E_T of the EM cluster (used over the range $0.8 < \eta < 1.37$)	R_{had}	e/γ
EM third layer	Ratio of the energy in the third layer to the total energy in the EM calorimeter	f_3	e
EM second layer	Ratio of the sum of the energies of the cells contained in a $3 \times 7 \eta \times \phi$ rectangle (measured in cell units) to the sum of the cell energies in a 7×7 rectangle, both centred around the most energetic cell	R_η	e/γ
	Lateral shower width, $\sqrt{(\sum E_i \eta_i^2)/(\sum E_i) - ((\sum E_i \eta_i)/(\sum E_i))^2}$, where E_i is the energy and η_i is the pseudorapidity of cell i and the sum is calculated within a window of 3×5 cells	w_{η_2}	e/γ
	Ratio of the sum of the energies of the cells contained in a $3 \times 3 \eta \times \phi$ rectangle (measured in cell units) to the sum of the cell energies in a 3×7 rectangle, both centred around the most energetic cell	R_ϕ	e/γ
EM first layer	Lateral shower width, $\sqrt{(\sum E_i (i - i_{\text{max}})^2)/(\sum E_i)}$, where i runs over all cells in a window of 3 cells around the highest-energy cell, with index i_{max}	$w_{s\ 3}$	γ
	Total lateral shower width, $\sqrt{(\sum E_i (i - i_{\text{max}})^2)/(\sum E_i)}$, where i runs over all cells in a window of $\Delta\eta \approx 0.0625$ and i_{max} is the index of the highest-energy cell	$w_{s\ \text{tot}}$	e/γ
	Fraction of energy outside core of three central cells but within seven cells	f_{side}	γ
	Difference between the energy of the cell associated with the second maximum, and the energy reconstructed in the cell with the smallest value found between the first and second maxima	ΔE_s	γ
	Ratio of the energy difference between the maximum energy deposit and the energy deposit in a secondary maximum in the cluster to the sum of these energies	E_{ratio}	e/γ
Track conditions	Ratio of the energy measured in the first layer of the electromagnetic calorimeter to the total energy of the EM cluster	f_1	e/γ
	Number of hits in the innermost pixel layer	$n_{\text{innermost}}$	e
	Number of hits in the pixel detector	n_{Pixel}	e
	Total number of hits in the pixel and SCT detectors	n_{Si}	e
	Transverse impact parameter relative to the beam-line	d_0	e
	Significance of transverse impact parameter defined as the ratio of d_0 to its uncertainty	$ d_0/\sigma(d_0) $	e
	Momentum lost by the track between the perigee and the last measurement point divided by the momentum at perigee	$\Delta p/p$	e
	Likelihood probability based on transition radiation in the TRT	eProbabilityHT	e
Track–cluster matching	$\Delta\eta$ between the cluster position in the first layer of the EM calorimeter and the extrapolated track	$\Delta\eta_1$	e
	$\Delta\phi$ between the cluster position in the second layer of the EM calorimeter and the momentum-rescaled track, extrapolated from the perigee, times the charge q	$\Delta\phi_{\text{res}}$	e
	Ratio of the cluster energy to the track momentum	E/p	e

Segment-tagged muons (ST) are seeded from the ID and are accepted as muon tracks if the ID extrapolation matches one MDT or CSC segment in at least one muon station. They are mainly used in areas of low muon acceptance or when only one station is crossed.

Calorimeter-tagged muons (CT or Calo-tagged) are seeded from an ID track that matches a calorimeter energy deposit compatible with a MIP crossing. Such muons are primarily used in the region $|\eta| < 0.1$ where the MS has a reduced coverage.

Extrapolated (ME) or standalone (SA) muons only use information from the MS with a loose requirement on emerging from the interaction point, and additionally requiring that three stations have matching segments (only two required for $|\eta| < 2.5$). They mainly target the region $2.5 < |\eta| < 2.7$ which is beyond the ID acceptance. Yet, if some tracker information is available (one or two silicon hits are possible even for $|\eta| > 2.5$, see Figure 3.3), this information is included in the track fit.

If two muons share the same ID track, they undergo an overlap removal with a preference of order $CB > ST > CT$. In all cases the muon energy is corrected for its losses in the calorimeter.

As for electrons and photons, muons also undergo an identification step to suppress the background, mainly composed of non prompt muon from pion and kaon decays. Since the light hadrons decay in-flight, their track parameter in the ID is likely to not match the one determined from the MS. Three variables are therefore used to discriminate between the prompt signal and non prompt muons:

- the ID and MS charge momentum q/p difference significance defined as

$$\sigma_{p/q} = \frac{|(q/p)_{\text{MS}} - (q/p)_{\text{ID}}|}{\sqrt{\sigma((q/p)_{\text{MS}})^2 + \sigma((q/p)_{\text{ID}})^2}}, \quad (4.3)$$

- the p_{T} imbalance between the ID and the MS: $\rho = |p_{\text{T}}^{\text{MS}} - p_{\text{T}}^{\text{ID}}|/p_{\text{T}}^{\text{combined}}$,
- the normalised χ^2 of the combined track fit.

Four isolation working points are defined, using a combination of these variables.

The *medium* muons are the standard muons for analyses in ATLAS, and are used in the calibration analysis of Chapter 6. The *medium* criterion only retains combined muons, and standalone muons tracks are allowed in the region $2.5 < |\eta| < 2.7$. They are required to have at least three hits in at least two stations (except for $|\eta| < 0.1$) and $\sigma_{p/q} < 7$. The *Loose* muons are dedicated to the $H \rightarrow 4\ell$ analysis and are therefore used in the analysis described in Chapters 8 and 9. They include the *medium* muons and allow calo-tagged and segment-tagged muons in the region $|\eta| < 0.1$. Identification efficiencies for each of the working points are detailed in Table 4.2.

Table 4.2 – Muon and background identification efficiencies for all working points. [134]

	$4 < p_{\text{T}} < 20 \text{ GeV}$		$20 < p_{\text{T}} < 100 \text{ GeV}$	
Selection	$\epsilon_{\mu}^{\text{MC}} [\%]$	$\epsilon_{\text{Hadrons}}^{\text{MC}} [\%]$	$\epsilon_{\mu}^{\text{MC}} [\%]$	$\epsilon_{\text{Hadrons}}^{\text{MC}} [\%]$
Loose	96.7	0.53	98.1	0.76
Medium	95.5	0.38	96.1	0.17
Tight	89.9	0.19	91.8	0.11
High- p_{T}	78.1	0.26	80.4	0.13

4.6 Jets and b -tagging

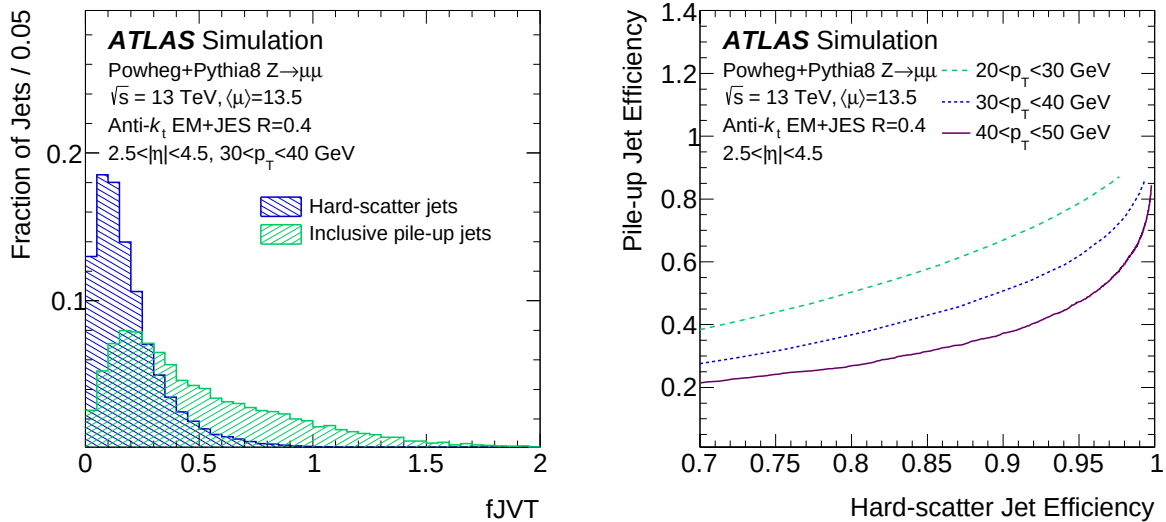
Final-state quarks or gluons created during collisions cannot be directly measured in the detectors: indeed, due to the confinement property of QCD, partons hadronise shortly after creation so only QCD bound states can be observed. The partons can also radiate gluons, that can in turn split into two gluons or two quarks, creating a shower like structure (an example is given in Figure 8.2). At sufficiently high energy, the particles get boosted along one direction, and are roughly contained in a cone, which is called a *jet*. That jet reconstruction and calibration algorithm aims at defining the area covered by a jet in the calorimeter and to measure the energy and momentum of the bunch of particles created, providing a proxy measurement for the underlying parton.

Jet reconstruction is seeded from topo-clusters [133] in a similar way as defined in Section 4.4, which are fed to the anti- k_t algorithm [135] with parameter $R = 0.4$. The flow of this algorithm is as follows: for a given cluster of index i , a distance-like variable $d_i = 1/p_{T,i}^2$ is computed as well as the distances to all other cluster with

$$d_{ij} = \min \left(\frac{1}{p_{T,1}^2}, \frac{1}{p_{T,2}^2} \right) \frac{\Delta R_{ij}^2}{R^2}. \quad (4.4)$$

Then, if $\forall j \neq i, d_i < d_{ij}$, then the cluster i forms a new jet. Otherwise, its cluster is merged with the closest cluster.

Due to the hadronic environment of the proton collisions, additional collisions in the same bunch crossing (*pileup*) create numerous other low p_T jets. Algorithms such as the jet vertex tagger (JVT and forward JVT) have been developed [136, 137] to mitigate the contamination of pileup jets reconstructed as hard-scatter jets. An example efficiency curve is shown in Figure 4.7.



(a) fJVT distribution for pileup (green) and hard-scatter (blue) jets for $30 < p_T < 40$.

(b) Pileup jet efficiency versus hard-scatter jet efficiency for various p_T ranges.

Figure 4.7 – Example performance of the forward jet vertex tagger (fJVT) in the region $2.5 < |\eta| < 4.5$ for an environment for an average number of collision per bunch crossing $\langle \mu \rangle = 13.5$. [137]

Another particular type of jet is the so-called b -jet, *i.e.* a jet that stem from the decay of a b -hadron. Since b -hadrons have a long enough lifetime to fly over hundreds of microns to millimetres (depending on their energy) before decaying, their decay vertex is displaced with respect to the primary vertex. Although this provides a measurable signature to tag a b -decay, the dense track environment linked to the high luminosity at the LHC does not allow for an easy b -tagging. However b -tagging is of prime importance to reconstruct top quark decays needed in dedicated top quark analyses as well as in Higgs boson production associated to top quarks.

Experiments have therefore developed b -tagging algorithm using machine learning methods, such as the ATLAS MV2 BDT-based algorithm [138], to discriminate between prompt b -jets and backgrounds made of c -jets, displaced vertices from τ^- decays and other light-flavour induced jets. This algorithm nominally provides four working points described in Table 4.3 which can be combined into a pseudo-continuous b -tagging efficiency. Its efficiency curve is displayed in Figure 4.8.

Table 4.3 – b -jet efficiency and c -jet, τ -jet and light-flavour jet rejection factor using the MV2C10 algorithm at its four nominal working points. [138]

ϵ_b	MV2			
	Selection	Rejection		
		c -jet	τ -jet	Light-flavour jet
60%	> 0.94	23	140	1200
70%	> 0.83	8.9	36	300
77%	> 0.64	4.9	15	110
85%	> 0.11	2.7	6.1	25

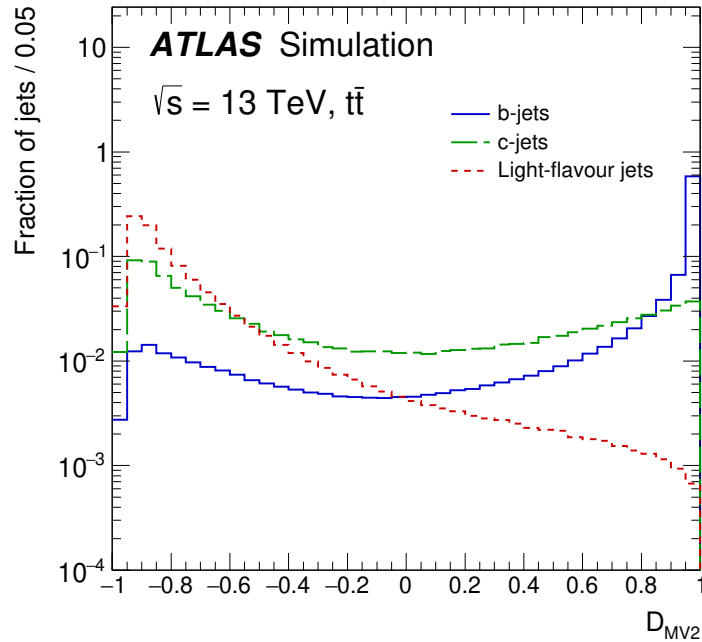


Figure 4.8 – MV2C10 BDT score for b -jets, c -jets and light-flavour jets. [138]

4.7 Missing transverse energy

Neutrinos only interact by weak interaction, making them hard to detect. In practice, in an experiment like ATLAS, they pass through the detector without interacting and therefore cannot be directly detected. This means that all disintegrations containing a neutrino, such as W boson or τ -lepton decays, cannot be fully measured. However, in a collision at the LHC the partons participating in the collision can be considered as moving only in the z direction. Therefore the momentum and energy in the transverse direction vanishes and by conservation of momentum and energy, the vectorial sum of momenta and energies of the final-state objects should also amount to zero. As a consequence, a neutrino, totally escaping the detection, can be inferred from the imbalance of energy in the final state.

The missing transverse energy (MET) denoted E_T^{miss} is defined as the opposite of the sum of all energies measured in the detector: $\vec{E}_T^{\text{miss}} = -\sum \vec{E}_T^{\text{visible}}$. These include reconstructed physics objects such as electrons, photons, jets and muons, but also tracks associated to primary vertices and not included in physics objects (called *soft-term*, for $|\eta| < 2.5$) and is corrected for the pileup. It provides a proxy for all undetectable particles, *i.e.* neutrinos but also some weakly interacting hypothetical BSM particles. The E_T^{miss} however cannot provide information on single escaping particles, only on their sum.

Given its definition, the E_T^{miss} is sensitive to all uncertainties in the electron, photon, jet and muon measurements, as well as to the pileup contribution. Moreover, the ATLAS detector only has an approximate 4π coverage, so particles passing through non instrumented regions also give a contribution to the E_T^{miss} . The Run 2 performance of the ATLAS E_T^{miss} algorithm is described in Ref. [139].

Calibration of the electromagnetic calorimeter

Contents

5.1	Electron and photon calibration overview	102
5.2	MVA-based calibration	103
5.3	E_1/E_2 intercalibration	106
5.4	Pre-Sampler scales	106
5.4.1	Overview of the PS scale determination	106
5.4.2	Determination of the correlation factor $A(\eta)$	107
5.4.3	Determination of the correction factor $b_{1/2}(\eta)$	107
5.4.4	Results	109
5.4.5	PS scales with muons	110
5.5	Third layer energy calibration	111
5.6	Material estimation in front of the calorimeter	111
5.6.1	Reasons for a precise material estimate and methodology	111
5.6.2	Results and validation	112
5.6.3	Nominal geometry for Run 2	113
5.7	$Z \rightarrow ee$ energy scales	115
5.7.1	Overview	115
5.7.2	Estimation of α_{ij} and c'_{ij} using the template method	118
5.7.3	Estimation of α_i and c'_i using the Z lineshape method	119
5.7.4	$Z \rightarrow ee$ scale results and systematics	120
5.8	Energy scale and resolution extrapolation over the energy range	121
5.9	Cross-checks of energy scales and resolution extrapolation	125
5.9.1	Methodology	125
5.9.2	Determination of the J/ψ prompt fraction	125
5.9.3	Extraction of residual scales using the lineshape method	126
5.9.4	Results and systematics	126

The full calibration procedure if the LAr calorimeter was established during Run 1, and the results of the final Run 1 calibration are detailed in Ref. [140]. Although the Run 2 procedure remains similar in Run 2, several optimisations and updates of the analyses using the 2015 and 2016 datasets are described in Ref. [141]. This chapter recalls the major aspects of the

electromagnetic calorimeter calibration, as used for the electron and photon energy calibration with $|\eta| < 2.5$.

The photon calibration is similar to the electron calibration and most of its steps, analyses, and results are actually common with the electron calibration. However since photons are not of central interest in the $H \rightarrow 4\ell$ analysis, photon calibration will not be discussed in detail. Forward electrons with $|\eta| > 2.5$, where no tracking information is available, are also calibrated but are beyond the scope of this Thesis since they are outside the $H \rightarrow 4\ell$ analysis acceptance. Finally, dedicated calibrations are applied to other physics objects that are used in the $H \rightarrow 4\ell$ analyses (muons, jets and b -tagged jets, missing transverse energy). However, as those calibrations radically differ from the LAr electromagnetic calorimetry considerations and are not part of this Thesis work, they will not be described here.

5.1 Electron and photon calibration overview

The previous Chapter described how the energy at the cell level is reconstructed from the electronics signal and how physics objects such as electrons or photons are built from higher level information given by tracks and cell clusters. These objects are the basic input to physics analyses, from searches of new particles to precision measurements. However such analyses aim for precision measurements, and using physics objects with energy only barely reconstructed will not achieve the precision required. To achieve such precision, physics objects after reconstruction need to be further calibrated.

The analysis work included in this Thesis (Chapters 8 and 9) focuses on Higgs-boson measurements in the four-lepton decay channel. This analysis makes a central use of low and medium energy electrons (90 % are below 100 GeV) located in the precision region of the ATLAS experiment ($|\eta| < 2.5$). The standard calibration for such electrons include several steps, one of which is a large part of this Thesis work. The main steps are summarised in this Chapter, and the following is dedicated to the calibration work achieved in this Thesis.

The calibration procedure starts from the raw cluster energy, reconstructed as described in Chapter 4. The first step is to derive a first estimate of the energy through an MC-trained MVA which is applied to both data and MC (Section 5.2). Some corrections are applied to the data in order to correct for the uniformity and homogeneity of the detector (for example a ϕ -dependent correction to correct for the structural material of the calorimeter, or gain-specific corrections). These are not described here. Then, the calorimeter longitudinal response is calibrated in the so-called layer intercalibration steps (Chapter 6, Sections 5.4 and 5.5). Section 5.6 describes the material measurements that are needed to ensure a good MVA calibration, and which needs a proper longitudinal calibration. Finally the global energy scale, applied to the data, and resolution corrections, applied to the MC, are derived (Section 5.7), and extrapolated to other energy ranges (Section 5.8). Cross-checks of these extrapolations are eventually performed and are described in Section 5.9. The steps of the calibration procedure are summarised in Figure 5.1. At the end of the procedure, electron energy precision reaches a level within the requirements of the precision analyses run by the ATLAS experiment.

All the analyses described here use 36.1 fb^{-1} of data collected in 2015 and 2016, except for the gain study which used special runs taken in 2017, and for the final energy scale and resolution corrections which can absorb most the residual miscalibration and are therefore derived every year.

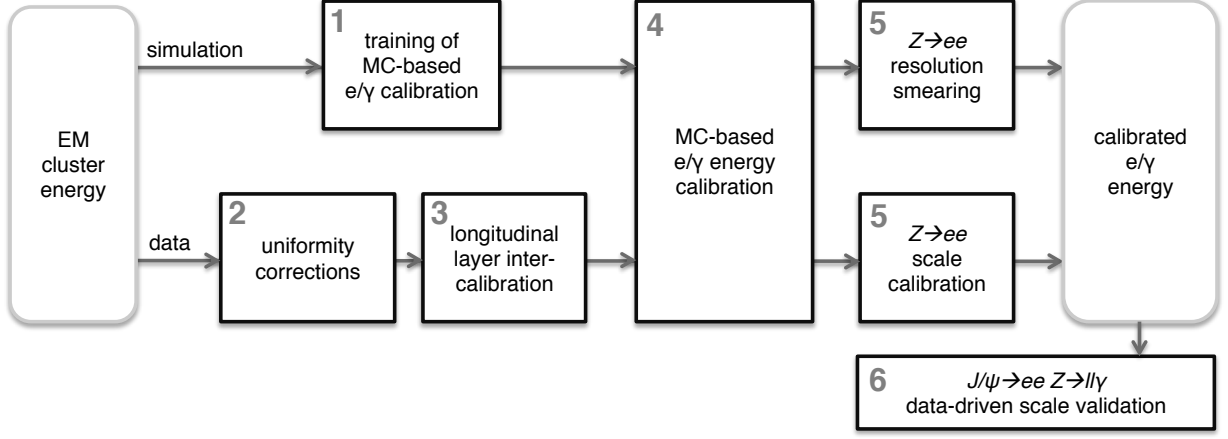


Figure 5.1 – Summary of the steps from uncalibrated cells/cluster energy to a calibrated physics object. [140]

5.2 MVA-based calibration

The first step in the calibration procedure consists in estimating the energy of electron or photon given the energies in each layer of the calorimeter. As the shower development depends on the amount and position of the material traversed by the particle before hitting the calorimeter as well as small variations of the electrode or absorber geometry, the calibrated energy reconstruction is not as simple as $E_{\text{calibrated}} = E_0 + E_1 + E_2 + E_3$ and depends on the position, the energy, and the type of particle (*e.g.* unconverted photons are much less sensitive to the upstream material for example). To account for all these subtle variations in a simple way, the calibrated energy reconstruction is performed with a multivariate analysis (MVA). This type of algorithm makes use of machine learning techniques to optimise the combination of variables given as input in order to minimise the standard deviation of the target output variable. In this case, a regression Boosted Decision Tree (BDT) was used to target the true simulated energy of particles (in fact the correction factor to the summed energy in the accordion $E_{\text{acc}} = E_1 + E_2 + E_3$, *i.e.* $E_{\text{true}}/E_{\text{acc}}$), and thus to minimise the energy resolution [142].

To account for large scale variations of the calorimeter geometry and detector inhomogeneities, the MVA is trained separately in bins of $|\eta|$. Moreover, since the electromagnetic shower development changes with energy, the training is done separately in bins of transverse energy $E_{T,\text{acc}}$. Finally, the MVA is trained separately for electrons, converted and unconverted photons: the electromagnetic shower shape will differ in each case. In particular unconverted photons have showers starting later than electrons due to a first interaction with material occurring much later (*i.e.* in the calorimeter instead of upstream), while electrons lose part of their energy before reaching the calorimeter. Converted photons have a behaviour between the two, depending on the length travelled by the photon before converting. This amounts to 111 bins in $E_{T,\text{acc}} \times |\eta|$, for each particle type (electron, converted photon, unconverted photon) [143].

To accommodate for all of the effects previously mentioned (both the detector uniformity and the energy dependence), the following variables are used for the BDT training [143]:

- $E_{\text{acc}}^{\text{raw}} = E_1^{\text{raw}} + E_2^{\text{raw}} + E_3^{\text{raw}}$: the total energy collected in the cluster summed over the three layers of the accordion.
- $E_0^{\text{raw}}/E_{\text{acc}}^{\text{raw}}$: the ratio of energy deposited in the pre-sampler over the total energy deposited in the accordion.
- $E_1^{\text{raw}}/E_2^{\text{raw}}$: the ratio of energy in the first and second sampling of the calorimeter.
- η_{cluster} : the pseudo-rapidity in the ATLAS frame. This variables allows to take into account

the misalignment of the calorimeter and the material in front of the calorimeter, for example.

- $\eta_{\text{calo}}/\Delta\eta$: the cluster central cell index in the η direction in the calorimeter frame ($\Delta\eta = 0.025$ is the cell width in the second sampling). This variable is more sensitive to inhomogeneity of the calorimeter like variation of cell width or depth of the third sampling.
- η_{calo} modulo $(\Delta\eta)$, with $\Delta\eta = 0.025$: distance of the cluster centre to the edge of the central cell in the η direction. This variable helps to correct for the finite cluster size.
- ϕ modulo $2\pi/1024$ in the barrel, $2\pi/768$ in the endcap. The value of the modulus corresponds to the periodicity of the lead absorbers in the calorimeter, this variable is thus used to correct for slight variations of the sampling fraction seen by incident particles.

In addition to these 7 variables, the converted photon BDT is trained with 3 additional variables if the conversion radius is $R_{\text{conv}} < 800$ mm:

- R_{conv} , the conversion radius. This is used only if the sum of transverse momentum of the conversion tracks is $p_T > 3$ GeV.
- $E_{T,\text{acc}}^{\text{raw}}/p_T$, with $E_{T,\text{acc}} = E_{\text{acc}}/\cosh(\eta_{\text{cluster}})$: ratio of the transverse energy in the accordion to the conversion tracks transverse momentum.
- $p_T(\text{leading track})/p_T$: fraction of the sum of the conversion tracks momenta carried by the highest p_T conversion track.

In order to improve the calibration in the barrel-endcap transition region (“crack”: $1.4 < |\eta| < 1.6$), an additional energy measurement from coarse scintillators of the Intermediate Tile Calorimeter (“ E_4 ” with $\Delta\eta \times \Delta\phi$ granularity of 0.1×0.1 (Figure 3.14a) was added to the training [144]. An example of energy resolution improvement in this region is shown in Figure 5.2, and reaches 20 %.

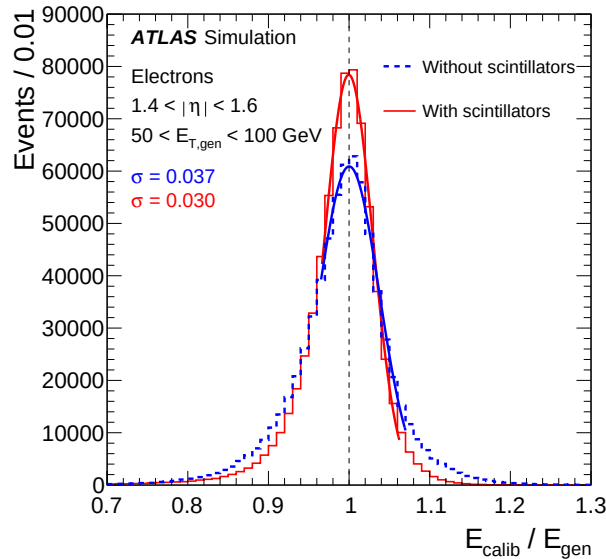


Figure 5.2 — Distribution of the calibrated energy divided by the truth energy for electrons with generated energy between 50 and 100 GeV within $1.4 < |\eta| < 1.6$. The blue histogram shows the resolution without including the E_4 scintillators, the red histogram after including them. The curves correspond to gaussian fits on the core of the distribution, from which the widths are extracted. [141]

The energy resolution achieved for electrons after the MVA is applied on simulated electrons is shown in Figure 5.3.

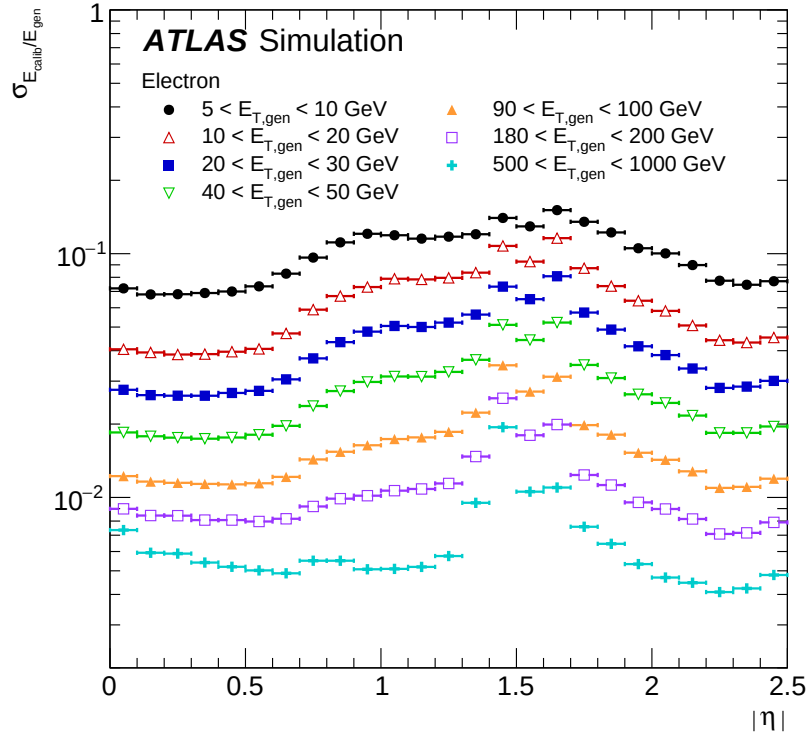


Figure 5.3 — Energy resolution of electrons as a function of their pseudo-rapidity, for various generated energies. The resolution is defined as the interquartile range (range of the second and third quartile) divided by 1.35 (to get closer to a gaussian width interpretation). The energy resolution is degraded in the $|\eta|$ range $[1.2, 1.8]$ due to the presence of more passive material in front of the calorimeter. [141]

5.3 E_1/E_2 intercalibration

The electronics and readout performance of the electromagnetic calorimeter was first measured during test beams [119, 128]. This calibration can however be cross-checked using real data recorded during standard LHC runs [126]. The layer intercalibration is designed to correct for residual miscalibration of the readout electrodes in the calorimeter, including slight mismodelling of the cross-talk or residual electronics miscalibration.

The intercalibration is primarily performed using muons, benefiting from their behaviour as minimum ionising particles (MIP): muon energy deposits will not depend on the material in front of the calorimeter and its possible mismodelling, but only on the length of active material (liquid argon) traversed. A secondary method uses electrons to perform a cross-check of the muon result, but suffers from additional electromagnetic shower shape mismodelling. As this has been a large part of this thesis work, the description of this method is given in Chapter 6.

5.4 Pre-Sampler scales

5.4.1 Overview of the PS scale determination

As described in Chapter 3.4.3, the main purpose of the Pre-Sampler (PS) is to estimate the energy lost by EM showers in front of the accordion, therefore its response must also be calibrated to accurately correct for the electromagnetic showers energy scale. The PS scale α_{PS} is defined as the ratio of energy deposited in the PS in the data to the one of MC. It is computed using electrons from the well-known $Z \rightarrow ee$ decay [145] which provides a controlled sample of electrons. A similar method to the one used for the E_1/E_2 intercalibration (with muons) was also tested and is used as a cross-checked (see Section 5.4.5). The difficulty of using muons in the PS though, is that the signal-to-noise ratio is very low: the PS is less than $1 X_0$ (see Section 3.4). Using electrons, issues in estimating the energy deposit following a Landau distribution are discarded, and since the shower extends over several cells, the pileup noise has a much smaller effect.

Since this method uses electrons, it is sensitive to material present in front of the PS. Consequently, a mismodelling of the material induces a bias in the measurement and this effect has to be corrected before the α_{PS} scale is interpreted as an energy correction. The material affecting the PS measurement is located in two places which will have a different impact. First, the material located in front of the PS directly impact the PS scale and the calibration of electrons. Corrections for this effect are derived in a set of electron samples simulated with various distorted geometry, profiting from the correlations between the PS energy (E_0) and the ratio of energy in the first and second sampling of the calorimeter ($E_{1/2}$). This step is described in Section 5.4.2. Second, the material located between the PS and the accordion impact the measurement of electrons, but not the PS scales. This correction has to be evaluated independently of the material in front of the PS and is thus performed using unconverted photons from an inclusive sample or from radiative $Z \rightarrow \ell\ell\gamma$ decays. This second step is described in Section 5.4.3.

Once these two corrections ($A(\eta)$ and $b_{1/2}(\eta)$) are derived, the PS scales α_{PS} is defined separately in several $|\eta|$ bins (to cover for detector and material inhomogeneities) as in [146] by

$$\alpha_{\text{PS}}(\eta) = \frac{E_0^{\text{data}}(\eta)}{E_0^{\text{corr}}(\eta)}, \quad (5.1)$$

where $E_0^{\text{corr}}(\eta)$ is PS energy from MC after all material correction applied. It is defined as

$$\frac{E_0^{\text{corr}}(\eta)}{E_0^{\text{nom}}(\eta)} = 1 + A(\eta) \left(\frac{E_{1/2}^{\text{data}}(\eta)}{E_{1/2}^{\text{nom}}(\eta)} \times \frac{1}{b_{1/2}(\eta)} - 1 \right). \quad (5.2)$$

Here,

- $E_0^{\text{data}}(\eta)$ and $E_0^{\text{nom}}(\eta)$ is the average energy deposit in the PS from electrons, in the data and in simulation with nominal geometry respectively;
- $A(\eta)$ is the correlation factor between $E_{1/2}$ and E_0 from varying the material budget in front of the PS (Section 5.4.2);
- $E_{1/2}^{\text{data}}(\eta)$ and $E_{1/2}^{\text{nom}}(\eta)$ is the average of the ratio of energy deposited in the first and second layer of the calorimeter by electrons, in the data and in the simulation with nominal geometry respectively;
- $b_{1/2}(\eta)$ is the ratio $E_{1/2}^{\text{data}}(\eta)/E_{1/2}^{\text{nom}}(\eta)$ using unconverted photons with low activity in the PS (Section 5.4.3).

To benefit from a good estimate of $E_{1/2}$, the $E_{1/2}$ intercalibration described in Chapter 6 is applied first.

5.4.2 Determination of the correlation factor $A(\eta)$

The first step in the computation of the PS scales is the determination of the correlation factor $A(\eta)$, needed to correct for material mismodelling and used to overcome the dependency on material located in front of the PS.

The determination of this factor makes use of the correlations between the PS energy (E_0) and the ratio of energy in the first and second samplings ($E_{1/2}$) as a function of the material [146]. Indeed, if more material is present in front of the PS, the electromagnetic shower is likely to start earlier, and the maximum of the energy deposit (or equivalently the shower energy barycentre) will be located closer to the interaction point. The $E_{1/2}$ ratio is therefore a good proxy of the shower extent in the longitudinal direction. Thus, more material induces an earlier shower development, increasing both E_0 and the $E_{1/2}$ ratio.

Several $Z \rightarrow ee$ samples have been produced with various modifications in the detector material description. These so-called *distorted geometries* are listed in Table 5.1 and examples of correlation plots are shown in Figure 5.4. The various correlations points are well aligned, except for two points: those correspond to geometries with an increased amount of material both before the PS and between the PS and the accordion, and the $E_{1/2}$ ratio gets therefore additional shifted.

The $A(\eta)$ correlation parameter is finally determined by fitting such correlation plot separately in each $|\eta|$ bin by

$$\frac{E_0^{\text{dist}}(\eta)}{E_0^{\text{nom}}(\eta)} = 1 + A(\eta) \left(\frac{E_{1/2}^{\text{dist}}(\eta)}{E_{1/2}^{\text{nom}}(\eta)} - 1 \right), \quad (5.3)$$

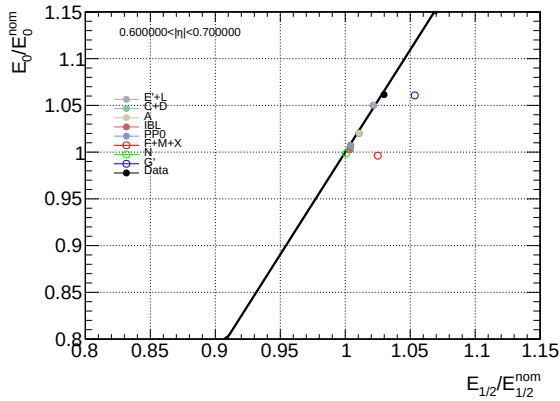
excluding the geometries with material variation between the PS and the accordion to be only sensitive to material in front of the calorimeter. This is similar to Eq. 5.2, setting $b_{1/2}$ to 1 thus imposing no correction for the nominal geometry.

5.4.3 Determination of the correction factor $b_{1/2}(\eta)$

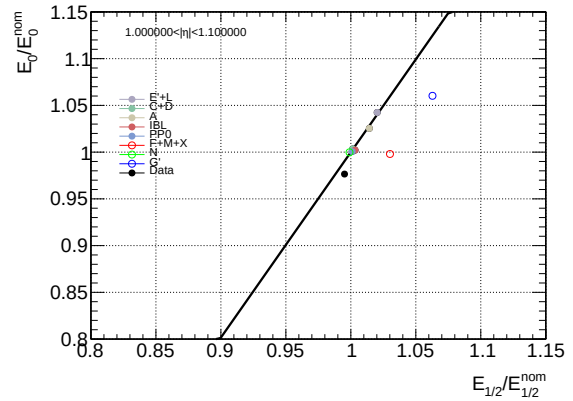
The second step prior to the determination of the PS scales is the computation of a $E_{1/2}$ correction factor to account for material difference located between the PS and the accordion and therefore affecting the measurement of electron energy but not the actual PS scales. To probe such material effect, a sample of particles sensitive to material between the PS and the accordion but not sensitive to material in front of the PS is needed. Such requirements are fulfilled by unconverted photons, *i.e.* photons that do not convert in the material between their production vertex and

Table 5.1 – List of distorted geometries used in the evaluation of the correlation factor $A(\eta)$. The “ID” geometries include changes in the Inner Detector, “Pixel S” and “SCT S” have increased services material, the “PS/S1” configuration have additional material between the PS and the accordion, the “Cryo 1” adds material from the cryostat in front of the calorimeter in the barrel and the “Calo-EC” adds material specific to the endcap calorimeter. All variations are an absolute change of X_0 with respect to the nominal MC geometry, except for the configuration A where the change is a 5% relative material scaling to the whole ID. A pictorial representation of the areas affected by various ID material change are included in Appendix A. [146]

Config	ID	ID-EC	Pixel S	SCT S	SCT/TRT-EC	PS/S1-B	PS/S1-EC	Cryo 1	Calo-EC
$ \eta $ region	[0, 2.4]	[1.8, 2.4]			[1.6, 2.2]	[0, 1.5]	[1.5, 1.8]	[0, 1.6]	[1.4, -]
A	5 %	-	-	-	-	-	-	-	-
N	-	-	-	-	-	-	0.05	-	-
C'+D'	-	-	0.1	0.1	-	-	-	-	-
E'+L'	-	-	-	-	0.075	-	-	0.05	-
F'+M+X	-	0.075	-	-	-	0.05	-	-	0.3
G'	5 %	0.075	0.1	0.1	0.075	0.05	0.05	0.05	0.3
IBL	Improved IBL geometry								
PP0	50 % increase in IBL + pixel services								



(a) For $|\eta| \in [0.6, 0.7]$



(b) For $|\eta| \in [1.0, 1.1]$

Figure 5.4 – Example correlation plots between $E_{1/2}^{\text{dist}}/E_{1/2}^{\text{nom}}$ and $E_0^{\text{dist}}/E_0^{\text{nom}}$ for various distorted geometries and linear fit among the various distorted geometries (Eq. 5.3), (a) for $|\eta| \in [0.6, 0.7]$ and (b) for $|\eta| \in [1.0, 1.1]$. The data point is shown for comparison but not used in the fit. The geometry variations are described in Table 5.1. [146]

the calorimeter. Photons are further required to have a low energy deposit in the PS to ensure they had no interaction in the cryostat walls or the solenoid in front of the calorimeter.

This part of the analysis selects photons coming from both a radiative $Z \rightarrow \mu\mu\gamma$ decays sample and an inclusive sample to cover the full p_T phase-space [146]. In both cases, a low activity in the PS is required to assess that the shower did not start before the PS: $E_0 < 0.5$ GeV. The results for each samples are defined as the mean of the double ratio $E_{1/2}^{\text{data}}/E_{1/2}^{\text{MC}}$. The results from the two samples are then combined as their weighted average, as shown in Figure 5.5. The $b_{1/2}(\eta)$ used in the final derivation of α_{PS} are provided after averaging $b_{1/2}$ over several $|\eta|$ bins: a trade off must be found between using a fine binning, reducing the extrapolation effect from the muon $E_{1/2}$ correction to electrons, and a coarser binning, with higher statistics stabilising the α_{PS} result. The averaging is then performed module-wise, *i.e.* with $\Delta\eta = 0.2$, as shown in Figure 5.5.

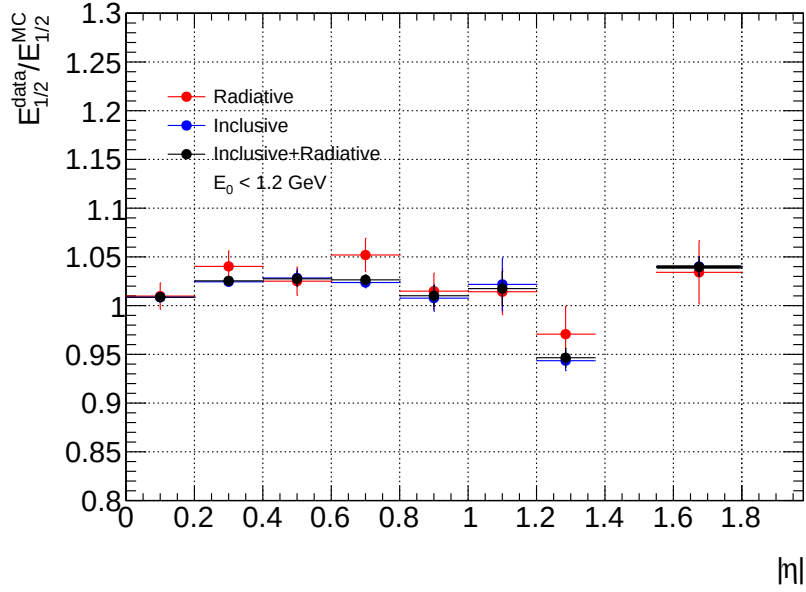


Figure 5.5 – Data-to-MC ratio of $E_{1/2}$ as a function of $|\eta|$ for photons from radiative $Z \rightarrow \mu\mu\gamma$ decays sample (red) and from the inclusive sample (blue). The two results are statistically combined (error-weighted average) in black. Since photons are poorly calibrated in the crack region ($1.37 < |\eta| < 1.55$) no $b_{1/2}$ measurement is provided for this region. A PS energy veto of 1.2 GeV is applied instead of the nominal 0.5 GeV, corresponding to a systematic uncertainty (see Section 5.4.4). [146]

5.4.4 Results

The final PS scales are finally derived using Eqs. 5.1 and 5.2 and are presented in Figure 5.6. While the analysis is derived in bins of $\Delta\eta = 0.05$, the final scales are averaged in each PS module, of size $\Delta\eta = 0.2$ (0.25 in the endcap): no significant variation is indeed expected within one single module of the PS.

The total uncertainty on the PS scales α_{PS} amounts to 1.5 % to 3 % depending on $|\eta|$, except for the last barrel bin in which $b_{1/2}$ has unexplained deviations depending on the averaging choice (per-module or over all barrel), leading to an increased uncertainty [141, 146]. Systematic uncertainties include the spread over a module, the systematics from the measurement of $A(\eta)$, and $b_{1/2}(\eta)$ related systematics: averaging choice and impact of changing the unconverted photon PS veto value from 0.5 to 1.2 GeV [146].

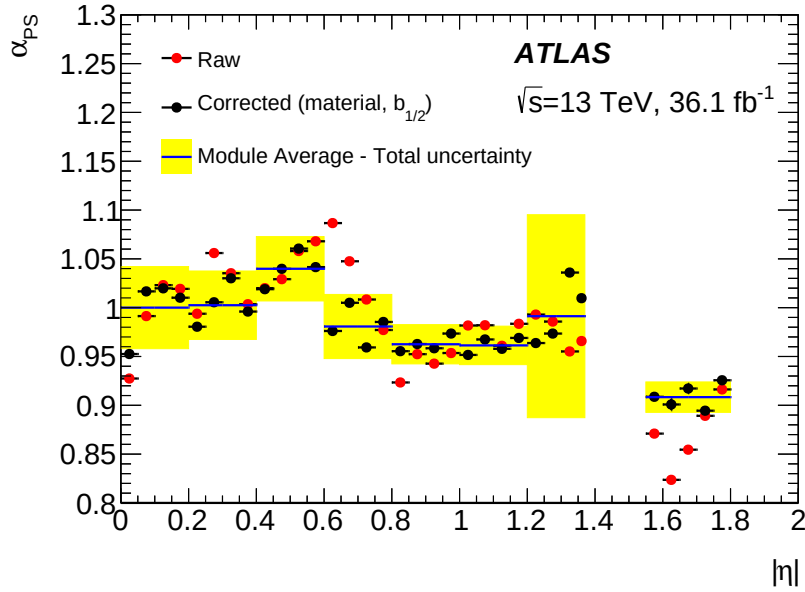


Figure 5.6 – PS scale α_{PS} and its uncertainty across the $|\eta|$ range, in red before the material and $b_{1/2}$ correction are applied and in black after these corrections are applied. The uncertainty on the measurement is within 1.5 to 3% except in the last barrel bin where it amounts to 10% due a large deviation of $b_{1/2}$ in this region. [141]

5.4.5 PS scales with muons

As an alternative to using electrons, the PS scales are derived using muons following a procedure similar to the one described in Chapter 6. On the one hand, since the signal-to-noise ratio for muon energy deposit in the PS is low (0.2-0.3), using muons over electrons makes it hard to extract a reliable value for the deposited energy. Yet, the available statistics are still sufficient to provide meaningful results. On the other hand, muons are not sensitive to material effects, so there is no need to compute corrections such as the $A(\eta)$ and $b_{1/2}(\eta)$ described above, simplifying the procedure.

The muon energy deposit $\langle E_0 \rangle$ is evaluated using the truncated mean method described in Section 6.5, with 2 iterations and restricting the range to compute the mean in $\mu_0 \pm 2 \cdot \text{RMS}$. This estimation is performed in bins of pileup and finally linearly extrapolated to $\langle \mu \rangle = 0$, similarly in both data and MC (Section 6.4.4). The PS scales are then simply defined as $\alpha_{\text{PS}} = E_0^{\text{extr, data}} / E_0^{\text{extr, MC}}$.

The results are found to be in agreement with the electron method within uncertainties. Systematics of the muon method include various definitions of the truncated mean (number of iterations, size of the core used to compute the mean), the number of cell used to recover the deposited muon energy, and alignment uncertainties [146]. The total systematic defined as the quadrature sum of the previously mentioned sources amounts to 3% to 5% depending on the $|\eta|$ region probed.

Another method using an analytical fit (Landau convolved with a gaussian) as defined in Section 6.4 instead of the truncated mean is investigated. This analytical fit method is however found to provide less stable results than the truncated mean method as the fits suffer from the low signal-to-noise ratio (SNR) in the distribution. The two methods remain in reasonable agreement over the $|\eta|$ range.

5.5 Third layer energy calibration

The third layer of the EM calorimeter (so-called *back* layer) was primarily design to measure the leakage energy of very high energy showers into the tile calorimeter. Thus electrons and photons of energy in the range studied in this global calibration procedure (≈ 10 GeV to a few hundreds of GeV) will only deposit a negligible amount of energy in this back layer [140]. For this reason, no dedicated calibration of the third layer is performed.

A sanity check was however performed using Run 1 data using electrons and simply looking at the $E_3^{\text{data}}/E_3^{\text{MC}}$ ratio [147]. The study found a possible bias of +10 to 30 % (and up to 250 % for $|\eta| > 2$) of the third layer energy in data. While this might seem a lot, the third layer only accounts for $\sim 2X_0$ (up to 8 depending on $|\eta|$, see Chapter 3.4.3) compared to the rest of the calorimeter (first and second sampling) which is more than $20X_0$ deep. Moreover the study in Ref. [147] concluded that the difference between data and MC is more likely to be explained by larger cross-talk effects than expected rather than by a genuine calibration issue. The induced bias on the total energy was found to be at worse 0.1 % for electrons with $p_T < 100$ GeV, confirming that no dedicated calibration is needed for such energy range. This bias on total energy could however grow up to a few percents for endcap electrons in the TeV range, indicating that the calibration of very high energy electrons and photons would require a dedicated analysis of the back layer in the future.

5.6 Material estimation in front of the calorimeter

5.6.1 Reasons for a precise material estimate and methodology

A good description of the material is paramount to achieve a performing calibration and reconstruction. On the reconstruction and identification side, an inaccurate material estimate leads to an inaccurate estimate of the fraction of converting photons and of electron bremsstrahlung. On the calibration side, the first energy estimate provided by the MC-based MVA calibration heavily relies on a good knowledge of the geometry and material to finely account for such effects. Measurements of material budget were therefore performed during the detector construction and installation, with an accuracy of 5 % [97].

As explained for the PS calibration (Section 5.4), a natural proxy of the material budget is found in E_1/E_2 : electromagnetic showers start closer to the interaction point in case of additional material [148]. Since the MC is simulated with the best geometry knowledge (*i.e.* material estimate) of the detector, a comparison of $E_{1/2}$ in data and simulation provides an estimate of the material budget difference in front of the calorimeter. One caveat of this method though is that it requires a well calibrated E_1/E_2 since genuine material effects cannot be distinguished from inaccurate layer intercalibration. This layer intercalibration is fully described in the Chapter 6.

In a similar way to the PS scales analysis, E_1/E_2 of electrons and unconverted photons are used to probe possible additional material from different parts of the detector. Unconverted photons with low PS activity do not start their shower before reaching the calorimeter, and thus are only sensitive to material located between the PS and the accordion. On the contrary, electrons start their shower earlier and are sensitive to material both in front of the PS and between the PS and the accordion.

The material distribution is assumed to be fairly symmetric as a function of ϕ , so the analysis is integrated along this direction [148]. The main dependency comes from the η direction as passive material such as cryogenic feed-through, service cables *etc.* are located as specific locations in η (see Chapter 3). The detector is also considered symmetric for $\eta < 0$ and $\eta > 0$, so the analysis

is performed as a function of $|\eta|$. In order to account for local variations of material budget, the analysis uses a binning of $\Delta\eta = 0.05$ for the electron part, but of $\Delta\eta = 0.2$ for the photon part due to lower statistic [148].

The E_1/E_2 distribution is reduced to its mean, computed in a truncated range (typically $[0, 2]$) to reduce the bias of long tails induced by background events. Once the mean of the E_1/E_2 distribution (later simply denoted E_1/E_2) is performed in both data and MC, two quantities are needed to convert the E_1/E_2 ratio into an actual measurement of material difference. First, the difference of E_1/E_2 in data and MC, noted

$$\Delta E_{1/2}^{\text{data}} = \frac{E_{1/2}^{\text{data}} - E_{1/2}^{\text{nom}}}{E_{1/2}^{\text{nom}}}, \quad (5.4)$$

and second, a *sensitivity factor* relating a variation of $E_{1/2}$ to a variation of material defined as

$$\frac{\delta X/X_0}{\delta_{\text{rel}} E_{1/2}}, \quad (5.5)$$

with $\delta_{\text{rel}} E_{1/2} = \frac{E_{1/2}^{\text{dist}} - E_{1/2}^{\text{nom}}}{E_{1/2}^{\text{nom}}}$, where the “dist” label denotes the quantity measured in a MC sample with distorted geometry (see Section 5.4). The sensitivities from the various distorted geometries are averaged to provide the final sensitivity factor [148]. The E_1/E_2 differences and sensitivity factors are evaluated separately for electrons and unconverted photons in each $|\eta|$ bin of the analysis. Example plots for the two quantities computed using electrons are given in Figure 5.7.

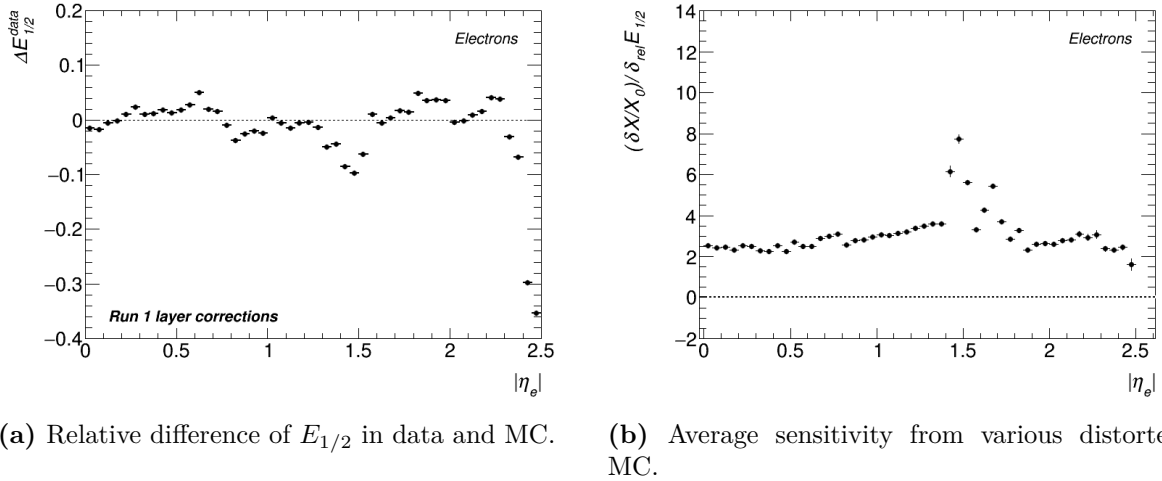


Figure 5.7 — Relative difference of $\Delta E_{1/2}^{\text{data}}$ in 2015+2016 data after Run 1 layer corrections applied (Run 2 were not available yet at that time) and MC with nominal geometry (a), and (b) weighted-average sensitivity (as defined in Eq. 5.5) over various MC with distorted geometries, as a function of $|\eta|$. [148]

The two equations 5.4 and 5.5 are then combined to estimate the material difference between the data and the nominal simulation:

$$\Delta X/X_0 = \Delta E_{1/2}^{\text{data}} \times \frac{\delta X/X_0}{\delta_{\text{rel}} E_{1/2}}. \quad (5.6)$$

5.6.2 Results and validation

The final material discrepancy estimation is shown in Figure 5.8. A value above 0 is interpreted as missing material in the nominal geometry, while areas below 0 are interpreted as an overestimate of the material in the simulation.

The total material estimate (shown in red points) shows a slight deficit in the simulation of about $0.15 X_0$ around $|\eta| = 0.6$ and a similar deficit between in $|\eta| \in [0.8, 1]$. In the crack region 1.3-1.6, the simulation overestimates the material by up to $0.8 X_0$ while in the endcap, a moderate deficit of material in the simulation up to $|\eta| = 2.3$ and a larger excess after are seen.

Between the PS and the accordion (difference of red and blue points), a slight deficit for $|\eta| < 0.8$, a slight excess for $0.8 < |\eta| < 1.37$, and a larger larger deficit for $1.5 < |\eta| < 1.8$ are observed. Due to the low statistics of the photon sample used to determine the material only in front of the PS, the results come with a large uncertainties and the material difference between the PS and the accordion is compatible with 0.

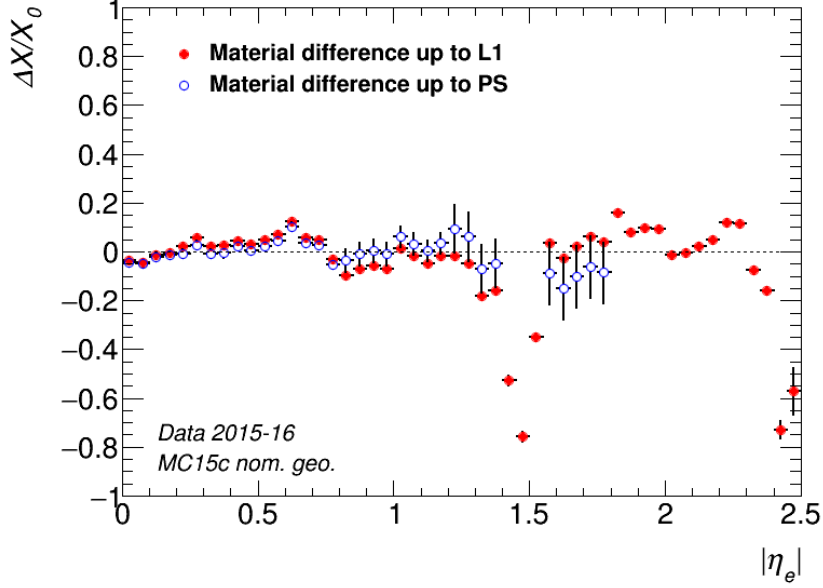


Figure 5.8 – Difference of material estimate up to the presampler (blue) and up to the first layer of the calorimeter (red), in the early Run 2 nominal simulation including (not improved) IBL and in the data. [148]

The systematic uncertainties associated to the results are roughly the same as in Run 1 as shown in Figure 5.9. These systematic uncertainties are dominated by the liquid argon modelling in the GEANT4 simulation, assessed by varying the models of electron scattering, bremsstrahlung and photon conversion, electrode geometry and cross-talk effects [149]. An additional Run 2 systematic uncertainty is added in the endcap, where additional material from the IBL and pixel services is badly accounted for in the early Run 2 simulation. Thus an additional systematic uncertainty was added in the endcap which is likely to be recovered using the improved Run 2 geometry (see Section 5.6.3).

As a final cross-check, the study is also performed as a function of ϕ on top of η using the high statistics electron samples. No large variation along ϕ was found, except for a periodic deficit of material in the simulation at $|\eta| = 0.6$. This effect is interpreted as a missing ID structure element in the simulation.

5.6.3 Nominal geometry for Run 2

The nominal geometry used in this study is the same as the one developed at the end of Run 1, with a retuning of the material in the transition region (crack) and with the addition of a bare material due to the installation of the IBL at the start of Run 2 (denoted *improved Run 1 plus IBL*). The material budget in the improved Run 1 geometry (without IBL) is shown in Figure 5.10,

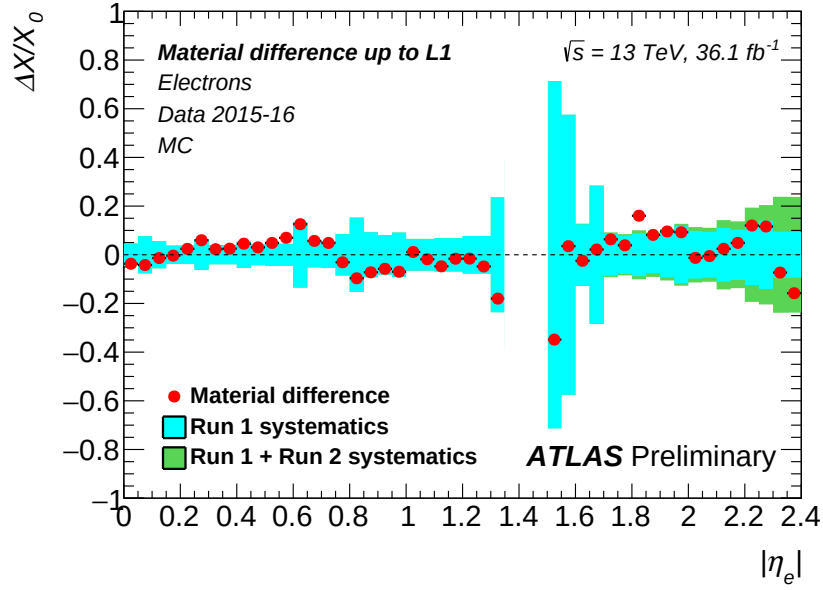
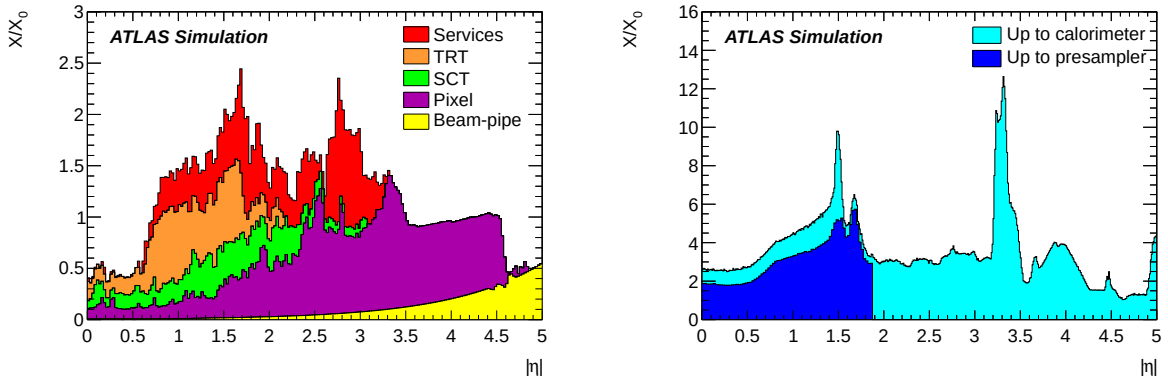


Figure 5.9 – Total material difference between data and nominal early Run 2 simulation as well as its associated systematic uncertainty. [150]

and expected modifications to the Inner Detector material due to the installation of the IBL is shown in Figure 5.11.

The first IBL material estimate was found inaccurate (see above) and a new Run 2-improved geometry was therefore prepared with a better IBL material modelling, especially in the endcap [112]. This improved-Run 2 geometry is the one used for the physics analyses described in Chapters 8 and 9.



(a) Material from the Inner Detector.

(b) Material up to the calorimeter and pre-sampler.

Figure 5.10 – Material budget in units of X_0 from the improved Run 1 geometry, as a function of $|\eta|$. These plots do not include the additional budget due to the insertion of the IBL in 2015. The material between the ID and the PS includes *e.g* the solenoid and the wall of the cryostat. Since the PS stops at about $|\eta| = 1.8$, no estimation of the material up to the PS can be done after this boundary. [140]

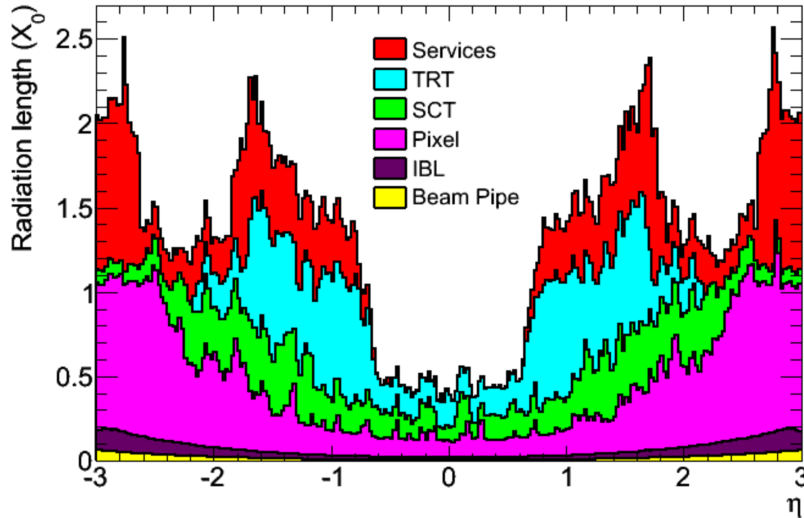


Figure 5.11 – Material budget in the Inner Detector in units of X_0 as a function of $|\eta|$ after installation of the IBL (initial Run 2 = improved Run 1 + IBL geometry). [151]

5.7 $Z \rightarrow ee$ energy scales

5.7.1 Overview

Once all the previously described corrections have been applied to the data, small differences still remain between the data and MC, as seen in Figure 5.12 for the $Z \rightarrow ee$ lineshape. These residual differences can come from effects difficult to model (such as the OFC optimisation impact, or the slightly varying LAr temperature), or corrections that do not conserve the total energy of the shower such as the layer intercalibration. This remaining bias in the average reconstructed energy must be corrected to achieve the permil-level accuracy required in precision physics analyses. Moreover it is found that the resolution in data is worse than in MC, and the latter is therefore smeared in order to improve the data/MC agreement.

The method was first introduced in Ref. [114] and later developed in Ref. [153]. It has been implemented for the Run 1 calibration [140] and updated for the Run 2 [141]. To achieve such precision, a high statistics sample of a well-known process is needed. Such needs are fulfilled by the $Z \rightarrow ee$ resonance [145] whose dataset for the 2015 and 2016 runs of the LHC contains about 20 million events after selection. Events with two opposite sign electrons falling in the mass window [75, 105] GeV are thus selected. Electrons are further required to pass the medium identification criteria and to be loosely isolated to reduce QCD backgrounds. This provide a clean control sample (with less than 1 % of background) of electrons with transverse momentum typically around $p_T \approx 40$ GeV.

As the simulation is an ideal model of the detector free from misalignment, gain, material or other effects, the energies and position in the MC are perfectly calibrated and are used as a reference. The method thus consists in comparing the $Z \rightarrow ee$ lineshape in both data and simulation to derive the correction coefficients needed to make the two distributions better agree.

Given two electrons, the reconstructed invariant mass of the electron pair (the electron mass is negligible) is computed as

$$M = \sqrt{2E_1E_2(1 - \cos(\theta_{12}))}, \quad (5.7)$$

where E_1 and E_2 are the energies of the two selected electrons, and θ_{12} the opening angle between the two electrons. We see that the reconstructed mass only depends on the energy and position

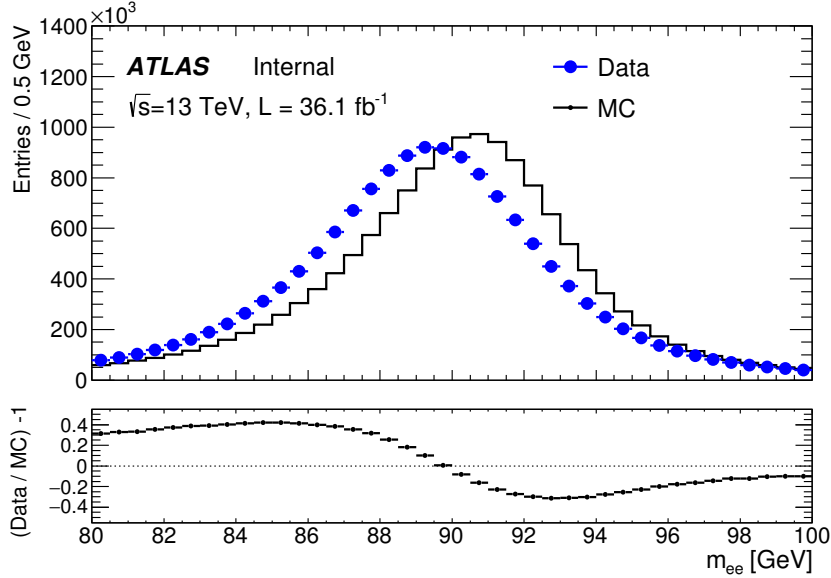


Figure 5.12 – Di-electron invariant mass distribution from $Z \rightarrow ee$ events in data before correction (blue) and in MC before smearing (black). The data distribution is shifted with respect to the MC one, and the resolution in the data is worse than in MC. [152]

(through θ_{12}) of the electrons. The position measurement mainly comes from the tracker which is independently calibrated, so supposing that the tracking measurements are perfectly calibrated, we can assume that the reconstructed invariant mass only depends on the electron energies which are measured by the calorimeter. Adjusting the invariant mass in data to match the one in the simulation therefore amounts to scaling the energy of the electrons. We can thus define an energy scale factor α that must be applied to the measured energy of an electron to get its “true” energy

$$E^{\text{meas}} = E^{\text{true}}(1 + \alpha). \quad (5.8)$$

In practice, the energy scale depends on its phase-space (momentum, position in the detector, *etc.*) so we denote α_i the scale factor in the phase-space of index i instead of just α .

If we plug Eq. 5.8 into Eq. 5.7, we obtain:

$$\begin{aligned} M^{\text{meas}} &= \sqrt{2E_1^{\text{meas}}E_2^{\text{meas}}(1 - \cos(\theta_{12}))} \\ &= \sqrt{2E_1^{\text{true}}(1 + \alpha_i)E_2^{\text{true}}(1 + \alpha_j)(1 - \cos(\theta_{12}))} \\ &= \sqrt{2E_1^{\text{true}}(1 + \alpha_i)E_2^{\text{true}}(1 + \alpha_j)(1 - \cos(\theta_{12}))} \\ &= \sqrt{2E_1^{\text{true}}E_2^{\text{true}}(1 - \cos(\theta_{12}))} \sqrt{1 + \alpha_i + \alpha_j + \alpha_i\alpha_j} \end{aligned}$$

and expanding the square root at leading order:

$$\approx M^{\text{true}} \cdot \left(1 + \frac{\alpha_i + \alpha_j}{2}\right). \quad (5.9)$$

Since the invariant mass depends on the energy of both electrons, α_i and α_j will always be measured simultaneously for a given configuration and we will rather denote

$$\alpha_{ij} = \frac{\alpha_i + \alpha_j}{2}. \quad (5.10)$$

We have therefore:

$$M^{\text{true}} = \frac{M^{\text{meas}}}{1 + \alpha_{ij}}. \quad (5.11)$$

Moreover, on top of the energy scale needed to shift the reconstructed invariant mass, another factor that modifies the energy resolution (*i.e.* the width of the invariant mass shape) is needed to make the simulation match the data. Despite best efforts to make the simulation as close as possible to the real detector, residual inhomogeneities or imperfect modelling (even after previously mentioned corrections) remain. The effect is to degrade the energy resolution of the calorimeter seen in the data, compared to the MC. The energy resolution of a calorimeter can be modelled using the following formula

$$\frac{\sigma(E)}{E} = \frac{a}{\sqrt{E}} \oplus \frac{b}{E} \oplus c, \quad (5.12)$$

where the $\oplus$ operator means a quadrature sum. The a factor represents the sampling (or stochastic) term, linked to the development of the shower in the absorbers (10 % in ATLAS), the b factor is the noise term, from the electronics chain used to read out the signal, and the c factor is the constant term, due to *e.g.* inhomogeneities, material effects *etc.* The detector design is to keep the latter below 1 % which is the purpose of this calibration procedure. The sampling term is about 10 % $\sqrt{\text{GeV}}$ with a relative uncertainty known to be about 10 % from test beam [154]. The noise term is measured to be of the order of 300 MeV in ATLAS with a subdominant effect at energies relevant for the $Z \rightarrow ee$ decays. The a and b factors are assumed to be the same in data and simulation.

As previously mentioned, since the simulation is generated with a better resolution than the one found in the data, the MC resolution is smeared by so-called *additional constant terms*, denoted c' . Similarly to the scale factors α , the c' depend on the phase-space region i and are defined as

$$\left(\frac{\sigma(E)}{E}\right)_i^{\text{data}} = \left(\frac{\sigma(E)}{E}\right)_i^{\text{MC}} \oplus c'_i, \quad (5.13)$$

which will also be extracted for both electrons of the $Z \rightarrow ee$ decays at the same time and therefore combined in two phase-space regions. Therefore we have

$$\left(\frac{\sigma(M)}{M}\right)_{ij}^{\text{data}} = \left(\frac{\sigma(M)}{M}\right)_{ij}^{\text{MC}} \oplus c'_{ij} \quad (5.14)$$

on the one hand, and on the other hand

$$\begin{aligned} \left(\frac{\sigma(M)}{M}\right)_{ij}^{\text{data}} &= \frac{1}{2} \left[\left(\frac{\sigma(E)}{E}\right)_i^{\text{data}} \oplus \left(\frac{\sigma(E)}{E}\right)_j^{\text{data}} \right] \\ &= \frac{1}{2} \left[\left(\frac{\sigma(E)}{E}\right)_i^{\text{MC}} \oplus c'_i \oplus \left(\frac{\sigma(E)}{E}\right)_j^{\text{MC}} \oplus c'_j \right] \\ &= \left(\frac{\sigma(M)}{M}\right)_{ij}^{\text{MC}} \oplus \frac{c'_i \oplus c'_j}{2}, \end{aligned} \quad (5.15)$$

hence

$$c_{ij}^{\prime 2} = \frac{c_i^{\prime 2} + c_j^{\prime 2}}{2}. \quad (5.16)$$

Assuming a gaussian resolution dependency, the c'_i are then used to smear the electron energy in the MC as $E_i^{\text{MC}} \left(1 + c'_i \times \mathcal{N}(0, 1)\right)$, where $\mathcal{N}(0, 1)$ is a random number drawn from a gaussian distribution of mean 0 and RMS 1. As such, the impact of the resolution correction cannot be seen from a single electron, but only on the distribution of some physical process before and after correction.

Definition of the phase-space regions: the energy scales α_i are extracted in 68 bins along the (signed) η direction (measured in the calorimeter), as a tradeoff between statistical requirement to achieve a sufficient precision and fine coverage of detector to catch inhomogeneities along this direction. On the other hand, the additional constant terms are extracted in 24 symmetrised bins along $|\eta_{\text{calo}}|$ to maximize the statistics in each configuration. Since no distinction is made between the two electrons in the decay, the configurations $(\eta_{\text{calo},i}, \eta_{\text{calo},j})$ and $(\eta_{\text{calo},j}, \eta_{\text{calo},i})$ are the same. This is summarised in Table 5.2

Table 5.2 – η binning used for the determination of the energy scale factors α_i and the additional constant terms c'_i in the in-situ calibration analysis. The c'_i are extracted in bins of $|\eta|$ (*i.e.* symmetrised for $\eta < 0$ and $\eta > 0$) while the α_i are extracted in bins of signed η . [152]

Barrel and crack																	
α_i	0	0.1	0.2	0.3	0.4	0.5	0.6	0.7	0.8	0.9	1	1.1	1.2	1.285	1.37	1.42	1.47
c'_i	0		0.2		0.4		0.6		0.8		1		1.2		1.37		1.55
Endcap																	
α_i	1.55	1.59	1.63	1.6775	1.725	1.7625	1.8	1.9	2	2.05	2.1	2.2	2.3	2.35	2.4	2.435	2.47
c'_i	1.55						1.8		2				2.3				2.47

5.7.2 Estimation of α_{ij} and c'_{ij} using the template method

The nominal method to find the α_i and c'_i in a given phase-space bin is to compare histograms of the distribution of di-electron invariant mass in data to the same distribution in MC after scaling and smearing its energy. The method acts as follows: given a configuration (i, j) , the distribution of di-electron invariant using electrons corrected with α_{ij} and c'_{ij} is generated such as

$$M_{ee}^{\text{corr}} = M_{ee}^{\text{MC}} \cdot \sqrt{(1 + \alpha_i) \left(1 + c'_i \times \mathcal{N}(0, 1)\right) (1 + \alpha_j) \left(1 + c'_j \times \mathcal{N}(0, 1)\right)}. \quad (5.17)$$

The likelihood between this generated *template* distribution and the data is then estimated with a χ^2 fit. This amounts to scanning the (α_{ij}, c'_{ij}) space for each configuration (i, j) . We can therefore denote $\chi^2(\alpha_{ij}, c'_{ij})$. The χ^2 value as a function of the correction applied in the template for a given configuration is shown in Figure 5.13.

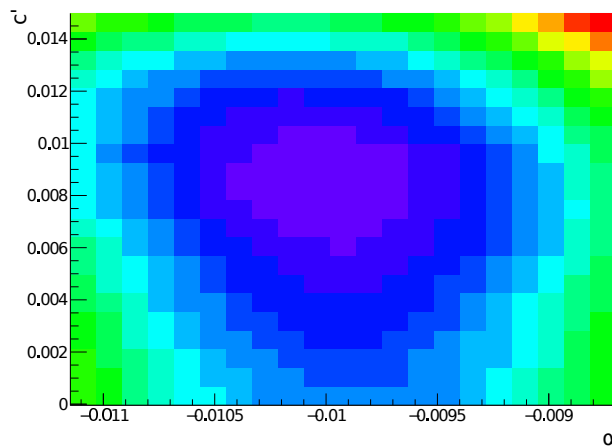


Figure 5.13 – Example of χ^2 compatibility between generated template and data for M_{ee} distribution as a function of the scale factor α_{ij} and the additional constant term c'_{ij} . [152]

Although the most straightforward way would be to perform a 2D fit over the α_{ij} and c'_{ij} simultaneously, this method was shown to give unstable results [155] and has been replaced by two 1D fits over α_{ij} and c'_{ij} successively. The exact procedure is thus as follows: first, compute

the χ^2 as a function of α_{ij} for a fixed value of c'_{ij} , *i.e.* scan along a row of Figure 5.13. This χ^2 distribution is parametrised as

$$\chi^2(\alpha_{ij}, c'_{ij}) = \chi^2_{\min}(c'_{ij}) + \frac{(\alpha_{ij} - \alpha_{ij,\min}(c'_{ij}))^2}{(\delta\alpha_{ij}(c'_{ij}))^2}, \quad (5.18)$$

where $\delta\alpha_{ij}(c'_{ij})$ is the error on the determination of $\alpha_{ij}(c'_{ij})$ defined by $\Delta\chi^2 = 1$. Second, once all scan has been done for each c'_{ij} (each line), the distribution of $\chi^2_{\min}(c'_{ij})$ as a function of c'_{ij} is fitted with a third order polynomial. The minimum of this distribution determines the best value of c'_{ij} , denoted $\widehat{c'_{ij}}$, and its uncertainty is also defined by $\Delta\chi^2 = 1$. Finally, $\alpha_{ij,\min}(c'_{ij})$ is plotted as a function of $\widehat{c'_{ij}}$ and a linear fit is performed around $\widehat{c'_{ij}}$. The best α_{ij} value $\widehat{\alpha_{ij}}$ is the one corresponding to $\widehat{c'_{ij}}$ in the linear fit. Each of these three steps is illustrated in Figure 5.14.

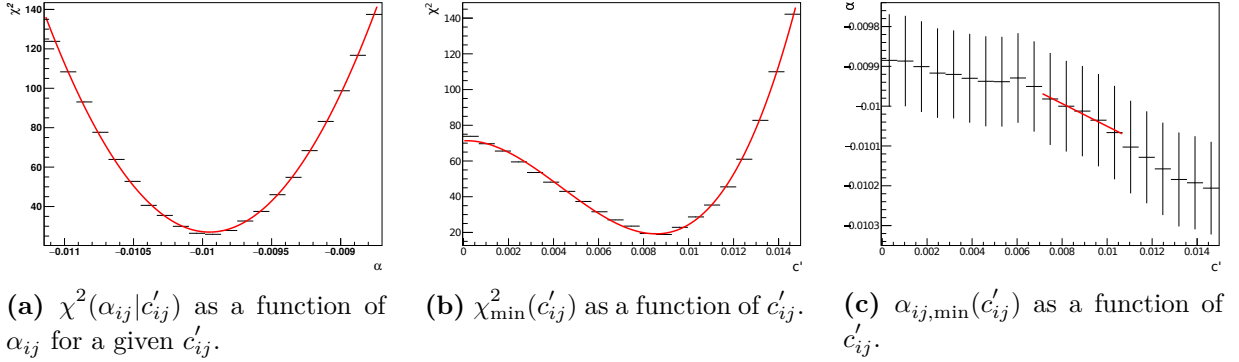


Figure 5.14 – Illustration of the various steps in the determination of $\widehat{\alpha_{ij}}$ and $\widehat{c'_{ij}}$ for a given (i, j) configuration. Step 1 (a): determination of $\chi^2_{\min}(c'_{ij})$ and $\alpha_{ij,\min}(c'_{ij})$ for a given c'_{ij} . Step 2 (b): determination of $\widehat{c'_{ij}}$. Step 3 (c): determination of $\widehat{\alpha_{ij}}$. [152]

Since both electrons are modified when the template is created, only the α_{ij} and c'_{ij} are accessible. To recover the single phase-space corrections α_i and c'_i , the information from the effective $Z \rightarrow ee$ corrections α_{ij} and c'_{ij} must be combined. A χ^2 quantity is built separately for α and c' as described in Eqs. 5.19 and 5.20, and are minimised against the $\{\alpha_i\}$ or $\{c'_i\}$:

$$\chi^2_{\alpha} = \sum_{i,j \leq i} \frac{\left(\frac{\alpha_i + \alpha_j}{2} - \alpha_{ij}\right)^2}{(\delta\alpha_{ij})^2}, \quad (5.19)$$

$$\chi^2_{c'} = \sum_{i,j \leq i} \frac{\left(\sqrt{\frac{c'^2_i + c'^2_j}{2}} - c'_{ij}\right)^2}{(\delta c'_{ij})^2}. \quad (5.20)$$

The uncertainties are here taken using $\Delta\chi^2 = 1$.

In some configurations the number of events is very low and does not allow for a reliable determination of the α_{ij} and c'_{ij} , leading to instabilities in the inversion procedure. As long as a limited amount of such configuration exist, they can be safely removed from the χ^2 minimisations.

5.7.3 Estimation of α_i and c'_i using the Z lineshape method

This second method, used as a cross-check, performs an analytical fit of the di-electron invariant mass distribution in each configuration, instead of generating many templates with various

injected corrections which is time-consuming. This method also allows to avoid the complicated inversion step of the template method, since the single scale factors and additional constants terms can be fitted directly. The di-electron invariant mass is modelled with an empiric shape composed of three Gaussians:

$$\text{Model}(m_{12}) = n_1 \cdot \mathcal{G}_1(m_{12}|\mu_1, \sigma_1) + n_2 \cdot \mathcal{G}_2(m_{12}|\mu_2, \sigma_2) + n_3 \cdot \mathcal{G}_3(m_{12}|\mu_3, \sigma_3) \quad (5.21)$$

and their parameters μ_k and σ_k fitted independently in each configuration. The scale factors α_i and additional constant terms c'_i are then expressed as a function of the ones in MC: for $k = 1, 2, 3$,

$$\mu_k^{\text{data}} = \mu_k^{\text{MC}} \cdot \sqrt{(1 + \alpha_i)(1 + \alpha_j)} \cdot (1 - \Delta_{\text{bias}}(c'_i, c'_j)) \quad (5.22)$$

and

$$(\sigma_k^{\text{data}})^2 = (1 + \alpha_i)(1 + \alpha_j) \cdot \left[(\sigma_k^{\text{MC}})^2 + \frac{(\mu_k^{\text{MC}})^2}{4} (c_i'^2 + c_j'^2) \right]. \quad (5.23)$$

The data distributions are then simultaneously fitted in all configurations against the $\{\alpha_i\}$ and $\{c'_i\}$, fixing all μ_k and σ_k to their values fitted in MC. The term $(1 - \Delta_{\text{bias}}(c'_i, c'_j))$ in Eq. 5.22 corrects for non-gaussianity of the smearing [152]. The sum of three gaussian was empirically chosen instead of a natural Breit-Wigner * Crystal-ball convolution due to the difficulty to tune the parameters of the Crystal-ball over a large range of configurations. Furthermore, the invariant mass distribution can be distorted when restricting the phase space (for example binning in p_T) adding even more tuning.

The two methods give compatible results, and the difference is assigned a systematic on the final result.

5.7.4 $Z \rightarrow ee$ scale results and systematics

The results presented in this section include only the template method, since the lineshape method is only used as a cross-check.

The energy scale factors and additional constant terms derived independently for each year of the Run 2: despite the calorimeter not changing in the course of Run 2, the α and c' are derived separately for every year to account for variations of the running condition of the LHC (see Chapter 2.10). As the luminosity increases, more particles fly through the calorimeter depositing more energy, possibly changing the energy scales. The consequences are two-fold: first an increase of the LAr temperature, which changes the energy response by $\approx -2\%/K$ [156]. Second, with higher amounts of energy deposited in the liquid argon gaps, larger currents I are created in the HV lines. Since the HV passes through a large resistance R between the power supply and the gaps, the effective HV is reduced by $R \cdot I$, reducing the drift velocity of the electrons in the liquid argon, and thus the energy response [157, 158].

For the resolution term, the values is found to decrease as a function of the year. Investigations showed that this effect is due to pileup noise mismodelling in the MC: while the electronic noise is constant, the pileup noise increases linearly with $\sqrt{\langle \mu \rangle}$. This behaviour is however not well modelled in the simulation, which predicts a larger pileup noise than what is seen in the data [132]. Since the operating instantaneous luminosity, and therefore the pileup, is higher in 2017 and 2018 than in 2016, itself higher than in 2015, the resolution decreases more in the MC than in data, which is absorbed by a smaller additional constant term. In 2017, the effect is even more important than in 2018 due to a part of the run using the 8b4e filling scheme, inducing more pileup for the same luminosity. For sake of simplicity, the additional constant term provided to external analyses is however combined using a weighted average of all years. These results are displayed in Figure 5.15.

A dozen sources of systematics have been identified [152, 155, 159], ranging from experimental factors such as reconstruction, identification, isolation and trigger efficiencies, to electron quality, bremsstrahlung veto or not requiring isolated electrons, and to methodological systematics such as the M_{ee} range used for the fit, the impact of QCD and EW background (neglected in the nominal procedure, below 1 % in the mass range), the impact of removing “bad” configurations (low number of events, *etc.*), or a possible bias of the method [155]. Finally a systematic is assigned to take into account the difference between the template and the lineshape method. The systematics are dominated by the choice of the M_{ee} mass window, and subleading contributions include the method difference (template vs. lineshape), the electron identification criterion, and the bremsstrahlung veto. The systematics are symmetrised in $|\eta|$.

After scaling the electron energies with the α in the data and smearing the electron energies in the MC with the c' , the $Z \rightarrow ee$ lineshape as presented in Figure 5.12 can be compared to check the improvement. This is shown in Figure 5.16: the residual differences amount to less than 2 % under the peak and up to 4 % in the tails of the $Z \rightarrow ee$ distribution.

Stability versus pileup: the dependency of the $Z \rightarrow ee$ invariant mass to the pileup $\langle\mu\rangle$ was also checked to assess the calibrations are correct as a function of $\langle\mu\rangle$. Such dependency is shown in Figure 5.17, where it can be seen that the variation of the $Z \rightarrow ee$ invariant mass as a function of the $\langle\mu\rangle$ is remarkably stable, below the 0.05 % level. The small slope seen in both data and simulation is due to the dynamical clustering algorithm (see Chapter 4.4) [161].

5.8 Energy scale and resolution extrapolation over the energy range

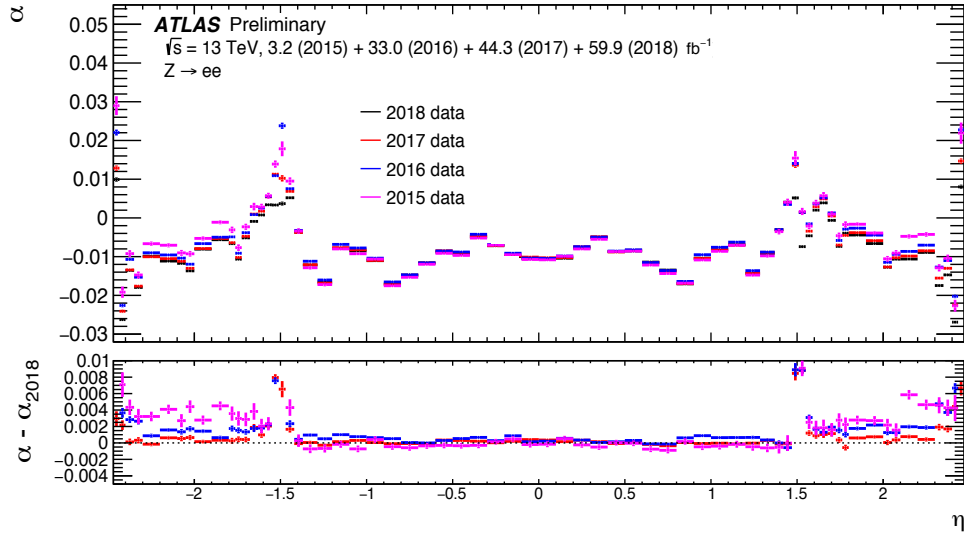
As described in the previous Section, energy scales and additional constant terms are only derived as a function of η and are integrated over the energy. Since the method is using electrons from $Z \rightarrow ee$ decays, the typical energy range probed is $E_T \approx 40 - 45$ GeV. By construction, the $Z \rightarrow ee$ scale factors absorb all other systematic uncertainties at $E_T \approx 40 - 45$ GeV, leaving only its own uncertainty as systematic [141]. To be used in other energy domains, the energy scales and resolution corrections must be extrapolated to lower and higher energies. This extrapolation is tainted with some uncertainties depending on η and the transverse energy range considered. Systematic uncertainties mainly come from the layer intercalibration and its extrapolation to electrons instead of muons, from material, and from gain systematics.

The energy scale extrapolation uncertainty for electrons ranges from a few permil at $E_T = 40$ GeV, to less than 1 % at the extremes of the transverse energy extrapolation range ([5, 200] GeV). Examples in a few bins of $|\eta|$ are given in Figure 5.18. All the uncertainties are considered uncorrelated and are summed in quadrature to give the total uncertainty. They are dominated

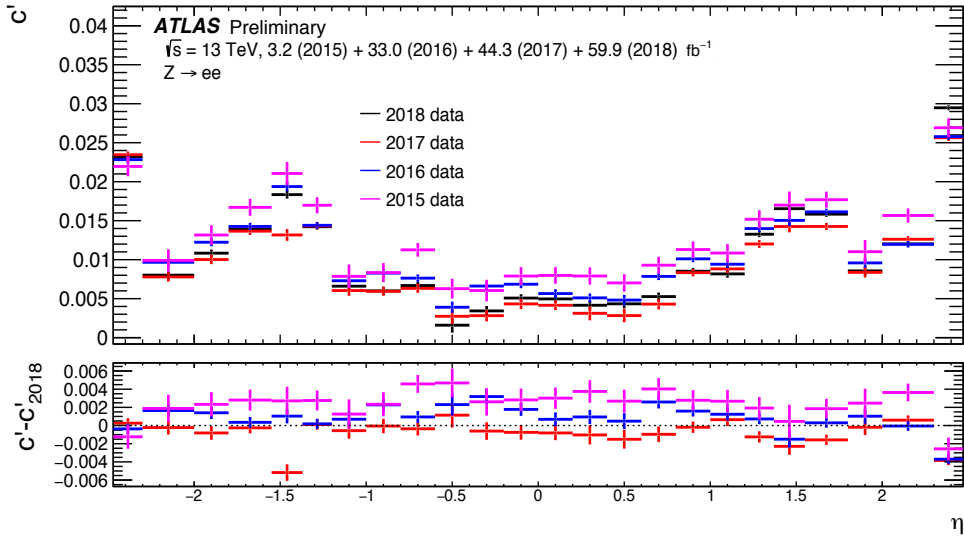
- in the barrel, by the gain for electrons above 45 GeV, and by the material for electron below 45 GeV,
- in the endcap, by various material contributions for all energies, with a non-negligible impact of the intercalibration scale α_{12} below 45 GeV.

On the other hand, the resolution extrapolation uncertainty for electrons ranges from 0.05 % at low transverse energy to 1.5–2.5 % at ~ 200 GeV, depending on $|\eta|$. Examples in a few bins of $|\eta|$ are given in Figure 5.19. The uncertainty is dominated

- in the barrel, by the sampling term uncertainty and the material,



(a) Scale factor.



(b) Additional constant term.

Figure 5.15 – Scale factors α (a) and additional constant term c' (b) as a function of η in the various years of Run 2 data taking. The α are extracted in 68 η bins while the c' are extracted in 24 bins of $|\eta|$. [160]

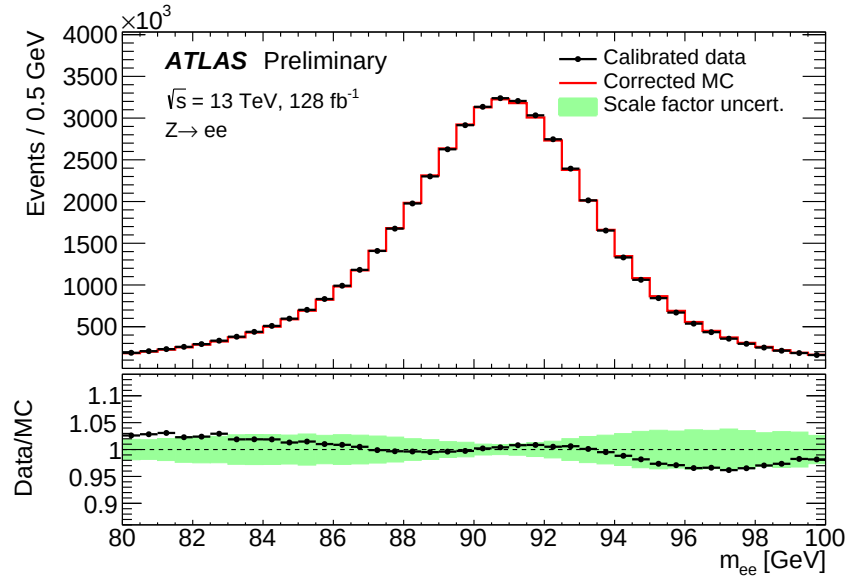


Figure 5.16 – Comparison of the di-electron invariant mass under the $Z \rightarrow ee$ peak in data (black) and simulation (red) after scaling the data and smearing the MC. The agreement after correction is within the uncertainty band and exhibits yet unexplained features. [161]

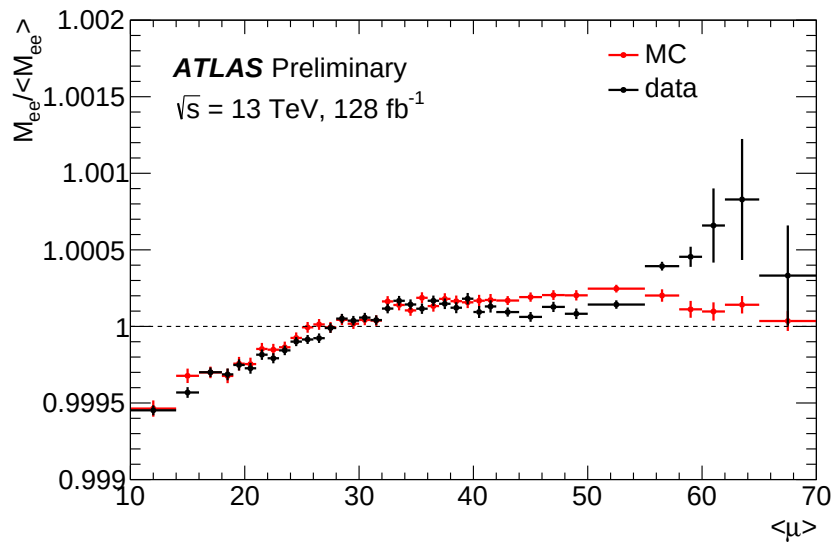


Figure 5.17 – Di-electron invariant mass per slice of $\langle \mu \rangle$ normalised to the average over Run 2 (using data up to 18 September 2018). [161]

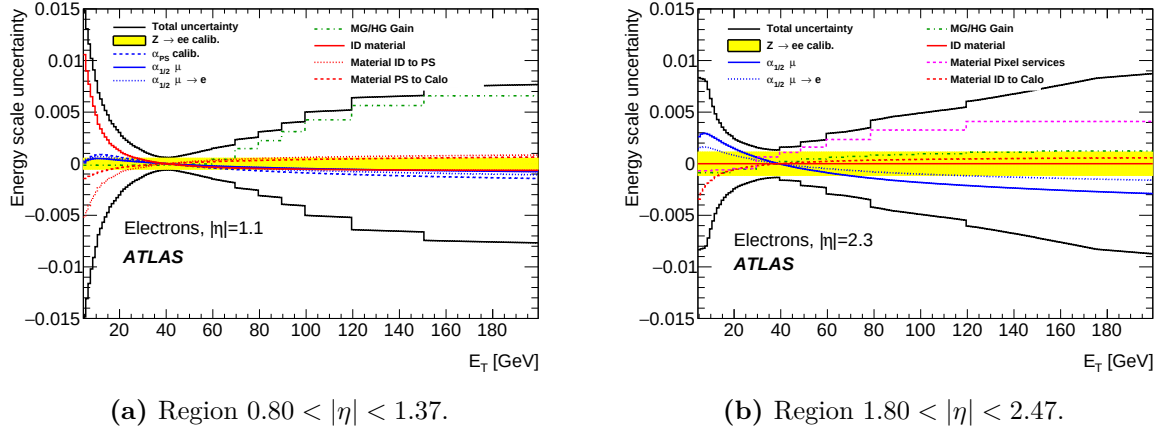


Figure 5.18 – Extrapolation uncertainty of the energy scale factors α for electrons as a function of their transverse energy E_T (2015+2016 calibration analysis). [141]

- in the endcap, by the sampling term and the intrinsic $Z \rightarrow ee$ procedure uncertainty.

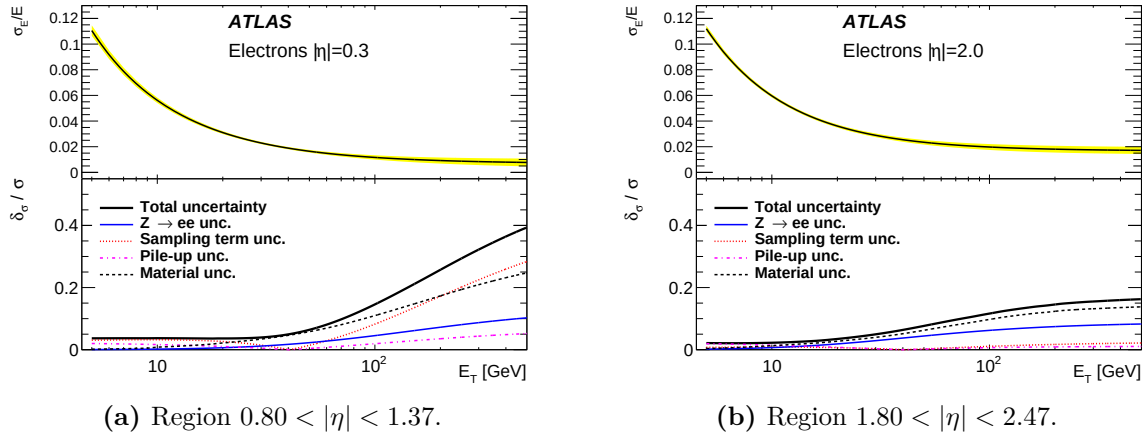


Figure 5.19 – Extrapolation uncertainty of the additional constant term for electrons as a function of their transverse energy E_T (2015+2016 calibration analysis). [141]

The extrapolations are expected to be valid up to $E_T \approx 500$ GeV [140]: above this value, electrons starts to have a larger deposit in the third sampling and a significant fraction of showers are recorded with at least one cell in low gain. As most electrons used in the $H \rightarrow 4\ell$ analysis (Chapters 8 and 9) have a transverse energy below 100-150 GeV, the case of such very high E_T electrons will not be discussed here.

On the other side of the E_T range, the extrapolation at low transverse energy can be cross-checked using another well-known decay: $J/\psi \rightarrow e^+e^-$. These cross-checks are described in the next Section (5.9).

Remark about photon scales: the scales and resolutions corrections are also extrapolated to photons of various energies, separately for converted and unconverted. The extrapolation to photons include specific systematics related to the shower width development or photon conversion. The uncertainty ranges from 0.5% in the barrel (similar to electrons) to 3% in the endcap with a smaller dependency on the energy than for the electrons. The extrapolation from ~ 45 GeV electrons to photons of various energies can also be cross-checked, this time using a control sample of photons from radiative $Z \rightarrow \ell\ell\gamma$ decays. This validation will not be detailed here since photons are not the main objects of the $H \rightarrow 4\ell$ analysis described later in Chapter 8.

5.9 Cross-checks of energy scales and resolution extrapolation

5.9.1 Methodology

The energy scales and resolution corrections described in Section 5.7 are derived for ~ 45 GeV electrons, and have to be extrapolated to other energy domains as seen in Section 5.8. In order to validate the procedure and the resulting extrapolations, the results are validated at lower energy thanks to another well-known resonance in this region: the J/ψ , at a mass of 3.097 GeV [41].

Despite a large J/ψ production cross-section at the LHC, the low E_T of its decay electrons require dedicated triggers that must be strongly prescaled. These triggers have a minimum E_T cut of 5, 9 and 14 GeV, yielding electrons from the $J/\psi \rightarrow e^+e^-$ decay with typically a transverse energy of $E_T \approx 10$ GeV. This part is much more relevant for the $H \rightarrow 4\ell$ analysis since the requirement on the electron minimal E_T is as low as 7 GeV, as will be described in Chapter 8.

Another consequence of the prescaling is that the $J/\psi \rightarrow e^+e^-$ sample is much smaller than the $Z \rightarrow ee$, so the $|\eta|$ binning use is much coarser than for the $Z \rightarrow ee$ analysis (see Table 5.3). This has been chosen to keep statistical uncertainties below 1 % in the barrel [162].

Table 5.3 – Boundaries of the η binning used for the residual energy scales determination with $J/\psi \rightarrow e^+e^-$ electrons.

-2.40 -1.52 -1.37 -1.10 -0.80 -0.40 0 0.40 0.80 1.10 1.37 1.52 2.40

Once all the previously described corrections and in particular the energy scales derived from the $Z \rightarrow ee$ analysis are applied, the residuals at low energy can be extracted. The method is quite similar to the $Z \rightarrow ee$ lineshape fit analysis used in Section 5.7.3. If the calibration procedure is correct, then the residuals should be compatible with 0, within uncertainties from the J/ψ procedure and $Z \rightarrow ee$ systematics.

5.9.2 Determination of the J/ψ prompt fraction

One specificity of J/ψ production with respect to the Z production is that they can be produced directly at the interaction vertex (“prompt” production), or within a b -hadron decay (“non-prompt”). Due to these two different production modes, it is expected that the two contributions have different kinematics distributions, isolation properties, *etc.* Since the MC are generated separately for the prompt and non-prompt part, the first step in this analysis is to estimate the fraction of prompt and non-prompt J/ψ in the data to properly weight the two MC sets and avoid a bias in the kinematical distributions. The original method is described in Ref. [163] and is summarised below.

Starting from the idea that non-prompt J/ψ have a displaced vertex due to the b -hadron flying, one can define a “pseudo proper-time” τ as

$$\tau = \frac{\mathbf{L} \cdot \mathbf{p}_T(J/\psi)}{p_T(J/\psi)} \frac{m(J/\psi)}{p_T(J/\psi)}, \quad (5.24)$$

with $\mathbf{L}_T$ the distance vector between the primary vertex and the J/ψ decay vertex, which is expected to be different for prompt and non-prompt, and therefore can serve as discriminating variable. The pseudo proper-time distribution is described by a Dirac function for the prompt contribution plus an exponential for the non-prompt part, convolved with a resolution term modelled with a sum of three gaussian [162]. An illustration is given in Figure 5.20.

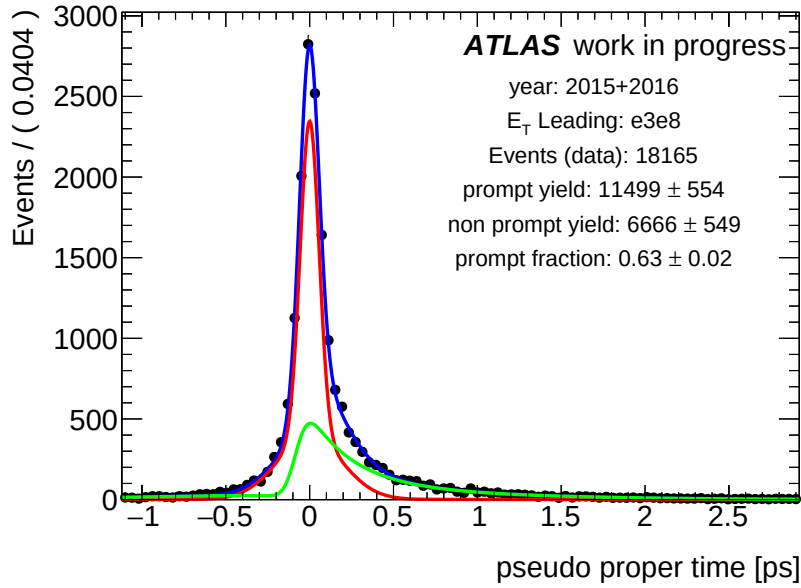


Figure 5.20 – Fit of the pseudo proper-time of J/ψ candidate events in the 2015 and 2016 data. The red curve corresponds to the Dirac function convolved with the resolution (modelling the prompt contribution), and the green curve to the exponential convolved with the resolution (modelling the non-prompt contribution). The total model is shown in blue. [162]

5.9.3 Extraction of residual scales using the lineshape method

Once the prompt fraction is extracted, the rest of the procedure is very similar to the lineshape method from the $Z \rightarrow ee$ analysis: the $J/\psi \rightarrow e^+e^-$ lineshape is fitted in the MC using a double-sided Crystal-Ball (DSCB), independently in each configuration (i, j) . The data model is a bit more complicated in the J/ψ case though: it is built using the sum of two DSCB, one for the J/ψ (around 3.097 GeV [41]) and another for the nearby $\psi(2S)$ peak (around 3.686 GeV [41]), and a second order polynomial to model the background (Drell-Yan processes). The DSCB describing the $\psi(2S)$ is assumed to have the same parameters as the one describing the J/ψ , scaled by a factor of $m(\psi(2S))/m(J/\psi) = 3.686/3.097$. Before fitting the data, the means of the DSCB are expressed as a function of the $\{\alpha_i\}$ in a similar way as for the $Z \rightarrow ee$ (Eq. 5.22):

$$\mu^{\text{data}} = \mu^{\text{MC}} \cdot \sqrt{(1 + \alpha_i)(1 + \alpha_j)}, \quad (5.25)$$

with μ^{MC} fixed. The J/ψ , $\psi(2S)$ and background normalisations, as well as the polynomial parameters describing the continuous background (not inferred from MC) are left free and independent in each data fit. In the end, the data are fitted simultaneously over all the (i, j) configurations, each of them having 7 free parameters (3 normalisations, 2 parameters for the second order polynomial, α_i, α_j). An example of MC and data fit for one configuration is shown in Figure 5.21.

Due to its relatively low mass, the J/ψ production is boosted by a large factor, leading to close-by electrons in the decay (Figure 5.22). As a consequence, the η_i, η_j plane is mainly diagonally populated and many configurations are empty. The retained categories for the simultaneous fit thus include only configurations with more than 600 events to avoid instabilities.

5.9.4 Results and systematics

Several sources of systematic uncertainties have been investigated for the J/ψ residual energy scales measurement, which include

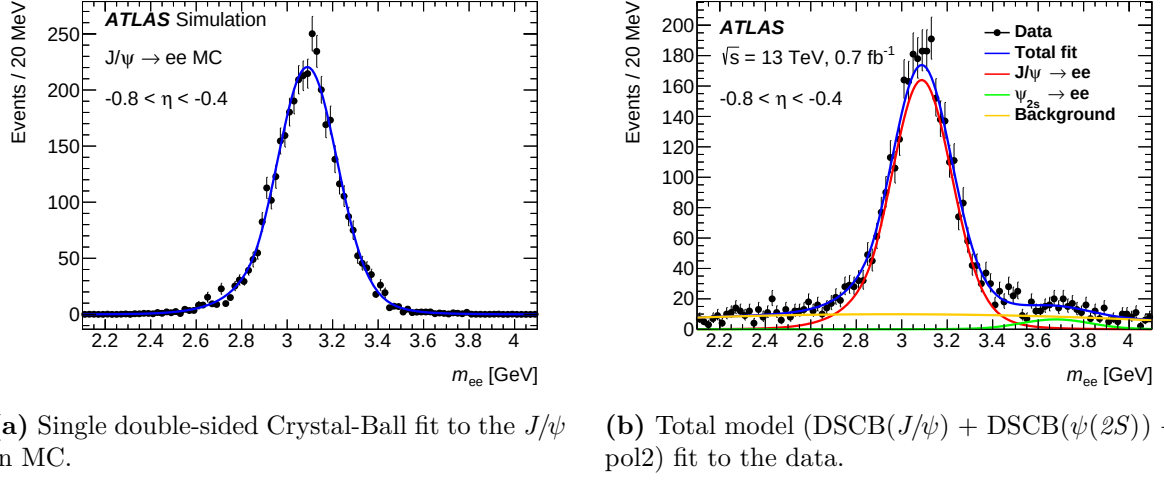


Figure 5.21 – Example fit of the J/ψ shape in the MC (a) and in the data (b) in the configuration $(i, j) = ([-0.8, -0.4], [-0.8, -0.4])$. [141]

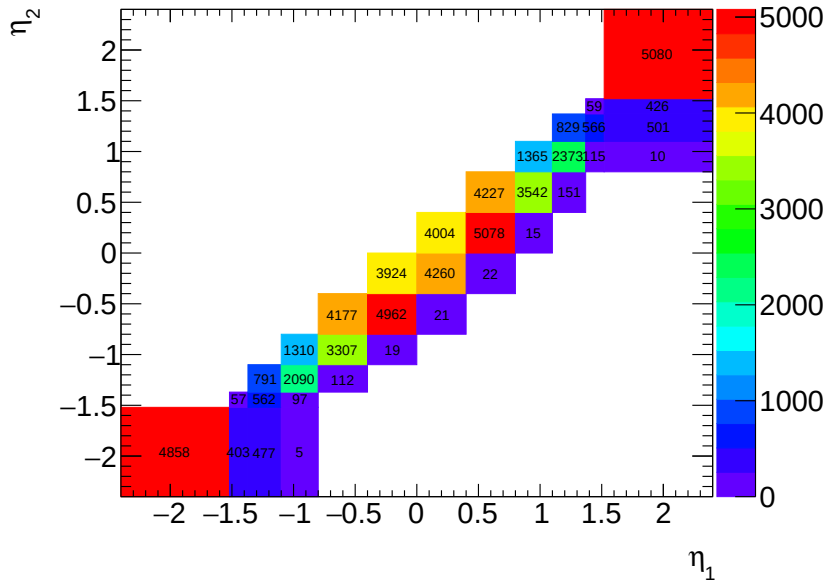


Figure 5.22 – Distribution of electrons from candidate J/ψ decay in the data. $\eta_1 = \max(\eta_{\text{leading}}, \eta_{\text{sub-leading}})$ and $\eta_2 = \min(\eta_{\text{leading}}, \eta_{\text{sub-leading}})$, where “leading” and “sub-leading” refer to the electron E_T ordering. [162]

- the model for the J/ψ distribution (sum of CB and gauss instead of DSCB),
- fixing or freeing the tails parameters of the DSCB,
- the model for the background (exponential instead of second order polynomial),
- the requirement on the minimal number of events to include the configuration in the fit (550 or 650 instead of 600),
- the mass window used in the lineshape fit,
- the modelling of the η distribution of electrons in the MC,
- and pseudo proper-time extraction specific systematics.

The systematics are dominated by the changes of the background and signal model, but the overall result is largely dominated by statistical uncertainty.

The final results are shown in Figure 5.23. The residual scales are all within the uncertainty band of the $Z \rightarrow ee$ energy scale extrapolation at low energies. This confirms the validity of this extrapolation to other energies than the typical 40 GeV from $Z \rightarrow ee$ electrons.

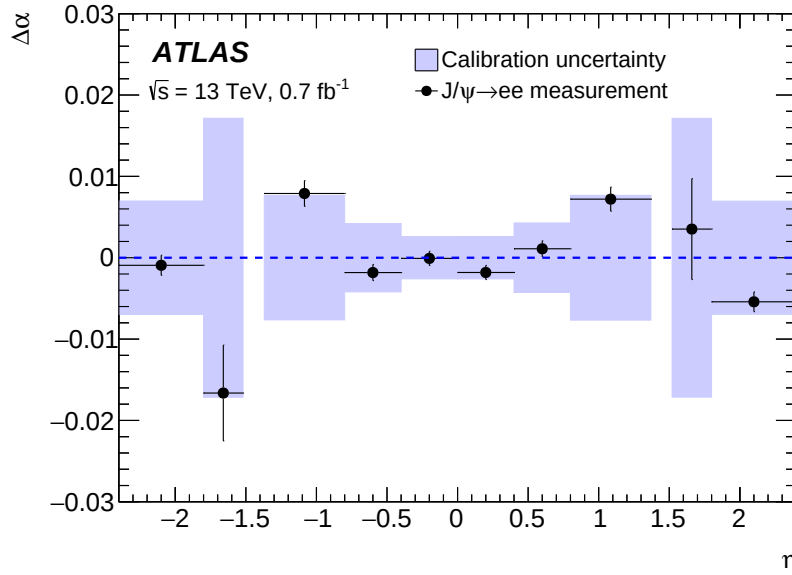


Figure 5.23 – Residual scales $\Delta\alpha$ extracted with the $J/\psi \rightarrow e^+e^-$ method as a function of η , after the $Z \rightarrow ee$ central scales are applied. The black point errors include both statistical and systematics uncertainties. The band shows the uncertainty on the $Z \rightarrow ee$ scale extrapolation to energies relevant for $J/\psi \rightarrow e^+e^-$ electrons. No point is reported in the crack region due to a very poor precision. [141]

ECal first and second layers intercalibration

Contents

6.1	Overview of the layer intercalibration	130
6.2	Geometry considerations: η definition and binning	133
6.2.1	Relation between cell η and track η	133
6.2.2	Calorimeter geometry impact in the barrel for $1.30 < \eta_{\text{corr}} < 1.48$	134
6.2.3	Calorimeter geometry impact in the endcap for $1.375 < \eta_{\text{corr}} < 1.50$. . .	135
6.2.4	Final binning	136
6.3	Number of cells summed	137
6.4	Extraction of the muon energy deposit using an analytical fit	138
6.4.1	Energy distribution fit with Landau * Gaussian convolution	138
6.4.2	Systematic investigation	139
6.4.2.1	Number of cells in S2	139
6.4.2.2	Bunch train effects	139
6.4.2.3	Tighter isolation requirements	141
6.4.2.4	Ambient noise subtraction	143
6.4.2.5	First hint: pileup dependency	145
6.4.3	Noise modelling with templates	148
6.4.4	Extrapolating the layer intercalibration to $\langle\mu\rangle = 0$	148
6.4.5	Results	153
6.4.6	Systematic uncertainties of the MPV method	153
6.4.6.1	Validation of the extrapolated results	155
6.4.6.2	Choice of the second cell in S2	155
6.4.6.3	Leakage energy in S1	157
6.4.6.4	Leakage energy in S2	159
6.4.6.5	Combined plot of MPV systematic uncertainties	161
6.5	Estimation of the muon energy deposit using the TM method	161
6.5.1	Truncated mean definition	161
6.5.2	Truncated mean results	164
6.5.3	Systematics of the truncated mean method	164
6.6	Comparison of muon MPV and TM method, and to Run 1 result . .	165
6.6.1	Comparison of MPV and TM uncertainties in Run 2	166

6.6.2	Final result: Run 2 combination of the MPV and TM methods	168
6.6.3	Run-1 / Run-2 comparison	169
6.7	Cross-check: electron method	169
6.7.1	E_1/E_2 follow up with $Z \rightarrow ee$ electrons	170
6.7.2	Results and compatibility with the muon method	172
	Conclusion of the layer intercalibration analysis	173

This Chapter describes the layer intercalibration of the first and second layers of the ATLAS electromagnetic calorimeter using muons and its cross-check with electrons, with the data collected in 2015 and 2016 at the LHC. Due to a change in the running condition of ATLAS and the LHC between Run 1 and Run 2 (electromagnetic calorimeter readout with 4 instead of 5 samples, different OFC, larger pileup, double the bunch crossing frequency), the method developed for Run 1 did not behave as expected with Run 2 data, and a new method had to be developed. These investigations have been a major task of the present Thesis and are presented in Sections 6.2, 6.3, 6.4 and 6.6,

6.1 Overview of the layer intercalibration

As described in the previous chapter, an accurate description of the ratio of energy deposited in the first and second layers of the electromagnetic calorimeter (ECal) is required for several steps of the electron and photon calibration. These include the training of the MVA energy estimation, the Pre-Sampler correction, and the material determination. Despite best efforts to calibrate the calorimeter electronics [126, 164, 165], imperfect knowledge of the electronics readout and cross-talk effects [119] induce residual miscalibration effects on the reconstructed energy. The purpose of intercalibrating the first and second layers (also called samplings, hence denoted S1 and S2, respectively) is to correct for such residual miscalibrations.

As electrons interact with the upstream detector material, they cannot provide an accurate probe of the layer intercalibration (at least not directly). Muons on the other hand have a behaviour close to a minimum ionising particle (MIP), corresponding to a uniform response throughout the detector without much dependence on the material in front of the calorimeter. Another consequence is that muons, in contrary to electrons, do not produce an electromagnetic shower in the calorimeter, enabling to probe very finely the calorimeter. This method has already been used to probe the energy response of the calorimeter in Run 1, as described in Ref. [166].

The difficulty is that muons also have a small energy deposit in the calorimeter, which only depends on the length of active material traversed (*i.e.* liquid argon): expressed at the same cell-level energy scale as in Eq. 4.2, the average energy deposit is thus typically 30 MeV to 60 MeV in the first sampling and 240 MeV to 300 MeV in the second, depending on $|\eta|$. This should be compared to the average noise in the calorimeter (Figure 4.3), about 15 MeV to 40 MeV in the first sampling and 40 MeV to 80 MeV in the second sampling, leading to a signal-to-noise ratio (SNR) of 0.5 to 2 and 3 to 4 respectively. In such conditions, the difficulty is to reliably estimate the muon energy deposit.

The main need for this analysis thus is to have a clean sample of muons with an homogeneous p_T distribution to fulfil the approximate MIP requirement. The event selection is therefore naturally targeted toward muons from $Z \rightarrow \mu^+ \mu^-$ decays satisfying the following:

- at least two opposite-sign muons in event,
- Medium quality for both muons,

- $80 < m_{\mu\mu} < 105$ GeV, keeping the opposite-sign muon pair with invariant mass closest to the PDG Z mass if several pairs are possible,
- $p_T > 27$ GeV for at least one of the two muons,

For selected events, a muon is used if it passes through at least one cell in both the first and second samplings. Such muon sample ends up with a p_T spectrum of $[27, \sim 100]$ GeV after kinematical cuts and can therefore be considered as MIP (Figure 6.1).

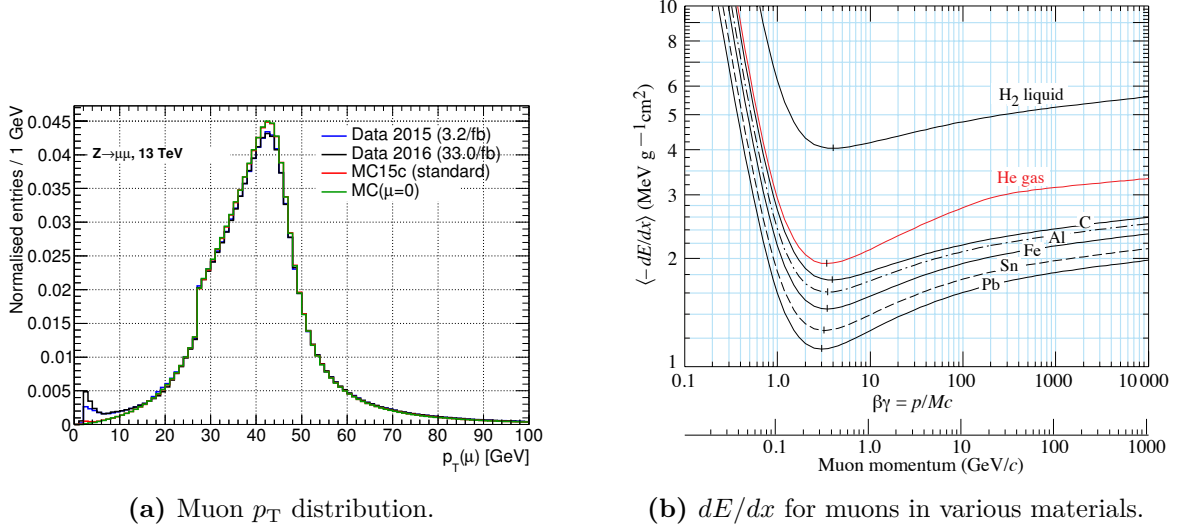


Figure 6.1 — p_T distribution of muons used in this analysis (a) and dE/dx as a function of muon momentum (b) [41].

As seen in Chapter 3.4.3, the longitudinal depth of each readout cell varies as a function of $|\eta|$ (Figures 3.11 and 3.12) to ensure a constant radiation length in the first and second layers for electromagnetic particles. This implies a corresponding variation of the liquid argon depth seen by the muons, and thus their signal response. This requires a binning in η for the analysis, discussed in Section 6.2.

Moreover, from Figure 3.13 we also know that large cross-talk happens in the first sampling. Consequently, the muon energy deposit is built from the energy contained in 3 neighbouring cells in the η direction in S1 to recover the signal induced in neighbouring cells. In S2, due to the accordion fold of the electrodes, a straight going muon most often passes through two neighbouring cells in ϕ , so the deposited muon energy is the sum of two consecutive cells in ϕ . More details are discussed in Section 6.3.

Example muon energy deposit distributions are displayed in Figure 6.2. One may notice the characteristic Landau distribution of a MIP which is largely asymmetric with a long tail at high energy. The distributions are also shifted by a few MeV comparing data to MC.

Since Run 1, two methods are used to estimate the muon energy deposit: a Landau fit and a truncated mean (TM). The first method takes advantage of the simple modelling of the energy loss of the muon to perform an analytical fit on the deposited energy to extract a relevant physical value from the distribution (Section 6.4). The fit model is defined as a Landau distribution, modelling the energy loss by ionisation [167], convolved with a noise distribution, extracting the most probable value (MPV) of the underlying Landau distribution. However given the low SNR of the muon energy deposit, this fit method is sensitive to the noise contribution in the tails, distorting the energy distributions. The second method avoids this issue by computing the mean of the energy distribution, but over a truncated range which cuts out the long tails of the distribution (hence *truncated* mean, in Section 6.5). This however does not directly relate to a physical quantity and is more phenomenological.

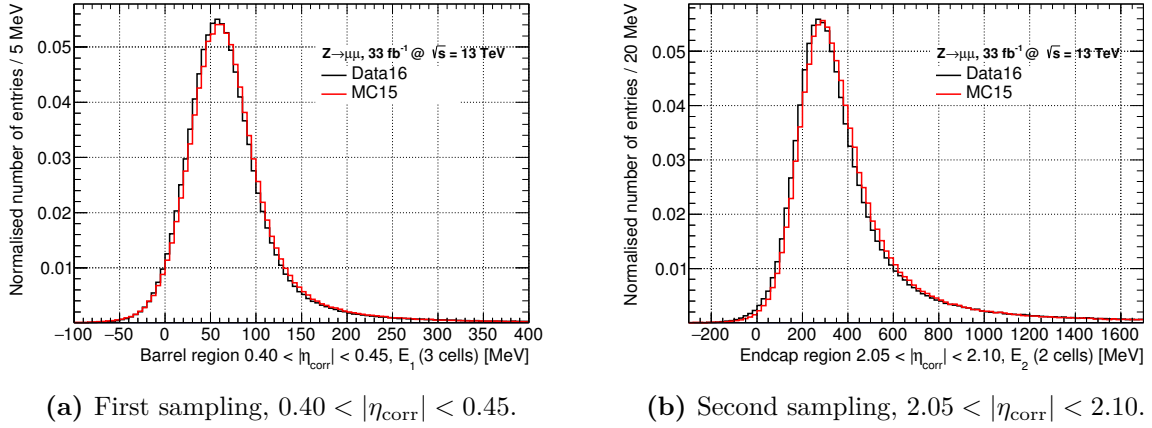


Figure 6.2 – Example energy distribution for data (black) and MC (red) in (a) S1 and (b) S2.

The same analysis is performed on both data and MC, and the final layer intercalibration scale factors are defined as

$$\alpha_{1/2} = \frac{E_1^{\text{data}}/E_2^{\text{data}}}{E_1^{\text{MC}}/E_2^{\text{MC}}}, \quad (6.1)$$

where E_1 (resp. E_2) denotes the measured energy deposited in the first (resp. second) sampling of the calorimeter by the muon, estimated with the MPV of the underlying Landau in the fit, or with the Truncated Mean. The reason for taking the ratio to MC is to account for the fact that the expected response varies as a function of η primarily due to the varying length of LAr traversed by the muons in each sampling. The layer bias corrections $\alpha_{1/2}$ are then applied to the data. At this stage, 3 options are available:

1. correct E_1 as $E_1^{\text{corr}} = E_1/\alpha_{1/2}$;
2. correct E_2 as $E_2^{\text{corr}} = E_2 \times \alpha_{1/2}$;
3. correct both E_1 and E_2 as a mix of the above, conserving the total energy $E_1 + E_2$.

While all possibilities are equivalent in terms of energy after applying the $Z \rightarrow ee$ calibration [152], The second possibility is retained as it was shown to provide the best linearity [140].

The final $\alpha_{1/2}$ scale factors are provided as the combination of the MPV and TM methods as described in Section 6.6. As the procedure is similar to the one used in Run 1, and since the calorimeter has not changed between 2012 and 2016, the results were not expected to differ by a large amount from the Run 1 result recalled in Figure 6.3. This however proved to be an incorrect assumption, and large differences were seen between the Run 2 and the Run 1 results. The differences was eventually found to be due to pileup which is significantly higher in Run 2.

Independently of the method used to extract the MIP signal, this study has shown that pileup has a large impact in the energy estimation (Section 6.4.1). As a consequence, the baseline method used in Run 1 provided different results, especially in the endcap. A method to extrapolate the muon response to zero pileup was developed (Section 6.4.4), and provided good agreement with the Run 1 result. An extrapolation across the pileup has therefore been developed, and is described in Section 6.4.4. The extrapolated results showed to recover the pileup dependency, giving result in good compatibility with the Run 1 expectations (Section 6.6.3).

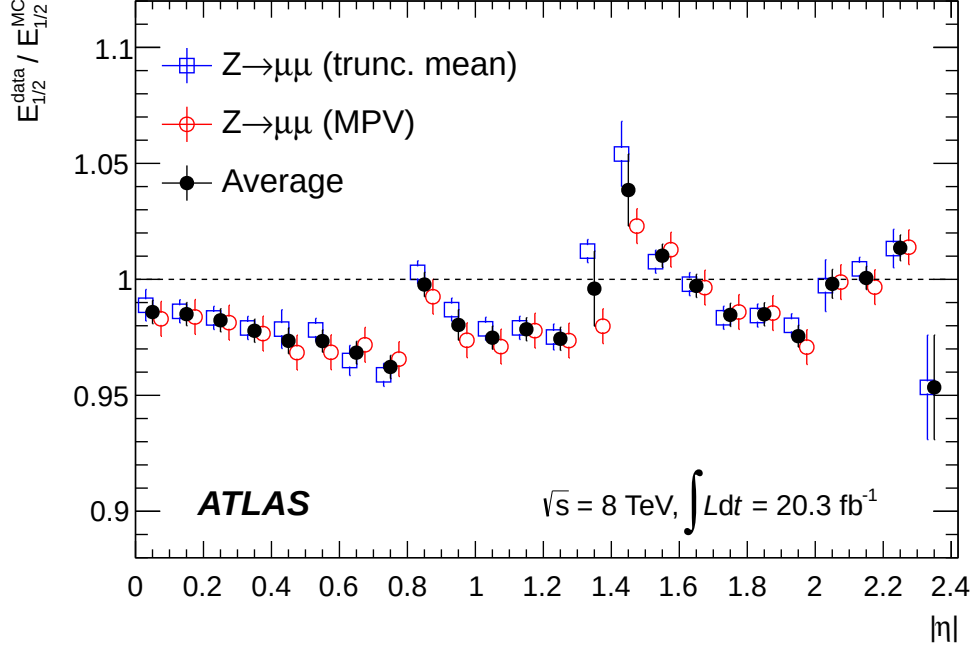


Figure 6.3 – Layer intercalibration scales $\alpha_{1/2}$ as a function of $|\eta|$ derived in Run 1. The blue (Truncated Mean) and red (MPV) points are slightly shifted left and right to improve readability. [140]

6.2 Geometry considerations: η definition and binning

The region of interest for this study is the same as in the previous calibration Chapter, *i.e.* the ATLAS “precision region” $|\eta| < 2.5$. Since the layer intercalibration corrects for electronics and cross-talk, the results will mainly depend on the electrode probed (geometry, electronics, *etc.*). In other words, it will depend on the cell localisation in the calorimeter. Given the granularity of the electrodes layout presented in Chapter 3.4.3, the most adequate binning would be the size of a second sampling cell ($\Delta\eta = 0.025$ and $\Delta\phi = 0.025$). However, since all electrode and modules have been built identically, a ϕ -dependency of the layer scales is not expected. In addition, the statistics available from the 2015+2016 dataset is large but not infinite, so the analysis is integrated over ϕ as the most relevant dependency should come from the electrode depth for each layer, which depends on η .

The $|\eta|$ range of $[0, 2.5]$ is thus roughly split in bins of width $\Delta\eta = 0.10$, as a trade-off between sufficient statistics for the analysis and sensitivity to the detector geometry effects. In the transition region between the barrel and endcap calorimeters ($1.375 < |\eta| < 1.48$) however, the muons can pass through both barrel and endcap cells. Since these do not have a similar energy response (Sections 6.2.2 and 6.2.3), merging them would seriously distort the energy distribution which the simple Landau * noise fit cannot handle. The barrel and endcap regions are therefore studied separately. Furthermore in the region $1.3 < |\eta| < 1.4$ of the barrel, distorted energy distributions are also seen and the region has to be further split, as explained in Section 6.2.2).

6.2.1 Relation between cell η and track η

To recover the position of the cell traversed by the muon, its track must be extrapolated from the ID and MS to the calorimeter. This extrapolation takes into account the deviation due to the magnetic field inside ATLAS, but must also account for the geometry of the calorimeter and its relative misalignment with respect to the ID and MS. To achieve this transformation, the η

quoted in the following plots is defined as

$$\eta_{\text{corrected}}^{\text{Sampling X}} = \eta_{\text{Track extrapolation}}^{\text{Sampling X}} - \left(\eta_{\text{Cell, ATLAS}}^{\text{Sampling X}} - \eta_{\text{Cell, Calo}}^{\text{Sampling X}} \right), \quad (6.2)$$

where $\eta_{\text{Track extrapolation}}^{\text{Sampling X}}$ is the value provided by the ID/MS measurement and $\eta_{\text{Cell, ATLAS}}^{\text{Sampling X}} - \eta_{\text{Cell, Calo}}^{\text{Sampling X}}$ the ATLAS $\rightarrow$ Calorimeter frame transformation, with X meaning strips (S1) or middle (S2).

A typical example of energy profile along η_{corr} for the first and second samplings in the region $0.7 < \eta_{\text{corr}} < 0.9$ in data is shown in Figure 6.4. The first striking points in these plots are the regular structures of a drop in energy appearing with a η period of 0.003 in the first sampling and 0.025 in the second sampling. These correspond to the etching of the electrodes which forms the cells: such areas, however fine, do not have conductive material and therefore can not transmit the ionisation signal. While the signal is recovered in one cell, the leakage in the neighbouring cell is large, which translates into some “energy inefficiency”. This is however a nice illustration of the finely localised muon energy deposit in the calorimeter. Thus, the change of strip width at $|\eta| = 1.8, 2.0$ and 2.4 is also clearly seen in such plots (not included here).

The second remark on this plot is the discontinuity at $\eta_{\text{corr}} = 0.8$, particularly visible in S1. This is due to the physical transition between the two barrel parts $0 < \eta < 0.8$ and $0.8 < \eta < 1.48$: the total liquid argon depth in each layer varies, which affects the muon energy deposit, as well as the lead thickness, and therefore the sampling fraction, which affects the energy reconstruction in the cell. The deposit is otherwise very regular, except for the transition region described in Sections 6.2.2 and 6.2.3.

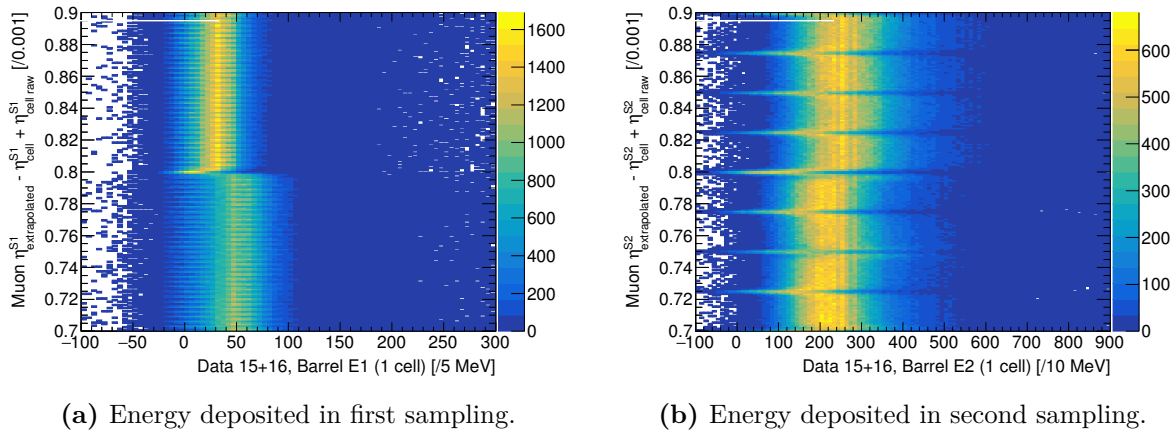


Figure 6.4 – Muon energy deposit as a function of η_{corr} in data for the barrel region $0.7 < \eta_{\text{corr}} < 0.9$, in (a) S1 and (b) S2. The step at $\eta_{\text{corr}} = 0.8$ corresponds to the physical transition of the two barrel parts (see Chapter 3.4.3). The regular dips are due to the etching of the electrodes to form the cells. The colour represents the number of entry in each bin. Each horizontal slice (line of constant $|\eta_{\text{corr}}|$) would give a plot similar to the ones shown in Figure 6.2.

6.2.2 Calorimeter geometry impact in the barrel for $1.30 < |\eta_{\text{corr}}| < 1.48$

The energy deposit as a function of η_{corr} for data in the region $1.3 < \eta_{\text{corr}} < 1.5$ in the barrel is shown in Figure 6.5. The equivalent plot for the MC is similar and all the following discussion can be directly translated to the simulation.

In the second sampling (Figure 6.5b), the energy deposit is stable before 1.325 while a slow decrease is seen starting around $\eta_{\text{corr}} = 1.35$ and continuing up to $\eta_{\text{corr}} = 1.40$. Looking back at Figure 3.11, this can be explained by the depth of the barrel S2 cells which decreases after

$\eta_{\text{corr}} = 1.325$ and is cut off at the “edge” of the modules, corresponding to a decrease in the length of argon traversed and thus deposited energy. To limit the distortion in the energy deposit distribution, this region is split into $[1.30, 1.35]$ and $[1.35, 1.405]$ instead of just $[1.30, 1.40]$. After $\eta_{\text{corr}} = 1.40$ there is only one S2 cell, in which the energy deposit distribution is highly deformed due its size similarly varying as a function of η due to the module edge.

In the first sampling (Figure 6.5a), the energy deposit follows a regular pattern until $\eta_{\text{corr}} = 1.4$. Between 1.4 and ~ 1.48 , three larger areas are distinguishable: looking again at Figure 3.11, these three regions correspond to the 3 much larger “strips” positioned at $[1.40, 1.425, 1.45, 1.475]$ in S1. In this region, the S2 readout is taken to the front instead of the back of the barrel module requiring a thin conducting strip to be removed from the S1 strips. The S1 cells are more widely separated, resulting in the thin white stripes at $\eta = 1.425$ and $\eta = 1.45$ in Figure 6.5a. One single analysis bin is nevertheless used for the region $1.405 < \eta_{\text{corr}} < 1.48$ since the three cells have a similar response.

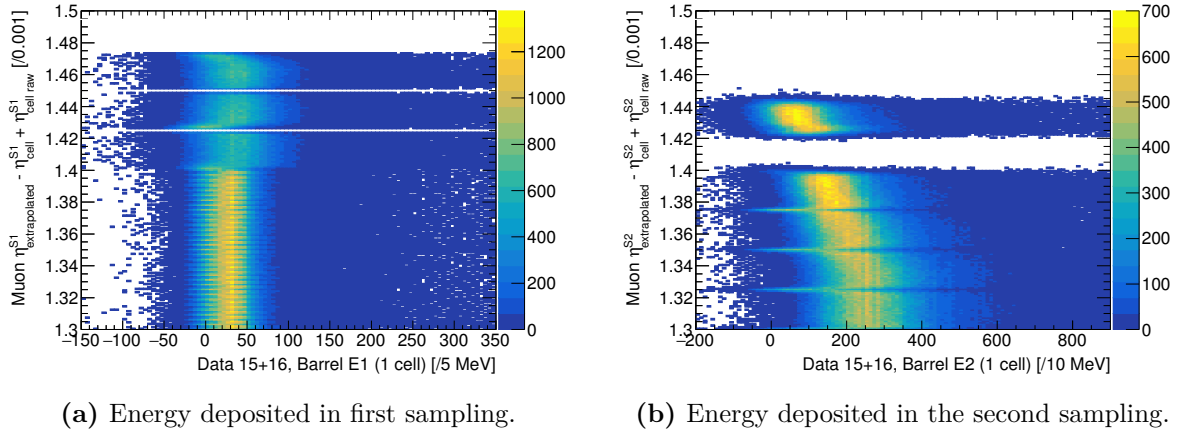


Figure 6.5 – Muon energy deposit as a function of η_{corr} in data for the barrel region $1.3 < \eta_{\text{corr}} < 1.5$, in (a) S1 and (b) S2. The colour represents the number of entry in each bin.

6.2.3 Calorimeter geometry impact in the endcap for $1.375 < |\eta_{\text{corr}}| < 1.50$

Similarly, the energy deposit as a function of η_{corr} for data in the region $1.35 < \eta_{\text{corr}} < 1.55$ in the endcap is shown in Figure 6.6. The corresponding MC plot demonstrates the exact same characteristics and the following directly applies to simulation as well.

As for this region in the barrel, a similar effect is seen in the second sampling (Figure 6.6b): its energy increases slowly between $\eta_{\text{corr}} = 1.38$ and $\eta_{\text{corr}} = 1.50$, and is stable after $\eta_{\text{corr}} = 1.5$. This is again explained by looking at Figure 3.12: the second sampling contains one large and distorted cell in $[1.375, 1.425]$ due to the edge of the module. The three next cells between $\eta_{\text{corr}} = 1.425$ and $\eta_{\text{corr}} = 1.50$ do not have a regular shape and their depth increases as the truncation from the module edge decreases. The muon response thus slowly increases with η_{corr} as the length of liquid argon increases.

The first sampling cells in the region $[1.375, 1.50]$ have the same width as their second sampling counterpart, so large continuous areas are seen in this region (Figure 6.6a). For the same readout reason explained above in Section 6.2.2, the regions are separated by a thin white stripe representing the conducting strips associated to S2 readout in the front of the module instead of the back. After $\eta_{\text{corr}} = 1.5$ the cells resume their standard shape and the energy deposited is smoother.

The larger white bands might correspond to locations where the extrapolated muon track to

a given sampling could not be assigned to a calorimeter cell, due to algorithmic limitations.

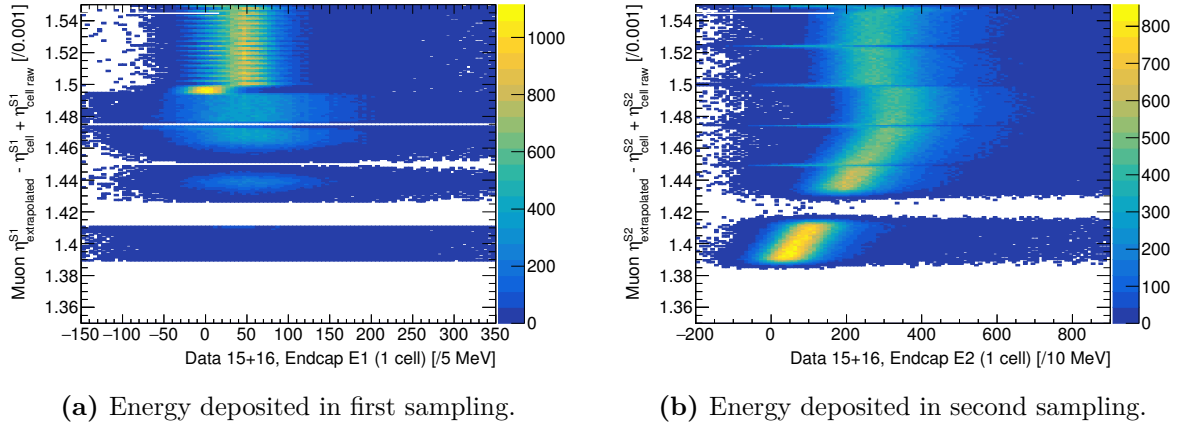


Figure 6.6 – Muon energy deposit as a function of η_{corr} in data for the endcap region $1.35 < \eta_{\text{corr}} < 1.55$, in (a) S1 and (b) S2. The colour represents the number of entry in each bin.

6.2.4 Final binning

The same structures are seen in both the $\eta_{\text{corr}} < 0$ and $\eta_{\text{corr}} > 0$ regions, and after checking that the $\alpha_{1/2}$ results are compatible left and right it was decided to do the analysis as a function of $|\eta_{\text{corr}}|$ instead of η_{corr} to improve the statistical precision. Moreover due to the different response of the calorimeter in the barrel and endcap modules in the crack region, the analysis is conducted separately for the barrel and the endcap. Consequently, the results presented in the rest of this Chapter (and especially in Section 6.4) show multiple points in the crack region, coming from the barrel and endcap analyses. This results in 28 $|\eta|$ regions dispatched as follow in Table 6.1.

Table 6.1 – Binning used for the layer intercalibration analysis with muons.

Barrel	0.00	0.10	0.20	0.30	0.40	0.50	0.60	0.70	0.80	...				
	...	0.90	1.00	1.10	1.20	1.30	1.35	1.405	1.48					
Endcap	1.375	1.425	1.50	1.60	1.70	1.80	1.90	2.00	2.10	2.20	2.30	2.40	2.50	

Due to the peculiar structures in the crack region, more fine tuning of the bin boundaries would probably be necessary if an increase in the precision of the layer intercalibration is needed in this region¹. For the present analysis, the final calibration value in the crack region $1.4 < |\eta| < 1.5$ is taken from the last barrel bin $1.405 < |\eta| < 1.48$. With similar considerations, the $[1.30, 1.35]$ and $[1.35, 1.40]$ results from the analysis are later averaged at the final stage and only one value is quoted.

Remark: the binning used in the first studies, described in Section 6.4.2, is twice finer ($\Delta\eta = 0.05$ in the nominal range). However, after it was understood that additional extrapolation across the pileup is needed, see Section 6.4.4, a coarser nominal η_{corr} binning of 0.1 was adopted. As a consequence, plots displayed in Sections 6.4.1 and 6.4.2 which were done before the extrapolation method was implemented use a $|\eta_{\text{corr}}|$ binning of 0.05.

¹↑ For example, the Higgs boson mass measurement in the $H \rightarrow \gamma\gamma$ channel removes photons in the crack region while the $H \rightarrow 4\ell$ analysis described in Chapters 8 and 9 keeps electrons falling in this region. This nevertheless has only little impact on the measurement.

6.3 Number of cells summed

As already mentioned, an interesting feature of the muon interaction in the calorimeter is a very localised energy deposit. As a consequence, most of the energy is deposited in one or two cells in ϕ , depending on the impact point. The muon energy deposit is therefore considered as being the energy of the cell closest to the track extrapolation to a given layer of the calorimeter, which is referred to as *central cell* in the following.

However, since the first layer is composed of thin strips, the ionisation current can leak into the readout of the first and second neighbouring cells in η around the central one. This cross-talk effect has been shown to be about 10 % for each of the first neighbouring cells and less than 1 % for the second neighbouring cells in data (Figure 3.13), but is poorly described in the simulation [147]. Therefore the comparison of the central cell energy only in S1 is not sufficient. The remedy is to consider the muon energy deposit as being the sum of the energies in the central cell and its two neighbour cells in η (“left and right cells”).

It should be noticed that due to varying η width of the strips at the module edges, the number of cell actually used for the energy computation in the first layer is not always three. The number of cells used is shown in Figure 6.7 for the data and is well reproduced in the MC. In the crack region ($1.4 < |\eta_{\text{corr}}| < 1.5$ for both barrel and endcap) and at the end of the endcap ($|\eta_{\text{corr}}| > 2.4$), the number of cell used to compute the energy is most often limited to the central one only. Around $\eta = 0$, the number of cell summed is also sometimes below three. This is explained by a little gap between the two half barrels which is not instrumented (Figure 3.11), so muons with $|\eta| < 0.006$ corresponding to four strip width are removed from the analysis.

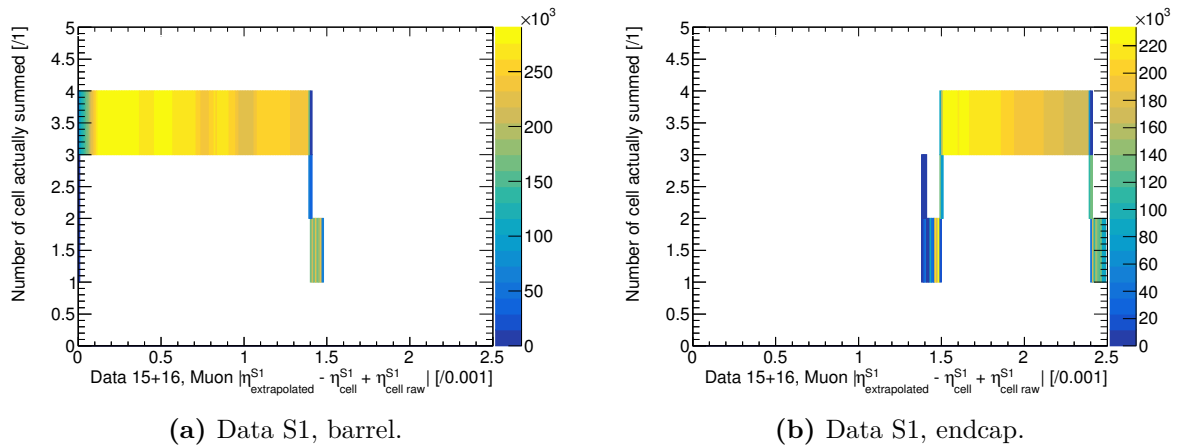


Figure 6.7 – Number of cells actually used to compute the muon deposited energy in the 2015 and 2016 dataset for S1.

In the second layer, due to the accordion geometry of the calorimeter the muon energy is most often deposited over two adjacent cells in the ϕ direction [166]. The full muon energy deposit is therefore built from the sum over the central cell and the most energetic among its two neighbours in ϕ (“above” or “below” cell). An alternative selection of the second S2 cell is to take the second closest to the track extrapolation among the same two neighbours instead of the most energetic. The default choice fell to the first proposal after considering that the signal-to-noise Ratio (SNR) in S2 is about 3-4: the most energetic one should be the one getting the muon energy leakage. However, since the alternative using the track extrapolation is also a reasonable choice, this is used to define a systematic uncertainty as described in Section 6.4.6.2, and is found to have a negligible impact on the result in any case.

Remark: the plots of energy deposit shown in the previous Section (6.2) use one cell for both the first and second samplings. While adding the second cell in S2 does not radically change the plot (only shifts the energy toward higher values), adding up three cells for S1 blurs the structures due to the cell etching and smooths the distribution. In this case indeed, the energy of a muon hitting S1 between two cells will be recovered by the summation of the neighbouring cells.

6.4 Extraction of the muon energy deposit using an analytical fit

6.4.1 Energy distribution fit with Landau * Gaussian convolution

As described in Section 6.1, the observed muon energy deposit distribution can be represented by a Landau convolved with a noise distribution. Assuming a gaussian noise model without energy bias, the energy distribution in each $|\eta_{\text{corr}}|$ bin in S1 and S2 for data and MC can be fitted independently with

$$\text{Model}(E; \text{MPV}, \Gamma, \sigma) = \text{Landau}(E; \text{MPV}, \Gamma) * \text{Gaussian}(E; \mu = 0, \sigma). \quad (6.3)$$

Examples of energy distributions with such fits is given in Figure 6.8 for the data (the MC fits look similar).

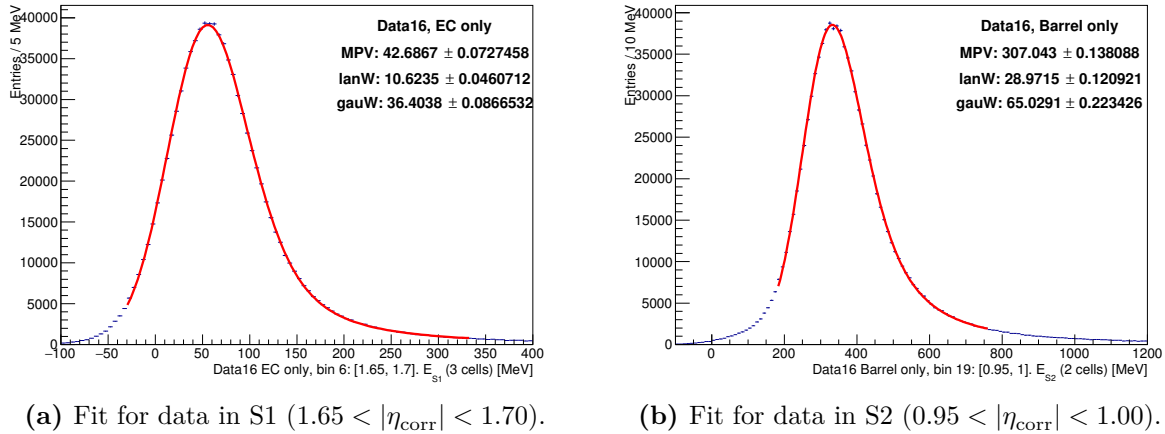


Figure 6.8 – Example distribution of energy deposited in (a) S1 and (b) S2 in various $|\eta_{\text{corr}}|$ region using 2016 data. The distributions are fitted with a Landau * Gaussian convolution and the resulting fit superimposed.

The fitted MPV as a function of $|\eta_{\text{corr}}|$ for both the first and second samplings in data and MC is shown in Figure 6.9. In the first sampling (Figure 6.9a), data and MC follow a similar pattern with an offset of 5 % in the barrel, while in the endcap their behaviour largely differ after $|\eta_{\text{corr}}| = 1.7$. The large step at $|\eta_{\text{corr}}| = 0.8$ is explained by the change in the length of liquid argon traversed due to a shorter S1 electrode. At $|\eta| = 2.3$ the MPV jumps from around 70 to around 50 MeV, which is similarly explained by a decrease of the cells length as shown in Figure 3.12.

In the second sampling (Figure 6.9b), no dramatic effect is seen but the data/MC ratio is less stable, varying between 0.95 to 1. After $|\eta| = 1.35$ the MPV drops by 80 MeV due to the cell shortening (see explanation given in subsection 6.2). At $|\eta| = 2.3$, the depth of the cells increases, giving a 20 MeV increase on the MPV.

All the other large structures (“V” shape, oscillations, *etc.*) in these plots can also be qualitatively explained by the finely varying cell length as a function of $|\eta|$: the S1 and S2 cells

were designed to have an approximately constant number of radiation lengths for electromagnetic showers across the whole $|\eta|$ range (see Chapter 3.4.3). Since the liquid argon and lead thickness is constant along z at a given radius, more forward particle see more material to reach a given radius and therefore a higher number of X_0 . To keep this number fixed, the depth of the cells shrinks as a function of $|\eta|$. However, the muon signal only depends on the length of active material (liquid argon) traversed and not of passive material (lead). The muon signal therefore decreases as the cell depth decreases, *i.e.* as within increasing $|\eta|$.

Finally the double ratio $E_{1/2}^{\text{Data}}/E_{1/2}^{\text{MC}}$ is computed in each $|\eta_{\text{corr}}|$ bin from the values found above, as shown in Figure 6.10. The overlay with the Run 1 result shows a good agreement within 1 % until $|\eta_{\text{corr}}| \approx 1.7$, but the end of the endcap show a totally different behaviour: the double ratio in Run 1 was rising, reaching values above 1, while this analysis exhibits a falling behaviour after $|\eta_{\text{corr}}| = 2.0$.

The following Section (6.4.2) details some of the studies that were performed to further examine the origin of the difference seen in the Run 2 data with respect to Run 1.

6.4.2 Systematic investigation

This Section describes the studies that have been conducted to investigate possible sources for the difference seen between Run 1 and Run 2.

6.4.2.1 Number of cells in S2

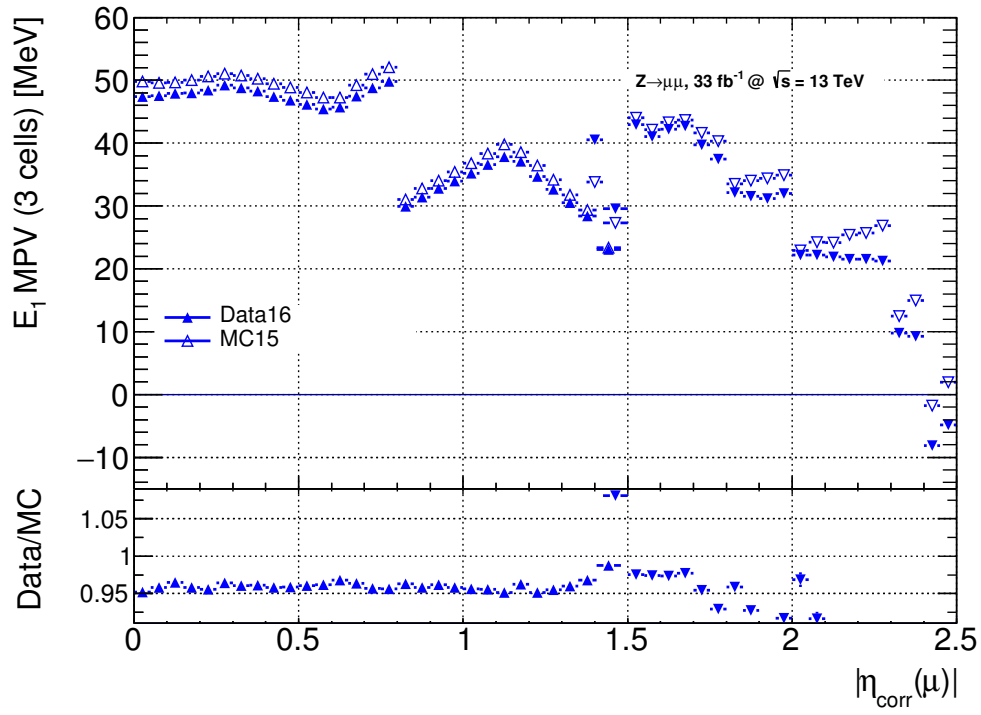
The Run 1 analysis uses only one cell in the second sampling [147], so in order to have a fair comparison, the 2016 analysis was redone changing just the number of cells summed to produce the S2 energy distribution. The results for one and two cells used in the S2 energy computation are shown in Figure 6.11, along with the 2012 result. The impact on the fitted MPV value is obviously a decrease, ranging from 60 to 80 MeV, but the data/MC ratio and consequently the double ratio $E_{1/2}^{\text{Data}}/E_{1/2}^{\text{MC}}$ do not change by much: the barrel changes by a few permil while the endcap undergoes a 1 % to 2 % variation at most. The expected 2012 behaviour is therefore not recovered in the endcap, excluding the number of cell in S2 as responsible for the observed difference.

6.4.2.2 Bunch train effects

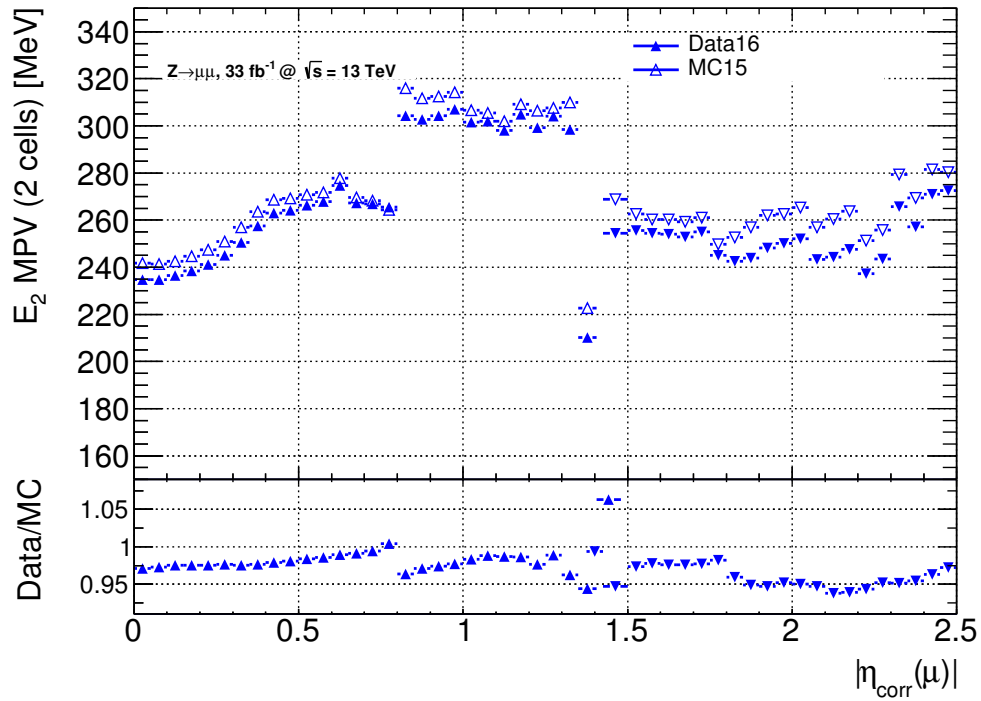
As described in Chapter 4.3, the cell energy needs to be corrected for pileup effects. However, the correction derived using the 2015 and 2016 dataset was not perfect and residuals were found to be up to 40 MeV depending on the $|\eta|$ region and the location of the bunch inside the train, as exemplified in Figure 6.12 (the correction was improved for the full Run 2 dataset processing).

As seen from Figure 6.12, this effect is mainly impacting the measurement of events in the first 20 bunches of a LHC train. It was therefore decided to test the analysis after cutting out the events occurring in this range of each train. The resulting MPV extracted from the Landau * Gaussian fit are shown in Figure 6.13. A difference below 0.5 % is observed in the barrel for S1 in both data and MC. In the endcap, the difference is up to 1 % for the MC but can reach 3 % in the data. This difference in data and MC corresponds to up to 3 % difference in the $E_1^{\text{Data}}/E_1^{\text{MC}}$ ratio. In the second sampling, the variation is negligible over all the $|\eta_{\text{corr}}|$ range, except for $|\eta_{\text{corr}}| > 2.4$.

The impact on the double ratio is lower than 1 % in the barrel (Figure 6.14). In the endcap however, the double ratio is on average 3 % higher, slightly reducing the observed discrepancy with the Run 1 result for the region after $|\eta_{\text{corr}}| > 2.0$, but largely increasing the difference in



(a) First sampling.



(b) Second sampling.

Figure 6.9 – MPV extracted from the Landau * Gaussian fit in each $|\eta_{\text{corr}}|$ region for (a) S1 and (b) S2. Filled points show the data MPV, open points the MC MPV.

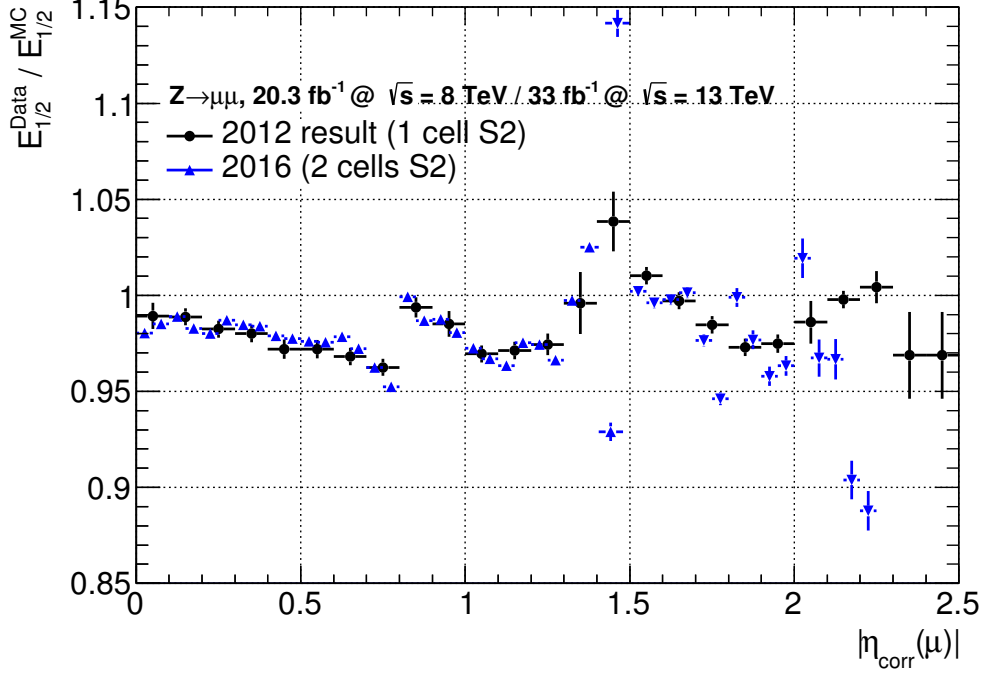


Figure 6.10 – Layer intercalibration scales derived from the Landau * Gaussian analytical fit with the 2016 dataset (blue points), compared to the Run 1 result [140] (black points, average of MPV and TM methods, see Section 6.6).

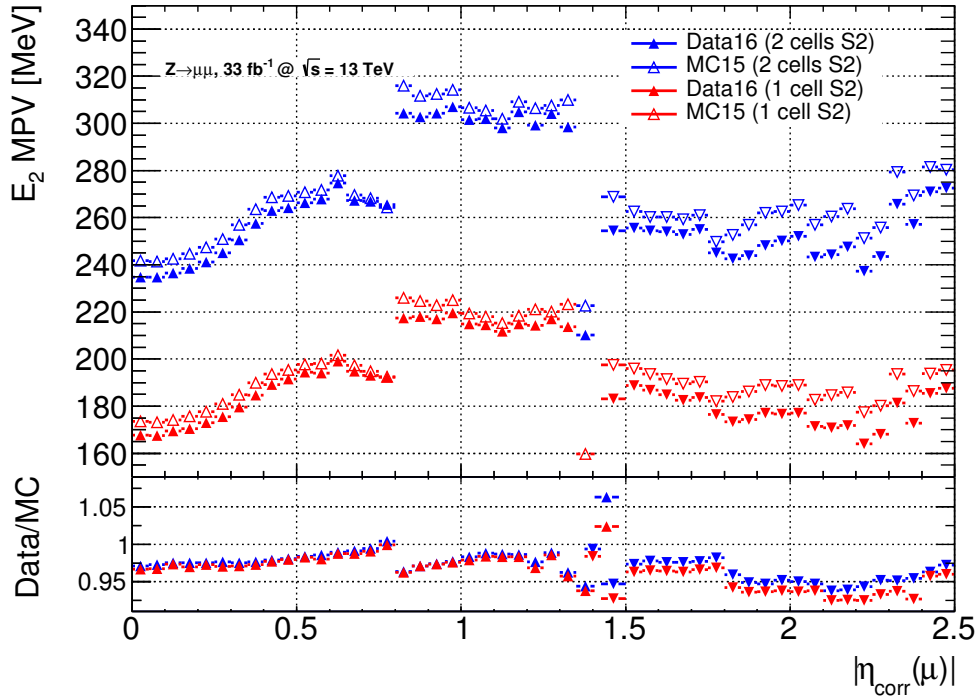
the region $1.5 < |\eta_{\text{corr}}| < 2.0$. This bunch-dependent miscorrection effect is thus ruled out as the cause of the observed Run 1/Run 2 difference.

6.4.2.3 Tighter isolation requirements

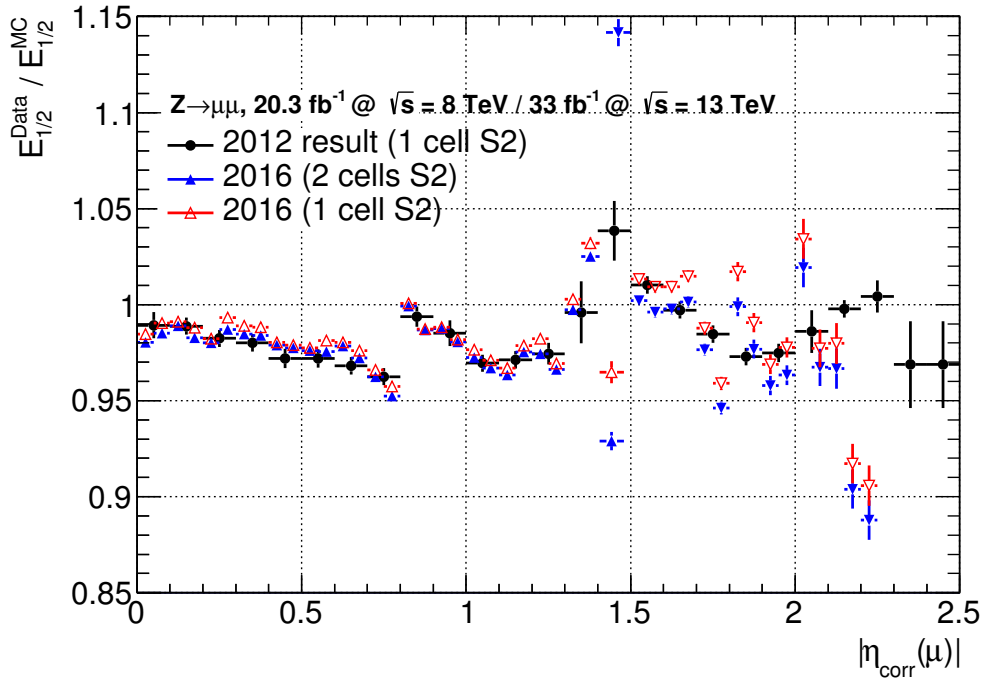
The muon selection criteria for this analysis are not strict, and while the $m_{\mu\mu}$ window selection should ensure that most selected muons actually come from the $Z \rightarrow \mu\mu$ decay, the QCD backgrounds are simply ignored. Since the energy deposit relies on the MIP assumption, which is only fulfilled by muons, any uncontrolled contamination by jets could lead to a bias in the tails of the energy distributions to which the fit is sensitive. In order to ensure that the selected muons are not affected by surrounding QCD jet activity, a tight isolation criterion can be added to the muon selection.

The impact of imposing an isolation requirement on the fitted MPV is shown for data in Figure 6.15. In the strips, the effect in the barrel is barely visible, but is of the order of 5–10% after $|\eta_{\text{corr}}| = 2.0$. The data/MC ratio using isolation in this region changes also but by a slighter amount (1% to 2%) as compared to not using isolation criteria. In the second sampling, the data/MC agreement does not change over the whole range with or without the isolation requirement.

The impact on the double ratio $E_{1/2}^{\text{Data}} / E_{1/2}^{\text{MC}}$ is shown in Figure 6.16. In the barrel, the impact is found to be a few permil, and increases up to 1.5% before $|\eta_{\text{corr}}| = 2.0$. After $|\eta_{\text{corr}}| = 2.0$, the impact is $\sim 2\%$ but increasing the disagreement with the Run 1 analysis. The isolation is therefore excluded from being responsible for the disagreement between Run 1 and Run 2.



(a) MPV of energy deposit in S2 (filled: data, open: MC). The bottom panel shows the data/MC for 1 and 2 cells.



(b) $E_{1/2}^{\text{Data}} / E_{1/2}^{\text{MC}}$ with 1 or 2 cells in S2. Black points recall the 2012 analysis.

Figure 6.11 – Comparison of the MPV in S2 (a) and of the layer scales $\alpha_{1/2}$ (b) as a function of $|\eta_{\text{corr}}|$ when using one or two cells to reconstruct the muon deposit energy in S2. The blue points correspond to the standard analysis using two cells in S2, and red points to the analysis using only one cell. The S1 energy distribution is not impacted and remains the same as in Figure 6.9a.

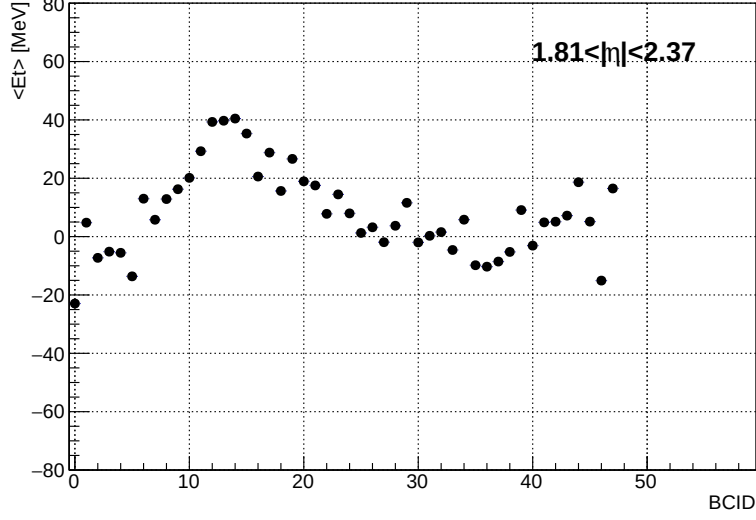


Figure 6.12 — Average transverse energy in clusters built from zero-bias events of a single data run in the $|\eta|$ region $[1.81, 2.37]$ as a function of the bunch crossing number, before correction of the residuals. The BCID identifies the position of a bunch in a train (all trains are folded together here). The residuals amount to up to 40 MeV. [168]

6.4.2.4 Ambient noise subtraction

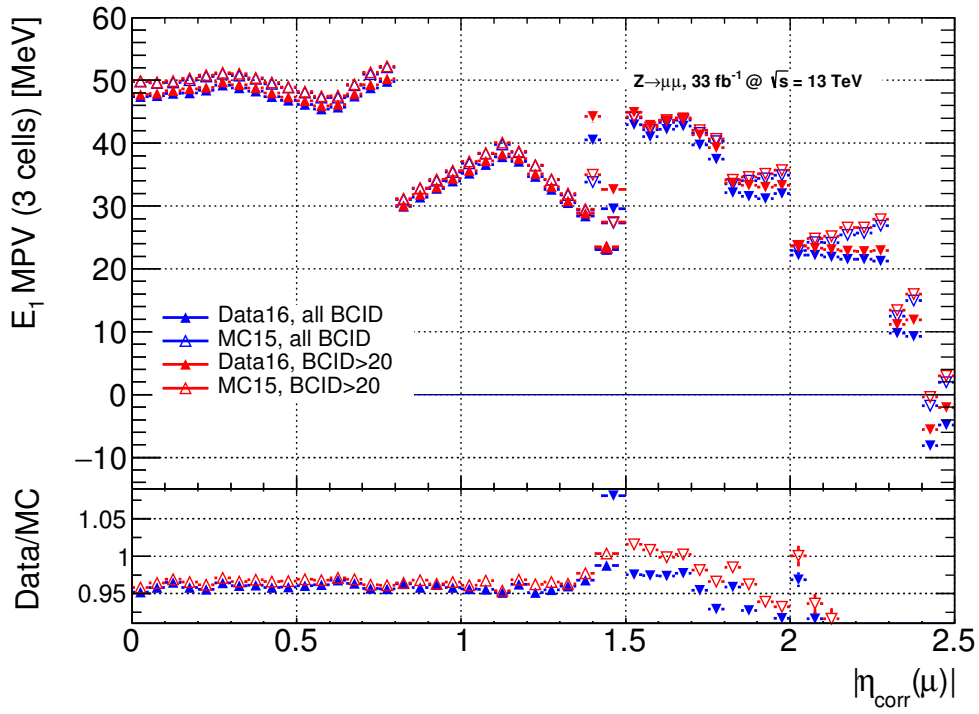
Another primary concern when dealing with MIP energy deposits is that noise fluctuations might have a significant impact on the measurement. This relates in this study to the gaussian used to model the noise behaviour possibly being too imprecise. In order to reduce such a possibility, the test developed here is to estimate the ambient noise around the cells hit by the muon and subtract it from the reconstructed muon energy deposit. To do so, a 3×7 (5×5) cell window is opened in S1 (S2) centred around the expected cell crossed by the muon track². One can therefore estimate the ambient noise by averaging the energy in each cell forming the external border of the window, providing an “average noise per cell” (Figure 6.17). This value is thereafter subtracted from the previously reconstructed muon energy deposit, having beforehand multiplied by the actual number of cell used to reconstruct the muon deposited energy.

The noise distributions in large $|\eta|$ regions are shown in Figure 6.18. One noticeable feature is that the noise distribution peaks at negative values, which is especially visible in the endcap region. Indeed in case of large pileup, the average contribution will be negative due to the LAr pulse shape.

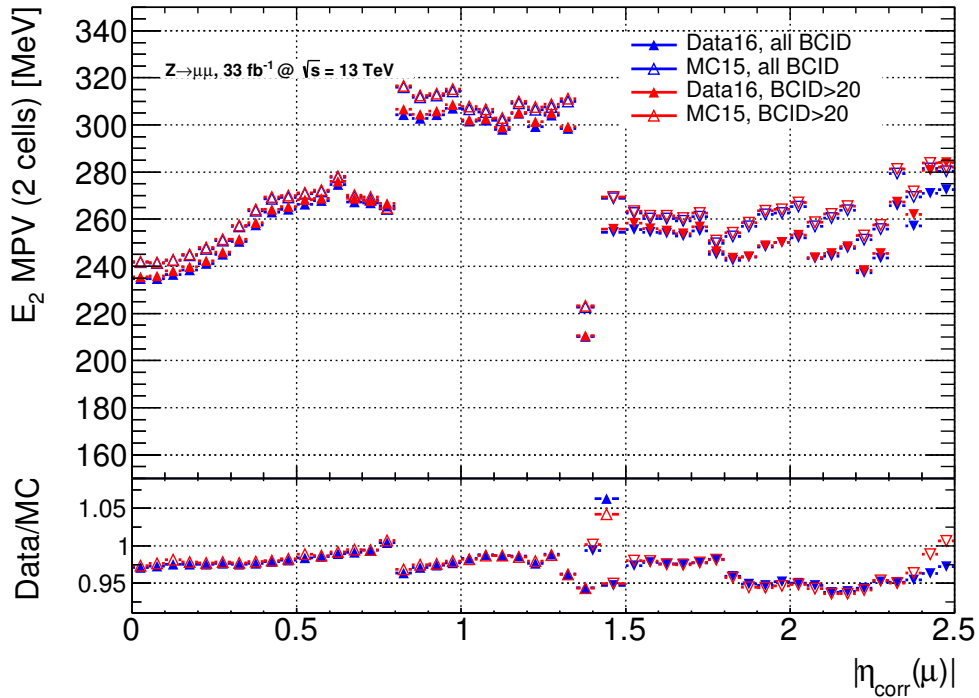
After subtracting the noise, the fitted MPV as a function of $|\eta_{\text{corr}}|$ is shown in Figure 6.19. In the first sampling, only the MC is affected by 1 % to 2 % in the barrel while the data is impacted by just a few permil. In the endcap, both data and MC MPV increase by up to 20 %, but the MC is still more impacted. In the second sampling, the impact in the barrel is negligible while the endcap shows a moderate difference of 2 %.

The total impact on the $E_{1/2}^{\text{Data}}/E_{1/2}^{\text{MC}}$ double ratio is a shift of -1% in the barrel, and a more dramatic impact amplifying the falling behaviour in the endcap, as presented in Figure 6.20. This effect is understood as the noise in the cells surrounding the muon track not being correlated

²↑ These window sizes are valid only in regular regions, *i.e.* excluding the crack and $|\eta| > 1.8$ for the first sampling: for the latter, since the strip width increases, the window decreases to 5×3 and even 3×3 cells. The ambient cell noise procedure is still achievable, only fewer cells are used in the average. In edge cases for which a cell used for the muon deposit computation also falls in the border cells, this cell is used in the muon energy reconstruction but removed from the noise estimate so as to not introduce a bias.



(a) MPV in the first sampling.



(b) MPV in the second sampling.

Figure 6.13 – MPV extracted from the Landau * Gaussian fit in each $|\eta_{\text{corr}}|$ region for S1 (a) and S2 (b) with the standard analysis (no BCID cut, blue) and after removing events in the first 20 bunches of the train (red).

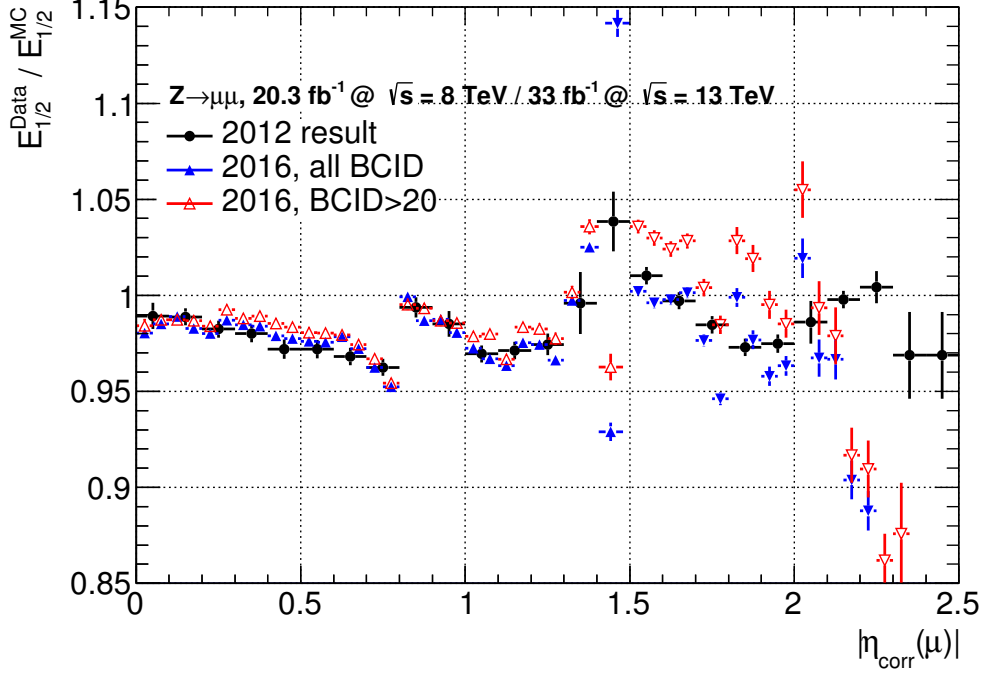


Figure 6.14 – Double ratio $E_{1/2}^{\text{Data}} / E_{1/2}^{\text{MC}}$ extracted with all events (blue) or with event located after the 20th bunch of each train (red). The black points recall the Run 1 analysis.

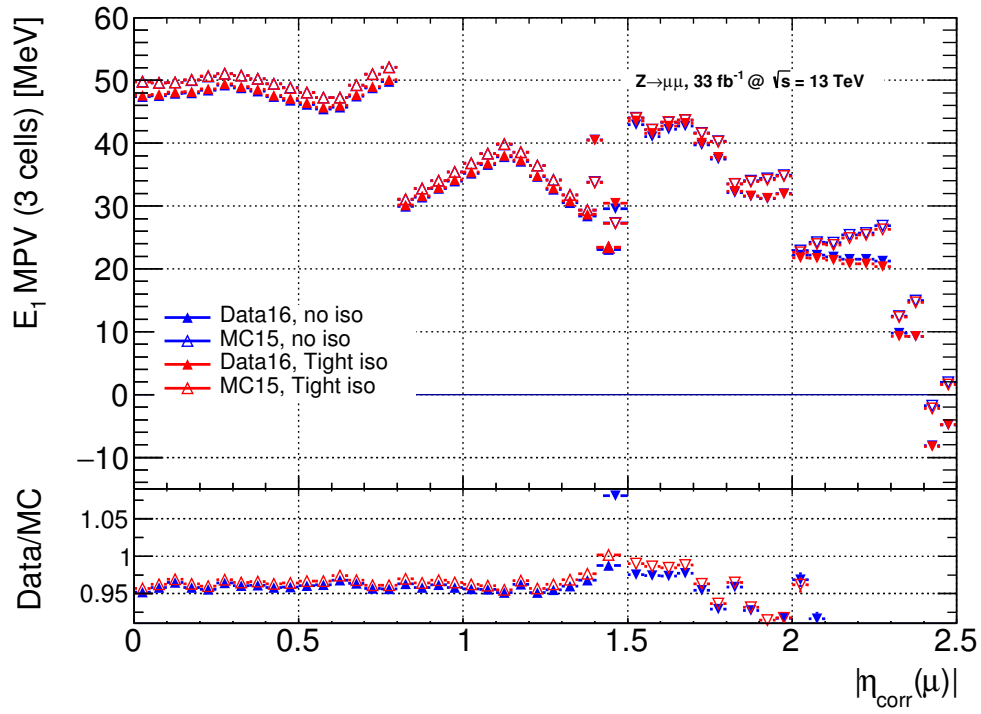
enough with the noise of the cell crossed by the muon.

6.4.2.5 First hint: pileup dependency

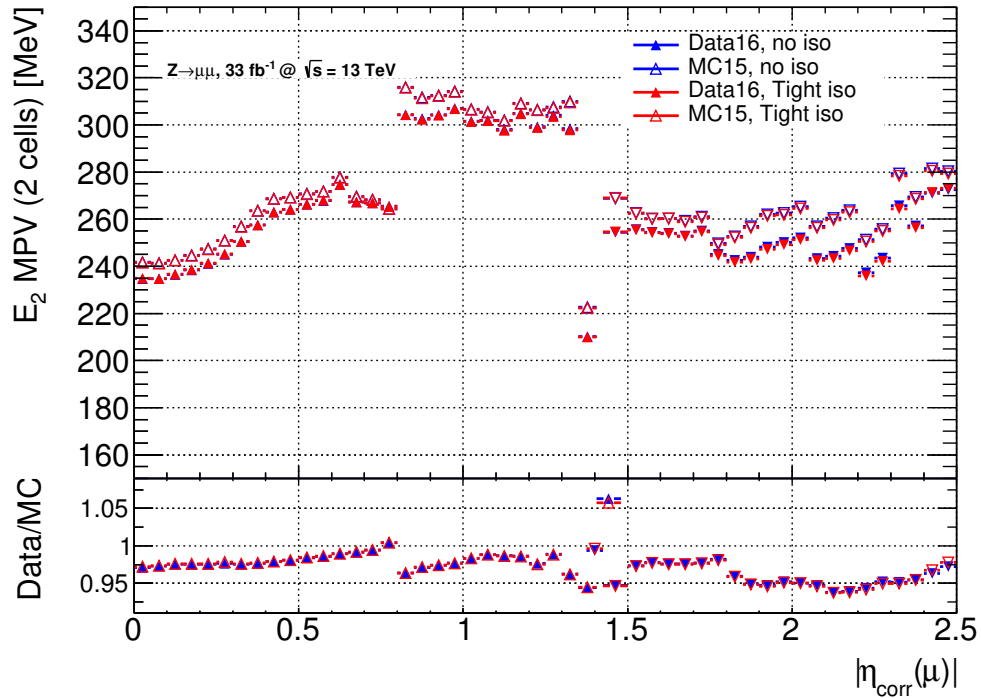
The main factor studied so far has focused on reducing the noise affecting the muon energy reconstruction. This allowed to measure an impact of the order of 1 MeV but all ideas showed a limited impact on the fitted MPV, of the order of 1 MeV. However, most of the noise in this study comes from pileup jets that bias the muon energy response in the calorimeter toward lower values. Moreover as seen in Chapter 2.10, the running conditions of the LHC and particularly the average number of interactions per bunch crossing (the *pileup*) $\langle \mu \rangle$ changed in a sizeable manner between Run 1 and Run 2. Pileup therefore constitutes a main suspect for the origin of the Run 1/Run 2 discrepancy. In order to probe this dependency, the MPV is fit as a function of $\langle \mu \rangle$: the 2016 dataset is thus split in 3 bins of equivalent statistics covering the pileup range ($[0, 19]$, $[19, 25]$, $[25, 40]$) and the analysis performed in each of these three bins.

The impact on both S1 and S2 fitted MPV in data is shown in Figure 6.21. The medium pileup bin matches well the pileup inclusive analysis, showing that the latter really follows an average behaviour regarding the pileup. It must be noticed that the higher the pileup value, the lower the fitted MPV. The most spectacular effect occurs in the first sampling where differences of more than 10 % among the low and high pileup bins are seen in the barrel, and even much larger in the endcap. In the second sampling the effect is more moderate but still noticeable with differences of the order of 2 % among the low and high pileup bins.

This dependency indicates that the scales derived inclusively in pileup also include a correction for pileup on the energy response, which they are not intended to do. The direct consequence is a resulting bias in the layer scales. A method to overcome the pileup impact is described in Section 6.4.4 and constitutes the main development of this analysis realised during this thesis.



(a) MPV in the first sampling.



(b) MPV in the second sampling.

Figure 6.15 – MPV extracted from the Landau * Gaussian fit in each $|\eta_{\text{corr}}|$ region for S1 (a) and S2 (b) with the standard analysis (no isolation requirement, blue) or after requiring a **Tight** isolation cut (red). The bottom panels show the data/MC value with or without isolation requirement.

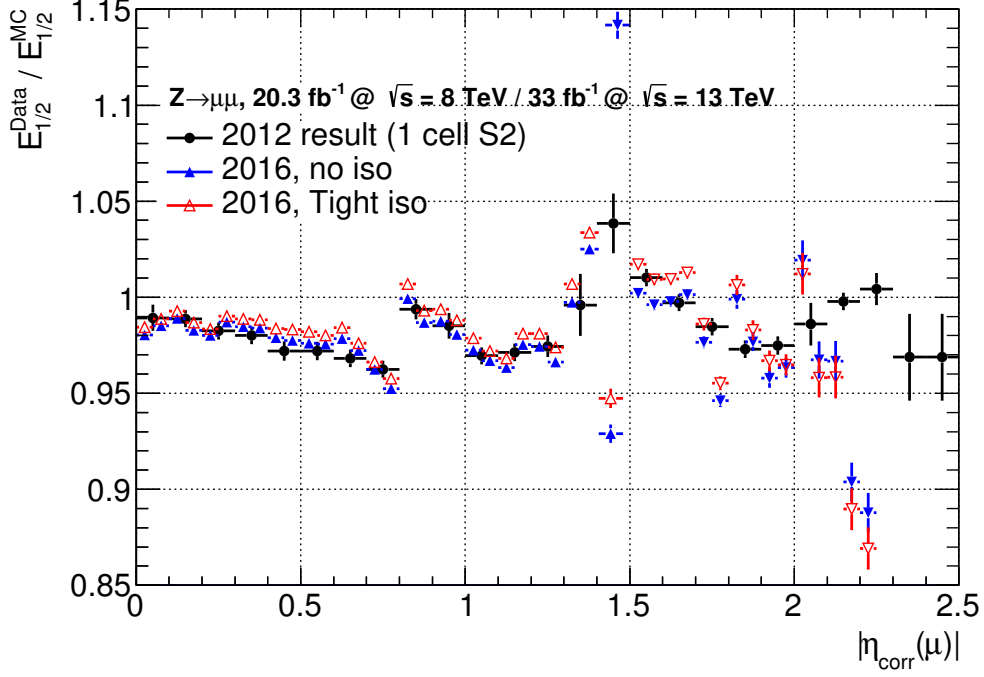
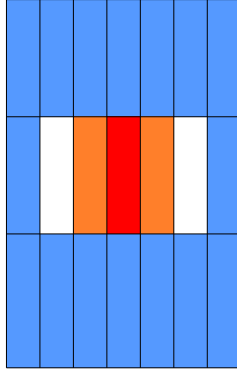
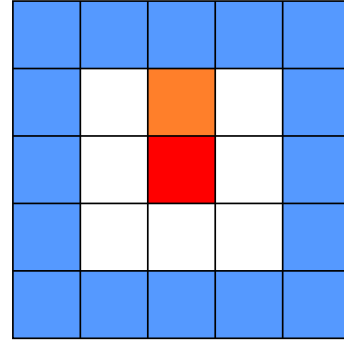


Figure 6.16 – Double ratio $E_{1/2}^{\text{Data}} / E_{1/2}^{\text{MC}}$ extracted with (red) and without (blue) isolation cuts. The black points recall the Run 1 analysis.



(a) 7×3 cell window in S1.



(b) 5×5 cell window in S2.

Figure 6.17 – Schematic of the cell window opened by the reconstruction algorithm around the expected muon track (a) in S1 and (b) in S2. The x axis stands for the η direction, the y axis for the ϕ direction. The red cell is the one closest to the muon track extrapolation, and the orange ones the additional cells summed for the muon energy deposit computation (see Section 6.3). The cells coloured in blue represent the cells used for the ambient noise computation. The schemes do not respect the relative proportions of the cells.

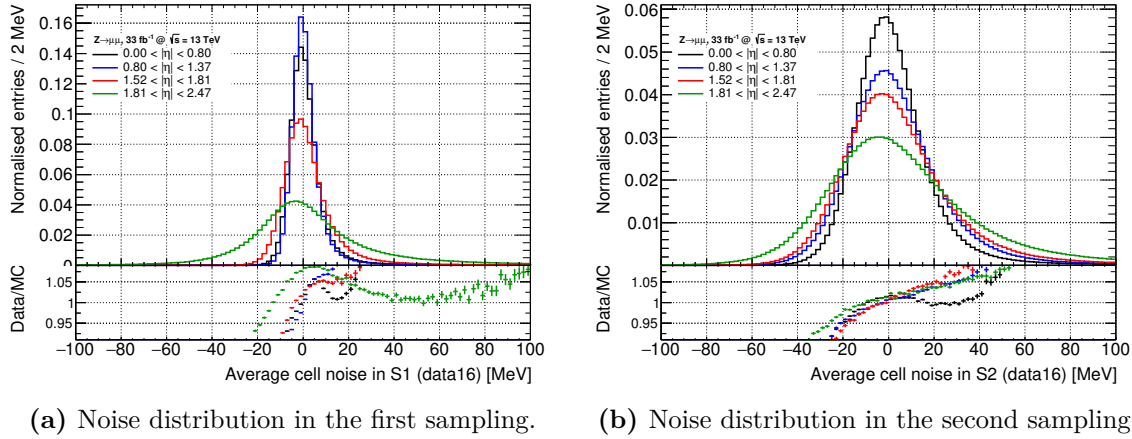


Figure 6.18 – Noise distribution in (a) S1 and (a) S2 computed with the average cell energy taken from the window border.

Remark: all previously mentioned studies were again performed in these three $\langle\mu\rangle$ bins and did not show a significant additional reduction on the Run 1 disagreement in the endcap, pointing to the pileup as the primary cause of the disagreement.

6.4.3 Noise modelling with templates

Since all the previous studies showed a dependency on the noise, it was investigated whether the gaussian model for the noise may not be sufficient. The fit model was therefore changed from a *Landau * Gaussian convolution* to a *Landau * Template convolution* to try to get rid of any noise modelling dependency. Such templates are built separately in data and MC in each $|\eta|$ bin of the analysis from randomly created clusters in events not containing any hard-scatter events (“pileup-only” or “zero-bias” events). The templates were also extracted separately in each pileup bin since a significant fraction of the noise is expected to come from the pileup.

Examples of templates in the data are given in Figure 6.22. The best Gaussian fit is superimposed and can be seen to not fully match the template. This is contrary to the previous assumptions: the template mean is not 0, and the noise distribution is asymmetric with a more extended tail toward high energy. This strongly supports the use of such template for the fitting procedure instead of a gaussian function.

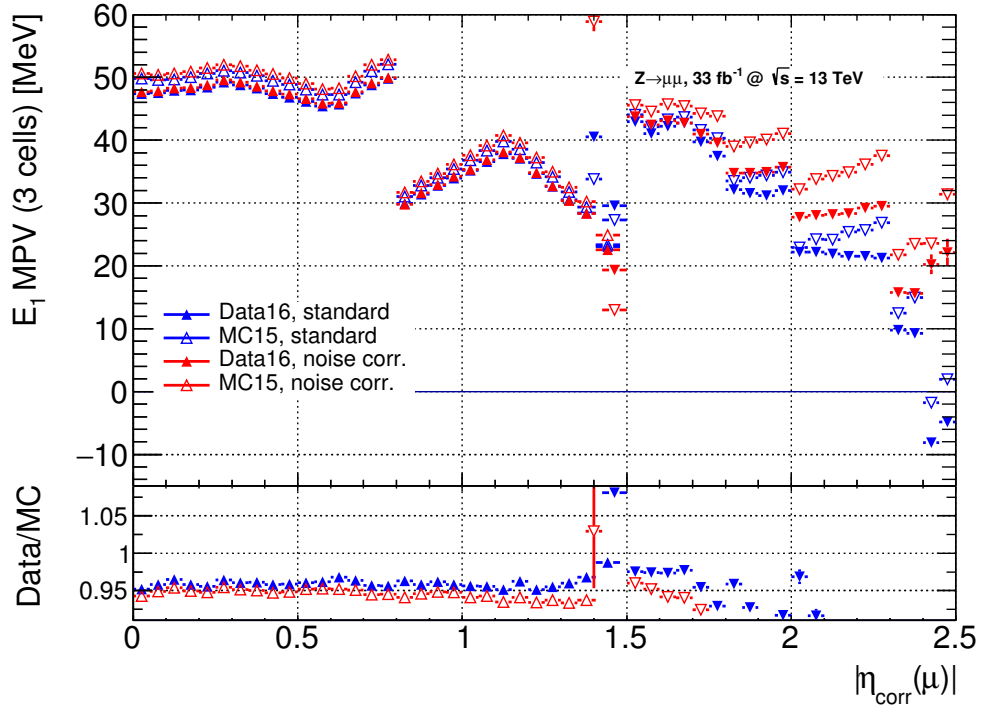
6.4.4 Extrapolating the layer intercalibration to $\langle\mu\rangle = 0$

Since the layer intercalibration is foreseen to correct for intrinsic electronics and cross-talk effects, it must be insensitive to pileup. The best method would thus be to measure the layer scales from $\langle\mu\rangle \approx 0$ collisions. Unfortunately there was no low pileup data recorded in 2015 nor 2016, and an alternative method has been developed.

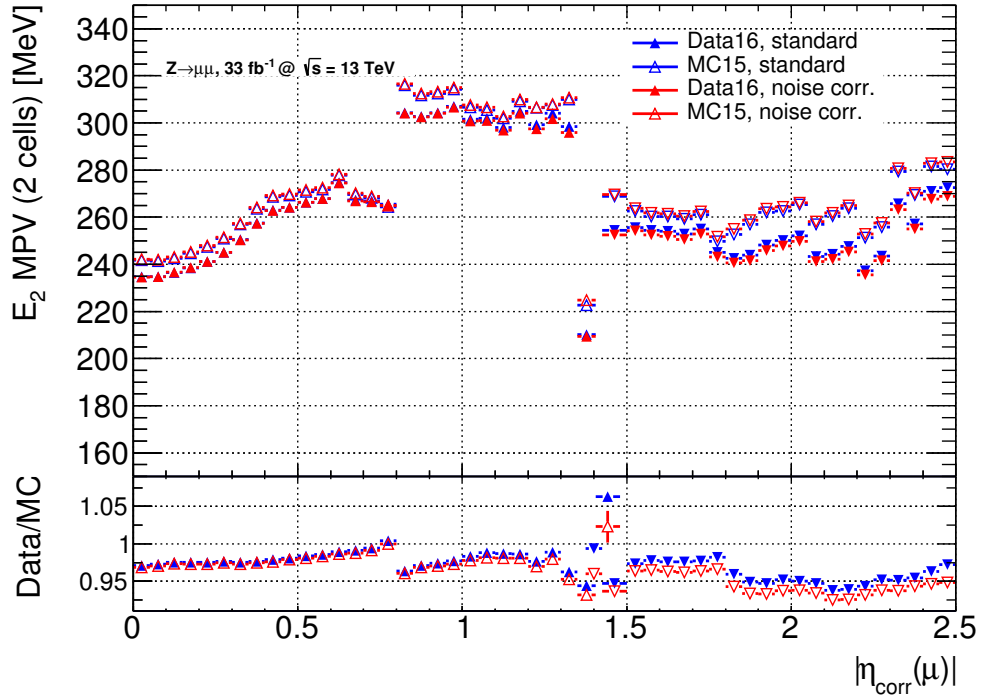
Given the high statistics available in the “high- $\langle\mu\rangle$ ” data taking, the considered solution relies on extrapolating across the pileup range from high $\langle\mu\rangle$ values to $\langle\mu\rangle = 0$. The high-pileup data and MC samples are therefore split in several pileup bins and the MPV fit procedure is repeated in each of these pileup bins, for both data and MC.

The chosen pileup binning uses 13 regions, which is the result of a balance between a sufficient number of points in the linear fit and statistical uncertainty. The bins are defined as

$$[0, 10, 12, 14, 16, 18, 20, 22, 24, 26, 28, 30, 34, 44], \quad (6.4)$$



(a) MPV in the first sampling.



(b) MPV in the second sampling.

Figure 6.19 – MPV extracted from the Landau * Gaussian fit in each $|\eta_{\text{corr}}|$ region for S1 (a) and S2 (b) with (red) and without (blue) applying the ambient noise subtraction procedure. The bottom panels show the data/MC value for each case.

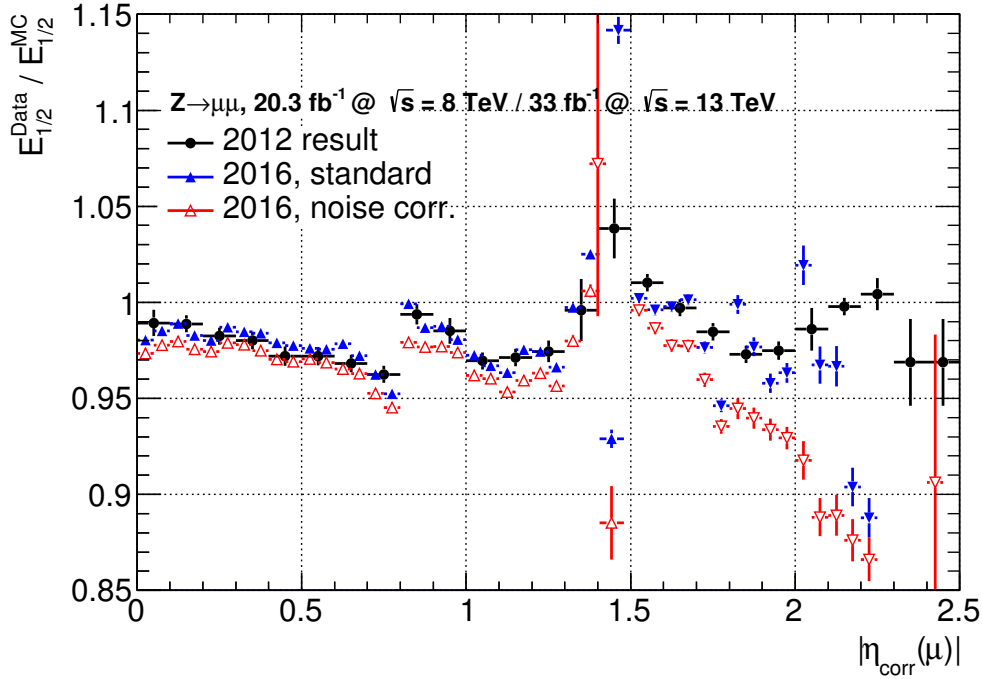


Figure 6.20 – Double ratio $E_{1/2}^{\text{Data}}/E_{1/2}^{\text{MC}}$ extracted with (red) and without (blue) ambient noise subtraction. The black points recall the Run 1 analysis.

covering the whole pileup range in both the data and MC. However as shown in Figure 6.23, the simulated pileup profile available in the MC (“MC15c”) lacks in statistics in the region above $\langle\mu\rangle > 30$. On the contrary, only a small amount of data was recorded with $\langle\mu\rangle < 12$. Consequently, the first and last two bins are removed from the extrapolation linear fit.

The muon energy deposit model is also updated to include the noise template modelling described in Section 6.4.3. The templates are built for each $|\eta|$ and $\langle\mu\rangle$ bin in order to have an accurate modelling across all the $|\eta|$ and $\langle\mu\rangle$ range.

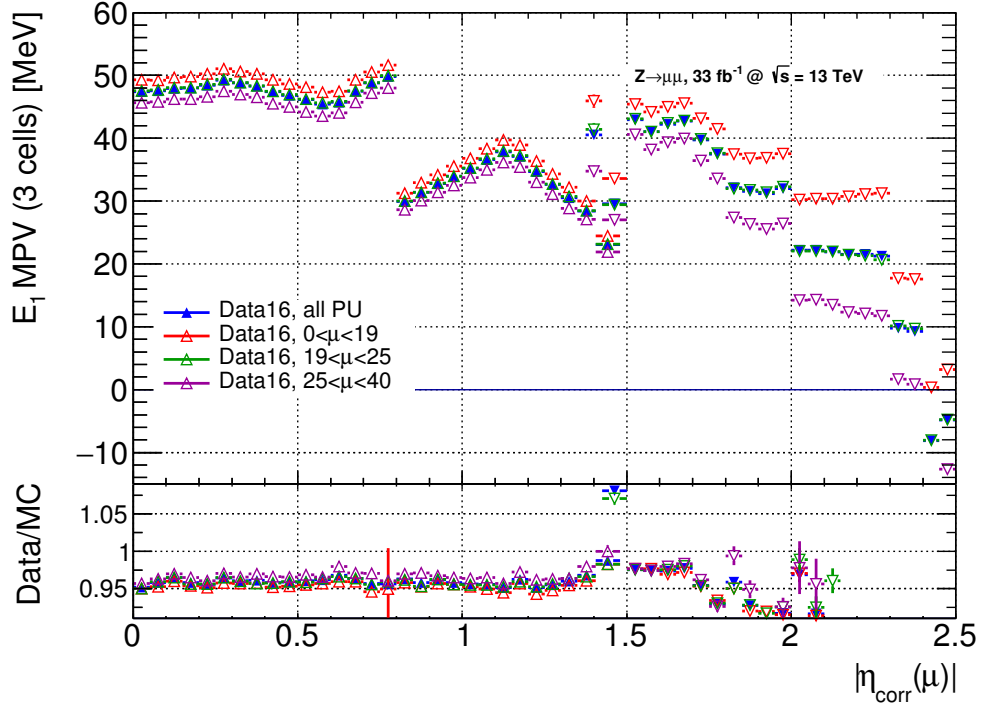
Examples of fitted MPV as a function of the pileup for data and MC are shown in Figure 6.24. As can be seen from these plots, the behaviour of the MPV as a function of the pileup is roughly linear, so a simple linear fit over the pileup range is performed in each $|\eta_{\text{corr}}|$ bin of the analysis and for S1 and S2, in data and MC. The intercept of these linear fits at $\langle\mu\rangle = 0$ gives the “extrapolated MPV” value in each $|\eta_{\text{corr}}|$ bin. The layer intercalibration scales are then computed similarly to Eq. 6.1, but with extrapolated E_1 and E_2 MPV values:

$$\alpha_{1/2} = \frac{E_1^{\text{data, extr}}/E_2^{\text{data, extr}}}{E_1^{\text{MC, extr}}/E_2^{\text{MC, extr}}}, \quad (6.5)$$

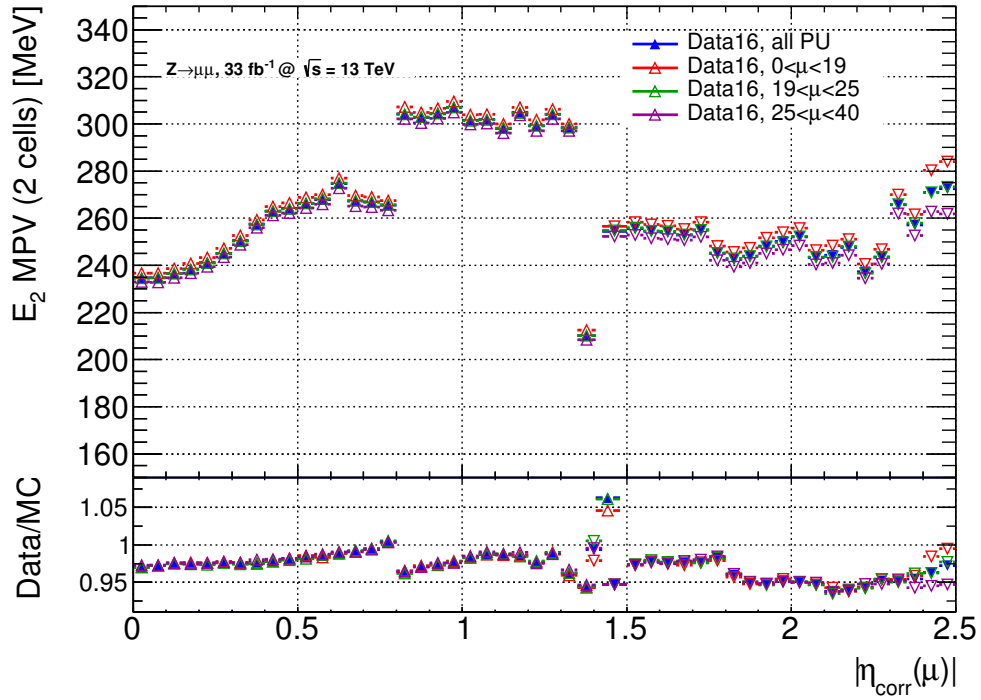
providing the final scale results discussed in Section 6.4.5.

To validate the extrapolation method, the extrapolated results from the MC are compared to the energy fit results from a MC sample produced without any pileup. The agreement is found to be at the percent level in the first sampling, and a few permil in the second sampling. Residual differences are assigned a systematic uncertainty, later described in Section 6.4.6.1.

Note: the need for the extrapolation validation in data on top of the MC validation has been a significant supporting point for the preparation of “low- $\langle\mu\rangle$ ” recording runs in ATLAS in 2017 and 2018. The analysis of such datasets is currently ongoing within the calibration group.

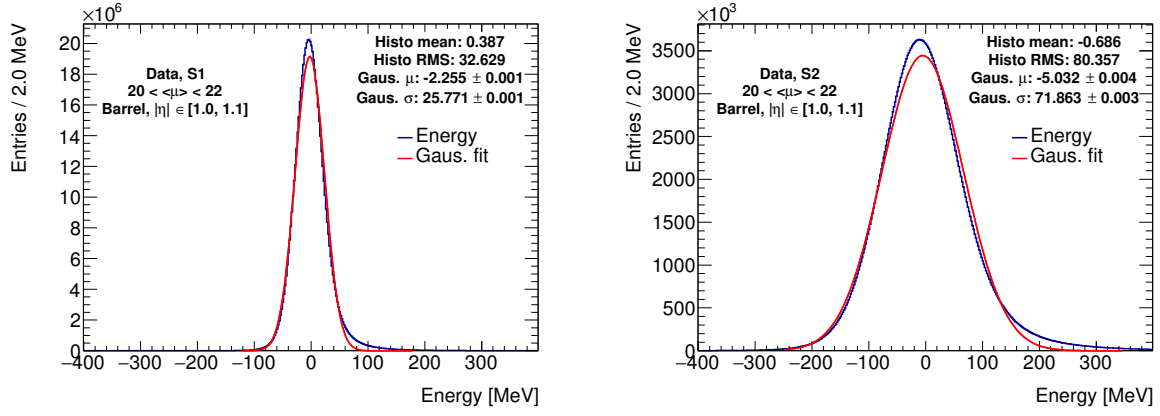


(a) MPV in the first sampling.



(b) MPV in the second sampling.

Figure 6.21 – MPV extracted from the Landau * Gaussian fit in each $|\eta_{\text{corr}}|$ region for (a) S1 and (b) S2 in several pileup bins: “low” ($0 < \langle \mu \rangle < 19$) in red, “medium” ($19 < \langle \mu \rangle < 25$) in green, and “high” ($25 < \langle \mu \rangle < 40$) in pink. The blue point is inclusive over pileup. The bottom panels show the data/MC value for each case.



(a) Noise template and fit in the first sampling. (b) Noise template and fit in the second sampling.

Figure 6.22 – Noise distribution in the data and best gaussian fit (in red) for the pileup bin $20 < \langle \mu \rangle < 22$ and the $|\eta_{\text{corr}}|$ bin [1.00, 1.10], in (a) S1 and (b) S2. These distribution are used as noise templates in the convolution for the fit of the total energy distribution.

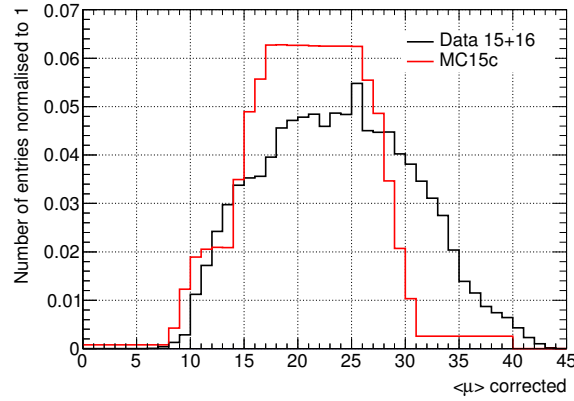
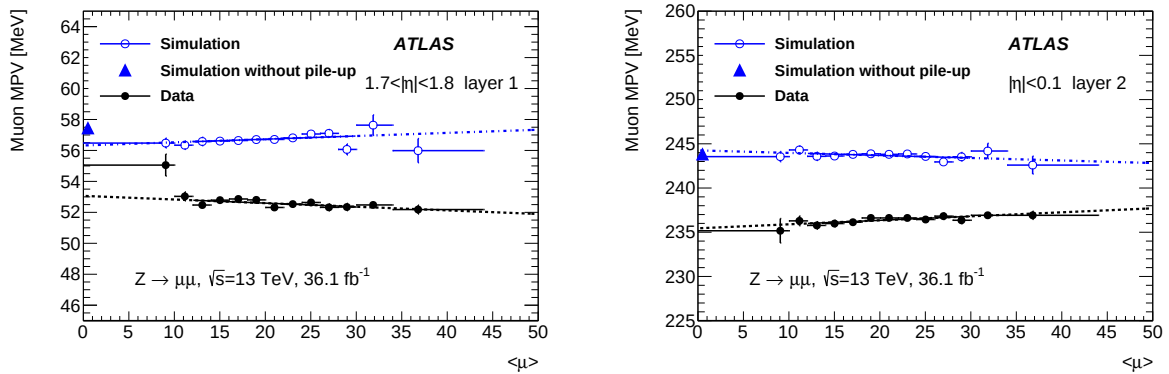


Figure 6.23 – Pileup profiles in the 2015+2016 dataset (black line) compared to the generated profile in the MC15c (red line) used in this study. The distributions are built after applying the selection cuts described in Section 6.1.



(a) MPV and fit as a function of $\langle \mu \rangle$ for $1.7 < |\eta_{\text{corr}}| < 1.8$ in S1. (b) MPV and fit as a function of $\langle \mu \rangle$ for $0 < |\eta_{\text{corr}}| < 0.1$ in S2.

Figure 6.24 – Example of MPV distribution as a function of $\langle \mu \rangle$ for (a) the endcap in S1 and (b) the barrel in S2. The data points are shown in black and the MC points in blue. The solid lines show the fit in the fitting range, and the dashed line its extrapolation. The blue triangle shows the fitted MPV in the MC sample simulated without pileup ($\langle \mu \rangle = 0$). [141]

6.4.5 Results

The extrapolated MPV value as a function of $|\eta_{\text{corr}}|$ using the full 2015+2016 dataset are presented in Figure 6.25. In the first sampling (Figure 6.25a), the qualitative behaviour of the MPV as a function of $|\eta_{\text{corr}}|$ in the endcap is quite different when using or not the extrapolation: instead of the steps seen at $|\eta_{\text{corr}}| = 1.8, 2.0$ and 2.4 , the extrapolated MPV shows a steady increase from $|\eta_{\text{corr}}| = 1.5$ to 2.3 . These pre-extrapolation steps can be explained by the increase of the strip width at these values³, which implies that the cells become more sensitive to the pileup which consequently decreases their energy. The intrinsic muon deposit is not sensitive to the width of the cells (only to their length), so these steps should not be seen in absence of pileup. The absence of such structure after extrapolation demonstrate that the pileup effects are indeed largely reduced using this technique. The strip length decreases by about 25 % for $|\eta| \in [2.3, 2.4]$ and a bit more for $|\eta| > 2.4$, which explains the remaining difference seen in the last two S1 bins after the pileup extrapolation. In the barrel the results with or without the extrapolation show a consistent behaviour, shifted by ~ 5 MeV up for the data and by ~ 7 MeV up for the MC, changing their ratio by 2 % to 3 %.

Contrary to S1, the pileup sensitivity in S2 is much lower so no large behaviour difference is found with or without the extrapolation. A global shift of 5 to 15 MeV is seen across the $|\eta_{\text{corr}}|$ range, but this shift is similar in data and MC as attested by the similar data/MC ratio for both methods.

The layer scales $\alpha_{1/2}$ derived with the muon MPV method are displayed in Figure 6.26. The shape of the extrapolated results (red) is much closer to the one seen from the Run 1 result (black), and a good agreement between the two is found in the endcap. In the barrel, the agreement is however better without extrapolation, but the results remain largely compatible within statistical errors. This result indicates the origin of the discrepancy with the Run 1 result as being due to pileup, and validates the extrapolation method developed here. In order to further assess the validity, a discussion of uncertainties attached to this measurement is detailed in the next Section.

Remark: all of the studies checking the impact of the noise described in Section 6.4.2 were redone after applying the extrapolation, without indicating any significant effect. This also supports that the original Run 1/Run 2 difference is directly linked to the pileup and not some other effect.

6.4.6 Systematic uncertainties of the MPV method

After the derivation of the layer intercalibration scales (Section 6.4.5), the uncertainties still have to be assessed. The following sources of uncertainties covering the pileup extrapolation procedure, effects of the calorimeter geometry and of the detector alignment have been studied:

- closure of the extrapolation method for the first and second samplings, evaluated in MC,
- variation of the choice of second cell in S2,
- leakage energy in S1 if the muon impact point is near the edge of the cell,
- leakage energy in S2 if the muon impact point is biased.

³↑ The strip width in η is 0.025/8 for $1.5 < |\eta| < 1.8$, 0.025/6 for $1.8 < |\eta| < 2.0$, 0.025/4 for $2.0 < |\eta| < 2.4$ and 0.025/1 for $2.4 < |\eta| < 2.5$. The full details can be found in Figure 3.3.

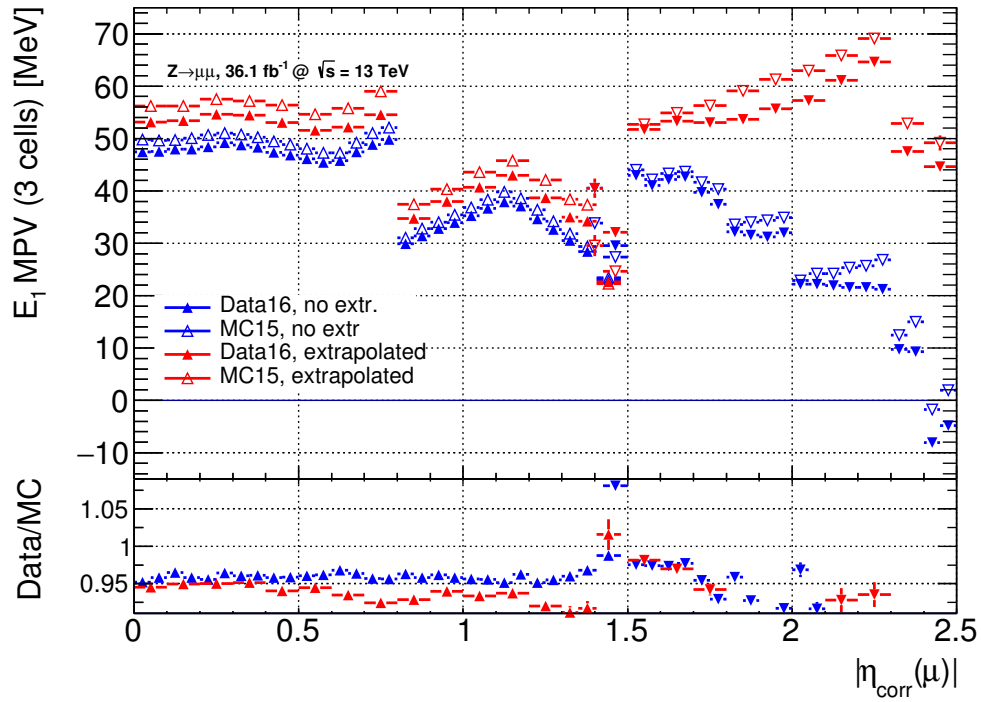
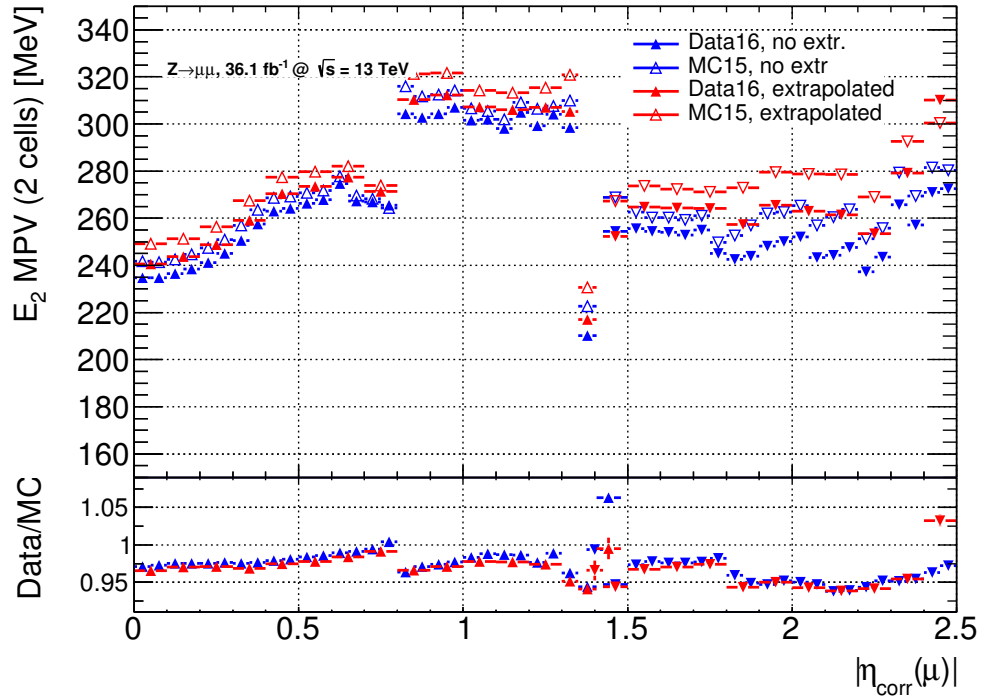

 (a) MPV extrapolated to $\langle\mu\rangle = 0$ in the first sampling.

 (b) MPV extrapolated to $\langle\mu\rangle = 0$ in the second sampling.

Figure 6.25 – Comparison of the energy deposit MPV with (red) and without (blue) applying the extrapolation method in (a) S1 and (b) S2. The filled points correspond to the data, the open points to the MC. in the data (black points) and in the MC (red points). The bottom panels show the data/MC ratio with and without extrapolation.

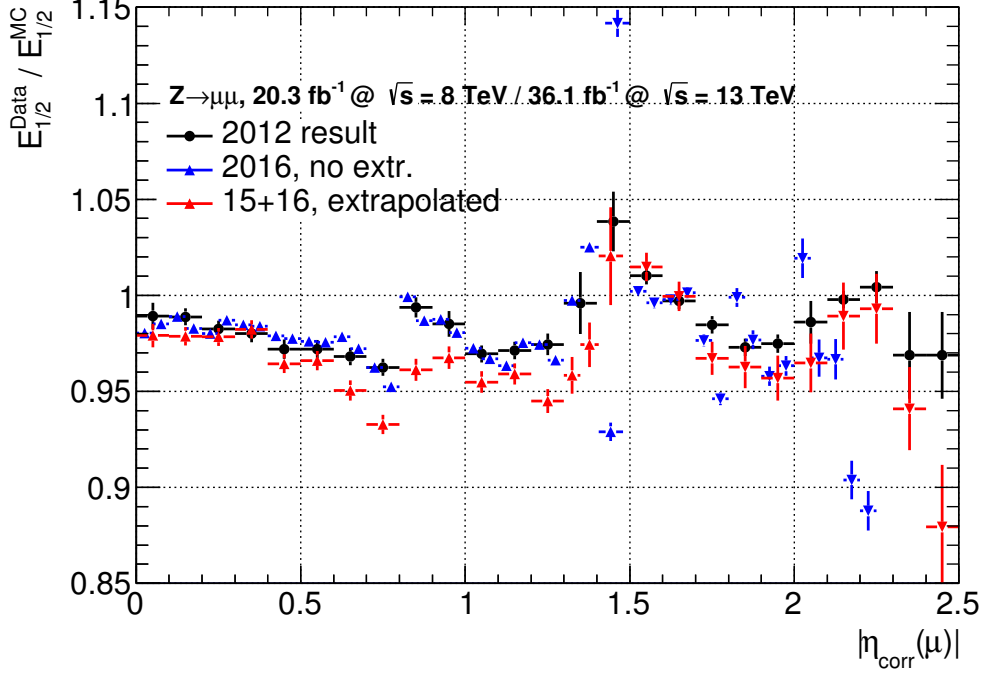


Figure 6.26 – Double ratio $E_{1/2}^{\text{Data}}/E_{1/2}^{\text{MC}}$ using the muon MPV method as a function of $|\eta_{\text{corr}}|$ using the 2015+2016 dataset. The results without extrapolation are shown in blue and with extrapolation in red. Errors are statistical only. The Run 1 result (combining MPV and TM methods) is overlaid in black.

6.4.6.1 Validation of the extrapolated results

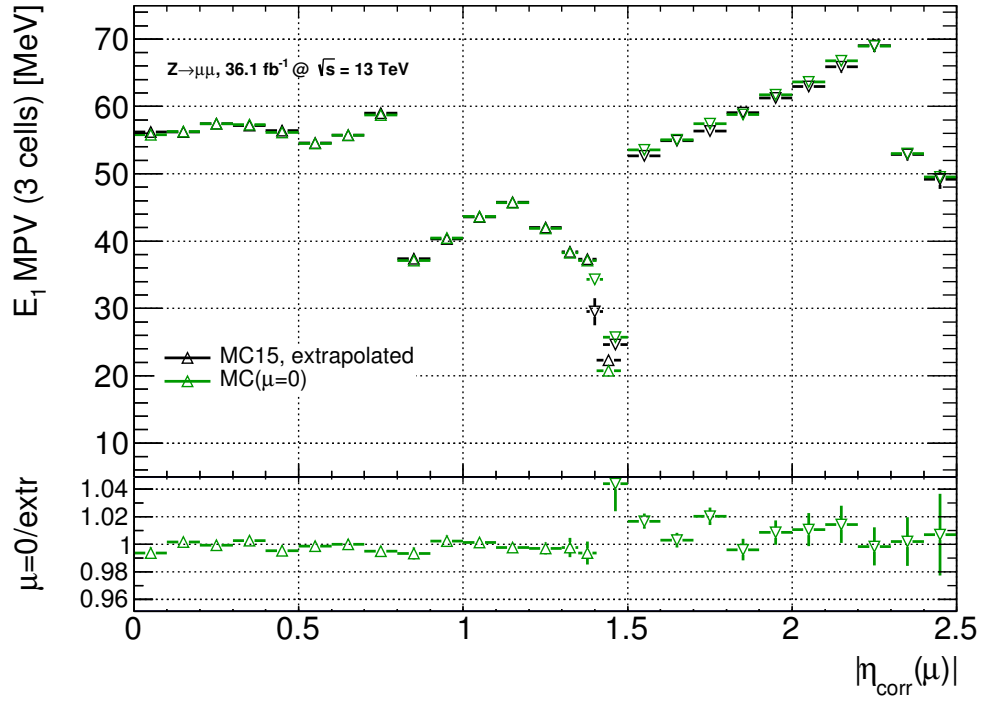
The most straightforward way to validate the extrapolation procedure is to check the extrapolated MPV results against the MPV extracted from a sample with no pileup. Since the 2015 and 2016 datasets did not include very low pileup runs (contrary to 2017 and 2018), the extrapolation cannot be checked in the data. However, the MC can be simulated with any value of pileup. A new simulation with $\langle\mu\rangle = 0$ (labelled “MC($\mu = 0$)”) was produced, and the MPV derived from this sample.

The closure test is performed comparing the extrapolated MPV in the MC to the MPV directly fitted from the MC sample with $\mu = 0$. The results of this closure test are shown in Figure 6.27. For the first sampling, the agreement is within 1 % in the barrel until $|\eta| = 1.4$ and after $|\eta| = 2.2$ (with larger statistical uncertainties in the latter case). Between 1.5 and 2.2, the agreement still stays within 2 %. In the second sampling, the agreement is within a few permil over the full range. Remaining differences in the closure test of S1 and S2 are taken as a systematic uncertainty on the final result.

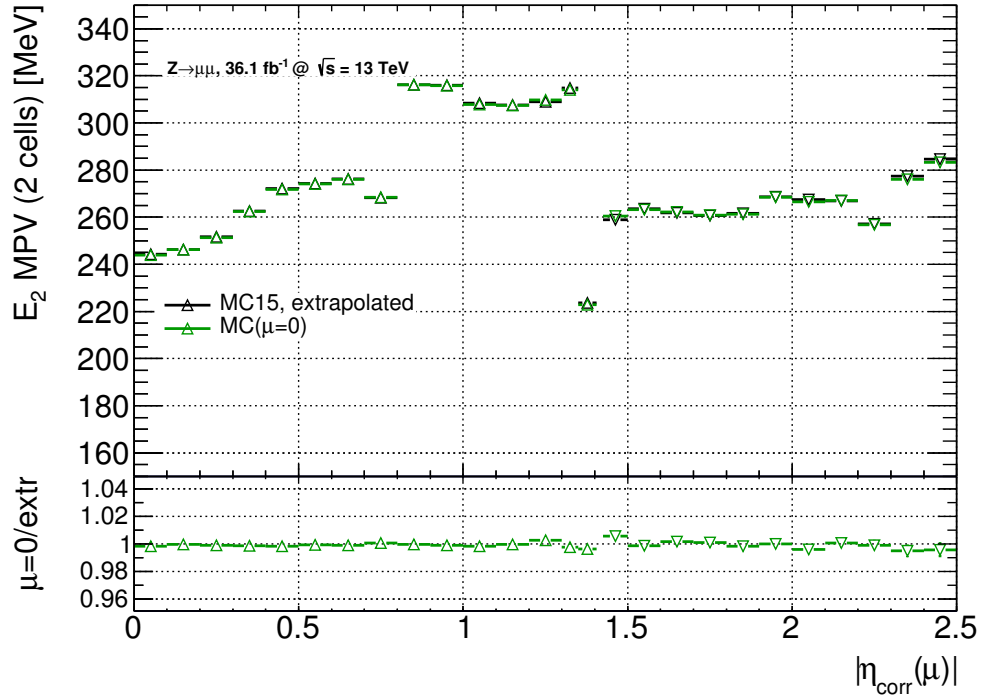
6.4.6.2 Choice of the second cell in S2

As already discussed in Section 6.3, due to the accordion form of the absorbers and electrodes a muon will most often deposit its energy in S2 in two neighbouring cells in the ϕ direction (Figure 6.28). The baseline choice for the second cell is then to add the most energetic neighbour cell in the ϕ direction to the central cell (the one closest to the track extrapolation)

Another reasonable choice of second cell is the second closest cell in ϕ to the extrapolated track. This second cell is actually more likely to be hit by the muon than the highest energetic. Figure 6.29 shows the distribution of energy difference induced by the different choice of second



(a) S1 comparison.



(b) S2 comparison.

Figure 6.27 – Result of the extrapolation method closure test in the MC, for (a) S1 and (b) S2 as a function of $|\eta_{\text{corr}}|$. The agreement is within 0.5 % in the barrel of S1, and 2 % in the endcap. For S2, the agreement is within a few permil over the whole range.

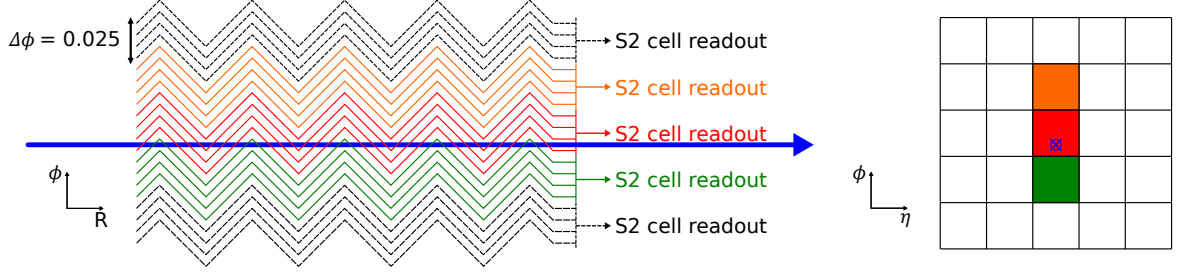


Figure 6.28 – Scheme of the folded electrodes in S2 in the (R, ϕ) plane representing 5 cells in the ϕ direction (left), and its associated opened 5×5 window in the (η, ϕ) plane around the muon track used in the analysis (right). One S2 cell is defined by the readout of 4 consecutive electrodes. The red electrodes and the red cell denote the one crossed by the muon (shown as a blue arrow). The orange denotes the most energetic cell among the two neighbouring cell in the ϕ direction, while the green denotes the closest to the track among the same two cells. In this example, the muon hits both red and green electrodes, so its deposit will be shared among the red and green cell. However the cell with highest energy deposit (orange) among the two neighbour is not the closest (green) to the track, so the original choice of cell (red + orange) will overestimate the muon energy. The second choice of cell (red + green) recovers the real energy share in the second cell. While the folding angle decreases as R increase in the actual calorimeter, the folding angle shown here is 90° everywhere for illustration purpose but the reasoning is the same.

S2 cell. The difference largely peaks at 0 (more than half the events), meaning that for half of the muons the largest energy deposit cell is the same as the second closest.

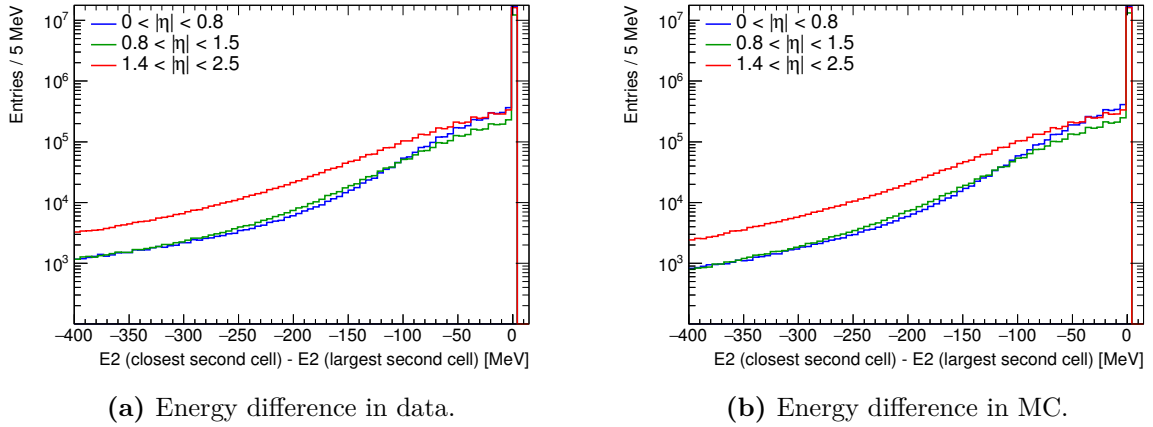
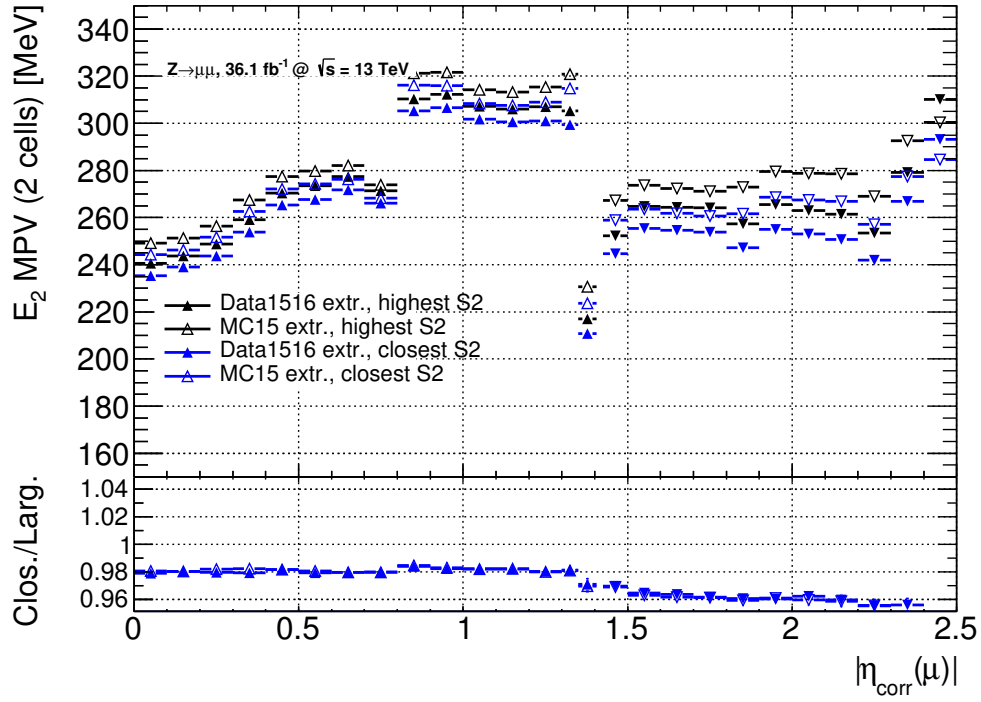


Figure 6.29 – Difference of energy between the highest energetic cell and the second closest cell in S2, in the data (a) and in the MC (b).

Results showing the difference in fitted energy for the two choices of cell are shown in Figure 6.30. In the barrel the reconstructed muon energy deposit is 2% lower when using the second closest cell, and above 4% in the endcap. The impact is similar in data and MC, resulting in an almost negligible difference on the double ratio of the order of a few permil. This systematic uncertainty is taken as the difference of the double ratios selecting one or the other cell.

6.4.6.3 Leakage energy in S1

As in S2, some energy is shared with a neighbouring cell in ϕ when the muon passes near the cell edge due to the accordion geometry (Figure 6.31). However, contrary to S2, S1 cells are four times wider in ϕ ($\Delta\phi = 0.1$) and the signal addition from a neighbouring ϕ cell would be overwhelmed by the increase in noise. To estimate the effect of this leakage, only muons falling completely within a single ϕ cell are selected. A cut was applied keeping only muons with



(a) MPV comparison in S2.

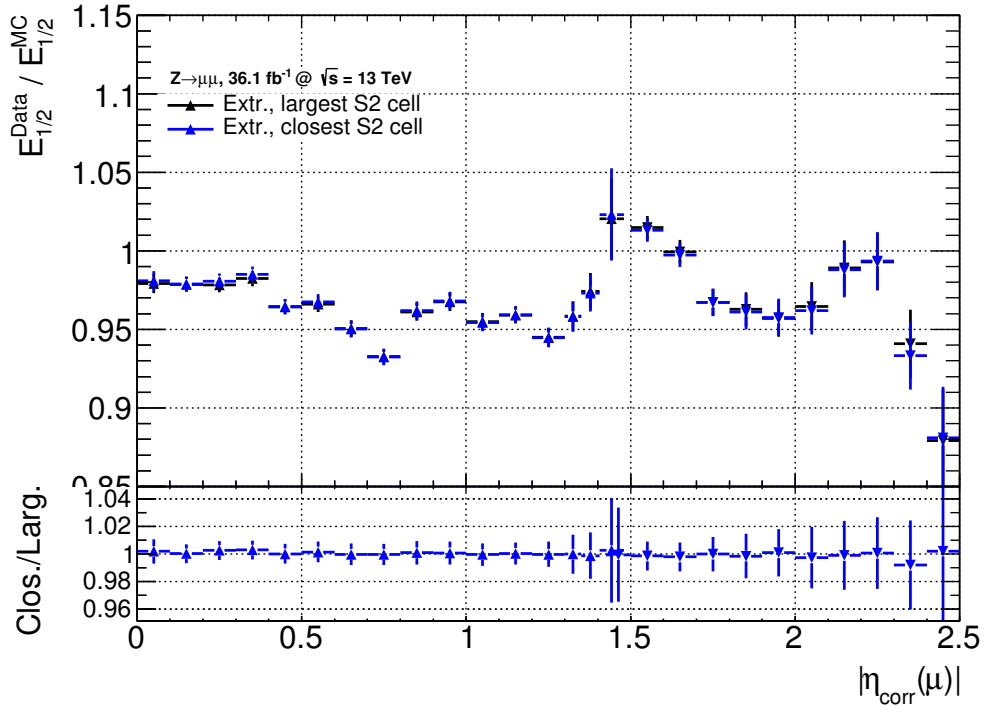

 (b) $\alpha_{1/2}$ comparison.

Figure 6.30 – Comparison of the extrapolated MPV in S2 (a) and of the double ratio $E_{1/2}^{\text{Data}} / E_{1/2}^{\text{MC}}$ (b) as a function of $|\eta_{\text{corr}}|$ when using either the second most energetic or second closest cell as the second cell choice for S2. The systematic uncertainty associated for this effect is of the order of a few permil over the whole $|\eta|$ range.

$$|\phi^{\text{extr.}} - \phi^{\text{cell}}| < 0.04.$$

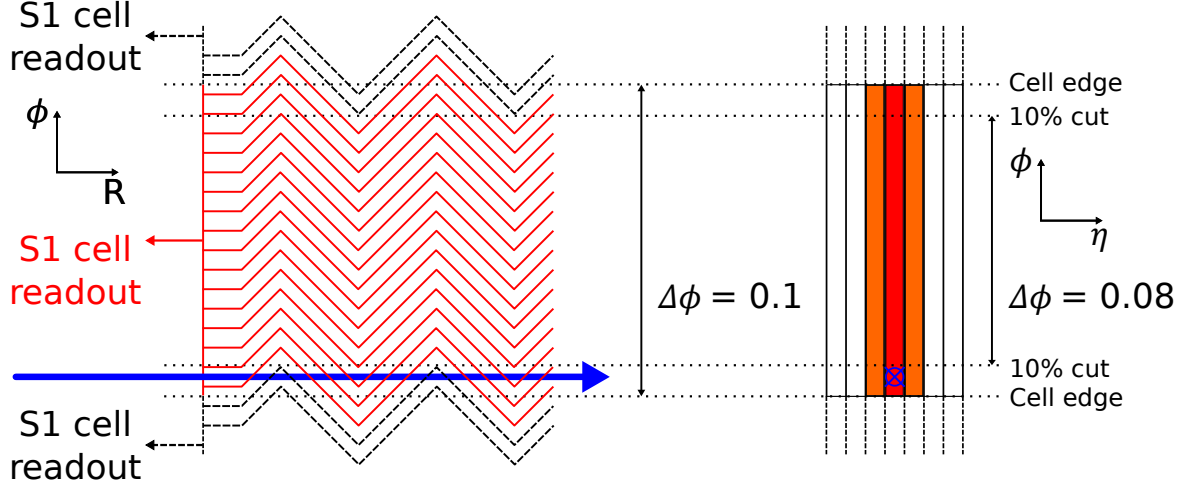


Figure 6.31 – Scheme of the folded electrodes in S1 in the (R, ϕ) plane representing one cell in the ϕ direction (left), and its associated scheme in the (η, ϕ) plane (right). The η/ϕ ratio is not at real scale (dilated by roughly a factor of 2). One S1 cell is defined by the readout of 16 consecutive electrodes. The red cells denotes the one crossed by the muon (shown as a blue arrow), and the orange cells the ones added to recover from the cross-talk effects. In this example, the muon hits the cell in the “bottom-most” part of the cell, and electrodes read by the previous cell in the ϕ direction also collect a bit of the muon energy. Modulo variations of the folding angle (90° used everywhere here for illustration purpose), the leakage occurs when the muon falls in the “top-most” or “bottom-most” 12.5 % of the cell, corresponding to a deposit in the first or last two electrodes of the cell (over a total of 16). The systematic uncertainty described here uses a cut of 10 % on the top and bottom of the cell as represented by the dotted lines.

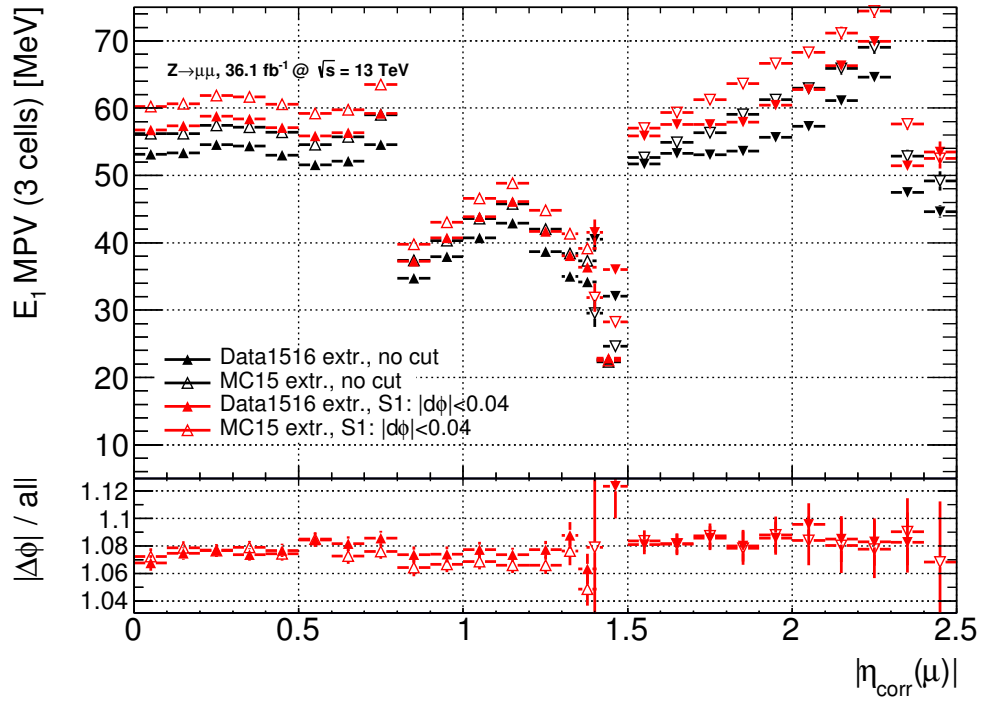
A comparison of the results applying or not the $\Delta\phi$ cut to S1 is shown in Figure 6.32. The energy deposited in the first sampling increases by about 6 MeV over the whole $|\eta|$ range after the cut is applied. This is expected since events losing energy in neighbouring cells are cut, so the average energy increases⁴. Qualitatively the effect is similar in the data and MC, with an increase of about 8 %. The only noticeable difference is in the range $|\eta| \in [0.6, 1.4]$ where the impact in the data is about 7.5 % and 6.5 % for MC. The associated systematic uncertainty is evaluated as the difference of the double ratios, as displayed in Figure 6.32b. The variation is statistically compatible with 1 over the full range, as this has similar behaviour in both data and MC. The systematic uncertainty is contained within a 0.5 % band for $|\eta_{\text{corr}}| \in [0, 0.6]$ and $[1.5, 2]$, and a 1 % band for $|\eta_{\text{corr}}| \in [0.6, 1.4]$.

6.4.6.4 Leakage energy in S2

As two cells are summed in the ϕ direction for S2, the effect described above for S1 is already taken into account. However, due to some misalignment in data, a muon near the cell boundary can share its energy deposit in two neighbouring cells in η . In order to estimate the size of this effect, a $|\eta^{\text{extr.}} - \eta^{\text{cell}}| < 0.008$ cut is applied on the muon selection.

A comparison of the results applying or not the $\Delta\eta$ cut in S2 is shown in Figure 6.33. In the barrel the effect is similar as for the $\Delta\phi$ cut in S1: a moderate increase of around 5 MeV, resulting in a $\sim 2\%$ difference. The effect is similar in both data and MC, leading to an almost negligible systematic uncertainty of a few permil on the double ratio. In the endcap, the energy difference switches sign at $|\eta| = 2.1$ (1.8 in the MC), the energy being higher with the cut in

⁴↑ From a simple geometrical argument, the 20 % events removed are expected to lose between 0 and 50 % of their energy in the neighbouring cell (25 % on average). This lead to an expected increase of $\frac{1 - (80\% \times 100\% - 20\% \times 75\%)}{80\% \times 100\%} \sim 6\%$.



(a) MPV comparison in S1.

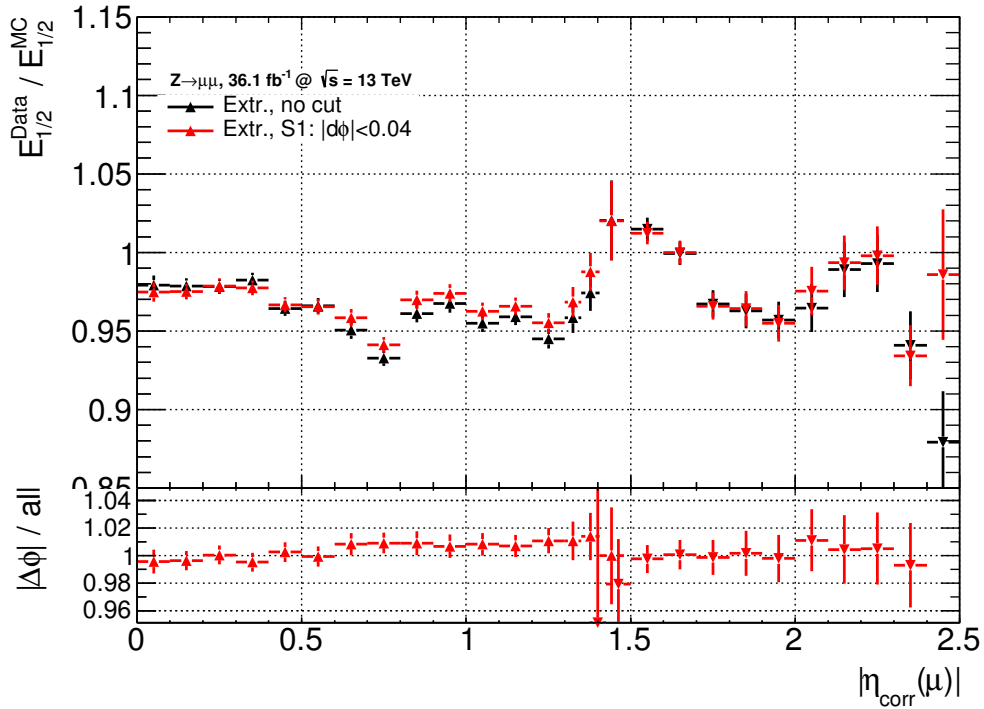

 (b) $\alpha_{1/2}$ comparison.

Figure 6.32 – Comparison of the fitted MPV (a) and of the double ratio $E_{1/2}^{\text{Data}}/E_{1/2}^{\text{MC}}$ (b) as a function of $|\eta_{\text{corr}}|$ with or without a $|\phi^{\text{extr.}} - \phi^{\text{cell}}| < 0.04$ cut applied. The systematic uncertainty associated to this leakage effect is within 0.5% in the whole range, except in $[0.6, 1.4]$ where it is around 1%.

the low $|\eta|$ region, but smaller after applying the cut in the high $|\eta|$ region⁵. This results in a systematic uncertainty on the double ratio of about 1 % to 2 % in the endcap, but this effect is not completely understood. The associated systematic uncertainties is evaluated as the difference of the double ratios, as displayed in Figure 6.33b.

6.4.6.5 Combined plot of MPV systematic uncertainties

These systematic uncertainties for the MPV method are summarised in Figure 6.34. The statistical component is taken as the quadratic sum of the statistical uncertainties from the extrapolation fit of E_1^{data} , E_2^{data} , E_1^{MC} and E_2^{MC} . The closure systematic in S2 is small enough to be neglected and was not propagated to the final systematics model. All the sources are considered uncorrelated in each bin and the final systematic uncertainty is defined as the quadratic sum of all the sources. This results in a 0.7 % to 1.3 % total uncertainty in the barrel and a 1.5 % to 3 % total uncertainty in the endcap.

For the final systematic uncertainties model, the bins 1.30-1.35 and 1.35-1.405 were averaged to give the bin 1.30-1.40, and the crack region 1.40-1.50 was taken to be last barrel bin 1.40-1.48.

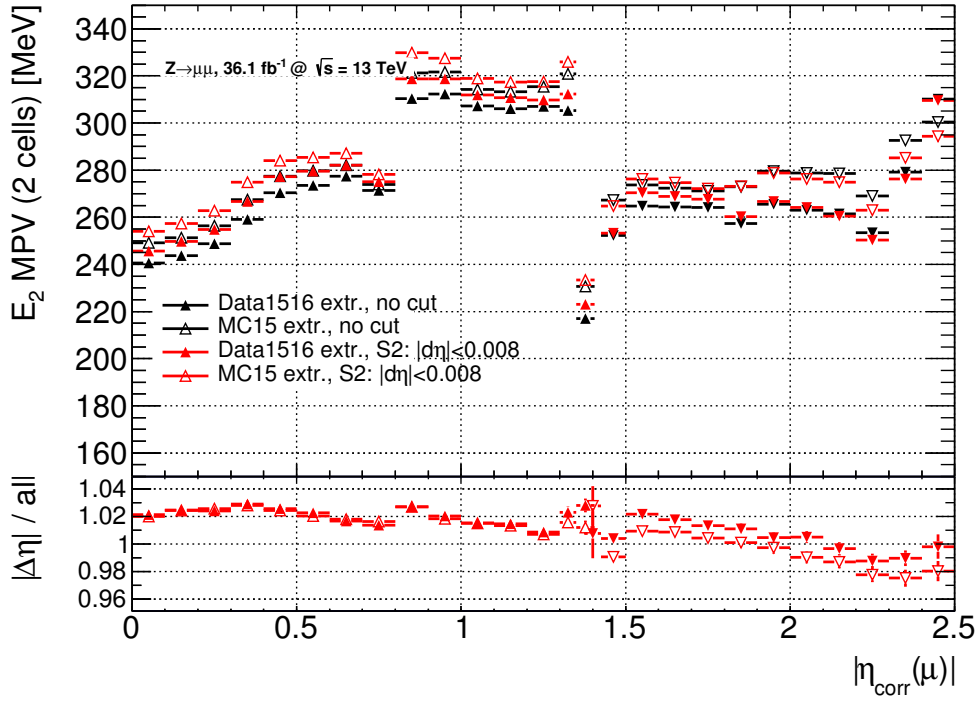
6.5 Estimation of the muon energy deposit using the truncated mean method

6.5.1 Truncated mean definition

One major issue in the fit method described above is the stability of the fit. Due to the low SNR of muon energy deposit in the calorimeter, the fitted MPV value can fluctuate and a significant amount of time and care has to be spent to control these instabilities. As a cross-check of the fitted MPV method, another method not involving a fit can be used. In this case, the parameter used to quantify the muon energy deposit is the mean of the distribution. However, since the muon deposit resembles a Landau distribution, its mean is not properly defined. The mean is therefore computed on a restricted range of the distribution, removing the tails, hence the denomination *truncated mean* (TM). This helps reducing the background contributions from noise and pileup making the energy estimate more robust against pileup effects. As a final advantage, the statistical uncertainties are shown to be reduced compared to the fit method. The drawback though is that the “mean” cannot be interpreted as the MPV of a Landau energy deposit anymore: increasing the range also increases the bias with respect to the MPV [146]. Still, the bias should be approximately the same in both data and MC, and should therefore largely cancel out in the data/MC ratio.

Two procedures have been tested to extract the truncated mean of the muon energy deposits: the first one, also used in the Run 1 analysis, is to restrict the window to the smallest range containing some fraction of the total number of events (typically 90 %). The second one is an iterative procedure: first the mean μ_0 and RMS of the distribution are computed over the range $[0, 600]$ MeV for the first sampling and $[0, 1000]$ MeV for the second sampling, then the truncated mean is computed in the range $[\mu_0 - n \cdot \text{RMS}, \mu_0 + n \cdot \text{RMS}]$ with $n = 2$ typically. These values are chosen as they provide the best closure for the MC generated with $\langle\mu\rangle = 0$. Examples of the two definitions are given in Figure 6.35.

⁵↑ The endcaps are not at their exact nominal position, ensuring projectivity with respect to the interaction point, but actually shifted by 45 mm in the z axis. As a consequence, muons hitting the endcaps wheels at higher radii (*i.e.* smaller $|\eta|$) are more likely to share energy with a η neighbouring cell.



(a) MPV comparison in S2.

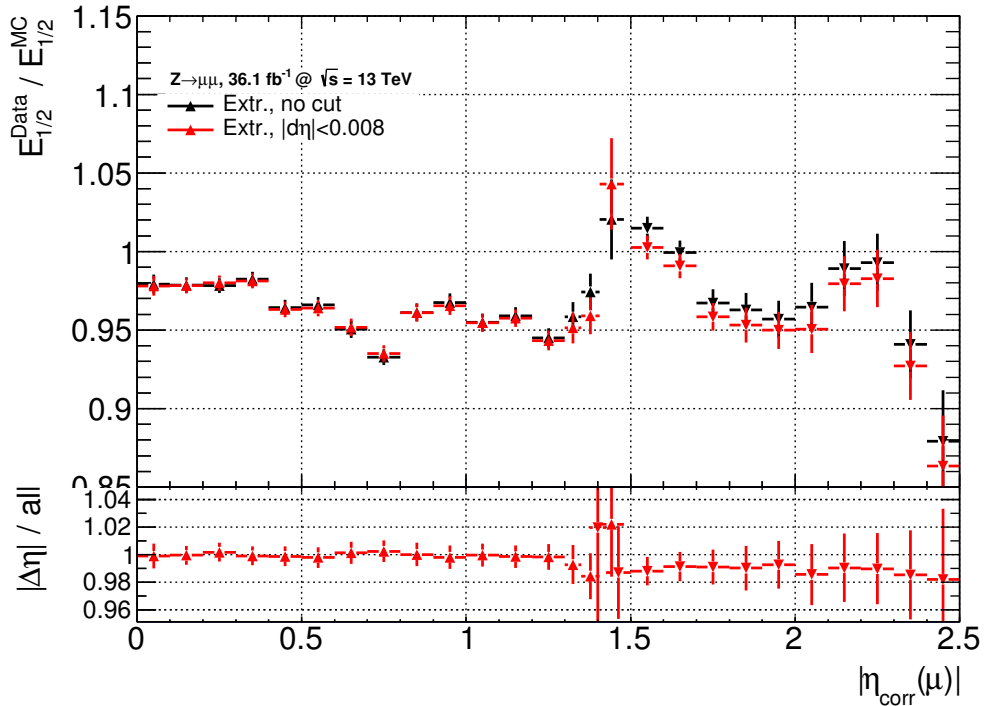

 (b) $\alpha_{1/2}$ comparison.

Figure 6.33 – Comparison of the fitted MPV (a) and of the double ratio $E_{1/2}^{\text{Data}}/E_{1/2}^{\text{MC}}$ (b) as a function of $|\eta_{\text{corr}}|$ with or without a $|\eta^{\text{extr.}} - \eta^{\text{cell}}| < 0.008$ cut applied. The systematic uncertainty associated to this leakage effect in the barrel is almost negligible at the level of a few permil, and around 1% in the endcap.

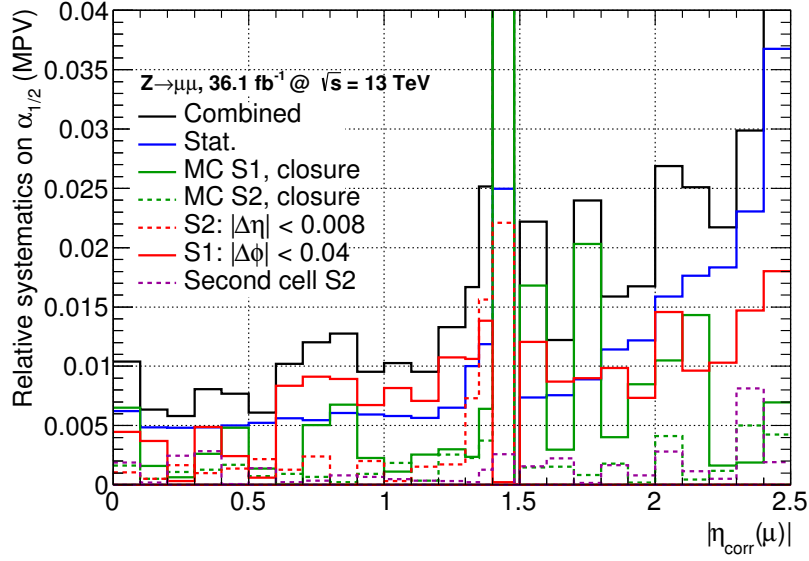
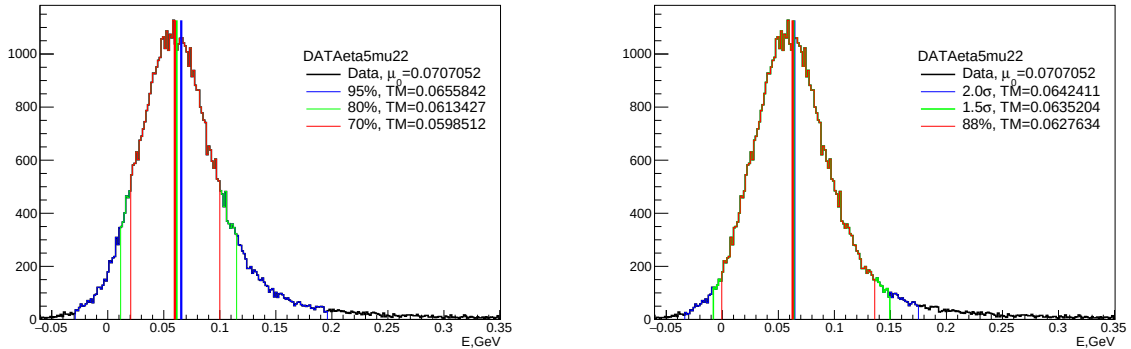


Figure 6.34 – Summary of systematic uncertainties affecting the relative layer intercalibration of the first and second layers of the electromagnetic calorimeter from the muon measurement using the MPV method as a function of $|\eta_{\text{corr}}|$. The closure systematic uncertainty are shown in green, the leakage energy in S1 and S2 in red, and the choice of second cell in S2 in purple. The quadratic sum of the extrapolation statistical uncertainties is plotted in blue. Finally the quadratic sum of all these uncertainty sources defining the total systematics associated to the measurement is in black.



(a) Truncated mean using the smallest interval definition. (b) Truncated mean using the iterative method.

Figure 6.35 – Example of muon energy distribution in the first layer of the electromagnetic calorimeter for $0.125 < |\eta| < 0.150$ and $21 < \langle \mu \rangle < 22$ using 2015 and 2016 data. Several definitions of the truncated mean are illustrated: (a), using the smallest interval containing 70 % (red), 80 % (green) or 95 % (blue) of the data, (b), using the iterative procedure with a cut on $n \cdot \text{RMS}$ for $n = 1.5$ (green) or $n = 2$ (blue), compared to the smallest interval containing 88 % of the distribution (red). [146]

6.5.2 Truncated mean results

Following the same procedure as described in Section 6.4.4, the values E_1 and E_2 in data and MC are extracted separately in several bins of pileup and then linearly extrapolated to $\langle\mu\rangle = 0$. Thanks to a higher statistical precision, the $\langle\mu\rangle$ binning used here is finer as compared to Eq. 6.4, splitting the range in bins of unit 1 over $[12, 30]$ which contains the bulk of events. An example of the pileup extrapolation is shown in Figure 6.36, and the extrapolated results as a function of $|\eta|$ in data and MC are displayed in Figure 6.37. The $\alpha_{1/2}$ scale is then computed as previously as the double ratio $E_{1/2}^{\text{extr, data}}/E_{1/2}^{\text{extr, MC}}$ (see Eq. 6.1).

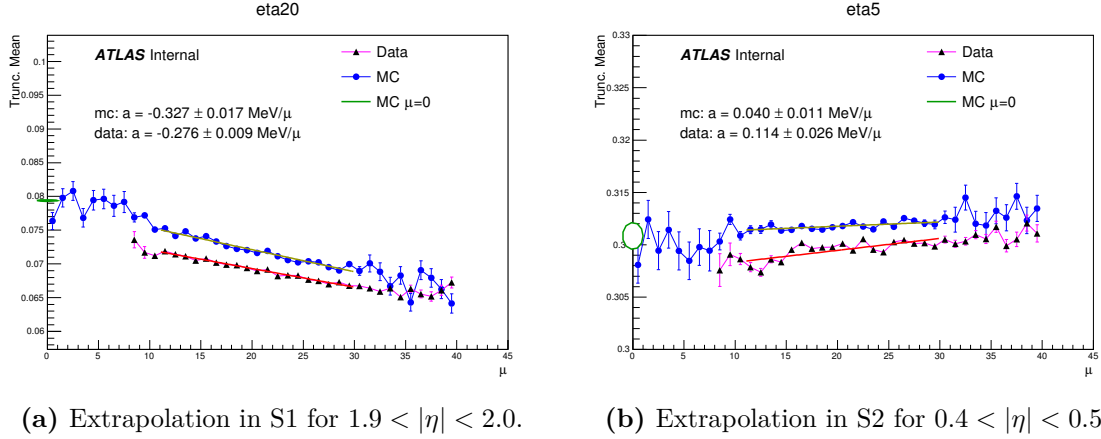


Figure 6.36 – Examples of truncated mean values as a function of $\langle\mu\rangle$ in (a) S1 and (b) S2 for both data (blue) and MC (red). Their linear extrapolation is also plotted, as well as the value found from the MC generated without any pileup ($\langle\mu\rangle = 0$, green). The TM is using the definition $\mu_0 \pm 2 \cdot \text{RMS}$. [146]

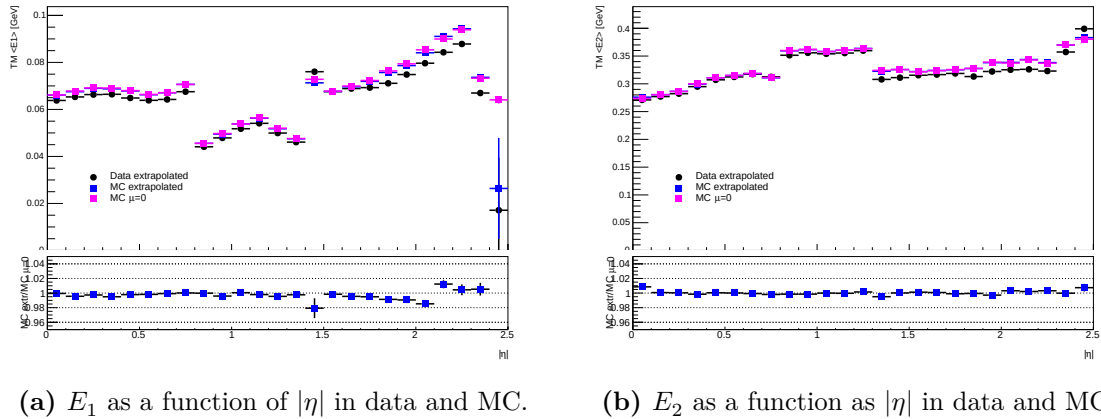


Figure 6.37 – Truncated mean pileup extrapolation results as a function of $|\eta|$ in (a) S1 and (b) S2, for data (black) and MC (blue) compared to the (non-extrapolated) result using the MC generated at $\langle\mu\rangle = 0$. The bottom insert shows the extrapolated MC result divided by the MC result generated at $\langle\mu\rangle = 0$. The TM is using the definition $\mu_0 \pm 2 \cdot \text{RMS}$. [146]

This procedure is repeated for three estimators of the mean: the iterative method with $n = 1.5$ and $n = 2$, and for the smallest range containing 90% of the total distribution. The final TM result for $\alpha_{1/2}$ is given as the average of these three definitions, as shown in Figure 6.38.

6.5.3 Systematics of the truncated mean method

Uncertainty sources for the truncated mean method are the same as the analytical MPV fit described in Section 6.4.6:

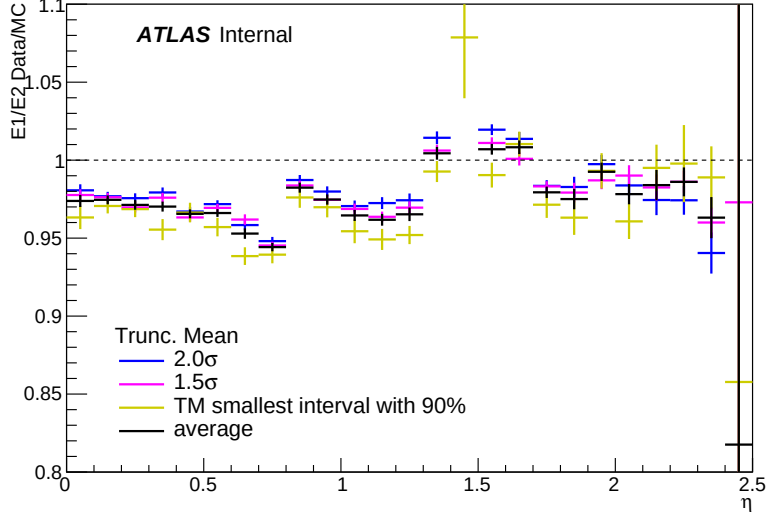


Figure 6.38 – Layer intercalibration scales $\alpha_{1/2}$ as a function of $|\eta|$ extracted with various definitions of the truncated mean: in blue using the iterative method computing the mean in the 2σ core of the distribution, in pink using the iterative method computing the mean in the 1.5σ core of the distribution, in yellow using the smallest interval containing 90% of the total distribution. The average of the three methods is shown in black and is quoted as the final $\alpha_{1/2}$ with truncated mean method. [146]

- the extrapolation closure in MC for the first and second samplings,
- the S1 leakage cuts $|\Delta\phi| = 0.1$ (this value is taken from the MPV method),
- the S2 alignment effects $|\Delta\eta| < 0.008$,
- and choice of second cell in S2.

Yet, two additional items specific to the TM method are included:

- the choice of the initial range for the truncation, varying the upper bound by ± 200 MeV,
- and the envelope of the TM definition, defined as the distance between the average and the most distant point of the 3 definitions described above.

The summary of the TM systematic uncertainties as a function of $|\eta|$ is displayed in Figure 6.39 and show a total systematic uncertainty of 1% to 2% in the barrel and 1.5% to 3% in the endcap. The systematic uncertainties are dominated by the TM definition envelope and the alignment systematics in the second layer, with a non-negligible contribution from the leakage systematic in the first layer (around 1% each). The pileup extrapolation closure between the extrapolated MC and the MC generated at $\langle\mu\rangle = 0$ is good and therefore has a subdominant effect, as well as the choice of initial range which only has a moderate impact in the endcap. The choice of second cell in S2 is almost negligible over all the range.

A comparison of the MPV and TM methods is provided in the next Section (6.6).

6.6 Comparison of muon MPV and TM method, and to Run 1 result

In order to provide one single result for the layer intercalibration to the central calibration procedure, the two methods (analytical MPV fit and truncated mean estimate) are combined as a

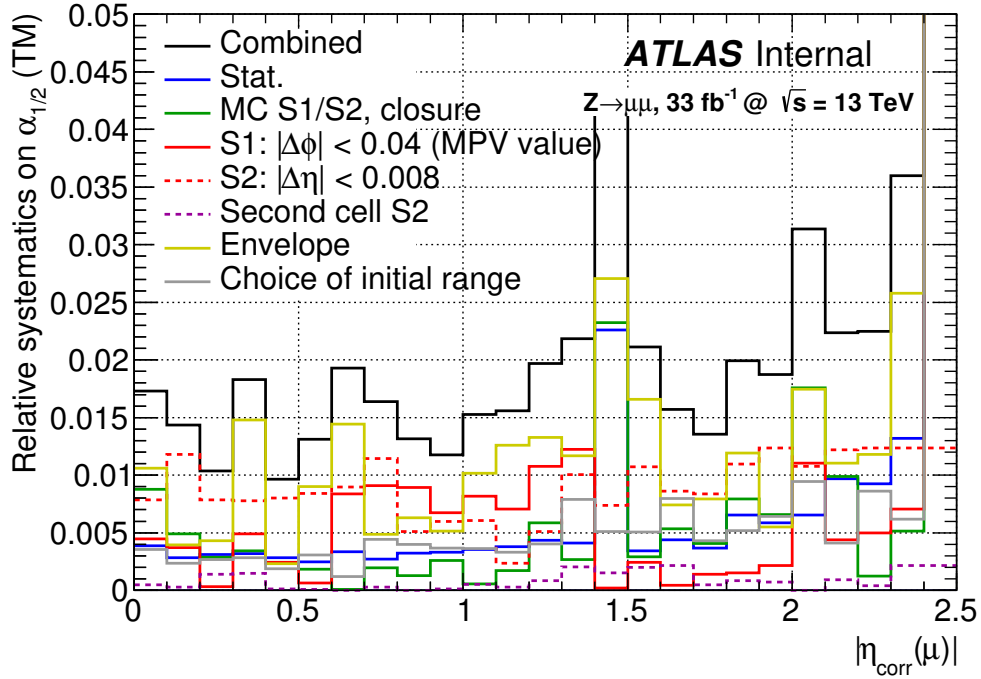


Figure 6.39 – Breakdown of layer intercalibration uncertainties with TM method. [146]

simple average. The systematic uncertainty is a combination of those of the individual methods (Section 6.6.1). The final $\alpha_{1/2}$ scales are then derived in Section 6.6.2, and a comparison to the Run 1 result is shown in Section 6.6.3.

6.6.1 Comparison of MPV and TM uncertainties in Run 2

This Section provides a comparison of the various the systematic uncertainty sources for the MPV (Section 6.4.6) and TM (Section 6.5.3) methods, which are recalled in Table 6.2.

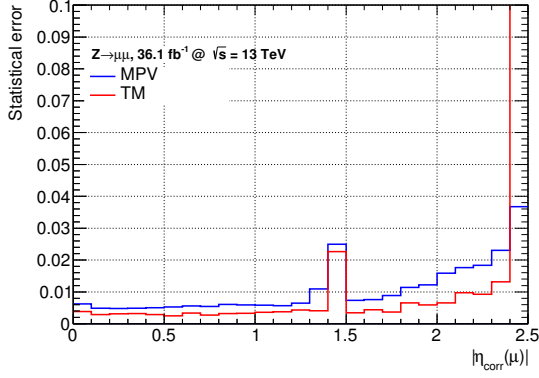
The S2 pileup extrapolation closure is evaluated but found to be negligible and is therefore pruned from the systematic uncertainties combination in the MPV and TM methods. The S1 leakage uncertainty evaluated through the $\Delta\phi$ cut was evaluated only with the MPV method and is therefore taken directly for the combination (see Figure 6.32b). The TM method has two additional specific systematic uncertainty sources consisting in variations of the truncation definition.

The plots in Figure 6.40 show comparisons of the systematic uncertainties having effects in both MPV and TM methods. The statistical uncertainty is larger for the MPV than for the TM method since a two-parameter fit is performed (the template has a negligible impact on the statistical uncertainty). This uncertainty is around 0.6 % (0.4 %) in the barrel for the MPV (TM) method, and varies from 0.7 % to 2.3 % (0.5 % to 1.3 %) in the endcap. The closure in S1 has a similar impact for both method, below 0.5 % in the barrel and 1 % to 2 % in the end cap (crack and last bin excluded). The choice of second cell is small (within 0.3 %) and has a similar impact in both methods. The S2 alignment systematic ($|\Delta\eta|$ cut) is 0.5 % to 1 % for the TM but below 0.2 % for the MPV method in the barrel, and compatible at more than 1 % to 1.5 % in the endcap.

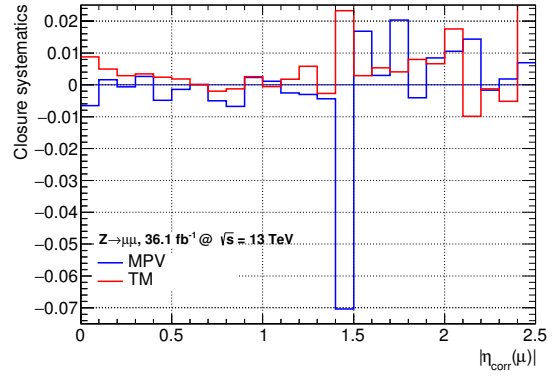
A final source of systematic uncertainty representing the difference of the MPV and TM central values is taken into account. The value of this systematic uncertainty is shown in Figure 6.41.

Table 6.2 – Summary of systematic uncertainties included in the MPV and TM evaluation. The S2 closure is found to have negligible contribution and is pruned from the two methods. The leakage systematic in S1 is included in both methods but was only evaluated with the MPV.

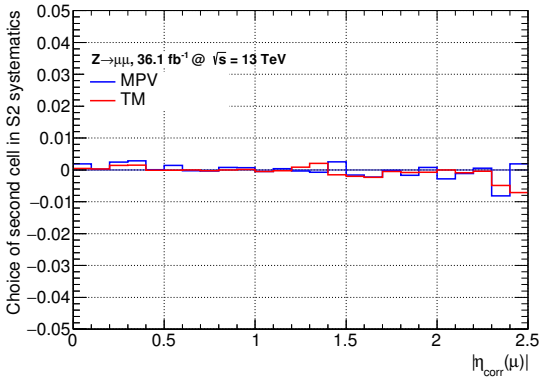
Source of uncertainty	Included in MPV	Included in TM
Statistics	yes	yes
Closure in S1	yes	yes
Closure in S2	no	no
Choice of second cell in S2	yes	yes
Leakage in S1 ($ \Delta\phi < 0.04$)	yes	same as MPV
Alignment in S2 ($ \Delta\eta < 0.008$)	yes	yes
Truncation range variation	-	yes
Truncation method envelope	-	yes



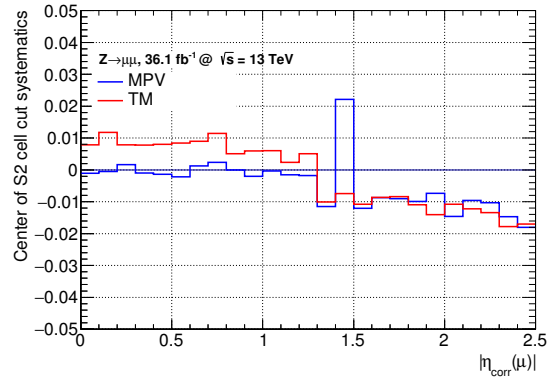
(a) Statistical uncertainty.



(b) Closure of the S1 extrapolation.



(c) Choice of second cell in S2.



(d) $|\Delta\eta| < 0.008$ cut in S2.

Figure 6.40 – Comparison of the common sources of uncertainty for the MPV and TM methods. For other sources, see Sections 6.4.6 and 6.5.3 for the MPV and TM methods respectively. The $|d\phi| < 0.04$ cut in S1 systematic is applied in both the MPV and TM methods, but was only evaluated with the MPV so not shown here. [146]

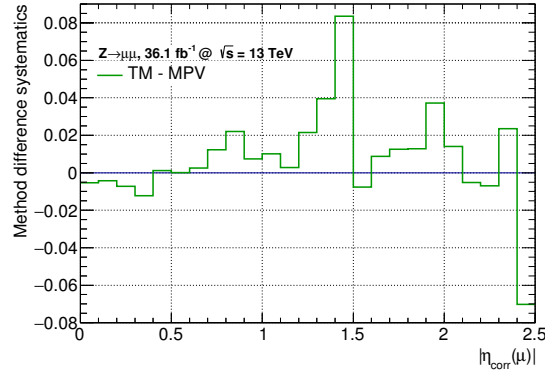


Figure 6.41 – Difference TM - MPV of $\alpha_{1/2}$ central value. [146]

6.6.2 Final result: Run 2 combination of the MPV and TM methods

Since the final $\alpha_{1/2}$ scales are computed as the average of the MPV and TM results, the final $\alpha_{1/2}$ scale uncertainty is the quadratic sum of the averaged uncertainties of the MPV and TM methods, and half the difference of the TM and MPV scale results. This is given by:

$$\alpha_{1/2} = \frac{\alpha_{1/2}^{\text{MPV}} + \alpha_{1/2}^{\text{TM}}}{2} \quad (6.6)$$

$$\sigma = \frac{\sigma_{\text{MPV}} + \sigma_{\text{TM}}}{2} \oplus \frac{|\alpha_{1/2}^{\text{TM}} - \alpha_{1/2}^{\text{MPV}}|}{2}. \quad (6.7)$$

The individual MPV and TM systematic uncertainties, the method difference, and the final combined uncertainty are shown in Figure 6.42.

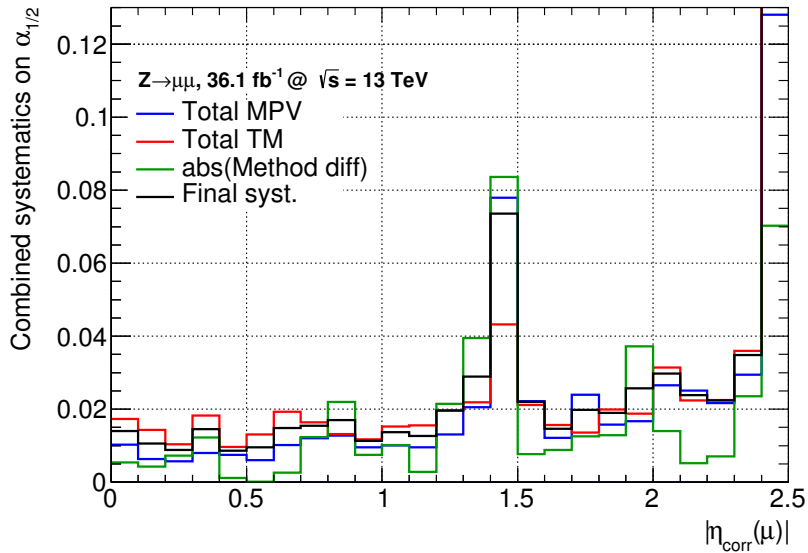


Figure 6.42 – Summary of systematic uncertainties for the MPV method (blue) and the TM method (red). The difference between the MPV and TM central value is shown in green. The final systematics applied on the $\alpha_{1/2}$ measurement amounts to the black curve, which is computed combining the three contributions above with the Eq. 6.7.

The final results along with the MPV and TM individual results are shown in Figure 6.43,

where the individual MPV and TM have only their statistical uncertainties shown. A fair agreement between the two methods is found across most of the whole $|\eta|$ range.

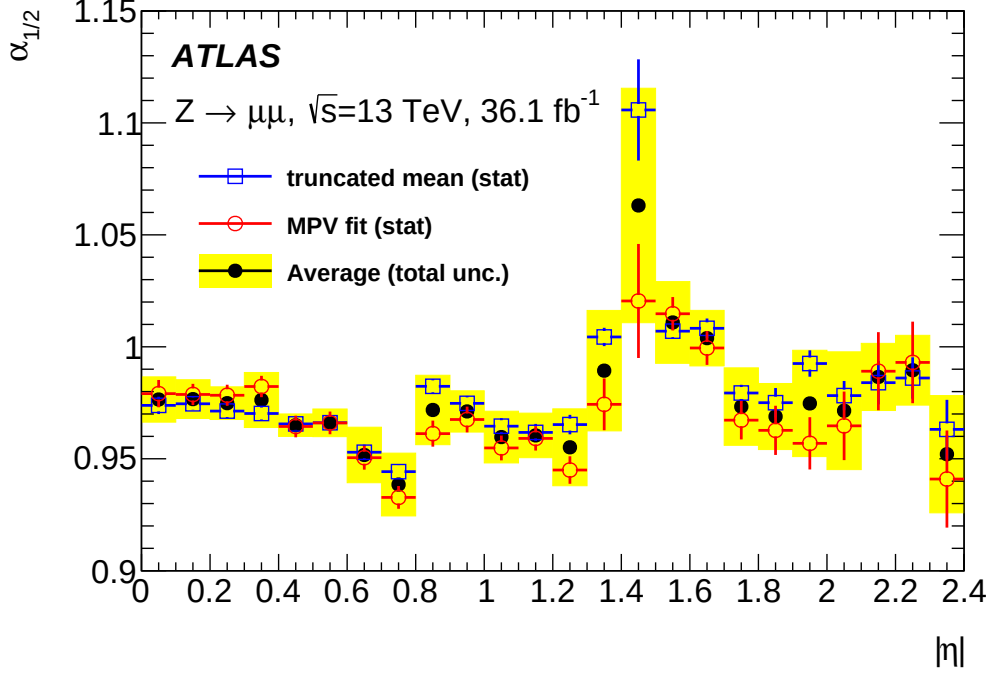


Figure 6.43 – Comparison of the layer intercalibration scales $\alpha_{1/2} = E_{1/2}^{\text{Data}}/E_{1/2}^{\text{MC}}$ using the 2015+2016 dataset, as a function of $|\eta_{\text{corr}}|$ for the MPV method (red, see Section 6.4), the TM method (blue, see Section 6.5), and their average (black). The error bars on the MPV and TM points are statistical only. The yellow band represents the total systematic uncertainties on the combined MPV and TM measurement, described in Section 6.6.1. [141]

6.6.3 Run-1 / Run-2 comparison

A comparison between the Run 2 and the Run 1 results (both averaging over MPV and TM methods) is provided in Figure 6.44. The two results display the same behaviour as a function of $|\eta|$ with a global offset of about 1 %, but are largely compatible within uncertainties. The larger uncertainty in the 2015+2016 result emerges from the extrapolation method. The offset can be explained by a different number of cells used in S2 for the Run-1 (1 cell) and the Run-2 (2 cells) analyses: this difference proved to give results 0.5 % higher in the barrel and 1 % higher in the endcap, as demonstrated in Figure 6.45.

6.7 Cross-check: electron method

Contrary to muons, electrons produce an electromagnetic shower in the calorimeter and are sensitive to upstream material effects which this is precisely why they are used in the material determination analysis described in Section 5.6. Measuring E_1/E_2 with electrons in the same way as it is done with muons therefore does not yield accurate conclusions on the layer intercalibration. The idea is to instead make use of a standard candle observable that depends on E_1/E_2 , and whose dependency should be well modelled in the MC. In this analysis the well-known $Z \rightarrow ee$ decay [145] is used to provide an electron pair whose invariant mass should statistically be the same as a function of $E_{1/2}$ in both data and MC.

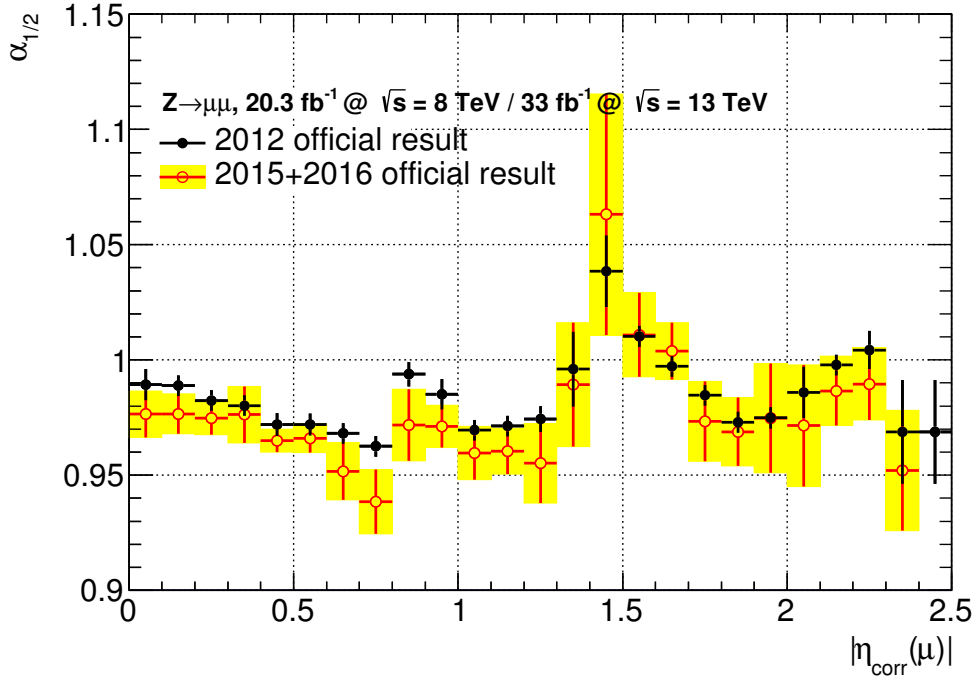


Figure 6.44 – Comparison the layer intercalibration scales $\alpha_{1/2} = E_{1/2}^{\text{Data}} / E_{1/2}^{\text{MC}}$ as a function of $|\eta_{\text{corr}}|$ in the Run-1 and the 2015+2016 analyses. The open red points show the result for the 2015+2016 analysis, averaging over the MPV and TM methods (same as the black point of Figure 6.43). The yellow band correspond to the total uncertainty (statistic $\oplus$ systematics) on the layer scales. The filled black points remind the Run-1 muon layer intercalibration scales, also averaged over the MPV and TM methods.

6.7.1 E_1/E_2 follow up with $Z \rightarrow ee$ electrons

The event selection in this analysis requires two opposite sign electrons of medium quality with transverse energy $E_T > 27$ GeV, and whose invariant mass falls in the window $[75, 105]$ GeV around the Z mass. To account for fine variations of the detector geometry, the analysis is performed in bins of $\Delta\eta = 0.05$ up to $|\eta| = 2.5$, folding the $\eta < 0$ and $\eta > 0$ sides together. In order to remove a potential bias from a PS miscalibration, the corrections derived in Section 5.4 are applied on the E_0 energy prior to this analysis.

The average di-electron invariant mass⁶ is analysed in bins of $E_{1/2}$ for both the data and MC (Figure 6.46a). The ratio $\langle M_{ee}^{\text{data}} \rangle / \langle M_{ee}^{\text{MC}} \rangle$ (Figure 6.46b) is then fitted as a linear function of $E_{1/2}$: $\frac{\langle M_{ee}^{\text{data}} \rangle}{\langle M_{ee}^{\text{MC}} \rangle} = a \times E_{1/2} + b$. If the layers are properly intercalibrated, the average mass dependence on $E_{1/2}$ should be the same in data and MC, and the linear slope fit should yield $a = 0$. Otherwise, the data is then modified by scaling the first sampling energy (E_1) by some amount and recalibrating the total electron energy assuming this modified E_1 energy. The fit procedure is then repeated with this modified data (and unchanged MC) for several values of the E_1 energy scaling (Figure 6.46b). The last step to get the $\alpha_{1/2}$ scales from the slopes is to plot the slope as a function of the correction and to interpolate between points to find the point where the slope vanishes.

As a validation of the independence on the geometry, the ratio of average mass in MC with distorted geometry⁷ against the MC with nominal geometry can be fitted. This is illustrated by

⁶↑ The same results are obtained using the peak of the E/p distribution instead of the average invariant mass, and applying the same procedure.

⁷↑ The distorted geometries are different simulations with additional material in front of the calorimeter (see

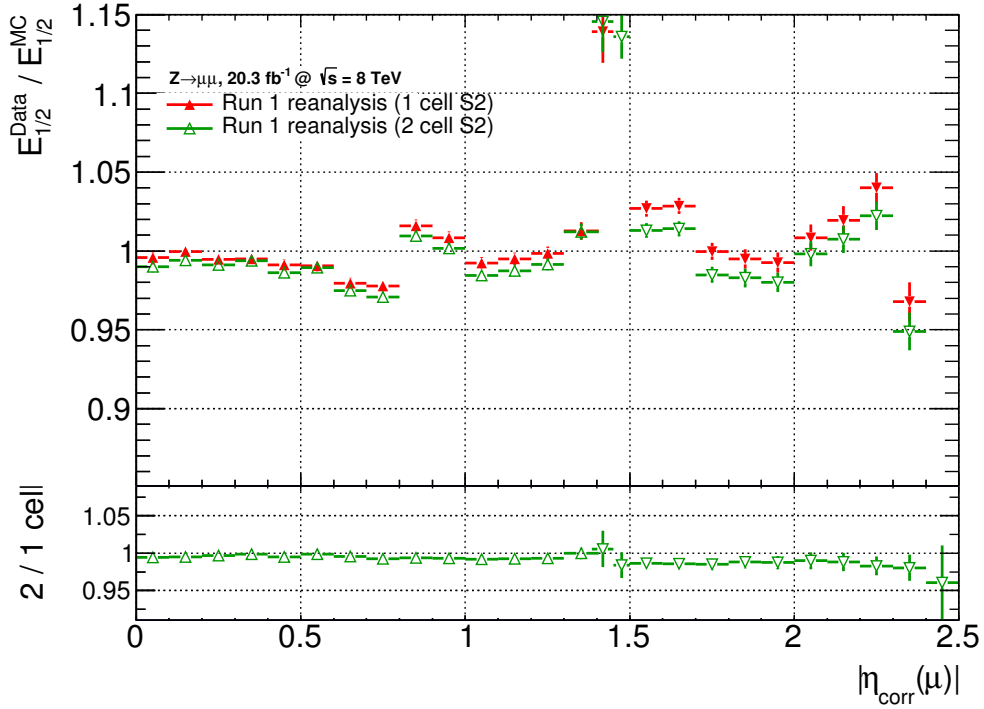


Figure 6.45 – Comparison of the layer intercalibration scales $\alpha_{1/2}$ using the MPV method on Run 1 data, with one (red points) or two (green points) cells in the energy determination of the muon deposit in S2. The result using two cells is 0.5 % to 1 % lower than the one-cell analysis in the barrel region, and 1.5 % to 2 % lower in the endcap region. This plot is for illustration purpose only: this a “re-analysis” so nor the same data nor the same analysis code as for the actual Run 1 analysis were used, therefore a difference with the official Run 1 result is expected (the two results presented in this plot form a consistent set though).

the pink curve in Figure 6.46: a change in the material upstream of the calorimeter, shifts the average M_{ee} energy and therefore the energy ratio in each bin of $E_{1/2}$. However the slope of the ratio to the nominal MC distribution is compatible with 0, which indicates that the E_1/E_2 dependency of the di-electron invariant mass is the same independently of the material [140].

Another possible reason for the non-vanishing slope could be a mismodelling of the cross-talk between the first and second layer. However such cross-talk would bias the E_1/E_2 ratio, but not the total energy $E_1 + E_2$. This can be checked by scaling E_1 , but keeping $E_1 + E_2$ constant (so decreasing E_2). This effect is ruled out as seen from the black and red points of Figure 6.46: the slope does not change if $E_1 + E_2$ is renormalised after the E_1 scaling. In conclusion, the slope present in the m_{ee} data/MC ratio as a function of E_1/E_2 is due to a genuine $E_{1/2}$ miscalibration, as shown by the green curve in Figure 6.46b.

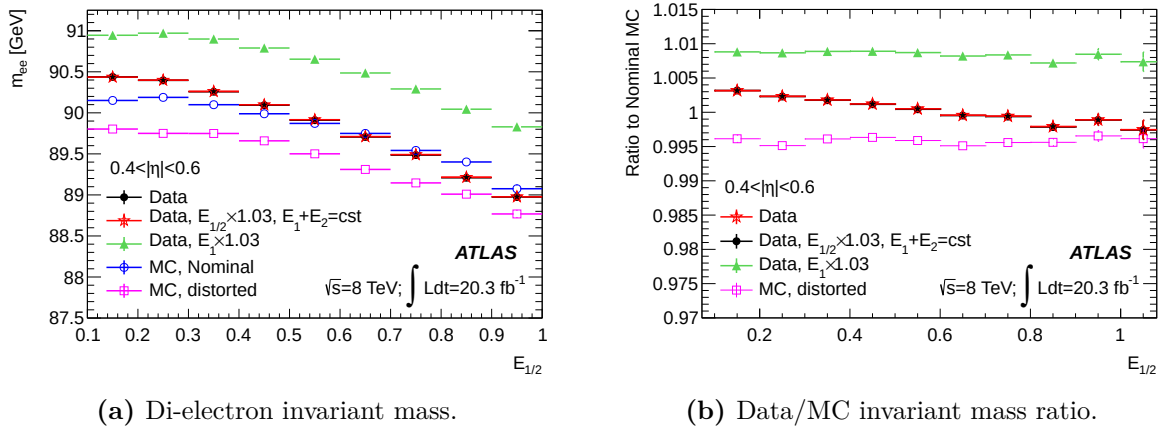


Figure 6.46 – (a) average di-electron invariant mass (b) and ratio of average invariant mass to average invariant mass in nominal MC, as a function of the energy ratio in S1 and S2 ($E_{1/2}$) for $0.4 < |\eta| < 0.6$. The next step in the procedure is to linearly fit the slope in (b) for various values of the scaling applied to E_1 . This plot is for illustration of the method and was taken from the Run 1 result in [140].

6.7.2 Results and compatibility with the muon method

The scales extracted with the electron and muon methods from above are compared in Figure 6.47. The two methods are in fair agreement over the whole $|\eta|$ range, except for the region $|\eta| \in [1.2, 1.8]$ (called *extended crack*), similarly to the Run 1 result presented in Ref. [140].

Several investigations were performed, and their summary is presented in Figure 6.48. As the shower shape modelling in MC is suspected to be not perfectly accurate in this region, the analysis was performed again after applying a cut on shower shape variables such as R_ϕ , R_η or $w_{s1,\text{tot}}$ (see Chapter 4.4). As seen in bins 3–6 of Figure 6.48, the slope of the invariant mass ratio is not much improved after applying these cuts and the needed bias on E_1 to achieve a vanishing slope remains at unreasonable levels. Similarly a requirement on electron energy lost by bremsstrahlung (bin 8) or on the fraction of energy in the third sampling (bin 10) has been tested without showing significant improvement on the result. Such differences have also been noticed during the Run 1 analysis but their impact might have been underestimated.

Applying a tighter cut on the mass window used to compute the invariant mass mean seems to impact the dip in the extended crack. Comparing the bins two and three with the first bin in Figure 6.48, the slope gets closer to 0 as the window becomes tighter, moving from the standard [75, 105] to [80, 100] then to [85, 96] GeV. This improvement can be further seen in Figure 6.49: in the region $1.2 < |\eta| < 1.8$, the slopes moves closer to 0 while in the rest of the $|\eta|$ range, none or

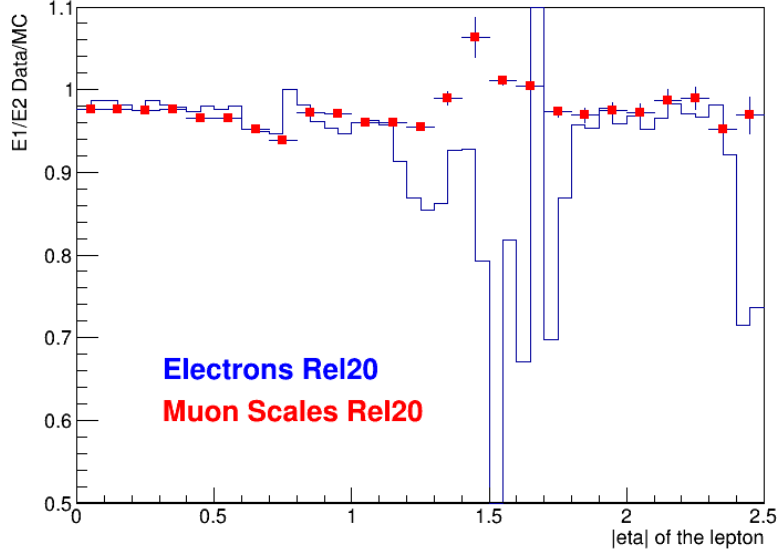


Figure 6.47 – Comparison of the $\alpha_{1/2}$ scales derived with the electron (blue) and muon methods (red). The muon points are the same as the central values in Figure 6.43. The agreement between the two methods is correct over most of the $|\eta|$ range except in the region $1.2 < |\eta| < 1.8$ and for $|\eta| > 2.4$. [146]

little change is seen. This indicates that the cause of this method’s discrepancy is a mismodelling of the electromagnetic shower of electrons in the tails of the $Z \rightarrow ee$ peak, and might be linked to the oscillation seen in $Z \rightarrow ee$ lineshape after central scale corrections seen in Figure 5.16.

The exact origin of the difference between the muon and the electron method are not yet understood, and further investigations are needed.

Conclusion of the layer intercalibration analysis

The layer intercalibration of the electromagnetic calorimeter is necessary to correct residual electronics and cross-talk effects in the calorimeter. The procedure defined during Run 1 was applied on 2015 and 2016 data, but showed an unexpected difference in the results, especially in the endcaps. The cause of this difference was tracked down to the pileup, higher during Run 2 than Run 1.

As the intercalibration scale factors should be independent of the pileup, a new procedure was implemented to mitigate its effects. First, the modelling of the noise in the calorimeter was changed from an analytical Gaussian to an actual noise template. Second, the extracted muon energies were extrapolated across the pileup to mimic a dataset without pileup. The layer intercalibration scales were then computed using these values. As in Run 1, both the MPV fit and truncated mean methods were explored, and both used this extrapolation method. The extrapolation method was also validated using simulations without pileup, and a good agreement was found.

Systematic uncertainties have been defined according to this new procedure, yielding an uncertainty of the same order of magnitude as in the Run 1 analysis despite a more complex analysis. The central values are in agreement with the expectations from Run 1.

As in Run 1, a difference with the electron method is seen in the region $1.2 < |\eta| < 1.8$. The origin of this difference is not yet explained but seems to be linked to the modelling of the tails in the $Z \rightarrow ee$ invariant mass distribution.

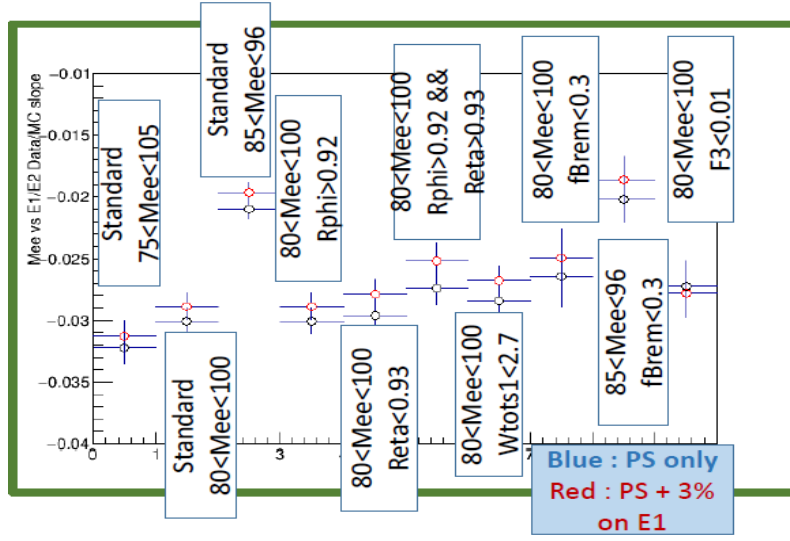


Figure 6.48 – Slope of the data/MC distribution in the inclusive bin $1.5 < |\eta| < 1.8$ for various cuts applied on top of the standard analysis selection (bin 1). Bins 2 and 3: after applying a tighter $Z \rightarrow ee$ mass window cut ($[80, 100]$ and $[85, 96]$ GeV respectively). Bins 4 and 5: after applying a cut on the shower shape with $R_\phi > 0.92$ or $R_\eta < 0.93$, or both of them (bin 6). Bin 7: after applying a cut on the shower shape with $w_{\text{tot},s1} < 2.7$. Bin 8: after requiring a low fraction of the electron energy lost by bremsstrahlung. Bin 9: after requiring the fraction of electron energy deposited in the third layer is negligible. The blue points include the PS calibration only, while the red points include both the PS calibration and a 3% bias on E_1 in the data. [146]

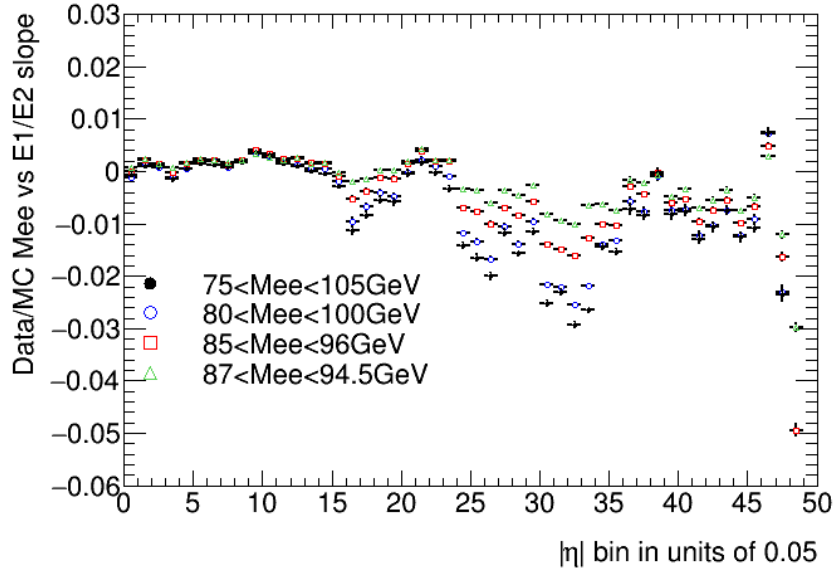


Figure 6.49 – Slope of the data/MC distribution ratio after applying a 3% bias on E_1 in data, as a function of $|\eta|$. The various curves show different M_{ee} ranges applied in the event selection, ranging from $[75, 100]$ GeV (black) to $[87, 94.5]$ GeV (green). The region $1.2 < |\eta| < 1.8$ is impacted by a large amount while other region are not or scarcely impacted. [146]

Higgs boson phenomenology and EFT

Contents

7.1	Theoretical aspects of the SM Higgs boson	176
7.1.1	Higgs boson production modes at pp colliders	176
7.1.1.1	Gluon fusion: ggF	176
7.1.1.2	Vector-boson fusion: VBF	177
7.1.1.3	Associated production with a vector boson: VH	178
7.1.1.4	Associated production with a $t\bar{t}$ or $b\bar{b}$ pair: $t\bar{t}H/b\bar{b}H$	179
7.1.2	Decay channels	179
7.1.3	Spin and parity	183
7.2	Beyond the Standard Model	184
7.2.1	The κ -framework	185
7.2.2	Simplified Template Cross-Sections (STXS)	186
7.3	Effective Field Theories: EFT	186
7.3.1	Energy power counting	186
7.3.2	A simple example: the Fermi theory of β decay	188
7.3.3	Standard Model EFT and the Warsaw basis	189
7.3.4	The Higgs basis	191
7.3.5	Basis translation and the SMEFTSIM package	192
7.3.6	Higgs Characterisation framework	192
7.4	Higgs couplings and CP-mixing results from LHC Run 1 and Run 2	194
7.4.1	STXS results	194
7.4.2	κ -framework	194
7.4.3	EFT interpretation	196
7.4.4	Direct CP measurements and couplings structure functions	198
7.4.5	Direct CP measurement with single parameter parametrisation of the effective Lagrangian	200

7.1 Theoretical aspects of the SM Higgs boson

An extended review of the SM Higgs boson properties such as production modes, decay channels, spin-parity and couplings is compiled in the CERN LHC Higgs Cross Section Working Group Yellow Reports [34, 72, 169, 170]. This Section recalls the essential information for the following $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ analysis described in Chapters 8 and 9.

7.1.1 Higgs boson production modes at pp colliders

Since protons are composite particles formed of so-called *partons* (see Chapter 2.9), particles as the Higgs boson do not strictly speaking emerge from a direct proton interaction but from the partons. The main Higgs boson production processes are therefore labelled differently, depending on what were the actual initial interacting partons. They are commonly referred to as gluon-gluon fusion (ggF, Section 7.1.1.1), vector-boson fusion (VBF, Section 7.1.1.2), associated production with a vector boson (VH , Section 7.1.1.3), and associated production to a top quark pair (ttH , Section 7.1.1.4).

Each production mode has a specific associated cross-section, which depends on the centre-of-mass energy of the collision (see Figure 1.18) and on the Higgs boson mass, as shown in Figure 7.1. Since the analysis developed in the following Chapters uses Run 2 data, the rest of this Section assumes a centre-of-mass energy of $\sqrt{s} = 13$ TeV for quantities depending on $\sqrt{s}$, unless explicitly stated. Their corresponding cross-sections, uncertainty and fraction of the total $pp \rightarrow H$ cross-section at $\sqrt{s} = 13$ TeV are summarised in Table 7.1.

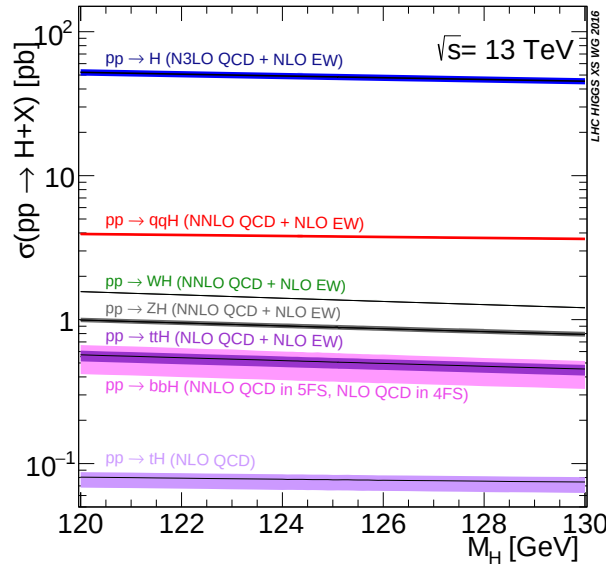


Figure 7.1 – SM Higgs boson production cross-section at pp colliders at $\sqrt{s} = 13$ TeV as a function of its mass, with main production mode breakdown. From top to bottom: (blue) ggF, (red) VBF, (green and grey) VH , (purple, pink and light violet) ttH . [72]

7.1.1.1 Gluon fusion: ggF

The dominant contribution to the Higgs boson production at pp colliders is the gluon fusion (ggF), accounting for about 87 % of the total $pp \rightarrow H$ cross-section. While gluons are massless and therefore do not couple to the Higgs boson at tree level, they can indirectly interact through loops, as pictured in Figure 7.2. Any quark can enter the loop, but since the coupling to the Higgs

Table 7.1 – Individual production mode cross-sections contributing to the $pp \rightarrow H$ process and their related uncertainties, for a SM Higgs boson of mass $m_H = 125$ GeV at $\sqrt{s} = 13$ TeV. Additional production with one associated top quark (tH) accounting for a few permil is not included. The VH cross-section is split into WH and ZH , and the sub-contributions corresponding to their leptonic decay modes ($VH\text{-lep}$) are detailed ($\ell = e, \mu, \tau$). $gg \rightarrow ZH$ and $Z(\rightarrow \ell^+ \ell^-)H/Z(\rightarrow \nu \bar{\nu})H$ are not mutually exclusive. The theory uncertainties include the factorisation and renormalisation scale variation for all production mode, while the ggF mode include more uncertainty sources [171]. [72, 172]

	σ [fb]	Uncertainties		Accuracy	$\sigma/\sigma_{\text{tot}}$
		Theory	PDF + α_s		
ggF	48 580	+4.56 % −6.72 %	±3.19 %	N ³ LO QCD + NLO EW	87 %
VBF	3782	+0.4 % −0.3 %	±2.1 %	approx. NNLO QCD + NLO EW	6.8 %
WH	1373	+0.5 % −0.7 %	±1.9 %	NNLO QCD + NLO EW	2.5 %
$W^+(\rightarrow \ell^+ \nu)H$	94.26	+0.5 % −0.7 %	±1.8 %		0.17 %
$W^-(\rightarrow \ell^- \bar{\nu})H$	59.83	+0.4 % −0.7 %	±2.0 %		0.11 %
ZH	883.9	+3.8 % −3.1 %	±1.6 %	NNLO QCD + NLO EW (ggZH: NLO QCD)	1.6 %
$gg \rightarrow ZH$	122.7	+25.1 % −19.8 %	±2.4 %		0.22 %
$Z(\rightarrow \ell^+ \ell^-)H$	29.82	+3.8 % −3.1 %	±1.6 %		0.05 %
$Z(\rightarrow \nu \bar{\nu})H$	177.6	+3.8 % −3.1 %	±1.6 %		0.31 %
ttH	507.1	+5.8 % −9.2 %	±3.6 %	NLO QCD + NLO EW	0.91 %
bbH	488.0	+20.2 % −23.9 %		NNLO (NLO) QCD for 5FS (4FS)	0.88 %
Total	55 690				

boson is linear with their mass, leading contributions come from the top quark, and subleading from the bottom quark.

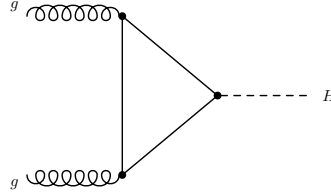


Figure 7.2 – Leading order Feynman diagram for the $gg \rightarrow H$ process.

At higher orders, the two gluon lines and the quark loop can emit other gluons that will produce additional jets. One can therefore differentially tag the ggF production mode by ggF + 0 jet, ggF + 1 jet, ggF + 2 jets, *etc.* The cross-section value provided in Table 7.1 is inclusive in the number of jets. In the ggF + 0 jet case, the Higgs boson is produced almost at rest (low p_T), but in case of higher jet multiplicity, the Higgs boson recoils against the additional jets and acquires a larger transverse momentum.

Similarly to the VBF and VH -had modes described below, the ggF + 2 jets has two additional jets in addition to the Higgs boson and can mimic their experimental signature. Since, for example, Higgs boson couplings measurements depend on the production mode, having a performing tagging between same-signature modes will be a challenge in many analyses.

7.1.1.2 Vector-boson fusion: VBF

The first subleading contribution to the total $pp \rightarrow H$ cross-section, is vector-boson fusion (VBF). Despite being a tree level process, because it occurs through electroweak vertices, its cross-section

is one order of magnitude lower than for the ggF process. Its topology with two forward jets and a large rapidity is a characteristic signature, and since there is no colour exchange (no gluon involved), reduced central jet activity is expected. An example diagram is shown in Figure 7.3.

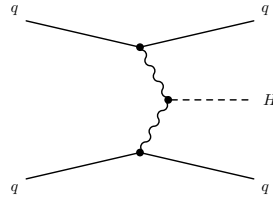


Figure 7.3 – Leading order diagram for the $qq \rightarrow qqH$ process.

While its topology should be enough to distinguish it from other production modes, higher order QCD corrections to the ggF process lead to ggF + 2 jets with jet topologies similar to VBF cases, diluting the VBF category purity. This production mode is also sensitive to additional pileup jets created in the forward region of the detector by the additional minimum bias interactions at each bunch crossing. A good VBF purity can nevertheless be achieved by selection on the jet p_T , invariant mass of the two jets and their separation in rapidity. This information is used in machine learning techniques as discussed in Chapter 8.6.2.

7.1.1.3 Associated production with a vector boson: VH

Another subdominant electroweak process is the production associated to a Z or W vector boson, together referred to as VH or “Higgs-strahlung”, accounting about 4 % of the total $pp \rightarrow H$ cross-section. This process corresponds at leading order to the Drell-Yan production of a vector boson, subsequently radiating off a Higgs boson (Figure 7.4a).

For experimental detection purposes, the VH production mode is typically split into “ VH -lep” including $Z(\rightarrow \ell^+ \ell^-)H$, $Z(\rightarrow \nu \bar{\nu})H$ and $W(\rightarrow \ell \nu)H$, and “ VH -had” including $Z(\rightarrow q \bar{q})H$ and $W(\rightarrow q \bar{q}')H$, where the latter is the s -channel version of the VBF diagram. The VH process is characterised by either additional leptons or missing transverse energy, or two additional jets in the final state whose invariant mass is compatible with the one of an on-shell Z or W boson. It should be noticed that two additional diagrams involving gluons in the initial state (and commonly named “ $ggZH$ ”) appear at higher order for the ZH production, as shown in Figures 7.4b and 7.4c.

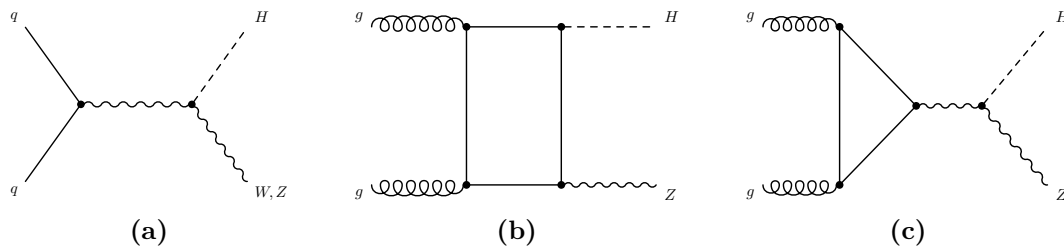


Figure 7.4 – Leading order diagram for the $qq \rightarrow VH$ process (a), and additional loop-level diagrams for the $gg \rightarrow ZH$ process (b), (c).

The VH -lep production mode is of high interest for hadronic final-state Higgs boson decays that will be further described in Section 7.1.2 (e.g. $H \rightarrow b \bar{b}$). Indeed, the leptons or missing transverse energy are clearly identified in the busy QCD environment, and the associated vector and Higgs bosons recoil back-to-back, providing a clear signature for such events.

Similarly to the VBF production mode, VH -had also suffers from contamination by the ggF process radiating additional gluons forming additional jets. The VH -had is therefore distinguished from the VBF and ggF + 2 jets by applying a requirement on the jet invariant mass, imposing it to be close to the weak bosons masses.

Furthermore, since VBF and VH -had share the same initial and final states, they might interfere. However, the above invariant mass requirement, typically around $m_{jj} < 130 \text{ GeV}$, highly constrains the phase-space and reduces interference effects by a large amount.

7.1.1.4 Associated production with a $t\bar{t}$ or $b\bar{b}$ pair: ttH/bbH

Finally, the Higgs boson production associated to a top quark pair ttH (Figure 7.5), while contributing to less than 1 % to the total cross-section, still is of prime interest. Since the top quark is much heavier than the Higgs boson, the $H \rightarrow t\bar{t}$ decay is not kinematically allowed. The ttH production is therefore the only way to directly probe the top Yukawa coupling at leading order. The other easily accessible processes containing couplings to the top quark involve loops, for example in the ggF production or $H \rightarrow \gamma\gamma$ decay under some assumptions.

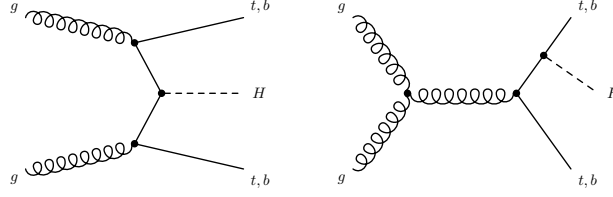


Figure 7.5 – Examples of LO diagrams for the $pp \rightarrow ttH$ process.

Since the top quark decays before hadronisation, the final state is not reconstructed as “Higgs boson + 2 top quarks”. The dominant quark branching ratio is $t \rightarrow b + W$ ($\sim 100\%$), with the W boson decaying hadronically or leptonically and the b -quark hadronising and also later decaying. This gives a large number of combinations of possible final states for the associated particles, and almost as many possible analysis categories.

A similar production mode with associated bottom quarks instead of top quarks is found with a similar cross-section, but its interest is more limited since it is experimentally difficult to separate from ggF production and the bottom quark Yukawa can be probed through the $H \rightarrow b\bar{b}$ decay.

7.1.2 Decay channels

Since the Higgs boson is an unstable particle, it will decay soon after being produced. Rather than measuring or computing its mean lifetime τ , the usual associated physics property is the *width* Γ representing the width of the resonance in the mass spectrum and linked to the mean lifetime by $\Gamma = \hbar/\tau$. For a Higgs boson with mass $m_H = 125 \text{ GeV}$, its *total width* is $\Gamma_H = 4.088 \text{ MeV}^{+0.73\%}_{-0.73\%} (\text{theory})^{+0.99\%}_{-0.98\%} (m_q)^{+0.61\%}_{-0.63\%} (\alpha_s)$ [72].

Given the structure of the interactions summarised in Section 1.1.9, the Higgs boson can decay at tree level into pairs of massive particles. To each of these *decay channels* (final state) is associated a *partial width* Γ_f such that $\sum_f \Gamma_f = \Gamma_H$, and a *branching ratio* (BR) defined as $\text{BR}_f = \Gamma_f/\Gamma_H$ representing the probability to have a Higgs boson decaying to the final state f . The branching fractions for each possible Higgs boson decay channels as well as their uncertainties are listed in Table 7.2 for $m_H = 125 \text{ GeV}$. A graphical representation of the branching ratios as a function of the Higgs boson mass is also shown in Figure 7.6.

Given that for $m_H \approx 125 \text{ GeV}$, $m_H \ll 2 \times m_t$ and $m_H < 2 \times m_Z, 2 \times m_W$, the $H \rightarrow t\bar{t}$, $H \rightarrow WW$ and $H \rightarrow ZZ$ are kinematically suppressed and the largest branching ratio is held by the $H \rightarrow b\bar{b}$ decay, contributing to $\sim 58\%$ of the total width. However, despite the largest branching fraction, this channel suffers from high QCD backgrounds and difficulties to firmly identify a b -jet. As a consequence, this channel is mainly studied in the VH -lep production mode

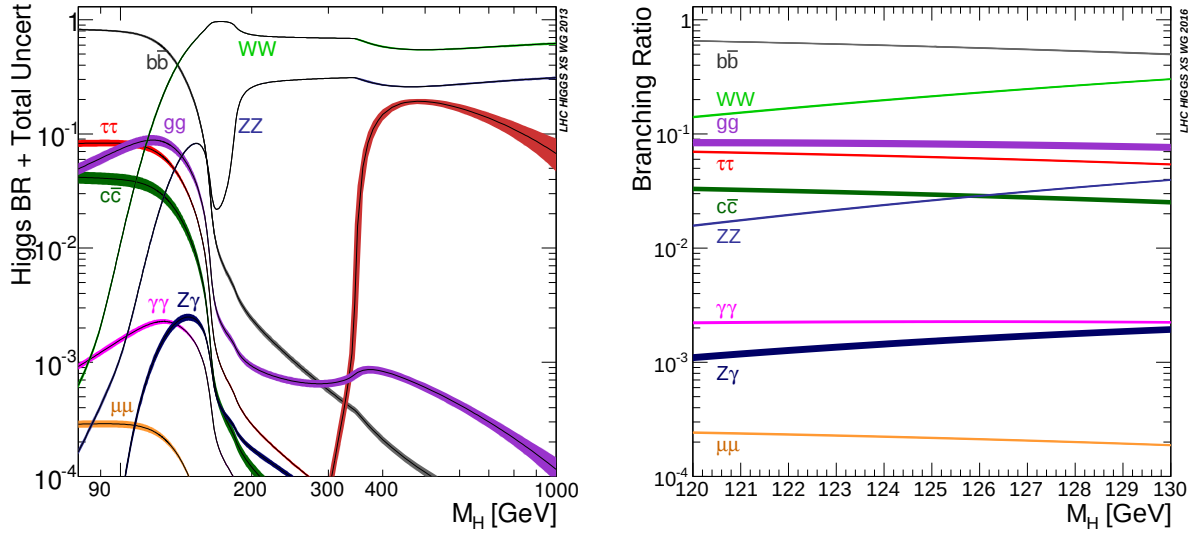


Figure 7.6 – SM Higgs boson branching ratios as a function of its mass (the right plot is a zoom around the discovered Higgs boson mass). [34, 72]

Table 7.2 – Branching ratio and their uncertainties for the Higgs boson of mass $m_H = 125 \text{ GeV}$. The theory uncertainties represent the systematic coming from varying the renormalisation and factorisation scales. m_q denotes the uncertainties due to the quark masses experimental measurements, and α_s the impact of varying the choice of strong coupling constant (precision of a few percents). The uncertainties on the left represent uncertainties on the branching fraction. However, the primary quantity computed are the partial widths, which are used to deduce the branching ratios. Since the total width is dominated by the $H \rightarrow b\bar{b}$ contribution, its uncertainties are largely propagated to other branching ratios. The perturbative order in QCD and EW used for the theoretical computation of the partial widths and their estimated error due to missing higher-order terms are shown on the right. [72, 170, 173]

	BR [%]	Relative uncertainty [%]			Accuracy on partial width			
		Theory	m_q	α_s	QCD		EW	
$H \rightarrow b\bar{b}$	58.24	+0.65 % -0.65 %	+0.72 % -0.74 %	+0.78 % -0.80 %	N ⁴ LO	< 0.1 %	NLO	1-2 %
$H \rightarrow WW$	21.37	+0.99 % -0.99 %	+0.99 % -0.98 %	+0.66 % -0.63 %	NLO	< 0.5 %	NLO	< 0.5 %
$H \rightarrow gg$	8.187	+3.40 % -3.41 %	+1.12 % -1.13 %	+3.69 % -3.61 %	N ³ LO	3 %	NLO	1 %
$H \rightarrow \tau^+\tau^-$	6.272	+1.17 % -1.16 %	+0.98 % -0.99 %	+0.62 % -0.62 %			NLO	1-2 %
$H \rightarrow c\bar{c}$	2.891	+1.20 % -1.20 %	+5.26 % -0.98 %	+1.25 % -1.25 %	N ⁴ LO	< 0.1 %	NLO	1-2 %
$H \rightarrow ZZ$	2.619	+0.99 % -0.99 %	+0.99 % -0.98 %	+0.66 % -0.63 %	NLO	< 0.5 %	NLO	< 0.5 %
$H \rightarrow \gamma\gamma$	0.227	+1.73 % -1.72 %	+0.93 % -0.99 %	+0.61 % -0.62 %	NLO	< 0.1 %	NLO	< 0.1 %
$H \rightarrow Z\gamma$	0.153	+5.71 % -5.71 %	+0.98 % -1.01 %	+0.58 % -0.65 %	NLO	< 1 %	LO	5 %
$H \rightarrow \mu^+\mu^-$	0.0218	+1.23 % -1.23 %	+0.97 % -0.99 %	+0.59 % -0.64 %				

to reduce the background and only passed the 5σ threshold using up to 2017 data [174, 175]. This channel provides unique sensitivity to the bottom quark Yukawa, but due to its fully hadronic final state it has a poor resolution and cannot contribute to a Higgs boson mass measurement.

The second contribution nevertheless comes from the $H \rightarrow WW$ channel ($\sim 21\%$). Since $m_H < 2 \times m_W$, one or the two W bosons must be off-shell. In practice, one is mostly on-shell and the other off-shell. Since the W boson is also unstable, they will subsequently decay as well, either hadronically giving a quark–antiquark pair ($\sim 70\%$) or leptonically giving a lepton and its related neutrino. The fully hadronic ($qqqq$) final state has the largest branching ratio, but as for the $H \rightarrow b\bar{b}$ case suffers from high multi-jets backgrounds. The semi-leptonic channel $\ell\nu qq$ also suffers from QCD backgrounds, so the most promising for couplings studies is the fully leptonic channel. Using the opposite flavour $H \rightarrow W(\rightarrow e^-\bar{\nu})W(\rightarrow \mu^+\nu)$ channel (or its charge conjugate) provides a clean signature, reducing by a large amount the same flavour backgrounds from $Z \rightarrow e^+e^-$ or $Z \rightarrow \mu^+\mu^-$ decays in the $H \rightarrow W(\rightarrow e^-\bar{\nu})W(\rightarrow e^+\nu)$ or $H \rightarrow W(\rightarrow \mu^-\bar{\nu})W(\rightarrow \mu^+\nu)$. However due to the presence of unmeasurable neutrinos in the final state, the resolution is insufficient to provide a measurement of the Higgs boson mass lineshape for example. An example is given in Figure 7.7a.

The $H \rightarrow gg$ and $H \rightarrow c\bar{c}$ channels are experimentally similar to the $H \rightarrow b\bar{b}$ one, with large QCD backgrounds but with lower statistics due to lower branching ratio, and much lower capabilities of gluon- or c -tagging. It should also be noted that the gluon being massless, the $H \rightarrow gg$ does not occur at tree level, but through loops (the leading order diagram is the reverse ggF one).

With a branching ratio around 6%, $H \rightarrow \tau^+\tau^-$ still constitutes a sizeable contribution to the Higgs boson width. Since the τ lepton decays quickly after being produced, emitting a neutrino and a W boson, its final state will contain at least two neutrinos yielding a poor resolution and a high contamination from the nearby $Z \rightarrow \tau\tau$ decay as seen in Figure 7.7b. This decay channel was nevertheless confirmed during Run 2 by CMS [180] and ATLAS [177], and can provide performing measurements of spin and CP -mixing effects thanks to its many-particle final state.

Since the Higgs boson is further away from the double Z boson threshold than the double W , its branching ratio is about one order of magnitude below the one of $H \rightarrow WW$. Similarly to $H \rightarrow WW$, the $H \rightarrow ZZ$ decay channel offers a variety of final states due to the unstable Z bosons decaying either to a quark–antiquark pair ($\sim 70\%$), a neutrino pair ($\sim 20\%$) or a charged-lepton pair ($\sim 3\%$ each family). The hadronic final states $q\bar{q}q\bar{q}$, $q\bar{q}\nu\bar{\nu}$ and $\ell^+\ell^-q\bar{q}$ suffer from the same QCD backgrounds as previously mentioned and are not much used for the low mass Higgs boson studies. The $\ell^+\ell^-\nu\bar{\nu}$ is mostly used for high mass Higgs boson searches due to the presence of neutrinos. Finally, the 4ℓ final state, limited to $\ell = e, \mu$ to avoid issues mentioned in the $H \rightarrow \tau^+\tau^-$ decay, constitutes one of the two “precision channels”: its final state is completely reconstructed and only composed of light leptons, which are well measured in both ATLAS and CMS. Despite a tiny branching ratio of 1.240×10^{-4} [72]¹, this channel comes with a largely reduced background to have a signal-to-background ratio (S/B) better than 2 in the mass window around 125 GeV. An illustration is given in Figure 7.7c. Furthermore and similarly to $H \rightarrow \tau^+\tau^-$, the complex final state with four particles allows to study spin and CP effects through various angles between the leptons.

The second precision channel is the $H \rightarrow \gamma\gamma$ decay: despite a low branching ratio of about 0.2%, its final state is made of a clean, high energy, isolated photon pair, which is easy to tag and well measured in ATLAS and CMS. This channel has a large di-photon combinatorial background as seen in Figure 7.7d, but provides a good resolution and dominates the precision of many cross-

¹↑ Taking the PDG values [41] for the branching ratios of a Z boson decaying either to electron or muon pairs: $\text{BR}(Z \rightarrow ee) = 3.3632$ and $\text{BR}(Z \rightarrow \mu\mu) = 3.3662$, we obtain $\text{BR}(H \rightarrow 4\ell) = 0.02619 \times (0.033632 + 0.033662)^2 = 1.186 \times 10^{-4}$. The actual value (1.240×10^{-4} [72]) is a bit higher than this naive computation due to interference effects in the 4μ and $4e$ final states.

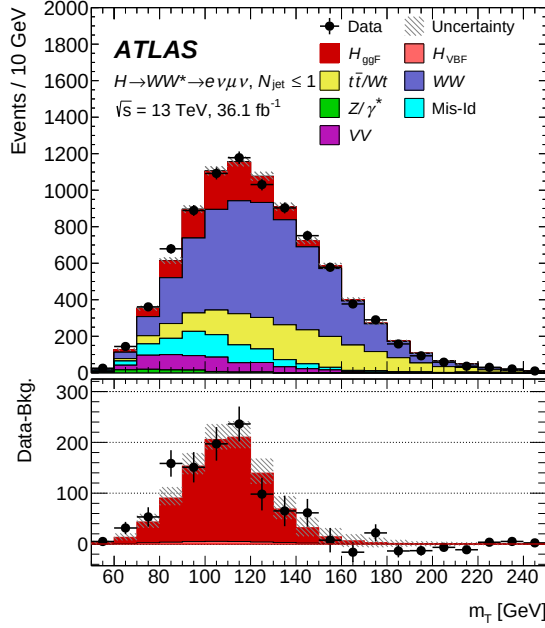
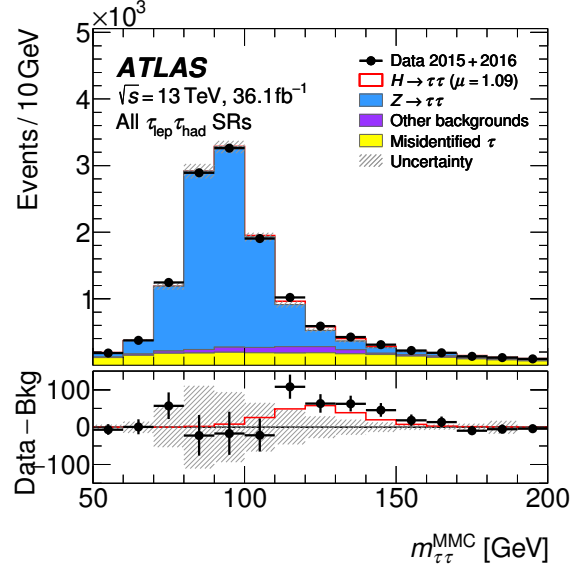
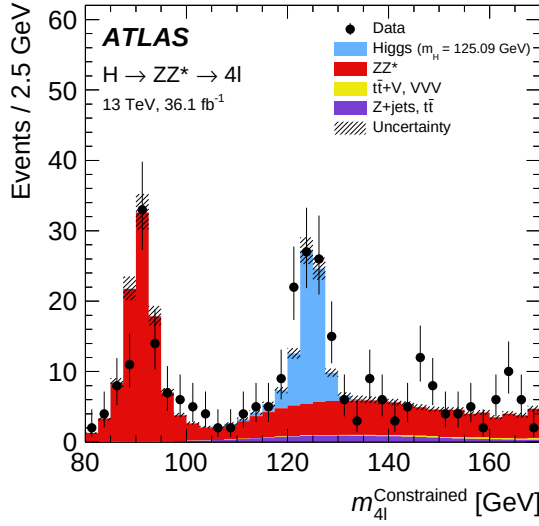
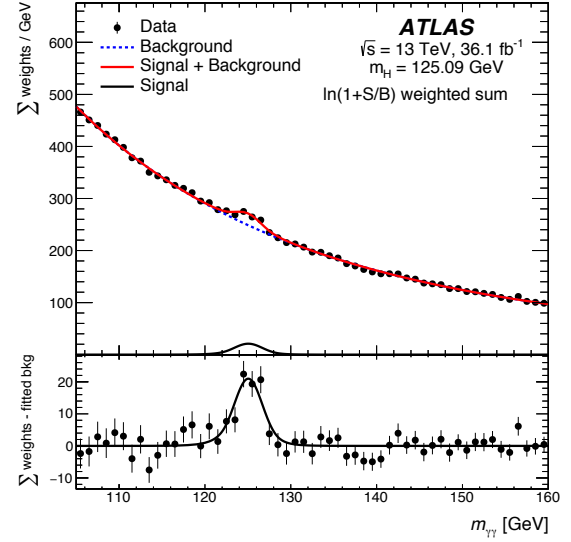

 (a) $H \rightarrow WW^{(*)} \rightarrow e\nu\mu\nu$

 (b) $H \rightarrow \tau^+\tau^-$

 (c) $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$

 (d) $H \rightarrow \gamma\gamma$

Figure 7.7 – Mass distribution for the (a) $H \rightarrow WW$ analysis [176], (b) $H \rightarrow \tau^+\tau^-$ analysis [177], (c) $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ analysis [178], (d) $H \rightarrow \gamma\gamma$ analysis [179], using 36.1 fb^{-1} of data collected at $\sqrt{s} = 13\text{ TeV}$ at the ATLAS experiment.

sections and couplings analyses as long as the $H \rightarrow 4\ell$ remains in the low statistics regime. Since the photon, as for the gluon, is massless, it does not directly couple to the Higgs boson. Therefore, the $H \rightarrow \gamma\gamma$ decays occurs through loops, as presented in Figure 7.8, explaining its low branching ratio. The combined mass measurement from ATLAS and CMS using the $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ and $H \rightarrow \gamma\gamma$ channels reached a precision of a few permil with the Run 1 data [55]:

$$125.09 \pm 0.21 \text{ (stat)} \pm 0.11 \text{ (syst)} \text{ GeV.} \quad (7.1)$$

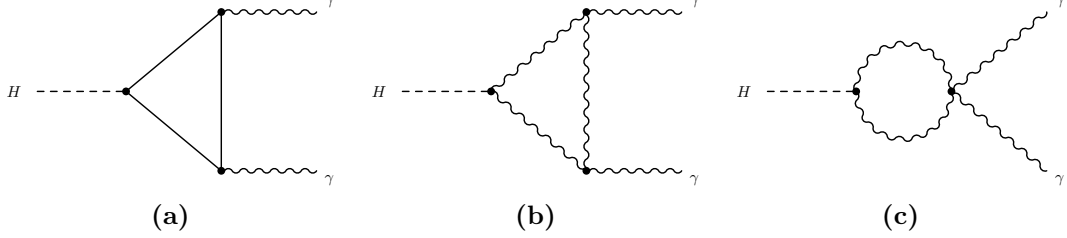


Figure 7.8 – Examples of diagrams contributing to the $H \rightarrow \gamma\gamma$ decay at leading order. Since the photon is massless and does not directly couple to the Higgs boson, its interaction must occur through fermionic loops (a), dominated by the top quark since it is the heaviest and has largest coupling to the Higgs boson, or through W boson loops involving two triple gauge couplings (b) or one quartic gauge coupling (c). The Z boson cannot enter the loop since the $ZZ\gamma$ vertex is forbidden in the Standard Model.

Experimentally, the $H \rightarrow \mu^+\mu^-$ analysis is similar to $H \rightarrow \gamma\gamma$, with a large continuous combinatorial background but a branching fraction one order of magnitude lower due to its low mass. Consequently, this channel still lacks statistics at the end of Run 2 [181].

7.1.3 Spin and parity

One fundamental property of the predicted Higgs boson is its spin-parity. The Standard Model predicts that the Higgs boson should be a scalar (*i.e.* having a spin $J = 0$) of positive parity (*i.e.* $P = +1$), so assessing the spin-parity of the discovered Higgs boson is therefore paramount.

The observation of the Higgs boson decay into two photons [57, 62] rules out the pure spin-1 state by virtue of the Landau-Yang [182, 183] theorem, yet under two assumptions. The first requires that the resonance seen in the $H \rightarrow \gamma\gamma$ channel is the same as the one seen by other channels. The second requires that the $H \rightarrow \gamma\gamma$ decay does not occur through light intermediate states: if the Higgs boson first decay to a yet unknown light particle a , which can thereafter decay into two photons, then the decay would be $H \rightarrow aa \rightarrow 4\gamma$. Under the assumption of a light a , its two decay photons could be sufficiently boosted to be misidentified as a single one, mimicking a genuine $H \rightarrow \gamma\gamma$ decay.

The pure spin-2 hypothesis is another interesting possibility since such resonances are predicted by many beyond the Standard Model theories. This could originate from gravitons, but poses multiple experimental and theoretical problems (see Section 11 of Ref. [34]).

The Run 1 data of the LHC already pushed the constraints enough to rule out at better than the 99% CL every pure spin-parity state other than 0^+ , as seen previously in Figure 1.17 and demonstrated in Refs. [66–68] and [61, 62, 69, 70]. However, while the parity admixture already started to be investigated, large parts of the phase-space are only loosely constrained. The analysis presented in Chapter 9 focuses on such spin-0 CP admixture.

An efficient probe of CP effects lies in the angular distributions of the decay products of a particle, as first noticed in the π^0 parity measurements from the $\pi^0 \rightarrow e^+e^-e^+e^-$ decay [184]. The many-body final states are hence of prime interest also for the Higgs boson, so natural candidate decays are $H \rightarrow WW^{(*)} \rightarrow e\nu\mu\nu$, $H \rightarrow \tau^+\tau^-$ (with the τ lepton subsequently decaying) and of

course $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ with four well-reconstructed leptons ($\ell = e, \mu$) in the final state. On top of the decay channels, CP -sensitive observables are also found in the angular distributions of jets in Higgs boson production modes associated to jets, in particular in the VBF production as noted in Refs. [185–187].

The difficulty of probing CP -odd states in the bosonic channels comes from their couplings to the 0^- particle not occurring at tree level. Since the CP -odd-state couplings to the vector bosons (denoted AVV) only happens through loops, their amplitude is suppressed compared to the CP -even-state couplings to the vector bosons (denoted HVV). As a consequence, in case of a CP admixture, the CP -sensitive observables still look like SM hypothesis even for sizeable values of the CP -odd couplings. Yet, with an increased coupling value, the event yield also increases but this effect is indistinguishable from other sources such as the presence of non-SM CP -even couplings.

In the fermion sector however, both the CP -odd couplings Aff and the CP -even couplings Hff can occur at tree level, offering a fair sensitivity in both cases. However as seen above, the decay to fermions (mainly $H \rightarrow b\bar{b}$, $H \rightarrow \tau^+\tau^-$) suffer from large backgrounds, making the analysis more difficult: the $H \rightarrow \tau^+\tau^-$ decay channel already provided a CP analysis [71, 188], but the $t\bar{t}H$ production and the decay to bottom quarks is barely established (Refs. [189, 190] and [174, 175] respectively), and the decay to second generation fermions is not yet observed.

7.2 Beyond the Standard Model

Since we suspect that the Standard Model cannot be the ultimate theory of particles physics (see Section 1.2.3), the current focus of analyses is to search for deviations from the theory, unpredicted effects that could be a hint toward new physics (NP). This sector not described by the Standard Model is called *Beyond the Standard Model*, commonly abridged BSM. In order to systematically investigate possible BSM effects, several frameworks have been introduced. Given the sensitivity expected with the Run 1 data, a simple interpretation framework was put in place: deviations were interpreted in terms of Higgs boson couplings modifiers in the so-called κ -framework (Section 7.2.1).

Thanks to a successful Run 2 data taking, the statistical precision of analyses improves by a sizeable factor, and more complete frameworks have been elaborated with various degrees of model dependency. Two types of frameworks must be distinguished: those providing phase space regions for the measurements, and those providing a global framework to interpret these measurements. The first ones mainly include the fiducial cross-sections, essentially targeting the Higgs boson decay phase-space, and the Simplified Template Cross-Sections (STXS), essentially targeting the Higgs boson production phase-space. To have an handle on various couplings, the coupling analyses rely on separating the various Higgs boson production modes and therefore mostly use the STXS framework. The interpretation frameworks mainly include the Pseudo-Observables (PO) and the Effective Field Theory (EFT) approaches.

The legacy κ -framework results from Run 1 are recalled in Section 7.2.1 and an overview of the STXS framework is presented in Section 7.2.2. Since the EFT approach is chosen to interpret the results of the CP analysis of Chapter 9, it is more extensively described in Section 7.3. Additional information and complete description of other measurement and interpretation frameworks can be found in Ref. [72].

7.2.1 The κ -framework

The first interpretation framework put in place was the so-called κ -framework (see Section 10 of Ref. [34]). In this framework, the tensor structure of the Higgs boson couplings are kept as in the SM case, so the Higgs is assumed to be a $J^P = 0^+$ scalar boson. The couplings modifiers κ are introduced as scale factors of the cross-sections or partial widths such that

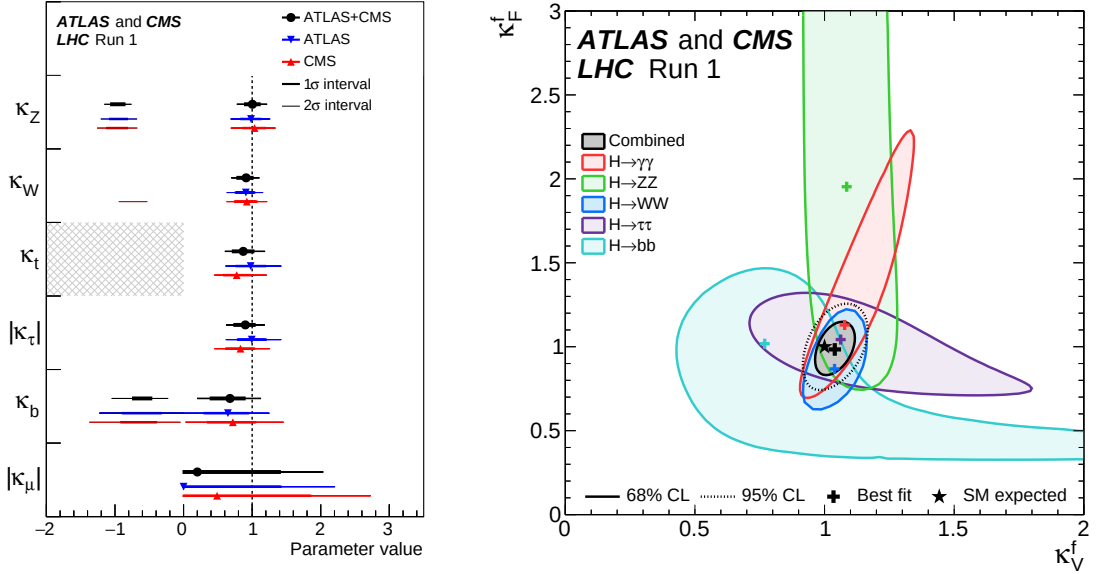
$$\sigma_i \times \text{BR}^f = \frac{\sigma_i(\vec{\kappa})\Gamma^f(\vec{\kappa})}{\Gamma_H}. \quad (7.2)$$

For example, one introduces $\kappa_W^2 = \Gamma_{WW}/\Gamma_{WW}^{\text{SM}}$ and $\kappa_Z^2 = \Gamma_{ZZ}/\Gamma_{ZZ}^{\text{SM}}$ which enter the $H \rightarrow WW$ and $H \rightarrow ZZ$ decay, and similarly for the $H \rightarrow \gamma\gamma$ and fermionic decays. But such couplings are also present in the production side in VBF and VH , so one also has $\kappa_W^2 = \sigma_{WH}/\sigma_{WH}^{\text{SM}}$ and $\kappa_Z^2 = \sigma_{ZH}/\sigma_{ZH}^{\text{SM}}$. Then, κ_{VBF} can be built as

$$\kappa_{\text{VBF}} = \frac{\kappa_W^2 \cdot \sigma_{\text{WBF}} + \kappa_Z^2 \cdot \sigma_{\text{ZBF}}}{\sigma_{\text{WBF}} + \sigma_{\text{ZBF}}}. \quad (7.3)$$

The Standard Model is recovered when all κ are 1.

This framework has widely been used during Run 1, but since the available dataset was limited, some simplifications had been made. Due to limited sensitivity when doing a simultaneous multi-parameter scan, experiments merged all the vector-boson-related scales factors into one single κ_V and all fermion-related scale factors in one single κ_F . The ATLAS and CMS combined result of this analysis is presented in Figure 7.9.



(a) Best fit value for the κ assuming $\text{BR}(\text{BSM}) = 0$.

(b) 68 % CL exclusion regions in the (κ_V, κ_F) plane, assuming all κ are positive.

Figure 7.9 – κ -framework result from combined ATLAS and CMS measurements using the Run 1 data of the LHC. [65]

Despite providing a first unified framework for theoreticians and experimentalists, this framework quickly suffers from limitations: on top of the above simplification due to limited experimental sensitivity, the scale factors only embrace cross-section modifications but cannot account for kinematics changes, for example the distribution of the Higgs boson transverse momentum. Moreover, the κ -framework only includes leading order EW corrections which is accurate at the 10 % level, while experimental results are expected to reach the percent level at the end of Run 2.

As a consequence, more generic frameworks have been developed in preparation for the LHC Run 2.

7.2.2 Simplified Template Cross-Sections (STXS)

With the larger dataset acquired during Run 2, more precise measurements have become possible requiring an extended binning of the phase-space for measurements. In order to provide both experiments and theoreticians with a unified definition of such phase-space region, the Higgs boson STXS was developed between the two communities [72].

The main requirements for this framework are to maximise the experimental sensitivity, in particular to the search of BSM effects by defining BSM-sensitive bins, and to minimise the theoretical uncertainties. The former is ensured by providing a phase-space splitting that can be easily experimentally matched, *i.e.* analyses should be able to define one or two bins that correspond to a large extent to one of the targeted phase-space regions. The latter is achieved by splitting the Higgs boson production modes into individual categories, and then to further split as a function of the number of jets or of the Higgs boson momentum, which carry most of the theoretical dependency. The bin boundary definitions also care for a uniform signal acceptance within a single bin. The signal acceptance inside each bins was also a determinant criterion to decide for the bin boundaries.

These requirements however compete with the finite amount of data available and therefore the limited sensitivity if the number of phase-space regions is too high compared to the experimental statistical accuracy. The framework was therefore additionally designed in an evolutive manner, called *stages*: the Stage-0 is mainly built for the Run 1 studies and only splits in the production modes (ggF, VBF, VH , ttH) without further splitting. The Stage-1 is built over the Stage-0 and targets the full Run 2 analyses, further splitting each production mode with respect to the number of jets, or Higgs boson, or jet momentum. The current recommended splitting is called *Stage-1.1* [191] (see Figure 7.10) and is a redefinition of the Stage-1 to better match analyses categories. It also introduced a possible sub-splitting of each category to better account for theory uncertainties if needed, these sub-bins being remerged during the result interpretation phase.

The ttH category is not currently split further due limited sensitivity, even with the full Run 2 dataset. It should be noticed that the bbH and tH contributions are not explicitly mentioned: indeed, these two modes cannot be probed with the current dataset. Due to its similar acceptance with the ggF mode, the small bbH contribution is included in the ggF categories. The tH contribution is similarly merged within the ttH bin. In both cases, their relative contributions in these bins is fixed to the Standard Model prediction.

Since the binning targets a combination of all decay channels at the end of Run 2, single analyses might not have high enough statistics to populate each bin. It is therefore given the possibility for analyses to merge related bins in phase-space to match the binning to their sensitivity. As an example, the $H \rightarrow 4\ell$ analysis uses a “reduced Stage-1.1”, described in Chapter 8.6.1.

7.3 Effective Field Theories: EFT

7.3.1 Energy power counting

In natural units ($\hbar = c = 1$), every quantity can be expressed in terms of energy dimension (called *canonical dimension*). The dimension is denoted $[\cdot]$, and we therefore have $[\hbar] = [c] = 0$, $[\text{energy}(E)] = 1$, and the usual $[\text{mass}] = [E/\epsilon^2] = 1$, $[\text{momentum}(p)] = [E/\epsilon] = 1$. One can also

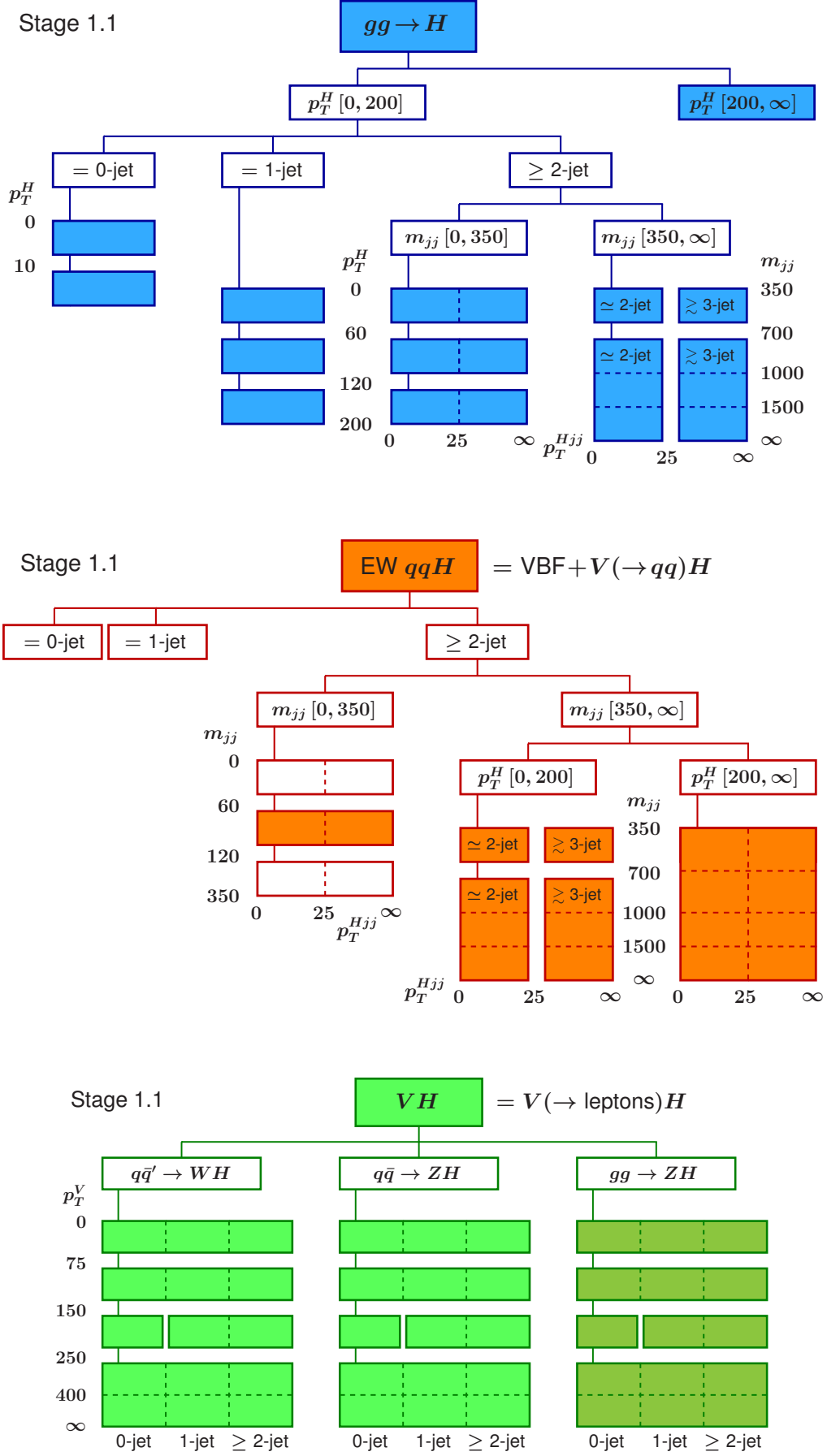


Figure 7.10 – Definition of truth-level STXS bins in the Stage-1.1. [191]

show that $[x_\mu] = [\text{length}] = [\text{time}] = -1$ and so $[\partial_\mu] = [\partial/\partial x_\mu] = 1$. From this and recalling that $[\text{action}] = [\hbar] = 0 = \int \mathcal{L} d^4x$, one deduces that $[\mathcal{L}] = 4$. The Lagrangian is said to be of canonical dimension 4.

As for the fields, one can show that the dimension of scalar and vector fields is $[\phi] = [A_\mu] = 1$, and that the canonical dimension of fermion fields is $[\psi] = 3/2$. It can be checked that $[D] = 1$ and $[F_{\mu\nu}] = 2$, and that the coupling constants are dimensionless. Finally, it can be shown (see for example Ref. [1]) that if the Lagrangian includes vertices with canonical dimension greater than 4, or equivalently if the coupling constant associated to that interaction has a negative energy dimension, then the theory is not renormalisable.

7.3.2 A simple example: the Fermi theory of β decay

An early example of effective theory is the Fermi theory of nuclear disintegration, postulated by Fermi in 1933 [192–194]. In this explanation of β decay, the disintegration occurs as a *point-like* interaction, as shown in Figure 7.11a: an incoming neutron splits into a proton, an electron and a neutrino. The *effective* Lagrangian describing the interaction is $\mathcal{L}^{\text{Fermi}} \supset G_F (\bar{d}\gamma_\rho u)(\bar{\nu}_e \gamma^\rho \frac{1-\gamma^5}{2} e)$. Since four fermions interact together, this term is of canonical dimension 6 and contains a constant $G_F \sim g^2/m_W^2$ called the Fermi constant, which is of dimension -2 . At the time, this was a sufficient modelling of the process, which could moreover be successfully applied to the muon decay. However, only the Fermi constant was measurable, not g nor m_W .

However, later attempts to compute the cross-section of this process spotted a divergence as the momentum transfer, *e.g.* between the muon and the electron, increased. This was later resolved by the electroweak theory [6–8] and the introduction of a new, heavy mediator regularising the cross-section: the W boson, discovered a few years later [78, 79]. Given the typical energy at play in the β or μ decays (a few MeV), this heavy and unknown degree of freedom (~ 80 GeV) has been integrated out in the Fermi theory. The *effective* theory approximation is valid as long as the typical energy of the process is not enough to “resolve” the heavy state, as illustrated in Figure 7.11b.

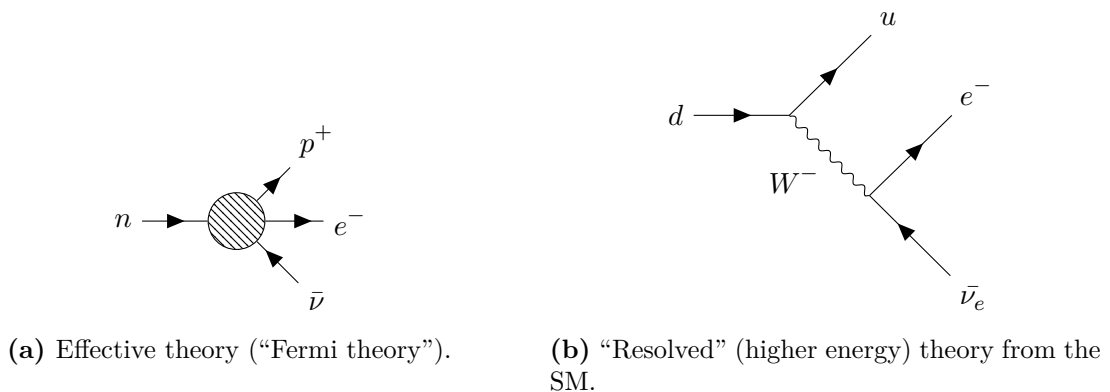


Figure 7.11 – Effective Fermi theory (a) and resolved (high energy) theory (b) of the β decay. In 1933, the quarks were not yet postulated so the interaction was thought to behave like a neutron n transforming into a proton p^+ and emitting an electron e^- and a (anti-)neutrino $\bar{\nu}$. In the Standard Model, it is rather a down quark d transforming into a up quark u through the emission of a W^- boson, subsequently decaying into an electron and an antineutrino.

7.3.3 Standard Model EFT and the Warsaw basis

The Standard Model Effective Field Theory (EFT) consists in a systematic expansion of the SM Lagrangian, introducing higher-dimension operators fulfilling Lorentz and gauge invariance. This defines the most generic Lagrangian in each dimension, the operators in each successive dimension being suppressed regarding the previous one by an high energy scale Λ . This scale Λ represents the typical mass of new particles and defines the domain where the EFT can be considered valid: the process probed through this EFT should happen at a relatively lower energy than the scale set by Λ . Such an EFT assumes that no new light particles enter the loops and that all the new physics occurs at energies above Λ .

Since the Standard Model Lagrangian $\mathcal{L}_{\text{SM}}$ operators are of dimension 4, the expansion starts at 5, suppressed by a factor Λ^1 , and continues as

$$\mathcal{L}_{\text{EFT}} = \mathcal{L}_{\text{SM}} + \sum_i \frac{c_i^{d=5}}{\Lambda} \mathcal{O}_i^{d=5} + \sum_i \frac{c_i^{d=6}}{\Lambda^2} \mathcal{O}_i^{d=6} + \sum_i \frac{c_i^{d=7}}{\Lambda^3} \mathcal{O}_i^{d=7} + \sum_i \frac{c_i^{d=8}}{\Lambda^4} \mathcal{O}_i^{d=8} + \dots \quad (7.4)$$

The c_i factors are the so-called *Wilson coefficients*, which are interpreted as new *coupling constants* for the $\mathcal{O}$ operators that introduce new *effective* interactions in the Lagrangian. A common convention is to redefine

$$\frac{c_i^d}{\Lambda^{d-4}} \rightarrow c_i^d \cdot \left(\frac{v}{\Lambda}\right)^{d-4} \quad \text{and} \quad \mathcal{O}_i^d \rightarrow \frac{\mathcal{O}_i^d}{v^{d-4}}. \quad (7.5)$$

For each given dimension, a complete and non-redundant set of operators must be computed for the expansion to be correct. Such work has been achieved for dimension-5 [195], dimension-6 [196], dimension-7 [197, 198] and dimension-8 [198, 199]. In higher dimension, the number of operators can be computed [198] but the exact form of each is not yet known. Hopefully, since each subsequent dimension is suppressed (by a factor Λ) compared to the previous one, the expansion can be truncated, bearing in mind that the more dimension included the more precise the result.

Moreover, it was shown in Ref. [195] that all dimension-5 operators violate the lepton number L and that all odd higher-dimension operators violate the $B - L$ number (baryon minus lepton number) in Ref. [200]. Since no such phenomenon is observed, all odd-dimension operators are usually discarded from the expansion and only B and L conserving terms are kept in even-dimension. Given that dimension-8 operators are suppressed by an additional Λ^2 factor compared to dimension-6 terms, the expansion is often limited to $\mathcal{L}_{\text{EFT}} = \mathcal{L}_{\text{SM}} + \sum_i c_i^{d=6} \mathcal{O}_i^{d=6}$, with dimension-8 operators serving as second-order corrections on dimension-6 when needed.

Focusing on the dimension-6, 59 B -conserving additional tensor structures (76 if also counting the hermitian conjugate of relevant operators) are added to the SM Lagrangian. Assuming three independent flavours, this results in 2499 independent real parameters, reduced to 76 if assuming complete flavour symmetry (*i.e.* only one flavour). This large number is far beyond reach of experimental sensitivity, so all the operators cannot be probed at the same time. Depending on the analysis, a fraction of them are left floating while all the others are fixed to their Standard Model values.

Once a complete and non-redundant set of operators has been derived, one is free to chose any other complete and non-redundant set resulting from a linear combination of these operators. Each of these sets forms an *EFT basis*, which are all equivalent since the Wilson coefficients can be translated between bases. Consequently, the choice of basis will rely on the convenience for a particular study. The baryon number conserving dimension-6 operators are listed in the historical *Warsaw basis* in Table 7.3. The operators include both CP -conserving and CP -violating terms, the latter being tagged with a tilde. Such operators include a tilded field defined by $\tilde{X}_{\mu\nu} = \frac{1}{2}\epsilon_{\mu\nu\rho\sigma}X^{\rho\sigma}$ where $\epsilon_{\mu\nu\rho\sigma}$ is the fully antisymmetric tensor in four dimensions with $\epsilon_{1234} = 1$.

Table 7.3 – The 76 dimension-6 baryon number conserving operators (noted Q) in the so-called Warsaw basis [196] as written in [201]. X stands for any of the field-strength tensor $B_{\mu\nu}, W_{\mu\nu}^I, G_{\mu\nu}^A$, D for a covariant derivative term, ψ for the fermion fields, and L and R for the left doublets and right singlets (l, q and e, u, d). The flavour labels of the form p, r, s, t on the Q operators are suppressed on the left hand side of the tables. τ^I and T^A are the generators of $SU(2)$ and $SU(3)$ (including a multiplicative factor of $1/2$). γ_μ are the usual Dirac matrices (see Chapter 1.1) and $\sigma_{\mu\nu} = i[\gamma_\mu, \gamma_\nu]/2$.

1 : X^3		2 : H^6		3 : $H^4 D^2$		5 : $\psi^2 H^3 + \text{h.c.}$	
Q_G	$f^{ABC} G_\mu^{A\nu} G_\nu^{B\rho} G_\rho^{C\mu}$	Q_H	$(H^\dagger H)^3$	$Q_{H\Box}$	$(H^\dagger H)\Box(H^\dagger H)$	Q_{eH}	$(H^\dagger H)(\bar{l}_p e_r H)$
$Q_{\tilde{G}}$	$f^{ABC} \tilde{G}_\mu^{A\nu} G_\nu^{B\rho} G_\rho^{C\mu}$			Q_{HD}	$(H^\dagger D_\mu H)^* (H^\dagger D_\mu H)$	Q_{uH}	$(H^\dagger H)(\bar{q}_p u_r \tilde{H})$
Q_W	$\epsilon^{IJK} W_\mu^{I\nu} W_\nu^{J\rho} W_\rho^{K\mu}$					Q_{dH}	$(H^\dagger H)(\bar{q}_p d_r H)$
$Q_{\tilde{W}}$	$\epsilon^{IJK} \tilde{W}_\mu^{I\nu} W_\nu^{J\rho} W_\rho^{K\mu}$						
4 : $X^2 H^2$		6 : $\psi^2 XH + \text{h.c.}$		7 : $\psi^2 H^2 D$			
Q_{HG}	$H^\dagger H G_{\mu\nu}^A G^{A\mu\nu}$	Q_{eW}	$(\bar{l}_p \sigma^{\mu\nu} e_r) \tau^I H W_{\mu\nu}^I$	$Q_{Hl}^{(1)}$	$(H^\dagger i \overleftrightarrow{D}_\mu H)(\bar{l}_p \gamma^\mu l_r)$		
$Q_{H\tilde{G}}$	$H^\dagger H \tilde{G}_{\mu\nu}^A G^{A\mu\nu}$	Q_{eB}	$(\bar{l}_p \sigma^{\mu\nu} e_r) H B_{\mu\nu}$	$Q_{Hl}^{(3)}$	$(H^\dagger i \overleftrightarrow{D}_\mu^I H)(\bar{l}_p \tau^I \gamma^\mu l_r)$		
Q_{HW}	$H^\dagger H W_{\mu\nu}^I W^{I\mu\nu}$	Q_{uG}	$(\bar{q}_p \sigma^{\mu\nu} T^A u_r) \tilde{H} G_{\mu\nu}^A$	Q_{He}	$(H^\dagger i \overleftrightarrow{D}_\mu H)(\bar{e}_p \gamma^\mu e_r)$		
$Q_{H\tilde{W}}$	$H^\dagger H \tilde{W}_{\mu\nu}^I W^{I\mu\nu}$	Q_{uW}	$(\bar{q}_p \sigma^{\mu\nu} u_r) \tau^I \tilde{H} W_{\mu\nu}^I$	$Q_{Hq}^{(1)}$	$(H^\dagger i \overleftrightarrow{D}_\mu H)(\bar{q}_p \gamma^\mu q_r)$		
Q_{HB}	$H^\dagger H B_{\mu\nu} B^{\mu\nu}$	Q_{uB}	$(\bar{q}_p \sigma^{\mu\nu} u_r) \tilde{H} B_{\mu\nu}$	$Q_{Hq}^{(3)}$	$(H^\dagger i \overleftrightarrow{D}_\mu^I H)(\bar{q}_p \tau^I \gamma^\mu q_r)$		
$Q_{H\tilde{B}}$	$H^\dagger H \tilde{B}_{\mu\nu} B^{\mu\nu}$	Q_{dG}	$(\bar{q}_p \sigma^{\mu\nu} T^A d_r) H G_{\mu\nu}^A$	Q_{Hu}	$(H^\dagger i \overleftrightarrow{D}_\mu H)(\bar{u}_p \gamma^\mu u_r)$		
Q_{HWB}	$H^\dagger \tau^I H W_{\mu\nu}^I B^{\mu\nu}$	Q_{dW}	$(\bar{q}_p \sigma^{\mu\nu} d_r) \tau^I H W_{\mu\nu}^I$	Q_{Hd}	$(H^\dagger i \overleftrightarrow{D}_\mu H)(\bar{d}_p \gamma^\mu d_r)$		
$Q_{H\tilde{W}B}$	$H^\dagger \tau^I H \tilde{W}_{\mu\nu}^I B^{\mu\nu}$	Q_{dB}	$(\bar{q}_p \sigma^{\mu\nu} d_r) H B_{\mu\nu}$	$Q_{Hud} + \text{h.c.}$	$i(\tilde{H}^\dagger D_\mu H)(\bar{u}_p \gamma^\mu d_r)$		
8 : $(\bar{L}L)(\bar{L}L)$		8 : $(\bar{R}R)(\bar{R}R)$		8 : $(\bar{L}L)(\bar{R}R)$			
Q_{ll}	$(\bar{l}_p \gamma_\mu l_r)(\bar{l}_s \gamma^\mu l_t)$	Q_{ee}	$(\bar{e}_p \gamma_\mu e_r)(\bar{e}_s \gamma^\mu e_t)$	Q_{le}	$(\bar{l}_p \gamma_\mu l_r)(\bar{e}_s \gamma^\mu e_t)$		
$Q_{qq}^{(1)}$	$(\bar{q}_p \gamma_\mu q_r)(\bar{q}_s \gamma^\mu q_t)$	Q_{uu}	$(\bar{u}_p \gamma_\mu u_r)(\bar{u}_s \gamma^\mu u_t)$	Q_{lu}	$(\bar{l}_p \gamma_\mu l_r)(\bar{u}_s \gamma^\mu u_t)$		
$Q_{qq}^{(3)}$	$(\bar{q}_p \gamma_\mu \tau^I q_r)(\bar{q}_s \gamma^\mu \tau^I q_t)$	Q_{dd}	$(\bar{d}_p \gamma_\mu d_r)(\bar{d}_s \gamma^\mu d_t)$	Q_{ld}	$(\bar{l}_p \gamma_\mu l_r)(\bar{d}_s \gamma^\mu d_t)$		
$Q_{lq}^{(1)}$	$(\bar{l}_p \gamma_\mu l_r)(\bar{q}_s \gamma^\mu q_t)$	Q_{eu}	$(\bar{e}_p \gamma_\mu e_r)(\bar{u}_s \gamma^\mu u_t)$	Q_{qe}	$(\bar{q}_p \gamma_\mu q_r)(\bar{e}_s \gamma^\mu e_t)$		
$Q_{lq}^{(3)}$	$(\bar{l}_p \gamma_\mu \tau^I l_r)(\bar{q}_s \gamma^\mu \tau^I q_t)$	Q_{ed}	$(\bar{e}_p \gamma_\mu e_r)(\bar{d}_s \gamma^\mu d_t)$	$Q_{qu}^{(1)}$	$(\bar{q}_p \gamma_\mu q_r)(\bar{u}_s \gamma^\mu u_t)$		
		$Q_{ud}^{(1)}$	$(\bar{u}_p \gamma_\mu u_r)(\bar{d}_s \gamma^\mu d_t)$	$Q_{qu}^{(8)}$	$(\bar{q}_p \gamma_\mu T^A q_r)(\bar{u}_s \gamma^\mu T^A u_t)$		
		$Q_{ud}^{(8)}$	$(\bar{u}_p \gamma_\mu T^A u_r)(\bar{d}_s \gamma^\mu T^A d_t)$	$Q_{qd}^{(1)}$	$(\bar{q}_p \gamma_\mu q_r)(\bar{d}_s \gamma^\mu d_t)$		
				$Q_{qd}^{(8)}$	$(\bar{q}_p \gamma_\mu T^A q_r)(\bar{d}_s \gamma^\mu T^A d_t)$		
8 : $(\bar{L}R)(\bar{R}L) + \text{h.c.}$		8 : $(\bar{L}R)(\bar{L}R) + \text{h.c.}$					
Q_{ledq}	$(\bar{l}_p^j e_r)(\bar{d}_s q_{tj})$	$Q_{quqd}^{(1)}$	$(\bar{q}_p^j u_r) \epsilon_{jk} (\bar{q}_s^k d_t)$				
		$Q_{quqd}^{(8)}$	$(\bar{q}_p^j T^A u_r) \epsilon_{jk} (\bar{q}_s^k T^A d_t)$				
		$Q_{lequ}^{(1)}$	$(\bar{l}_p^j e_r) \epsilon_{jk} (\bar{q}_s^k u_t)$				
		$Q_{lequ}^{(3)}$	$(\bar{l}_p^j \sigma_{\mu\nu} e_r) \epsilon_{jk} (\bar{q}_s^k \sigma^{\mu\nu} u_t)$				

7.3.4 The Higgs basis

In the following of this thesis (Chapters 8 and 9), focus will be put on probing Higgs boson physics and more specifically to its couplings to vector bosons (denoted “HVV”). The coefficients related to other modifications in the EFT will be dropped. Therefore, a more convenient formulation could be expressed as a function of the physical states W^+ , W^- , Z , γ (mass eigenstates after symmetry breaking) instead of the fields W^i and B . This can be achieved after some redefinitions as described in Section 3 of Ref. [202]². After such transformations, and focusing on the terms relevant for the Higgs boson to vector boson couplings, the effective Lagrangian reads

$$\begin{aligned} \mathcal{L}_{\text{HVV}}^{\text{SM}+d=6} = & \frac{h}{v} \left[(1 + \delta c_w) \frac{g^2 v^2}{2} W_\mu^+ W_\mu^- + (1 + \delta c_z) \frac{(g^2 + g'^2) v^2}{4} Z_\mu Z_\mu \right. \\ & + c_{ww} \frac{g^2}{2} W_{\mu\nu}^+ W_{\mu\nu}^- + \tilde{c}_{ww} \frac{g^2}{2} W_{\mu\nu}^+ \tilde{W}_{\mu\nu}^- + c_{w\Box} g^2 (W_\mu^- \partial_\nu W_{\mu\nu}^+ + \text{h.c.}) \\ & + c_{gg} \frac{g_s^2}{4} G_{\mu\nu}^a G_{\mu\nu}^a + c_{\gamma\gamma} \frac{e^2}{4} A_{\mu\nu} A_{\mu\nu} + c_{z\gamma} \frac{e\sqrt{g^2 + g'^2}}{2} Z_{\mu\nu} A_{\mu\nu} + c_{zz} \frac{g^2 + g'^2}{4} Z_{\mu\nu} Z_{\mu\nu} \\ & + c_{z\Box} g^2 Z_\mu \partial_\nu Z_{\mu\nu} + c_{\gamma\Box} g g' Z_\mu \partial_\nu A_{\mu\nu} \\ & \left. + \tilde{c}_{gg} \frac{g_s^2}{4} G_{\mu\nu}^a \tilde{G}_{\mu\nu}^a + \tilde{c}_{\gamma\gamma} \frac{e^2}{4} A_{\mu\nu} \tilde{A}_{\mu\nu} + \tilde{c}_{z\gamma} \frac{e\sqrt{g^2 + g'^2}}{2} Z_{\mu\nu} \tilde{A}_{\mu\nu} + \tilde{c}_{zz} \frac{g^2 + g'^2}{4} Z_{\mu\nu} \tilde{Z}_{\mu\nu} \right]. \end{aligned} \quad (7.6)$$

The first line contains both Standard Model terms ($1 + \dots$) and modifiers to this terms (δc_w , δc_z). The Standard Model is recovered when all Wilson coefficients c vanish.

From this Lagrangian, we can construct another basis that should be more convenient for catching effects on Higgs boson couplings, parametrising such effects by one coefficient instead of linear combination of other coefficients. Since this set of operators is built from the operators in the Warsaw basis, the set should also be complete and independent. The Lagrangian of Eq. 7.6 however includes more coefficients than the number of required operators, meaning that some are redundant and can be expressed as a function of others. The recommended choice [202] for the independent couplings is the following:

$$\delta c_z, c_{gg}, c_{\gamma\gamma}, c_{z\gamma}, c_{zz}, c_{z\Box}, \tilde{c}_{gg}, \tilde{c}_{\gamma\gamma}, \tilde{c}_{z\gamma}, \tilde{c}_{zz}. \quad (7.7)$$

The $\tilde{c}$ denote coefficients of the CP -violating operators, while the others relate to the CP -conserving ones. The independent couplings can be recovered by comparing the operator structure in Eq. 7.6 to the operators in Table 7.3:

$$\delta c_w = \delta c_z \quad (7.8)$$

$$c_{ww} = c_{zz} + 2s_\theta^2 c_{z\gamma} + s_\theta^4 c_{\gamma\gamma} \quad (7.9)$$

$$\tilde{c}_{ww} = \tilde{c}_{zz} + 2s_\theta^2 \tilde{c}_{z\gamma} + s_\theta^4 \tilde{c}_{\gamma\gamma} \quad (7.10)$$

$$c_{w\Box} = \frac{1}{g^2 - g'^2} \left[g^2 c_{z\Box} + g'^2 c_{zz} - e^2 s_\theta^2 c_{\gamma\gamma} - (g^2 - g'^2) s_\theta^2 c_{z\gamma} \right] \quad (7.11)$$

$$c_{\gamma\Box} = \frac{1}{g^2 - g'^2} \left[2g^2 c_{z\Box} + (g^2 + g'^2) c_{zz} - e^2 c_{\gamma\gamma} - (g^2 - g'^2) c_{z\gamma} \right], \quad (7.12)$$

with $s_\theta = g'/\sqrt{g^2 + g'^2}$ and $e = gg'/\sqrt{g^2 + g'^2}$. In the first line, the additional “ $+4\delta m$ ” term has been dropped since it does not depend on Higgs boson couplings, so it will be considered as in the Standard Model, *i.e.* 0.

²↑ Ref. [202] starts from the *SILH basis* instead of the Warsaw basis, but since all bases are equivalent, a similar approach is possible starting from the Warsaw basis as well.

7.3.5 Basis translation and the SMEFTsim package

Since the analysis presented in Chapter 9 is focused on studying Higgs boson couplings, a natural EFT basis choice is the Higgs basis. The analysis is performed in the $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ decay channel, and is primarily sensitive to the Higgs-boson couplings to Z bosons on the decay side, and to the Z and W bosons on the production side through the VBF and VH modes. The ggF production mode can also provide sensitivity to the effective Higgs–gluon couplings or the top Yukawa with the resolved loop, but this is marginal and therefore dropped from the analysis. The frame of such study would still include 8 Wilson coefficients (the 10 above minus the CP -conserving and violating Higgs–gluon couplings) which is too many for the current experimental sensitivity and must therefore be further constrained by other assumptions. Since the analysis presented in Chapter 9 is furthermore focusing on the CP -odd tensor structure of the HVV couplings, probing the VBF and VH production modes, only 3 operators will be included, all the others being set to 0: $\tilde{c}_{zz}, \tilde{c}_{z\gamma}, \tilde{c}_{\gamma\gamma}$.

Aside the Higgs basis, the Warsaw basis still provides a convenient set of operators to combine results outside Higgs boson physics. This basis is implemented in the SMEFTSIM package [203] which is interfaced with various Monte Carlo generators, and used in various EFT analyses in ATLAS. It is therefore useful to be able to translate between the two bases to include Higgs boson results in a wider frame. Limiting the translation formulae to the CP -violating operators mentioned above, the three coefficients in the Higgs basis ($\tilde{c}_{zz}, \tilde{c}_{z\gamma}, \tilde{c}_{\gamma\gamma}$) can be converted to the three Warsaw basis coefficients $c_{H\tilde{W}}, c_{H\tilde{B}}, c_{H\tilde{W}B}$ according to [204]:

$$\frac{v^2}{\Lambda^2} c_{H\tilde{W}} = \frac{g^2}{4(g^2 + g'^2)^2} \left((g^2 + g'^2)^2 \tilde{c}_{zz} + 2(g^2 + g'^2) g'^2 \tilde{c}_{z\gamma} + g'^4 \tilde{c}_{\gamma\gamma} \right) \quad (7.13)$$

$$\frac{v^2}{\Lambda^2} c_{H\tilde{B}} = \frac{g'^2}{4(g^2 + g'^2)^2} \left((g^2 + g'^2)^2 \tilde{c}_{zz} - 2(g^2 + g'^2) g^2 \tilde{c}_{z\gamma} + g^4 \tilde{c}_{\gamma\gamma} \right) \quad (7.14)$$

$$\frac{v^2}{\Lambda^2} c_{H\tilde{W}B} = \frac{gg'}{2(g^2 + g'^2)^2} \left((g^2 + g'^2)^2 \tilde{c}_{zz} - (g^4 - g'^4) \tilde{c}_{z\gamma} - g^2 g'^2 \tilde{c}_{\gamma\gamma} \right). \quad (7.15)$$

The inverse translation Warsaw to Higgs is performed with the following:

$$\tilde{c}_{zz} = 4 \left(\frac{g^2 c_{H\tilde{W}} + g'^2 c_{H\tilde{B}} + gg' c_{H\tilde{W}B}}{(g^2 + g'^2)^2} \right) \frac{v^2}{\Lambda^2} \quad (7.16)$$

$$\tilde{c}_{z\gamma} = 4 \left(\frac{c_{H\tilde{W}} - c_{H\tilde{B}} - \frac{g^2 - g'^2}{2gg'} c_{H\tilde{W}B}}{g^2 + g'^2} \right) \frac{v^2}{\Lambda^2} \quad (7.17)$$

$$\tilde{c}_{\gamma\gamma} = 4 \left(\frac{1}{g^2} c_{H\tilde{W}} + \frac{1}{g'^2} c_{H\tilde{B}} - \frac{1}{gg'} c_{H\tilde{W}B} \right) \frac{v^2}{\Lambda^2}. \quad (7.18)$$

The translation formula between the Higgs and Warsaw bases are implemented in the ROSETTA framework [205]. These include both the CP -conserving and violating operators, for any flavour scheme (one, three, or three with mixing). Furthermore, other bases not mentioned in this Chapter are also included, allowing for translations between sets of parameters in a general way.

7.3.6 Higgs Characterisation framework

Finally, another EFT-related framework proves useful in the context of studying the Higgs boson couplings: the Higgs Characterisation (HC) framework [206]. Its approach is similar to the

construction of the EFT bases from the previous Sections, using the same set of operators, except that all “Wilson coefficients” are kept, even the redundant ones. Consequently, the HC cannot be considered as an EFT basis mathematically speaking, and forms a superset above them.

Yet, it is one of the first Higgs effective Lagrangian approach that has been implemented in term of `FeynRules` [207], usable by generators such as `MADGRAPH`. As such, it is still used to compute the matrix elements needed for the analysis of Chapter 9.1.

Its effective Lagrangian is implemented in terms of physical states (*i.e.* A_μ , Z_μ , W_μ^+ and W_μ^- , and not B_μ , W_μ^1 , W_μ^2 , W_μ^3). The relevant part for HVV couplings is as follows:

$$\begin{aligned} \mathcal{L}_{HVV(HC)}^{\text{SM}+d=6} = & \left\{ c_\alpha \kappa_{\text{SM}} \left[\frac{1}{2} g_{HZZ} Z_\mu Z^\mu + g_{HWW} W_\mu^+ W^{-\mu} \right] \right. \\ & - \frac{1}{4} [c_\alpha \kappa_{H\gamma\gamma} g_{H\gamma\gamma} A_{\mu\nu} A^{\mu\nu} + s_\alpha \kappa_{A\gamma\gamma} g_{A\gamma\gamma} A_{\mu\nu} \tilde{A}^{\mu\nu}] \\ & - \frac{1}{2} [c_\alpha \kappa_{HZ\gamma} g_{HZ\gamma} Z_{\mu\nu} A^{\mu\nu} + s_\alpha \kappa_{AZ\gamma} g_{AZ\gamma} Z_{\mu\nu} \tilde{A}^{\mu\nu}] \\ & - \frac{1}{4} [c_\alpha \kappa_{Hgg} g_{Hgg} G_{\mu\nu}^a G^{a,\mu\nu} + s_\alpha \kappa_{Agg} g_{Agg} G_{\mu\nu}^a \tilde{G}^{a,\mu\nu}] \\ & - \frac{1}{4} \frac{1}{\Lambda} [c_\alpha \kappa_{HZZ} Z_{\mu\nu} Z^{\mu\nu} + s_\alpha \kappa_{AZZ} Z_{\mu\nu} \tilde{Z}^{\mu\nu}] \\ & - \frac{1}{2} \frac{1}{\Lambda} [c_\alpha \kappa_{HWW} W_{\mu\nu}^+ W^{-\mu\nu} + s_\alpha \kappa_{AWW} W_{\mu\nu}^+ \tilde{W}^{-\mu\nu}] \\ & \left. - \frac{1}{\Lambda} c_\alpha [\kappa_{H\partial\gamma} Z_\nu \partial_\mu A^{\mu\nu} + \kappa_{H\partial Z} Z_\nu \partial_\mu Z^{\mu\nu} + (\kappa_{H\partial W} W_\nu^+ \partial_\mu W^{-\mu\nu} + h.c.)] \right\} X_0. \end{aligned} \quad (7.19)$$

Note that the various operators found are the same as in Eq. 7.6 (as expected), with just the coefficients in front of the them changed. The notation follow the same conventions (in particular $\tilde{X}_{\mu\nu} = \frac{1}{2} \epsilon_{\mu\nu\rho\sigma} X^{\rho\sigma}$), except that the scalar field is here denote X_0 instead of h . The g_i are the dimensionful coupling constants and the κ_i the coupling modifiers, equivalent to the Wilson coefficients. The CP -even coupling modifiers are denoted $\kappa_{H\dots}$ while the CP -odd coupling modifiers are denote $\kappa_{A\dots}$. α is the mixing angle between the CP -even and the CP -odd states: $\alpha = 0$ means a pure CP -even scalar field, while $\alpha = 1$ yields a pure CP -odd state. Any value between 0 and 1 implies a CP -admixture state, and its value is conventionally taken as $1/\sqrt{2}$ (*i.e.* $\alpha = \pi/4$). c_α and s_α are short notations for $\cos \alpha$ and $\sin \alpha$ respectively. Λ denotes the high-energy scale as defined earlier (typically, $\Lambda = 1 \text{ TeV}$).

Since the HC forms a superset of the EFT bases, the translation from the HC to the Higgs or Warsaw bases is not straightforward and requires some assumptions. However the other direction is possible. The translation formulae toward the Higgs basis for the relevant CP -odd operators are recalled below [202]:

$$s_\alpha \kappa_{A\gamma\gamma} = -3\pi^2 \tilde{c}_{\gamma\gamma} \quad (7.20)$$

$$s_\alpha \kappa_{AZZ} = -\frac{\Lambda}{v} (g^2 + g'^2) \tilde{c}_{zz} \quad (7.21)$$

$$s_\alpha \kappa_{AWW} = -\frac{\Lambda}{v} g^2 (\tilde{c}_{\gamma\gamma} + 2s_\theta^2 \tilde{c}_{z\gamma} + s_\theta^4 \tilde{c}_{\gamma\gamma}) \quad (7.22)$$

$$s_\alpha \kappa_{AZ\gamma} = -\frac{24\pi^2}{8c_\theta^2 - 5} \tilde{c}_{z\gamma}, \quad (7.23)$$

where $s_\theta = g'/\sqrt{g^2 + g'^2}$ is the same as defined above, and $c_\theta = \sqrt{1 - s_\theta^2} = g/\sqrt{g^2 + g'^2}$.

7.4 Higgs couplings and CP -mixing results from LHC Run 1 and Run 2

Couplings measurements have already been performed by the ATLAS and CMS experiments using a partial Run 2 dataset. These results provide absolute limits on the cross-section for some processes (Section 7.4.1). Some of them are further interpreted in terms of κ -framework (Section 7.4.2) or EFT (Section 7.4.3). Finally, some measurements directly target CP -odd effects by building dedicated observables (Sections 7.4.4 and 7.4.5).

Since the Higgs boson couplings has a different structure for fermions and vector bosons couplings, CP -mixing effects could be different in each sector and it is therefore important to test both. However until now, focus has been set to the couplings to vector bosons, for reasons of accessibility: on the one hand, the coupling to vector bosons can be naturally probed with $H \rightarrow ZZ$ (dominated by the $H \rightarrow 4\ell$) and $H \rightarrow WW$ decay, but also through the VBF and VH production modes. Consequently, the $H \rightarrow \gamma\gamma$ and $H \rightarrow \tau^+\tau^-$ channels also provide a complementary sensitivity to the $H \rightarrow 4\ell$ and $H \rightarrow WW$ analyses in this sector (using both production- and decay-side information).

On the other hand, the couplings to fermions was only driven by the $H \rightarrow \tau^+\tau^-$ decay until recently. Since the $H \rightarrow b\bar{b}$ decay mode [174, 175] and the $t\bar{t}H$ production mode [189, 190] have now been observed, new measurements will be possible. However, given the low cross-section of the $t\bar{t}H$ production mode, the only statistically significant analyses are at the moment the $t\bar{t}H(\gamma\gamma)$ and the $t\bar{t}H(b\bar{b})$ channels. While this seems promising, these observations are too recent for such CP results to be ready. Consequently the results presented in this Section relate to the Higgs boson coupling to vector bosons.

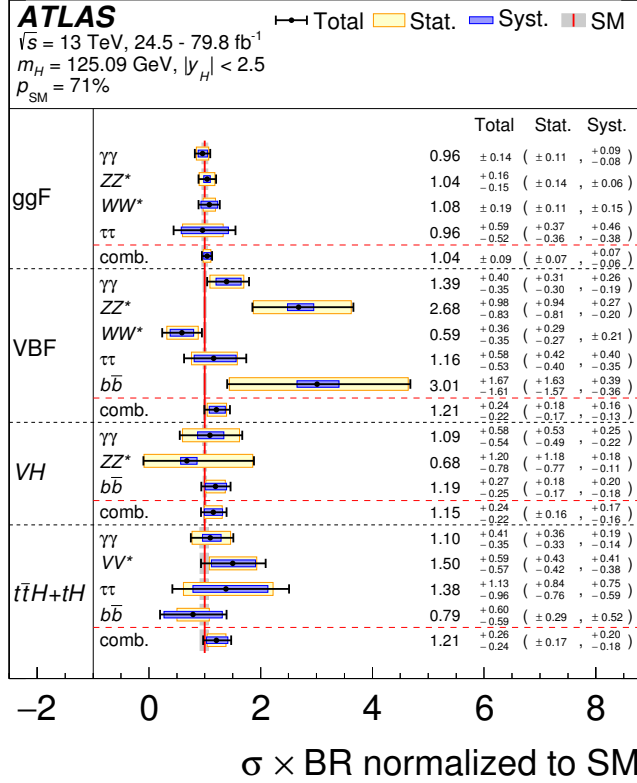
Finally, due to a lower momentum transfer, the ggF production mode only has a marginal contribution to the CP effects sensitivity.

7.4.1 STXS results

The STXS measurements using an early Run 2 dataset is presented in Figures 7.12a and 7.12b for the ATLAS [208] and CMS [209] experiments respectively. The ggF production cross-section is measured with a better than 10 % precision, while the VBF, VH , and $t\bar{t}H$ cross-sections are measured to a 20 %-30 % precision. The branching ratios of the $H \rightarrow \gamma\gamma$, $H \rightarrow 4\ell$, $H \rightarrow WW$, $H \rightarrow \tau^+\tau^-$ and $H \rightarrow b\bar{b}$ are measured at a 15 %-30 % precision level. All the results are compatible with the Standard Model expectation within their uncertainties.

7.4.2 κ -framework

The early Run 2 analyses mainly use the dataset recorded at 13 TeV in 2015 and 2016. This amounts to approximately 36 fb^{-1} , only yielding a factor of about three in Higgs boson statistics compared to the Run 1 dataset. Consequently, the κ -framework remained a baseline over EFT for the interpretation of the first Run 2 analyses in both ATLAS and CMS. The result, presented in Figure 7.13 updates the result shown in Figure 7.9b. The combined result using the $H \rightarrow \gamma\gamma$, $H \rightarrow 4\ell$, $H \rightarrow WW$, $H \rightarrow \tau^+\tau^-$ and $H \rightarrow b\bar{b}$ decay channels are shown in the (κ_F, κ_V) plane. The compatibility with the Standard Model expectation is within the 68 % CL interval for the ATLAS result, and within the 95 % CL interval for the CMS experiment. As such, no significant deviation from the Standard Model is observed.



(a) ATLAS result.

Production process	Best fit value	Uncertainty	
		stat.	syst.
ggH	1.22	$+0.14$ -0.12 $(+0.11, -0.11)$	$+0.08$ -0.08 $(+0.07, -0.07)$
VBF	0.73	$+0.30$ -0.27 $(+0.29, -0.27)$	$+0.17$ -0.15 $(+0.16, -0.15)$
WH	2.18	$+0.58$ -0.55 $(+0.53, -0.51)$	$+0.34$ -0.32 $(+0.30, -0.29)$
ZH	0.87	$+0.44$ -0.42 $(+0.43, -0.41)$	$+0.20$ -0.18 $(+0.19, -0.17)$
ttH	1.18	$+0.30$ -0.27 $(+0.28, -0.25)$	$+0.16$ -0.21 $(+0.16, -0.20)$
Decay mode	Best fit value	Uncertainty	
		stat.	syst.
$H \rightarrow b\bar{b}$	1.12	$+0.29$ -0.29 $(+0.28, -0.27)$	$+0.22$ -0.22 $(+0.21, -0.20)$
$H \rightarrow \tau\tau$	1.02	$+0.26$ -0.24 $(+0.24, -0.22)$	$+0.21$ -0.19 $(+0.19, -0.17)$
$H \rightarrow WW$	1.28	$+0.17$ -0.16 $(+0.14, -0.13)$	$+0.14$ -0.13 $(+0.09, -0.10)$
$H \rightarrow ZZ$	1.06	$+0.19$ -0.17 $(+0.18, -0.16)$	$+0.11$ -0.08 $(+0.10, -0.08)$
$H \rightarrow \gamma\gamma$	1.20	$+0.18$ -0.14 $(+0.14, -0.12)$	$+0.12$ -0.09 $(+0.10, -0.07)$
$H \rightarrow \mu\mu$	0.68	$+1.25$ -1.24 $(+1.20, -1.17)$	$+0.13$ -0.11 $(+0.19, -0.03)$

(b) CMS result.

Figure 7.12 – Signal strength (cross-section, branching ratio or $\sigma \times \text{BR}$ normalised to the Standard Model expectation) from partial Run 2 STXS results by the (a) ATLAS and (b) CMS experiments. The ATLAS result [208] includes 79.8 fb^{-1} of 13 TeV Run 2 data for the $H \rightarrow \gamma\gamma$, $H \rightarrow 4\ell$ and $H \rightarrow b\bar{b}$ channels, and 36.1 fb^{-1} for the $H \rightarrow WW$ and $H \rightarrow \tau^+\tau^-$ channels. The global compatibility of the results with the SM hypothesis has a p -value of 71 %. The CMS result [209] includes 35.9 fb^{-1} of 13 TeV Run 2 data for the $H \rightarrow \gamma\gamma$, $H \rightarrow 4\ell$, $H \rightarrow b\bar{b}$, $H \rightarrow WW$ and $H \rightarrow \tau^+\tau^-$ decay channels. The brackets show the expected uncertainty.

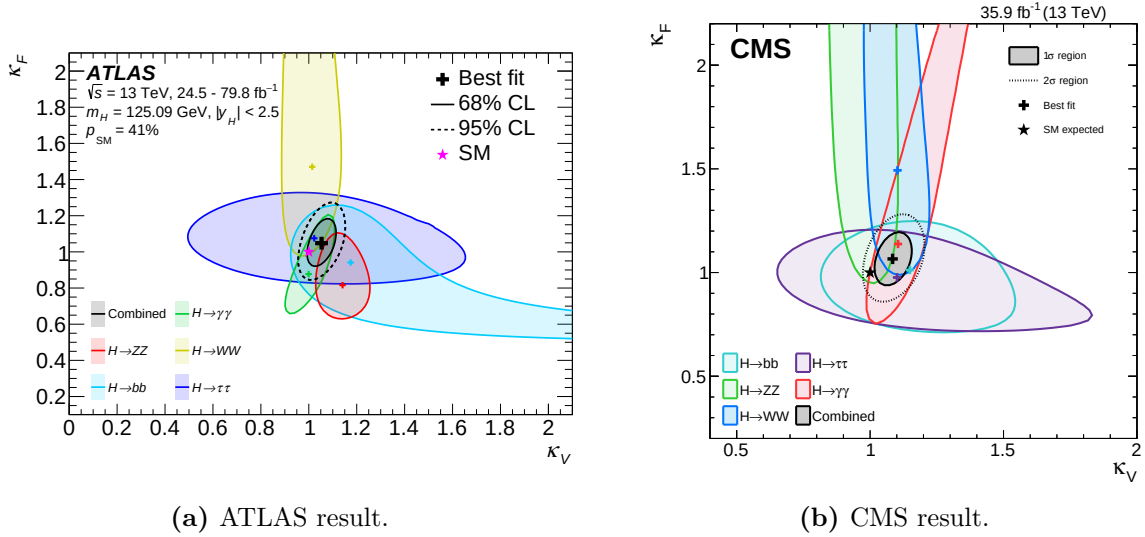


Figure 7.13 – Couplings results interpreted in terms of couplings modifiers in the κ -framework. The couplings modifiers are merged as $\kappa_F = \kappa_t = \kappa_b = \kappa_\tau = \kappa_\mu$ and $\kappa_V = \kappa_Z = \kappa_W$. The contours show the 68 % CL limits in the (κ_F, κ_V) plane for individual channels as well as their combination. The ATLAS result (a) [208] includes 79.8 fb^{-1} of 13 TeV Run 2 data for the $H \rightarrow \gamma\gamma$, $H \rightarrow 4\ell$ and $H \rightarrow b\bar{b}$ channels, and 36.1 fb^{-1} for the $H \rightarrow WW$ and $H \rightarrow \tau^+\tau^-$ channels. The p -value of the best fit point with the SM hypothesis is 41 %. The CMS result (b) [209] includes 35.9 fb^{-1} of 13 TeV Run 2 data for the $H \rightarrow \gamma\gamma$, $H \rightarrow 4\ell$, $H \rightarrow b\bar{b}$, $H \rightarrow WW$ and $H \rightarrow \tau^+\tau^-$ decay channels. The SM result lies within the 95 % CL area around the best fit value.

7.4.3 EFT interpretation

An EFT interpretation of cross-section measurements has been performed by the ATLAS experiment using 36.1 fb^{-1} of 13 TeV data collected during Run 2, in the $H \rightarrow 4\ell$ decay channel [178] using the Higgs Characterisation framework, and in the $H \rightarrow \gamma\gamma$ decay channel [179] using the SILH basis³.

The $H \rightarrow 4\ell$ result is presented in Figure 7.14. Given the limited amount of data available, it is assumed that $k_{HVV} = k_{HZZ} = k_{HWW}$ and $k_{AVV} = k_{AZZ} = k_{AWW}$. The observed (resp. expected) 95 % CL interval for k_{HVV} is $[0.8, 4.5]$ (resp. $[-2.9, 3.2]$) with a best fit value of 2.9. A slight deviation from the Standard Model of 2.3σ is observed. For the CP -violating couplings k_{AVV} , the observed (resp. expected) 95 % CL interval is $[-5.2, 5.2]$ (resp. $[-3.5, 3.5]$) with a best fit value at ± 2.9 . This yields a compatibility with the Standard Model of about 1.4σ . Since this measurement is based on cross-sections only, the sign of the CP -odd contribution cannot be determined. The equivalent Run 1 result from Ref. [68] used observables defined similarly as in Chapter 9.1.2. Combining the sensitivity from the $H \rightarrow 4\ell$ and $H \rightarrow WW$ decay channels, the expected sensitivity is around 1.5 times better than this partial Run 2 cross-sections only result, showing the power of the matrix element observables.

The $H \rightarrow \gamma\gamma$ result [179] is shown in Figure 7.15, providing an update of the Run 1 result of Ref. [213]. The limits are shown in the $(\tilde{c}_{HW}, \tilde{c}_{HW})$ plane of the SILH basis. The $\tilde{c}_{HW}$ direction corresponds to a CP -violating coupling to the vector bosons. This result uses a combination of differential cross-sections in events with two jets, building a discrimination from the five following variables:

- the di-photon transverse momentum,

³↑ The Strongly Interacting Light Higgs (SILH) basis is another EFT basis, which can be mapped to the Warsaw (Section 7.3.3) or Higgs basis (Section 7.3.4). However, as this basis is not used for the analysis described in Chapter 9, it will not be detailed here. See for example Refs. [210, 211] and [212].

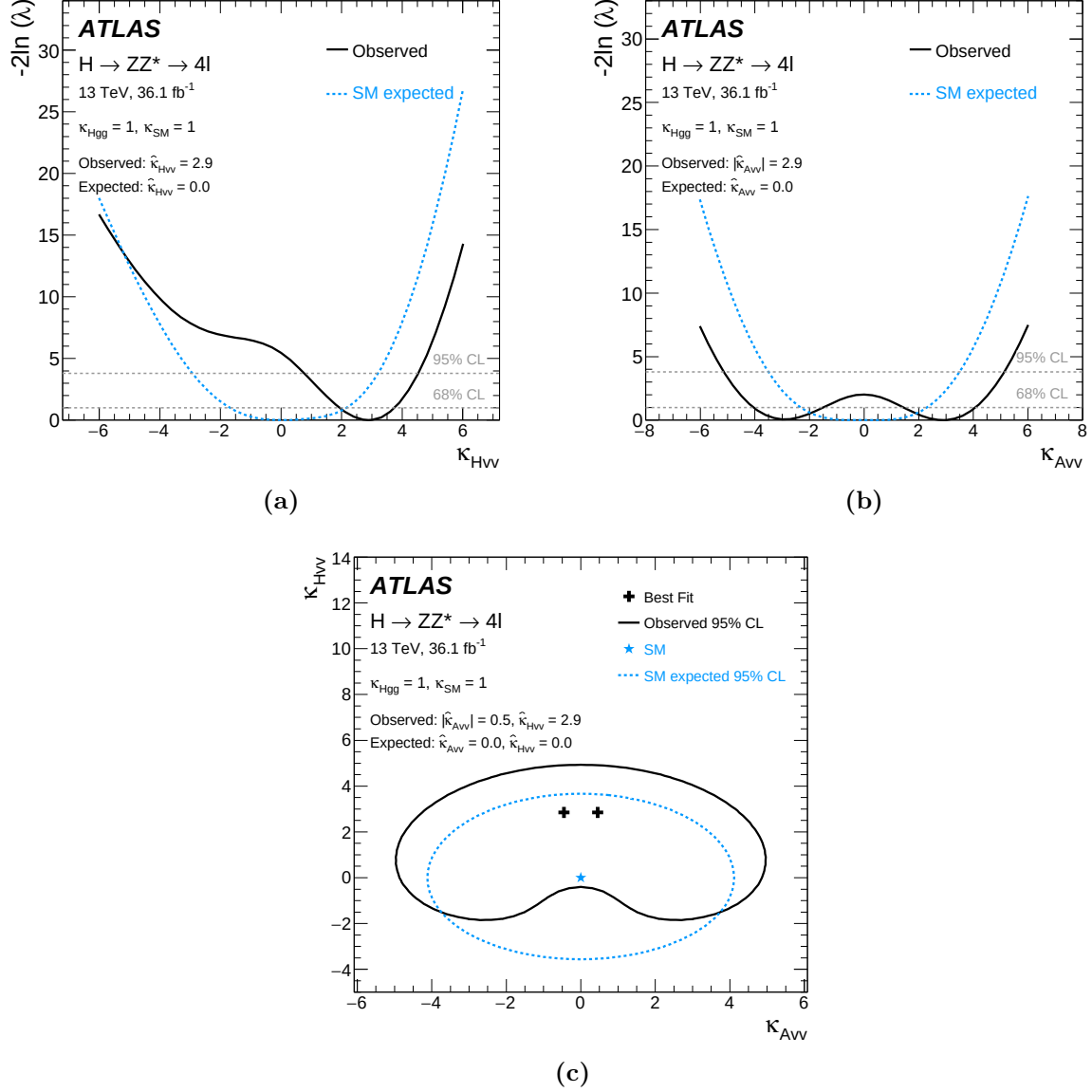


Figure 7.14 – ATLAS couplings results in the $H \rightarrow 4\ell$ decay channel using 36.1 fb $^{-1}$ of 13 TeV data collected during Run 2 and interpreted in the Higgs Characterisation framework. The likelihood scans yielding the one-dimensional limits for (a) κ_{HVV} and (b) κ_{AVV} , and the two-dimensional limits for (c) κ_{HVV} against κ_{AVV} are shown. The blue curve represents the expected sensitivity under the Standard Model hypothesis, while the black curve represents the observed sensitivity. A non-significant deviation of 2.3σ is seen for the κ_{HVV} 1D scan. Cross-section only measurements cannot distinguish the sign of CP effects, so the κ_{AVV} must be symmetrical. This result is compatible with the SM expectation within 2σ . [178]

- the number of jets in the event,
- the jet highest transverse momentum,
- the dijet invariant mass,
- and the angular separation in the ϕ plane between the two jets.

Both the yields and the shape of the observables are used in the discrimination. The observed (resp. expected) 95 % CL interval on $\bar{c}_{HW}$ is $[-0.057, 0.051]$ ($[-0.050, 0.050]$), and $[-0.16, 0.16]$ (resp. $[-0.14, 0.14]$) on $\tilde{c}_{HW}$. If the yields are dropped, the expected limits weaken by 20 % to 50 %. The results are compatible with the Standard Model expectation within the 95 % CL.

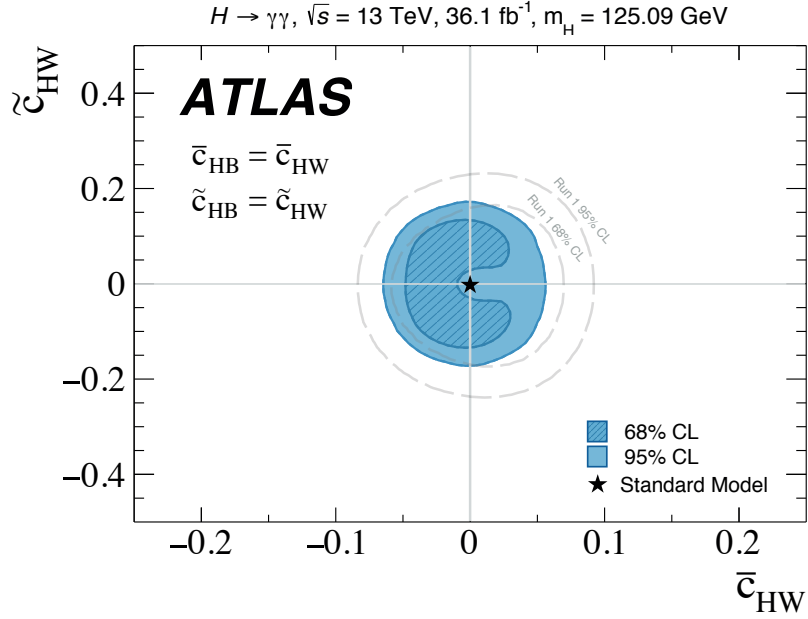


Figure 7.15 – ATLAS couplings results in the $H \rightarrow \gamma\gamma$ decay channel plane using 36.1 fb⁻¹ of 13 TeV data collected during Run 2. The 68 % CL and 95 % CL interval in the $(\bar{c}_{HW}, \tilde{c}_{HW})$ plane (SILH basis) are shown and compared to the Run 1 result [213]. It is assumed that $\bar{c}_{HW} = \bar{c}_{HB}$ and $\tilde{c}_{HW} = \tilde{c}_{HB}$, and that all other Wilson coefficients vanish. [179]

7.4.4 Direct CP measurements and couplings structure functions

The CMS experiment uses an approach slightly different to EFT interpretations, rather relying on the parametrisation of the Higgs boson to vector boson amplitude. This approach is for example developed in Refs. [214–216] and [217]. This parametrisation reads [188]:

$$\mathcal{A} \sim \left(a_1^{VV} + \frac{\kappa_1^{VV} q_1^2 + \kappa_2^{VV} q_2^2}{(\Lambda_1^{VV})^2} \right) \cdot m_{V1}^2 \epsilon_{V1}^* \epsilon_{V2}^* + a_2^{VV} f_{\mu\nu}^{*(1)} f^{*(2)\mu\nu} + a_3^{VV} f_{\mu\nu}^{*(1)} \tilde{f}^{*(2)\mu\nu}, \quad (7.24)$$

where

- q_1 and q_2 are the four momenta of the vector bosons,
- m_{V1} is the pole mass of the vector boson,
- ϵ_{Vi} is the polarisation vector of the vector boson i ,
- $f^{(i)\mu\nu} = \epsilon_i^\mu q_i^\nu - \epsilon_i^\nu q_i^\mu$,

- $\tilde{f}_{\mu\nu}^{(i)} = \frac{1}{2}\epsilon_{\mu\nu\rho\sigma}f^{(i)\rho\sigma}$ with $\epsilon_{\mu\nu\rho\sigma}$ the fully anti-symmetric Levi-Civita tensor,
- Λ_1^{VV} and Λ_Q^{VV} are the energy scales of BSM physics,
- the a_i^{VV} are the coupling strength modifiers (complex numbers),
- and the $\kappa_{1,2,3}^{VV}$ are complex numbers such that $|\kappa_{1,2,3}^{VV}| = 0$ or 1.

If the coupling constants are constant and real, then this is equivalent to an effective Lagrangian.

In order to have a better control on systematic uncertainties, it is easier to measure ratios of quantities of interest. Therefore, one defines [188]:

$$f_{a3} = \frac{|a_3|^2 \sigma_3}{|a_1|^2 \sigma_1 + |a_2|^2 \sigma_2 + |a_3|^2 \sigma_3 + \tilde{\sigma}_{\Lambda 1}/(\Lambda_1)^4}, \quad \phi_{a3} = \arg\left(\frac{a_3}{a_1}\right), \quad (7.25)$$

$$f_{a2} = \frac{|a_2|^2 \sigma_2}{|a_1|^2 \sigma_1 + |a_2|^2 \sigma_2 + |a_3|^2 \sigma_3 + \tilde{\sigma}_{\Lambda 1}/(\Lambda_1)^4}, \quad \phi_{a2} = \arg\left(\frac{a_2}{a_1}\right), \quad (7.26)$$

$$f_{\Lambda 1} = \frac{\tilde{\sigma}_{\Lambda 1}/(\Lambda_1)^4}{|a_1|^2 \sigma_1 + |a_2|^2 \sigma_2 + |a_3|^2 \sigma_3 + \tilde{\sigma}_{\Lambda 1}/(\Lambda_1)^4}, \quad \phi_{\Lambda 1}. \quad (7.27)$$

Thus, the f_i and ϕ_i are interpreted as effective fractional cross-sections and phases. In particular f_{a3} is interpreted as an effective fraction of CP -odd cross-section in the Higgs boson couplings to the vector bosons.

The analyses construct specific discriminating variables targeting each effective fraction of interest, in a pure- or mixed- CP state hypothesis. These discriminants statistically combine the angular information built from the vector bosons and their attached fermion lines, as well as the invariant masses of these bosons. These variables appear both in the decay side in the $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ decay and in the production side in the VBF production mode, in a similar fashion as described in Section 9.1.1. The discriminants targeting the f_{a3} parameter in case of a mixed- CP state are shown in Figure 7.16.

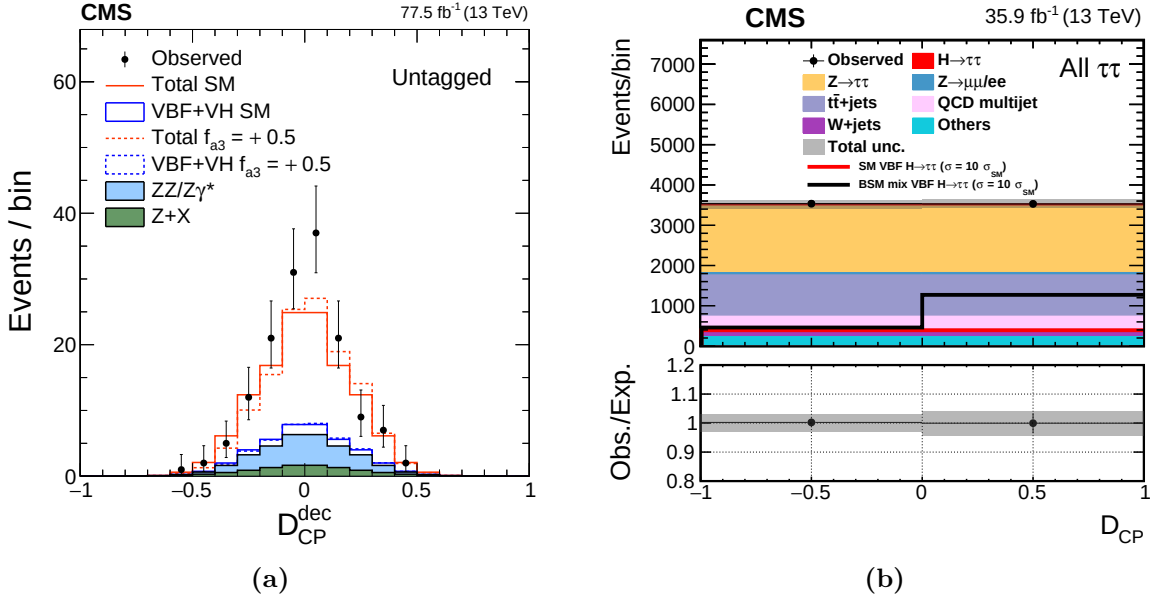


Figure 7.16 – Distributions of CP discriminants used in the measurement of the f_{a3} parameter by the CMS experiment (a) in the $H \rightarrow 4\ell$ analysis [218] using 77.5 fb⁻¹ of 13 TeV data collected in 2016 and 2017, and (b) in the VBF $H \rightarrow \tau^+\tau^-$ analysis [188] using 35.9 fb⁻¹ of 13 TeV data collected in Run 2.

After the Run 1 results developed in Refs. [70] and [64], results using a partial Run 2 dataset are developed in Refs. [218] and [188]. They are recalled in Figure 7.17 and summarised in Table 7.4.

All the scanned parameters $f_{a3} \cdot \cos(\phi_{a3})$, $f_{a2} \cdot \cos(\phi_{a2})$, $f_{\Lambda 1} \cdot \cos(\phi_{\Lambda 1})$ and $f_{\Lambda 1}^{Z\gamma} \cdot \cos(\phi_{\Lambda 1}^{Z\gamma})$ are compatible with the Standard Model expectation at 0.

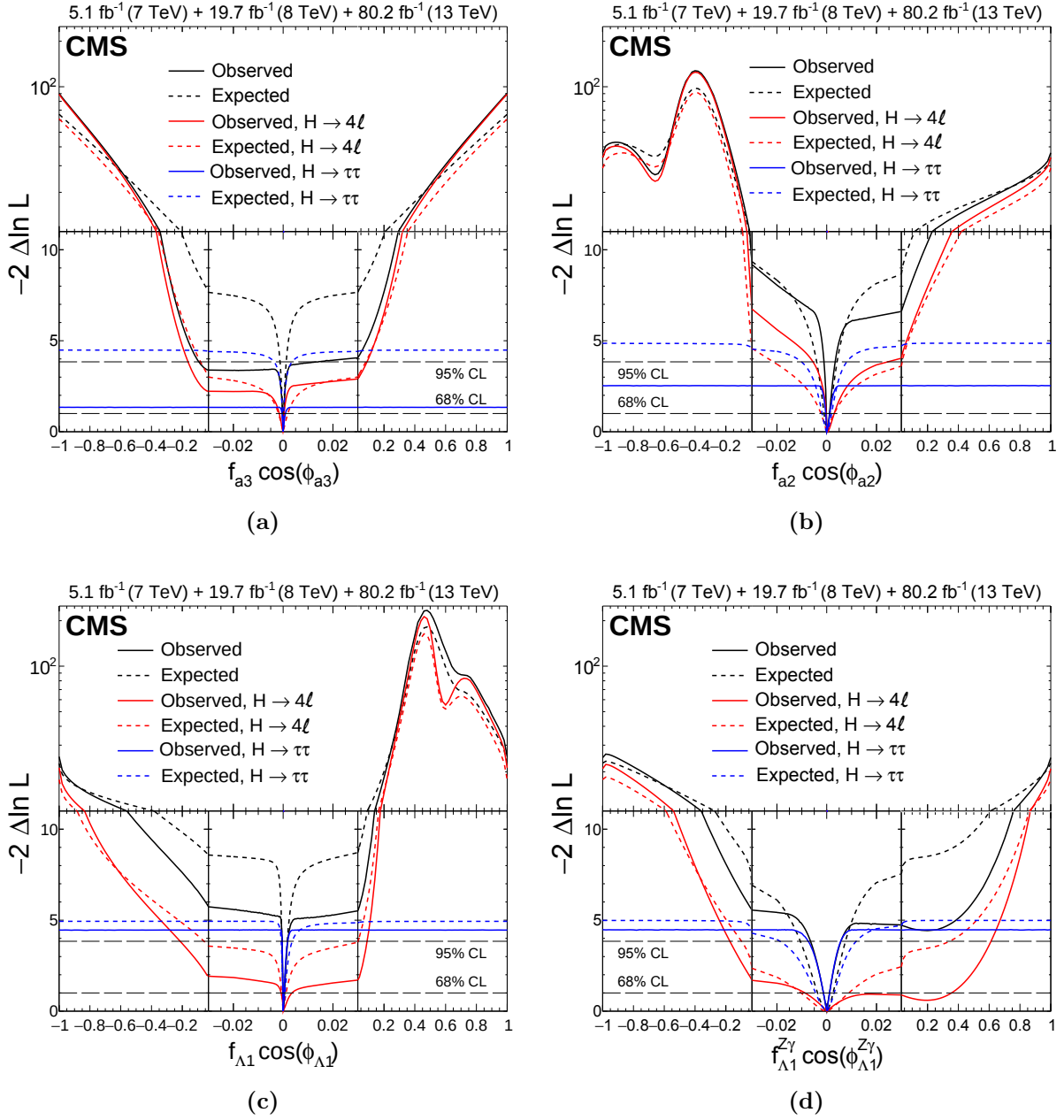


Figure 7.17 – Combined likelihood scan of the CMS $H \rightarrow \tau^+\tau^-$ and $H \rightarrow 4\ell$ results on effective couplings for (a) $f_{a3} \cdot \cos(\phi_{a3})$, (b) $f_{a2} \cdot \cos(\phi_{a2})$, (c) $f_{\Lambda 1} \cdot \cos(\phi_{\Lambda 1})$, (d) $f_{\Lambda 1}^{Z\gamma} \cdot \cos(\phi_{\Lambda 1}^{Z\gamma})$. The dataset includes all the Run 1 data, plus 35.9 fb^{-1} of 13 TeV data for the $H \rightarrow \tau^+\tau^-$ analysis [188] or 80.2 fb^{-1} of 13 TeV data for the $H \rightarrow 4\ell$ analysis [218]. The $H \rightarrow \tau^+\tau^-$ analysis uses only VBF-like events to constrain the couplings to vector bosons.

7.4.5 Direct CP measurement with single parameter parametrisation of the effective Lagrangian

Finally, another approach is used by the ATLAS $H \rightarrow \tau^+\tau^-$ analysis, parametrising an Higgs effective Lagrangian with a single CP -odd parameter $\tilde{d}$ (under some simplifying assumptions) [71]:

$$\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{SM}} + \frac{g}{2m_W} \tilde{d} \cdot \left(H \tilde{A}_{\mu\nu} A^{\mu\nu} + H \tilde{Z}_{\mu\nu} Z^{\mu\nu} + 2H \tilde{W}_{\mu\nu}^+ W^{-\mu\nu} \right). \quad (7.28)$$

Table 7.4 – Combined 68 % CL and 95 % CL intervals for the CMS effective coupling analysis using $H \rightarrow \tau^+\tau^-$ and $H \rightarrow 4\ell$ results. The dataset includes all the Run 1 data, plus 35.9fb^{-1} of 13 TeV data for the $H \rightarrow \tau^+\tau^-$ analysis [188] or 80.2fb^{-1} of 13 TeV data for the $H \rightarrow 4\ell$ analysis [218].

Parameter	Observed/(10^{-3})		Expected/(10^{-3})	
	68% CL	95% CL	68% CL	95% CL
$f_{a3} \cos(\phi_{a3})$	0.00 ± 0.27	$[-92, 14]$	0.00 ± 0.23	$[-1.2, 1.2]$
$f_{a2} \cos(\phi_{a2})$	$0.08^{+1.04}_{-0.21}$	$[-1.1, 3.4]$	$0.0^{+1.3}_{-1.1}$	$[-4.0, 4.2]$
$f_{\Lambda 1} \cos(\phi_{\Lambda 1})$	$0.00^{+0.53}_{-0.09}$	$[-0.4, 1.8]$	$0.00^{+0.48}_{-0.12}$	$[-0.5, 1.7]$
$f_{\Lambda 1}^{Z\gamma} \cos(\phi_{\Lambda 1}^{Z\gamma})$	$0.0^{+1.1}_{-1.3}$	$[-6.5, 5.7]$	$0.0^{+2.6}_{-3.6}$	$[-11, 8.0]$

This approach is a restricted case of the Higgs Characterisation interpretation, but contrary to the $H \rightarrow 4\ell$ result in Section 7.4.3 only interpreting cross-section measurements, the present $H \rightarrow \tau^+\tau^-$ analysis provides a direct measurement of possible CP effects.

The analysis uses a matrix element observable dedicated to the measurement of CP effects, similarly to the procedure described in Chapter 9.1.2. The distribution of this *optimal observable* is shown in Figure 7.18a for simulated VBF events with various values of $\tilde{d}$. This shows that the distribution becomes asymmetric with non-vanishing $\tilde{d}$ values. The distribution with both signal and background events included (with VBF-like topologies) is shown in Figure 7.18b and compared to the data. The observed and expected distribution under hypothesis $\tilde{d} = 0$ are in agreement, and it is concluded that no sign of CP -violation is expected. The observed and expected distributions under hypothesis $\tilde{d} = 0$ are in agreement, and it is concluded that no sign of CP violation is expected.

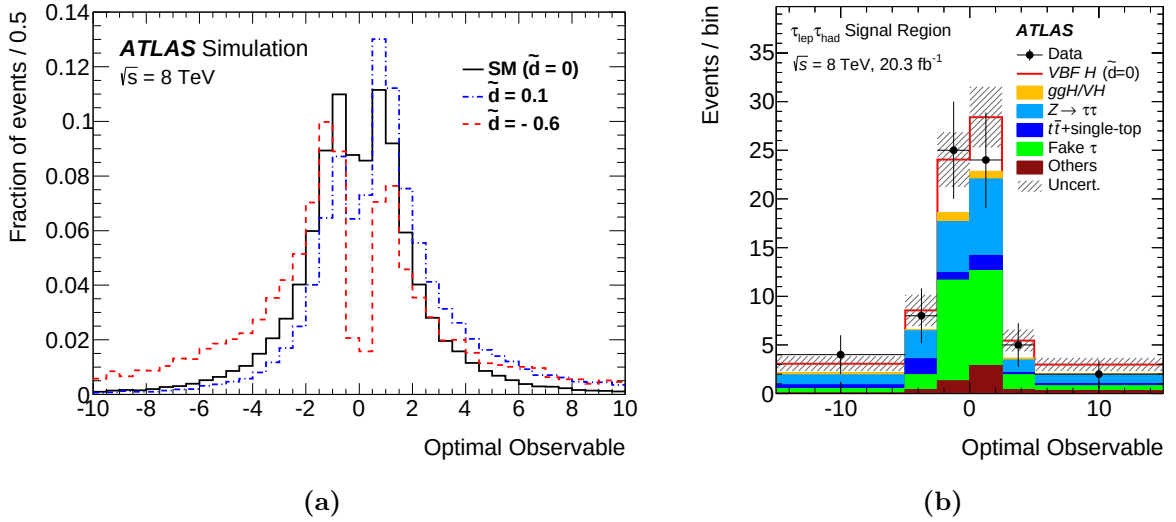


Figure 7.18 – Distribution of CP -odd *optimal observable* for (a) simulated VBF events with various $\tilde{d}$ hypotheses, and (b) simulated events from various processes (scaled to the observed signal strength $\mu = 1.55^{+0.87}_{-0.76}$) and compared to the data in the $\tau_{\text{lep}}\tau_{\text{had}}$ signal region. The analysis uses 20.3fb^{-1} of 8 TeV data collected by the ATLAS experiment in 2012. Events with a VBF-like topology are selected. [71]

A likelihood scan using this optimal observable is used to put limits on the value of $\tilde{d}$: The observed (expected) 68 % CL interval on $\tilde{d}$ is $[-0.11, 0.05]$ ($[-0.08, 0.08]$). using 20.3fb^{-1} of 8 TeV data collected in Run 1, the observed (resp. expected) 68 % CL interval on $\tilde{d}$ is $[-0.11, 0.05]$ (resp.

$[-0.08, 0.08]$), as shown in Figure 7.19. The results are compatible with the Standard Model expectation, confirming that no hint of CP -violation in the Higgs boson couplings to vector boson is seen. using 20.3 fb $^{-1}$ of 8 TeV data collected in Run 1, the observed (resp. expected) 68 % CL interval on $\tilde{d}$ is $[-0.11, 0.05]$ (resp. $[-0.08, 0.08]$), as shown in Figure 7.19. The results are compatible with the Standard Model expectation, confirming that no hint of CP violation in the Higgs boson couplings to vector boson is seen. This analysis yields a 68 % CL constraint on CP -odd parameters a factor of 10 better than the Run 1 $H \rightarrow ZZ + H \rightarrow WW$ result in the Higgs Characterisation framework of Ref. [68].

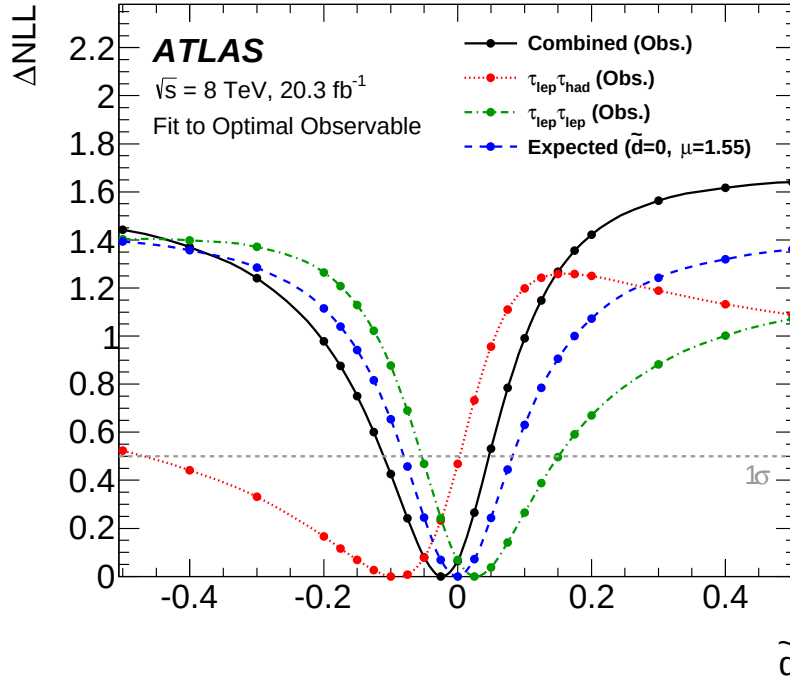


Figure 7.19 – Expected (blue) and observed (black) likelihood scans for the $\tilde{d}$ in the using the optimal observable as discriminating variable. The expected result uses an Asimov dataset scaled with the observed signal strength ($\mu = 1.55^{+0.87}_{-0.76}$). [71]

Experimental aspects of the Higgs to Four Leptons decay

Contents

8.1	Overview of the analysis	204
8.2	Data and simulated samples	205
8.2.1	Data samples	205
8.2.2	Steps of an event generation	206
8.2.2.1	Hard-scatter	206
8.2.2.2	Additional radiations	207
8.2.2.3	Hadronisation	207
8.2.3	Simulated signal samples	207
8.2.4	Simulated background samples	208
8.3	Event selection	208
8.3.1	Trigger requirements	209
8.3.2	Physics objects selection	209
8.3.3	Lepton quadruplet building	210
8.3.4	Lepton quadruplet selection	211
8.4	Backgrounds estimation	212
8.4.1	Non-resonant background	212
8.4.2	Reducible $\ell\ell + \mu\mu$ background	213
8.4.3	$\ell\ell + \mu\mu$ transfer factor systematics uncertainties	218
8.4.4	Reducible $\ell\ell + ee$ background	220
8.4.5	Extrapolation to the SR using the $Z+X$ CR	222
8.4.6	Categorised background yield and determination of the background shapes	224
8.5	Systematic uncertainties common to all $H \rightarrow 4\ell$ analyses	225
8.5.1	Experimental uncertainties	226
8.5.2	Theory uncertainties on the signal modelling	227
8.5.3	Background uncertainties	228
8.6	Event categorisation for the couplings analysis	228
8.6.1	STXS categorisation	229
8.6.2	Improved category separation using neural networks	230
8.7	Expected yields with the full Run-2 dataset	232

This Chapter describes the general features of the $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ analyses. The simulation samples required to perform the analyses are described in Section 8.2, and the general event selection is detailed in Section 8.3, along with the background estimations in Section 8.4. An overview of the experimental and theoretical systematic uncertainties affecting all the analyses is given in Section 8.5. A further categorisation for couplings analyses using multivariate analyses and the STXS framework is given in Section 8.6. The observed and expected inclusive yields are presented in Section 8.7, along with some distributions.

8.1 Overview of the analysis

The four-lepton decay channel is one of the discovery modes of the Higgs boson by the ATLAS [37] and CMS [38] experiments in 2012, along with the two other bosonic decay channels $H \rightarrow \gamma\gamma$ and $H \rightarrow WW$. After the discovery during the Run 1 of the LHC, the Run 2 analyses focus on the precise determination of its properties.

The $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ ($\ell = e$ or μ) decay channel (Figure 8.1), on top of providing a good discovery potential at low mass, is also able to conduct precise measurements of mass, cross-sections, couplings, spin-parity (denoted CP), *etc.* To make all $H \rightarrow 4\ell$ measurements consistent, the event selection described in this Chapter is common to all $H \rightarrow 4\ell$ analyses, in particular the CP analysis described in Chapter 9. The good sensitivity of this channel is ensured by a fully-reconstructed leptonic final state: jets are not part of the signal (except for tagging specific production modes), thus there is no significant QCD background contamination, and taus are excluded from the final state to avoid the presence of undetected neutrinos.

New physics searches in the $H \rightarrow 4\ell$ channel benefit from the four-body final state, in which five angles can be defined, some of them providing anomalous spin-parity sensitivity. Since the Higgs boson mass is below the double of the Z mass, one of the two Z boson is off-shell in the decay. Its mass can provide a proxy for spin and parity measurements. Moreover, the jet kinematics in the production mode with associated jets (VBF) provide additional sensitivity to BSM effects. At large momentum transfer, observables on the production side surpass the sensitivity provided by the decay side observables. Consequently, the CP analysis detailed in Chapter 9 is focused on the VBF production mode. This channel stands as one of the best candidates for the CP measurements targeted in this Thesis.

Due to a low branching ratio of 1.24×10^{-4} , measurements performed in the $H \rightarrow 4\ell$ channel are statistically limited with the Run 2 dataset. This low yield is however compensated by a low background as well, with a signal-to-background ratio almost reaching 2 in the signal region $115 < m_{4\ell} < 130$ GeV. While the background from quark- or gluon-induced Z boson pair production is small and follows a smooth shape in the mass signal region, it also provides a resonant contribution from the Z boson directly decaying to four leptons not far from the Higgs boson resonance. This resonant background can be used to validate some of the analysis steps independently from the signal region, for example when comparing to the simulation

Thanks to its versatility in providing clean measurements, the $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ is nicknamed *golden channel*. The rest of this chapter describes the various steps leading to the Higgs boson CP analysis, described in Chapter 9.

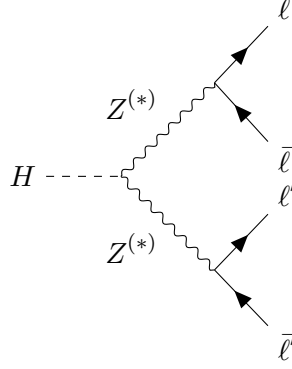


Figure 8.1 – Leading order diagram for the $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ decay.

8.2 Data and simulated samples

In order to estimate the signal and backgrounds shape and yields, and to extract meaningful quantities, the data are compared to simulated event distributions. The experimental data samples are described in Section 8.2.1 while the generation of simulated events, is briefly described in Section 8.2.2. The generated signal and background samples used in ATLAS are described in Sections 8.2.3 and 8.2.4. These generator, or *truth-level*, samples are then passed through a full simulation of the ATLAS detector [219] using the GEANT4 framework [220], which is not further described here.

Additional minimum bias interactions are separately generated using PYTHIA 8 [221] with the A3 tune [222] and the NNPDF 2.3 LO PDF set [223]. These are superimposed on the simulated signal and background events during the digitisation step to account for the pileup in the same and neighbouring bunch crossings¹. The number of interactions superimposed is done randomly, following the measured pileup distribution seen in the data at the beginning of each year of data-taking. Small differences between the predicted and actually recorded pileup distributions for the full data-taking period are accounted for by the weighting of the simulated events according to these distributions.

Finally, both real and simulated data are reconstructed using the same central ATLAS reconstruction software described in Chapter 4.

8.2.1 Data samples

The standard ATLAS physics campaign during the full Run 2 amounts to around 147 fb^{-1} of data recorded at 13 TeV, with an average data taking efficiency of $\sim 94\%$. Each ATLAS *run*, roughly corresponding to one LHC *fill* with stable beams, is split into short periods of about 1 minute, called *lumi-block*, during which the run conditions such as luminosity, triggers and sub-detector status are stable. Before being used by analyses, the recorded data undergo quality requirements that veto lumi-blocks where the detector underwent a problem. Additionally, distributions of relevant quantities (number of clusters, hits, average energy, trigger rates) are compared to reference distributions to detect long term drifts or transient issues.

After applying these data quality requirements, the selected “good for physics” data amount to 139 fb^{-1} with an uncertainty of 1.7% [124]. The main run characteristics for each years of the Run 2 campaign were recalled in Chapter 2.

¹↑ Neighbouring bunch crossings are needed to approximately simulate out-of-time pileup in the muon system (Chapter 3.5), the TRT (Chapter 3.3.3), and the 400 ns LAr calorimeter signals (Chapter 3.4.3). Bunch crossings occur every 25 ns (see Chapter 2).

8.2.2 Steps of an event generation

The generation of a simulated event follows a step-by-step procedure: generation of the process matrix element, addition of initial and final-state radiations, showering of the partons, hadronisation and decay.

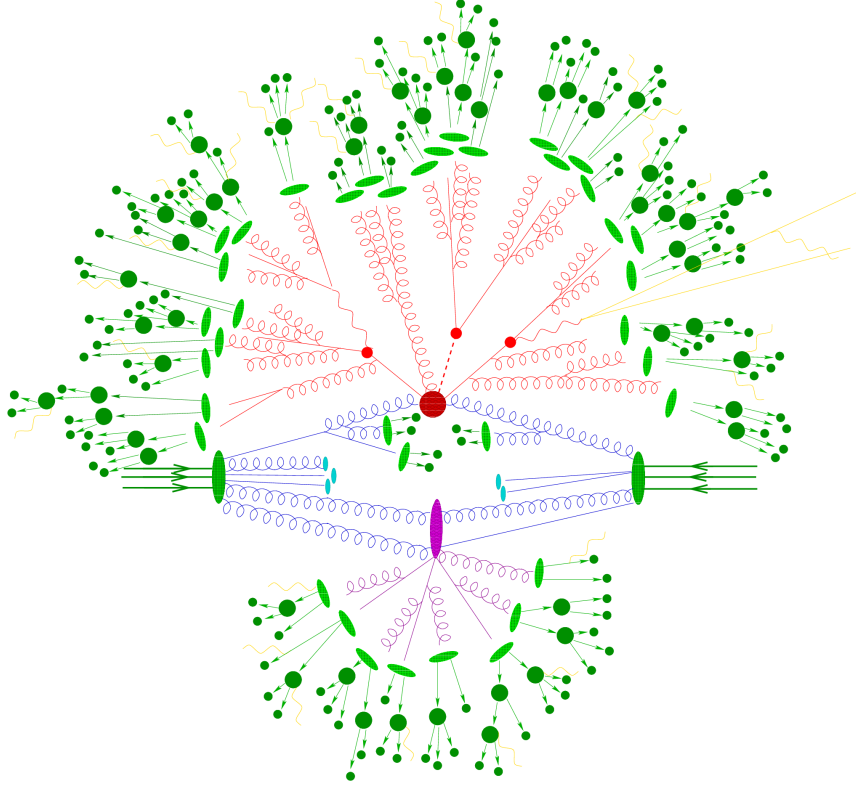


Figure 8.2 – Schematic of an event layout after all generation steps (hard-scatter/underlying event, ISR/FSR, parton showering and hadronisation). The hard-scatter is represented by the red blob, the parton showering by the smaller red dots. The hadronisation step is represented by the light green blobs, and the hadron decay by the dark green dots (in cascade). Final-state radiations are shown by the thin yellow lines. An additional underlying event is represented by the purple blob and undergoes the same steps. [224]

8.2.2.1 Hard-scatter

Due to the probabilistic nature of collisions, it would not be efficient to generate events proportionally to their actual cross-section: rare processes of interest such as the Higgs boson production would only be generated a few times while many events of “uninteresting” processes would be generated. Instead, event generators target a specific process, relying on perturbative calculations from the QFT to compute cross-sections and event kinematics.

This first step of the event generation is called the *hard-scatter* process or *matrix element* generation, which corresponds to the highest momentum transfer part of the collision, and depends on the PDFs as mentioned in Chapter 2.9. However, as QFT is a perturbative theory, the number of orders included in the hard scatter computation is often a compromise between required accuracy and available computing time. As a consequence, even though the inclusive cross-section (the probability of occurrence) of many processes is known up to NNLO or N³LO (see Chapter 7.1), the event generation (the kinematics of the outgoing particles) is mostly available at only NLO. Commonly used generators include for instance POWHEG, MADGRAPH or SHERPA.

8.2.2.2 Additional radiations

After the hard process has been generated, initial and final state radiations (ISR and FSR respectively) still remain to be simulated. The ISR and FSR include emission of additional photons by any charged fermion (electrons, muons and quarks), and of gluons by quarks or other gluons, which affect the kinematics of both the incoming (corrections must be back-propagated in time) and outgoing kinematics.

Moreover, since the hard process is generated at NLO, it can also provide additional jets in case of real emission lines. In order to avoid double counting between the matrix element and the parton shower jets and therefore to bias the cross-sections and kinematics, the additional parton emission has to be *matched* and *merged*. Several prescriptions for this procedure are available, as MINLO [225] or MEPS @ NLO [226] for example. The matching/merging is a coordinated procedure between the matrix element generator and the parton shower (*e.g.* PYTHIA).

8.2.2.3 Hadronisation

One interesting property of QCD is that its potential is increasing at large radius, leading to a phenomenon called *colour confinement*. This phenomenon has the following feature: given a pair of coloured states (quarks or gluons) whose separation is increasing, the energy of the system will increase up to the point that this binding energy is sufficient to create a quark–antiquark pair. Consequently, quarks and gluons cannot be observed as free particles and will instead manifest as QCD bound states, *i.e.* hadrons.

This phenomenon is mainly modelled with the Lund *string model* [227] and with the *cluster model* [228, 229], implemented for example in PYTHIA. This can be complemented by the decay of the *c*- and *b*-hadrons just formed in some specialised software as in EVTGEN [230].

8.2.3 Simulated signal samples

All SM Higgs boson production modes are taken into account in the simulated samples: ggF, VBF, *VH*, *ttH*, *bbH* and *tH*. The generators settings for each production mode are detailed below and summarised in Table 8.1. Additional samples for BSM processes will be described in Chapter 9.3.1.

Table 8.1 – Summary of the generators and settings used for the generation of Standard Model Higgs boson signal samples. The showering and Higgs boson decay are always performed with the PYTHIA 8 software, using however different PDF tuning. The final *c*- and *b*-hadron decays are performed using EVTGEN v1.2.0.

Process	Generator	Accuracy in QCD	PDF set	PYTHIA 8 tune
ggF	POWHEG-BOX v2 + MINLO + NNLOPS	NNLO in y^H	PDF4LHC NNLO	AZNLO
VBF	POWHEG-BOX v2	NLO	PDF4LHC NLO	AZNLO
<i>VH</i>	POWHEG-BOX v2 + MINLO	NLO	PDF4LHC NLO	AZNLO
<i>ttH</i>	POWHEG-BOX v2	NLO	PDF4LHC NLO	A14
<i>bbH</i>	MADGRAPH5_aMC@NLO 2.3.3	NLO	CT10 NLO	AZNLO
<i>tH</i>	MADGRAPH5_aMC@NLO 2.6.0	NLO	NNPDF 3.0	AZNLO

The ggF, VBF, *VH* and *ttH* are simulated using the POWHEG-BOX v2 generator [231–235]. The associated PDF is chosen to be the LHC-targeted PDF provided by PDF4LHC [236], using the NNLO set for ggF, and the NLO set for VBF, *VH* and *ttH*. The ggF generation accuracy reaches NNLO in QCD thanks to two additional correction procedures: first, the jet merging with the parton shower using POWHEG and MINLO [225] procedure, and second, a reweighting proce-

dured according to the Higgs boson rapidity distribution using the HNNLO [237, 238] programme implementing the NNLOPS [239] procedure.

The matrix elements for the VBF, $qq \rightarrow VH$ and ttH processes are generated at NLO. In addition, the $qq \rightarrow VH$ benefits from the MINLO [225] method to merge events with additional 0 or 1 jets. On the other hand, the small contribution from the $gg \rightarrow ZH$ process is modelled at LO only.

Finally, turning to the small additional contributions from associated bbH and tH production, the bbH process is generated at NLO [240] with the MADGRAPH5_aMC@NLO 2.3.3 [240] software, using the CT10 NLO PDF set [241]. The $tH + jb$ and $tH + W$ (commonly referred to as tH) are generated with MADGRAPH5_aMC@NLO 2.6.0, using the NNPDF 3.0 PDF set [242].

Once the parton level events have been generated, the Higgs boson decay and the jet QCD showering is performed using the PYTHIA 8 generator [221] with the AZNLO PDF set [243] (except for ttH using the A14 PDF set [244]). The last step before passing the events to the ATLAS simulation is to simulate the c - and b -hadron decays, which is performed with EVTGEN v1.2.0 [230].

The simulation assumes a Higgs boson mass of $m_H = 125 \text{ GeV}$, and the cross-sections are scaled to their best-known values, described in Chapter 7.

8.2.4 Simulated background samples

The backgrounds of the $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ decay channel are small and cannot be easily estimated from data-driven only methods. Their estimation therefore relies on the simulation of several processes to have an as complete as possible description of the backgrounds.

These include the dominant *non-resonant* background from double Z boson production ($ZZ^{(*)}$ through the gg and $q\bar{q}$ channels) and smaller *reducible* backgrounds due to misidentified or non-prompt objects from WZ , triple-boson production (ZZZ , WZZ and WWZ commonly referred to as VVV), double top production ($t\bar{t}$), the Z boson production in association with jets ($Z + \text{jets}$), and the top-associated production ($t\bar{t}Z$ and the smaller tWZ , $t\bar{t}W^+W^-$, $t\bar{t}t$, $t\bar{t}\bar{t}$, tZ , commonly referred to as tXX). Given the variety of samples using various generators and settings, the details concerning the generation of each background are not recalled in this Thesis but can be found in Refs. [245] and [246].

The normalisation of the dominant contributions (for the non-resonant, and for the $t\bar{t}$ and $Z + \text{jets}$ contributions to the reducible background) are estimated using the data-driven methods detailed in Section 8.4. The other smaller contributions, for which eventual simulation mis-modelling would only have a tiny impact on the yield estimates, are directly taken from MC simulation.

8.3 Event selection

The $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ event selection presented in this Section focuses on the low mass analysis. The high mass analysis ($m_{4\ell} > 190 \text{ GeV}$) uses slightly different mass cuts on the lepton pairs and FSR-recovery considerations due to the two Z being on-shell, but will not be described here.

8.3.1 Trigger requirements

The triggers used in the $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ selection correspond to the lowest unprescaled² triggers with one, two or three leptons. Due to increasing peak luminosity across the Run 2 data taking, the lowest energy thresholds changed from 20 GeV to 26 GeV for the single muon triggers and from 24 GeV to 26 GeV for the single electron triggers [247–249]. Other single lepton triggers with looser isolation requirements but higher thresholds are additionally used. Double and triple lepton trigger with lower energy thresholds for each of the leptons provide additional efficiency in the low p_T regime. The overall trigger efficiency for the $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ selection is around 98 %.

8.3.2 Physics objects selection

The main limitation of the $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ analysis is its low rate. However, thanks to a small background in the signal region, the acceptance can be increased as much as possible without harming the sensitivity. The object selection reflects this need to maximise the acceptance by using loose reconstruction and identification criteria, and by imposing the lowest cuts on energy or momentum provided by the reconstruction and identification performance.

The selection starts with events containing at least four candidate leptons that can be grouped into two same-flavour opposite-sign lepton pairs. In case more than four leptons are present, more than one lepton *quadruplet* can be formed, and the final quadruplet selection is described below in Sections 8.3.3 and 8.3.4.

Electrons are required to pass loose likelihood criteria as defined in Chapter 4.4, leading to a combined reconstruction and identification efficiency of $\sim 95\%$. After calibration, they are kinematically required to have a $p_T > 7\text{ GeV}$ and $|\eta| < 2.5$. Note that the “crack” region between 1.37 to 1.52 is included in the analysis despite a less precise calibration to increase the analysis acceptance.

Depending on the η coverage of the ID and the MS, muons are required to pass different reconstruction criteria (see Chapter 4.5 for details). In the range $0.1 < |\eta| < 2.5$, muons are required to pass the nominal reconstruction algorithm using combined information from the ID and the MS. Due to the partial coverage of the muon stations in the range $|\eta| < 0.1$, muons are allowed to be segment-tagged (full ID + partial MS track) or calorimeter-tagged (full ID + MIP-like calorimeter deposit). In the region $|\eta| > 2.5$ where the ID cannot provide a fully reconstructed track, muons are allowed to have standalone reconstruction (requiring only three MS stations) or be associated with silicon hits (ID tracklet). At most one calorimeter-tagged or standalone or silicon-associated muon is allowed per quadruplet. The p_T requirement on the reconstructed muons is $p_T > 5\text{ GeV}$, except for the calorimeter-tagged muons reaching $p_T > 15\text{ GeV}$.

Jets reconstructed from topological calorimeter clusters with the anti- k_t algorithm of parameter 0.4 as described in Chapter 4.6 are pre-selected. After calibration, they are required to have $p_T > 30\text{ GeV}$ and $|\eta| < 4.5$. The pileup jet background is further suppressed using the algorithms described in Chapter 4.6. These selected jets are finally evaluated using the MV2_c10 algorithm mentioned in Chapter 4.6 to assign them with a b -tagging weight.

Before moving to the quadruplet selection, objects undergo an overlap removal procedure. First, if two electron clusters overlap, only the electron with higher E_T is kept. Second, if an electron and a muon share the same ID track, the muon is rejected in case of a calorimeter-tagged muon, otherwise the electron is rejected. Third, jets within $\Delta R = 0.2$ of electrons or 0.1 of muons

²↑ A prescaled trigger is one that is accepted only a fixed fraction of the time so as to not overflow the readout system. The successful unprescaled triggers are guaranteed to be passed to the rest of the acquisition system.

are removed.

Events are required to have at least one vertex with at least two tracks, where the vertices from the collision are built from ID track with $p_T > 0.5$ GeV. The primary vertex is then defined as the vertex with largest sum of p_T in the event. Since it is not expected that the collisions yield more than one hard scatter process per bunch crossing, the primary vertex is assumed to be the one where the Higgs boson decayed. Furthermore, as neither the Higgs nor the Z bosons live long enough to produce displaced vertices, the tracks of all leptons should emerge not farther than $|z_0 \sin(\theta)| < 0.5$ mm from the primary vertex along the beam direction. Additionally, muons with impact parameter $d_0 > 1$ mm are rejected to avoid contamination from cosmic muons.

8.3.3 Lepton quadruplet building

At this stage, all the physics objects (electrons, muons, jets, and the missing transverse energy E_T^{miss} not mentioned above) have been selected and are ready to undergo the lepton pairing and event categorisation procedure.

The quadruplets are required to have the three first p_T -ordered leptons with a p_T greater than 20, 15, and 10 GeV respectively. In each quadruplet, the lepton pair with invariant mass closest to the PDG Z boson mass is the *leading* pair, and the other is the *subleading*. The quadruplets can then be labelled according to their final-state composition: 4μ , $2e2\mu$, $2\mu2e$, $4e$, where the first lepton pair corresponds to the leading pair.

The leading lepton pair is required to have a mass comprised between $50 \text{ GeV} < m_{12} < 106 \text{ GeV}$, while the subleading pair invariant mass must fall in the range $12 \text{ GeV} < m_{34} < 115 \text{ GeV}$. In the $4e$ and 4μ channels, the quadruplets are removed if the alternative cross-pairing yield a pair with invariant mass below 5 GeV to suppress the J/ψ background.

Within a quadruplet, all leptons must be geometrically separated by at least $\Delta R > 0.1$. In order to reject the heavy flavour background whose hadrons have a lifetime long enough to produce (slightly) displaced vertices, an additional cut on the lepton impact parameter significance is applied. Muons are required to have $d_0/\sigma_{d_0} < 3$ while electrons, due to poorer track resolution from bremsstrahlung, are required to have $d_0/\sigma_{d_0} < 5$. Similarly, the Z + jets and $t\bar{t}$ backgrounds are further suppressed with a vertex cut of the four leptons. A requirement on the fit quality is placed with 99.5 % signal efficiency, requiring $\chi^2/N_{\text{d.o.f.}} < 6$ for the 4μ quadruplets, and $\chi^2/N_{\text{d.o.f.}} < 9$ for the other types ($2e2\mu$, $2\mu2e$, $4e$) to account for the lower vertex resolution of electrons. This results in a 20 % to 30 % rejection of the Z + jets and $t\bar{t}$ events, as seen in Figure 8.3.

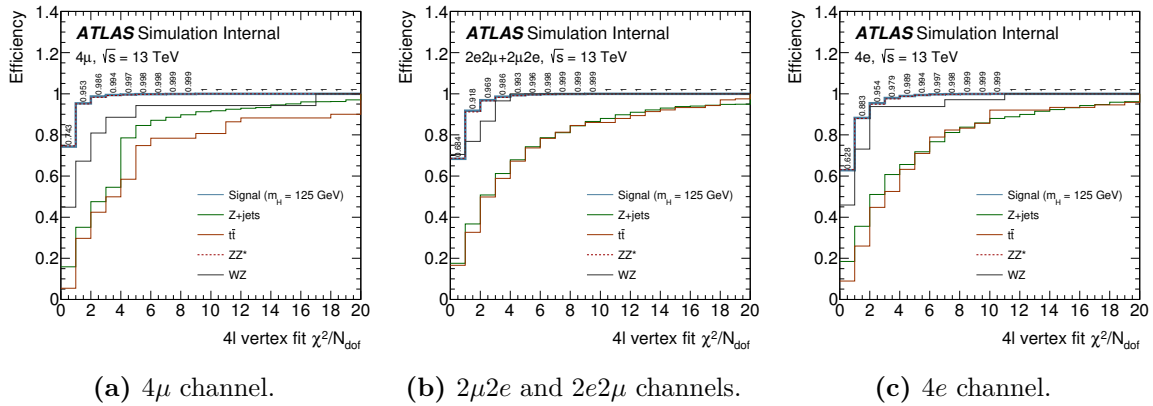


Figure 8.3 – Selection efficiency for the signal, irreducible and reducible background as a function on the cut applied on the vertex fit χ^2 in the (a) 4μ channel, (b) $2\mu2e$ and $2e2\mu$ channels, (c) $4e$ channel. The cut is chosen to keep a 99.5 % signal efficiency, leading to a 20 % to 30 % reducible background rejection. [246]

Further track and calorimeter isolation requirements are applied to both electrons and muons to reject jets misidentified as leptons. The track isolation is defined as the sum of the tracks with $|p_T| > 0.5 \text{ GeV}$ and within $|z_0 \sin(\theta)| < 3 \text{ mm}$ of the primary vertex not associated with any other vertex. The tracks must additionally lie in a cone around the electron or muon track whose size varies from $\Delta R < 0.3$ for $p_T < 33 \text{ GeV}$ to $\Delta R < 0.2$ for $p_T > 50 \text{ GeV}$ linearly between the two thresholds. The calorimeter isolation is defined as the E_T sum of the positive energy clusters which fall within $\Delta R < 0.2$ from the electron or muon and are not associated to a track. The calorimeter isolation is corrected for electron shower leakage, pileup and underlying event contributions. Both the track and the calorimeter isolations are corrected for the contributions of the other three leptons. The leptons pass the isolation criteria if the sum of their track and calorimeter isolation fulfils $\text{Iso}_{\text{track}} + 0.4 \times \text{Iso}_{\text{calo}} < 0.16 \times p_T$, achieving a signal efficiency of 80 %.

As final-state electrons and muons can radiate a photon in the final state, missing out the FSR degrades the resolution on the four-lepton mass. To recover a part of this resolution, collinear FSR photons are considered and it is allowed to add at most one FSR photon to the event. Since for electrons the collinear FSR should already be taken into account thanks to the reconstruction dynamical clustering method, only muons are candidates for this FSR-recovery. Only muons from the leading Z boson if its mass is $m_{12} < 89 \text{ GeV}$ can be candidate to avoid selecting initial-state radiations. The photon FSR candidate are considered if their transverse energy is $E_T > 1$ (resp. 3.5) GeV , their energy deposited in the first sampling of the calorimeter is $f_1 > 0.2$ (resp. 0.1), and the distance between the cluster and the muon is $\Delta R_{e,\mu} < 0.08$ (resp. 0.15). In case several FSR candidates are found, only the one with highest transverse energy is kept. If after adding the FSR photon to the leading Z boson its mass exceeds 100 GeV , the FSR is eventually rejected. The FSR recovery affects about 3 % of the events. The impact on corrected events is shown in Figure 8.4.

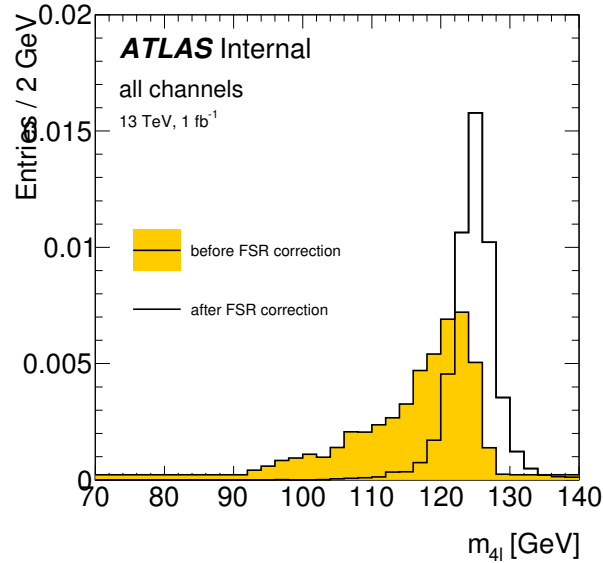


Figure 8.4 – Impact of the FSR recovery on the four-lepton invariant mass for events benefiting from FSR recovery. 3 % of events are impacted. [246]

8.3.4 Lepton quadruplet selection

At this point, all possible quadruplets have been formed. There might be several final-state possibilities or even several quadruplets candidates for a single final state. This is naturally the case in the $4e$ and 4μ decay channels, but can also be due to additional leptons in case of $t\bar{t}H$ production with top quarks decaying leptonically or VH -lep topologies. This section describes

the procedure for ranking the quadruplets and finally select one Higgs boson candidate per event.

If after the quadruplet building, the event contains at least one additional lepton fulfilling the same selection, identification, isolation and separation with the others leptons as the standard selection, additionally requiring its transverse momentum be $p_T > 12 \text{ GeV}$, then lepton pairing might be ambiguous. For such events, the ambiguity is resolved, if needed for each final state separately, using a matrix element pairing: for each ambiguous quadruplet, the matrix element for the decay process $H \rightarrow 4\ell$ is computed at leading order using MADGRAPH5_aMC@NLO [250]. The quadruplet with highest matrix element is chosen.

After this step, if several quadruplets remain in each decay channel, priority is given to the quadruplet with m_{12} closest to m_Z leading to at most one quadruplet per decay channel. If several channels yield one quadruplet, the one with highest signal efficiency is retained: the selection efficiencies reach 31 %, 21 %, 17 % and 16 % for the 4μ , $2e2\mu$, $2\mu2e$ and $4e$ decay channels, respectively.

Finally, the signal region (SR) includes events passing the selection and falling in the mass range $115 \text{ GeV} < m_{4\ell} < 130 \text{ GeV}$. The side band region (SB) follows the same selection but the mass range is selected as $105 \text{ GeV} < m_{4\ell} < 160 \text{ GeV}$ with the range $115 \text{ GeV} < m_{4\ell} < 130 \text{ GeV}$ excluded. Finally the control regions (CR) for the background estimations use a slightly different event selections which are described in Sections 8.4.2 and 8.4.4.

8.4 Backgrounds estimation

The backgrounds to the $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ analysis are split in two types: the processes that yield four prompt leptons, called *non-resonant backgrounds*, and the processes whose final state contains at least one jet, photon or lepton stemming from a hadron decay that can be misidentified as prompt leptons, called *reducible backgrounds*. The former are described in Section 8.4.1, the latter in Sections 8.4.2 and 8.4.4 depending on what final state they impact most.

8.4.1 Non-resonant background

The non-resonant background is dominated by the Standard Model $pp \rightarrow ZZ^{(*)}$ production. Due to the di-lepton invariant mass cuts, the main contributions come from the Z boson, so this background is abridged $ZZ^{(*)}$ in the following.

The dominant process is the quark radiation as pictured in Figure 8.5a, giving a continuous background over the whole mass range. The quark annihilation shown in Figure 8.5b yields a resonant (peak) contribution at $m_{4\ell} \approx m_Z$ which does not contribute much in the $115 \text{ GeV} < m_{4\ell} < 130 \text{ GeV}$ region, but can be used to perform resolution or selection validation studies. Since the $gg \rightarrow ZZ^{(*)}$ production as seen in Figure 8.5c only occurs at loop level, its contribution is suppressed compared to the $qq \rightarrow ZZ^{(*)}$ production. The shape of $ZZ^{(*)}$ background is provided by the simulation, but its normalisation is adjusted from the side band region around $115 \text{ GeV} < m_{4\ell} < 130 \text{ GeV}$.

An additional subleading contribution comes from the tri-boson production (VVV) with at least 4 leptons in the final state ($ZZZ \rightarrow 4$ or 6 leptons, $WZZ \rightarrow 5$ leptons, $WWZ \rightarrow 4$ leptons) and the $t\bar{t}+V$ production with the t -quark decaying leptonically. The shape and normalisation of these subleading contributions are estimated from the simulation.

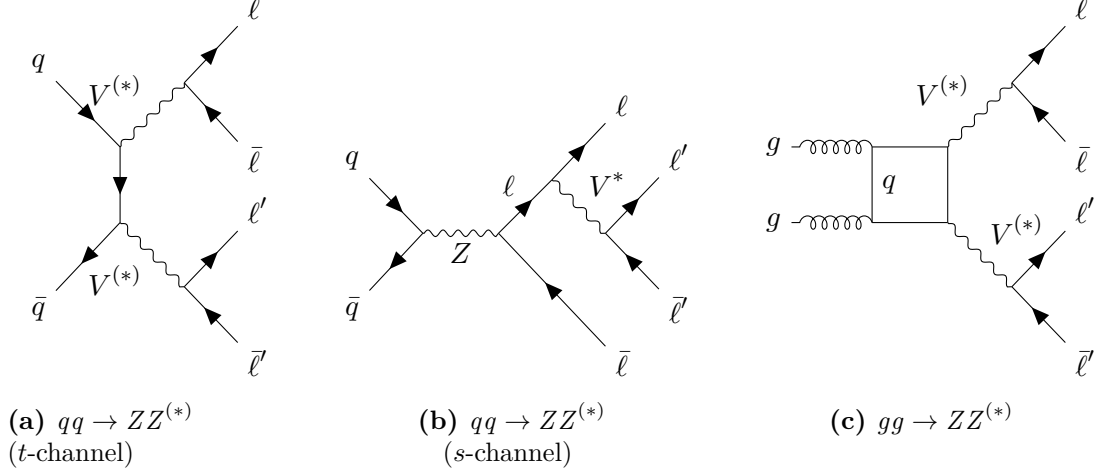


Figure 8.5 – Feynman diagrams corresponding to the dominant background topologies for the $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ decay channel. V corresponds to either γ or Z .

8.4.2 Reducible $\ell\ell + \mu\mu$ background

The muon backgrounds mostly come from non-prompt muons created in the semi-leptonic decays of heavy flavoured hadrons formed in the $t\bar{t}$ of Z + jets production. An additional smaller contribution from light hadron decays is also taken into account.

Four control regions (CR) are defined by relaxing or removing cuts, ensuring their orthogonality to the signal region and between them. Only one quadruplet is selected per event but due to non-standard pairing, the ME ambiguity resolver cannot be applied. This was however checked to not bias the shape of the backgrounds.

Inverted d_0/σ_{d_0} CR (heavy flavour (HF) jets and $t\bar{t}$ enriched): in this CR, at least one of the leptons in the subleading pair is required to have its impact parameter significance inverted ($d_0/\sigma_{d_0} > 3$ (resp. 5) mm for a muon (resp. electron) pair). The vertex and isolation requirements are also removed but the rest of the selection remains the same, leading to a region enriched in heavy flavour jets (Z +HF and $t\bar{t}$) due to their displaced vertex.

Inverted isolation CR (light flavour (LF) enriched): similarly to the HF region above, inverting the isolation requirement for at least one of the two subleading leptons enriches the region in light hadron (mainly charged pion and kaon) decays. The rest of the standard selection is applied, including the d_0 and vertex requirements to reject HF contributions.

$e\mu + \mu\mu$ CR ($t\bar{t}$ enriched): by requiring the leading lepton pair to be of opposite sign but different flavour, the pair is ensured to not stem from a Z boson decay but most likely from the two top quarks. The subleading pair is allowed to be of opposite or same sign to increase the statistics. The selection is otherwise similarly applied, with the exception of the impact parameter and isolation requirements.

Same-sign (SS) CR: this CR does not target a specific contribution and includes significant contributions from both HF and LF, and $t\bar{t}$ to help constraining the LF part. It is defined by requiring the subleading lepton pair to be of same flavour but also of same sign. The rest of the selection is applied normally, with the exception of the d_0 -significance and the isolation criteria.

Relaxed VR: an additional “relaxed” validation region (VR) is defined with all the standard cuts but the isolation, d_0 -significance and vertex fit. This VR is not orthogonal to the SR nor the different CR and is therefore not used in the background estimations, but rather forms a superset of the SR used to derive the transfer factors to the SR and to validate the CR estimates.

The idea of this background estimation is to fit the data simultaneously in the four CR using constrained shapes from the MC for each background component (Z +HF, Z +LF, $t\bar{t}$). The transfer factors of the yields from the CR to the SR are estimated using the MC simulation. The observable is chosen to be m_{12} as its shape for the Z + jets contribution (resonant at the Z boson peak) and the $t\bar{t}$ contribution (smooth over all of the m_{12} range) allows for a correct discrimination. In order to have a better handle on the Z +LF normalisation relative to the Z +HF, the inverted isolation and same-sign CR are further divided in two subregions with respect to the p_T balance between the ID and MS: a LF-enriched region for pion and kaon decay in flight can be built by requiring $(p_T^{\text{ID}} - p_T^{\text{MS}})/p_T^{\text{ID}} > 0.2$ for at least one of the two muons of the subleading pair.

The first step is to build the shapes for the $t\bar{t}$, Z + jets and other small contributions (residual $ZZ^{(*)}$, WZ , $t\bar{t}+V$, VVV and Higgs boson all merged together under “diboson”), which are chosen to be analytical. The $t\bar{t}$ contribution is modelled by a second-order polynomial, shared across all four CR. The Z + jets contribution is modelled by a Breit-Wigner convolved with a Crystal-Ball for the inverted d_0 -significance, inverted isolation and same sign CR, but since the $e\mu + \mu\mu$ CR is purposely non-resonant for m_{12} , it cannot share the same PDF and its shape is rather modelled with a first order polynomial. The diboson has a similar shape to the Z + jets, but with a larger tail. The relative fraction between the two PDF is taken as the ratio of events with m_{12} below or above 75 GeV. These shapes are individually fitted to the MC distributions in each CR and for each background type. The analytical models used to define the PDF in the final fit are summarised in Table 8.2, and the MC shapes used to fix their parameters are shown in Figure 8.6. The data points overlaid on this figure show a large difference in normalisation, but similar shape. This demonstrates the need of a data-driven estimate of the background yields.

Table 8.2 – Modelling of each contribution entering the $\ell\ell + \mu\mu$ background estimation. The yields of all WZ sub-contributions are taken directly from the CR in MC.

Bkg. type	Inverted d_0/σ_{d_0}	Inverted isolation	Same-sign	$e\mu + \mu\mu$
$t\bar{t}$	Second order polynomial			
Z + jets	BW $\times$ CB			First order pol.
WZ	$f \cdot G(; \mu = 60, \sigma = 10) + (1 - f) \cdot \text{BW} \times \text{CB}$ (from relaxed VR in MC)			

Once the parameters fixing the shapes of the various background contributions in each CR are fit to the MC, the Z +HF, Z +LF and $t\bar{t}$ normalisations are expressed in terms of the number of events N_i in the relaxed VR, and the total PDF in each CR is defined by

$$\text{Model}_{\text{CR}} = \sum_i N_i \cdot f_{i,\text{CR}} \cdot \text{AnaShape}_{i,\text{CR}}, \quad (8.1)$$

where i runs over each background type, $f_{i,\text{CR}} = N_i^{\text{CR}}/N_i^{\text{VR}}$ are defined as the fraction of events corresponding to the background i in the given CR divided by the fraction of events from background type i in the relaxed VR (evaluated from MC), and the $\text{AnaShape}_{i,\text{CR}}$ is the analytical shapes defined above. The parameters of interest of the simultaneous fit are the N_i . The method is validated by a closure test performed on the MC, which shows a good agreement as demonstrated in Table 8.3.

During the fit to the data CR, the parameters governing the shape parameters are constrained to their best MC fit within their uncertainties ($\pm 3\sigma$). Similarly, the transfer factors $f_{i,\text{CR}}$ are constrained to their MC value within their uncertainties ($\pm 3\sigma$ including systematics from $Z + \mu$

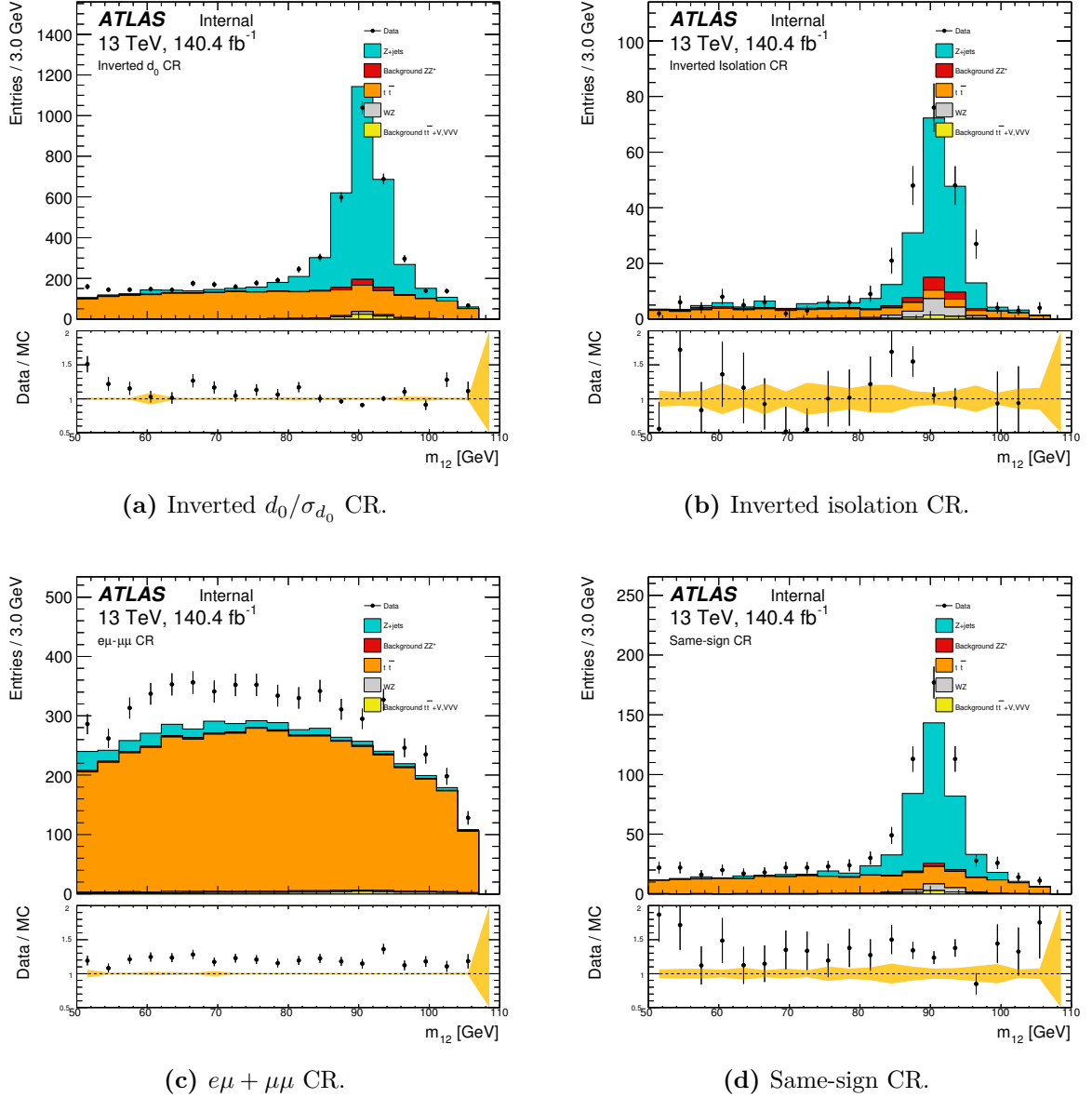


Figure 8.6 – m_{12} distributions generated from MC-simulated events for different processes and in the various $\ell\ell + \mu\mu$ control regions: (a) inverted d_0/σ_{d_0} , (b) inverted isolation, (c) $e\mu + \mu\mu$ and (d) same-sign. These distributions are used to define the PDFs for the background estimate simultaneous fit. The full Run 2 data in the same CR are overlaid, the differences indicating the need for a data-driven estimate instead of relying on pure MC predictions. [246]

Table 8.3 – Results of the MC closure test for the $\ell\ell + \mu\mu$ background modelling. The number are the yields expressed in the MC relaxed VR for each background component and the respective estimation from the fit (diboson here incorporates signal from Higgs boson and the minor $t\bar{t}V$, VVV , and is MC fixed). A good agreement is observed. [246]

Sample	MC Events	Fit Result
$t\bar{t}$	2586 ± 9	2590 ± 42
Z+HF	3300 ± 29	3306 ± 102
Z+LF	142 ± 23	133 ± 37
Diboson	3365 ± 5	3365 ± 5

described in Section 8.4.3 as well). The resulting simultaneous fit over each data CR is shown in Figure 8.7. A comparison of the expected background composition in the relaxed VR after the simultaneous CR fit in data is shown in Figure 8.8 and exhibits a good agreement.

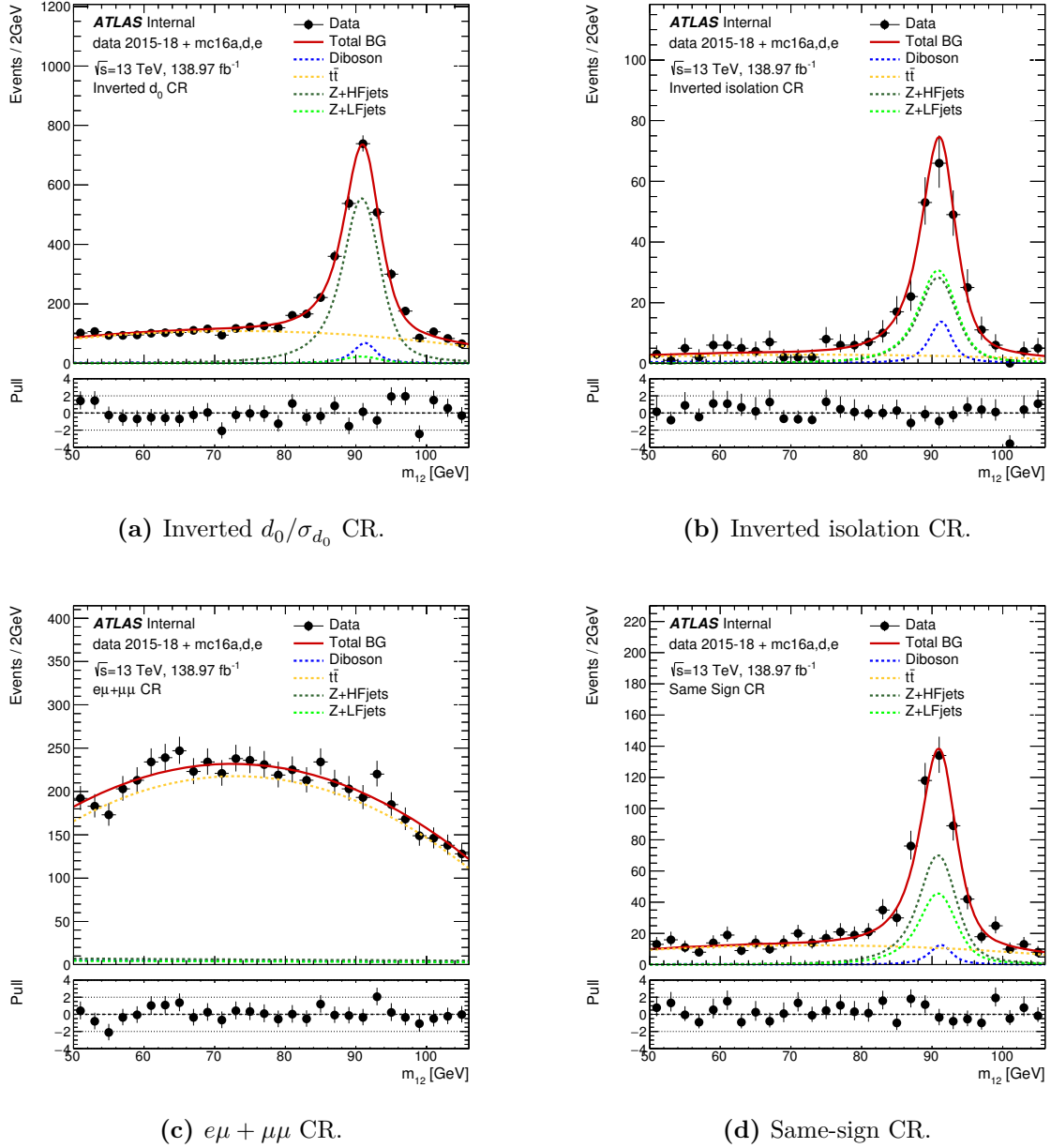


Figure 8.7 – Background component simultaneous fit to the full Run 2 data using m_{12} as observable in each $\ell\ell + \mu\mu$ control region: (a) inverted d_0/σ_{d_0} , (b) inverted isolation, (c) $e\mu + \mu\mu$ and (d) same-sign. The lower panels show the fit pulls. [246]

The final yields for each background component in the SR (mass range [115, 130] GeV) are deduced from the yields in the relaxed VR (full mass range) using transfer factors estimated from the MC for each background component. The $t\bar{t}$ and HF Z + jets transfer factor are computed as the event yield in the SR divided by the event yield in the relaxed VR. The LF Z + jets transfer factor is computed from the isolation, the d_0 -significance, and the vertexing cut efficiencies, which are derived from the $Z + \mu$ sample (Section 8.4.3). The WZ sample is directly estimated from the MC in the signal region and therefore does not need a transfer factor.

The yields expressed in the data VR, the transfer factors from MC and the extrapolated results to the data SR are listed in Table 8.4. The values for the exclusive channels 4μ and $2e2\mu$

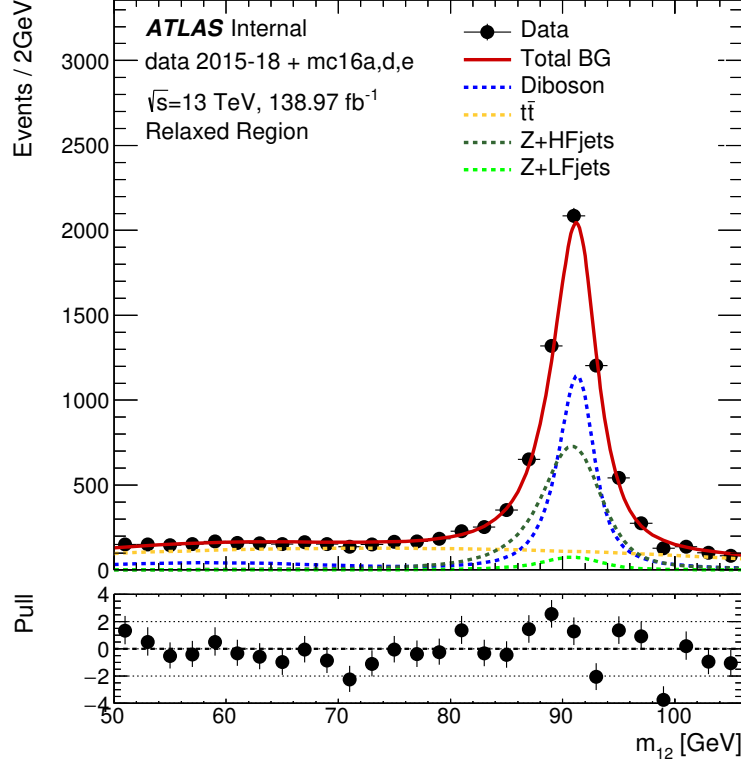


Figure 8.8 – Distribution of m_{12} in the relaxed isolation and d_0 -significance validation region of data, compared to the background yields extrapolated from the data control regions. [246]

are derived from the inclusive $\ell\ell + \mu\mu$ by a simple fraction taken from the SR in the MC for each background component. Uncertainties on the SR yield arise from the limited statistics of the samples used to determine the transfer factors, and from additional systematics assigned to the background muon selection efficiencies which are described in Section 8.4.3. The systematic uncertainties are derived separately for each channel ($\ell\ell + \mu\mu$, 4μ and $2e2\mu$). The procedure described here determines the background yields, the background shape determination is described in Section 8.4.6.

Table 8.4 – Final $\ell\ell + \mu\mu$ background estimates in the “relaxed” VR in the full $m_{4\ell}$ range for each background component. The second column shows the extrapolation factors to the SR and their corresponding statistical uncertainties. The last column shows the estimates for the SR yields with both statistical and systematic uncertainties. [246]

Bkg type	Data fit	TF [%]	SR yield		
			$4\mu + 2e2\mu$	4μ	$2e2\mu$
$t\bar{t}$	3118 ± 47	0.25 ± 0.03	$7.65 \pm 0.12 \pm 0.84$	$2.06 \pm 0.04 \pm 0.30$	$5.59 \pm 0.09 \pm 0.65$
$Z + \text{jets (HF)}$	2905 ± 113	0.43 ± 0.04	$12.47 \pm 0.49 \pm 1.17$	$7.12 \pm 0.37 \pm 0.69$	$5.35 \pm 0.26 \pm 0.53$
$Z + \text{jets (LF)}$	299 ± 68	1.00 ± 0.10	$3.10 \pm 0.70 \pm 0.28$	$1.75 \pm 0.49 \pm 0.30$	$1.35 \pm 0.39 \pm 0.30$
$Z + \text{jets (HF+LF)}$	3204 ± 132		$15.57 \pm 0.85 \pm 1.45$	$8.87 \pm 0.61 \pm 0.99$	$2.21 \pm 0.47 \pm 0.83$
WZ	MC-based estimation		4.52 ± 0.32	2.13 ± 0.50	2.39 ± 0.51

The method is validated using two additional methods: the first simply merges the Z+HF and Z+LF together, which ensures that the single Z+LF and Z+HF are not biased by the $t\bar{t}$ single contribution. The second starts with fitting only two CR to extract the Z+HF and $t\bar{t}$ components, and then fit another CR adding the Z+LF in order to extract this last component. The results from the nominal and two cross-check methods are shown in Table 8.5 and display a good agreement.

Table 8.5 – Comparison of the $\ell\ell + \mu\mu$ background estimation in the signal region obtained from the three methods. [246]

Method	$Z+\text{HF}$	$Z+\text{LF}$	$t\bar{t}$	WZ
Standard	$12.47 \pm 0.4 \pm 1.17$	$3.10 \pm 0.70 \pm 0.28$	$7.65 \pm 0.12 \pm 0.84$	4.53 ± 0.32
Merged $Z+\text{jets}$	$14.44 \pm 0.32 \pm 1.35$		$7.66 \pm 0.10 \pm 0.85$	
2CR+1	$12.74 \pm 0.38 \pm 1.19$	$2.91 \pm 0.63 \pm 0.29$	$7.74 \pm 0.10 \pm 0.84$	

8.4.3 $\ell\ell + \mu\mu$ transfer factor systematics uncertainties

The main systematic uncertainties in the background yields arise from the uncertainty on the transfer factors, *i.e.* the background muon selection efficiencies. These are studied using samples containing one on-shell Z boson plus one additional muon, denoted $Z+\mu$.

The selection for this sample is close to the nominal one, dropping the requirements on the subleading Z boson. The tri-lepton and $e\mu$ triggers are also dropped to avoid any bias on the additional muon quality. The Z boson is built from an opposite pair of electrons or muons passing the nominal selection (d_0 -significance and isolation included), and it is required that the two leptons have a p_T greater than 20 GeV and 15 GeV respectively to follow the nominal selection. The invariant mass of the pair must be in the range $76 \text{ GeV} < m_{\ell\ell} < 106 \text{ GeV}$, retaining the one with mass closest to the mass of the Z boson in case several pairs are possible. Only events with exactly one additional muon with $p_T > 5 \text{ GeV}$ are kept, and further requiring that the three leptons are well separated with $\Delta R_{\ell,\ell} > 0.1$. The event is however killed if any alternative pairing yields an invariant mass lower than 5 GeV, to follow the J/ψ cut of the standard analysis.

Such requirements yield a sample enriched in $Z+\text{jets}$, and the efficiencies of the isolation and d_0 -significance cuts (number of events passing these selections divided by the total number of events) are compared in data and MC, as shown in Table 8.6. The small contributions from WZ are removed from both data and MC by subtracting the MC distribution for such events. The impact parameter cut gives consistent efficiencies for both data and MC, but the isolation cut shows a 5 % to 10 % difference. To understand this discrepancy, the HF and LF are studied separately.

Table 8.6 – Efficiency of isolation and impact parameter selections for background muons selected in $Z+\mu$ events. Data-MC efficiency differences divided by MC efficiency are also shown. [246]

Selection applied	Data [%]	MC [%]	(Data - MC)/MC
Isolation	15.4 ± 0.03	14.7 ± 0.1	5%
d_0 -significance	64.40 ± 0.04	64.1 ± 0.1	0.4%
Iso. + d_0 -sig.	10.09 ± 0.03	9.1 ± 0.1	11%

In order to separate the effects on the LF and HF contributions, a LF enriched sample is built by requiring the p_T imbalance to be greater than 0.1 (Figure 8.9). From truth level information, this sample is known to be 67%-pure in LF and contains $\sim 32\%$ of HF ($Z+\text{jets}$ and $t\bar{t}$), the remaining 1 % being composed of WZ and $ZZ^{(*)}$ events. The HF, WZ and $ZZ^{(*)}$ contributions are subtracted from both data and MC using the MC prediction. The isolation and d_0 -significance efficiencies are computed in this data sample, while the vertexing efficiency is taken from the HF-enriched $Z+\mu$ below. The cut efficiencies in the data and in the simulations are shown in Table 8.7. The systematics uncertainties on the LF efficiencies are derived by varying the fraction of HF and $t\bar{t}$ subtracted by a factor of 2, by varying the fraction of WZ subtracted by a factor of 1.5, and by varying the value of the p_T -imbalance cut used to define the LF-enriched region by a factor of 2.

The HF-enriched control sample is built by inverting the d_0/σ_{d_0} cut ($d_0/\sigma_{d_0} > 3$) on the

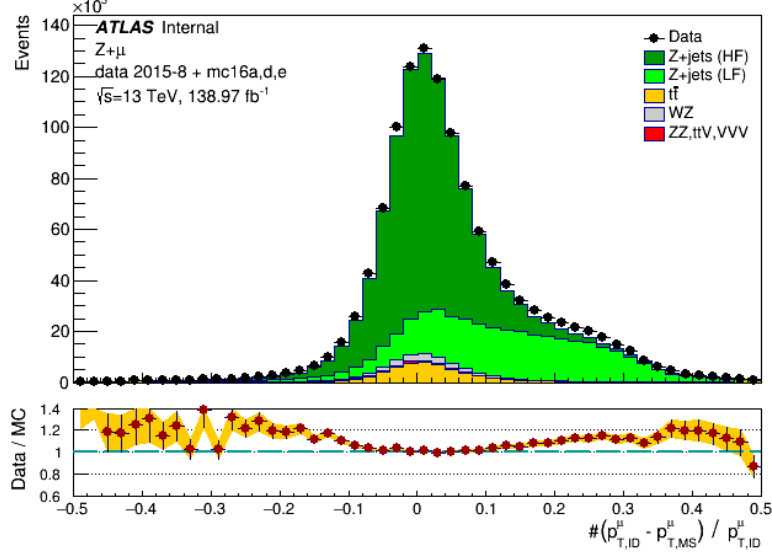


Figure 8.9 – Fractional p_T balance between the ID and MS measurements for additional muons in $Z+\mu$ events shown for data and MC simulation. [246]

Table 8.7 – Isolation and impact-parameter efficiencies for background muons in the light-flavour enriched and heavy-flavour enriched $Z+\mu$ samples, after subtraction of the remaining contributions (WZ , HF and $t\bar{t}$ in the LF sample, WZ , LF and $t\bar{t}$ in the HF sample). Since the HF-enriched region is defined using the d_0 significance cut, only isolation efficiency is computed. [246]

Cut	Data	MC	(Data - MC)/Data
Isolation (LF)	8.0 ± 0.2	11.66 ± 0.01	-45%
d_0 significance (LF)	92.4 ± 0.2	90.7 ± 0.1	1.1%
Isolation + d_0 significance (LF)	7.4 ± 0.2	10.9 ± 0.1	-32%
Isolation after d_0 sig. (LF)	8.1 ± 0.2	12.1 ± 0.1	-33%
Isolation (HF)	17.1 ± 0.1	16.3 ± 0.1	4.9%

additional muons of the $Z+\mu$ sample, as shown in Figure 8.10. From truth level information, this sample is known to contain about 93 % of HF and ~ 6 % of LF, the latter being subtracted from both data and MC using the MC predictions. The small additional WZ and $t\bar{t}$ contributions are also subtracted from data and MC using the MC prediction. The systematic uncertainty on the HF Z + jets transfer factor is taken as the difference of the isolation cut efficiency in the data and in the simulation (Table 8.7). As the d_0 -significance cut is used to define the HF-enriched region, no specific uncertainty can be assigned. However, as shown in Table 8.6, the difference of the d_0 -significance cut in data and simulation is small, no such systematic uncertainty is assigned.

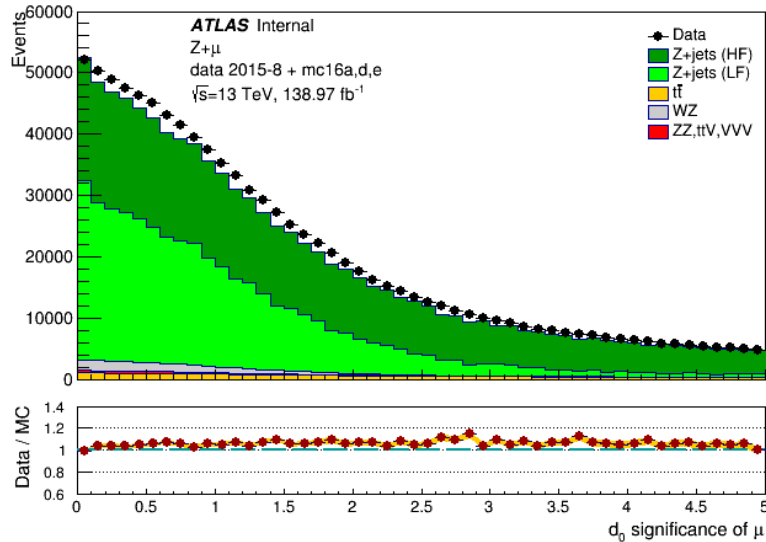


Figure 8.10 – Distribution of d_0 significance for additional muons in $Z+\mu$ events for data and MC simulation. [246]

The $t\bar{t}$ transfer factor systematic uncertainty is assumed to be the same as the HF Z + jets systematic uncertainty.

8.4.4 Reducible $\ell\ell + ee$ background

The electron background is mainly composed of light jets whose deposits in the calorimeter is misidentified as electrons clusters (denoted f). Other contributions include electrons from heavy-quark decays (denoted q) and from photon conversions (denoted γ).

This background is estimated using a dedicated control region where the standard selection is applied except for the lower- p_T lepton of the subleading pair, denoted X . This fourth lepton is only required to pass the number of silicon hits demanded for an electron, but its identification is otherwise not applied. The standard quadruplet selections are applied, including the vertex cut, except for the ME ambiguity resolver. In order to reduce the contribution from the $ZZ^{(*)}$ background, the subleading pair is further required to be composed of same sign (and still same flavour) leptons. However even after this requirement, about 10 % of events in this $3\ell+X$ CR originate from $ZZ^{(*)}$, and this contribution must be accounted for in the following. On the contrary to the $\ell\ell + \mu\mu$ above, all quadruplets sharing the same Z bosons are selected to estimate the contribution of each possible X to the SR. Thanks to the same-sign requirement, the $3\ell+X$ CR is ensured to be orthogonal to the SR.

Of the three main background components, the conversions γ can be distinguished from the q and f components by means of the n_{InnerPix} variable. This variable is defined as the number of hits in the IBL, or, if the electron crossed an IBL dead area, as the number of hits in the next-to-

innermost pixel layer. The q component from the semileptonic heavy-quark decays has the same shape as the f component, but amounts to about 1%. These two are therefore merged in the fit. Electrons from conversion stem from a photon and will therefore populate the $n_{\text{InnerPix}} = 0$ bin. The $ZZ^{(*)}$ component has a small impact from charge-misidentified electrons, which is subtracted in the data by taking the expected contribution from MC. The data distribution of n_{InnerPix} is fitted using templates built from MC. Since the $3\ell+X$ suffers from low statistics, the templates are rather built from the $Z+X$ CR described in Section 8.4.5. The result of the fit to data is shown in Figure 8.11. Each of the component yields in the CR are then estimated in bins of electron p_T and number of jets (N_{jet}) from this inclusive distribution using the $s\mathcal{P}lot$ method [251].

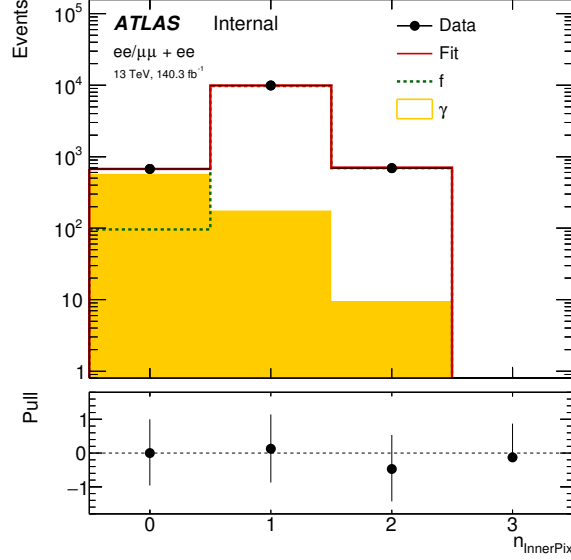


Figure 8.11 – Data fit of the n_{InnerPix} variable in the $3\ell+X$ CR, combining the $2\mu 2e$ and $4e$ decay channels. [246]

The transfer factor from the $3\ell+X$ CR to the SR are computed as the efficiency for a background electron to pass the standard selection (identification, isolation, and n_{InnerPix} requirements). These transfer factors are estimated separately for the f and γ components and in each bin of p_T and N_{jet} from the $Z+X$ CR in the MC described in Section 8.4.5. To better reproduce the efficiencies found in the data, data/MC scale factors are applied on the MC efficiency estimates. These scale factors are assigned a systematic uncertainty of about 23% for the f component, taking into account the MC modelling, the uncertainties on the heavy flavour contribution and the statistical limitations of the samples. For the γ component, even though it is less impacted by the heavy flavour contamination, a 20% uncertainty is assigned coming mainly from the MC modelling using different generators.

For each background component c , the final yield in the SR (N_{SR}^c) is obtained by summing the contributions in each of the p_T and N_{jet} bins weighted by their respective efficiency:

$$N_{\text{SR}}^c = \sum_i s_i^c \sum_j \epsilon_{ij}^c \cdot N_{ij}^{c,s\mathcal{P}lot}, \quad (8.2)$$

where i runs over the p_T bins, j runs over the N_{jet} bins, s_i^c is the data/MC efficiency scale factor in the p_T region i for the background component c , ϵ_{ij}^c is the efficiency for the background component c in the p_T bin i and N_{jet} bin j , and $N_{ij}^{c,s\mathcal{P}lot}$ is the yield in the p_T bin i and N_{jet} bin j for the background component c from the $s\mathcal{P}lot$ method.

Since the heavy flavour component is subtracted from the $3\ell+X$ CR, it is not extrapolated to the SR. Its contribution is instead taken from $t\bar{t}$ and Z +jets MC. The data/MC agreement of

this estimate is assessed in a specific CR defined as the $3\ell+X$ but applying $2.5 < d_0/\sigma_{d_0} < 5$ and removing the isolation criteria on the third lepton. A comparison of data and MC in this HF CR is shown in Figure 8.12 with good agreement: no further scaling is needed for this component.

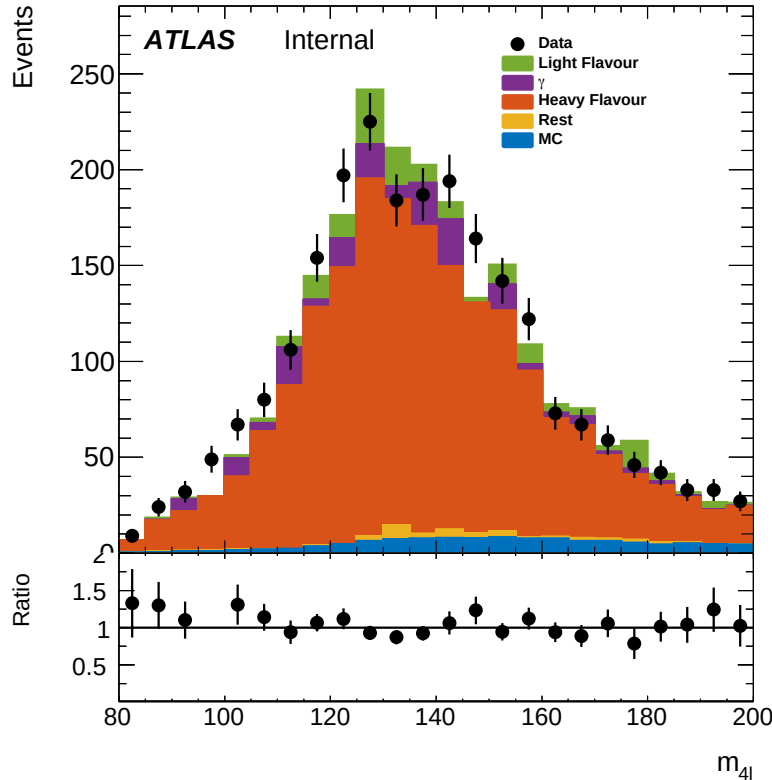


Figure 8.12 – Comparison of the FSR-corrected $m_{4\ell}$ distribution in the HF-enriched $3\ell+X$ CR between data and the various background contributions from MC. The “MC” category includes the $ZZ^{(*)}$, VVV , $t\bar{t}+V$ and residual Higgs boson contributions. The “Rest” category includes residual contributions from Z + jets where the third lepton is a prompt electron. [246]

The $t\bar{t}$ contribution to the HF component is taken as the $\ell\ell + \mu\mu$ estimate, scaled by the difference among the fit estimate of this component in the relaxed VR and the yield returned by the MC simulation also in the VR. The disagreement between the two is $\sim 20\%$ using the full Run 2 dataset, and the scale factor is therefore taken to be 1.2.

The final HF (Z + jets and $t\bar{t}$) contribution to the SR is obtained from the HF $3\ell+X$ CR after applying the isolation and impact parameter cuts to the subleading lepton pair, as well as the vertex fit requirement. An overall 30 % systematic uncertainty is assigned to this estimate to take into account possible variation of selection efficiency in the simulation.

The results of the template fit to data and the residual contributions, of the transfer factors and of the resulting yields in the SR are shown in Table 8.8, with their uncertainties. The separate estimate in the $4e$ and $2\mu 2e$ regions are performed by doing the procedure separately the two regions.

8.4.5 Extrapolation to the SR using the $Z+X$ CR

Due to the statistical limitations of the $3\ell+X$ sample, another sample with similar properties is constructed. A high statistics sample can be built from a Z boson sample with an additional electron candidate, called $Z+X$. The Z boson selection is the same as for the Z boson in the $Z+\mu$ samples (Section 8.4.3), and X candidate electron must follow the same requirements as the

Table 8.8 – Fit result for the various background component yields in the $3\ell+X$ CR (statistical errors only) using the full Run 2 dataset, shown together with the $ZZ^{(*)} + \text{HF}$ contamination and the efficiencies used to extrapolate the yields to the SR. The extrapolation are also shown separately for the $4e$ and $2\mu 2e$. The sum of the two contribution slightly differ from the total $\ell\ell + ee$ fit but the difference is covered by the uncertainties. The f and γ SR yields show both the statistical uncertainty from the data fit and the systematic uncertainty of the efficiency. The q component is not fitted to the data and its yield in the SR is taken directly from MC, with the $t\bar{t}$ and $Z + \text{jets}$ components scaled to match the data. Only the total uncertainty for this component is shown. [246]

Bkg. type	Data fit	$ZZ^{*} + \text{HF}$	Efficiency [%]		SR yield	
	$\ell\ell + ee$	$\ell\ell + ee$	$\ell\ell + ee$	$\ell\ell + ee$	$4e$	$2\mu 2e$
f	10503 ± 105	1342 ± 9	0.174 ± 0.05	$15.9 \pm 0.59 \pm 2.50$	$7.63 \pm 0.41 \pm 1.19$	$8.24 \pm 0.43 \pm 1.31$
γ	762 ± 34	121 ± 1.2	0.638 ± 0.03	$4.09 \pm 0.71 \pm 0.82$	$1.98 \pm 0.49 \pm 0.40$	$2.11 \pm 0.51 \pm 0.42$
q	(MC-based estimation)			12.72 ± 3.92	4.91 ± 1.69	7.81 ± 2.37

X from the $3\ell+X$ sample (only basic track quality and impact-parameter cuts). The additional electron must similarly be well separated with $\Delta R > 0.1$ from the other leptons and not yield an invariant mass below 5 GeV (J/ψ cut) in any alternative opposite-sign pairing. The 4ℓ vertex cut is not applied as a fit has not been performed.

The template for the γ and fake contributions used in the $3\ell+X$ n_{InnerPix} fit are built from the distribution in this $Z+X$ CR in MC, looking at the truth origin of the additional electron candidate. The resulting shapes are shown in Figure 8.13, and compared to the q distribution taken directly from the $3\ell+X$ CR (this contribution is not included in the fit, but is used to subtract its component in both data and MC).

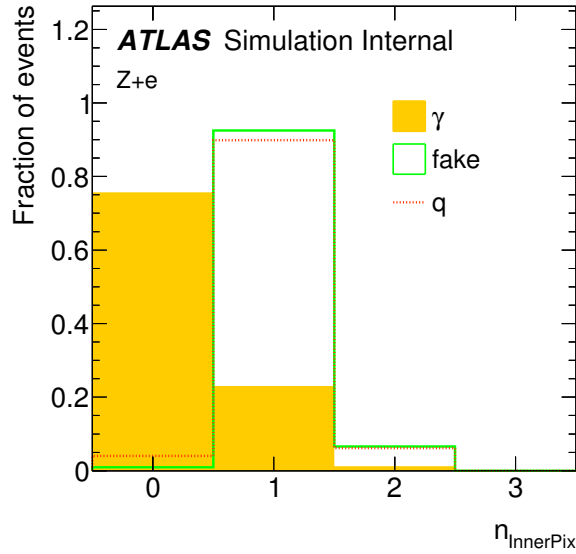


Figure 8.13 – Template shapes built from the $Z+X$ sample used in the n_{InnerPix} fit in the $3\ell+X$ CR (γ and f), compared to the shape from the q component directly extracted from the $3\ell+X$ CR (not used in the fit). All distribution are normalised to unity. [246]

The second point needed for the electron background estimation is the transfer factor from the CR to the SR, so-called selection efficiencies. These are computed in each bin of p_T and N_{jet} , by comparing the number of X electrons in the $Z+X$ CR passing or not passing the isolation and identification requirements defined for the SR. The efficiencies for the background components γ and f are computed using the truth record of the X candidates. However, differences of efficiency as a function of p_T are seen between data and MC, the latter are therefore corrected using scale factors. These differences are seen as a function of p_T , but a good agreement is found as a

function of N_{jet} . Consequently, the scale factors are derived only as a function of p_T .

The scales factors are determined separately for the γ component and the f component of the background by splitting the $Z+X$ CR into a γ -enriched region (by requiring $n_{\text{InnerPix}} = 0$) and a f -enriched region (by requiring $n_{\text{InnerPix}} > 0$), and using the iterative method described below. The purity of these regions is $\sim 90\%$ for the γ -enriched region, and $\sim 98\%$ for the f -enriched region. The differences between data and MC arise from two sources: firstly, from the selection efficiency modelling of the X in the MC, and secondly, from the efficiencies of the impurities being different from the main component (γ or f) in the two $Z+X$ sub-regions.

In the γ -enriched $Z+X$ region, the γ efficiency is much higher than the impurity efficiency thus, despite the impurities accounting for $\sim 10\%$ of the sample, the scale factor is not much affected by such impurities. On the contrary, the f -enriched $Z+X$ region is pure at $\sim 98\%$ but the efficiency of impurities is higher than the f efficiency. In order to restore the correct efficiencies, scale factors are defined according to the following procedure. The γ and f efficiencies (ϵ_γ and ϵ_f) are defined as

$$\epsilon_\gamma = \mathcal{S}_\gamma \cdot \frac{N_\gamma^{\text{pass}}}{N_\gamma^{\text{tot}}} = \frac{s_\gamma^{\text{pass}}}{s_\gamma^{\text{tot}}} \cdot \frac{N_\gamma^{\text{pass}}}{N_\gamma^{\text{tot}}} \quad \text{and} \quad \epsilon_f = \mathcal{S}_f \cdot \frac{N_f^{\text{pass}}}{N_f^{\text{tot}}} = \frac{s_f^{\text{pass}}}{s_f^{\text{tot}}} \cdot \frac{N_f^{\text{pass}}}{N_f^{\text{tot}}}, \quad (8.3)$$

where N_γ^{tot} is the number of true γ events in the $Z+X$ CR, N_γ^{pass} is the number of true γ events in the $Z+X$ CR with the X additionally passing the nominal selection requirements, $\mathcal{S}_\gamma$ is the efficiency scale factor for the γ component, s_γ^{tot} and s_γ^{pass} are the scale factors to the true γ yield in the γ -enriched $Z+X$ CR defined below, and similarly for the f part. The yield scale factors are simultaneously defined (separately for the “pass” and “tot” cuts) by

$$s_\gamma = \frac{N_{\text{data}} - N_e - s_f \cdot N_f}{N_\gamma} \quad \text{and} \quad s_f = \frac{N_{\text{data}} - N_e - s_\gamma \cdot N_\gamma}{N_f}, \quad (8.4)$$

where N_{data} is the data yield in the enriched $Z+X$ CR, and N_e , N_f and N_γ are the yields for true electrons, fakes and photons, respectively, in the enriched $Z+X$ CR. A first iteration is done computing s_γ assuming $s_f = 1$, then the value of s_γ is injected into the definition of s_f , which is in turned used to recompute s_γ , *etc.* After 3 such cycles, s_γ and s_f converge toward stable values, and the efficiency scales factors can be computed as $\mathcal{S} = s^{\text{pass}}/s^{\text{tot}}$ for both of the γ and f components. The efficiency scale factors for the γ and f components in each p_T bin are summarised in Table 8.9.

Table 8.9 – Data/MC efficiency scale factors for the f and γ component of the background evaluated on the X electron candidate of the $Z+X$ CR. [246]

p_T [GeV]	scale factor [%]	
	f component	γ component
[7, 10]	0.93 ± 0.04	1.02 ± 0.03
[10, 15]	1.10 ± 0.05	1.11 ± 0.05
[15, 70]	1.37 ± 0.13	1.89 ± 0.11

8.4.6 Categorised background yield and determination of the background shapes

The background yield estimations from Sections 8.4.2, 8.4.3, 8.4.4 and 8.4.5 were described in the “inclusive” case, *i.e.* when the target analysis is not split into bins of $p_T(4\ell)$, N_{jet} , production mode, or any other variable or combination of variables. As most of the analyses are performed

differentially (*i.e.* split in kinematical bins), the background yields in each of these bins are a fraction of the total background.

For the muon background, the estimate is simply performed in each bin required by the analysis. The procedure described in this Section can be performed inclusively, or per category bin as defined in Section 8.6 if required.

For the electron background, the heavy-flavour component yield is directly taken from the MC as in the inclusive case. The fake and conversion components yields are computed using the $sPlot$ method to obtain the yield in each category.

The non-resonant $ZZ^{(*)}$ background yield can be straightforwardly constrained from the side bands in each needed category. Its shape is always taken from the MC.

Finally, the reducible background estimate presented here only deals with the expected yields for the full $m_{4\ell}$ spectrum, but the shape must also be known for complete accuracy. The muon background $m_{4\ell}$ shape is simply taken from the relaxed VR in data. For the electron background, the heavy flavour $m_{4\ell}$ shape is directly taken from the MC, while the light flavour and conversion electron $m_{4\ell}$ shapes are taken from the $3\ell+X$ CR.

8.5 Systematic uncertainties common to all $H \rightarrow 4\ell$ analyses

This Section gives an overview of the method for the systematic uncertainty evaluation and an estimate of the typical uncertainty values. The precise estimates for the CP -odd couplings analysis are carried out in Chapter 9.4. Uncertainties are usually split into three groups related by a common origin: uncertainties linked to experimental conditions, theoretical uncertainties affecting the signal modelling, and uncertainties affecting the background estimations.

The experimental uncertainties affect the measurement of the primary objects such as electrons, muons and jets, and also include contributions from the pileup modelling and, of course, the uncertainty on the recorded luminosity. These constitute a set of approximately 200 nuisance parameters (NP) that can depend on the phase-space region. They are summarised in Section 8.5.1.

Theoretical uncertainties include uncertainties on the cross-sections of the Higgs boson signal and background production due to uncertainties on several QCD parameters, missing higher order in the calculations, and migration effects between production modes due to a changing number of jets in the event. These are described in Sections 8.5.2 and 8.5.3 for the signal and background modelling respectively.

Systematic uncertainties are evaluated with the MC simulation and have three types of effect:

- account for bin migration between different regions of the phase-space, which does not change the total event yield,
- change the properties of the physics objects (energy scale factor, resolution, *etc.*), adding or removing events from the selection,
- and change the weight of events (reconstruction or identification efficiencies, *etc.*).

The second type is evaluated by redoing the analysis with the changed properties, The third is evaluated by changing the weight of each event after varying each nuisance parameter (NP), one at a time. Each NP comes with a “up” and “down” variation corresponding to an uncertainty of plus or minus 1σ (respectively) on the quantity varied. If only a one-sided variation is available, the one provided is taken to be symmetric and account for both the up and down variations.

The systematic uncertainty σ_{NP} associated to a single nuisance parameter NP is defined as the relative yield variation with and without the $(\pm)1\sigma$ variation of the nuisance parameter:

$$\sigma_{\text{NP}} = \left| \frac{\sum_{k \in \text{events}} (w^k + \delta w_{\text{NP}}^k) - \sum_{k \in \text{events}} w^k}{\sum_{k \in \text{events}} w^k} \right|, \quad (8.5)$$

where w^k is the nominal event weight (*i.e.* without variation) of event numbered k and δw_{NP}^k is the weight variation associated to the nuisance parameter NP for event numbered k . This procedure is repeated for both the “up” and “down” variation of each nuisance parameter included in the analysis. For leptons and jets, the individual object weights are multiplied to give the total event weight:

$$w^k = \prod_{i \in \text{lep/jet}} w_i^k \quad \text{and} \quad w^k + \delta w^k = \prod_{i \in \text{lep/jet}} (w_i^k + \delta w_i^k). \quad (8.6)$$

Note that this definition correlates the variations among leptons or jets, providing a conservative estimate of the uncertainty.

A summary of the typical impact of these systematic uncertainties on production and decay cross-section measurements can be found in Table 8.10. Additional systematic uncertainties specific to the CP -odd coupling analysis are described in Chapter 9.

Table 8.10 – Typical impact of the dominant systematic uncertainties (in percent rounded to the nearest 0.5 %) on the measured inclusive fiducial cross-section and on the STXS Stage-0 (see Section 8.6.1 production mode cross-sections for the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ Run 2 analysis. Similar sources of systematic uncertainties are grouped together: luminosity (Lumi.), electron and muon reconstruction and identification efficiencies and pileup modelling (e , μ , pileup), jet energy scale and resolution and b -tagging efficiencies (jets, flavour tagging), uncertainties on reducible background, theoretical uncertainties on $ZZ^{(*)}$ background and tXX background, and theoretical uncertainties on the signal due to PDF, QCD renormalisation and factorisation scales, and parton showering algorithm. The luminosity uncertainty (nominally 1.7 %) can increase for a measurement due to simulated background. [245]

Measurement	Experimental uncertainties [%]				Theory uncertainties [%]				
	Lumi.	e , μ , pileup	Jets, flavour tagging	Red. bkg.	$ZZ^{(*)}$ bkg.	tXX bkg.	Signal		
							PDF	QCD scale	PS
Inclusive fiducial cross-section									
σ_{comb}	1.7	2.5	—	< 0.5	1	< 0.5	< 0.5	2	1
Per decay final state fiducial cross-sections									
4μ	1.7	2.5	—	0.5	1	< 0.5	< 0.5	2	1
$4e$	1.7	7	—	0.5	1.5	< 0.5	< 0.5	2	0.5
$2\mu 2e$	1.7	5.5	—	0.5	1	< 0.5	< 0.5	2	1.5
$2e 2\mu$	1.7	2.0	—	0.5	1	< 0.5	< 0.5	2	1
STXS Stage-0 production bin cross-sections									
ggF	1.7	1.5	1	0.5	1.5	< 0.5	0.5	1	2
VBF	1.7	1	4.5	0.5	2	0.5	1.5	8	6
VH	1.8	1.5	3.5	1	5	0.5	2	12	8
ttH	1.7	1	4.5	1	1	0.5	0.5	8	4

8.5.1 Experimental uncertainties

One important experimental uncertainty that reflects straightforwardly on the cross-section measurements is the uncertainty on the recorded luminosity. This uncertainty has been measured for the full Run 2 in Ref. [124] and amount to 1.7 %. It affects the normalisation of backgrounds,

except when inferred from the side bands (as for the non-resonant $ZZ^{(*)}$ component) or other data-driven methods.

Systematics uncertainties on the leptons are two fold: firstly, the trigger, reconstruction, identification and isolation efficiencies, secondly, their energy or momentum scale and resolution. All these uncertainties are derived from $Z \rightarrow \ell^+\ell^-$ and $J/\psi \rightarrow \ell^+\ell^-$ as described in Refs. [130, 134, 141] and an overview of the electron energy scale and resolution can be found in Chapter 5. They are dominated by the electron reconstruction and identification (1% to 2%) and muon reconstruction and identification (less than 1%), with an additional contribution from the electron and muon isolation efficiencies ($\sim 1\%$). The electron energy scale and resolution, and the muon momentum scale and resolution are found to have a negligible impact for the cross-section measurements.

The jet uncertainties are described in Refs. [252] and [253] and mostly affect the categories with jets, *i.e.* VBF, VH , and ttH . In such categories, each of the jet energy scale and resolution brings a systematic uncertainty in the range of 1% to 3%. Since these production modes only account for 10% of the total signal yield, the jet uncertainties have a negligible impact on the inclusive measurements. The b -tagging uncertainty [254] is only relevant in the ttH category, in which it adds an additional 1% systematic uncertainty. The missing transverse energy reconstruction uncertainties are found to have a negligible impact in all measurements.

The limited accuracy of the pileup modelling in the simulation is taken into account by varying the number of pileup interaction per event in order to cover the differences between the simulation and the measurement from inelastic collision [255]. The pileup rejection algorithms (see Section 4.6) uncertainties are also included.

An additional systematic is defined to take into account the small dependency of the neural network (defined in Section 8.6.2) scores used to discriminate the $ZZ^{(*)}$ background and the various Higgs signals on the Higgs boson mass. The only non-negligible impact is found in the ggF 0- and 1-jet categories, and is below 2% in the high score (“signal-like”) bins for a m_H uncertainty of 0.27 GeV (uncertainty from the ATLAS+CMS combined Run 1 measurement [55]). As an additional cross-check, the neural network scores for ggF are tested against an independent sample generated using MADGRAPH5_aMC@NLO. The scores agree within statistical accuracy of the samples so no additional systematic uncertainty is considered.

8.5.2 Theory uncertainties on the signal modelling

The common theory systematics cover for the incomplete knowledge of QCD effect on cross-sections: missing higher orders in the matrix element computations, choice of the QCD renormalisation and factorisation scales, modelling of the parton showers and of the underlying event, and uncertainties from the experimental results used in the determination of the strong coupling constant α_s and of the parton distribution functions (PDFs).

Missing higher order and QCD scales variations. One of the dominant uncertainty comes from the contribution of ggF in higher jet multiplicity categories. This is estimated using the method recommended by the Yellow Report 4 [72] (from developments performed in Refs [256–260]) and includes the impact of missing higher QCD orders in the generation, of varying the renormalisation and factorisation scales (μ_R and μ_F respectively) by a factor of 2 in the cross-section calculations, and of event migration between jet bins. The variations are considered uncorrelated among the various production modes.

Parton shower uncertainties. The main parton shower software used during the event generation is PYTHIA 8 with the AZNLO tune (see Section 8.2). This tune comes with a set of uncertainties which are independently varied, each providing one nuisance parameter for the analyses. Additional uncertainties are computed by changing PYTHIA 8 to HERWIG 7 [261] for all Higgs boson production modes. The AZNLO tune variations are taken as correlated among all production modes, while the systematic uncertainties from the Herwig comparison are considered uncorrelated.

PDF variations. Similarly to the parton shower uncertainties, the nominal PDF set (PDF4LHC 3.0, see Section 8.2) is provided with a set of uncertainties, each of them being assigned to a single nuisance parameter in the analysis. They are computed in the same way for each of the production modes and are therefore considered correlated.

8.5.3 Background uncertainties

The shape of the non-resonant $ZZ^{(*)}$ background is studied using an alternative generator (POWHEG instead of the nominal SHERPA). A systematic uncertainty is assigned as the difference of normalisation in the signal region and the side bands between the two generators. An uncertainty is also assigned on the shape of the $ZZ^{(*)}$ background from the nominal SHERPA generator itself similarly to the method used in the signal samples: the renormalisation and factorisation scales are varied, and the PDF and PS internal variation sets are assigned to individual nuisance parameters.

Since the $gg \rightarrow ZZ^{(*)}$ component of the irreducible background is only generated at leading order, a scale factor of 1.7 is applied to match the NLO cross-section. An uncertainty of 100 % is applied on this scale factor, varying the relative contributions from $qq \rightarrow ZZ^{(*)}$ and $gg \rightarrow ZZ^{(*)}$ to the total $ZZ^{(*)}$ background.

The reducible background uncertainties on $Z + \text{jets}$ and $t\bar{t}$ are propagated from the uncertainties found in Section 8.4. They include the statistical uncertainties from the normalisation fit ($\sim 3\%$) and the systematics on the various transfer factors and efficiencies ($\sim 6\%$). The statistical limitation from the finite number of events in the simulated samples (8 % to 70 %) is also included.

The $t\bar{t}Z$ background is estimated from the side bands similarly to the non-resonant background. An additional systematic uncertainty on the shape is evaluated by using a $t\bar{t}+V$ sample generated with MADGRAPH5_aMC@NLO instead of SHERPA.

Lastly, the VVV background yield is also assigned an uncertainty based on varying the PDF and QCD scales.

8.6 Event categorisation for the couplings analysis

The CP analysis (Chapter 9) relies on the measurement of kinematics that are specific to the VBF production mode. As a consequence, a category as pure as possible in VBF events needs to be defined. Such categories are experimentally defined to match the STXS framework (Chapter 7.2.2) which provides a truth level categorisation. The purity in each category is subsequently improved by the use of MVA discriminants.

8.6.1 STXS categorisation

In order to have a handle on various Higgs boson couplings, reconstruction categories sensitive to the various couplings are defined. Since the various Higgs boson production modes do not exhibit the same couplings of the Higgs boson (for instance ggF probes the coupling to gluons while VBF and VH probe the couplings to Z and W bosons), one possible categorisation is to split according to the various production modes. The STXS framework has been set up to target such couplings measurements, allowing to probe different potential BSM effects with best sensitivity while controlling theoretical uncertainties due to the jet multiplicity.

This framework was described in Section 7.2.2 for the full Stage-1.1 splitting targeting the full Run 2 analyses. However, given the high number of bins proposed in this framework and the limited available statistic in the $H \rightarrow 4\ell$ channel, using all the defined bins would result in large statistical errors in the measurements. A “reduced” Stage-1.1 was therefore developed by the $H \rightarrow 4\ell$ analysis group to provide a categorisation in agreement with the expected statistical sensitivity. Each category is called a *production bin*.

Compared to the full Stage-1.1, all ggF + 2 jets bins are merged, yielding a total 7 bins for the ggF process: six of them have a $p_T^{4\ell} < 200$ GeV, the 0-jet category is split into $p_T^{4\ell} \in [0, 10, 200]$ GeV, the 1-jet category being is into $p_T^{4\ell} \in [0, 60, 120, 200]$ GeV, and the 2-jets category is inclusive. The last category includes all events with $p_T^{4\ell} > 200$ GeV independent on the number of jets. The VBF and VH category are not split into bins of N_{jet} or m_{jj} with VBF and VH -had merged together, but rather simply split the VBF only categories into $p_T^{4\ell}$ lower or greater than 200 GeV, and the VH categories into hadronically or leptonically decayed. This reduced splitting is summarised in the second from the left panel of Figure 8.14 along with the nomenclature for all these bins.

Analyses then define *reconstruction categories*, designed to closely match the “*truth production bins*”. The event reconstruction categorisation flow is as follow: first the event is tested for the ttH production mode. It is classified as ttH -Lep-enriched (targeting semi-leptonic and fully leptonic top decays) if the event has at least one additional lepton (with $p_T > 12$ GeV) and satisfies one of the following three requirements on jets: at least two b -tagged jet (at 85 % b -tagging efficiency), or at least five jets with at least one b -tagged jet (at 85 % b -tagging efficiency), or at least two jets with at least one b -tagged jet (at 60 % b -tagging efficiency). It is instead classified as ttH -Had-enriched (targeting fully hadronic top quark decays) if the event has at least five jets with at least two b -tagged jets (at 85 % b -tagging efficiency), or at least four jets with at least one b -tagged jet (at 60 % b -tagging efficiency). Otherwise, if the event still has at least one additional lepton (with $p_T > 12$ GeV), the event is classified as VH -Lep-enriched.

The classification flow then continues, splitting with respect to the number of jets: if the event has at least two jets with invariant mass $m_{\text{jj}} > 120$ GeV and $p_T^{4\ell} > 200$ GeV, it is sent to the 2j-BSM-like category. The remaining events with at least two jets are classified at 2j, which contains VBF, VH but also ggF events. Events with exactly one jet are categorised as for the truth categories, splitting in $[0, 60, 120, 200]$ GeV $p_T^{4\ell}$ bins (1j- $p_T^{4\ell}$ -Low, 1j- $p_T^{4\ell}$ -Med, 1j- $p_T^{4\ell}$ -High), plus an additional one for $p_T^{4\ell} > 200$ GeV (1j- $p_T^{4\ell}$ -BSM-Like). Events without jets are classified closely to the truth categories, splitting in $p_T^{4\ell} \in [0, 10, 100]$ and greater than 100 GeV (0j- $p_T^{4\ell}$ -Low, 0j- $p_T^{4\ell}$ -Med, 0j- $p_T^{4\ell}$ -High).

The $m_{4\ell}$ sidebands $[105, 115] \cup [130, 160]$ (or $[130, 350]$) are also classified, according to the number of additional leptons (SB- VH -Lep-enriched with at least one additional $p_T > 12$ GeV lepton) or to the number of jets (SB-0j if $N_{\text{jet}} = 0$, SB-1j if $N_{\text{jet}} = 1$, and SB-2j if $N_{\text{jet}} \geq 2$). An additional category is defined to match the ttH signal region: SB- tXX -enriched, gathering events with at least two jets, of which at least one is b -tagged (at 60 % b -tagging efficiency), and with $E_T^{\text{miss}} > 100$ GeV.

The signal and sideband regions reconstruction categories are presented in Figure 8.14, along with the STXS reduced Stage-1.1 truth categories.

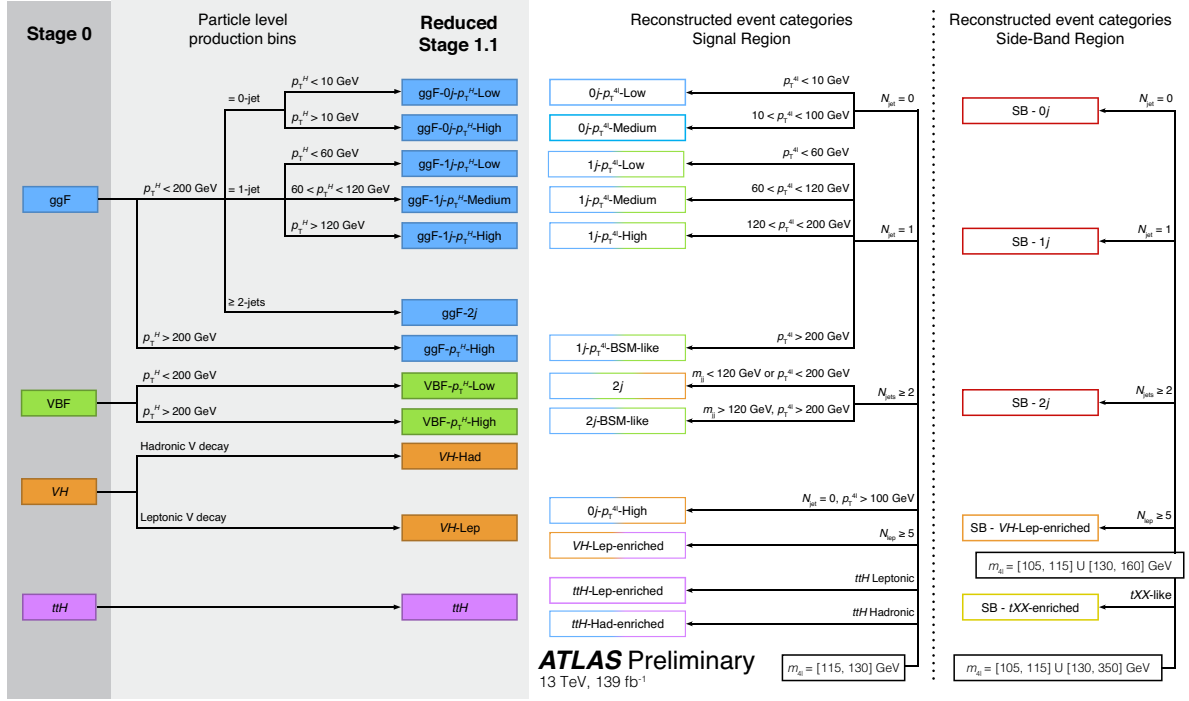


Figure 8.14 – Truth production categories as defined in the STXS Stage-0 and Stage-1.1 (left panel), and corresponding reconstruction categories in the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ analysis for in the signal (middle) and sidebands (right) regions. [245]

The expected yields of each reduced Stage-1.1 truth category in each reconstruction category is shown in Figure 8.15. Figure 8.16 displays the same information but in fractional form for each reconstruction category. The expected STXS stage-0 yields in each reconstructed is shown in Table B.1 of Appendix B.

In order to improve the purity of the targeted signal in each reconstructed category, multivariate analysis algorithms are trained in most reconstructed categories to provide a better signal and background separation. These methods are described in Section 8.6.2.

8.6.2 Improved category separation using neural networks

In order to improve the signal purity in most of the reconstruction categories, a set of multivariate algorithms (MVA) are trained, one per category. The MVA used in the analysis have long been based on boosted decision trees (BDT) but have recently been changed to deep neural networks (dNN). This change brings additional improvement in the signal discrimination by the use of multiple targets and the inclusion of lower level variables. Moreover, on the contrary to BDT, recurrent neural networks (rNN) accept a varying number of inputs which is convenient for the number of jets which varies event-by-event.

The general architecture of the neural network is the following: two recurrent neural networks (rNN), one for the kinematics of the four leptons, the other for the kinematics of up to three jets, and one multilayer perceptron (MLP), which adds higher level variables to the training, are combined using another MLP, the four of them constituting a dNN. One dNN per reconstruction category is trained to discriminate between two or three signal types, including various SM Higgs boson modes or $ZZ^{(*)}$ background. The dNN output can be interpreted as the probability for an

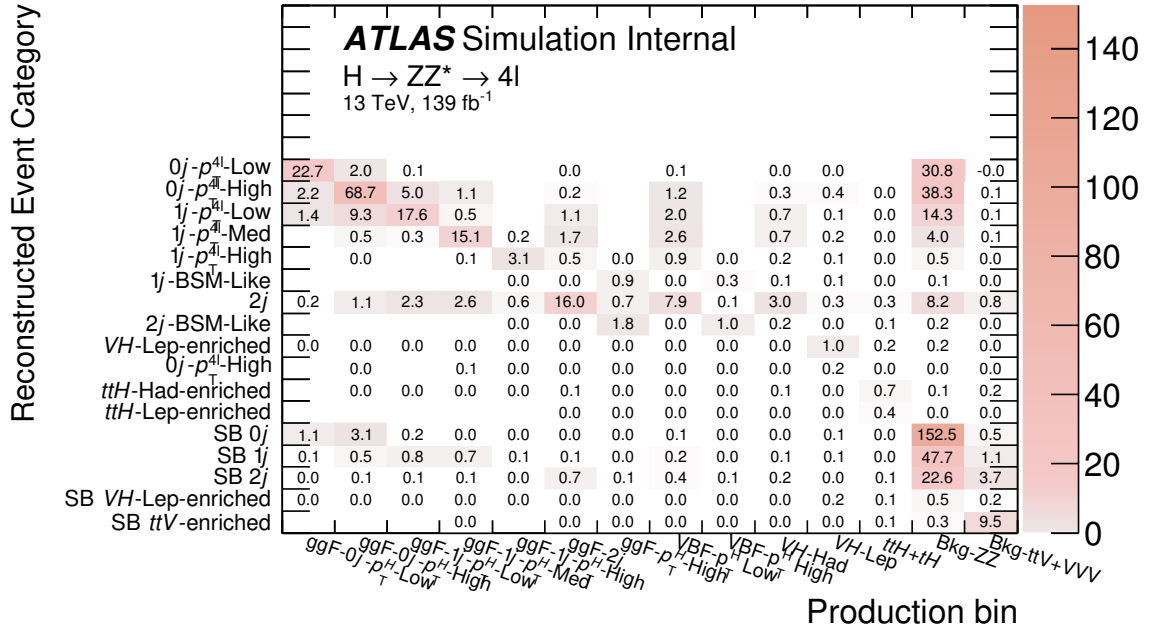


Figure 8.15 – Expected yield in each reconstruction category (signal and sideband) of the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ analysis in terms of STXS Stage1.1 truth category, assuming an integrated luminosity of 139 fb⁻¹ at $\sqrt{s} = 13$ TeV. The bbH contribution is included in the ggF ones, and the tH is included in the ttH one. [262]

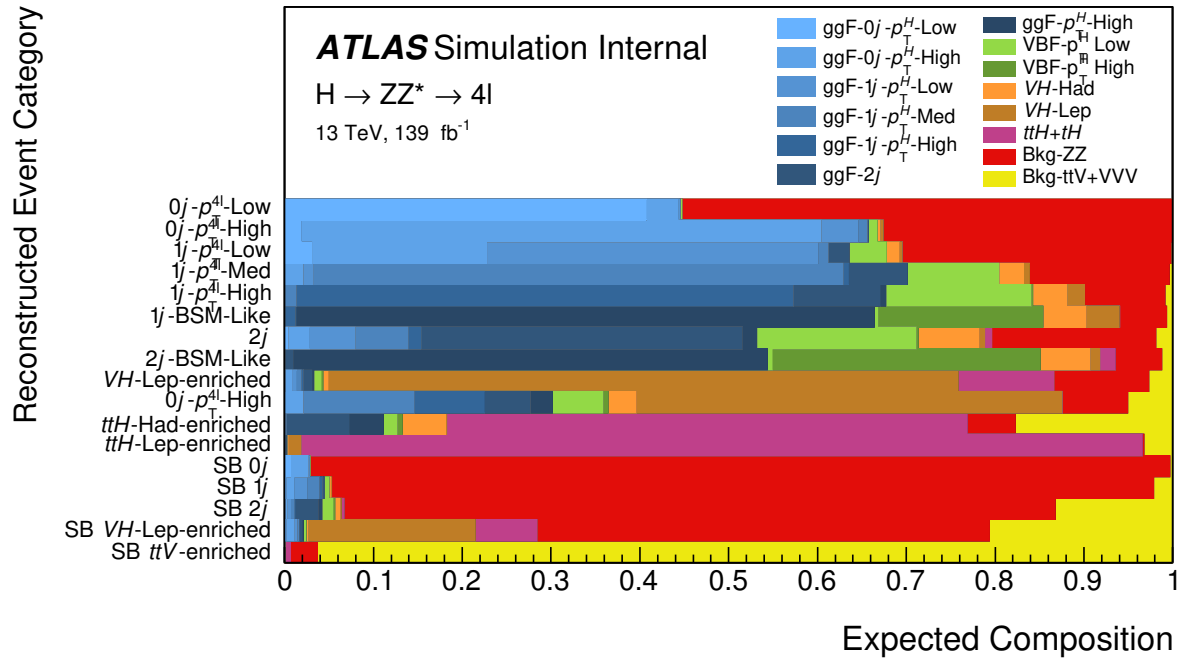


Figure 8.16 – Expected relative composition in each reconstruction categories of the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ analysis in terms of reduced STXS Stage1.1 truth category. The bbH contributions are included in the ggF production bins. [262]

event to be of one or the other type. For categories where two signals are discriminated against, only one output probability is given since the probabilities for the two signal types should add up to one. In categories where three signals are distinguished, two output probabilities are provided.

The different variables and the signal types targeted for the discrimination are presented in Table 8.11. An exhaustive description of the variables used in the NN training follows:

$p_{T,\ell}, \eta_\ell$	lepton transverse momentum and η coordinate,
$p_{T,j}, \eta_j$	jet transverse momentum and η coordinate,
$p_T^{4\ell}, \eta_{4\ell}$	transverse momentum and η coordinate of the four-lepton system,
m_{12}, m_{34}	leading and subleading lepton pair invariant masses,
$ \cos \theta^* $	leading Z boson production angle in the four-lepton rest frame,
$\cos \theta_1$	angle between the negatively charged lepton of the leading Z boson in the leading Z boson rest frame and the direction of flight of the leading Z boson in the four-leptons rest frame,
ϕ_{ZZ}	angle between two Z bosons decay planes in the four-lepton rest frame,
$\min(\Delta R_{4\ell j})$	angular separation of a jet from the 4ℓ system,
m_{jj}	invariant mass of the two leading jets system (j_1, j_2 respectively),
$\Delta\eta_{jj}$	angular separation in the η direction between the two leading jets ($ \eta_{j_1} - \eta_{j_2} $),
$p_{T,4\ell jj}$	transverse momentum of the system constituted of the four leptons and the two leading jets,
$\Delta\eta_{4\ell jj}$	$ \eta_{4\ell} - (\eta_{j_1} + \eta_{j_2})/2 $,
N_{jets}	number of jets in the event,
$N_{b\text{-jets}}$	number of b -jets,
E_T^{miss}	missing transverse energy in the event
HT	scalar sum of all energy and momentum vectors entering the E_T^{miss} computation,
$D_{ZZ^*}, \ln(\mathcal{M}_{\text{sig}} ^2)$	denoting $\mathcal{M}_{\text{sig}}$ and $\mathcal{M}_{ZZ}$ the leading order matrix elements (similar to the matrix element ambiguity resolver introduced in Section 8.3.4) for the ggF and $qq \rightarrow ZZ^{(*)}$ hypotheses, $D_{ZZ^*} = \log(\mathcal{M}_{\text{sig}} ^2/ \mathcal{M}_{ZZ} ^2)$. It is used to discriminate between ggF and the $ZZ^{(*)}$ background.

Note that not all these variables are used for each NN.

The expected distributions for the dNN scores in each category are presented in Figures 8.17 and 8.18. The observed data points are superimposed. Note that Figures 8.18g and 8.18h present the yields in counting only categories (not using dNN): $0j\text{-}p_T^{4\ell}\text{-High}$, $1j\text{-}p_T^{4\ell}\text{-BSM-Like}$ and $t\bar{t}H\text{-Lep-enriched}$.

The merging of the VBF and VH -had reconstruction categories mentioned in Section 8.6.1 and seen in Figure 8.14 is justified by the need of the NN to have sufficient statistics for both the signal and background components to discriminate against. If the two categories were split, for example by cutting on m_{jj} at 120 GeV, the low invariant mass region would not contain enough VBF events to allow the MVA algorithm to perform to its best.

8.7 Expected yields with the full Run-2 dataset

The number of expected events for the signal and background contribution in each of the four final states 4μ , $2e2\mu$, $2\mu2e$ and $4e$ including statistical and systematic uncertainties, as well as

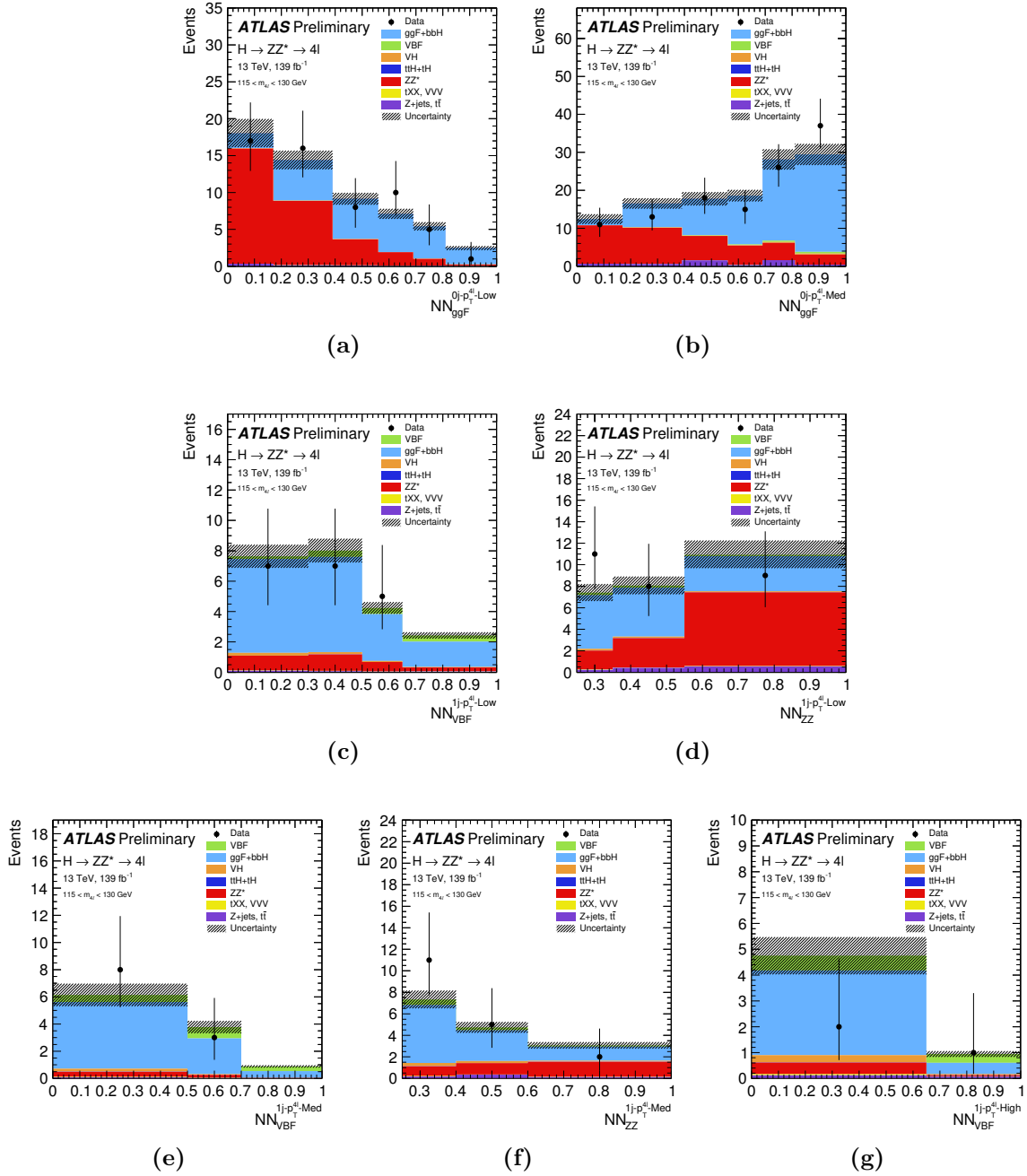


Figure 8.17 – Observed and expected NN output (pre-fit) distributions in the different reconstruction categories of the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ Run 2 analysis for an integrated luminosity of 139 fb^{-1} and at $\sqrt{s} = 13 \text{ TeV}$: (a) NN_{ggF} in $0j\text{-}p_{\text{T}}^{4\ell}\text{-Low}$, (b) NN_{ggF} in $0j\text{-}p_{\text{T}}^{4\ell}\text{-Med}$, (c) NN_{VBF} in $1j\text{-}p_{\text{T}}^{4\ell}\text{-Low}$ with $\text{NN}_{\text{ZZ}} < 0.25$, (d) NN_{ZZ} in $1j\text{-}p_{\text{T}}^{4\ell}\text{-Low}$ with $\text{NN}_{\text{ZZ}} > 0.25$, (e) NN_{VBF} in $1j\text{-}p_{\text{T}}^{4\ell}\text{-Med}$ with $\text{NN}_{\text{ZZ}} < 0.25$, (f) NN_{ZZ} in $1j\text{-}p_{\text{T}}^{4\ell}\text{-Med}$ with $\text{NN}_{\text{ZZ}} > 0.25$, (g) NN_{VBF} in $1j\text{-}p_{\text{T}}^{4\ell}\text{-High}$. The SM Higgs boson signal is assumed with a mass $m_H = 125 \text{ GeV}$. The bin boundaries are chosen to maximise the significance of the targeted signal in each category. The uncertainty in the prediction is shown by the hatched band, including also the theoretical uncertainties on the SM cross-section for the signal and the $ZZ^{(*)}$ background. [245]

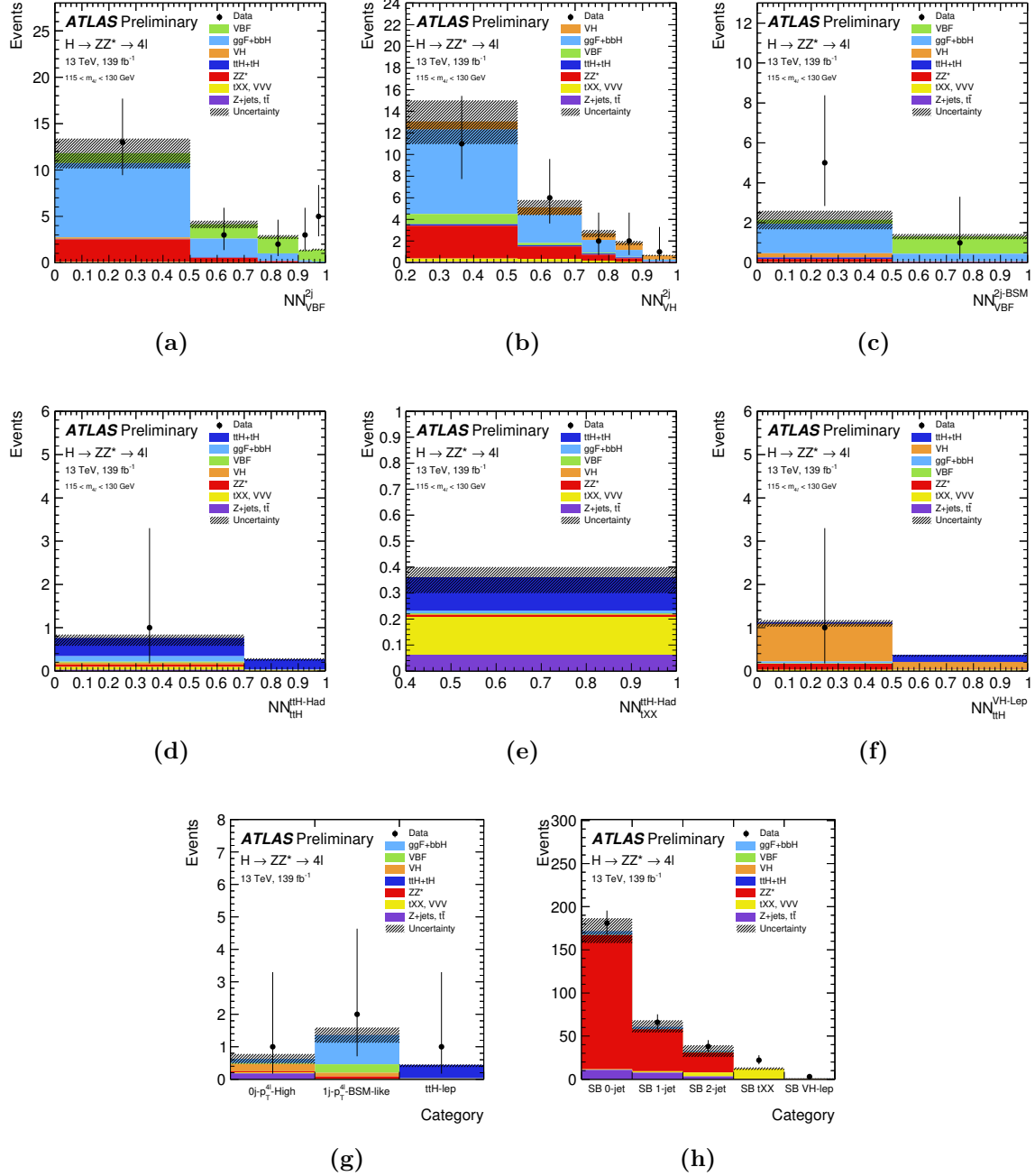


Figure 8.18 – Observed and expected NN output (pre-fit) distributions in the different reconstruction categories of the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ Run 2 analysis for an integrated luminosity of 139 fb^{-1} and at $\sqrt{s} = 13 \text{ TeV}$: (a) NN_{VBF} in $2j$ with $NN_{VH} < 0.2$, (b) NN_{VH} in $2j$ with $NN_{VH} > 0.2$, (c) NN_{VBF} in $2j$ -BSM-like, (d) NN_{ttH} in ttH -Had-enriched with $NN_{tXX} < 0.4$, (e) NN_{tXX} in ttH -Had-enriched with $NN_{tXX} > 0.4$, (f) NN_{ttH} in VH -Lep-enriched. The categories where no NN discriminant is used are also shown (g), together with the sidebands used to constrain the $ZZ^{(*)}$ and tXX backgrounds (h). The SM Higgs boson signal is assumed with a mass $m_H = 125 \text{ GeV}$. The bin boundaries are chosen to maximise the significance of the targeted signal in each category. The uncertainty in the prediction is shown by the hatched band, including also the theoretical uncertainties on the SM cross-section for the signal and the $ZZ^{(*)}$ background. [245]

Table 8.11 – Input variables used to train the MLP and the two rNNs (one for the four leptons, the other for up to three jets) for each reconstruction category of the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ Run 2 analysis. The processes trained against in each category are shown in the second column. The $0j$ - $p_T^{4\ell}$ -High, $1j$ - $p_T^{4\ell}$ -BSM-Like and ttH -Lep-enriched categories do not use MVA categorisation and are therefore not listed here. [245]

Category	Processes	MLP	Lep rNN	Jet rNN
$0j$ - $p_T^{4\ell}$ -Low, $0j$ - $p_T^{4\ell}$ -Med	ggF, ZZ	$p_T^{4\ell}, D_{ZZ^*}, m_{12}, m_{34},$ $\cos \theta^*, \cos \theta_1, \phi_{ZZ}$	$p_{T,\ell}, \eta_\ell$	None
$1j$ - $p_T^{4\ell}$ -Low	ggF, VBF, ZZ	$p_T^{4\ell}, p_{T,j}, \eta_j,$ $\Delta R_{4\ell j}, D_{ZZ^*}$	$p_{T,\ell}, \eta_\ell$	None
$1j$ - $p_T^{4\ell}$ -Med	ggF, VBF, ZZ	$p_T^{4\ell}, p_{T,j}, \eta_j, E_T^{\text{miss}}$ $\Delta R_{4\ell j}, D_{ZZ^*}, \eta_{4\ell}$	None	None
$1j$ - $p_T^{4\ell}$ -High	ggF, VBF	$p_T^{4\ell}, p_{T,j}, \eta_j,$ $\Delta R_{4\ell j}, \eta_{4\ell}, E_T^{\text{miss}}$	$p_{T,\ell}$	None
$2j$	ggF, VBF, VH	$m_{jj}, \Delta\eta_{jj}, p_{T,4\ell jj}$	$p_{T,\ell}, \eta_\ell$	$p_{T,j}, \eta_j$
$2j$ -BSM-like	ggF, VBF	$\Delta\eta_{jj}, \Delta\eta_{4\ell jj}, p_{T,4\ell jj}$	$p_{T,\ell}, \eta_\ell$	$p_{T,j}, \eta_j$
VH-Lep-enriched	ttH, VH	$N_{\text{jets}}, N_{b\text{-jets}}, E_T^{\text{miss}},$ $\text{HT}, \ln(\mathcal{M}_{\text{sig}} ^2)$	$p_{T,\ell}$	None
ttH -Had-enriched	ttH, tXX, ggF	$N_{\text{jets}}, N_{b\text{-jets}}, p_T^{4\ell}, m_{jj}, \Delta\eta_{jj},$ $p_{T,jj}, \min(\Delta R_{Zj}), \Delta\eta_{4\ell jj},$ $E_T^{\text{miss}}, \min(\Delta R_{4\ell j}), \text{HT}, \ln(\mathcal{M}_{\text{sig}} ^2)$	$p_{T,\ell}, \eta_\ell$	$p_{T,j}, \eta_j$

the number observed events in the signal region $115 < m_{4\ell} < 130$ GeV is given in Table 8.12. The distribution of the $m_{4\ell}$ invariant mass (including the FSR correction) is shown in Figure 8.19. Figure 8.20 show the distribution of four-lepton transverse momentum and number of jet for events falling in the signal region.

Table 8.12 – Number of expected and observed events by the ATLAS experiment in the four decay final states for an integrated luminosity of 139 fb^{-1} at $\sqrt{s} = 13$ TeV after the event selection, in the mass range $115 < m_{4\ell} < 130$ GeV. The $ZZ^{(*)}$ background includes both the $qq \rightarrow ZZ^{(*)}$ and the $gg \rightarrow ZZ^{(*)}$ components. The “other backgrounds” include the contributions from $Z + \text{jets}$, $t\bar{t}$, WZ , $t\bar{t} + V$, and VVV . The quoted errors include both statistics and systematics uncertainties. [245]

Final state	Signal	$ZZ^{(*)}$ background	Other backgrounds	Total expected	Observed
4μ	78 ± 5	38.1 ± 2.2	2.87 ± 0.18	119 ± 5	118
$2e2\mu$	52.8 ± 3.1	26.1 ± 1.4	3.01 ± 0.19	81.9 ± 3.4	98
$2\mu 2e$	40.0 ± 2.9	17.4 ± 1.3	3.5 ± 0.5	60.9 ± 3.2	57
$4e$	35.3 ± 2.6	15.1 ± 1.5	2.9 ± 0.4	53.3 ± 3.1	43
Total	206 ± 13	97 ± 6	12.3 ± 0.9	315 ± 14	316

The number of expected signal and background events for an integrated luminosity of 139 fb^{-1} , as well as the observed number in each reconstruction category is shown in Table 8.13. The detail of the signal per production mode can be found in Table B.1 of Appendix B. No significant deviation from the expectations is found across all the reconstruction categories.

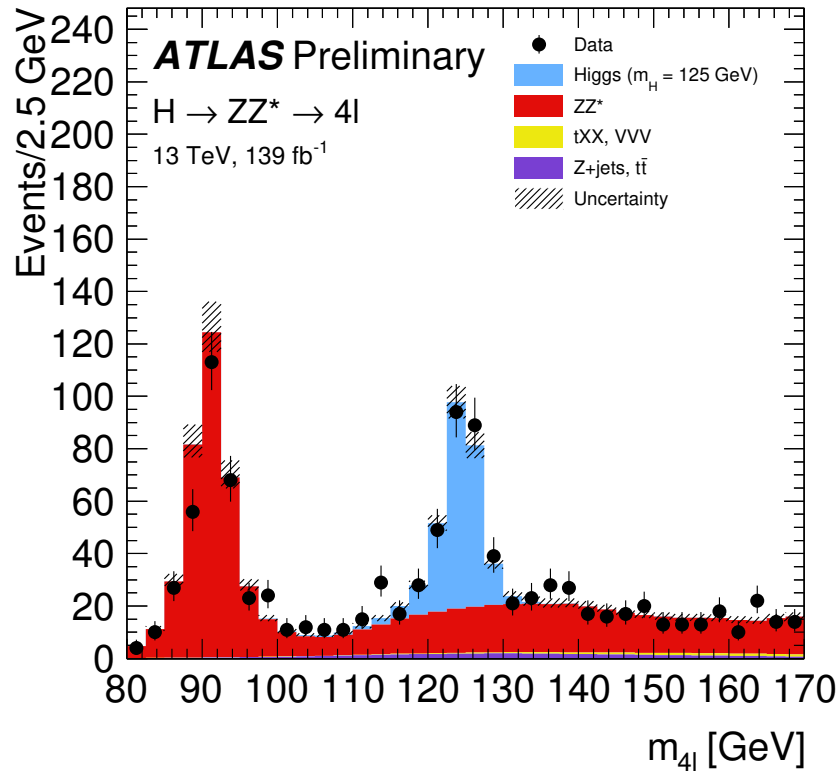


Figure 8.19 – Distribution of the four-lepton invariant mass seen by the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ Run 2 analysis in the low mass range, showing the data and compared to various contributions from the MC. [245]

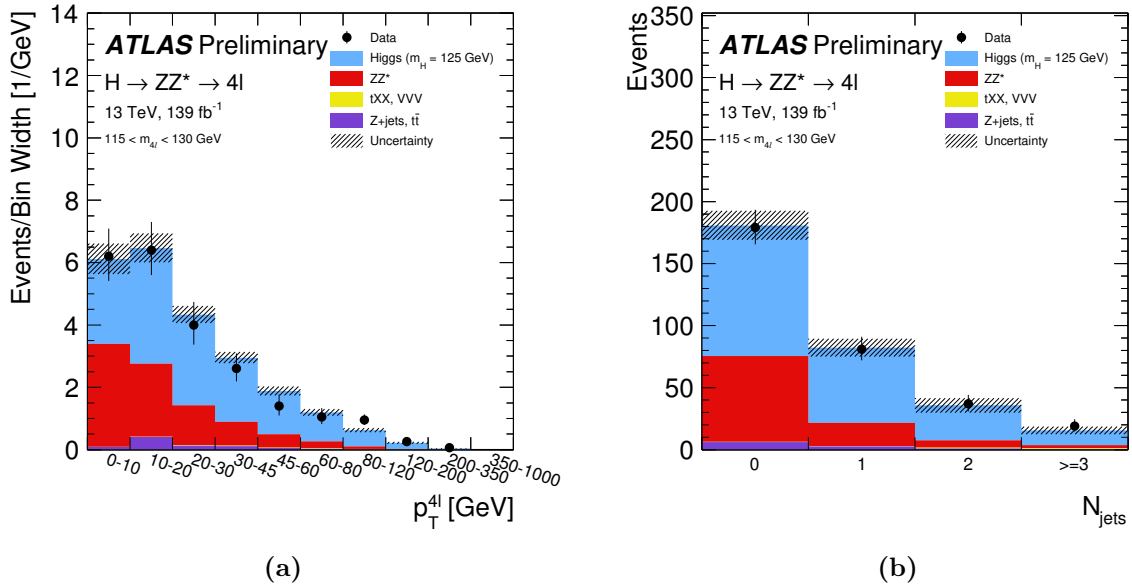


Figure 8.20 – Distributions of (a) the four-lepton system transverse momentum and (b) the number of jets shown in the signal region, as seen in the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ Run 2 analysis. [245]

Table 8.13 – Expected and observed Higgs boson signal and background yields in the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ Run 2 analysis for an integrated luminosity of 139 fb^{-1} and at $\sqrt{s} = 13 \text{ TeV}$ assuming the SM Higgs boson signal with a mass $m_H = 125 \text{ GeV}$. Results are given for each reconstructed category (each line) in the signal region $115 < m_{4\ell} < 130 \text{ GeV}$ or in the sideband region ($105 < m_{4\ell} < 115 \text{ GeV}$ or $130 < m_{4\ell} < 160 \text{ GeV}$) or ($130 < m_{4\ell} < 350 \text{ GeV}$) for tXX -enriched. The “total expected” corresponds to the sum of the expected number of SM Higgs boson events and the estimated background. The errors include both statistic and systematic uncertainties. Expected contributions below 0.2% of the total yield in each reconstructed category are not shown and replaced by “–”. [245]

Reconstructed event category	Signal	$ZZ^{(*)}$ background	tXX background	Other backgrounds	Total expected	Observed
Signal $115 < m_{4\ell} < 130 \text{ GeV}$						
$0j\text{-}p_{\text{T}}^{4\ell}\text{-Low}$	24.9 ± 3.1	31 ± 4	–	0.78 ± 0.11	57 ± 5	57
$0j\text{-}p_{\text{T}}^{4\ell}\text{-Med}$	79 ± 8	38 ± 5	0.047 ± 0.009	5.2 ± 0.5	123 ± 10	120
$0j\text{-}p_{\text{T}}^{4\ell}\text{-High}$	0.39 ± 0.04	0.033 ± 0.015	0.011 ± 0.004	0.182 ± 0.031	0.63 ± 0.05	1
$1j\text{-}p_{\text{T}}^{4\ell}\text{-Low}$	33 ± 4	14.3 ± 2.6	0.088 ± 0.007	1.52 ± 0.22	49 ± 5	47
$1j\text{-}p_{\text{T}}^{4\ell}\text{-Med}$	21.2 ± 2.8	4.0 ± 0.6	0.114 ± 0.010	0.77 ± 0.14	26.0 ± 2.8	29
$1j\text{-}p_{\text{T}}^{4\ell}\text{-High}$	5.0 ± 0.8	0.50 ± 0.09	0.047 ± 0.007	0.189 ± 0.030	5.7 ± 0.8	3
$1j\text{-}p_{\text{T}}^{4\ell}\text{-BSM-Like}$	1.28 ± 0.23	0.072 ± 0.033	0.006 ± 0.004	0.040 ± 0.008	1.41 ± 0.23	2
$2j$	35 ± 5	8.2 ± 2.4	0.96 ± 0.08	0.23 ± 0.06	45 ± 5	48
$2j\text{-BSM-like}$	3.2 ± 0.6	0.18 ± 0.06	0.032 ± 0.005	1.20 ± 0.11	4.6 ± 0.6	6
$VH\text{-Lep-enriched}$	1.25 ± 0.07	0.159 ± 0.020	0.037 ± 0.008	0.0052 ± 0.0031	1.47 ± 0.07	1
$tH\text{-Had-enriched}$	0.95 ± 0.16	0.063 ± 0.025	0.225 ± 0.031	0.09 ± 0.04	1.32 ± 0.17	1
$tH\text{-Lep-enriched}$	0.41 ± 0.04	–	0.0130 ± 0.0013	–	0.42 ± 0.04	1
Sideband $105 < m_{4\ell} < 115 \text{ GeV}$ or $130 < m_{4\ell} < 160 \text{ GeV}$						
SB-0j	4.6 ± 0.5	155 ± 14	0.22 ± 0.05	12.6 ± 1.8	173 ± 14	181
SB-1j	2.74 ± 0.29	49 ± 6	1.35 ± 0.18	6.4 ± 0.9	59 ± 6	66
SB-2j	1.90 ± 0.26	23 ± 6	4.3 ± 0.5	4.4 ± 0.6	34 ± 6	38
SB- $VH\text{-Lep-enriched}$	0.266 ± 0.015	0.49 ± 0.06	0.132 ± 0.020	0.07 ± 0.13	1.06 ± 0.14	3
Sideband $105 < m_{4\ell} < 115 \text{ GeV}$ or $130 < m_{4\ell} < 350 \text{ GeV}$						
SB- $tXX\text{-enriched}$	0.066 ± 0.008	0.31 ± 0.11	11.6 ± 1.3	0.46 ± 0.13	12.5 ± 1.3	22

Measurement of CP -odd Higgs-boson couplings

Contents

9.1	Optimal observables	240
9.1.1	CP -sensitive observables in the $H \rightarrow 4\ell$ decay channel and in the VBF production mode	240
9.1.2	Definition of Optimal Observables	242
9.1.3	Technical generation of matrix elements with MADGRAPH	243
9.1.4	Dependence on the PDF choice	245
9.1.5	Impact of changing the hypothesis provided to the matrix element computing	246
9.1.6	Development of a MADGRAPH-based simplified matrix element computation package	247
9.2	Morphing technique	250
9.2.1	Mathematical definition of the morphing method	250
9.2.2	Validation of the $\mathcal{OO}_{1,jj}$ morphed shapes	252
9.3	Preliminary studies	252
9.3.1	Additional BSM simulated samples	254
9.3.2	Sensitivity using VBF signal samples only	254
9.3.3	Limits from cross-section results	255
9.3.4	Validation of the morphing technique	257
9.3.5	Sensitivity including $ZZ^{(*)}$ and other Higgs boson production modes	257
9.4	Preliminary results	259
9.4.1	Expected sensitivity with fully-simulated samples	259
9.4.2	Systematic uncertainties assessment	261
9.4.3	Study of a potential bias from usage of neural networks	262
	Summary and conclusion of the CP-odd couplings analysis	264

This Chapter presents the CP -odd couplings analysis performed in the $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ channel using the full Run 2 dataset. Since the $pp \rightarrow H$ cross-section is modified by both CP -even and CP -odd operators in a similar way, a measurement of CP -odd effects cannot rely on inclusive cross-section only. An observable able to disentangle between CP -even and CP -odd effects must therefore be used. Such observables, called *optimal observables*, are studied in Section 9.1.

Provided the current computational resources, only a finite number of samples can be fully simulated. To cover the full phase-space of interest, points between the few generated points must be interpolated, using the method described in Section 9.2. Truth-level samples were generated covering the full phase-space of interest, and were used to determine the expected sensitivity range of several EFT Wilson coefficients targeting the CP -odd Higgs to weak-boson couplings. This as well as a comparison to expected limits from previous cross-section measurements, is presented in Section 9.3. Finally, systematic uncertainties impacting the measurements are assessed in Section 9.4.

9.1 Optimal observables

The $H \rightarrow 4\ell$ decay channel, with its high object multiplicity, yields eight independent observables characterising its final state, plus additional ones if jets are also considered for the final state as in the VBF production mode (Section 9.1.1). Some of these variables are sensitive to CP effects and provide natural discriminating variables for a CP study analysis. However, fitting multidimensional observables results in losses of fit quality, and a better approach is to combine all of this information into a single variable. For that purpose a matrix element-based (Section 9.1.6) method is used to define custom so-called *optimal observables* (OO, Section 9.1.2). This method was first introduced during the LEP era, and first applied to CP measurements in Ref. [263]. It was also used for Run 1 LHC measurement, for example in Ref. [68].

The goal of the present analysis is to test for a CP component of the Higgs boson couplings, using an observable (the optimal observable) which is directly sensitive to CP effects. However, if a statistically significant deviation is seen, it would have to be confronted with constraints coming from cross-section measurements.

9.1.1 CP -sensitive observables in the $H \rightarrow 4\ell$ decay channel and in the VBF production mode

The complete $H \rightarrow 4\ell$ final state can be characterised by three masses and five angles, represented in Figure 9.1 and described below:

- $m_{4\ell}$: the invariant mass of the four-lepton system,
- m_{12} and m_{34} : the invariant masses of the leading and subleading lepton pairs (respectively),
- θ_1 (θ_2): the angle between the negatively charged lepton of the (sub)leading lepton pair in the Z -boson rest frame and the flight direction of the (sub)leading Z boson in the four-lepton rest frame,
- ϕ : the angle between the decay planes of the two lepton pairs in the four-lepton rest frame,
- ϕ_1 : the angle between the decay plane of the leading Z boson, and a plane formed by the leading Z boson momentum in the four-lepton rest frame and the beam axis,
- θ^* : the angle of the leading Z boson in the four-lepton rest frame with respect to the beam axis.

In the production side, and especially in the VBF mode, some variables describing the jet kinematics are also sensitive to CP -odd contributions to the Higgs boson couplings. The most sensitive one is the signed difference of azimuth between the two jets, defined as

$$\Delta\phi_{jj} = \text{sign}(\eta_{j_1} - \eta_{j_2}) \cdot ((\phi_{j_1} - \phi_{j_2})[2\pi]). \quad (9.1)$$

Examples of distributions of CP -sensitive observables are shown in Figure 9.2.

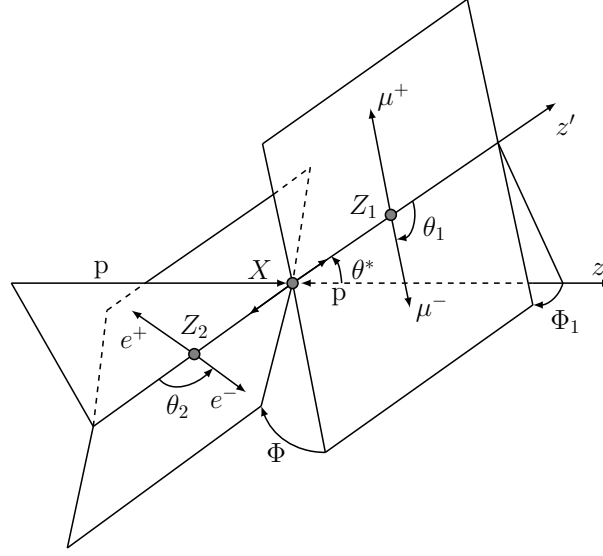


Figure 9.1 – Illustration of the decay angles defined in the $H \rightarrow 4\ell$ final state. [68]

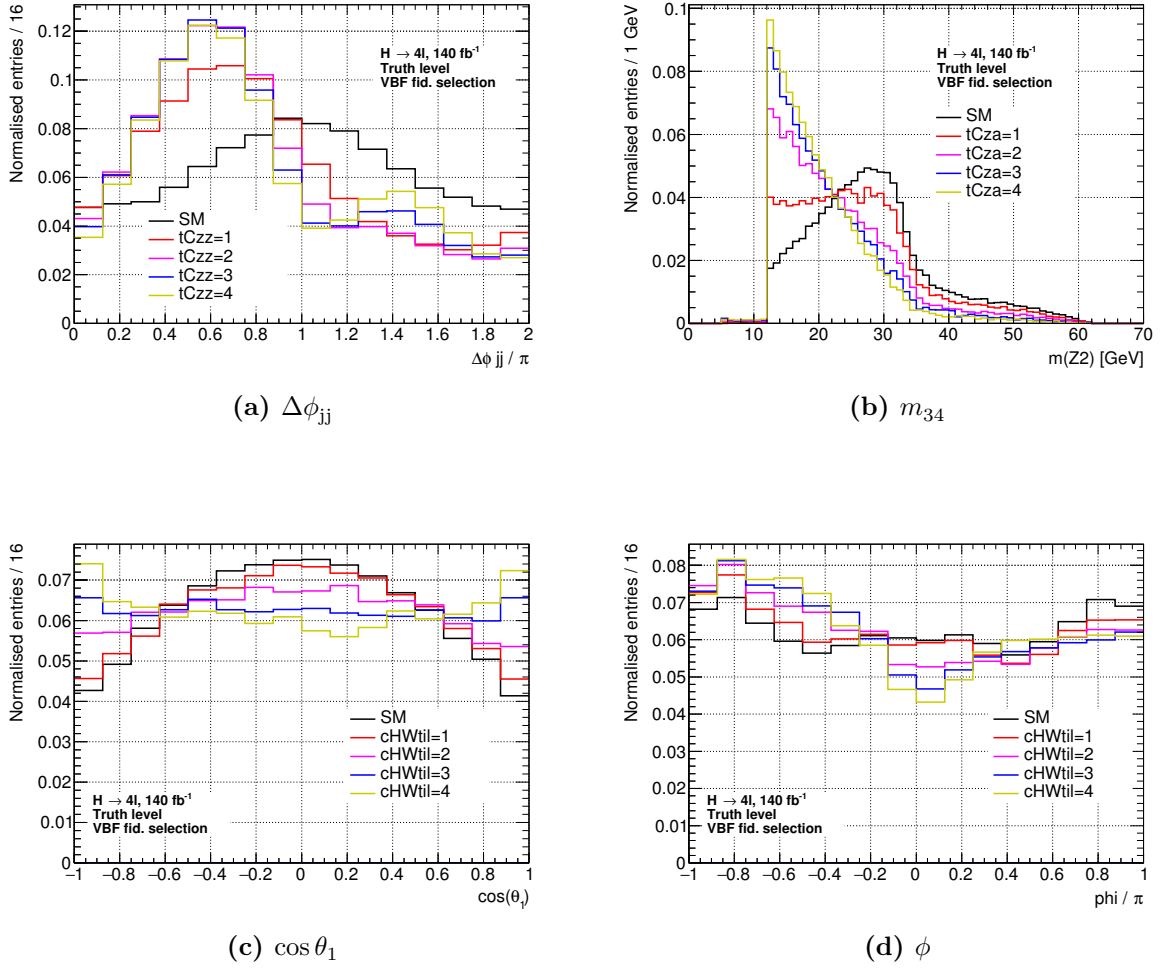


Figure 9.2 – Example distribution of (a) $\Delta\phi_{jj}$, (b) m_{34} , (c) $\cos\theta_1$, and (d) ϕ for various values of the CP -odd couplings and compared to the SM distribution (all normalised to unit integral).

9.1.2 Definition of Optimal Observables

Since a matrix element (ME) is the most elementary descriptor of a physics process, all the kinematic information of an event is contained in such an object. As a consequence, a matrix element-based observable should provide maximal information for a processes, combining all low-level quantities (lepton or jet kinematics) into a higher-level observable. Its value can however hardly be interpreted as a physical observable, so other high level quantities, such as the ones presented in the previous Section, are useful to understand its behaviour.

Matrix elements are commonly computed internally by event generators (POWHEG, PYTHIA, MADGRAPH5_aMC@NLO, *etc.*), but not all of them provide an interface to compute the matrix element for a given set of kinematics. The MADGRAPH [250] software, on top of being an event generator, can also be used to compute the matrix element of a given a set of kinematics, assuming the event is produced from a specific process (for example: VBF). Its technical use is briefly described in Section 9.1.3, and for the remainder of this Section, it will be assumed that the matrix elements of any required process is available for each event.

Starting from such ME results, it is possible to define several observables built on the ratios of ME. One of particular interest for this study is the *optimal observables* (denoted $\mathcal{OO}$) of first order in the BSM amplitude, defined as

$$\mathcal{OO}_1(\mathbf{c}) = \frac{\text{Interference term}}{|\mathcal{M}_{\text{SM}}|^2} = \frac{2\Re(\mathcal{M}_{\text{SM}}\mathcal{M}_{\text{BSM}}^*(\mathbf{c}))}{|\mathcal{M}_{\text{SM}}|^2}. \quad (9.2)$$

The distribution of this observable has the interesting property of being symmetric for a Standard Model-like set of events, and asymmetric if a CP -odd contribution is present. It should be noted that $\mathcal{M}_{\text{BSM}}$ and the $\mathcal{OO}$ depend on a parameter $\mathbf{c}$ which parametrises a BSM *hypothesis* under which the matrix elements are computed. As the matrix elements are generated using the Higgs Characterisation (HC) model [206], $\mathbf{c}$ will take the form of a set of HC couplings modifiers. This set of couplings modifiers can equivalently take the form of a set of Wilson coefficients from the Higgs basis (HB) or Warsaw basis (WB) since the translation HB/WB $\rightarrow$ HC is straightforward (see Chapter 7.3.6). For instance, in the first examples given in this Section, the simple $\mathbf{c} = \{\tilde{c}_{zz} = 2.65\}$ is chosen (where $\tilde{c}_{zz}$ is the Wilson coefficient for the CP -odd HZZ operator, see Chapter 7.3) and it will be assumed that this Higgs basis Wilson coefficient is implicitly translated into its corresponding set of HC couplings modifiers. The impact of the choice of this set of coefficients is discussed in detail in Section 9.1.5 below.

However, matrix elements are not observable physical quantities, only their squared amplitude is experimentally accessible. As a consequence, the matrix elements computed by MADGRAPH are only provided as squared amplitudes, from which the interference term must be deduced. This is done by computing the squared matrix elements for the SM process, the BSM process, and the process involving both of them (called *Mix* in the following). Since the mixed ME is defined as

$$\mathcal{M}_{\text{Mix}}(\mathbf{c}) = \mathcal{M}_{\text{SM}} + \mathcal{M}_{\text{BSM}}(\mathbf{c}), \quad (9.3)$$

the squared amplitude reads

$$\begin{aligned} |\mathcal{M}_{\text{Mix}}(\mathbf{c})|^2 &= (\mathcal{M}_{\text{SM}} + \mathcal{M}_{\text{BSM}}(\mathbf{c}))(\mathcal{M}_{\text{SM}}^* + \mathcal{M}_{\text{BSM}}^*(\mathbf{c})) \\ &= |\mathcal{M}_{\text{SM}}|^2 + 2\Re(\mathcal{M}_{\text{SM}}\mathcal{M}_{\text{BSM}}^*(\mathbf{c})) + |\mathcal{M}_{\text{BSM}}(\mathbf{c})|^2, \end{aligned} \quad (9.4)$$

from which the interference term can be extracted:

$$2\Re(\mathcal{M}_{\text{SM}}\mathcal{M}_{\text{BSM}}^*(\mathbf{c})) = |\mathcal{M}_{\text{Mix}}(\mathbf{c})|^2 - |\mathcal{M}_{\text{SM}}|^2 - |\mathcal{M}_{\text{BSM}}(\mathbf{c})|^2. \quad (9.5)$$

Replacing in Eq. 9.2, the first optimal observable can be computed as

$$\mathcal{OO}_1(\mathbf{c}) = \frac{|\mathcal{M}_{\text{Mix}}(\mathbf{c})|^2 - |\mathcal{M}_{\text{SM}}|^2 - |\mathcal{M}_{\text{BSM}}(\mathbf{c})|^2}{|\mathcal{M}_{\text{SM}}|^2}. \quad (9.6)$$

An example distribution of this variable can be seen in Figure 9.3a for the SM and several values of BSM couplings. As expected, the distribution for samples including a CP -odd coupling (red and blue) show an asymmetric tail, extending toward negative values for negative value of the couplings, and toward positive values for positive value of the coupling. The distributions for samples including a BSM CP -even coupling remain symmetric (green and orange), with larger tails on both sides. In both cases, the larger the couplings, the larger the tails.

A second observable can be defined (second order in BSM amplitude optimal observable $\mathcal{OO}_2$) as the simple ratio of the pure BSM contribution matrix element normalised by the SM matrix element:

$$\mathcal{OO}_2(c) = \frac{|\mathcal{M}_{\text{BSM}}(c)|^2}{|\mathcal{M}_{\text{SM}}|^2}, \quad (9.7)$$

which is often represented after taking its logarithm. This observable does not present an asymmetry in case of CP -odd contribution and therefore cannot be used to distinguish between CP -even and CP -odd contributions. However, it brings additional information regarding the magnitude of the coupling. Example distributions are shown in Figure 9.3b for various BSM couplings values. The expected behaviour is that for a given Wilson coefficient, increasing values increases the $\mathcal{M}_{\text{BSM}}$ contribution so the $\mathcal{OO}_2$ shifts toward higher values. This is verified as the $c_{zz} = -0.24$ sample (orange) gets higher tails on the positive side, and the $c_{zz} = 1.2$ (green) sample is largely shifted with respect to the SM sample. The two CP -odd samples (red and blue) have similar distributions, as expected given they both have $|\tilde{c}_{zz}| = 1.3$, and are shifted toward positive values. The impact for $c_{zz} = 1.2$ is larger than for $\tilde{c}_{zz} = 1.3$ due to different coupling constants in front of their respective operator.

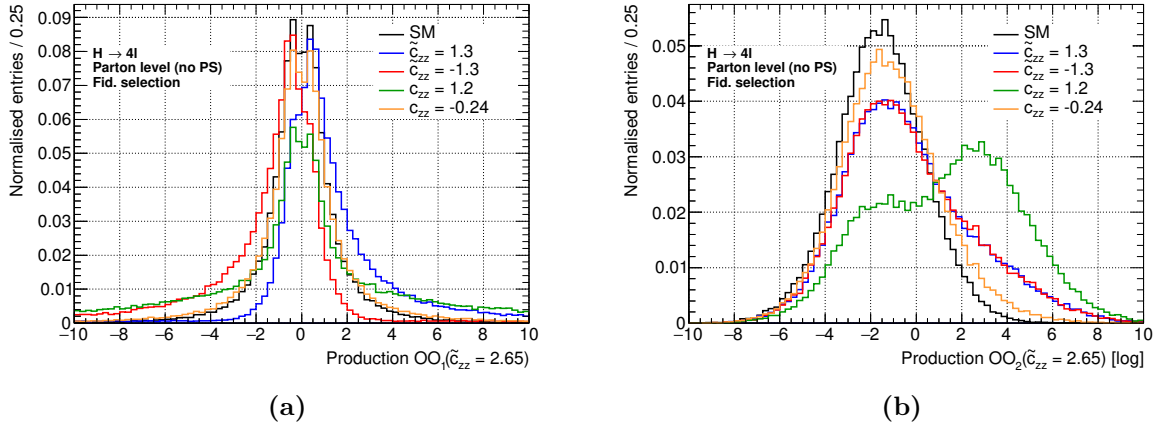


Figure 9.3 – Distribution of (a) first order and (b) second order optimal observable for the process $pp \rightarrow jjH$ (normalised to unit integral) in truth-level MADGRAPH samples generated assuming a pure SM hypothesis (black), or SM with an additional CP -odd contribution (blue: $\tilde{c}_{zz} = 1.3$, red: $\tilde{c}_{zz} = -1.3$), or SM with an additional CP -even contribution (green: $c_{zz} = 1.2$, orange: $c_{zz} = -0.24$). For $\mathcal{OO}_1$, a CP -odd contribution makes the shape looking asymmetric, while a CP -even contribution leaves it symmetric. In $\mathcal{OO}_2$, only the absolute amplitude of the BSM contribution can be determined, couplings with opposite value yielding close distributions, *e.g.* $\tilde{c}_{zz} = 1.3$ and $\tilde{c}_{zz} = -1.3$.

9.1.3 Technical generation of matrix elements with MadGraph

While MADGRAPH is a versatile and powerful tool for event generation and matrix element computation, its settings must be carefully checked in order to ensure the validity of matrix element computations. This Section and the following describe the general method used to compute the matrix elements entering the optimal observables and the tests performed to ensure

the correctness of the results. It should be noted that only leading order matrix element computing code can be generated at the moment.

With MADGRAPH, such code can be generated with just a few lines. For example, to generate the leading order matrix element computing code for the process of Higgs boson production through the VBF mode in the presence of BSM operators, the lines listed in Listing 9.1 are enough.

Listing 9.1 – Generation of a leading order VBF process without Higgs boson decay in the presence of BSM operators using HC with MADGRAPH. `$$ a z w+ w- QCD=0` imposes that no vector boson can appear in the s -channel, `QCD=0` imposes that no QCD vertex can appear in the process.

```
import model HC_UFO
generate p p > x0 j j $$ a z w+ w- QCD=0
output standalone_cpp procVBF
```

This generates a set of `c++` classes, each of them capable of computing a single partonic process (for instance $u\bar{u} \rightarrow u\bar{u}H$). Since for a given collision it is impossible to know what partons effectively collided, these *partonic matrix elements* must be combined in order to have the correct result. To get the total matrix element, they must be summed, each of them being weighted by the PDF of their partons:

$$\mathcal{M}(pp \rightarrow jjH) = \sum_{i,j,k,l} f_i(x_1) f_j(x_2) \mathcal{M}(ij \rightarrow klH), \quad (9.8)$$

where the sum runs over all possible combination of incoming (i, j) and outgoing (k, l) partons, f_i and f_j are the PDF for partons i and j , and x_1 and x_2 are the proton momentum fraction carried by each partons. The latter is not directly experimentally measured (as above the exact colliding partons cannot be determined), but can be inferred from the outgoing kinematics through

$$x_{1,2} = \frac{m_{4\ell,jj}}{\sqrt{s}} \cdot \exp(\pm y_{4\ell,jj}), \quad (9.9)$$

where $\sqrt{s}$ is the center-of-mass energy of the collision (*i.e.* 13 TeV), $m_{4\ell,jj}$ is the invariant mass of the system formed by the four leptons and the two jets, and $y_{4\ell,jj}$ is the rapidity of this system. The PDF are provided by the LHAPDF-6 package [264], and the dependency of the matrix element result on the PDF choice is studied in Section 9.1.4.

In addition to generating the VBF process without the subsequent Higgs boson decay (hereafter referred to as *production only* ME or $\mathcal{OO}$ and denoted $\mathcal{OO}_{1,jj}$), other processes of interest for the $H \rightarrow 4\ell$ study are the *decay only* $\mathcal{OO}$ ($\mathcal{OO}_{1,4\ell}$), which can be generated using the snippet of Listing 9.2, and the total matrix element (referred to as *production+decay* ME and $\mathcal{OO}$ and denoted $\mathcal{OO}_{1,4\ell,jj}$). Assuming electrons and muons are massless, the former only produces 2 computing classes: one for $H \rightarrow 4\mu/4e$ and the other for $H \rightarrow 2e2\mu$. The production+decay process on the other hand creates a large number of classes that must be handled in the same way as the production only process: using PDF to weight the sum of all partonic matrix elements.

Listing 9.2 – Generation of a leading order Higgs boson decay to four leptons in the presence of BSM operators using HC with MADGRAPH.

```
import model HC_UFO
generate x0 > l+ l- l+ l-
output standalone_cpp procH4l
```

Listing 9.3 – Generation of a leading order VBF process with Higgs boson decay in the presence of BSM operators using HC with MADGRAPH.

```
import model HC_UFO
generate p p > x0 j j $$ a z w+ w- QCD=0, x0 > l+ l- l+ l-
output standalone_cpp procVBF4l
```

9.1.4 Dependence on the PDF choice

As seen in Eq. 9.8, the summation of the partonic matrix elements depends on the choice of the PDF used as a weighting function. The PDFs are retrieved using the LHAPDF-6 library [264], which offers a large choice of PDF sets. The default PDF used in MADGRAPH is the NNPDF 2.3 computed at leading order with $\alpha_s = 0.130$ (later referred to by its LHAPDF index: 247000).

To assess the dependence of the total ME and $\mathcal{OO}$ computation on the PDF, several other sets are used. The impact of changing the PDF on the total ME evaluation is shown in Figure 9.4a and display a systematic bias toward higher ME values comparing any other PDF set to the default. This bias is however less than 10 % across the ME range. Moreover, while this bias is not negligible, it has a similar behaviour for all ME hypotheses (SM, pure BSM and Mix) and vanishes in the ME ratios (Eqs. 9.6 and 9.7). Consequently such a bias is not observed on the optimal observables as in Figure 9.4b.

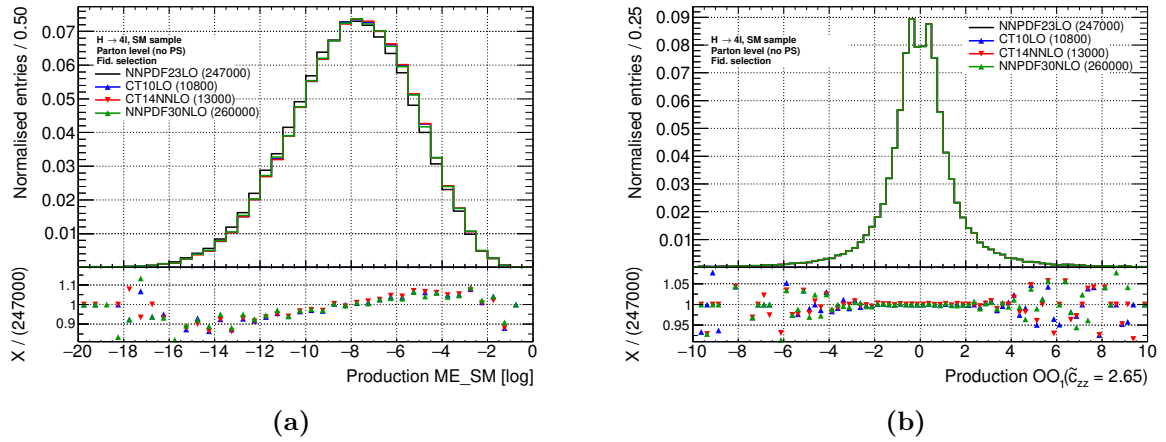


Figure 9.4 – Distribution of (a) $|\mathcal{M}_{\text{SM}}|^2$ and (b) $\mathcal{OO}_1(\tilde{c}_{zz} = 2.65)$ computed for the process $pp \rightarrow jjH$ in a SM-like generated sample. The distributions are shown for various choices of PDFs used in the weighting of the partonic processes of Eq. 9.8. The codes correspond to the LHAPDF-6 set index. For $|\mathcal{M}_{\text{BSM}}(\tilde{c}_{zz} = 2.65)|^2$ and $|\mathcal{M}_{\text{Mix}}(\tilde{c}_{zz} = 2.65)|^2$, the results are similar to $|\mathcal{M}_{\text{SM}}|^2$, and the impact on $\mathcal{OO}_2(\tilde{c}_{zz} = 2.65)$ is similar to the one on $\mathcal{OO}_1(\tilde{c}_{zz} = 2.65)$. Similar results arise for samples other than SM [265].

Moreover, as seen in Chapter 2.9, the PDF evaluation depends on the momentum transfer Q^2 inside the collision. Since the processes studied involve the production of a Higgs boson with $m_H \approx 125 \text{ GeV}$, a natural choice of Q^2 is m_H^2 , and the tests on PDF set dependence were carried out using this value. However, the total momentum transfer in a VBF event, represented by the invariant of the system composed of the four leptons and the two jets $m_{4\ell jj}$ can reach on average hundreds of GeV.

In order to assess for a possible dependency on the chosen Q^2 value, a similar test as above is performed selecting the default PDF set (247000) but changing the Q^2 value used in the evaluation. The results of this study are presented in Figure 9.5. The impact of changing the Q^2 on the ME computation is within 5 %, slightly shifting the ME result toward lower values (only $|\mathcal{M}_{\text{SM}}|^2$ is shown in Figure 9.5a but $|\mathcal{M}_{\text{BSM}}(\tilde{c}_{zz})|^2$ and $|\mathcal{M}_{\text{Mix}}(\tilde{c}_{zz})|^2$ show similar results), and completely cancels in the ratios used to build the optimal observables as seen in Figure 9.5b (only $\mathcal{OO}_1(\tilde{c}_{zz})$ is shown in Figure 9.5b but a similar behaviour is observed for $\mathcal{OO}_2(\tilde{c}_{zz})$).

Other reasonable choices of Q^2 would be to choose the event-dependent four-lepton invariant mass, or since the main interest is for VBF production, the invariant mass of the system formed by the four leptons and the two jets. However given the negligible difference of the result for the optimal observables between $Q^2 = m_Z^2$ and $Q^2 = (1 \text{ TeV})^2$, such possibility is not expected to

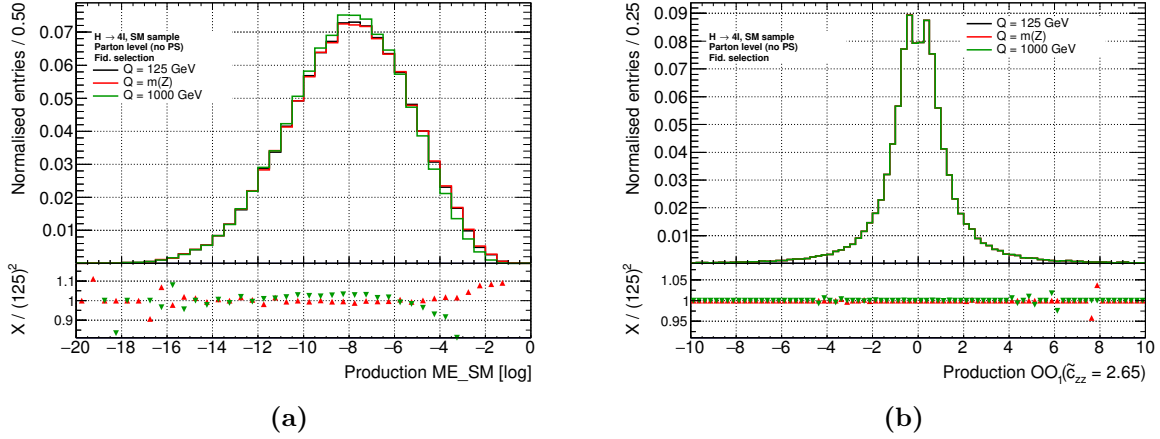


Figure 9.5 – Distribution of (a) $|\mathcal{M}_{\text{SM}}|^2$ and (b) $\mathcal{OO}_1(\tilde{c}_{zz} = 2.65)$ for the process $pp \rightarrow jjH$, and for various choices of Q^2 in the evaluation of the PDFs used in the weighting of the partonic processes of Eq. 9.8. The PDF set is constant across the three trials and set to the default (247000). For $|\mathcal{M}_{\text{BSM}}(\tilde{c}_{zz})|^2$ and $|\mathcal{M}_{\text{Mix}}(\tilde{c}_{zz})|^2$, the results are similar to $|\mathcal{M}_{\text{SM}}|^2$, and the impact on $\mathcal{OO}_2(\tilde{c}_{zz})$ is similar to the one on $\mathcal{OO}_1(\tilde{c}_{zz})$. Similar results arise for samples other than SM [265].

have a larger impact and was therefore not studied.

It is concluded that neither the choice of PDF nor the choice of Q^2 used for the weighting in the partonic ME summation has a sizeable impact on the $\mathcal{OO}$ computation. As a consequence, the default PDF is chosen (NNPDF 2.3 computed at leading order with $\alpha_s = 0.130$, code 247000) and is evaluated with $Q^2 = (125 \text{ GeV})^2 \approx m_H^2$. The results in this Section were only shown for ME and $\mathcal{OO}$ computed for the process $pp \rightarrow jjH$ and on a SM-like sample, but similar results are obtained for other processes ($H \rightarrow 4\ell$ or $pp \rightarrow jjH(\rightarrow 4\ell)$) and applied on non-SM samples.

9.1.5 Impact of changing the hypothesis provided to the matrix element computing

As seen in Section 9.1.2, the definition of optimal observables is parametrised by a set of BSM couplings $\mathbf{c}$ (the *hypothesis*). In the studies above, the set of couplings $\mathbf{c} = \{\tilde{c}_{zz} = 2.65\}$ was assumed. In this Section, the impact of changing the hypothesis given to the ME computation is studied. Such a verification is important because if two observables with different hypotheses give different discriminating power of CP -odd contribution against the SM case, this might require an extra step to optimise the hypothesis used to compute the observable.

Figure 9.6 shows the $\mathcal{OO}_1$ distribution for the process $pp \rightarrow jjH$ for two sets of BSM couplings provided to the computing classes during the ME evaluation: $\tilde{c}_{zz} = 1$ and $\tilde{c}_{zz} = 2.65$. In Figure 9.6a, the two distributions appear different: the distribution with $\tilde{c}_{zz} = 2.65$ has much wider spread than the distribution with $\tilde{c}_{zz} = 1$. However, the ratio of the spread is (almost) exactly 2.65, and by dilating the distribution (equivalent to a rebinning and a redefinition of the range), the two distributions become close in shape, as seen in Figure 9.6b. Residual differences are of the order of 2% as seen by the ratio of the two distributions after rebinning, which is likely to stem from numerical errors of the calculations.

From this study it can be concluded that for a given couplings axis (*e.g.* $\tilde{c}_{zz}$), the value of the coupling used in the $\mathcal{M}_{\text{BSM}}$ generation can be chosen arbitrarily: the distribution for a given value is equivalent to the distribution for another coupling value after a shift and a dilation, *i.e.* after rebinning the distribution. A similar result can be obtained for $\log(\mathcal{OO}_2)$, for which the distribution for two given coupling values are equivalent after a simple shift. As a consequence,

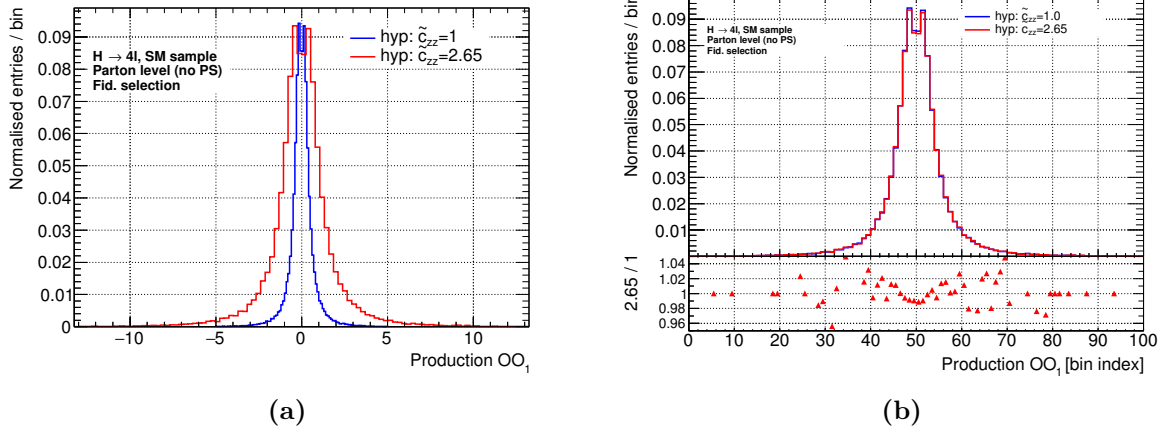


Figure 9.6 – Comparison of $\mathcal{O}_1$ computation for the process $pp \rightarrow jjH$ using either the $\tilde{c}_{zz} = 1$ (blue) or $\tilde{c}_{zz} = 2.65$ (red) hypotheses, and applied on a SM-like generated sample. In (a), both distributions are represented using 100 bins, distributed over the range $[-13.25, 13.25]$ for hypothesis $\tilde{c}_{zz} = 2.65$ and over the range $[-5, 5]$ for hypothesis $\tilde{c}_{zz} = 1$ (the range lengths have a ratio of 2.65). (b) shows a bin-by-bin comparison of the two distributions: the distributions are represented as a function of their bin index from (a).

the discriminating power for BSM contribution will be the same whatever the hypothesis used to build the $\mathcal{O}$, after adequate rebinning. However, the $\mathcal{O}$ computed along different axes (for instance along $\tilde{c}_{zz}$ and along $\tilde{c}_{z\gamma}$) yield non-equivalent shapes, and the discriminating power might differ between observables built along two different axes.

9.1.6 Development of a MadGraph-based simplified matrix element computation package

During the generation of the matrix element computation classes with MADGRAPH using the snippet of Listing 9.1 (process $pp \rightarrow jjH$), it is assumed that quarks are massless and that the colliding protons are only composed of gluons, u -, d -, s - and c -quarks and antiquarks. With these assumptions, symmetries appear in the processes and only 77 computing classes (*i.e.* 77 *partonic processes*) are necessary to get the correct result after the PDF-weighted summation. For example, with these assumptions, the $u\bar{u} \rightarrow u\bar{u}H$ and $c\bar{c} \rightarrow c\bar{c}H$ partonic processes are identical and a common class can be used to compute both of them. The resulting matrix element of this partonic process still have to be weighted by both the u - and c -quark PDFs during the combination.

The evaluation of matrix elements is however CPU-intensive, and the computation of the observables ends up taking a large fraction of the simulated data processing time. Consequently, the possibility of removing some of the computing classes has been investigated. It should also be noted that due to internal technicalities, MADGRAPH can output up to two matrix elements for the same process¹ (called *sub-ME* hereafter). The output of each computing class and for each of the two sub-ME is shown in Figure 9.7 for a given event. From this plot it appears that two groups of eight processes (one for the first sub-ME, the other for the second sub-ME) have a similar value, so it might be possible to only compute one of them, then adequately weighted. The observed behaviour is moreover consistent for all three hypotheses needed to compute the $\mathcal{O}$ (SM, BSM, Mix). The validity of such simplification is assessed in the following, focusing on the ME computation with SM hypothesis (this extends to the other hypotheses straightforwardly).

¹↑ In short, this is due to the reuse of the same computation for the t - and u -channels thanks to crossing symmetry. In case only one channel is physical meaningful, the second matrix element result is always set to 0 in the computation. Further details are not relevant for the present study.

The first group of eight partonic processes for $pp \rightarrow jjH$ is composed of

$$\begin{array}{llll} cd \rightarrow Hus, & cs \rightarrow Hcs, & \bar{c}\bar{d} \rightarrow H\bar{u}\bar{s}, & \bar{c}\bar{s} \rightarrow H\bar{c}\bar{s}, \\ ud \rightarrow Hud, & us \rightarrow Hcd, & \bar{u}\bar{d} \rightarrow H\bar{u}\bar{d}, & \bar{u}\bar{s} \rightarrow H\bar{c}\bar{d}, \end{array}$$

yielding a value around 0.05×10^{-3} , while the second is composed of

$$\begin{array}{llll} c\bar{c} \rightarrow Hs\bar{s}, & c\bar{u} \rightarrow Hs\bar{d}, & d\bar{d} \rightarrow Hu\bar{u}, & d\bar{s} \rightarrow Hu\bar{c}, \\ s\bar{s} \rightarrow Hc\bar{u}, & s\bar{s} \rightarrow Hc\bar{c}, & u\bar{c} \rightarrow Hd\bar{s}, & u\bar{u} \rightarrow Hd\bar{d}, \end{array}$$

yielding a value around 0.056×10^{-3} . Looking at the diagrams associated to each of these partonic processes, it appears that the first (resp. second) group is composed of diagrams exchanging a W boson in the u - (resp. t -)channel. More specifically, for both groups, the diagrams correspond to W boson exchanges in which none of the W boson changes the quark family between incoming and the outgoing quark ($u \leftrightarrow d$ and $c \leftrightarrow s$ but no $u \leftrightarrow s$ nor $c \leftrightarrow d$). This behaviour can be checked for all events, and it can be shown that each of these two groups can be represented with only one partonic process among the eight [266].

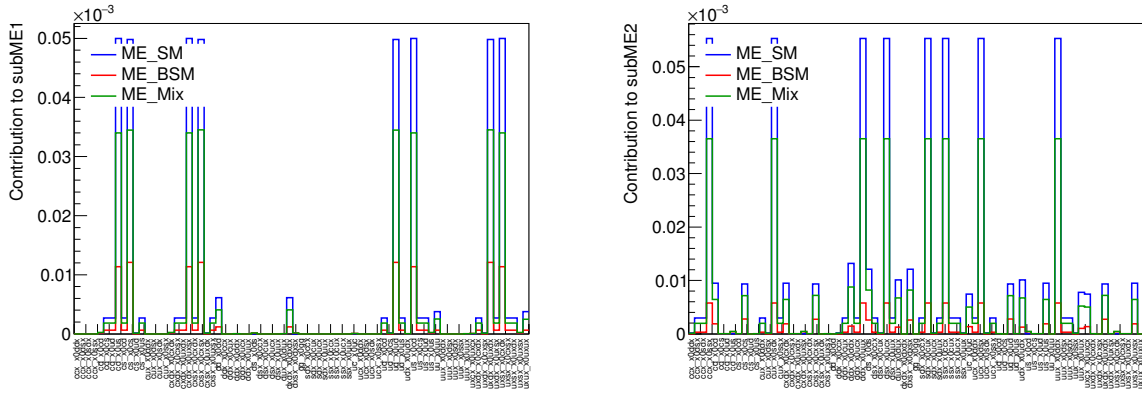


Figure 9.7 – Breakdown of the matrix element result returned by each partonic process computing class from the process defined in Listing 9.1 for one event. The left hand side plot show the values for the first sub-ME, the right hand side plot for the second sub-ME. The y -scale is slightly different between the two plots. The values for the ME computation under the SM hypothesis is represented with the blue histogram, under a pure BSM hypothesis ($\tilde{c}_{zz} = 2.65$) with the red histogram, and under the mix hypothesis with the green histogram. All three histograms display the same features, with an overall scaling between them.

This reasoning can be extended to other structures seen in the plot: for example, another population of 32 matrix element computations with values around 0.003×10^{-3} arise for the SM hypothesis in this specific event. A more zoomed plot would show as previously that these can be split into two subgroups with slightly different computed value, and that they also correspond to partonic processes with exchange of W bosons (one subgroup of 8 for t -channel exchange, the other subgroup of 8 for u -channel), this time one of them changing the quark family. This contribution is expected to be suppressed by a factor of $\sin^2 \theta_C \approx 0.22^2$ due to the quark family transition (θ_C is the Cabbibo angle). The observed ratio between the leading and subleading W boson exchange contributions qualitatively matches the expectation, validating the approach. These two subgroup can thus also be represented by only one partonic process each.

Another similar group of twice 8 additional contributions arise, but at a too low level to be seen in from the plot. This contribution is suppressed by a factor $\approx \sin^4 \theta_C$ and correspond to partonic processes with W boson exchange in which both W bosons change the quark family. It could also be represented by two additional processes only, but given its high suppression factor, their contribution is dropped from the total contribution.

Finally, one last group of 20 ME remains, corresponding to computing classes yielding a result between 0.006×10^{-3} to 0.013×10^{-3} for the SM hypothesis in this event. The diagrams associated to these processes correspond to the exchange of Z bosons, but the spread is larger than in the case of W boson exchange². These can nevertheless be represented by 5 partonic processes instead of 20.

The study presented here takes the SM sample as an example but the same results were obtained with other samples and other matrix element hypotheses. As a consequence, it is possible to represent the complete ME computation with only 9 subprocesses computing classes, down from the 77 generated par MADGRAPH. An eventual bias from removing some of the contributions has been checked and is shown in Figure 9.8. While removing from the final ME computation all partonic processes contributing with a Z boson exchange or a W boson exchange with one or two quark family transition (or any combination of such removal) obviously biases the final ME value (Figure 9.8a), the bias cancels in the ME ratios defining the optimal observables (Figure 9.8b). The choice of removing only the diagrams contributing with a double quark family transition W boson exchange is validated and thus adopted.

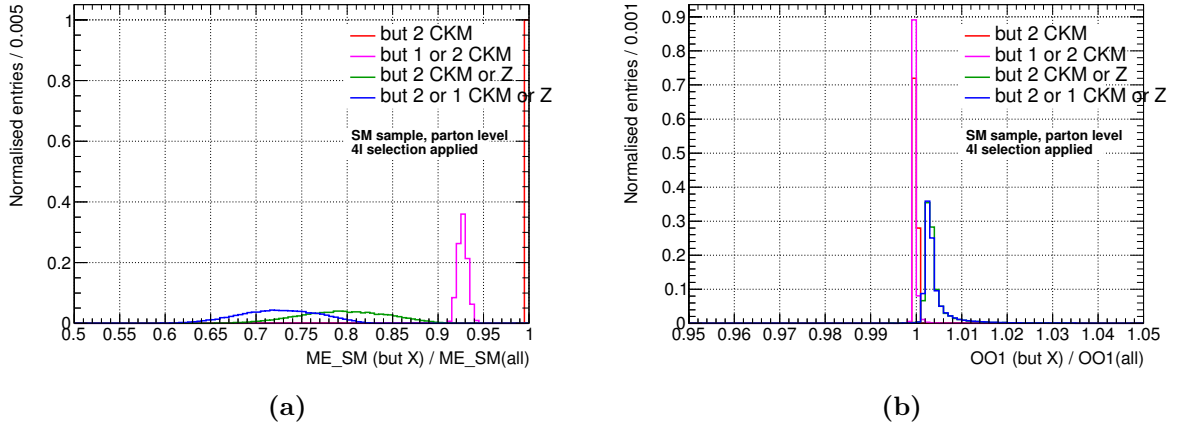


Figure 9.8 – Bias introduced on (a) $|\mathcal{M}_{\text{SM}}|^2$ and (b) $\mathcal{OO}_1(\tilde{c}_{zz} = 2.65)$ for the process $pp \rightarrow jjH$ by removing all partonic processes with diagrams contributing by a W boson exchange with two quark family transitions (red, no bias seen), all partonic processes with diagrams contributing by a W boson exchange with one or two quark family transitions (pink, 7 % bias and small spread on ME computation but compensated on $\mathcal{OO}$ computation), all partonic processes with diagrams contributing with a two quark family transitions W boson exchange or a Z boson exchange (green, 20 % bias on ME with large spread, but less than 0.5 % bias and small spread on $\mathcal{OO}$ computation), all sub-subprocesses with diagrams contributing with one or two quark family transitions W boson exchange or a Z boson exchange (blue, 25 % bias on ME with large spread, but less than 0.5 % bias and small spread on $\mathcal{OO}$ computation). The biases in $|\mathcal{M}_{\text{BSM}}|^2$ and $|\mathcal{M}_{\text{Mix}}|^2$ are similar to the bias in $|\mathcal{M}_{\text{SM}}|^2$, and the bias on $\mathcal{OO}_2$ is similar to the bias on $\mathcal{OO}_1$.

The drawback of this method is its loss of generality: for each new generated process, one has to explore the values for all matrix element generated to spot similarities, select one representative matrix element per similar group, and assign the correct weight structure to each of the selected ME. However, after such setup is performed, a significant improvement in computing time is achievable (factor ~ 8 in this case) from the reduced number of computing classes actually used.

²↑ Differences are seen between processes that can occur *e.g.* through both Z and W boson exchange (such as $cd \rightarrow Hcd$ in the t - and u -channels respectively) and those who can occur only through Z boson exchange ($uu \rightarrow Huu$). Details of this discussion are however beyond the scope of this Thesis.

9.2 Morphing technique

In order to evaluate the compatibility between the Standard Model and a certain point in the phase-space, the shape of the observable used to discriminate between the two must be generated at this particular phase-space point. Consequently, a Monte-Carlo simulation (MC) has to be produced at this point of the couplings phase-space. Assuming the number of couplings of interest is limited to one, the number of points to generate is proportional to the number of compatibility tests to perform. However, as the dimensionality of the target phase-space grows (*i.e.* the number of couplings of interest increases), the number of required points grows exponentially, so the required number of MC to generate becomes quickly unmanageable. As a consequence, a method to extrapolate between a relatively small number of points to cover for the full phase-space of interest is needed.

Several methods can be elaborated, the one used for this analysis being the *morphing technique*. This method was first introduced in Ref. [267], and adapted to the case of EFT in Ref. [268]. A brief description of this method in the context of EFT is given in Section 9.2.1, while Section 9.2.2 validates its usage for the present analysis.

9.2.1 Mathematical definition of the morphing method

The basic assumption for the morphing method is that the morphed quantity T (which can be a cross-section as well as distribution, *e.g.* a differential cross-section) must directly relate to the matrix elements:

$$T(\mathbf{g}) \propto |\mathcal{M}(\mathbf{g})|^2 \quad \text{with} \quad \mathcal{M}(\mathbf{g}) = \sum_i g_i \mathcal{M}_i, \quad (9.10)$$

where $\mathbf{g}$ is a set of couplings $\{g_1, \dots, g_i, \dots\}$ corresponding to a set of matrix elements $\mathcal{M}_i$. The squared matrix element $|\mathcal{M}(\mathbf{g})|^2$ is therefore a multidimensional polynomial in the g_i . An additional assumption is that the target distribution $T_{\text{target}}(\mathbf{g})$ can be written as a linear combination of some input distributions $T_{\text{input},j}$:

$$T_{\text{target}}(\mathbf{g}) = \sum_j w_j(\mathbf{g}) T_{\text{input},j}. \quad (9.11)$$

The $T_{\text{input},j}$ will often be given by the MC simulation. Given Eq. 9.10, the weights $w_j(\mathbf{g})$ are therefore also multidimensional polynomials in the couplings g_i .

The analysis presented in this Chapter uses EFT coefficients to interpret eventual modification of the Higgs boson couplings and focuses on the VBF production mode and the $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ decay channel. Moreover, it is assumed that only one EFT coefficient is modified at a time. This reduces to a single BSM parameter appearing in the morphing, but this parameter appears twice: once in the production side, and once in the decay side. In this case, $\mathbf{g}$ is limited to $\{g_{\text{SM}}, g_{\text{BSM}}\}$ where g_{BSM} is computed from the value of the Wilson coefficient being evaluated: $\tilde{c}_{zz}$, or $\tilde{c}_{z\gamma}$, or *etc.* The matrix element for the VBF $H \rightarrow 4\ell$ process reads

$$\begin{aligned} \mathcal{M}_{\text{Mix}}^{pd}(g_{\text{SM}}, g_{\text{BSM}}) &= \mathcal{M}_{\text{Mix}}^p(g_{\text{SM}}, g_{\text{BSM}}) \cdot \mathcal{M}_{\text{Mix}}^d(g_{\text{SM}}, g_{\text{BSM}}) \\ &= (g_{\text{SM}} \mathcal{M}_{\text{SM}}^p + g_{\text{BSM}} \mathcal{M}_{\text{BSM}}^p) \cdot (g_{\text{SM}} \mathcal{M}_{\text{SM}}^d + g_{\text{BSM}} \mathcal{M}_{\text{BSM}}^d). \end{aligned} \quad (9.12)$$

Squaring $\mathcal{M}_{\text{Mix}}^{pd}(g_{\text{SM}}, g_{\text{BSM}})$:

$$\begin{aligned}
|\mathcal{M}_{\text{Mix}}^{pd}(g_{\text{SM}}, g_{\text{BSM}})|^2 &= |\mathcal{M}_{\text{Mix}}^p(g_{\text{SM}}, g_{\text{BSM}})|^2 \cdot |\mathcal{M}_{\text{Mix}}^d(g_{\text{SM}}, g_{\text{BSM}})|^2 \\
&= (|g_{\text{SM}}\mathcal{M}_{\text{SM}}^p + g_{\text{BSM}}\mathcal{M}_{\text{BSM}}^p|^2) (|g_{\text{SM}}\mathcal{M}_{\text{SM}}^d + g_{\text{BSM}}\mathcal{M}_{\text{BSM}}^d|^2) \\
&= g_{\text{SM}}^4 |\mathcal{M}_{\text{SM}}^p|^2 |\mathcal{M}_{\text{SM}}^d|^2 \\
&\quad + g_{\text{SM}}^3 g_{\text{BSM}} \cdot (|\mathcal{M}_{\text{SM}}^p|^2 \Re(\mathcal{M}_{\text{SM}}^d \mathcal{M}_{\text{BSM}}^{d*}) + \Re(\mathcal{M}_{\text{SM}}^p \mathcal{M}_{\text{BSM}}^{p*}) |\mathcal{M}_{\text{SM}}^d|^2) \\
&\quad + g_{\text{SM}}^2 g_{\text{BSM}}^2 \cdot (|\mathcal{M}_{\text{SM}}^p|^2 |\mathcal{M}_{\text{BSM}}^d|^2 + |\mathcal{M}_{\text{SM}}^p|^2 |\mathcal{M}_{\text{BSM}}^d|^2) \\
&\quad + g_{\text{SM}} g_{\text{BSM}}^3 \cdot (|\mathcal{M}_{\text{BSM}}^p|^2 \Re(\mathcal{M}_{\text{SM}}^d \mathcal{M}_{\text{BSM}}^{d*}) + \Re(\mathcal{M}_{\text{SM}}^p \mathcal{M}_{\text{BSM}}^{p*}) |\mathcal{M}_{\text{BSM}}^d|^2) \\
&\quad + g_{\text{BSM}}^4 |\mathcal{M}_{\text{BSM}}^p|^2 |\mathcal{M}_{\text{BSM}}^d|^2. \tag{9.13}
\end{aligned}$$

This is a fourth order polynomial with respect to the g_i , with 5 independent terms. To produce a $T_{\text{target}}(g_{\text{SM}}, g_{\text{BSM}}) \propto |\mathcal{M}_{\text{Mix}}^{pd}(g_{\text{SM}}, g_{\text{BSM}})|^2$, one therefore need 5 base samples, each of them verifying $T_{\text{input},j}(g_{\text{SM},j}, g_{\text{BSM},j}) \propto |\mathcal{M}_{\text{Mix}}^{pd}(g_{\text{SM},j}, g_{\text{BSM},j})|^2$ for $j \in [1, 5]$. Typically for this study, the input distributions will be built with

$$\left\{ (g_{\text{SM}}, g_{\text{BSM}}(\tilde{c}_{zz} = -2)), (g_{\text{SM}}, g_{\text{BSM}}(\tilde{c}_{zz} = -1)), (g_{\text{SM}}, g_{\text{BSM}}(\tilde{c}_{zz} = 0)), \right. \\
\left. (g_{\text{SM}}, g_{\text{BSM}}(\tilde{c}_{zz} = 1)), (g_{\text{SM}}, g_{\text{BSM}}(\tilde{c}_{zz} = 2)) \right\}. \tag{9.14}$$

From this, the morphing function can be written

$$\begin{aligned}
T_{\text{target}}(g_{\text{SM}}, g_{\text{BSM}}) &= \underbrace{(a_{11}g_{\text{SM}}^4 + a_{12}g_{\text{SM}}^3g_{\text{BSM}} + a_{13}g_{\text{SM}}^2g_{\text{BSM}}^2 + a_{14}g_{\text{SM}}g_{\text{BSM}}^3 + a_{15}g_{\text{BSM}}^4)}_{w_1(g_{\text{SM}}, g_{\text{BSM}})} T_{\text{input},1}(g_{\text{SM},1}, g_{\text{BSM},1}) \\
&\quad + \underbrace{(a_{21}g_{\text{SM}}^4 + a_{22}g_{\text{SM}}^3g_{\text{BSM}} + a_{23}g_{\text{SM}}^2g_{\text{BSM}}^2 + a_{24}g_{\text{SM}}g_{\text{BSM}}^3 + a_{25}g_{\text{BSM}}^4)}_{w_2(g_{\text{SM}}, g_{\text{BSM}})} T_{\text{input},2}(g_{\text{SM},2}, g_{\text{BSM},2}) \\
&\quad + \underbrace{(a_{31}g_{\text{SM}}^4 + a_{32}g_{\text{SM}}^3g_{\text{BSM}} + a_{33}g_{\text{SM}}^2g_{\text{BSM}}^2 + a_{34}g_{\text{SM}}g_{\text{BSM}}^3 + a_{35}g_{\text{BSM}}^4)}_{w_3(g_{\text{SM}}, g_{\text{BSM}})} T_{\text{input},3}(g_{\text{SM},3}, g_{\text{BSM},3}) \\
&\quad + \underbrace{(a_{41}g_{\text{SM}}^4 + a_{42}g_{\text{SM}}^3g_{\text{BSM}} + a_{43}g_{\text{SM}}^2g_{\text{BSM}}^2 + a_{44}g_{\text{SM}}g_{\text{BSM}}^3 + a_{45}g_{\text{BSM}}^4)}_{w_4(g_{\text{SM}}, g_{\text{BSM}})} T_{\text{input},4}(g_{\text{SM},4}, g_{\text{BSM},4}) \\
&\quad + \underbrace{(a_{51}g_{\text{SM}}^4 + a_{52}g_{\text{SM}}^3g_{\text{BSM}} + a_{53}g_{\text{SM}}^2g_{\text{BSM}}^2 + a_{54}g_{\text{SM}}g_{\text{BSM}}^3 + a_{55}g_{\text{BSM}}^4)}_{w_5(g_{\text{SM}}, g_{\text{BSM}})} T_{\text{input},5}(g_{\text{SM},5}, g_{\text{BSM},5}). \tag{9.15}
\end{aligned}$$

Then, by requiring that $T_{\text{target}}(g_{\text{SM}}, g_{\text{BSM}}) = T_{\text{input},j}(g_{\text{SM},j}, g_{\text{BSM},j})$ for $(g_{\text{SM}}, g_{\text{BSM}}) = (g_{\text{SM},j}, g_{\text{BSM},j})$, one obtains five sets of five equations:

$$\begin{aligned}
1 &= a_{11}g_{\text{SM},1}^4 + a_{12}g_{\text{SM},1}^3g_{\text{BSM},1} + a_{13}g_{\text{SM},1}^2g_{\text{BSM},1}^2 + a_{14}g_{\text{SM},1}g_{\text{BSM},1}^3 + a_{15}g_{\text{BSM},1}^4 \\
0 &= a_{21}g_{\text{SM},1}^4 + a_{22}g_{\text{SM},1}^3g_{\text{BSM},1} + a_{23}g_{\text{SM},1}^2g_{\text{BSM},1}^2 + a_{24}g_{\text{SM},1}g_{\text{BSM},1}^3 + a_{25}g_{\text{BSM},1}^4 \\
&\dots \\
0 &= a_{11}g_{\text{SM},2}^4 + a_{12}g_{\text{SM},2}^3g_{\text{BSM},2} + a_{13}g_{\text{SM},2}^2g_{\text{BSM},2}^2 + a_{14}g_{\text{SM},2}g_{\text{BSM},2}^3 + a_{15}g_{\text{BSM},2}^4 \\
1 &= a_{21}g_{\text{SM},2}^4 + a_{22}g_{\text{SM},2}^3g_{\text{BSM},2} + a_{23}g_{\text{SM},2}^2g_{\text{BSM},2}^2 + a_{24}g_{\text{SM},2}g_{\text{BSM},2}^3 + a_{25}g_{\text{BSM},2}^4 \\
&\dots,
\end{aligned} \tag{9.16}$$

which can be written in matrix form as

$$\begin{pmatrix} a_{11} & a_{12} & a_{13} & a_{14} & a_{15} \\ a_{21} & a_{22} & a_{23} & a_{24} & a_{25} \\ a_{31} & a_{32} & a_{33} & a_{34} & a_{35} \\ a_{41} & a_{42} & a_{43} & a_{44} & a_{45} \\ a_{51} & a_{52} & a_{53} & a_{54} & a_{55} \end{pmatrix} \cdot \begin{pmatrix} g_{\text{SM},1}^4 & g_{\text{SM},2}^4 & g_{\text{SM},3}^4 & g_{\text{SM},4}^4 & g_{\text{SM},5}^4 \\ g_{\text{SM},1}^3g_{\text{BSM},1} & g_{\text{SM},2}^3g_{\text{BSM},2} & g_{\text{SM},3}^3g_{\text{BSM},3} & g_{\text{SM},4}^3g_{\text{BSM},4} & g_{\text{SM},5}^3g_{\text{BSM},5} \\ g_{\text{SM},1}^2g_{\text{BSM},1}^2 & g_{\text{SM},2}^2g_{\text{BSM},2}^2 & g_{\text{SM},3}^2g_{\text{BSM},3}^2 & g_{\text{SM},4}^2g_{\text{BSM},4}^2 & g_{\text{SM},5}^2g_{\text{BSM},5}^2 \\ g_{\text{SM},1}g_{\text{BSM},1}^3 & g_{\text{SM},2}g_{\text{BSM},2}^3 & g_{\text{SM},3}g_{\text{BSM},3}^3 & g_{\text{SM},4}g_{\text{BSM},4}^3 & g_{\text{SM},5}g_{\text{BSM},5}^3 \\ g_{\text{BSM},1}^4 & g_{\text{BSM},2}^4 & g_{\text{BSM},3}^4 & g_{\text{BSM},4}^4 & g_{\text{BSM},5}^4 \end{pmatrix} = \mathbb{1}. \tag{9.17}$$

or in short: $A \cdot G = \mathbb{1}$. The coefficients a_{ji} of the weight functions w_j are obtained by $A = G^{-1}$, if $\det G \neq 0$.

9.2.2 Validation of the $\mathcal{OO}_{1,jj}$ morphed shapes

While the morphing technique theoretically can span the full phase-space, provided its realisation conditions are fulfilled, in practice it acquires a validity range due to numerical precision of the computation and to the finite statistics of the input distributions. A few examples of morphed shapes are compared to MC-generated shapes along the $\tilde{c}_{zz}$ axis in Figure 9.9. Two morphing bases are tested: $\tilde{c}_{zz} \in \{0, \pm 1, \pm 2\}$ and $\tilde{c}_{zz} \in \{0, \pm 1.5, \pm 4\}$ (referred to as small range and extended range, respectively), and give different validity ranges. Figure 9.9a shows the distributions for the SM sample. The morphed shapes with both bases perfectly agree with the generated shape, as expected since this is a basis point of the morphing.

The shapes at $\tilde{c}_{zz} = 1$ are compared in Figure 9.9b. As this is a basis point for the small range morphing, the morphed shape with this basis perfectly agrees with the generated shape. The morphed shape with extended range also shows a good agreement with the generated shape.

In Figure 9.9c, the tested coupling is $\tilde{c}_{zz} = 3$. This value is still within the range of the extended range basis, so the extended range morphed shape has a good compatibility with the generated shape. It is however outside the range of the small range morphing, consequently the morphed shape using the small range basis shows some significant deviation from the generated shape in some bins. This indicates the edge of the validity range of the small range morphing.

Finally, Figure 9.9d shows the morphed shape for $\tilde{c}_{zz} = 5$. The coupling value is far beyond the last basis sample of the small range morphing, which shows large differences with the generated shape. In the case of the extended range morphed sample however, $\tilde{c}_{zz} = 5$ is only slightly beyond the last basis point and the morphing manages to produce a shape consistent with the MC generated shape within less than 10 %, showing here as well the limit of its validity range. Given these results, “extended range” morphing bases are preferred in the following of this Chapter.

9.3 Preliminary studies

Since no sign of new physics has yet been seen around the electroweak scale (around 200 GeV), it is expected that BSM effects would rather manifest at higher energy. Consequently, BSM searches should mainly target processes with high momentum transfer. In the Higgs sector, most the Higgs boson production happens through the ggF mode ($\sim 87\%$). Since the ggF process occurs in the s -channel, the typical momentum transfer of Higgs boson production is m_H (~ 125 GeV), therefore not in the range where new physics would expectedly shows. However, the VBF production ($\sim 7\%$) occurs in t -channel, allowing for much larger momentum transfer, up to hundred’s of GeV. Moreover, the jet kinematics produced in VBF collisions carry some information on the spin and parity of the Higgs boson. Considering these arguments, it has been decided to focus the analysis on production-side observables as discriminating variables for the Higgs boson parity. In order to get a maximal sensitivity, the $\mathcal{OO}_{1,jj}$ as defined in Section 9.1 is used.

In order to prepare a limited set of samples needing a full detector simulation, an approximate sensitivity range estimation is performed. These studies have been conducted on truth-level samples (*i.e.* without any detector simulation, only hard-scatter and parton showering simulation) that can be generated with limited computing resources. It also allows to test the analysis framework to ensure the validity of the results. After having generated MC simulations covering the possible sensitivity range over the full phase-space (more than 250 samples with 200k events each, see Section 9.3.1), rough sensitivity estimates are derived according to each parametrisation of the phase-space (Section 9.3.2). The morphing method is validated on these results, as demonstrated in Section 9.3.4. In Section 9.3.5, the results are refined to get a more realistic estimate of the expected sensitivity after inclusion of detector effects.

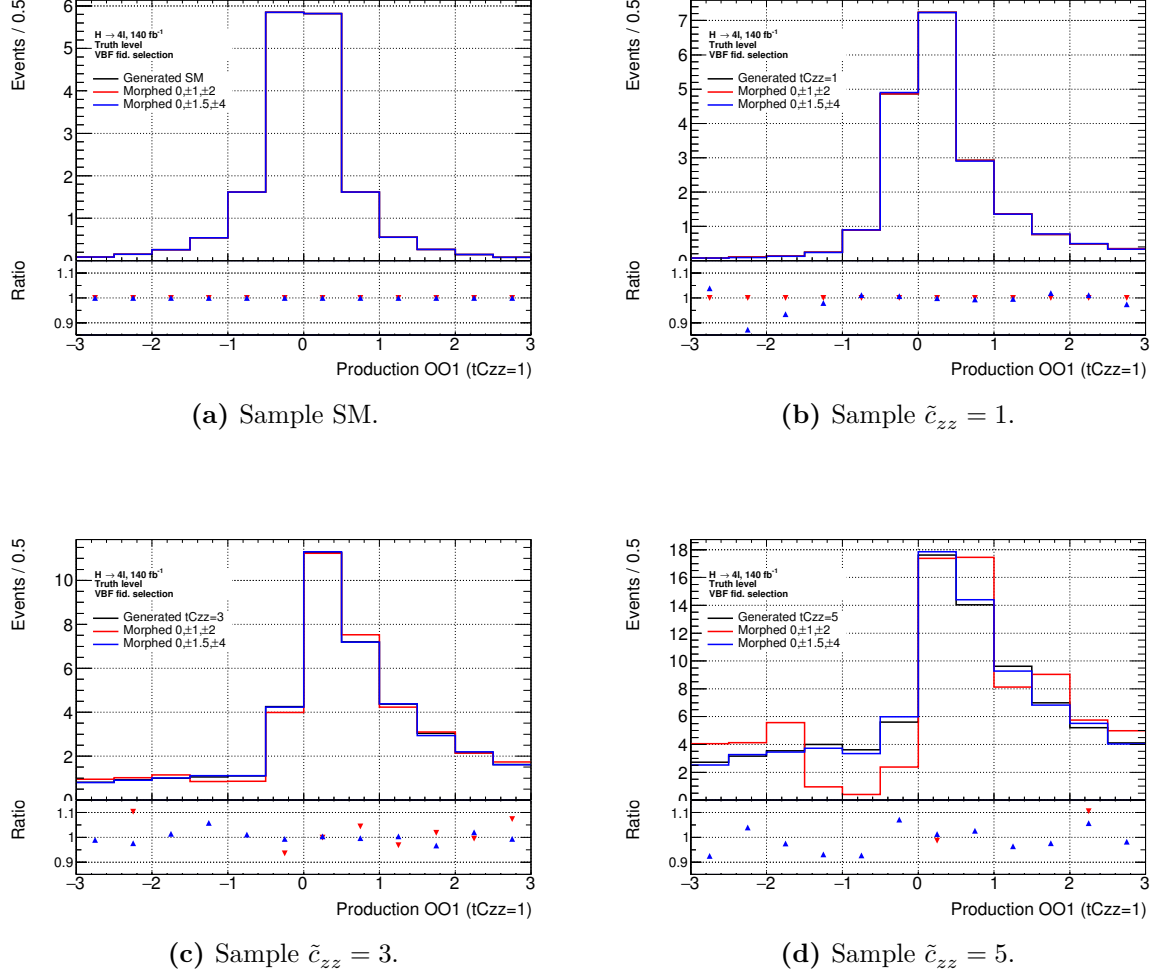


Figure 9.9 – Comparison of morphed shapes for two different morphing bases (red: $\tilde{c}_{zz} \in \{0, \pm 1, \pm 2\}$, and blue: $\tilde{c}_{zz} \in \{0, \pm 1.5, \pm 4\}$) to the MC generated shape (black) for $\mathcal{OO}_{1,jj}(\tilde{c}_{zz} = 1)$ for samples generated with various values of tCzz. (a) shows the $\mathcal{OO}_{1,jj}$ shape for the SM sample, which is a base point for both bases. It is perfectly morphed as expected. (b) shows the $\mathcal{OO}_{1,jj}$ shape for the sample $\tilde{c}_{zz} = 1$: it is a base point of the red basis but not for the blue basis. It is expectedly perfectly morphed in the red basis, and correctly morphed for the blue basis. (c) shows the $\mathcal{OO}_{1,jj}$ shape for the sample $\tilde{c}_{zz} = 3$: it is slightly beyond the last point of the red basis and some differences compared to the MC-generated shape start to be seen. The morphing using the blue basis works well. (d) shows the $\mathcal{OO}_{1,jj}$ shape for the sample $\tilde{c}_{zz} = 5$: this sample extends far beyond the last sample of the red basis, in which the morphing fails. For the blue basis however, this point is still not too far from one of the base point and the morphing still works correctly.

9.3.1 Additional BSM simulated samples

The analysis presented in this Chapter targets an interpretation in terms of limits on EFT Wilson coefficients, so in principle the results can be quoted in any EFT basis and later translated to other equivalent bases. The primary basis chosen for the interpretation is the Higgs basis, since it is more suited for Higgs boson measurements. Given the current yield in the VBF $H \rightarrow 4\ell$ analysis however, simultaneous constraints on several eigenaxes of a basis would not be adapted, and the limits on Wilson coefficients are only quoted along a single axis. Furthermore, the translation of two-dimensional constraints to other bases might not be as straightforward. Since the Warsaw basis is preferred for the combination with other non-Higgs analyses, the results should therefore also be quoted in this basis. From these arguments, the analysis presented here is performed along each eigenaxis of the Higgs and Warsaw bases ($\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$, $\tilde{c}_{\gamma\gamma}$ and $c_{H\tilde{W}}$, $c_{H\tilde{W}B}$, $c_{H\tilde{B}}$, respectively). Additionally, having measurements in both bases allows for a cross-validation of the results if needed.

To be able to compare a measured shape to an expected shape under a set of couplings, reference shapes have to be produced, hence some MC simulations are needed. To allow for detailed studies in both the Higgs and the Warsaw bases, samples were generated along each of the three CP -odd coefficients axes in both bases. Reasonable upper bounds for the Wilson coefficients are of the order 1 to 10, so the samples were generated across this range, with higher density around 0 (Standard Model). For some of the couplings (namely $\tilde{c}_{\gamma\gamma}$, $c_{H\tilde{W}B}$ and $c_{H\tilde{B}}$), this “reasonable” estimate revealed insufficient to provide an estimate, and additional points were generated extending as far as 50 to 100.

In order to generate events with modified couplings according to some EFT coefficients, EFT bases need to be implemented in the generator (in this case, the Higgs basis, the Warsaw basis, or the Higgs Characterisation). The MADGRAPH generator implements both the WB and HC so any set of Wilson coefficient from HB, WB or HC can be used to produce BSM-like samples, if necessary after a basis translation step. 200k events were produced per sample to have sufficient statistics in case of low acceptance due to strongly modified kinematics. The generation matrix elements are computed as leading order in QCD, and the samples were later showered using the standard PYTHIA 8 + EVTGEN chain.

It should be noted that no ggF or VH BSM samples were produced, so they are assumed to be SM-like for all the studies in this Section. Since modifying the HVV couplings only impacts the decay side of the ggF production mode and given that the analysis is focused on production-side observables, the impact of modified ggF final state can be neglected. For VH however, both the decay and production couplings are affected, so this argument does not hold. Yet, the VH production is suppressed compared to the VBF production and assuming SM-like VH production is reasonable for such preliminary studies. The final fully-simulated samples however include both VBF and VH BSM samples.

9.3.2 Sensitivity using VBF signal samples only

As mentioned above, the analysis primarily focuses on events produced by VBF process since it is expected to give the best sensitivity. Experimentally, this category is defined as events with at least two jets (with $p_{T,j1}, p_{T,j2} > 30$ GeV) whose invariant mass is $m_{jj} > 120$ GeV (to remove part of the VH -had events). Additionally, MVA methods (Chapter 8.6.2) can be used to improve the purity of this category.

As seen above in Section 9.1, the most sensitive variables that can be built to probe for a CP -odd contribution in the Higgs boson couplings are the so-called optimal observables. Since the second order optimal observable $\mathcal{O}_{2,jj}$ is only sensitive to the amplitude of the BSM contribution

but not its sign, only the first order optimal observable $\mathcal{OO}_{1,jj}$ is used. Moreover, these $\mathcal{OO}$ are built with a given hypothesis. Given the sensitivity range needs to be established along all six eigenaxes of the Higgs and Warsaw bases, optimal observables built along each of these six axes are tried.

In order to provide a first, back-of-the-envelope sensitivity range estimate, it is assumed that the VBF category is 100 % pure in VBF events (*i.e.* there is no backgrounds). A refinement of the estimate including backgrounds is discussed in Section 9.3.5. In Figure 9.10 are shown the sensitivity curves (likelihood scans) along each of the six eigenaxes of the Higgs and Warsaw bases, for each of the six possible $\mathcal{OO}_{1,jj}$. The numerical values are summarised in Table 9.1.

First, it is seen that the sensitivity curves reach a plateau at relatively low values of likelihood ratio for the scans along $\tilde{c}_{\gamma\gamma}$, $c_{H\tilde{W}B}$ and $c_{H\tilde{B}}$. As a consequence, it is unlikely that a result could be quoted when detector effects and background contamination are included, and these axes are not further studied. For the three other axes ($\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$, $c_{H\tilde{W}}$), similar conclusions can be drawn: firstly, the different observables $\mathcal{OO}_{1,jj}(\tilde{c}_{zz} = 1)$, $\mathcal{OO}_{1,jj}(\tilde{c}_{z\gamma} = 1)$, $\mathcal{OO}_{1,jj}(\tilde{c}_{\gamma\gamma} = 1)$ and $\mathcal{OO}_{1,jj}(c_{H\tilde{W}} = 1)$, have a similar performance, while $\mathcal{OO}_{1,jj}(c_{H\tilde{W}B} = 1)$ and $\mathcal{OO}_{1,jj}(c_{H\tilde{B}} = 1)$ perform worse. Secondly, the sensitivity does not reach a plateau in the given scanned range. An inflexion point starts to appear for $\tilde{c}_{zz}$ and $c_{H\tilde{W}}$ at the edges of this range, but the likelihood ratio values reached are high enough to expect a sufficient sensitivity, enabling a confidence interval at the 95 % CL to be quoted after backgrounds and detector effects are included.

9.3.3 Limits from cross-section results

The limits shown in Section 9.3.2 only use the information contained in the optimal observable and is therefore a purely *CP*-sensitive information. However, generic limits using the current knowledge of Higgs boson production cross-section also provides some constraints. Starting from the STXS result presented in Chapter 7.4.1, the current constraint on VBF production cross-section is $1.21^{+0.24}_{-0.22}$. In other words, the 68 % CL uncertainty on the VBF cross-section is about 25 %. As a consequence, any BSM coupling yielding a VBF cross-section 25 % higher than the Standard Model expectation is already excluded at the 68 % CL. It should be noted that this is not meant to replace this analysis but rather to give an idea of where the analysis can provide competitive results compared to other methods.

The MADGRAPH program was used to compute the VBF $H \rightarrow 4\ell$ cross-section at leading order for various $\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$, $\tilde{c}_{\gamma\gamma}$, $c_{H\tilde{W}}$, $c_{H\tilde{W}B}$ and $c_{H\tilde{B}}$ input values. The VBF $H \rightarrow 4\ell$ cross-section times branching ratio normalised by the Standard Model expectation (all computed at leading order) as a function of the coupling value is shown in Figure 9.11. As the couplings are found in both the production and the decay of the Higgs boson when considering the VBF $H \rightarrow 4\ell$ process, the cross-section times branching is expected to vary as the fourth power of the coupling. Moreover, since the cross-section of *CP*-even and *CP*-odd production do not interfere, the cross-section times branching is expected to be the same (within computing precision) for a given coupling value and its opposite. The variation is thus expected to be symmetrical, which is verified in the plots. The cross-section times branching ratio normalised by the Standard Model expected value can hence be fitted the symmetric quartic function $f(x) = 1 + a \cdot x^2 + b \cdot x^4$. The result of these fits are also shown on the plots.

The coupling value yielding a cross-section times branching ratio 25 % higher than the leading order Standard Model expectation are summarised in Table 9.1 for each axis. If the expected 68 % CL expected sensitivity is much stronger with the cross-section measurement than with the optimal observable method, further analysis of the sensitivity along this axis is of least importance as the constraint will not be improved using the optimal observable. Comparing the constraints on *CP*-odd coefficients from cross-section and from the optimal observable sensitivity, the sensitivity

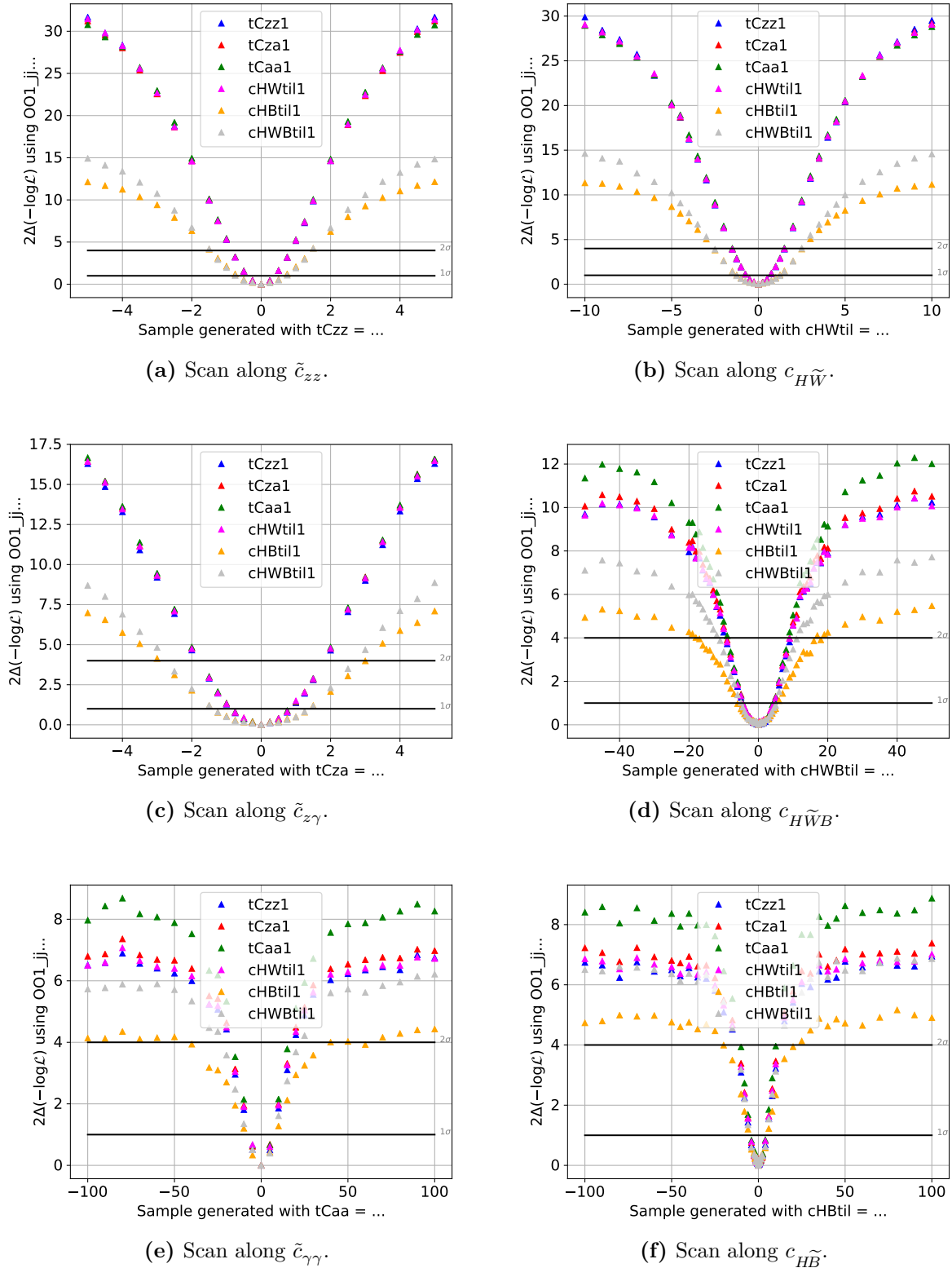


Figure 9.10 — Log-likelihood ratio of (left) Higgs basis CP -odd Wilson coefficients ($\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$, $\tilde{c}_{\gamma\gamma}$) and on (right) Warsaw basis CP -odd Wilson coefficients ($c_{H\tilde{W}}$, $c_{H\tilde{W}B}$, $c_{H\tilde{B}}$), using $\mathcal{O}_{1,jj}$ built with various hypotheses (one along each axis). The scans are performed on truth-level VBF samples, assuming an integrated luminosity of 140 fb^{-1} and no background contamination. The reference shape (“asimov dataset”) is assumed to be the Standard Model.

to $\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$ and $c_{H\tilde{W}}$ is better or of the same order of magnitude using the optimal observable and the cross-section result while the sensitivity to $\tilde{c}_{\gamma\gamma}$, $c_{H\tilde{W}B}$ and $c_{H\tilde{B}}$ is roughly ten-fold better from cross-section constraints than from optimal observable method. Consequently, the conclusion drawn in Section 9.3.2 is confirmed and the next steps will mainly focus on the former three parameters as they give best chances of improvement (especially $\tilde{c}_{zz}$).

Table 9.1 – Comparison of limits on the CP -odd Wilson coefficients obtained using truth-level pure VBF samples and from cross-section measurements (coupling values yielding a $\text{VBF } H \rightarrow 4\ell$ cross-section times branching ratio 25 % higher than the Standard Model prediction, see Chapter 7.4.1).

Wilson coefficient	Higgs basis			Warsaw basis		
	$\tilde{c}_{zz}$	$\tilde{c}_{z\gamma}$	$\tilde{c}_{\gamma\gamma}$	$c_{H\tilde{W}}$	$c_{H\tilde{W}B}$	$c_{H\tilde{B}}$
68 % CL limit with $\mathcal{O}\mathcal{O}_{1,jj}$ (VBF only)	± 0.38	± 0.85	± 6.2	± 0.70	± 5.2	± 5.4
68 % CL limit from cross-section	± 1.24	± 0.39	± 0.94	± 0.73	± 0.73	± 0.40

9.3.4 Validation of the morphing technique

In Section 9.2.2, the compatibility between MC-generated and morphed shapes for $\mathcal{O}\mathcal{O}_{1,jj}$ has been checked. In this Section, the validity of the scan results obtained using morphed shapes is assessed. Figure 9.12 shows the likelihood scan using the MC generates shapes (points) or using morphed shapes (lines). As the morphing result can depend on the set of basis samples chosen, the scan was also performed using two different morphing bases. The three trials show a good agreement within the validity range of the morphing: for scanning points beyond the last morphing point, the morphed shape is not compatible with the MC-generated shape and the results are not reliable.

9.3.5 Sensitivity including $ZZ^{(*)}$ and other Higgs boson production modes

The results in Section 9.3.2 and Section 9.3.3 demonstrated that this analysis would not be sensitive to $\tilde{c}_{\gamma\gamma}$ in the Higgs basis, nor $c_{H\tilde{W}B}$ and $c_{H\tilde{B}}$ in the Warsaw basis. Consequently, the following studies focus on $\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$, and $c_{H\tilde{W}}$ only. Also since the morphing is validated (Section 9.3.4), further results all use this method as interpolation between points.

The study presented in Section 9.3.2 however assumes that the analysis category is 100 % pure in VBF signal. However, from Figure 8.16 it is expected that in the data, a significant fraction of this category is polluted by ggF and $qq \rightarrow ZZ^{(*)}$ processes (referred to as backgrounds in the following). A further necessary refinement of the sensitivity estimates is therefore needed to assess the amount of sensitivity lost due to contamination of these backgrounds in the VBF analysis category. Moreover, while the results of Section 9.3.2 already included detector acceptance effects, it did not include the event reconstruction efficiency, which is around 45 %.

The impact of including the main background processes as well as the reconstruction efficiency effects in the sensitivity estimate is seen by comparing the black and red curves in Figure 9.13. After inclusion of the backgrounds, the precision of the measurement decreases by a factor of ~ 3 for the 68 % CL interval, and by a factor of 5 to 6 for the 95 % CL interval. In order to recover some of the lost sensitivity, MVA methods can however be applied to separate VBF events from ggF and $qq \rightarrow ZZ^{(*)}$. At the time these checks were performed, the neural network (NN) discriminants presented in Chapter 8.6.2 were being developed and the improved classification relied on Boosted Decision Trees (BDT) instead. These two methods however rely on the same assumptions and provide similar results.

The method using MVA consists in splitting the VBF category into sub-bins corresponding

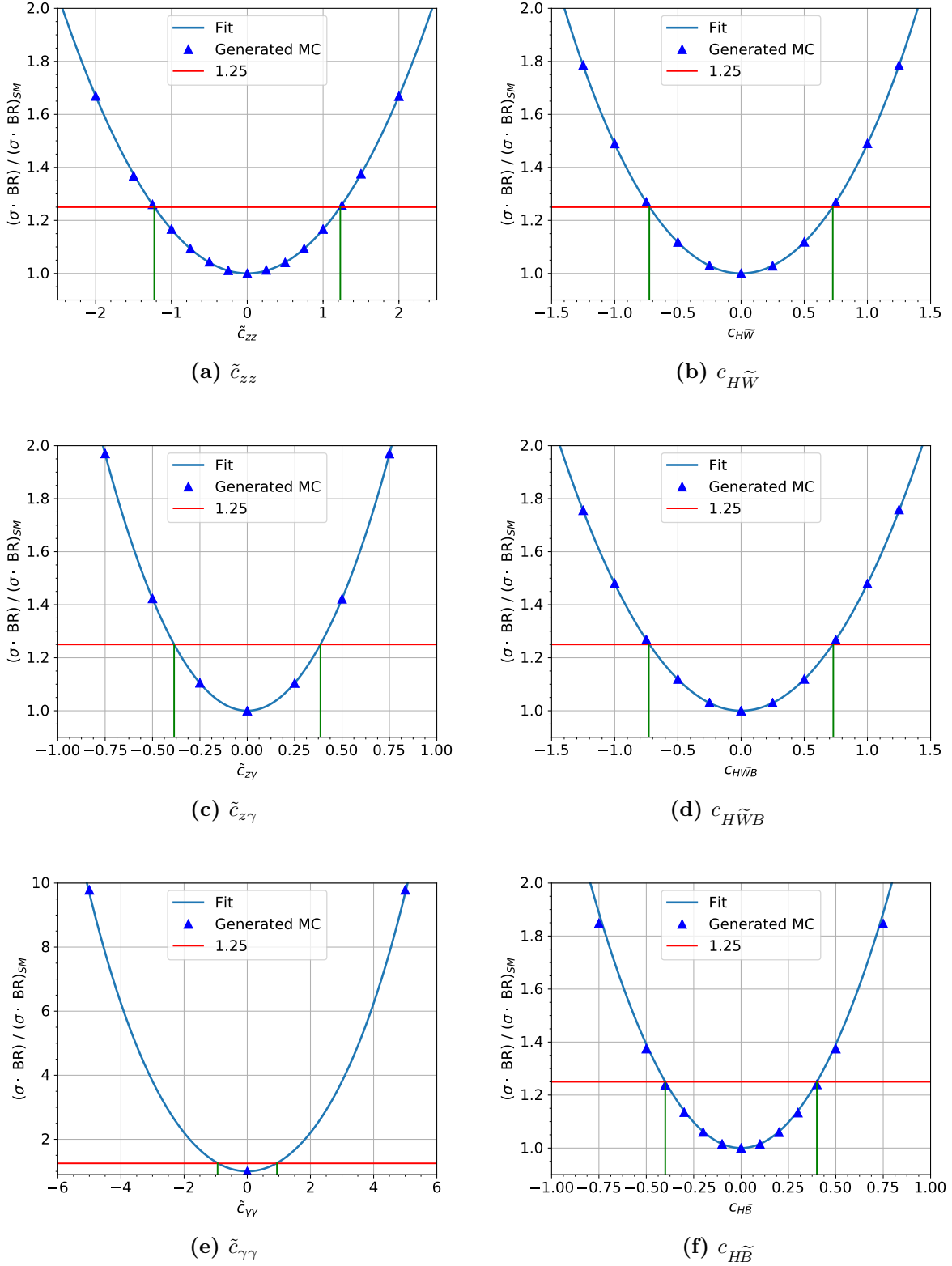


Figure 9.11 – $\text{VBF}H \rightarrow 4\ell$ cross-section times branching ratio normalised to the Standard Model value (computed at leading order) for the (left) Higgs basis CP -odd Wilson coefficients ($\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$, $\tilde{c}_{\gamma\gamma}$) and the (right) Warsaw basis CP -odd Wilson coefficients ($c_{H\tilde{W}}$, $c_{H\tilde{W}B}$, $c_{H\tilde{B}}$). The fit is performed with a quartic function $x \rightarrow ax^4 + bx^2 + 1$ over all the generated points while the plots only show points within a few units of the SM cross-section times branching ratio to improve readability. The values yielding a cross-section times branching ratio 25 % higher than the Standard Model prediction are summarised in Table 9.1.

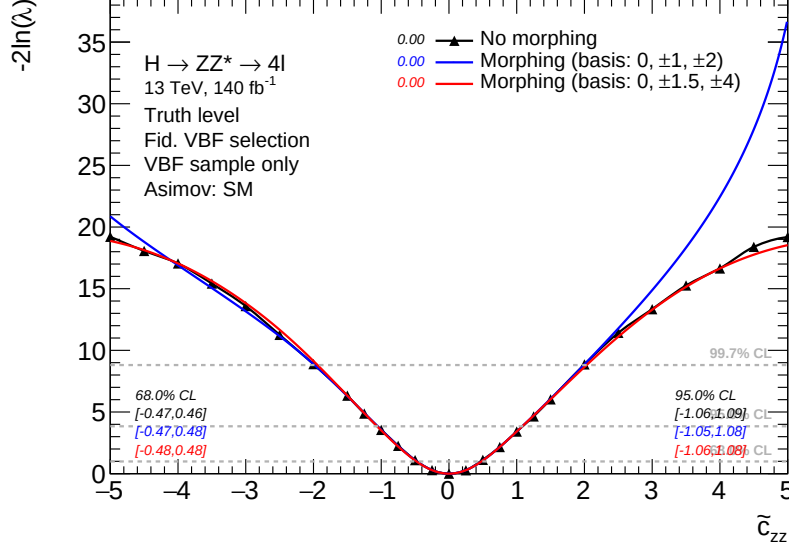


Figure 9.12 – Comparison of log-likelihood ratios to the Standard Model scanning along $\tilde{c}_{zz}$ with the $\mathcal{OO}_{1,jj}(\tilde{c}_{zz} = 1)$ observable, and using generated samples (black points) or morphed samples (blue and red curves). An integrated luminosity of 140 fb^{-1} is assumed. Only the truth-level generated VBF sample is used (no background included). Two morphing bases are tested: $\tilde{c}_{zz} \in \{-2, -1, 0, 1, 2\}$ (blue) and $\tilde{c}_{zz} \in \{-4, -1.5, 0, 1.5, 4\}$ (red). The agreement between the individually generated samples and the morphing with “short range” basis (blue) is good as long as the tested sample is not too far from the last morphing basis points (*i.e.* until $\tilde{c}_{zz} \approx 2.5$). Since the morphing with “large range” basis (red) validity range extends further (until $\tilde{c}_{zz} \approx 5$), the agreement between the generated points and this morphing basis is good across all the tested range.

to bins of BDT score (similarly to Figure 8.18a). This method acts like a weight which is applied to events, with a larger weight for events in BDT bins with higher VBF purity, and a smaller weights for events with a BDT score corresponding to background-like bins. The result of using this method is shown with the blue curves in Figure 9.13, that should be compared to the red ones. The method manages to improve the precision of the result by a factor 1.5 for the 68 % CL interval and around 2 for the 95 % CL interval. The results are summarised in Table 9.2.

9.4 Preliminary results

As the best sensitivity is expected to come from the $\tilde{c}_{zz}$ coefficient, this Section focuses on this coefficient only. A similar analysis can nevertheless be performed on the other axes of interest ($\tilde{c}_{z\gamma}$, $c_{H\tilde{W}}$ as well as $\tilde{c}_{\gamma\gamma}$, $c_{H\tilde{W}B}$ and $c_{H\tilde{B}}$). In Section 9.4.1 the results at truth level are compared to the one obtained using fully-simulated samples. In Section 9.4.2 a preliminary evaluation of systematic uncertainties is performed. Finally, the possibility of a bias in the results due to the use of neural networks in the selection is discussed in Section 9.4.3.

9.4.1 Expected sensitivity with fully-simulated samples

Other studies in the group have showed that only about 45 % of events seen at truth level are selected after a full reconstruction in the detector. Consequently, the results presented in Section 9.3.5 with truth-level samples include a 45 % scaling in order to get closer to the expected results from a fully-simulated sample. This factor approximately accounts for the loss of statistics from detector effects, but does not provide an estimate of sensitivity loss due to, for example,

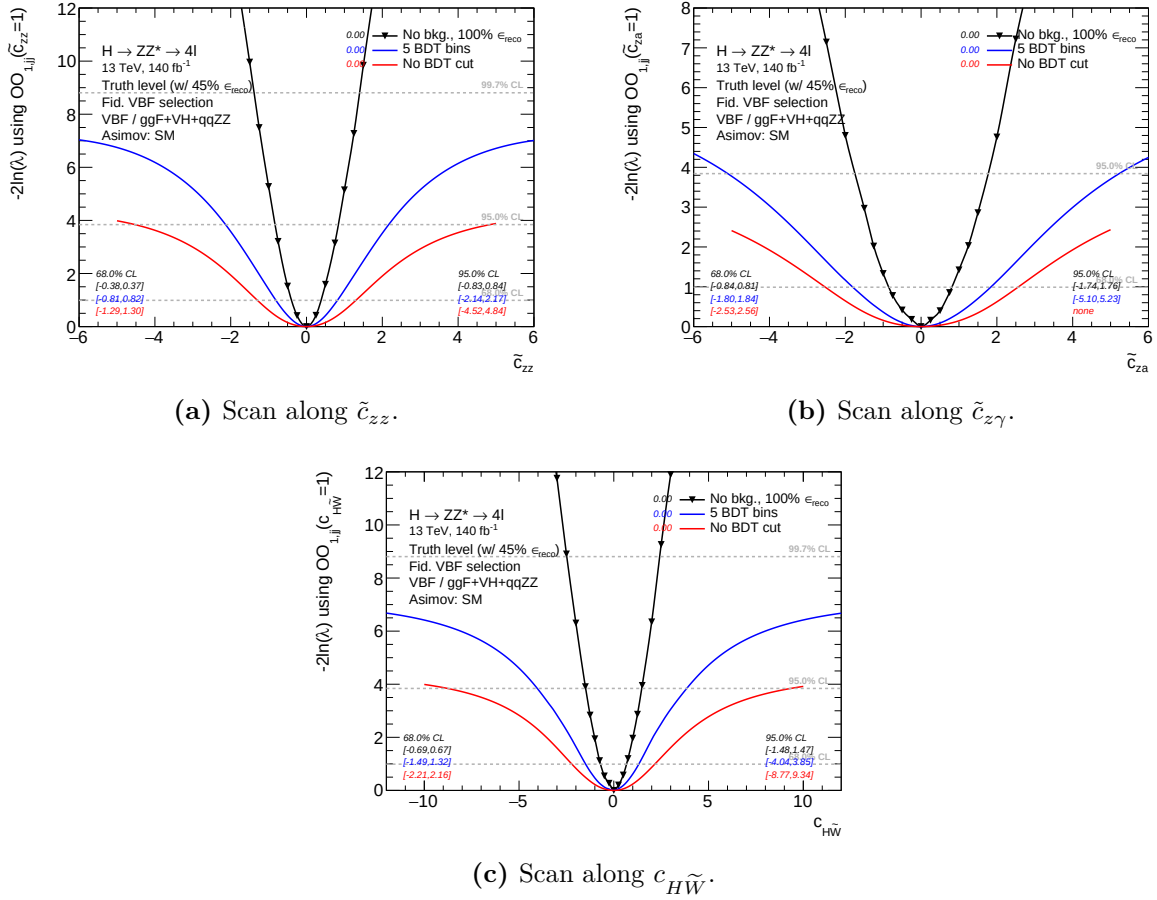


Figure 9.13 – Log-likelihood ratio to a SM asimov scanning along the (a) $\tilde{c}_{zz}$, (b) $\tilde{c}_{z\gamma}$ and (c) $c_{H\tilde{W}}$ axes (using $\mathcal{OO}_{1,jj}(\tilde{c}_{zz} = 1)$, $\mathcal{OO}_{1,jj}(\tilde{c}_{z\gamma} = 1)$ and $\mathcal{OO}_{1,jj}(c_{H\tilde{W}} = 1)$ respectively). The scan without background and assuming a 100 % reconstruction efficiency (black points, corresponding to values from Figure 9.10) are recalled and compared to the results including the ggF, VH and $qq \rightarrow ZZ^{(*)}$ backgrounds as well as a 45 % reconstruction efficiency (red curves). The blue curves show the improvement over the red curve by combining the analysis over 5 VBF sub-categories defined by 5 BDT bins (following a similar pattern as the NN in Figure 8.18a). An integrated luminosity of 140 fb^{-1} is assumed in each case.

Table 9.2 – Truth-level sensitivity scans along the $\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$ and $c_{H\tilde{W}}$ axes without background and assuming 100 % reconstruction efficiency, and with backgrounds (ggF and VH Higgs boson production, and non resonant $qq \rightarrow ZZ^{(*)}$ background) further assuming 45 % reconstruction efficiency. Two variations are given for this last category: with and without use of BDT to separate the VBF signal from the other processes. When used, the analysis is split in 5 bins of BDT score. The values correspond to the plots of Figure 9.13. An integrated luminosity of 140 fb^{-1} is assumed in each case. “–” indicates that no sensitivity at this confidence level was achieved in the range of the scan.

Wilson coefficient	Higgs basis				Warsaw basis	
	$\tilde{c}_{zz}$		$\tilde{c}_{z\gamma}$		$c_{H\tilde{W}}$	
Confidence level	68 % CL	95 % CL	68 % CL	95 % CL	68 % CL	95 % CL
No background, $\epsilon_{\text{Reco}} = 100 \%$	± 0.4	± 0.8	± 0.8	± 1.8	± 0.7	± 1.5
With bkg., $\epsilon_{\text{Reco}} = 45 \%$, 5 BDT bins	± 0.8	± 2.2	± 1.8	± 5.2	± 1.4	± 4.0
With bkg., $\epsilon_{\text{Reco}} = 45 \%$	± 1.3	± 4.7	± 2.6	–	± 2.2	± 9.0

resolution effects on lepton and jet energy and position. To give a reliable estimate of such effects, fully-simulated samples have been generated and the results compared to the truth-level sensitivity estimate. Given the much larger computing resources needed to generate such fully-simulated samples, only six samples with 200k events along each of the $\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$ and $c_{H\tilde{W}}$ axes have been generated in addition to the Standard Model sample. The previously mentioned morphing technique is used to interpolate between these samples.

Figure 9.14 shows a comparison of the same likelihood scan performed with truth-level samples (green), including the 45 % reduction of statistics, or with fully-simulated samples (blue). The two scans display a close behaviour and a good agreement between the expected 68 % CL intervals ($[-0.81, +0.82]$ for truth level, $[-0.82, +0.79]$ for fully-simulated analysis). It is concluded that the result does not significantly change using reconstructed jets instead of truth-level jets, despite the degradation of the jet position measurement and energy resolution due to the detector.

The 95 % CL interval differ slightly more between the truth-level ($[-2.1, +2.2]$) and the fully-simulated analysis ($[-2.3, +2.4]$) due to the likelihood scans slightly diverging. This can be due to small changes in the analysis setting, such as the choice of morphing basis ($\tilde{c}_{zz} \in \{0, \pm 1.5, \pm 4\}$ for truth samples, and $\tilde{c}_{zz} \in \{0, \pm 1, \pm 2\}$ for the fully-simulated samples), but is unlikely to be linked to reconstruction effects.

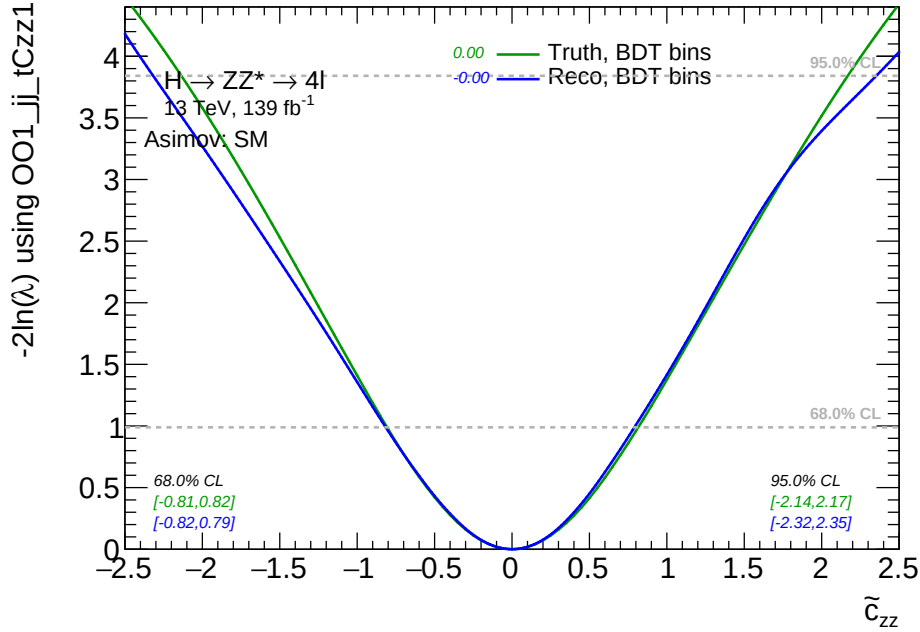


Figure 9.14 – Comparison of the log-likelihood scans along the $\tilde{c}_{zz}$ axis for a SM asimov using the $\mathcal{OO}_{1,jj}$ observable for truth (green) and fully reconstructed samples (blue) including both signal and backgrounds. The truth samples include a 45 % reconstruction efficiency. The green curve corresponds to the blue curve of Figure 9.13a (plotted over a shorter range). An integrated luminosity of 139 fb^{-1} is assumed.

9.4.2 Systematic uncertainties assessment

After assessing that reconstruction effects do not have an unexpected impact on the sensitivity, systematics uncertainties affecting the analysis can be studied. The results given in this Section are preliminary and not all sources of systematic uncertainties have been included yet. The systematics common to all $H \rightarrow 4l$ analyses could be evaluated, but more specific sources such as the variation of the fraction of ggF events selected in the VBF category need more careful studies before their impact is reliably estimated.

Following the classification of the various systematic uncertainty types described in Chapter 8.5, the following sources of weight-changing systematics have been taken into account (the number of nuisance parameters relating to each source is indicated in parenthesis):

- luminosity (1);
- electrons: identification (34), isolation (1) and reconstruction (1) efficiencies;
- muons: isolation (2), reconstruction (4) and track-to-vertex association (2) efficiencies;
- jets: flavour tagging efficiencies (12), pileup rejection efficiency (2);
- pileup reweighting (1);
- theory: parton shower (26), PDFs (150), scales (28), other (33);

amounting to a total 297 independent nuisance parameters. Their combined contribution is computed as the quadratic sum of each nuisance parameter contribution, and their impact on the expected 68 % CL sensitivity to $\tilde{c}_{zz}$ is found to be negligible.

The nuisance parameters changing the shape of distributions include 3 variations for electrons, 5 variations for muons and 31 variations for jets. As the fit framework is not yet ready to include shape systematics in the morphing process, the systematic uncertainties presented here are estimated by replacing the nominal shape by its varied shape for each morphing base sample and recomputing the upper and lower bounds of the 68 % CL interval using this configuration. The systematic on the upper (resp. lower) bound is defined as the difference of the upper (resp. lower) bound computed with the nominal shape and the varied shape in the morphing. The individual impact of the 39 shape variations are quadratically summed to give their combined impact on the 68 % CL interval. To be conservative, the largest absolute effect seen among the 1σ and -1σ was used. The impact is found to be less than 1 %.

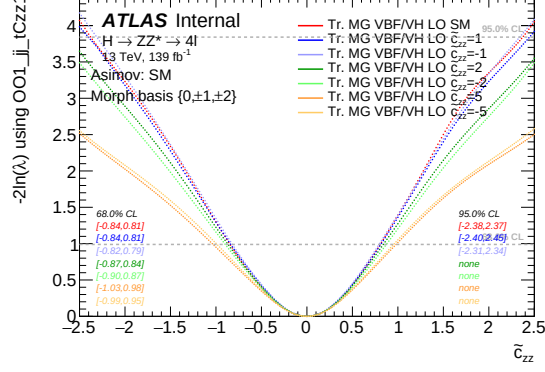
9.4.3 Study of a potential bias from usage of neural networks

In the previous Sections, more advanced event classification separating VBF from ggF, VH and $qq \rightarrow ZZ^{(*)}$ events used BDTs. Following the development of machine learning techniques, event classification is increasingly making use of neural networks, which are being implemented in the $H \rightarrow 4\ell$ analyses. This CP -odd couplings analysis is therefore also moving to using the commonly developed neural networks to produce analysis categories with a higher VBF purity.

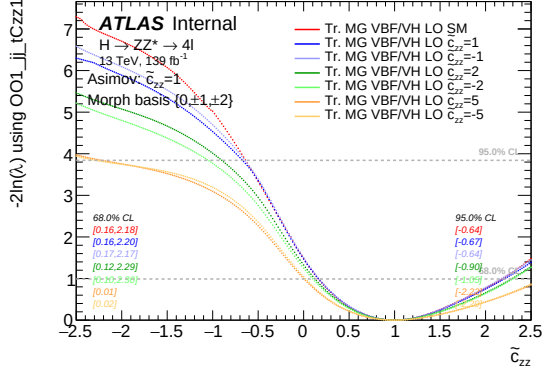
As any machine learning technique, the neural networks must first be trained on a known dataset (so-called *training sample*) to be able to classify unknown data (so-called *test sample*). To avoid biasing the classification algorithm, the training dataset must be statistically representative of the test sample. In the currently developed analysis, the VBF classification is performed with a neural network trained on Standard Model VBF, VH and ggF simulated events (see Chapter 8.6.2). However this analysis targets a measurement of BSM CP -odd couplings. Therefore, if a deviation from the Standard Model is found in the data, the test sample (*i.e.* the real data) might not be correctly represented by the training dataset, possibly inducing a bias in the CP -odd parameter measurement.

Such effect is checked by comparing the likelihood scans using classifiers trained on Standard Model and BSM samples. It should be noted that the ggF sample used for the training dataset is always the Standard Model one: since the VBF/ggF discrimination mainly relies on the jet kinematics and since those are modified by CP -odd BSM couplings in a lesser extent in the ggF than in the VBF case, such approximation is acceptable for the current sensitivity. The sensitivity curves using a classifier trained on VBF samples with $\tilde{c}_{zz} = 0$ (Standard Model, red), $\tilde{c}_{zz} = \pm 1$ (blue), $\tilde{c}_{zz} = \pm 2$ (green) or $\tilde{c}_{zz} = \pm 5$ (orange) are displayed in each plot of Figure 9.15.

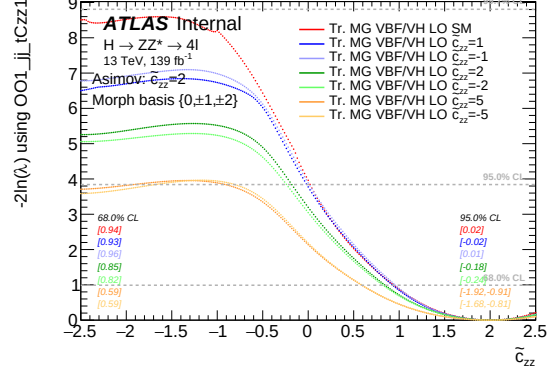
Figure 9.15a shows the likelihood scan for a Standard Model asimov and illustrates that whatever training is used, the best fitted value is the Standard Model. A similar conclusion is reached from Figure 9.15b ($\tilde{c}_{zz} = 1$ asimov), and Figure 9.15c ($\tilde{c}_{zz} = 2$ asimov): the best fitted value matches the value injected in the asimov dataset, discarding a bias from the neural network training on the central value.



(a) Scan for asimov $\tilde{c}_{zz} = 0$.



(b) Scan for asimov $\tilde{c}_{zz} = 1$.



(c) Scan for asimov $\tilde{c}_{zz} = 2$.

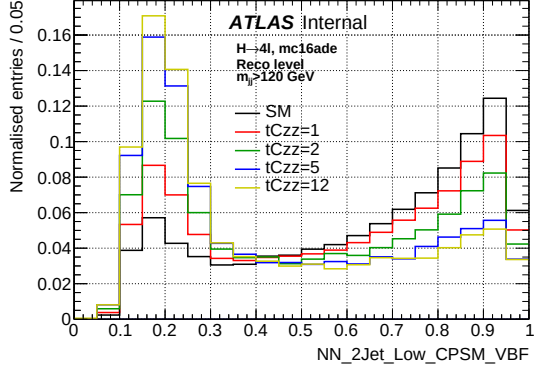
Figure 9.15 – Likelihood scans along the $\tilde{c}_{zz}$ axis using the $\mathcal{O}_{1,jj}$ observable for an asimov dataset with (a) a Standard Model hypothesis, (b) a $\tilde{c}_{zz} = 1$ hypothesis and (c) a $\tilde{c}_{zz} = 2$ hypothesis. In each case the scan is showed for an analysis using a neural network classifier trained on Standard Model VBF events (red), $\tilde{c}_{zz} = \pm 1$ VBF events (blue), $\tilde{c}_{zz} = \pm 2$ VBF events (green) and $\tilde{c}_{zz} = \pm 5$ VBF events (orange). An integrated luminosity of 139 fb^{-1} is assumed and backgrounds samples are included in the scans.

While the best fit value is insensitive to the training, the 68 % CL and 95 % CL interval are less constrained as the sample used in the training has an increasing BSM coupling. This is observed in Figure 9.15 where the curves keep the same shape but become wider as the coupling increases. Such behaviour is expected from comparing the response of various samples to classifiers trained with different VBF samples, as seen in Figure 9.16. For example, Figure 9.16a shows the response of various VBF samples to a classifier trained with a Standard Model VBF: while the response of the Standard Model VBF sample expectedly peaks toward higher values of the classifier, the response of the $\tilde{c}_{zz} = 2$ VBF sample rather peaks toward lower values. This means that $\tilde{c}_{zz} = 2$ VBF events are more likely to be classified as background-like by the Standard Model-trained classifier, leading to less pure VBF categories and decreasing the power of the analysis.

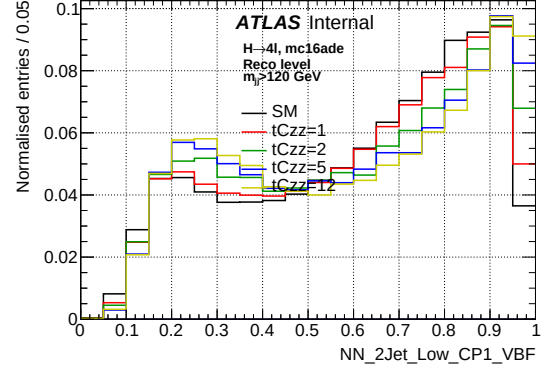
One nevertheless notices that the likelihood scan produced using the classifier trained on a $\tilde{c}_{zz} = 1$ VBF sample has no significant deviation from the scan produced using the classifier trained on a Standard Model VBF sample, and that the scan with the classifier trained on a $\tilde{c}_{zz} = 2$ VBF sample shows a less than 10 % variation. These observations hold whatever the

value injected in the asimov dataset. This indicates that for reasonable values of BSM parameters (e.g. $\tilde{c}_{zz}$), the 68 % CL and 95 % CL interval on CP -odd couplings remain reliable.

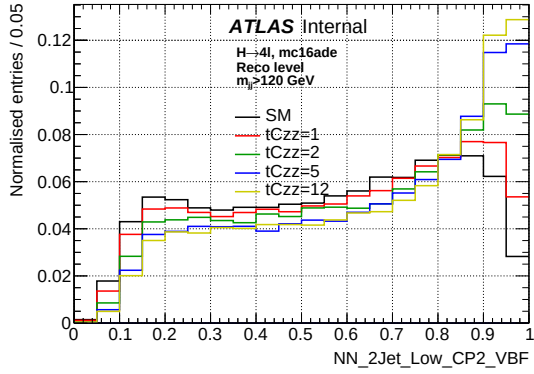
The same reasoning applies to BDTs and similar conclusions could be drawn concerning the former use of BDTs instead of the newer neural networks.



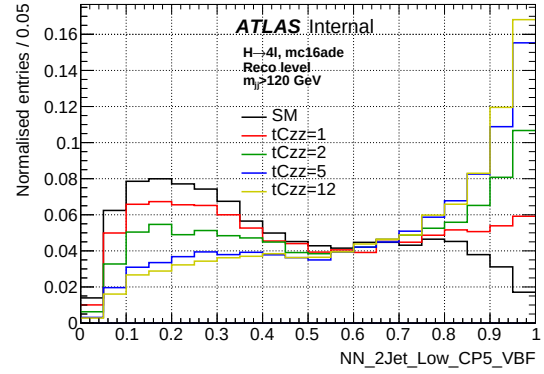
(a) Trained using VBF + VH with $\tilde{c}_{zz} = 0$.



(b) Trained using VBF + VH with $\tilde{c}_{zz} = 1$.



(c) Trained using VBF + VH with $\tilde{c}_{zz} = 2$.



(d) Trained using VBF + VH with $\tilde{c}_{zz} = 5$.

Figure 9.16 – Response of a VBF sample generated with $\tilde{c}_{zz} = 0$ (black), $\tilde{c}_{zz} = 1$ (red), $\tilde{c}_{zz} = 2$ (green), $\tilde{c}_{zz} = 5$ (blue) to a neural network trained on a VBF + VH sample generated with (a) $\tilde{c}_{zz} = 0$, (b) $\tilde{c}_{zz} = 1$, (c) $\tilde{c}_{zz} = 2$ and (d) $\tilde{c}_{zz} = 5$.

Summary and conclusion of the CP -odd couplings analysis

Since no large BSM effects were yet discovered, BSM physics is more likely to be seen at high energy. As a consequence, high momentum transfer processes such as the Higgs boson production via vector-boson fusion are of prime interest. CP -odd Higgs boson couplings in the bosonic sector are indirectly constrained from precise cross-section measurements. The current precision on the VBF cross-section is about 25 %, yielding the constraint $|\tilde{c}_{zz}| < 1.25$ at a 68 % CL.

However, production cross-section are modified by both CP -even and CP -odd BSM operators, so cross-section do not provide a pure CP -sensitive observable. As CP -odd operator modify the angular distribution of final state particles, angular variables provide unambiguous probes of CP effects. In the VBF process, such angular observable is built from the azimuthal angle between the two tagging jets. Yet, this variable only partly use the kinematic information available in the collision. Consequently observables using the full kinematic information should provide additional sensitivity to CP effects. Such observables are called optimal observables and are based on matrix element computations.

The matrix elements are defined at leading order using MADGRAPH, but are computationally expensive. The optimal observable definition and computation have been thoroughly checked, and an interface code decreasing the computing time by a factor 7–8 was developed. The sensitivity to EFT Higgs to weak-bosons CP -odd parameters ($\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$, $\tilde{c}_{\gamma\gamma}$ in the Higgs basis and $c_{H\tilde{W}}$, $c_{H\tilde{W}B}$, $c_{H\tilde{B}}$ in the Warsaw basis) using these optimal observables was assessed, and expected limits were compared to cross-section limits. It was found that the optimal observable provides the best sensitivity to the $\tilde{c}_{zz}$ Wilson coefficient while not improving the limits on $\tilde{c}_{\gamma\gamma}$, $c_{H\tilde{W}B}$ and $c_{H\tilde{B}}$ compared to cross-section measurements. The rest of the analysis therefore focused on $\tilde{c}_{zz}$, and marginally on $\tilde{c}_{z\gamma}$ and $c_{H\tilde{W}}$.

Many BSM simulated samples were produced to cover the full phase space of interest. However, such extensive event generation is only possible at truth level: samples with a full detector simulation require more computing resources and only a few samples can be generated. A means of interpolating between samples must therefore be used. This analysis uses the morphing method, which has been validated for this analysis.

As the analysis relies on VBF events, the sensitivity is spoiled due to background processes such as ggF (with two additional jets), hadronically decaying VH , and $qq \rightarrow ZZ^{(*)}$ with two additional jets. In order to purify the selection in VBF events, the analysis makes use of machine learning techniques such as boosted decision trees or neural networks. The addition of such multi-variate analysis tool allows to recover most of the sensitivity that would be achieved with a 100% pure analysis category. Since the classification algorithms are trained using Standard Model events, a concern of such classification introducing a bias in the results has been checked and no significant bias was found.

Finally, systematic uncertainties impacting the measurement of leptons and jets were evaluated and found to have a negligible impact on the sensitivity. The final expected result on the $\tilde{c}_{zz}$ coefficient is:

$$\tilde{c}_{zz} = 0.0 \pm 0.8 \text{ (stat)} \pm 0.01 \text{ (syst)}, \quad (9.18)$$

showing that the measurement is overwhelmingly statistically dominated.

Additional systematic effects must still be carefully assessed and the analysis must be completed with expected results on the other axes, in both EFT bases. Another completely model-independent measurement could be constructed by measuring the raw asymmetry in the optimal observable or jet angular distributions.

Conclusion

The search for physics beyond the Standard Model after the Higgs boson discovery focused on precision measurements. The ATLAS and CMS experiments at the LHC therefore scrutinise the Higgs boson properties to detect potential hints of new physics.

Precision physics analysis need precise measurements of physics objects (electrons, photons, muons, *etc.*) present in the final state of heavy-state disintegrations. The ATLAS electromagnetic calorimeter is central in the measurement of electron and photon energy, and needs to be carefully calibrated following several steps. One of these steps is the layer intercalibration, which was updated using the data collected during 2015 and 2016. The higher level of pileup showed to have a significant impact on the result. These effects were mitigated by developing an extrapolation method, and by changing the background modelling to using templates. The results for the 2015+2016 dataset are derived and the systematics are re-evaluated, yielding a precision comparable to the Run 1 uncertainty. The results are in agreement with the results derived during Run 1.

The inclusion of the 2017 and 2018 dataset, with even higher pileup levels, promises a great challenge for the layer intercalibration. However, the data recorded during low pileup runs could provide additional handle for the extrapolation. The possibility of only performing the intercalibration with this low-pileup dataset, and corresponding simulation, can be investigated. This would remove most of the pileup effects.

With the successful Run 2 of the LHC, ten times more Higgs data is available compared to Run 1. This allows for more delicate, specific, and precise analyses, shrinking the room for possible deviations from the Standard Model as more data are added. The uncertainties in the $H \rightarrow 4\ell$ analyses are largely dominated by the statistical component. Consequently, analyses performed in this channel benefit from such larger datasets.

One yet loosely constrained area is the CP violation in the Higgs sector, which, if found, would be a direct sign of new physics. This Thesis focused on the measurements of the Higgs boson couplings to the weak bosons in the $H \rightarrow 4\ell$ decay channel, and looked for a possible CP -odd contribution in the couplings tensor structure. As the VBF production mode is predicted to give the highest sensitivity, thanks to a high momentum transfer, the analysis focuses on VBF-like events. The reconstruction of this production modes is however polluted by the ggF production with additional jets, and machine learning techniques are employed to improve the separation. This improves the purity of the analysis category and consequently the expected sensitivity at the 68 % CL by a factor of about 1.5.

Previous limits from cross-section measurements are available, but such measurements cannot distinguish between CP -even and CP -odd contributions. In case of CP -odd effects, cross-sections are also blind to the sign of the effect. However, angular variable, such as the angle between the two associated jets in the transverse ϕ -plane, have a different behaviour under the presence of

CP -even or CP -odd effects. In order to improve the sensitivity, all the kinematic and angular information available can be used. An observable relying on matrix elements computations is therefore built. This *optimal observable*, dedicated to the measurement of CP -odd contributions, has been studied and validated for this analysis. The shape of this observable is symmetric under a CP -even contribution but gets asymmetrically distorted if CP -odd effects are present. This allows to unambiguously distinguish between beyond the Standard Model CP -even and CP -odd effects, and provides sensitivity to the sign of the contribution.

The results are interpreted in terms of effective field theory, putting limits on Wilson coefficients in the Higgs or Warsaw bases. The work performed in this Thesis has shown that the best sensitivity is achieved on the $\tilde{c}_{zz}$ coefficient from the Higgs basis, and that marginal limits could be placed on the $\tilde{c}_{z\gamma}$ coefficient (also in the Higgs basis) as well as on the $c_{H\tilde{W}}$ coefficient (in the Warsaw basis). The analysis proved to have no relevant sensitivity to other Higgs–weak boson CP -odd operators ($\tilde{c}_{\gamma\gamma}$ in the Higgs basis and $c_{H\tilde{W}B}$ and $c_{H\tilde{B}}$ in the Warsaw basis) using only the production-side information in the VBF $H \rightarrow 4\ell$ channel.

The expected 68 % CL sensitivity on the $\tilde{c}_{zz}$ coefficient is

$$\tilde{c}_{zz} = 0.0 \pm 0.8 \text{ (stat)} \pm 0.01 \text{ (syst)}.$$

This result should be compared to the cross-section only estimate, giving an expected limit of roughly $-1.25 < \tilde{c}_{zz} < 1.25$ assuming a 25 % precision on the VBF production cross-section.

Multiple ideas of improvement of this analysis can be discussed. First, the observable used in this analysis only uses the kinematic information of the associated jets and of the reconstructed Higgs boson (*i.e.* only the production-side kinematics). We have shown that the decay side kinematics carry significant information on the coefficients that could not be efficiently probed with a production-side observable ($\tilde{c}_{\gamma\gamma}$, $c_{H\tilde{W}B}$, $c_{H\tilde{B}}$). Building a matrix element observable using information from both the production and the decay (leptons) kinematics is therefore foreseen to enhance the sensitivity. Second, despite needing additional assumptions, the combination of information from the cross-sections and from the optimal observable shape would certainly boost the sensitivity. Third, only one dimensional limits were derived during this analysis, assuming other contributions are vanishing. A more robust result can be achieved by simultaneously constraining the three independent CP -odd coefficients ($\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$, $\tilde{c}_{\gamma\gamma}$, or $c_{H\tilde{W}}$, $c_{H\tilde{W}B}$, $c_{H\tilde{B}}$), after relaxing the above assumption.

On a longer term, a combination with other Higgs boson decay channels can be considered. A global combination with results from the top-quark sector, Standard Model, and flavour physics would also provide an global overview, with unprecedented accuracy on many parameters, and could undisclose hints of new physics.

Appendices

Appendix A

Layout of distorted geometries

This Appendix lists the regions impacted by a change in material budget in for various distorted geometry configurations. The impacted areas are highlighted in blue.

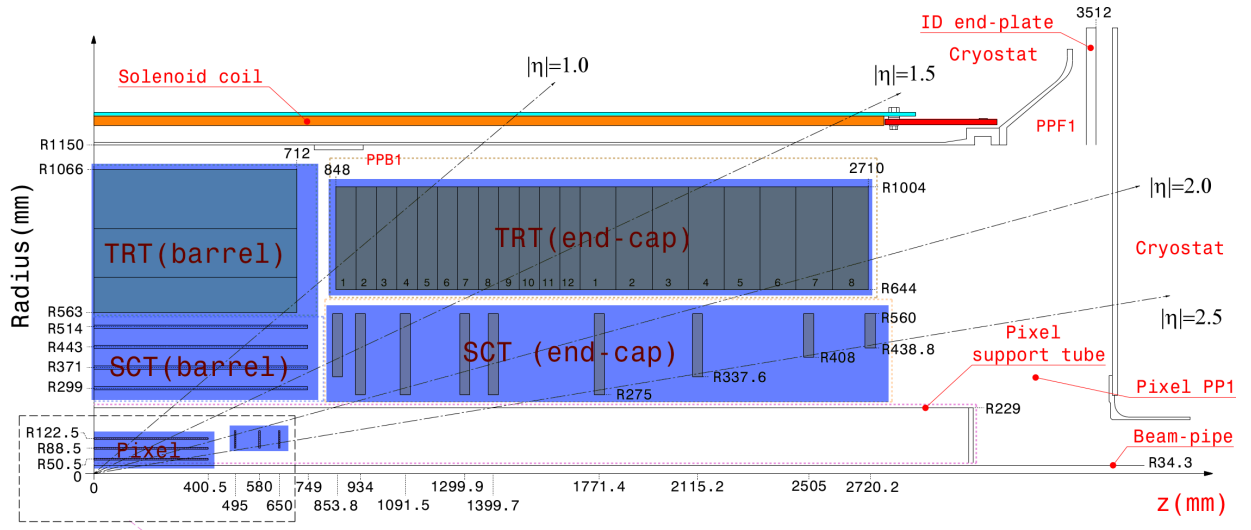


Figure A.1 – Impacted area of distorted geometry configuration “A”: scaling of material budget by 5 % in all the ID. [148]

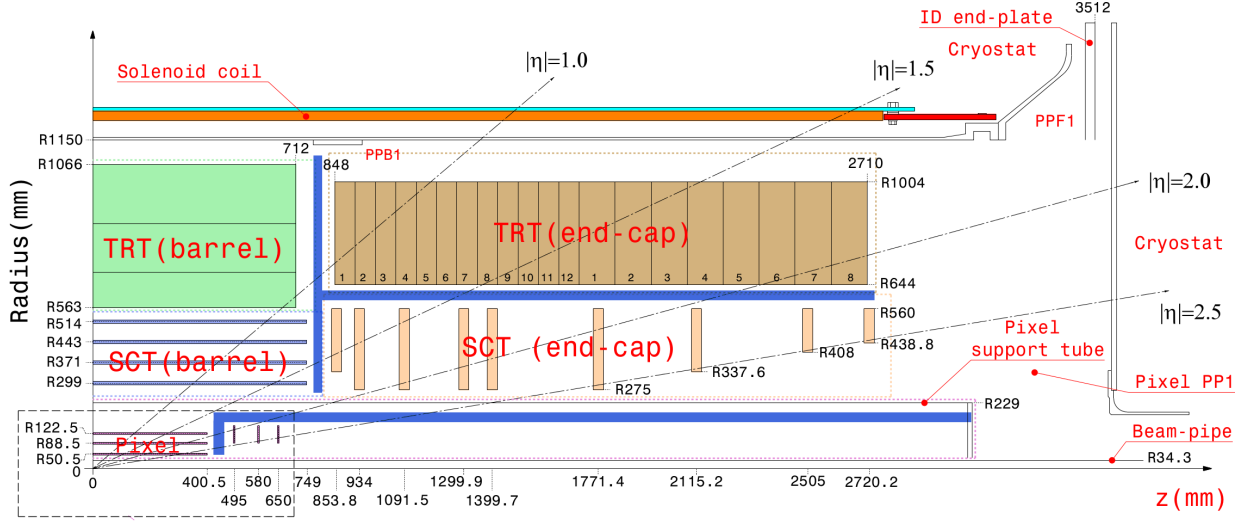


Figure A.2 – Impacted area of distorted geometry configuration “C”+D”: change of material budget by $0.1 X_0$ of the IBL and SCT services. [148]

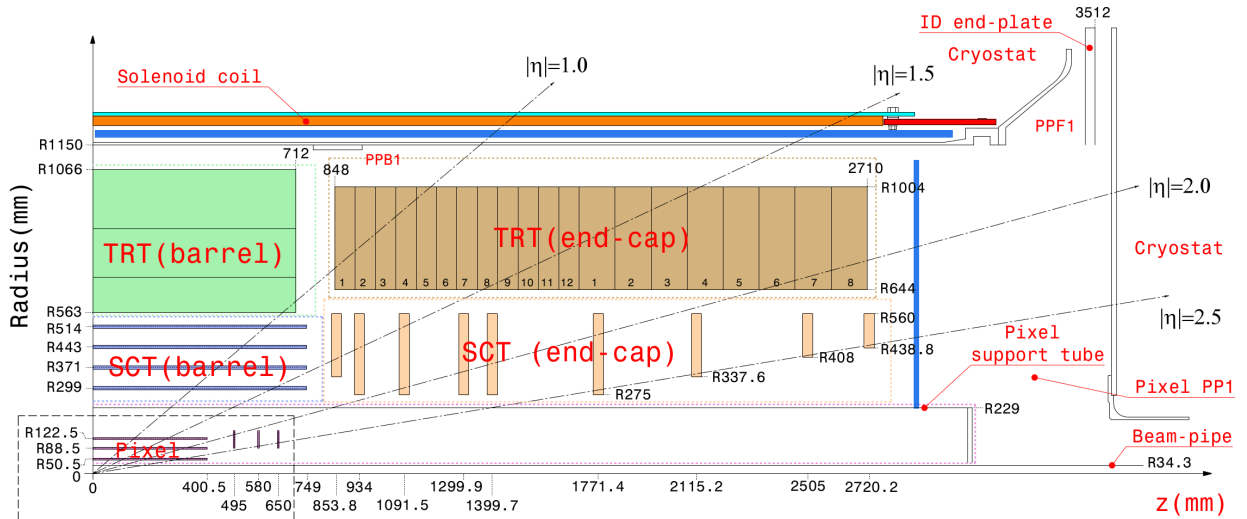


Figure A.3 – Impacted area of distorted geometry configuration “E”+L”: change of material budget by $0.075 X_0$ of the endcap SCT and TRT, and by $0.05 X_0$ of the cryostat in the barrel. [148]

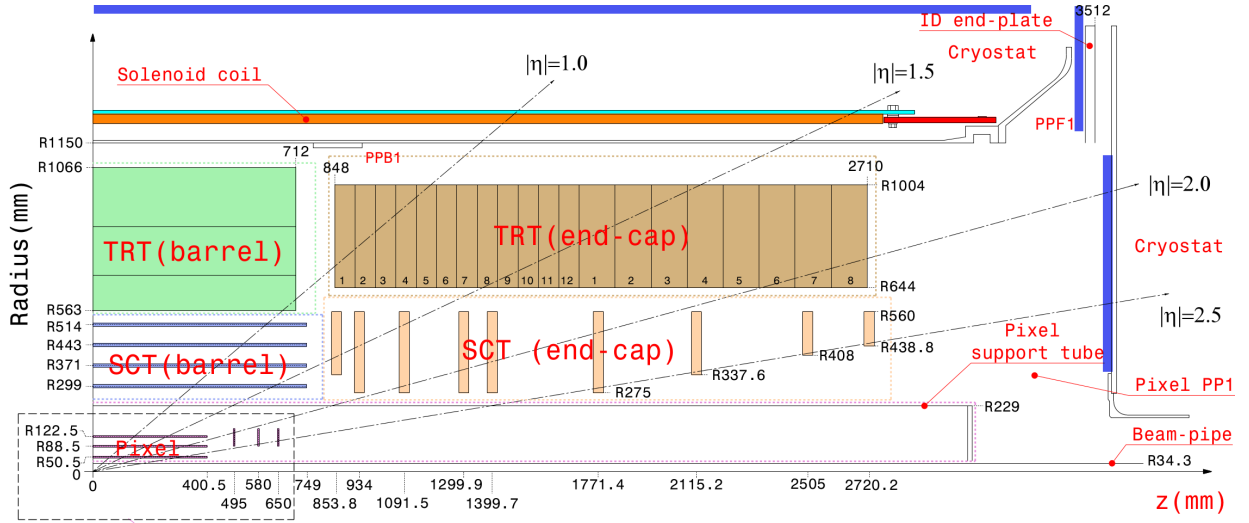


Figure A.4 – Impacted area of distorted geometry configuration “F+M+X”: change of material budget by $0.075 X_0$ of the ID support, by $0.05 X_0$ between the PS and first sampling in the barrel, and by $0.3 X_0$ in front of the endcap calorimeter. [148]

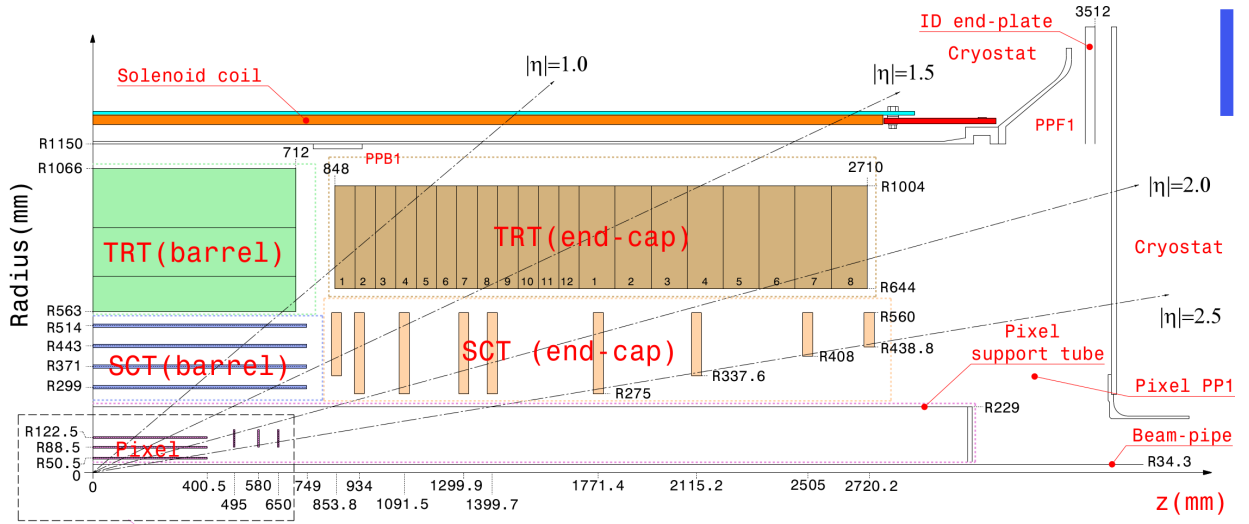


Figure A.5 – Impacted area of distorted geometry configuration “N”: change of material budget of $0.05 X_0$ between the PS and endcap calorimeter. [148]

Appendix B

STXS Stage-0 categorisation yields

Table B.1 – Expected number of SM Higgs boson events with $m_H = 125$ GeV at an integrated luminosity of 139 fb^{-1} and $\sqrt{s} = 13 \text{ TeV}$ in each reconstructed event signal ($115 < m_{4\ell} < 130 \text{ GeV}$) and sideband ($m_{4\ell}$ in $105\text{-}115 \text{ GeV}$ or $130\text{-}160 \text{ GeV}$ for $ZZ^{(*)}$, $130\text{-}350 \text{ GeV}$ for tXX) category, shown separately for each Stage-0 production bin. The ggF and bbH yield are shown separately but both contribute to the same (ggF) production bin, and ZH and WH are reported separately but are merged together for the final result. Statistical and systematic uncertainties, including those for theory, are added in quadrature. Contributions that are below 0.2 % of the total signal in each reconstructed category are not shown and replaced by “–”. [245]

Reconstructed event category	ggF	VBF	SM Higgs boson production mode			
			WH	ZH	$ttH + tH$	bbH
Signal	$115 < m_{4\ell} < 130 \text{ GeV}$					
$0j\text{-}p_T^{4\ell}\text{-Low}$	24.6 ± 3.1	0.077 ± 0.010	0.0194 ± 0.0035	0.0131 ± 0.0024	–	0.18 ± 0.09
$0j\text{-}p_T^{4\ell}\text{-Med}$	76 ± 8	1.18 ± 0.14	0.39 ± 0.05	0.36 ± 0.04	–	0.8 ± 0.4
$0j\text{-}p_T^{4\ell}\text{-High}$	0.132 ± 0.032	0.0302 ± 0.0033	0.064 ± 0.006	0.161 ± 0.015	0.00065 ± 0.00025	–
$1j\text{-}p_T^{4\ell}\text{-Low}$	30 ± 4	2.03 ± 0.11	0.52 ± 0.05	0.306 ± 0.031	0.0074 ± 0.0016	0.40 ± 0.20
$1j\text{-}p_T^{4\ell}\text{-Med}$	17.5 ± 2.8	2.65 ± 0.16	0.52 ± 0.05	0.354 ± 0.035	0.0087 ± 0.0020	0.09 ± 0.05
$1j\text{-}p_T^{4\ell}\text{-High}$	3.7 ± 0.8	0.93 ± 0.07	0.167 ± 0.014	0.154 ± 0.013	0.0047 ± 0.0011	0.012 ± 0.006
$1j\text{-}p_T^{4\ell}\text{-BSM-Like}$	0.90 ± 0.23	0.268 ± 0.019	0.065 ± 0.010	0.052 ± 0.008	0.0017 ± 0.0006	0.0008 ± 0.0004
$2j$	23 ± 5	8.0 ± 0.5	1.86 ± 0.14	1.44 ± 0.11	0.47 ± 0.05	0.28 ± 0.14
$2j\text{-BSM-like}$	1.9 ± 0.6	1.05 ± 0.05	0.119 ± 0.013	0.110 ± 0.012	0.078 ± 0.007	0.0027 ± 0.0014
$VH\text{-Lep-enriched}$	0.046 ± 0.017	0.0191 ± 0.0031	0.80 ± 0.06	0.211 ± 0.017	0.172 ± 0.015	0.0026 ± 0.0013
$ttH\text{-Had-enriched}$	0.13 ± 0.13	0.0162 ± 0.0033	0.0142 ± 0.0024	0.044 ± 0.007	0.73 ± 0.08	0.017 ± 0.009
$ttH\text{-Lep-enriched}$	0.0008 ± 0.0012	0.00019 ± 0.00014	0.0039 ± 0.0024	0.0023 ± 0.0014	0.40 ± 0.04	–
Sideband	$105 < m_{4\ell} < 115 \text{ GeV}$ or $130 < m_{4\ell} < 160 \text{ GeV}$					
$SB\text{-}0j$	4.4 ± 0.5	0.058 ± 0.009	0.103 ± 0.012	0.040 ± 0.005	–	0.046 ± 0.024
$SB\text{-}1j$	2.30 ± 0.29	0.256 ± 0.023	0.100 ± 0.011	0.060 ± 0.006	0.0056 ± 0.0012	0.021 ± 0.011
$SB\text{-}2j$	1.17 ± 0.25	0.40 ± 0.05	0.116 ± 0.014	0.089 ± 0.010	0.109 ± 0.010	0.016 ± 0.008
$SB\text{-}VH\text{-Lep-enriched}$	0.019 ± 0.008	0.0029 ± 0.0010	0.086 ± 0.008	0.090 ± 0.008	0.066 ± 0.007	0.0013 ± 0.0007
	$105 < m_{4\ell} < 115 \text{ GeV}$ or $130 < m_{4\ell} < 350 \text{ GeV}$					
$SB\text{-}tXX\text{-enriched}$	0.0009 ± 0.0015	0.00015 ± 0.00015	0.00042 ± 0.00016	0.00041 ± 0.00016	0.064 ± 0.008	0.00008 ± 0.00008
Total	186 ± 14	17.0 ± 0.8	5.0 ± 0.4	3.48 ± 0.25	2.12 ± 0.18	1.9 ± 1.0

Synthèse en français

C.1 Contexte expérimental et théorique

C.1.1 Le Modèle Standard et le boson de Higgs

Le Modèle Standard est, à ce jour, la meilleure théorie décrivant les propriétés des particules élémentaires et leurs interactions (Figure C.1). La découverte d’une nouvelle particule à une masse d’environ 125 GeV en 2012 par les expériences ATLAS [37] et CMS [38] au LHC, rapidement identifiée comme le boson de Higgs, est venue compléter le Modèle Standard, confortant encore sa validité. Cette particule, pierre angulaire du modèle et prédite dans les années 1960 [17–21], est associée au mécanisme de brisure spontanée de symétrie électro-faible, donnant une masse à toutes les autres particules massives.

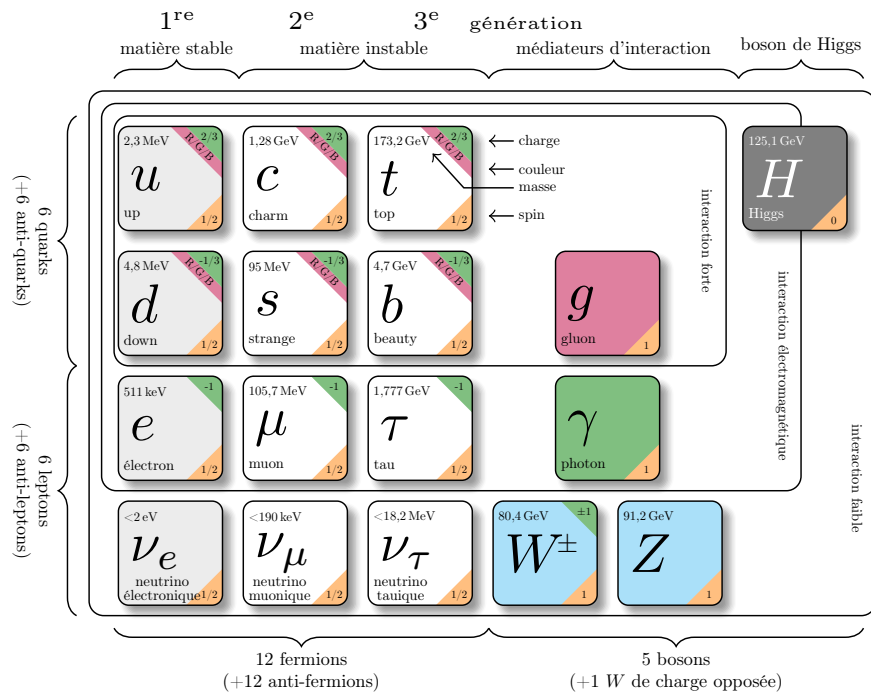


Figure C.1 – Les particules élémentaires connues à ce jour (2019) et leur classification. Les huit gluons sont représentés comme une seule particule. Adapté de [269].

Dans les mois qui suivirent, ses propriétés telles que les sections efficaces de ses différents modes de production ($\sigma_{i \rightarrow H}$), ses rapports d'embranchement¹ ($BR_{H \rightarrow f}$), ses couplages aux autres particules, sa masse et sa spin-parité (J^{PC}) ont été mesurées avec précision et se sont montrées compatibles avec les valeurs prédites par le Modèle Standard pour le boson de Higgs [53–70], confirmant que la particule découverte est bien le boson de Higgs du Modèle Standard.

Bien qu'il se soit révélé d'une redoutable précision pour de nombreuses mesures, le Modèle Standard ne peut être le modèle ultime de la physique des particules : il ne peut expliquer, par exemple, l'origine et la nature de la matière noire, cinq fois plus abondante que la matière ordinaire dans l'Univers, ou encore l'asymétrie entre matière et antimatière (asymétrie baryonique). Ces limites impliquent l'existence d'une théorie plus générale, *au-delà* du Modèle Standard (« Beyond the Standard Model », ou BSM en anglais). La recherche d'effets non prévus par le Modèle Standard est donc maintenant l'un des principaux points d'attention de la physique des particules. Récemment découvert, le boson de Higgs est un objet d'étude de choix et l'étude de plus en plus précise de ses propriétés pourrait permettre de détecter des variations de plus en plus fines par rapport aux prédictions théoriques.

C.1.2 Phénoménologie du boson de Higgs

Le boson de Higgs est produit au LHC par plusieurs processus. Le principal est la fusion de gluon (ggF, Figure C.2a), comptant pour 87 % de la section efficace totale de production. A cause de la présence de gluons dans l'état initial, émettant facilement un autre gluon, l'état final dans ce mode de production peut inclure un ou plusieurs jets additionnels. Le second est la fusion de bosons vecteurs (VBF, Figure C.2b) avec 6,8 %. Ce mode est caractérisé par deux jets à grande rapidité et une faible activité hadronique dans la région centrale. Le troisième est la radiation d'un boson de Higgs par un boson vecteur (VH, Figure C.2c) avec 4,1 %. Ce mode est caractérisé par la présence de deux jets ou de deux leptons de masse invariante proche de la masse des bosons Z ou W , et correspondant à la désintégration du boson vecteur associé. Le quatrième est la production associée à une paire de quarks top (ttH , Figure C.2d) avec 0,9 %. Une analyse telle que développée à la Section C.4 étudiant la production d'un boson de Higgs par le mode VBF aura donc un bruit de fond composé d'événements VH dans lesquels le boson associé se désintègre en deux jets (« VH -had »), mais aussi d'événements ggF dans lesquels les gluons ont émis deux jets supplémentaires.

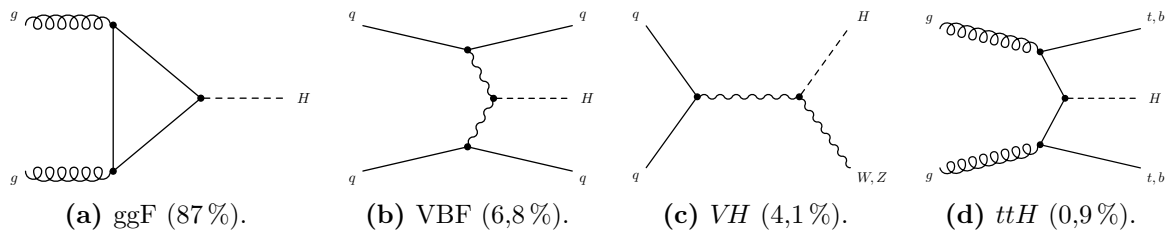


Figure C.2 – Diagrammes dominants pour la production du boson de Higgs.

Étant instable avec une durée de vie de l'ordre de 10^{-22} s, le boson de Higgs ne peut être directement détecté et est reconstruit à partir de ses produits de désintégration mesurés dans un détecteur. Bien que son mode de désintégration dominant soit la paire de quarks beauté (58 %), la présence de bruits de fond importants liés à l'environnement hadronique rend la reconstruction complexe et les analyses moins précises que des modes de désintégrations moins fréquents mais avec une signature claire. Les modes dits « de précision » sont ainsi les désintégrations $H \rightarrow \gamma\gamma$ (0,23 %) et $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ (0,0124 %), offrant un état final univoque et très bien

¹↑ Les rapports d'embranchement, ou largeurs partielles, mesurent la probabilité qu'une particule instable (comme le boson de Higgs) se désintègre en un état final donné.

reconstruit, malgré leur plus faible précision statistique. D'autres modes notables ont un rapport d'embranchement intermédiaire et une signature non-ambiguë ($H \rightarrow \tau^+\tau^-$, 6,3 %), mais la présence de neutrinos dans l'état final, non détectés, rend la reconstruction du boson de Higgs partielle. L'étude des couplages du boson de Higgs décrite à la Section C.4 se concentre ainsi sur la désintégration $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$, permettant une grande précision dans de nombreuses analyses.

C.1.3 Spin-parité (CP) du boson de Higgs : motivation et résultats récents

L'une des conditions nécessaires à l'asymétrie baryonique mentionnée plus haut est la violation de la symétrie CP , c'est-à-dire que la matière et l'antimatière aient un comportement différent par les lois de la physique. La violation de symétrie CP a déjà été observée dans le secteur des quarks [23, 24] mais son amplitude n'est pas suffisante pour expliquer l'asymétrie baryonique observée dans l'Univers. La recherche de violation de symétrie CP doit donc être explorée dans d'autres secteurs, en particulier celui du boson de Higgs.

Le Modèle Standard prédit la valeur $J^{PC} = 0^{++}$, et les études menées au Run 1 ont permis d'éliminer beaucoup d'autres valeurs hypothétiques [61, 62, 66–70], confirmant 0^{++} comme valeur la plus probable (Figure C.3). Cependant, à cause d'une précision statistique limitée, la plupart de ces études n'ont pu étudier et rejeter que des états purs de CP , laissant ouverte la possibilité d'un état mixte pair et impair par CP . Avec environ dix fois plus de données sur le Higgs qu'au Run 1, le Run 2 (2015 – 2018) permet de raffiner les études et d'étudier de tels états mixtes. La Figure C.4a montre l'exemple de l'analyse de couplages mixtes dans le canal $H \rightarrow \tau^+\tau^-$ d'ATLAS (Run 1), et la Figure C.4b l'exemple de l'analyse dans le canal $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ utilisant les données récoltées par ATLAS en 2015 et 2016.

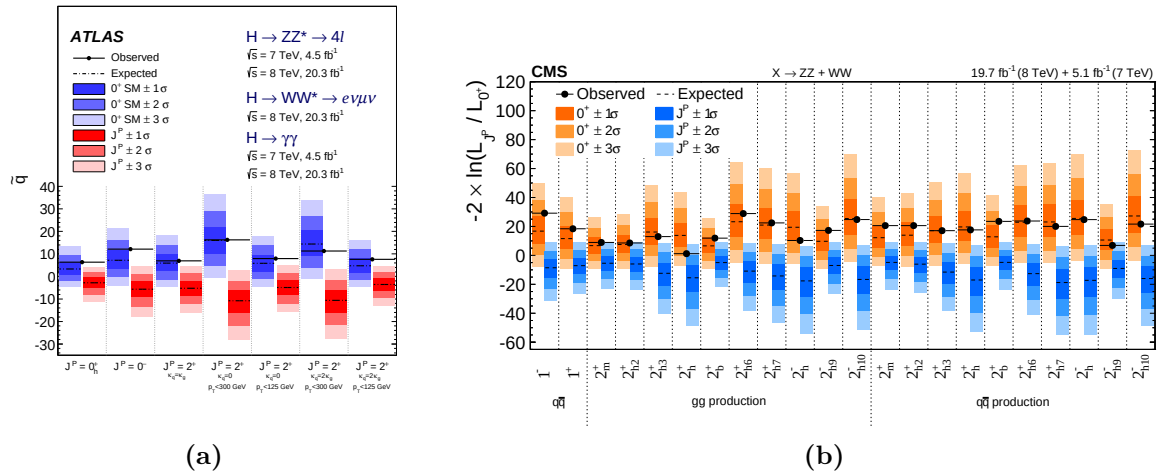


Figure C.3 – Test de vraisemblance de différentes hypothèses de spin-parité J^P pure par rapport à l'hypothèse du Modèle Standard ($J^P = 0^{++}$). Les analyses (a) d'ATLAS [68] et (b) de CMS [70] utilisent toutes deux l'intégralité des données enregistrées au Run 1.

Le travail mené dans cette Thèse et résumé à la Section C.4 utilise la totalité des données récoltées durant le Run 2 et améliore les résultats de ces précédentes analyses. L'analyse construit des observables dédiées aux mesures de CP (*observables optimales* similaires à l'analyse $H \rightarrow \tau^+\tau^-$, voir Section C.4.2) et interprète les résultats dans un cadre de théorie effective (Section C.1.4). Du fait de sa haute précision et de son état final à quatre objets permettant la construction de plusieurs variables angulaires sensibles aux effets de CP , le canal $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ est employé (voir Section C.4.1). La mesure d'un état significativement mixte serait immédiatement liée à un effet BSM, le Modèle Standard prédisant un état pur 0^{++} .

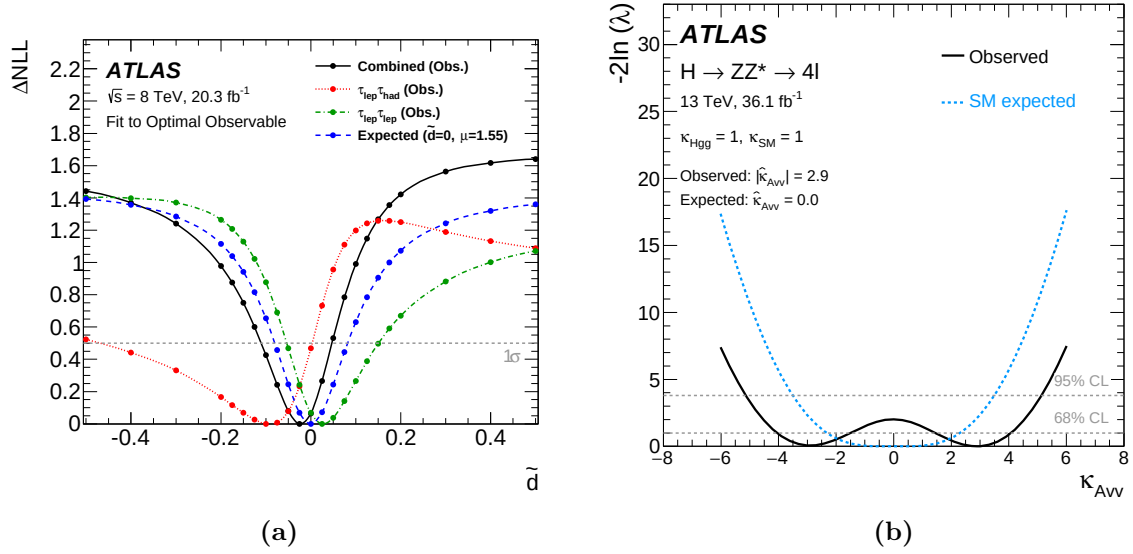


Figure C.4 – Exemples de limites posées sur le mixte de CP du boson de Higgs. (a) Scans de vraisemblance attendu (bleu) et observé (noir) pour le paramètre $\tilde{d}$, encodant de manière unique les effets impairs par CP entre le boson de Higgs et les bosons Z , W et γ . L’analyse dans le canal $H \rightarrow \tau^+\tau^-$ utilise les $20,3 \text{ fb}^{-1}$ de données récoltées en 2012 à $\sqrt{s} = 8 \text{ TeV}$ par ATLAS et utilise une *observable optimale* (voir Section C.4.2) dédiée pour les mesures de CP [71] comme variable discriminante. (b) Scans de vraisemblance attendu (bleu) et observé (noir) pour le paramètre κ_{AVV} du cadre « Higgs Characterisation » (semblable aux théories effectives, voir Section C.1.4) en utilisant uniquement les informations de sections efficaces. L’analyse dans le canal $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ utilise les $36,1 \text{ fb}^{-1}$ de données récoltées en 2015 et 2016 par ATLAS à $\sqrt{s} = 13 \text{ TeV}$ [178].

La structure des couplages du boson de Higgs aux fermions et aux bosons étant différente, ces deux catégories fournissent des analyses complémentaires à la recherche de couplages anomaux. Puisque des effets impairs par CP n’ont pas encore été observés, leur éventuelle observation est plus probable avec des événements à haut transfert d’impulsion (Q^2). Ainsi une production via les modes VBF, VH ou ttH (Q^2 de plusieurs centaines de GeV) sera préférée par rapport à ggF ($Q^2 \sim m_H \approx 125 \text{ GeV}$). Étant donné les sections efficaces de chacun de ces modes (Section C.1.2) et les contraintes expérimentales pour leur identification, l’analyse décrite à la Section C.4 se concentre sur les événements produits par le mode VBF. Cette analyse est donc essentiellement sensible aux couplages entre le boson de Higgs et les bosons Z , W et γ .

C.1.4 Au-delà du Modèle Standard : théories effectives

Afin de faciliter la comparaison et la combinaison des différentes analyses, en particulier du boson de Higgs, plusieurs cadres ont été mis en place. Au Run 1, le *cadre* κ donnait des interprétations des mesures en termes de modificateurs de couplages (κ). Au Run 2, grâce à une statistique dix fois plus élevée pour le boson de Higgs, des cadres comme les théories effectives, offrant plus de précision et de généralité, ont pu être élaborés. Une théorie effective consiste en une expansion systématique des opérateurs² du Lagrangien du Modèle Standard dans des dimensions canoniques³ supérieures : $\mathcal{L}_{\text{EFT}} = \mathcal{L}_{\text{SM}} + \sum_i \frac{c_i^{d=5}}{\Lambda} \mathcal{O}_i^{d=5} + \sum_i \frac{c_i^{d=6}}{\Lambda^2} \mathcal{O}_i^{d=6} + \dots$. Chaque opérateur

²↑ Les opérateurs définissent la structure des couplages entre particules, ou de manière illustrée les vertex d’interaction dans les diagrammes de Feynman. Par exemple le développement du terme $(D_\mu \phi)^\dagger (D^\mu \phi)$ (avec $D_\mu = \partial_\mu - ig \frac{\sigma_a}{2} W_\mu^a - ig' \frac{Y}{2} B_\mu$ où g et g' sont les constantes de couplage électro-faible et Y l’hypercharge) inclus dans le Lagrangien du Modèle Standard donne lieu aux opérateurs d’interaction HZZ , HWW et $H\gamma\gamma$.

³↑ Schématiquement, un champ fermionique a une dimension canonique (en unités d’énergie) $3/2$, un champ bosonique une dimension 1 et une dérivée d’espace-temps une dimension de 1. Les opérateurs du Modèle Standard

$\mathcal{O}_i$ rajoute une classe d'interaction entre particules (un nouveau vertex), dont le médiateur est de masse trop élevée pour être observable directement au LHC. L'amplitude de ces opérateurs est paramétrée par un *coefficient de Wilson* c_i et est réduite d'un certain facteur Λ (typiquement 1 TeV) supplémentaire à chaque ordre successif, assurant la convergence de l'expansion. Cette échelle d'énergie Λ définit également la limite de validité de la théorie en supposant qu'aucune nouvelle particule n'existe en dessous de celle-ci.

Les opérateurs de dimension impaire violent la conservation du nombre leptonique ou baryonique [195, 200]. Ceci n'ayant pas été observé à ce jour, ces opérateurs sont généralement négligés. Les termes dominants dans l'expansion sont donc les termes d'ordre 6, au nombre de 2499 et avec une amplitude diminuée d'un facteur Λ^2 par rapport aux opérateurs du Modèle Standard. Tout ensemble complet et non-redondant de 2499 coefficients de Wilson forme une *base* de la théorie effective. Les bases sont toutes équivalentes mais plus faciles à utiliser dans certains contextes. Dans l'étude menée à la Section C.4, deux bases sont utilisées : l'une dédiée à l'étude du boson de Higgs et nommée *Higgs* [202], l'autre plus générale (utilisée pour combiner les résultats avec d'autres secteurs d'étude) et nommée *Warsaw* [195–199].

Afin de simplifier l'analyse, deux hypothèses supplémentaires sont formulées : l'équivalence des familles fermioniques (*universalité*) est imposée, réduisant le nombre d'opérateurs à 76, et en se restreignant aux opérateurs donnant des couplages impairs par *CP* entre le boson de Higgs et les bosons Z , W et γ , ce nombre est réduit à 3. L'étude présentée dans cette Thèse se restreint donc à trois paramètres, dénotés $\{\tilde{c}_{zz}, \tilde{c}_{z\gamma}, \tilde{c}_{\gamma\gamma}\}$ dans la base *Higgs*

$$\mathcal{L}_{\text{Higgs}}^{\text{dim-6}} \supset \left[\tilde{c}_{\gamma\gamma} \frac{e^2}{4} A_{\mu\nu} \tilde{A}_{\mu\nu} + \tilde{c}_{z\gamma} \frac{e\sqrt{g^2 + g'^2}}{2} Z_{\mu\nu} \tilde{A}_{\mu\nu} + \tilde{c}_{zz} \frac{g^2 + g'^2}{4} Z_{\mu\nu} \tilde{Z}_{\mu\nu} \right] \frac{h}{v} \quad (\text{C.1})$$

et $\{c_{H\tilde{W}}, c_{H\tilde{B}}, c_{H\tilde{W}B}\}$ dans la base *Warsaw*, tous les autres coefficients étant supposé nuls (donc ne donnant pas de déviation par rapport au Modèle Standard). Ces deux ensembles de coefficients peuvent être interchangés au moyen d'une matrice de conversion 3×3 [202].

C.2 Le LHC et le détecteur ATLAS

C.2.1 Le LHC (Grand Collisionneur de Hadrons)

Le LHC [92] (Large Hadron Collider) est un collisionneur de hadrons (protons) d'environ 27 km de circonférence situé au CERN. Quatre expériences sont installées en différents points du LHC : ATLAS (A Toroidal LHC ApparatuS, détaillée à la Section C.2.2) [97] et CMS (Compact Muon Solenoid) [98], deux expériences généralistes avec un large spectre d'études (mesures de précisions du Modèle Standard, étude du quark top, du boson de Higgs, recherche de la super-symétrie, *etc.*) ; LHCb [99], dédié à l'étude du quark b ; et ALICE (A Large Ion Collider Experiment) [100] orienté vers l'étude du plasma quark-gluon.

Après une première période de prise de données à 7 TeV et 8 TeV entre 2010 et 2012 qui permit la découverte du boson de Higgs, le LHC a achevé une deuxième période (Run 2) entre 2015 et 2018 à l'énergie record de 13 TeV et délivrant six fois plus de collisions qu'au Run 1. Un autre paramètre d'importance au LHC est sa luminosité, c'est-à-dire le nombre de collisions produites (par seconde). Une haute luminosité permet la recherche de phénomènes plus rares, comme par exemple la production d'un boson de Higgs, mais nécessite plus de collisions à chaque croisement de faisceau, induisant un bruit de fond énergétique dans le détecteur et pénalisant les performances (Section C.3.2)). Le nombre d'interactions supplémentaires se nomme *empilement*

ont une dimension canonique de 4.

et avait une valeur moyenne d'environ 20 pendant le Run 1 avec un pic vers 35, atteignant une moyenne de 34 au Run 2 pour un pic à plus de 65 [108] (Figure C.5).

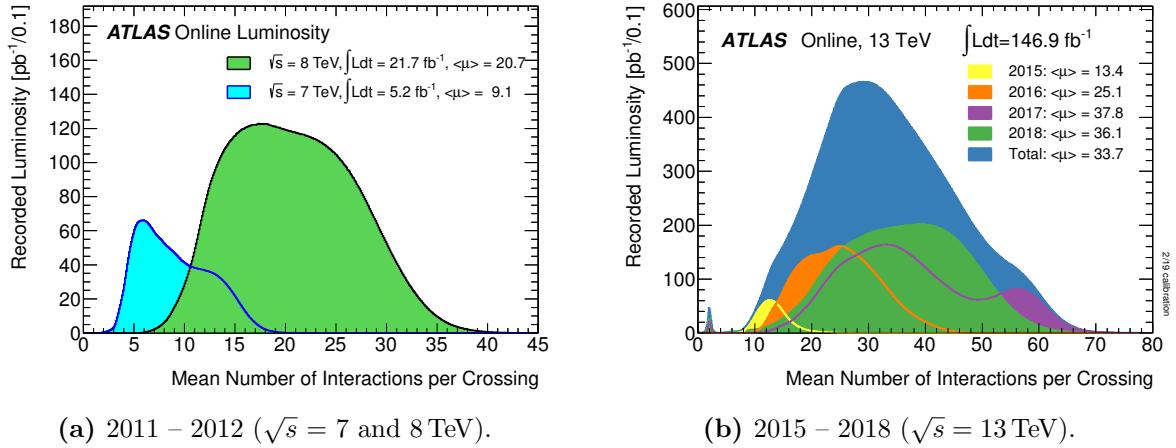


Figure C.5 – Nombre moyen d’interactions par croisement de faisceaux enregistré par ATLAS pour chaque année de prise de données et pondéré par la luminosité instantanée. [108]

C.2.2 Le détecteur ATLAS

Le détecteur ATLAS [97], représenté sur la Figure C.6, suit le schéma classique d’un détecteur de physique des particules. La partie la plus proche du point d’interaction est composée d’un trajectographe baignant dans un intense champ magnétique solénoïdal de 2 T et permettant de mesurer la position et l’impulsion des particules chargées produites lors de la collision. Viennent ensuite les calorimètres électromagnétique puis hadronique, mesurant l’énergie des électrons et photons pour le premier, et des hadrons (protons, neutrons, pions, kaons) pour le second en les absorbant complètement. Les couches les plus externes du détecteur sont constituées d’un spectromètre à muons (seules particules à pouvoir traverser les calorimètres) et sont baignées dans un champ magnétique toroïdal permettant de mesurer leurs trajectoires. Le détecteur est conçu pour couvrir la (quasi-)totalité de l’angle solide, permettant la mesure de toutes les particules détectables. Des systèmes dédiés (non représentés sur le schéma) permettent de mesurer la luminosité et donc le nombre d’événements produits.

Cet agencement permet de mesurer et d’identifier chaque particule stable (ou quasi-stable avec une durée de vie $\gtrsim 10 \text{ ns}$) : électrons, photons, muons et pions chargés. En raison des propriétés de l’interaction forte, les quarks et gluons ne peuvent être observés individuellement et « s’hadronisent » lors de leur propagation, formant ainsi des *jets* de particules. D’autres particules, telles que les neutrinos ou des particules hypothétiques, sont stables mais interagissent trop peu avec la matière pour pouvoir être détectées. Comme l’énergie et l’impulsion totales sont conservées lors des collisions, et puisque toutes les particules détectables sont mesurées, ces particules non-détectables se manifestent par un *manque d’énergie* dans une certaine direction dans le détecteur, indiquant leur présence.

Pour permettre des analyses précises (par exemple des propriétés du boson de Higgs), les particules détectées doivent également être mesurées précisément. Les différents systèmes du détecteur doivent donc être minutieusement étalonnés. Le calorimètre électromagnétique donne une mesure de l’énergie des photons et électrons, tous deux intervenant dans les canaux de désintégration du Higgs $H \rightarrow \gamma\gamma$ et $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$, et revêt donc une importance particulière. Une description plus détaillée de celui-ci est donné à la Section C.2.3.

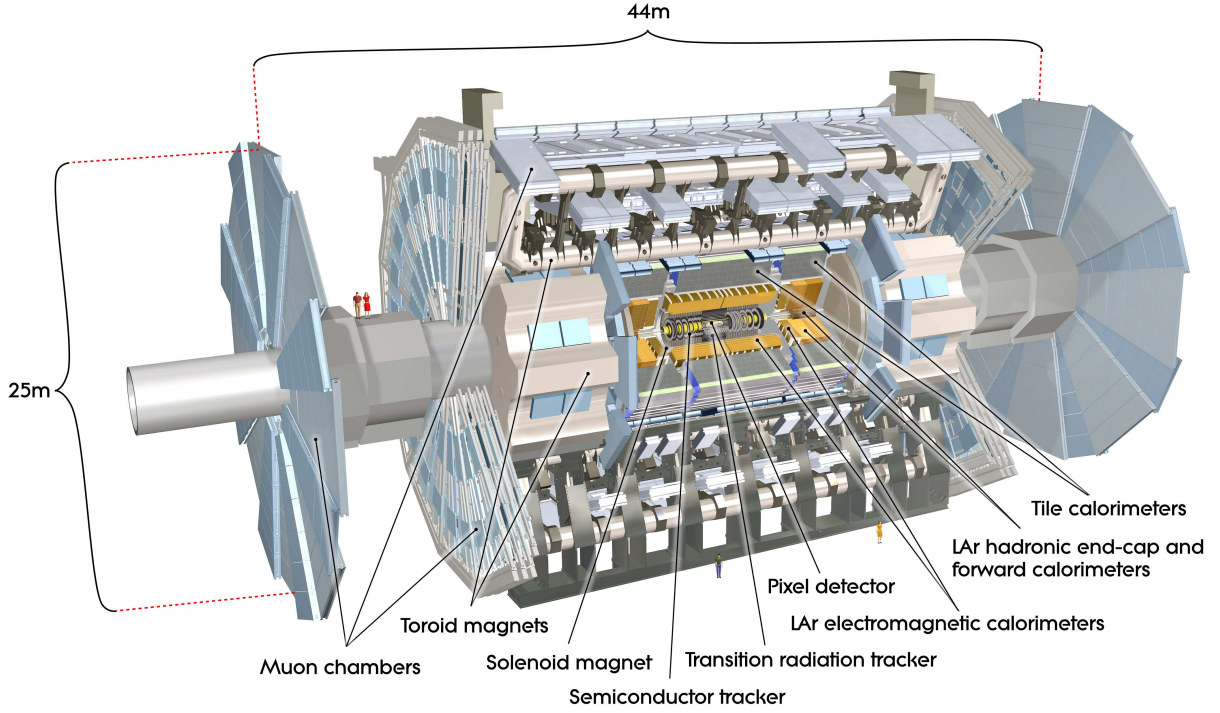


Figure C.6 – Vue éclatée du détecteur ATLAS. [97]

C.2.3 Aperçu du calorimètre électromagnétique

Le calorimètre électromagnétique d'ATLAS [113] est divisé en plusieurs éléments qui, combinés, couvrent l'intégralité de l'angle solide jusqu'à la pseudo-rapidité ⁴ $|\eta| < 4,9$. La région « de précision » (jusqu'à $|\eta| = 2,5$) est constituée d'une partie appelée *tonneau* ($|\eta| < 1,48$) complétée par une partie *bouchon* de chaque côté ($1,375 < |\eta| < 2,5$). La zone $1,35 < |\eta| < 1,50$ dans laquelle se fait la transition entre le tonneau et les bouchons est un peu moins précise à cause d'effets géométriques et d'une grande quantité de matériel passif lié aux câbles de services (alimentation, refroidissement, électronique de lecture) des détecteurs internes.

Le calorimètre électromagnétique est un calorimètre à échantillonnage utilisant de l'argon liquide comme matériel actif, enchâssé entre des plaques d'absorbeurs en plomb. Le rôle de l'absorbeur est d'induire une *cascade* (ou *gerbe*) d'électrons et de photons de plus en plus basse énergie lors du passage d'une particule électromagnétique, jusqu'à absorption totale. Le passage de ces particules va ioniser le matériel actif et, sous l'effet d'une haute tension, induire un courant formant un signal électrique mesurable. Sa forme typique en accordéon (Figure C.7) assure une homogénéité en fonction de l'angle azimutal ϕ . L'accordéon a une profondeur d'environ 47 cm, ce qui correspond à au moins 22 longueurs de radiation ⁵ (X_0).

Comme décrit à la Figure C.7, le calorimètre est segmenté en trois couches concentriques (dénotées L1 à L3, pour *layer*) ce qui permet de déduire plus précisément la direction des gerbes électromagnétiques non associées à des traces (comme dans le cas des photons). La deuxième couche (L2), avec des $16 X_0$, recueille la majeure partie du dépôt énergétique. Sa lecture est faite par cellule de taille $\Delta\eta \times \Delta\phi = 0,025 \times 0,025$, définissant l'unité de base dans le calorimètre. La première (L1) a une profondeur d'environ $4,3 X_0$ et est hautement segmentée en η ($0,025/8$) pour distinguer les vrais photons des paires de photons très collimatés venant des désintégrations de

⁴↑ Le référentiel d'ATLAS utilise un système de coordonnées dérivé des coordonnées cylindrique dans lequel l'axe z est orienté le long du faisceau et l'angle polaire $\theta \in [0, \pi]$ est remplacé par la pseudo-rapidité $\eta = -\ln(\tan \theta/2)$.

⁵↑ La longueur de radiation X_0 est la distance de parcours moyenne dans un milieu pour laquelle un électron n'a plus qu'une fraction $1/e$ de son énergie initiale.

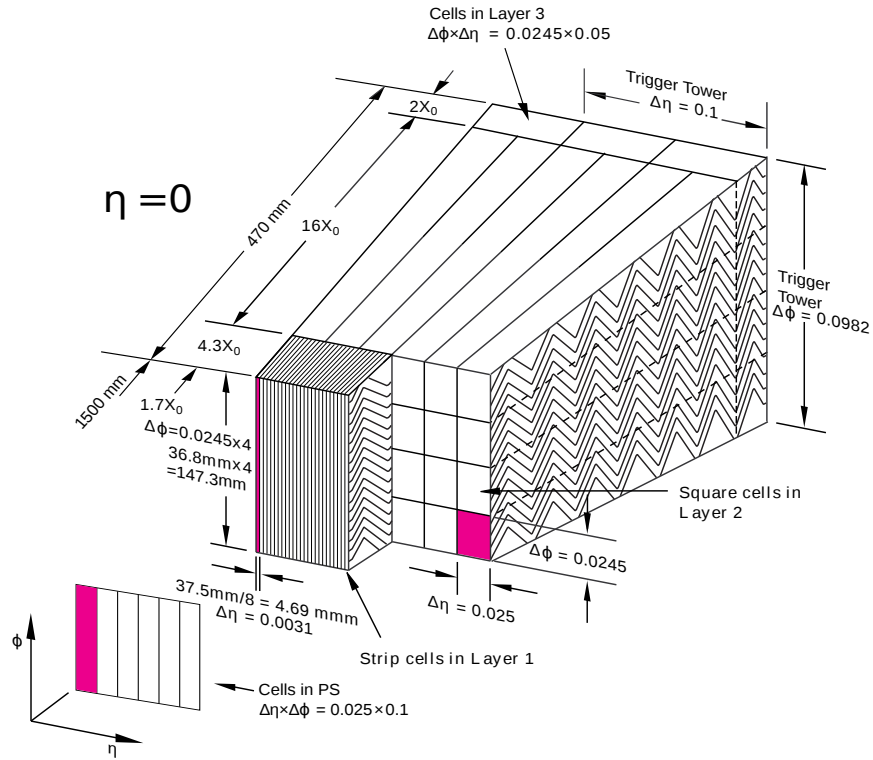


Figure C.7 – Schéma des cellules du calorimètre électromagnétique dans la région $0 < \eta < 0,15$. Le pré-échantillonneur à l'avant ainsi que les trois couches de l'accordéon sont visibles. [118]

pions neutres. En contrepartie, elle est moins segmentée en ϕ ($0,1$), pour attraper les photons émis par bremsstrahlung. La segmentation en η diminue à la fin des bouchons, passant de $0,025/8$ pour $1,5 < |\eta| < 1,8$ à $0,025/6$ pour $1,8 < |\eta| < 2,0$ et à $0,025/4$ pour $2,0 < |\eta| < 2,4$, et finalement à $0,025/1$ pour $2,4 < |\eta| < 2,5$. Sa haute segmentation en η la rend sujette à des effets de diaphonie électronique de l'ordre de $\sim 4\%$ entre cellules voisines. La troisième couche (L3) forme le reste de l'accordéon et sa profondeur est d'au moins $2X_0$ et est lue avec une granularité de $0,05 \times 0,025$. Un schéma des électrodes dans le tonneau est montré en Figure C.8. Une couche additionnelle nommée *pré-échantillonneur* et contenant uniquement de l'argon est placée devant le calorimètre et permet d'estimer l'énergie perdue dans le matériel passif (trajectographe, câbles, etc.) par les gerbes électromagnétiques avant d'entrer dans le calorimètre.

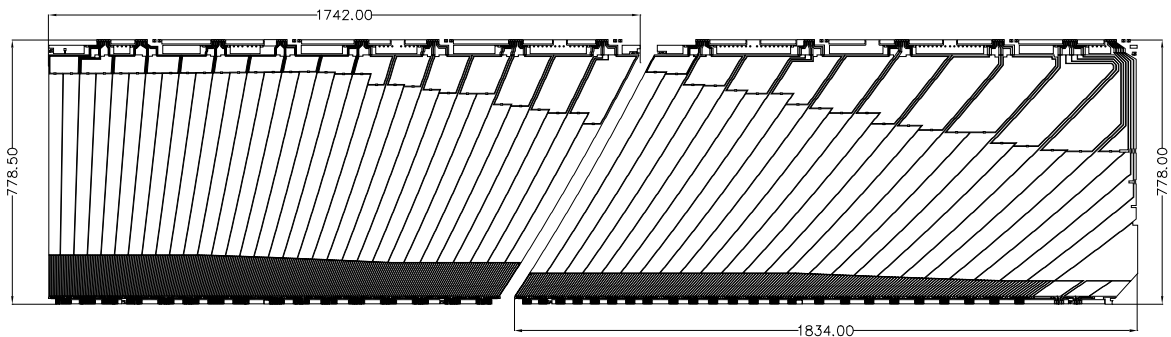


Figure C.8 – Schéma des électrodes du tonneau du calorimètre électromagnétique avant pliage en accordéon dans la région $0 < \eta < 1,48$. L'axe du faisceau est en bas de l'image. Le changement d'électrode se fait à $\eta = 0,8$. La profondeur des première et deuxième couches décroît lentement avec η pour garder une longueur d'absorption X_0 constante. [97]

C.3 Étalonage du calorimètre électromagnétique d'ATLAS

La procédure d'étalonnage en énergie des électrons a été développée au Run 1 [140] et a été mise à jour avec les premières données du Run 2 [141]. Elle inclut de nombreuses étapes, résumées dans la Figure C.9 : après l'entraînement sur des simulations d'un algorithme d'apprentissage déduisant l'énergie totale d'une gerbe à partir de ses propriétés dans chacune des couches et l'application de corrections sur les données (uniformité et homogénéité du matériel ainsi que l'étalonnage des différentes couches), les différences résiduelles observées entre les données et les simulations pour le processus $Z \rightarrow ee$ sont compensées en appliquant une échelle d'énergie aux électrons. Après cette dernière étape, l'énergie des électrons est connue entre le pour-cent et le pour-mille.

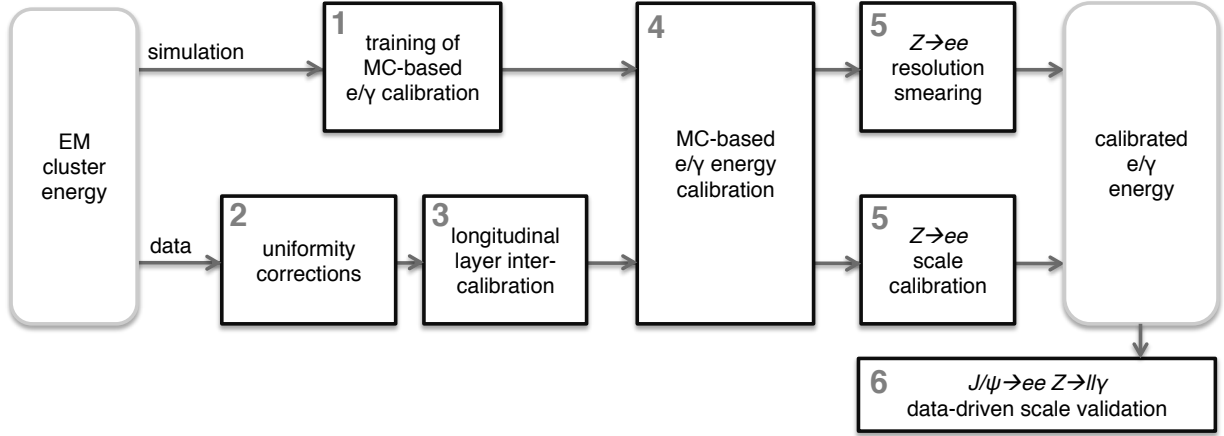


Figure C.9 – Schéma des différentes étapes de l'étalonnage en énergie des électrons. [140]

C.3.1 Étalonage relatif L1/L2 : objectifs et méthode

Le but de l'étalonnage relatif des couches L1 et L2 du calorimètre électromagnétique est de corriger les imprécisions d'étalonnage de l'électronique et des effets de diaphonie résiduels mal pris en compte, et ainsi de rééquilibrer l'énergie dans ces deux couches. Idéalement, cette correction doit donc être dérivée pour chaque cellule individuellement. Elle est également cruciale pour d'autres étapes de la calibration telle que l'entraînement de l'algorithme basé sur les simulations (dont l'une des variables d'entraînement est directement le ratio des énergies dans L1 et L2), l'étalonnage du pré-échantillonneur, ou la mesure du matériel passif en amont du calorimètre.

Au contraire des autres étapes de l'étalonnage, celle-ci ne peut que difficilement utiliser des électrons. En effet, ceux-ci forment une gerbe électromagnétique s'étalant sur plusieurs cellules et interagissent avec le matériel en amont du calorimètre, rendant l'extension longitudinale réelle de la gerbe inconnue. En conséquence, il est impossible de distinguer de réels effets d'étalonnage relatif entre L1 et L2 d'effets liés à la présence de matériel. Pour contourner ces deux problèmes, la méthode principale utilise des muons à la place des électrons. À des énergies de quelques dizaines de GeV, les muons peuvent être considérés comme des particules d'ionisation minimum (MIP) et présentent la propriété d'avoir un dépôt énergétique uniforme, et ce indépendamment de la présence de matériel en amont du détecteur ou de leur énergie. De plus, ces particules ne forment pas de gerbe et leur dépôt est donc très localisé spatialement.

Cette étude ne requiert donc qu'un échantillon raisonnablement pur de muons, ce qui peut être obtenu en sélectionnant des muons de la désintégration $Z \rightarrow \mu\mu$ (large résonance avec peu de bruits de fond). L'énergie déposée par ces muons et mesurée dans le calorimètre ne dépend que de la longueur de matériel actif (argon liquide) traversée. L'inconvénient des muons est que ce dépôt énergétique est faible (30 MeV à 60 MeV dans L1 et 240 MeV à 300 MeV dans L2, en fonction de

$|\eta|$) en comparaison du bruit du détecteur (15 MeV à 40 MeV dans L1 et 40 GeV à 80 MeV dans L2), donnant un rapport signal sur bruit entre 0,5 et 2 dans L1, et entre 3 et 4 dans L2. À cause d'une statistique finie, l'étude n'est pas réalisée pour chaque cellule, mais par région en $|\eta|$ (pas de différence attendue entre les parties $\eta < 0$ et $\eta > 0$, ce qui a été vérifié) de largeur 0,05 ou 0,1 (changé au cours de l'étude, voir Section C.3.3). Par ailleurs aucune différence significative n'est attendue en fonction de ϕ (tous les modules du calorimètre étant construits à l'identique), l'analyse est donc intégrée sur cette coordonnée.

Deux méthodes, identiques à l'analyse effectuée au Run 1 [140], sont utilisées pour extraire une valeur représentative de la distribution du dépôt énergétique des muons (Figure C.10a). La première, utilisée dans la suite, consiste à effectuer l'ajustement d'une distribution de Landau [167] (modélisant le dépôt d'une MIP dans une couche d'argon) convoluée à une gaussienne centrée en 0 (modélisant le bruit dans le calorimètre) sur la distribution en énergie (Figure C.10b) :

$$\text{Modèle}(E; \text{MPV}, \Gamma, \sigma) = \text{Landau}(E; \text{MPV}, \Gamma) * \text{Gaussienne}(E; \mu = 0, \sigma). \quad (\text{C.2})$$

On utilise ensuite la position du pic (*Most Probable Value*, MPV) de la Landau sous-jacente dans l'ajustement comme valeur d'énergie des muons. Cette méthode a l'inconvénient d'être sensible au bruit pouvant déformer la distribution. La seconde méthode s'affranchit de ce problème en calculant la moyenne de la distribution dans un intervalle tronqué (*Truncated Mean*, TM). Cette troncature est nécessaire pour éliminer les queues de distribution, plus sensibles au bruit. Cependant cette seconde méthode est plus *ad-hoc* et n'est pas directement reliée à une quantité physique bien définie.

Quelle que soit la méthode d'extraction de l'énergie des muons, une référence est nécessaire pour comparer les résultats. La même méthode est employée dans les données réelles et simulées pour extraire les énergies dans L1 et L2, et les facteurs d'étalonnage relatif sont ensuite calculés par le double ratio E_1/E_2 dans les données normalisé par cette même valeur dans la simulation :

$$\alpha_{1/2} = E_{1/2}^{\text{Data}} / E_{1/2}^{\text{MC}} \quad (\text{C.3})$$

où E_1 et E_2 sont les valeurs extraites comme mentionné ci-dessus, dans L1 et L2 respectivement.

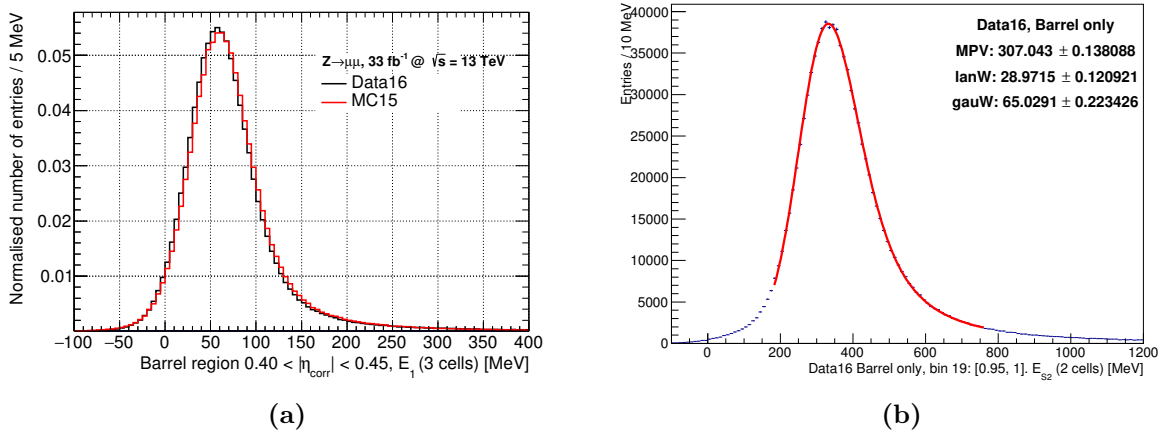


Figure C.10 – (a) Exemple de distribution du dépôt d'énergie des muons dans la première couche dans la région $0.40 < |\eta| < 0.45$ pour les données (noir) et la simulation (rouge). (b) Exemple d'ajustement Landau * Gaussien sur les données dans la deuxième couche dans la région $0.95 < |\eta| < 1.00$.

Il est nécessaire de noter qu'à cause des grands effets de diaphonie dans L1 (environ 5 % de chaque côté), l'énergie de chaque muon dans cette couche est déterminée par la somme des énergies dans trois cellules adjacentes dans la direction η . La cellule centrale est définie comme la plus proche de la trajectoire du muon reconstruite par les trajectographes et extrapolée à la première couche. Dans L2, à cause de la géométrie en accordéon, les muons partagent toujours

leur énergie entre deux cellules dans la direction ϕ . L'énergie associée au muon est donc celle de la cellule « centrale » (déterminée de manière similaire), à laquelle est ajoutée l'énergie de la cellule voisine avec la plus grande énergie dans la direction ϕ .

C.3.2 Étalonnage relatif L1/L2 : résultats initiaux et limites

La Figure C.11 présente les valeurs de MPV pour chacun des ajustements effectués dans chaque régions en $|\eta|$, pour L1, L2, données et simulation. Ces graphiques montrent plusieurs structures similaires entre données et simulations, pouvant s'expliquer par la géométrie des électrodes (Figure C.8). Par exemple, le saut dans les valeurs à $|\eta| = 0,8$ est lié au changement d'électrode : les cellules de L1 raccourcissent, donc la longueur d'argon liquide traversée par les muons est moindre, donc l'énergie collectée par le calorimètre diminue. Inversement dans L2, les cellules s'allongent, augmentant ainsi le dépôt dans le calorimètre. Dans les bouchons, les différents sauts dans L1 correspondent aux valeurs de $|\eta|$ pour lesquelles la granularité en η des cellules change.

En comparant données et simulation (Figure C.11), on note un comportement similaire pour les deux mesures, avec cependant un décalage stable d'environ 5 % dans L1 pour le tonneau, et variant entre 0 et 5 % dans L2 (tonneau et bouchons). Une différence plus large et irrégulière est notée dans les bouchons pour L1. On peut ensuite calculer le double ratio, et sa distribution en fonction de $|\eta|$ est montrée en Figure C.12. Alors que les résultats du Run 1 (noir) et ceux de 2016 (bleu) sont en bon accord dans le tonneau (la courbe rouge est décrite dans la Section C.3.3), de larges différences apparaissent dans les bouchons après $|\eta| = 2,0$: le résultat avec les données de 2016 diminue rapidement alors que le résultat du Run 1 augmente. Le calorimètre n'ayant pas changé entre le Run 1 et le Run 2, de telles différences ne sont pas attendues. Plusieurs études ont été menées pour essayer d'en trouver la cause et sont résumées dans la suite.

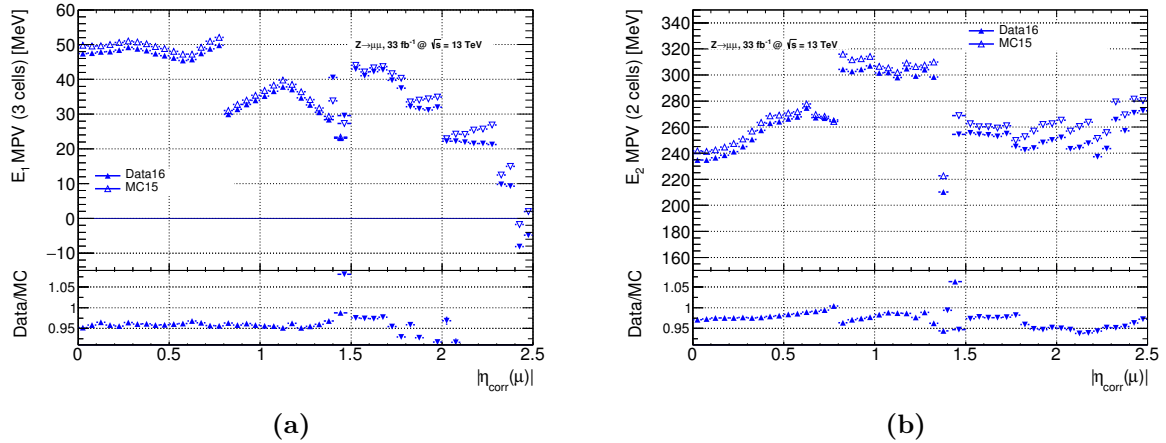


Figure C.11 – MPV extraite par la méthode Landau * Gaussian pour chaque région de $|\eta|$ dans L1 (a) et L2 (b).

Premièrement, l'analyse effectuée au Run 1 a révélé n'utiliser qu'une seule cellule pour l'énergie dans L2 (contrairement à deux dans l'analyse de 2016). N'utiliser qu'une seule cellule pour l'analyse de 2016 s'est conclu par un simple décalage en énergie entre 1 % et 2 % mais sans changer le comportement des valeurs. Cette raison a donc été exclue comme cause possible de la différence. Deuxièmement, des effets liés à la structure discrète des faisceaux au LHC et à leur impact sur le bruit dans le calorimètre ont été envisagés mais se sont également révélés insuffisants pour expliquer les différences observées. Troisièmement, l'isolation des muons a été étudiée : la sélection standard n'applique pas de coupure sur leur isolation. Ainsi que mentionné plus haut, le signal des MIP est du même ordre de grandeur que le bruit dans le calorimètre, la mesure de leur énergie est donc particulièrement sensible au bruit ambiant. Utiliser des muons non-isolés

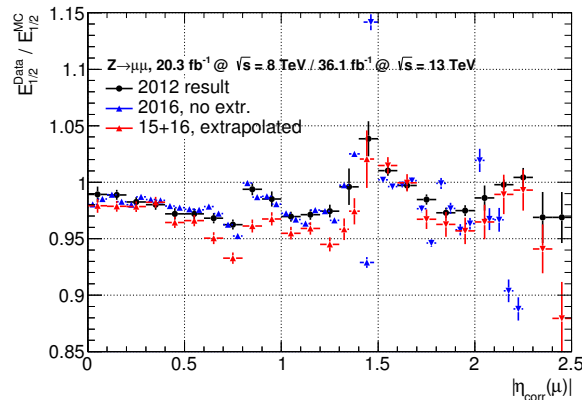


Figure C.12 — Comparaison des facteurs d'étalonnage relatifs en fonction de $|\eta|$ dérivés au Run 1 (en noir, combinaison des méthodes MPV et TM) et dérivés de manière similaire dans les données 2016 avec la méthode MPV (en bleu). La courbe rouge correspond au résultat de la méthode améliorée utilisant les données 2015 et 2016, et décrite à la Section C.3.3.

énergétiquement (comme par exemple dans un jet) pourrait donc biaiser cette mesure, et une coupure stricte sur leur isolation a été ajoutée sans changer significativement le résultat obtenu. Quatrièmement, une méthode a été développée pour estimer le bruit environnant les cellules du calorimètre proche du passage des muons afin d'en soustraire l'effet sur les cellules touchées par le muon. Cette étude n'a pas non plus permis d'expliquer les différences observées.

Les quatre pistes précédentes visaient à réduire le bruit dans le calorimètre. Or le bruit est fortement lié à l'empilement (voir Section C.2.1), particulièrement dans les bouchons. Comme l'empilement moyen est plus élevé dans les données de 2016 qu'au Run 1, son effet direct sur les valeurs d'étalonnage a également été étudié. En répétant et comparant l'étude en trois intervalles d'empilement ($\mu \in [0; 19], [19; 25], [25; 40]$), on observe que les MPV extraites dans chacun des intervalles diffèrent de 10 % dans le tonneau pour L1 et jusqu'à 50 % dans les bouchons. La variation est plus modérée (2 %) dans L2. Après avoir calculé les doubles ratios, on observe que l'analyse effectuée dans l'intervalle $[0; 19]$ est la plus proche du Run 1 alors que l'analyse effectuée dans l'intervalle $[25; 40]$ en est la plus éloignée. Ceci indique une dépendance des facteurs d'étalonnage à l'empilement. Cependant, l'étalonnage relatif L1/L2 ne doit corriger que des effets d'électronique ou de diaphonie, qui doivent être indépendants de l'empilement. Le développement d'une telle méthode a été un des points centraux du travail réalisé lors de cette Thèse, et est décrit à la Section C.3.3

C.3.3 Étalonnage relatif L1/L2 : méthode d'extrapolation et patrons de bruit

Les études précédemment réalisées n'ont pas été suffisantes pour expliquer la différence observée entre les résultats du Run 1 et de 2016. Cependant, toutes ont montré une dépendance au bruit dans le calorimètre. La première étape dans la résolution du problème a été de vérifier l'hypothèse gaussienne utilisée dans l'ajustement Landau * Gaussian pour modéliser le bruit. Les distributions d'énergie recueillies dans les cellules du calorimètre en l'absence de collision (ne contenant donc que du bruit) ont été construites pour chaque région en $|\eta|$, séparément pour L1 et L2 et dans les données et la simulation. Un exemple est présenté dans la Figure C.13a pour les données dans L2, ainsi que le meilleur ajustement d'une gaussienne à cette distribution. Cet ajustement montre que les distributions de bruit ne sont pas centrées en 0, et ne sont pas gaussiennes (les queues de distribution ne sont pas symétriques). Le modèle de bruit utilisé précédemment a donc été remplacé par ces patrons de distributions de bruit, améliorant la modélisation de ce dernier.

Comme vu précédemment, les facteurs d'étalonnage relatif L1/L2 ne doivent pas dépendre de l'empilement dans les collisions. L'idéal serait donc de pouvoir extraire ces facteurs dans un environnement sans empilement afin d'en éliminer tout effet. Puisqu'aucune donnée n'a été enregistrée sans ou à faible empilement, une détermination directe des facteurs dans ces conditions est impossible. En revanche, une grande statistique a été accumulée pour des valeurs d'empilement entre 15 et 35. La méthode développée est la suivante : la détermination des MPV est faite indépendamment dans 13 régions d'empilement ($[0; 10; 12; 14; 16; 18; 20; 22; 24; 26; 28; 30; 34; 44]$), et pour chacune des régions en $|\eta|$. Pour chaque région en $|\eta|$, les MPV en fonction de l'empilement sont ajustées par une fonction linéaire, permettant une extrapolation des MPV à une valeur d'empilement nulle. Un exemple est montré à la Figure C.13b. Le coefficient d'extrapolation (pente de la droite) diffère entre L1 et L2 et entre données et simulation mais reste faible dans tous les cas.

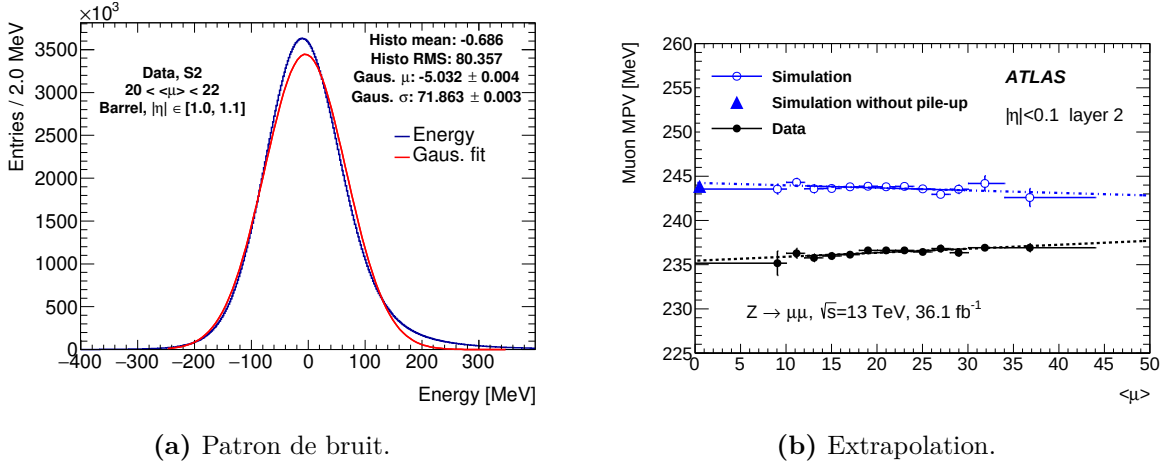


Figure C.13 – (a) Exemple de distribution de bruit dans les données pour la région $|\eta_{\text{corr}}| \in [1.00, 1.10]$ dans L2 et l'intervalle $20 < \langle \mu \rangle < 22$, et meilleur ajustement d'une gaussienne. (b) Exemple de distribution de MPV en fonction de l'empilement $\langle \mu \rangle$ pour la région $|\eta_{\text{corr}}| \in [1.70, 1.80]$ dans L2 et ajustement linéaire [141]. Les données sont en noir et la simulation en bleu. Les lignes pleines montrent l'intervalle d'ajustement et les lignes pointillées l'extrapolation. Les triangles bleus sont les valeurs extraites dans une simulation sans empilement ($\langle \mu \rangle = 0$).

Les MPV extrapolées en fonction de $|\eta|$ sont illustrées dans la Figure C.14 et montrent un simple décalage en énergie dans le tonneau pour L1 ainsi que dans tout L2. Le double ratio $\alpha_{1/2}$ est finalement calculé comme précédemment mais en utilisant les MPV extrapolées. Les résultats sont présentés dans la Figure C.12, ainsi qu'une comparaison aux valeurs sans utiliser l'extrapolation. La méthode avec extrapolation donne des résultats beaucoup plus proches des valeurs attendues dans les bouchons, confirmant ainsi la dépendance aux effets d'empilement dans les collisions.

C.3.4 Étalonnage relatif L1/L2 : résultats finals et systématiques

La dernière étape est d'estimer les erreurs systématiques pouvant affecter cette mesure. La première source d'incertitude se trouve dans la méthode d'extrapolation elle-même. Afin de valider la méthode, une simulation sans empilement a été générée et les MPV ont été extraites avec la méthode standard (sans extrapolation). La comparaison entre ces valeurs et les MPV après extrapolation sont en bon accord, validant ainsi la méthode. Les différences résiduelles dans L1 et L2 sont utilisées comme systématiques sur le résultat final. Cette validation n'a pu être effectuée dans les vraies données faute d'événement sans empilement. Cette étude a cependant contribué à la justification d'une demande de prise de données à faible empilement en 2017 et 2018.

Deuxièmement, la seconde cellule utilisée pour calculer l'énergie dans L2 a été choisie comme la plus énergétique des cellules voisines dans la direction η , ayant un meilleur rapport signal sur

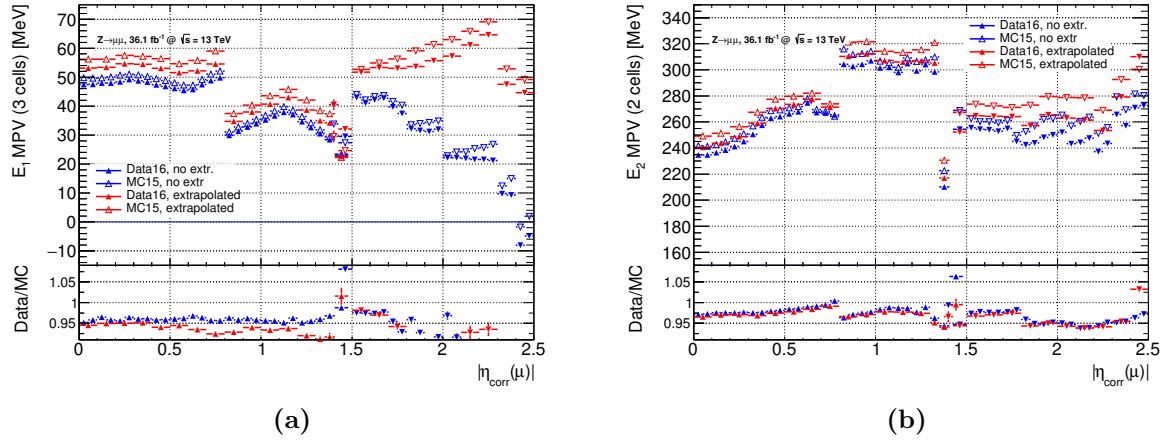


Figure C.14 – Comparaison des MPV dans (a) L1 et (b) L2, avec (rouge) ou sans (bleu) appliquer la méthode d'extrapolation. La partie basse montre le ratio données/simulation avec et sans extrapolation.

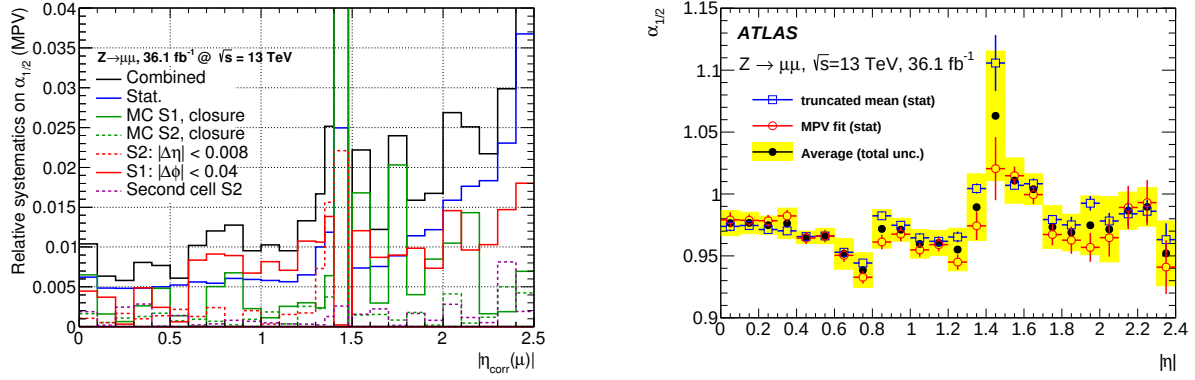
bruit. Cependant un autre choix raisonnable est de prendre la seconde cellule la plus proche de l'extrapolation de la trace du muon, toujours dans la direction ϕ . Ce changement affecte environ la moitié des muons ainsi que l'énergie reconstruite pour les muons (2 % à 4 %), mais cet effet est similaire dans les données et la simulation. L'effet sur le double ratio $\alpha_{1/2}$ est ainsi négligeable, de l'ordre de quelques pour-mille au maximum.

Troisièmement, les muons touchant les extrêmes en ϕ des cellules de L1 partagent leur énergie entre deux cellules en ϕ , comme pour les cellules de L2. Cependant, ajouter une autre cellule dans la direction ϕ dans L1 rajouterait beaucoup de bruit en comparaison du signal récupéré. L'estimation de cet effet a été fait en restreignant l'analyse à des muons passant au milieu des cellules dans L1 de manière à éliminer les muons partageant leur énergie. L'effet associé est d'environ 8 % et est similaire dans les données et la simulation, sauf dans la région $|\eta| \in [0,8; 1,4]$. En conséquence l'impact sur le double ratio $\alpha_{1/2}$ est de l'ordre de 1 % pour $0,8 < |\eta| < 1,4$, et de l'ordre de 0,5 % en dehors.

Quatrièmement, à cause de la grande profondeur des cellules dans L2, des effets liés à l'alignement du calorimètre peuvent être perceptibles, en particulier dans le tonneau. En effet, si celui-ci n'est pas à sa position nominale, les cellules ne sont pas parfaitement projectives par rapport au point d'interaction, et les muons peuvent déposer leur énergie dans plusieurs cellules en η . Cet effet a été estimé en contraignant les muons à passer dans le centre de la cellule dans S2. La différence est d'environ 2 % sur la MPV mais est à nouveau similaire entre les données et la simulation. L'effet sur le double ratio $\alpha_{1/2}$ est de l'ordre du pour-mille dans le tonneau, et du pour-cent dans les bouchons.

L'impact individuel ainsi que la combinaison de ces effets est présenté dans la Figure C.15a. Enfin, le résultat final des facteurs d'étalonnage relatif L1/L2 est la moyenne des valeurs extraites par la méthode MPV et par la méthode TM. Ces résultats sont montrés dans la Figure C.15b. L'erreur finale sur le résultat est la somme quadratique de la différence entre les deux méthodes (MPV et TM) et des erreurs individuelles de chacune des deux méthodes.

Conclusion : la dépendance à l'empilement des facteurs d'étalonnage relatif des couches 1 et 2 du calorimètre électromagnétique a été démontrée. Une nouvelle méthode améliorant la modélisation du bruit dans l'analyse et introduisant une extrapolation des résultats en fonction de l'empilement a été implémentée. Cette méthode a été validée à l'aide de simulations dédiées et les résultats extraits des données 2015 et 2016 donnent des facteurs d'étalonnage et des incertitudes comparables à l'analyse effectuée avec les données du Run 1.



(a) Résumé des systématiques de la méthode MPV.

(b) Résultat final d'étalonnage relatif L1/L2.

Figure C.15 – (a) Résumé des systématiques sur les facteurs d'étalonnage relatif L1/L2 pour la méthode MPV, en fonction de $|\eta|$. La valeur combinée est la somme quadratique toutes les autres contributions. L'incertitude statistique est calculée comme la somme quadratique des erreurs sur les MPV extrapolées venant des ajustements de Landau*Patron. (b) Résultat final des facteurs d'étalonnage relatif L1/L2 pour les méthodes MPV (rouge) et TM (bleu) ainsi que leur moyenne (noir). Les barres d'erreur sur les points représentent uniquement l'erreur statistique. La bande jaune représente l'incertitude globale du résultat, moyenne quadratique de la différence de résultat entre les deux méthode et de leurs erreurs individuelles. [141]

C.4 Étude de couplages CP -impairs entre le boson de Higgs et les bosons faibles dans le canal $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$

C.4.1 L'analyse $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$

A cause de son faible rapport d'embranchement ($1,24 \times 10^{-4}$), la précision des analyses dans le canal de désintégration $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ est limité statistiquement. Cependant, son état final est spécifique, éliminant dans une grande mesure les bruits de fond et donnant un rapport signal sur bruit d'environ 2 : 1 dans la région de signal ($115 \text{ GeV} < m_{4\ell} < 130 \text{ GeV}$). Afin de garder une haute précision dans les résultats, l'analyse limite les leptons aux électrons (e) et muons (μ), bien mesurés. L'état final ne possède également pas de jets (sauf pour marquer certains modes de productions spécifiques) et n'est donc pas sujet aux plus larges incertitudes de mesures les entâchant.

Le bruit de fond principal est constitué du processus non-résonnant $q\bar{q} \rightarrow ZZ^{(*)}$, donnant le même état final que $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$. Il est estimé à partir des simulations, et sa normalisation est déduite des données en dehors de la région de signal. D'autres contributions moins importantes sont liées à une mauvaise identification des électrons (conversion de photons, hadrons légers au dépôt semblable dans le calorimètre, ou désintégration d'un hadron lourd) ou des muons (désintégration des hadrons légers ou lourds, ou chaîne de désintégration d'un quark top). Ces contributions sont évaluées séparément pour les électrons et les muons dans des régions de contrôle dédiées et orthogonales à la région de signal. Une description détaillée est faite en Réf. [246].

Les incertitudes systématiques affectant les mesures concernent d'une part les électrons, muons, jets et l'empilement de collisions (efficacités de déclenchement, reconstruction, identification et isolation, et incertitudes d'étalonnage en énergie des objets mesurés) [130, 134, 141, 252–254] à hauteur de 2 % à 7 % en fonction des états finals, et d'autre part les simulations des processus physiques ou du détecteur (distributions partoniques, énergie de factorisation et renormalisation et couplage de l'interaction forte, simulation des gerbes partoniques) à hauteur de 3 % à

15 %. L'incertitude sur la luminosité, évaluée à 1,7 % pour l'ensemble des 139 fb^{-1} du Run 2 [124], est également prise en compte. Enfin, des incertitudes sont associées à l'estimation des bruits de fond mentionnés précédemment et sont propagées dans les analyses. Une description complète des incertitudes systématiques est développée en Réf. [245]. La distribution de masse invariante des différentes composantes est donné à la Figure C.16a. Dans l'intervalle $115 \text{ GeV} < m_{4\ell} < 130 \text{ GeV}$, 316 événements ont été observés pour 315 ± 14 attendus, dont 206 désintégrations de bosons de Higgs et 97 venant des processus non résonnants $qq \rightarrow ZZ^{(*)}$ ou $gg \rightarrow ZZ^{(*)}$.

Comme mentionné à la Section C.1.3, la recherche de couplages anormaux entre le boson de Higgs et les bosons faibles est la plus prometteuse dans l'étude du mode de production VBF, et ce en particulier dans le canal $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$. Cependant, comme décrit à la Section C.1.2, la sélection des événements avec deux jets ne suffit pas à obtenir une catégorie d'analyse pure en événements VBF. Ainsi, des 50 événements attendus dans la catégorie à 2 jets, 38 sont issus d'un boson de Higgs, et seuls 9 d'entre eux découlent d'une production VBF. Pour purifier cette sélection, un *réseau de neurones artificiels* (ou NN : *neural network*) est entraîné à séparer les événements Higgs des bruits de fond $qq \rightarrow ZZ^{(*)}$, et un autre à séparer les événements VBF des événements ggF avec deux jets additionnels et des événements VH -had. L'entraînement de ces algorithmes se base sur les variables cinématiques des leptons et des jets, ainsi que des propriétés du système des deux jets. La distribution des résultats de ce dernier, pouvant être assimilés à la probabilité qu'un événement donné soit réellement VBF, est montrée à la Figure C.16b.

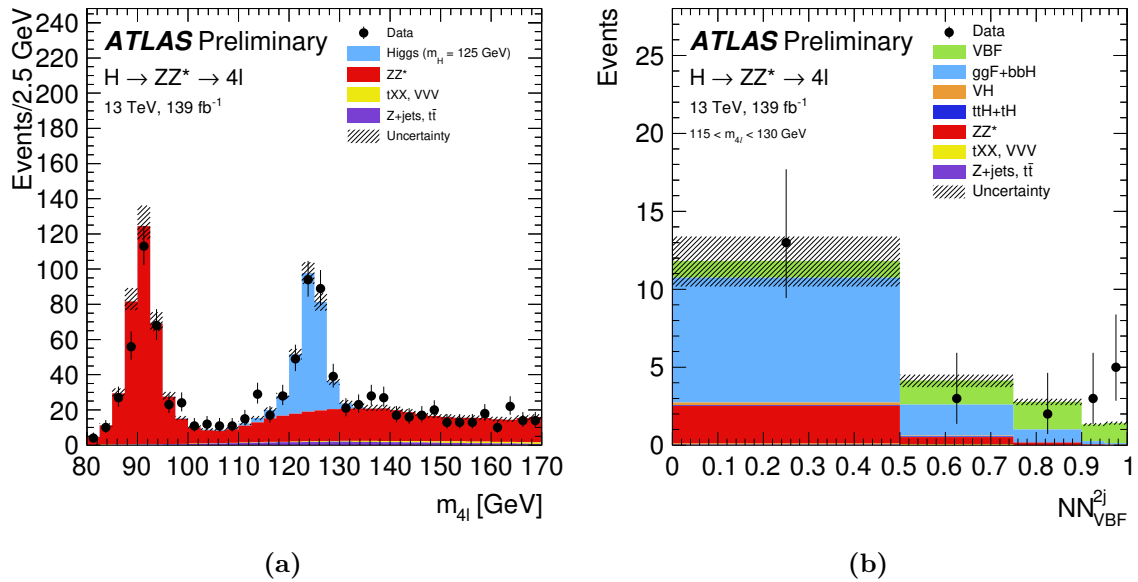


Figure C.16 — Distribution de masse invariante du système de quatre leptons (a) et distribution de résultat du réseau de neurone pour VBF (en pré-sélectionnant les événements avec deux jets) dans l'analyse $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ d'ATLAS au Run 2. [245]

C.4.2 Observables optimales

La présence de couplages anormaux entre le boson de Higgs et les bosons faibles modifie les sections efficaces de production et les rapports d'embranchement du boson de Higgs. La précédente recherche de couplages anormaux effectuée au sein de l'analyse $H \rightarrow 4\ell$ [178] utilisait ces quantités pour déduire des limites sur les couplages. Or l'effet de couplages pairs et impairs par CP modifie de manière similaire les sections efficaces, et une analyse ne reposant que sur la mesure de sections efficaces ne peut distinguer la présence de couplages pairs ou impairs. Pour séparer ces effets, des observables dédiées doivent être utilisées. Contrairement aux couplages pairs, les couplages impairs modifient les distributions angulaires entre les différents objets de l'état final (leptons

et jets). Les observables basées sur les distributions angulaires permettent ainsi une analyse minimisant l'impact de possibles effets pairs par CP , réduisant les hypothèses faites dans l'analyse.

Cependant, l'utilisation d'une variable telle que la séparation angulaire azimutale entre les jets ($\Delta\phi_{jj}$, Figure C.17a) ne prend en compte qu'une partie de l'information cinématique de l'état final. Une observable basée sur les éléments de matrice de l'évènement, nommée *observable optimale* et contenant la totalité de cette information cinématique, donne ainsi une meilleure sensibilité aux effets recherchés. Cette observable optimale (OO) est définie par :

$$\mathcal{OO}_1(\mathbf{c}) = \frac{\text{Terme d'interférence}}{|\mathcal{M}_{\text{SM}}|^2} = \frac{2\Re(\mathcal{M}_{\text{SM}}\mathcal{M}_{\text{BSM}}^*(\mathbf{c}))}{|\mathcal{M}_{\text{SM}}|^2} = \frac{|\mathcal{M}_{\text{Mix}}(\mathbf{c})|^2 - |\mathcal{M}_{\text{SM}}|^2 - |\mathcal{M}_{\text{BSM}}(\mathbf{c})|^2}{|\mathcal{M}_{\text{SM}}|^2}, \quad (\text{C.4})$$

où $\mathcal{M}_{\text{SM}}$, $\mathcal{M}_{\text{BSM}}$ and $\mathcal{M}_{\text{Mix}}$ sont les éléments de matrices calculés selon une hypothèse $\mathbf{c}$ de couplages du Modèle Standard, de couplages purement impairs, ou d'un mixte de couplages du Modèle Standard et impairs, respectivement. Seuls les modules des éléments de matrices ($|\mathcal{M}|^2$) sont des observables physiques, et sont ici calculés avec le programme MADGRAPH5_aMC@NLO au premier ordre. J'ai montré que la valeur du couplage (par exemple $\mathbf{c} = \{\tilde{c}_{zz} = 1\}$ ou $\mathbf{c} = \{\tilde{c}_{zz} = 2\}$) ne change pas la sensibilité de l'analyse (cela revient à une homothétie de la distribution de l'observable). La combinaison relative de $\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$ et $\tilde{c}_{\gamma\gamma}$ (c'est-à-dire sa direction), en revanche, importe ($\mathbf{c} = \{\tilde{c}_{zz} = 1\}$ et $\mathbf{c} = \{\tilde{c}_{z\gamma} = 1\}$ donneront des résultats différents).

Une telle observable a la propriété d'être symétrique pour les couplages du Modèle Standard ou pour des couplages pairs par CP , mais devient asymétrique en présence de couplages impairs, comme l'atteste la Figure C.17b (distributions paires par CP en vert et orange, et impaires par CP en rouge et bleu, à comparer au Modèle Standard en noir). Étant donnée une hypothèse de couplage, l'observable est calculée à partir des variables cinématiques des différents objets dans l'état final. Une observable peut être définie pour la partie production ($pp \rightarrow H + \text{jets}$) ou bien pour la partie désintégration ($H \rightarrow 4\ell$). L'analyse décrite ici étant portée sur la production VBF, l'observable optimale de production (notée $\mathcal{OO}_{1,jj}$) sera utilisée.

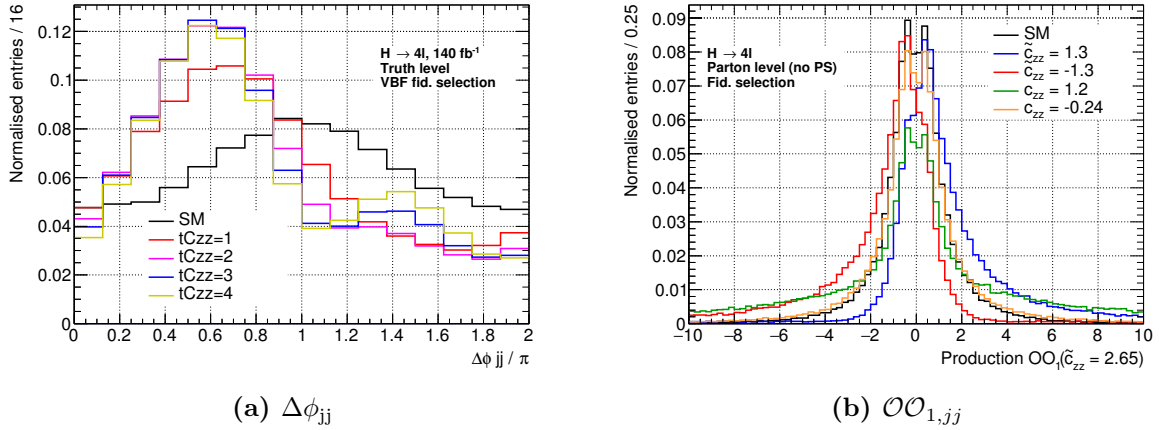


Figure C.17 – Distributions d'observables sensibles aux effets de CP : (a) $\Delta\phi_{jj}$ pour le Modèle Standard et plusieurs valeurs de $\tilde{c}_{zz}$, (b) $\mathcal{OO}_{1,jj}$ pour le Modèle Standard, et des couplages pairs (c_{zz}) ou impairs $\tilde{c}_{zz}$. Toutes les distributions sont normalisées à 1.

Les protons étant des particules composites, un élément de matrice doit être calculé pour chaque paire de partons donnant lieu à la collision. Comme les calculs d'éléments de matrices sont coûteux en temps, certaines contributions mineures au résultat total ont été négligées, et les contributions donnant des résultats similaires ont été fusionnées. Ces optimisations ont permis un gain d'un facteur 7 à 8 sur le temps de calcul tout en ayant un effet négligeable sur les observables optimales.

Les contributions individuelles doivent ensuite être combinées à l'aide des distributions de parton dans le proton (« *parton distribution function* » ou *PDF*) pour donner le résultat final. Différents calculs de PDF sont disponibles et leurs résultats sont paramétrés par le transfert d'impulsion Q^2 de la collision. Bien que faire varier la PDF et la valeur de Q^2 choisie ait un impact sur le calcul des éléments de matrice, ces différences se compensent lors du calcul de l'observable optimale (ratio d'éléments de matrice).

C.4.3 Étude des couplages impairs par CP du boson de Higgs

Une première étude préliminaire a été effectuée pour déterminer la gamme de couplages à laquelle l'analyse serait sensible. Environ 250 échantillons simulant le processus VBF avec des couplages impairs par CP ont été produits, couvrant une vaste gamme de valeurs de couplages sur les axes $\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$ et $\tilde{c}_{\gamma\gamma}$ dans la base Higgs ainsi que les axes $c_{H\tilde{W}}$, $c_{H\tilde{W}B}$ et $c_{H\tilde{B}}$ dans la base Warsaw. Chaque échantillon a été généré avec 200k évènements, mais sans inclure la simulation du détecteur (coûteuse en calcul). La distribution $\mathcal{OO}_{1,jj}$ pour chaque échantillon est comparée à celle donnée par un échantillon VBF « Modèle Standard », donnant une valeur de vraisemblance (ratio de *log-likelihood*) qui peut être utilisée pour calculer la sensibilité de l'analyse.

La vraisemblance en fonction de la valeur de couplage pour $\tilde{c}_{zz}$ est montrée à la Figure C.18a, formant un *scan de vraisemblance*. Plus le scan est resserré, plus les contraintes posées sur le paramètre d'intérêt sont fortes. La Figure C.18a montre que le scan sur $\tilde{c}_{zz}$ est similaire pour une observable construite avec $\mathbf{c} = \{\tilde{c}_{zz} = 1\}$ ou $\{\tilde{c}_{z\gamma} = 1\}$ ou $\{\tilde{c}_{\gamma\gamma} = 1\}$ ou $\{c_{H\tilde{W}} = 1\}$. Dans la suite, l'observable construite avec l'hypothèse $\mathbf{c} = \{\tilde{c}_{zz} = 1\}$ sera choisie. La plage de couplage pour laquelle le scan est en dessous de 1 (resp. 4) donne l'intervalle de confiance à 68 % (resp. 95 %) des valeurs probables pour ce paramètre. Pour $\tilde{c}_{zz}$, l'intervalle à 68 % de confiance est $[-0.38, 0.38]$.

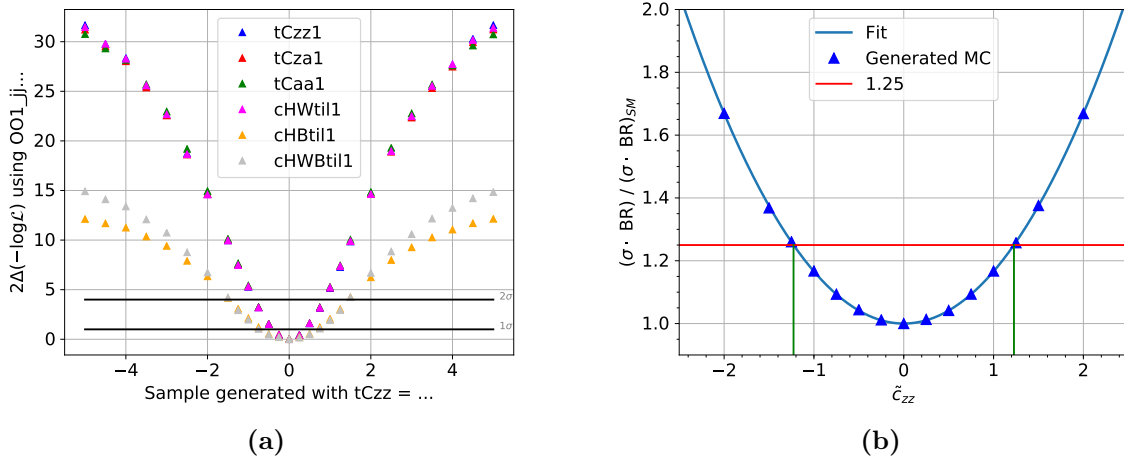


Figure C.18 – (a) Scan de vraisemblance pour le coefficient $\tilde{c}_{zz}$ en utilisant $\mathcal{OO}_{1,jj}$ (construite suivant différentes hypothèses), avec un échantillon de données simulées de 140 fb^{-1} pur en évènements VBF et généré sans effets du détecteur. L'échantillon de référence (asimov) utilisé est le Modèle Standard. (b) Section efficace VBF multipliée par de rapport d'embranchement $H \rightarrow 4\ell$ normalisée par la valeur donnée par le Modèle Standard en fonction de $\tilde{c}_{zz}$. L'ajustement à l'aide d'une fonction quartique $x \rightarrow ax^4 + bx^2 + 1$ est effectué sur plus de points qu'affichés. Les valeurs de couplage donnant une valeur 25 % inférieure au Modèle Standard sont listées en Table C.1.

Un précédent résultat n'utilisant que les informations de section efficace [208] montre que la section efficace de production par le mode VBF est connue à environ 25 %. L'analyse présentée ici ne sera donc compétitive que si elle peut placer des limites sur les valeurs de couplages qui ne sont pas déjà exclues par cette analyse. En utilisant la section efficace calculée par MADGRAPH5_aMC@NLO pour chacun des échantillons produits, la section efficace normalisée par

la section efficace VBF du Modèle Standard est tracée en fonction de la valeur de couplage (Figure C.18b). Toutes les valeurs de couplages donnant un ratio $(\sigma \times \text{BR})/(\sigma \times \text{BR})_{\text{SM}}$ plus grand que 1.25 sont déjà exclues au niveau de confiance 68 %. Les limites par l'analyse utilisant la comparaison des formes de distribution de $\mathcal{OO}_{1,jj}$ pour des échantillons VBF purs, et par l'analyse utilisant les sections efficaces uniquement pour chacun des paramètres sondés sont regroupées dans la Table C.1. L'amélioration possible sur le coefficient $\tilde{c}_{zz}$ se révèle prometteuse (± 0.38 contre ± 1.24), et la sensibilité aux coefficients $\tilde{c}_{z\gamma}$ et $c_{H\tilde{W}}$ mérite d'être étudiée. La sensibilité aux coefficients $\tilde{c}_{\gamma\gamma}$, $c_{H\tilde{B}}$ et $c_{H\tilde{W}B}$ est un ordre de grandeur plus faible que celle de l'analyse des sections efficaces. La suite de cette étude se concentre donc sur le coefficient $\tilde{c}_{zz}$.

Table C.1 – Comparaison des limites sur différents coefficients de Wilson impairs par CP ($\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$, $\tilde{c}_{\gamma\gamma}$, $c_{H\tilde{W}}$, $c_{H\tilde{B}}$, $c_{H\tilde{W}B}$) obtenus par un scan de vraisemblance avec un échantillon VBF pur, et des valeurs de couplage donnant une section efficace multipliée VBF par le rapport d'embranchement $H \rightarrow 4\ell$ 25 % supérieure à la valeur du Modèle Standard.

Coefficient de Wilson	Base Higgs			Base Warsaw		
	$\tilde{c}_{zz}$	$\tilde{c}_{z\gamma}$	$\tilde{c}_{\gamma\gamma}$	$c_{H\tilde{W}}$	$c_{H\tilde{W}B}$	$c_{H\tilde{B}}$
Limite à 68 % CL utilisant $\mathcal{OO}_{1,jj}$ (VBF uniquement)	± 0.38	± 0.85	± 6.2	± 0.70	± 5.2	± 5.4
Limite à 68 % CL d'après les sections efficaces	± 1.24	± 0.39	± 0.94	± 0.73	± 0.73	± 0.40

Cette analyse préliminaire se base sur un nombre élevé d'événements simulés (environ 50 millions). Ceci est possible tant que la simulation complète du détecteur n'est pas utilisée. L'analyse finale doit au contraire reposer sur des échantillons incluant la simulation complète afin de calculer précisément les erreurs systématiques affectant le résultat. Comme seulement un nombre restreint d'échantillons avec une simulation complète peuvent être produits, et donc un nombre limité de valeurs de couplage générés, une méthode d'interpolation entre les différents échantillons doit être utilisée. La méthode utilisée pour cette analyse est la technique de morphing [267, 268]. Dans le cas présent (interpolation d'un unique paramètre en plus du Modèle Standard), un minimum de cinq échantillons sont nécessaires, formant une base de morphing. Cette méthode d'interpolation a été validée pour cette analyse et a montré que le résultat est valide entre les points extrémaux de la base, voire légèrement en dehors.

Une fois cette étape préliminaire de sensibilité effectuée, l'impact des bruits de fond sur l'analyse a été étudiée. Le bruit de fond de cette analyse est principalement constitué de bosons de Higgs produits par le mode ggF avec deux jets additionnels, et d'événements $qq \rightarrow ZZ^{(*)}$ avec deux jets, présentant le même état final que les événements VBF constituant le signal. Cette étude est également réalisée avec des échantillons n'incluant pas les effets du détecteur. Pour avoir un résultat plus proche des conditions réelles (incluant les effets du détecteur), une efficacité de sélection des événements de 45 % est incluse dans l'analyse. Le scan de vraisemblance de cette étude pour le paramètre $\tilde{c}_{zz}$ est comparé au scan de l'étude précédente (sans bruit de fond et avec efficacité de 100 %) dans la Figure C.19a (courbes rouge et noire respectivement). Une forte baisse de la sensibilité à 68 % de confiance est observée, passant de $[-0.38, 0.37]$ à $[-1.3, 1.3]$ et confirmant le large impact des bruits de fond.

Afin d'augmenter la sensibilité de l'analyse, une meilleure séparation entre le signal et les bruits de fond est nécessaire. Dans ce but, un réseau de neurones artificiels, « NN » (ou des arbres de décision boostés, « BDT », dans les premières versions de l'analyse) est employé (Figure C.16b). L'analyse est donc effectuée simultanément dans cinq régions de score NN ou BDT, permettant d'avoir des catégories très pures en événements VBF et donc plus sensibles que l'analyse inclusive. Une amélioration significative de la sensibilité à 68 % CL de confiance est observée, passant à $[-0.81, 0.82]$, comme montré par la courbe bleue de la Figure C.19a. Il a été vérifié que l'utilisation de tels NN n'introduisait pas de biais significatif dans les scans, même en présence de couplages impairs.

Étant donnés ces résultats précédents, des échantillons incluant une simulation complète du

Table C.2 – Résultats de sensibilité aux effets des coefficients de Wilson $\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$, $c_{H\tilde{W}}$ dérivés des scans de vraisemblance utilisant $\mathcal{OO}_{1,jj}$, pour les niveaux de confiance 68 % et 95 %. Les valeurs comparent les résultats des analyses sans bruits de fond, avec bruit de fond (ggF, VBF et $qq \rightarrow ZZ^{(*)}$) et incluant un facteur d'efficacité de reconstruction de 45 %, et comme précédemment mais en séparant l'analyse en 5 régions de score de BDT. Dans tous les cas, le résultat est donné pour une luminosité de 140 fb^{-1} . « – » indique que le niveau de confiance n'est pas atteint pour le scan.

Coefficient de Wilson	Base Higgs				Base Warsaw	
	$\tilde{c}_{zz}$		$\tilde{c}_{z\gamma}$		$c_{H\tilde{W}}$	
Intervalle de confiance	68 % CL	95 % CL	68 % CL	95 % CL	68 % CL	95 % CL
Sans b.d.f., $\epsilon_{\text{Reco}} = 100 \%$	± 0.4	± 0.8	± 0.8	± 1.8	± 0.7	± 1.5
Avec b.d.f., $\epsilon_{\text{Reco}} = 45 \%$	± 1.3	± 4.7	± 2.6	–	± 2.2	± 9.0
Avec b.d.f., $\epsilon_{\text{Reco}} = 45 \%$, 5 régions de BDT	± 0.8	± 2.2	± 1.8	± 5.2	± 1.4	± 4.0
Avec simulation détecteur, 5 régions de BDT	± 0.8	± 2.3	Non étudié		Non étudié	

détecteur sont produits pour les valeurs $\tilde{c}_{zz} = -2, -1, 0, 1, 2$, et servent de base de morphing pour la suite. Afin de vérifier l'impact des effets du détecteur sur l'analyse, le scan de vraisemblance utilisant ces nouveaux échantillons est comparé au précédent résultat. Comme en atteste la Figure C.19b, les résultats sont proches aux niveaux de confiance 68 % ($[-0.81, 0.82]$ contre $[-0.82, 0.79]$) et 95 % ($[-2.1, 2.2]$ contre $[-2.3, 2.4]$), indiquant que les effets du détecteur ont un impact limité sur l'analyse. Tous ces résultats sont résumés à la Table C.2.

Finalement, la composante systématique des incertitudes est évaluée grâce à ces échantillons incluant la simulation complète du détecteur. Les différentes sources prises en compte sont celles listées à la Section C.4.1, et leurs impacts individuels sont sommés quadratiquement. L'incertitude systématique totale affectant les intervalles de confiance est estimée à moins de 1 %.

Conclusion : la sensibilité finale attendue de cette analyse pour le paramètre $\tilde{c}_{zz}$ est ainsi

$$\tilde{c}_{zz} = 0.0 \pm 0.8 \text{ (stat)} \pm 0.01 \text{ (syst)}. \quad (\text{C.5})$$

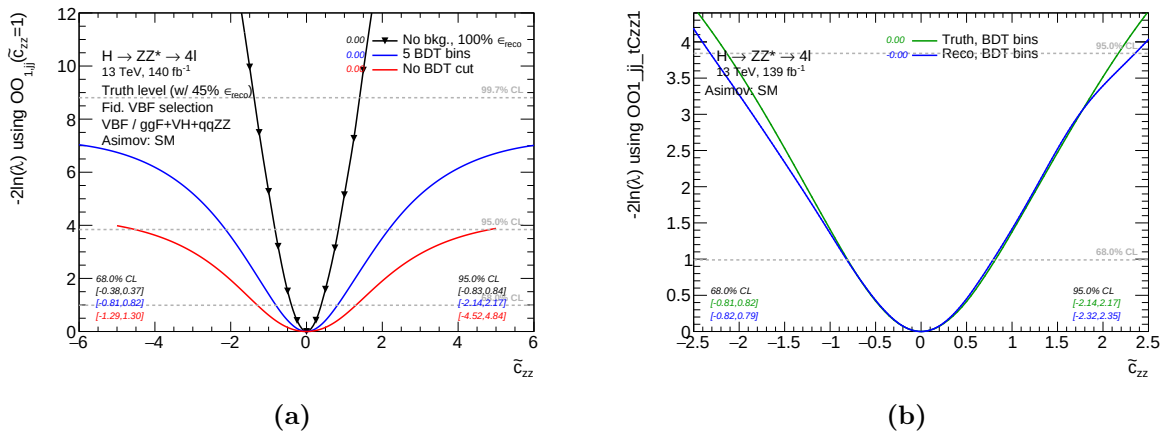


Figure C.19 – Scan de vraisemblance du paramètre $\tilde{c}_{zz}$ en utilisant l'observable $\mathcal{OO}_{1,jj}$ contre l'hypothèse du Modèle Standard (asimov), (a) avec des échantillons sans simulation du détecteur et sans bruit de fond (noir), avec bruits de fond ggF, VH et $qq \rightarrow ZZ^{(*)}$ et incorporant un facteur de 45 % d'efficacité de reconstruction (rouge), comme précédemment mais en séparant l'analyse en cinq régions délimitées par le score du BDT (bleu); (b) avec (bleu) ou sans (vert) effets du détecteurs mais incluant 45 % d'efficacité de reconstruction, dans les deux cas en considérant les bruits de fond et en séparant l'analyse en cinq régions délimitées par le score de BDT ou de NN. Une luminosité de 140 fb^{-1} (sans) ou 139 fb^{-1} (avec effets du détecteur) est utilisée.

List of Abbreviations

ATLAS	A Toroidal LHC ApparatuS
BCID	Bunch Crossing IDentifier
BDT	Boosted Decision Tree
BSM	Beyond the Standard Model
CERN	European Organisation for Nuclear Research
CL	Confidence Level
CMS	Compact Muon Solenoid
CP, CP	Charge-Parity (symmetry)
CR	Control Region
DIS	Deep Inelastic Scattering
DNN	Deep Neural Network
ECal, EMCal	Electromagnetic Calorimeter
EFT	Effective Field Theory
E_T	Transverse energy
EW	ElectroWeak
FSR	Final-State Radiation
ggF	gluon–gluon fusion
HB	Higgs Basis
HC	Higgs Characterisation
HCal	Hadronic Calorimeter
HF	Heavy Flavour
HLT	High-Level Trigger
IBL	Insertable B-Layer
ID	Inner Detector
ISR	Initial-State Radiation
LAr	Liquid Argon
LEP	Large Electron–Positron collider
LF	Light Flavour
LHC	Large Hadron Collider
LO	Leading Order
MC	Monte-Carlo
ME	Matrix Element
MET	Missing transverse energy
MIP	Minimum Ionising Particle
MLP	Multi-Layer Perceptron
MPV	Most Probable Value
MS	Muon Spectrometer
NLO	Next to Leading Order

NN	Neural Network
NP	New Physics
	Nuisance Parameter
OFC	Optimal Filtering Coefficient
OO	Optimal Observable
PDF	Parton Distribution Function
PS	Pre-Sampler
	Parton Shower
p_T	Transverse momentum
QED	Quantum ElectroDynamics
QCD	Quantum ChromoDynamics
RF	Radio-Frequency
RNN	Recurrent Neural Network
SCT	Strip semiConductor Tracker
SM	Standard Model
SMEFT	Standard Model Effective Field Theory
SNR	Signal-to-Noise Ratio
SR	Signal Region
TM	Truncated Mean
TRT	Transition Radiation Tracker
VBF	Vector-Boson Fusion
VR	Validation Region
WB	Warsaw Basis

List of Figures

1	The Standard Model of particle physics	15
1.1	Example trajectories before and after a global or a local space translation.	18
1.2	Feynman diagrams associated with the Dirac Lagrangian under $U(1)$ gauge symmetry.	21
1.3	Additional Feynman diagrams associated with the Dirac Lagrangian under non-abelian gauge symmetry $SU(n)$	23
1.4	Illustration of the Higgs potential $V(\phi^\dagger\phi) = -\mu^2\phi^\dagger\phi + \lambda \phi^\dagger\phi ^2$	26
1.5	Feynman diagrams associated with the propagation of the Higgs field.	27
1.6	Feynman diagram associated with the interaction between the symmetry breaking complex scalar field and the gauge bosons.	27
1.7	Feynman diagram associated with the interaction between the symmetry breaking complex scalar field and the fermions.	28
1.8	Theoretical Higgs mass constraints from same-sign $W_L^\pm W_L^\pm$ scattering, triviality and vacuum stability.	31
1.9	Interplay between the top quark, W and Higgs bosons masses in the Standard Model.	32
1.10	Evolution of the Higgs boson production cross-section at $\sqrt{s} = 8$ TeV proton–proton colliders, total width and branching ratios as a function of the Higgs boson mass.	33
1.11	95 % CL Higgs boson mass exclusion limits from LEP and Tevatron.	34
1.12	χ^2 of the fit to all electroweak observables as a function of the Higgs mass, and limits from LEP and Tevatron experiments.	35
1.13	Observed local p -value as a function of the hypothesised Higgs boson mass using the 2011 and the early 2012 data collected at the LHC.	35
1.14	Total and fiducial production cross-sections of SM processes at the LHC (theory and ATLAS/CMS measurements).	37
1.15	Higgs mass measurements from the ATLAS and CMS experiments in the $H \rightarrow \gamma\gamma$ and $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ decay channels with Run 1 data, and combination of these measurements.	39
1.16	Best fit values for signal strengths and couplings combining ATLAS and CMS measurements using the Run 1 dataset.	39
1.17	Alternative pure J^P states tested against the SM 0^+ hypothesis with Run 1 data in ATLAS and CMS.	40
1.18	Higgs production cross-section through different modes as a function of the centre-of-mass energy.	41

2	The Large Hadron Collider	43
2.1	The CERN accelerator complex and the LHC injection chain.	47
2.2	Schematic layout of the main LHC ring.	48
2.3	Cross-section of a main LHC dipole.	49
2.4	Nominal beam structure at the LHC.	50
2.5	Proton PDFs from the MSTW collaboration.	53
2.6	Fraction of the LHC uptime spent in the various mode during Run 2.	54
2.7	Cumulative delivered luminosity measurement by the ATLAS experiment as a function of time.	55
2.8	Recorded luminosity-weighted distribution of the mean number of interactions per crossing in the ATLAS experiment.	55
3	ATLAS: A Toroidal LHC ApparatuS	57
3.1	General cut-away view of the ATLAS detector.	59
3.2	Schematic of the ATLAS Inner Detector.	61
3.3	r - z cross-section view of an ATLAS Inner Detector quadrant with the localisation of the active layers and their structural envelope.	62
3.4	General view of the various calorimeter systems of ATLAS.	65
3.5	Mass stopping power of muons in copper as a function of $\beta\gamma = p/Mc$	66
3.6	Interaction of electrons and photons with lead.	67
3.7	Structure of the active/passive material sandwich for the LAr electromagnetic calorimeter.	70
3.8	Current produced in the liquid argon gap as a function of time and output signal from the electronics after bipolar shaping.	71
3.9	Scheme of the “accordion” electromagnetic calorimeter.	72
3.10	Nominal high-voltage settings in the outer and inner endcap wheels as a function of η	73
3.11	Scheme of the barrel electrodes before folding.	74
3.12	Scheme of the ECal electrodes in the endcap outer wheel before folding.	76
3.13	Cross-talk results from the test beam measurements of the LAr barrel module 0.	77
3.14	Layout of the tile calorimeter and scheme of HCal module.	77
3.15	HEC calorimeter general and electrode layout.	78
3.16	Schematic of the electrode structure in the first layer of the forward calorimeter (FCal1). The Moliere radius R_M is represented by a solid disk. [97]	80
3.17	Cross section of the muon system in the barrel seen in the x - y plane.	83
3.18	Cross section of a quadrant of the muon system in the endcap seen in the z - y plane.	83
3.19	Illustration of the toroid systems and produced magnetic field as a function of η	84
3.20	Simplified view of the ATLAS trigger system in Run 2.	85
3.21	Example of L1 and HLT trigger rates in one of the 2015 runs.	86
3.22	Distribution of HLT processing time per event in one 2015 run.	86
4	Reconstruction and identification of physics objects	89
4.1	Data acquisition chain for the ATLAS LAr calorimeters.	91
4.2	Average transverse energy in zero-bias events per unit $\eta \times \phi$ area and unit number of interactions per bunch crossing as a function of the distance from the beginning of the bunch train.	92

4.3	Expected combined electronic and pileup noise per cell as a function of $ \eta $ for each layer of the LAr calorimeter.	93
4.4	Schematic of the supercluster building algorithm.	94
4.5	Schematic of the electron/photon supercluster disambiguation algorithm.	94
4.6	Electron identification efficiency as a function of the pileup and electron p_T using partial Run 2 data with the ATLAS detector.	95
4.7	Example performance of the fJVT.	98
4.8	MV2C10 BDT score for b -jets, c -jets and light-flavour jets.	99
5	Calibration of the electromagnetic calorimeter	101
5.1	Summary of the steps from uncalibrated cells/cluster energy to a calibrated physics object.	103
5.2	Distribution of the calibrated energy divided by the truth energy for electrons within $1.4 < \eta < 1.6$	104
5.3	Energy resolution of electrons as a function of their pseudo-rapidity, for various generated energies.	105
5.4	Example correlation plot between $E_{1/2}^{\text{dist}}/E_{1/2}^{\text{nom}}$ and $E_0^{\text{dist}}/E_0^{\text{nom}}$ for various distorted geometries.	108
5.5	Data-to-MC ratio of $E_{1/2}$ as a function of $ \eta $ for photons from radiative $Z \rightarrow \mu\mu\gamma$ decays sample and from the inclusive sample.	109
5.6	PS scale α_{PS} and its uncertainty across the $ \eta $ range.	110
5.7	Relative difference of $\Delta E_{1/2}^{\text{data}}$ in 2015+2016 data after Run 1 layer corrections applied and MC with nominal geometry, and weighted-average sensitivity over various MC with distorted geometries, as a function of $ \eta $	112
5.8	Difference of material estimate in the early Run 2 nominal simulation including (not improved) IBL and in the data.	113
5.9	Total material difference between data and nominal early Run 2 simulation as well as its associated systematic uncertainty.	114
5.10	Material budget in units of X_0 from the improved Run 1 geometry, as a function of $ \eta $	114
5.11	Material budget in the Inner Detector in units of X_0 as a function of $ \eta $ after installation of the IBL.	115
5.12	Di-electron invariant mass distribution from $Z \rightarrow ee$ events in data before correction and in MC before smearing.	116
5.13	Example of χ^2 compatibility between generated template and data for M_{ee} distribution as a function of the scale factor α_{ij} and the additional constant term c'_{ij}	118
5.14	Illustration of the various steps in the determination of $\widehat{\alpha}_{ij}$ and $\widehat{c'_{ij}}$	119
5.15	Scale factors α and additional constant term c' as a function of η in the various years of Run 2 data taking.	122
5.16	Comparison of the di-electron invariant mass under the $Z \rightarrow ee$ peak in data and simulation after scaling the data and smearing the MC.	123
5.17	Di-electron invariant mass per slice of $\langle\mu\rangle$ normalised to the average over Run 2.	123
5.18	Extrapolation uncertainty of the energy scale factors α for electrons as a function of their transverse energy E_T (2015+2016 calibration analysis).	124
5.19	Extrapolation uncertainty of the additional constant term for electrons as a function of their transverse energy E_T (2015+2016 calibration analysis).	124

5.20	Fit of the pseudo proper-time of J/ψ candidate events in the 2015 and 2016 data.	126
5.21	Example fit of the J/ψ shape in the MC and in the data.	127
5.22	Distribution of electrons from candidate J/ψ decay in the data.	127
5.23	Residual scales $\Delta\alpha$ extracted with the $J/\psi \rightarrow e^+e^-$ method as a function of η , after the $Z \rightarrow ee$ central scales are applied.	128
6	ECal first and second layers intercalibration	129
6.1	p_T distribution of muons used in this analysis and dE/dx as a function of muon momentum.	131
6.2	Example energy distribution for data and MC in S1 and S2.	132
6.3	Layer intercalibration scales $\alpha_{1/2}$ as a function of $ \eta $ derived in Run 1.	133
6.4	Muon energy deposit as a function of η_{corr} in data for the barrel region $0.7 <$ $\eta_{\text{corr}} < 0.9$, in S1 and S2.	134
6.5	Muon energy deposit as a function of η_{corr} in data for the barrel region $1.3 <$ $\eta_{\text{corr}} < 1.5$, in S1 and S2.	135
6.6	Muon energy deposit as a function of η_{corr} in data for the endcap region $1.35 <$ $\eta_{\text{corr}} < 1.55$, in S1 and S2.	136
6.7	Number of cells actually used to compute the muon deposited energy in the 2015 and 2016 dataset for S1.	137
6.8	Example distribution of energy deposited in S1 and S2 in various $ \eta_{\text{corr}} $ region using 2016 data.	138
6.9	MPV extracted from the Landau * Gaussian fit in each $ \eta_{\text{corr}} $ region for S1 and S2.	140
6.10	Layer intercalibration scales derived from the Landau * Gaussian analytical fit with the 2016 dataset, compared to the Run 1 result.	141
6.11	Comparison of the MPV in S2 and of the layer scales $\alpha_{1/2}$ as a function of $ \eta_{\text{corr}} $ when using one or two cells to reconstruct the muon deposit energy in S2.	142
6.12	Average transverse energy in clusters built from zero-bias events of a single data run in the $ \eta $ region $[1.81, 2.37]$ as a function of the bunch crossing number, before correction of the residuals.	143
6.13	MPV extracted from the Landau * Gaussian fit in each $ \eta_{\text{corr}} $ region for S1 and S2 with the standard analysis and after removing events in the first 20 bunches of the train.	144
6.14	Double ratio $E_{1/2}^{\text{Data}}/E_{1/2}^{\text{MC}}$ extracted with all events or with event located after the 20 th bunch of each train.	145
6.15	MPV extracted from the Landau * Gaussian fit in each $ \eta_{\text{corr}} $ region for S1 and S2 with the standard analysis (no isolation requirement) or after requiring a Tight isolation cut.	146
6.16	Double ratio $E_{1/2}^{\text{Data}}/E_{1/2}^{\text{MC}}$ extracted with and without isolation cuts.	147
6.17	Schematic of the cell window opened by the reconstruction algorithm around the expected muon track.	147
6.18	Noise distribution in S1 and S2 computed with the average cell energy taken from the window border.	148
6.19	MPV extracted from the Landau * Gaussian fit in each $ \eta_{\text{corr}} $ region for S1 and S2 with and without applying the ambient noise subtraction procedure.	149
6.20	Double ratio $E_{1/2}^{\text{Data}}/E_{1/2}^{\text{MC}}$ extracted with and without ambient noise subtraction.	150
6.21	MPV extracted from the Landau * Gaussian fit in each $ \eta_{\text{corr}} $ region for S1 and S2 in several pileup bins.	151

6.22	Noise distribution in the data and best gaussian fit for the pileup bin $20 < \langle \mu \rangle < 22$ and the $ \eta_{\text{corr}} $ bin $[1.00, 1.10]$, in S1 and S2.	152
6.23	Pileup profiles in the 2015+2016 dataset compared to the generated profile in the MC15c used in this study.	152
6.24	Example of MPV distribution as a function of $\langle \mu \rangle$ for the endcap in S1 and for the barrel in S2.	152
6.25	Comparison of the energy deposit MPV with (red) and without (blue) applying the extrapolation method in S1 and S2.	154
6.26	Double ratio $E_{1/2}^{\text{Data}}/E_{1/2}^{\text{MC}}$ using the muon MPV method as a function of $ \eta_{\text{corr}} $ using the 2015+2016 dataset.	155
6.27	Result of the extrapolation method closure test in the MC, for S1 and S2 as a function of $ \eta_{\text{corr}} $	156
6.28	Scheme of the folded electrodes in S2 in the (R, ϕ) plane representing 5 cells in the ϕ direction, and its associated opened 5×5 window in the (η, ϕ) plane around the muon track used in the analysis.	157
6.29	Difference of energy between the highest energetic cell and the second closest cell in S2, in the data and in the MC.	157
6.30	Comparison of the extrapolated MPV in S2 and of the double ratio $E_{1/2}^{\text{Data}}/E_{1/2}^{\text{MC}}$ as a function of $ \eta_{\text{corr}} $ when using either the second most energetic or second closest cell as the second cell choice for S2.	158
6.31	Scheme of the folded electrodes in S1 in the (R, ϕ) plane representing one cell in the ϕ direction, and its associated scheme in the (η, ϕ) plane.	159
6.32	Comparison of the fitted MPV and of the double ratio $E_{1/2}^{\text{Data}}/E_{1/2}^{\text{MC}}$ as a function of $ \eta_{\text{corr}} $, with or without a $ \phi^{\text{extr.}} - \phi^{\text{cell}} < 0.04$ cut applied.	160
6.33	Comparison of the fitted MPV and of the double ratio $E_{1/2}^{\text{Data}}/E_{1/2}^{\text{MC}}$ as a function of $ \eta_{\text{corr}} $, with or without a $ \eta^{\text{extr.}} - \eta^{\text{cell}} < 0.008$ cut applied.	162
6.34	Summary of systematic uncertainties affecting the relative layer intercalibration of the first and second layers of the electromagnetic calorimeter from the muon measurement using the MPV method as a function of $ \eta_{\text{corr}} $	163
6.35	Example of muon energy distribution in the first sampling of the electromagnetic calorimeter for $0.125 < \eta < 0.150$ and $21 < \langle \mu \rangle < 22$ using 2015 and 2016 data. Several definitions of the truncated mean are illustrated.	163
6.36	Examples of truncated mean values as a function of $\langle \mu \rangle$ in S1 and S2 for both data and MC, as well as the linear extrapolation.	164
6.37	Truncated mean pileup extrapolation results as a function of $ \eta $ in S1 and S2, for data and MC compared to the (non-extrapolated) result using the MC generated at $\langle \mu \rangle = 0$	164
6.38	Layer intercalibration scales $\alpha_{1/2}$ as a function of $ \eta $ extracted with various definitions of the truncated mean.	165
6.39	Breakdown of layer intercalibration uncertainties with TM method.	166
6.40	Comparison of the common sources of uncertainty for the MPV and TM methods.	167
6.41	Difference TM - MPV of $\alpha_{1/2}$ central value.	168
6.42	Summary of systematic uncertainties for the MPV method and the TM method.	168
6.43	Comparison of the layer intercalibration scales $\alpha_{1/2} = E_{1/2}^{\text{Data}}/E_{1/2}^{\text{MC}}$ using the 2015+2016 dataset, as a function of $ \eta_{\text{corr}} $ for the MPV method, the TM method, and their average.	169
6.44	Comparison the layer intercalibration scales $\alpha_{1/2} = E_{1/2}^{\text{Data}}/E_{1/2}^{\text{MC}}$ as a function of $ \eta_{\text{corr}} $ in the Run-1 and the 2015+2016 analyses.	170

6.45	Comparison of the layer intercalibration scales $\alpha_{1/2}$ using the MPV method on Run 1 data, with one or two cells in the energy determination of the muon deposit in the S2.	171
6.46	Average di-electron invariant mass and ratio of average invariant mass to average invariant mass in nominal MC, as a function of the energy ratio in S1 and S2 ($E_{1/2}$) for $0.4 < \eta < 0.6$	172
6.47	Comparison of the $\alpha_{1/2}$ scales derived with the electron and muon methods.	173
6.48	Slope of the data/MC distribution in the inclusive bin $1.5 < \eta < 1.8$ for various cuts applied on top of the standard analysis selection.	174
6.49	Slope of the data/MC distribution ratio after applying a 3% bias on E_1 in data, as a function of $ \eta $ and for various di-electron invariant mass selection windows. . .	174
7	Higgs boson phenomenology and EFT	175
7.1	SM Higgs boson production cross-section at pp colliders at $\sqrt{s} = 13$ TeV as a function of its mass, with main production mode breakdown.	176
7.2	Leading order Feynman diagram for the $gg \rightarrow H$ process.	177
7.3	Leading order diagram for the $qq \rightarrow qqH$ process.	178
7.4	Leading order diagram for the $qq \rightarrow VH$ process and additional loop-level diagrams for the $gg \rightarrow ZH$ process.	178
7.5	Examples of LO diagrams for the $pp \rightarrow ttH$ process.	179
7.6	SM Higgs boson branching ratios as a function of its mass.	180
7.7	Mass distribution for the $H \rightarrow WW$, $H \rightarrow \tau^+\tau^-$, $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ and $H \rightarrow \gamma\gamma$ analyses using $36,1 \text{ fb}^{-1}$ of data collected at $\sqrt{s} = 13$ TeV at the ATLAS experiment.	182
7.8	Examples of diagrams contributing to the $H \rightarrow \gamma\gamma$ decay at leading order.	183
7.9	κ -framework result from combined ATLAS and CMS measurements using the Run 1 data of the LHC.	185
7.10	Definition of truth-level STXS bins in the Stage-1.1.	187
7.11	Effective Fermi theory and resolved (high energy) theory of the β decay.	188
7.12	Signal strength from partial Run 2 STXS results by the ATLAS and CMS experiments.	195
7.13	Couplings results from partial Run 2 ATLAS and CMS measurements interpreted in terms of couplings modifiers in the κ -framework.	196
7.14	ATLAS couplings results in the $H \rightarrow 4\ell$ decay channel using $36,1 \text{ fb}^{-1}$ of 13 TeV data collected during Run 2 and interpreted in the Higgs Characterisation framework.	197
7.15	ATLAS couplings results in the $H \rightarrow \gamma\gamma$ decay channel plane using $36,1 \text{ fb}^{-1}$ of 13 TeV data collected during Run 2.	198
7.16	Distributions of CP discriminants used in the measurement of the f_{a3} parameter in the $H \rightarrow 4\ell$ and VBF $H \rightarrow \tau^+\tau^-$ analyses using a partial Run 2 dataset collected by the CMS experiment.	199
7.17	Combined likelihood scan of the $H \rightarrow \tau^+\tau^-$ and $H \rightarrow 4\ell$ results on effective couplings using all the Run 1 plus partial Run 2 data collected by the CMS experiment.	200
7.18	Distribution of CP -odd <i>optimal observable</i> for simulated VBF events with various $\tilde{d}$ hypotheses and simulated events from various processes, and compared to the data in the $\tau_{\text{lep}}\tau_{\text{had}}$ signal region (using $20,3 \text{ fb}^{-1}$ of 8 TeV data collected by the ATLAS experiment in 2012).	201
7.19	Expected and observed likelihood scans using the optimal observable as discriminating variable.	202

8	Experimental aspects of the Higgs to Four Leptons decay	204
8.1	Leading order diagram for the $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ decay.	205
8.2	Schematic of an event layout after all generation steps.	206
8.3	Selection efficiency for the signal, irreducible and reducible background as a function on the cut applied on the vertex fit χ^2 .	210
8.4	Impact of the FSR recovery on the four-lepton invariant mass for events benefiting from FSR recovery.	211
8.5	Feynman diagrams corresponding to the dominant background topologies for the $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ decay channel.	213
8.6	m_{12} distributions generated from MC-simulated events for different processes and in the various $\ell\ell + \mu\mu$ control regions.	215
8.7	Background component simultaneous fit to the full Run 2 data using m_{12} as observable in each $\ell\ell + \mu\mu$ control region.	216
8.8	Distribution of m_{12} in the relaxed isolation and d_0 -significance validation region of data, compared to the background yields extrapolated from the data control regions.	217
8.9	Fractional p_T balance between the ID and MS measurements for additional muons in $Z+\mu$ events shown for data and MC simulation.	219
8.10	Distribution of d_0 significance for additional muons in $Z+\mu$ events for data and MC simulation.	220
8.11	Data fit of the n_{InnerPix} variable in the $3\ell+X$ CR.	221
8.12	Comparison of the FSR-corrected $m_{4\ell}$ distribution in the HF-enriched $3\ell+X$ control region between data and the various background contributions from MC.	222
8.13	Template shapes built from the $Z+X$ sample used in the n_{InnerPix} fit in the $3\ell+X$ CR (γ and f), compared to the shape from the q component directly extracted from the $3\ell+X$ CR.	223
8.14	Truth production categories as defined in the STXS Stage-0 and Stage-1.1, and corresponding reconstruction categories in the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ analysis for in the signal and sidebands regions.	230
8.15	Expected yield in each reconstruction category (signal and sideband) of the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ analysis in terms of STXS Stage1.1 truth category, assuming an integrated luminosity of 139 fb^{-1} at $\sqrt{s} = 13 \text{ TeV}$.	231
8.16	Expected relative composition in each reconstruction categories of the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ analysis in terms of reduced STXS Stage1.1 truth category.	231
8.17	Observed and expected NN output (pre-fit) distributions in the different reconstruction categories of the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ Run 2 analysis for an integrated luminosity of 139 fb^{-1} at $\sqrt{s} = 13 \text{ TeV}$ and assuming $m_H = 125 \text{ GeV}$ (1/2).	233
8.18	Observed and expected NN output (pre-fit) distributions in the different reconstruction categories of the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ Run 2 analysis for an integrated luminosity of 139 fb^{-1} at $\sqrt{s} = 13 \text{ TeV}$ and assuming $m_H = 125 \text{ GeV}$ (2/2).	234
8.19	Distribution of the four-lepton invariant mass seen by the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ Run 2 analysis in the low mass range, showing the data and compared to various contributions from the MC.	236
8.20	Distributions of (a) the four-lepton system transverse momentum and (b) the number of jets shown in the signal region, as seen in the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ Run 2 analysis.	236

9	Measurement of CP-odd Higgs-boson couplings	239
9.1	Illustration of the decay angles defined in the $H \rightarrow 4\ell$ final state.	241
9.2	Example distribution of CP -sensitive observables for various values of the CP -odd couplings and compared to the SM distribution.	241
9.3	Distribution of first and second order optimal observables for the process $pp \rightarrow jjH$ in truth-level MADGRAPH samples generated assuming a pure SM hypothesis, or SM with an additional CP -odd contribution, or SM with an additional CP -even contribution.	243
9.4	Distribution of $ \mathcal{M}_{\text{SM}} ^2$ and $\mathcal{OO}_1$ computed for the process $pp \rightarrow jjH$ in a SM-like generated sample using various PDFs for the weighting of the partonic processes.	245
9.5	Distribution of $ \mathcal{M}_{\text{SM}} ^2$ and $\mathcal{OO}_1$ for the process $pp \rightarrow jjH$, and for various choices of Q^2 in the evaluation of the PDFs used in the weighting of the partonic processes.	246
9.6	Comparison of $\mathcal{OO}_1$ computation for the process $pp \rightarrow jjH$ using either the $\tilde{c}_{zz} = 1$ or $\tilde{c}_{zz} = 2.65$ hypotheses, and applied on a SM-like generated sample.	247
9.7	Breakdown of the matrix element result returned by each partonic process computing class for one event.	248
9.8	Bias introduced on $ \mathcal{M}_{\text{SM}} ^2$ and $\mathcal{OO}_1$ for the process $pp \rightarrow jjH$ by removing some of the partonic diagrams.	249
9.9	Comparison of morphed shapes of the $\mathcal{OO}_{1,jj}$ observable for two different morphing bases and of the MC generated shape.	253
9.10	Log-likelihood ratio to the Standard Model with the using $\mathcal{OO}_{1,jj}$ observable along the $\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$, $\tilde{c}_{\gamma\gamma}$, $c_{H\tilde{W}}$, $c_{H\tilde{W}B}$ and $c_{H\tilde{B}}$ scanning axes, assuming an integrated luminosity of 140fb^{-1} and no background contamination.	256
9.11	VBFH $\rightarrow 4\ell$ cross-section times branching ratio normalised to the Standard Model value as a function of $\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$, $\tilde{c}_{\gamma\gamma}$, $c_{H\tilde{W}}$, $c_{H\tilde{W}B}$ or $c_{H\tilde{B}}$, and best quartic fits.	258
9.12	Comparison of log-likelihood ratios to the Standard Model scanning along $\tilde{c}_{zz}$ with the $\mathcal{OO}_{1,jj}$ observable, and using generated or morphed samples. An integrated luminosity of 140fb^{-1} and no background contamination is assumed.	259
9.13	Comparison of the log-likelihood ratios to a SM asimov scanning along the $\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$ and $c_{H\tilde{W}}$ axes using the $\mathcal{OO}_{1,jj}$ observable with and without background, and after splitting the analysis into 5 BDT bins. An integrated luminosity of 140fb^{-1} is assumed in each case.	260
9.14	Comparison of the log-likelihood scans along the $\tilde{c}_{zz}$ axis for a SM asimov using the $\mathcal{OO}_{1,jj}$ observable for truth and fully reconstructed samples.	261
9.15	Likelihood scans along the $\tilde{c}_{zz}$ axis using the $\mathcal{OO}_{1,jj}$ observable for various asimov datasets.	263
9.16	Response of VBF samples generated with various values of $\tilde{c}_{zz}$ to a neural network classifier trained on a VBF + VH sample generated with various values of $\tilde{c}_{zz}$	264
A	Layout of distorted geometries	271
A.1	Impacted area of distorted geometry configuration “A”.	271
A.2	Impacted area of distorted geometry configuration “C’+D’”.	272
A.3	Impacted area of distorted geometry configuration “E’+L’”.	272
A.4	Impacted area of distorted geometry configuration “F+M+X”.	273
A.5	Impacted area of distorted geometry configuration “N”.	273
B	STXS Stage-0 categorisation yields	275

C Synthèse en français	277
C.1 Les particules élémentaires et leur classification.	277
C.2 Diagrammes dominants pour la production du boson Higgs par les modes ggF, VBF, VH et ttH	278
C.3 Test de vraisemblance de différentes hypothèses de spin-parité J^P pure par rapport à l'hypothèse du Modèle Standard ($J^P = 0^+$) par ATLAS et CMS en utilisant les données collectées au Run 1.	279
C.4 Exemples de limites posées sur le mixte de CP du boson de Higgs : analyses $H \rightarrow \tau^+\tau^-$ 2012 d'ATLAS sur le paramètre $\tilde{d}$ utilisant une observable optimale et $H \rightarrow 4\ell$ 2016 d'ATLAS sur le paramètre κ_{AVV} dans le cadre « Higgs Characterisation » utilisant les informations de section efficaces.	280
C.5 Nombre moyen d'interactions par croisement de faisceaux enregistré par ATLAS pour chaque année de prise de données et pondéré par la luminosité instantanée.	282
C.6 Vue éclatée du détecteur ATLAS.	283
C.7 Schéma des cellules du calorimètre électromagnétique.	284
C.8 Schéma des électrodes du tonneau du calorimètre électromagnétique avant pliage.	284
C.9 Schéma des différentes étapes de l'étalonnage en énergie des électrons.	285
C.10 Exemple de distribution et d'ajustement de la distribution du dépôt d'énergie des muons dans le calorimètre.	286
C.11 MPV extraite par la méthode Landau * Gaussian pour chaque région de $ \eta $ dans L1 et L2.	287
C.12 Facteurs d'étalonnage relatifs en fonction de $ \eta $ et dérivés avec la méthode MPV en utilisant les données de 2016 sans extrapolation, les données 2015 et 2016 avec extrapolation et patrons de bruit, et comparés au résultat combiné du Run 1.	288
C.13 Exemple de distribution de bruit dans L2 avec meilleur ajustement d'une gaussienne et exemple de distribution de MPV en fonction de l'empilement $\langle\mu\rangle$ dans L2.	289
C.14 Comparaison des MPV dans L1 et L2 avec ou sans extrapolation.	290
C.15 Résumé des systématiques de la méthode MPV et résultat final des facteurs d'étalonnage relatif L1/L2.	291
C.16 Distributions de masse invariante du système de quatre leptons et du résultat du réseau de neurone pour VBF.	292
C.17 Distributions d'observables sensibles aux effets de CP pour le Modèle Standard et plusieurs couplages anomaux.	293
C.18 Scan de vraisemblance en utilisant $\mathcal{OO}_{1,jj}$ et section efficace VBF multipliée par de rapport d'embranchement $H \rightarrow 4\ell$ théorique en fonction du couplage $\tilde{c}_{zz}$	294
C.19 Scan de vraisemblance du paramètre $\tilde{c}_{zz}$ en utilisant l'observable $\mathcal{OO}_{1,jj}$ contre l'hypothèse du Modèle Standard (asimov), variant la présence de bruits de fond et d'effets du détecteur.	296

List of Tables

1	The Standard Model of particle physics	15
1.1	Examples of symmetries and their conserved quantities.	21
1.2	Summary table of the Standard Model fermion content and their properties in the electroweak theory.	24
2	The Large Hadron Collider	43
2.1	A few examples of the main accelerators filling scheme in Run 2.	51
2.2	LHC configuration parameters during Run 1 and Run 2.	56
3	ATLAS: A Toroidal LHC ApparatuS	57
3.1	A few examples of η - θ correspondence.	60
3.2	Atomic and nuclear parameters of some elements composing the ATLAS electromagnetic calorimeter.	69
3.3	Summary of the readout granularity and number of readout channels, the nominal high-voltage settings, and building component parameters in the LAr electromagnetic calorimeter (barrel and endcap).	75
3.4	Summary of the readout granularity and number of readout channels, the nominal high-voltage settings, and building component parameters for the tile hadronic central calorimeter, the LAr hadronic endcap calorimeter and the LAr forward calorimeters.	79
3.5	Summary of the parameters for each of the four types of muon chambers.	80
4	Reconstruction and identification of physics objects	89
4.1	Listing of the discrimination variables used in the electron and photon identification.	96
4.2	Muon and background identification efficiencies for all working points.	97
4.3	b -jet efficiency and c -jet, τ -jet and light-flavour jet rejection factor using the MV2C10 algorithm at its four nominal working points.	99
5	Calibration of the electromagnetic calorimeter	101
5.1	List of distorted geometries used in the evaluation of the correlation factor $A(\eta)$	108
5.2	η binning used for the determination of the energy scale factors α_i and the additional constant terms c'_i in the in-situ calibration analysis.	118

5.3	Boundaries of the η binning used for the residual energy scales determination with $J/\psi \rightarrow e^+e^-$ electrons.	125
6	ECal first and second layers intercalibration	129
6.1	Binning used for the layer intercalibration analysis with muons.	136
6.2	Summary of systematic uncertainties included in the MPV and TM evaluation. . .	167
7	Higgs boson phenomenology and EFT	175
7.1	Individual production mode cross-sections contributing to the $pp \rightarrow H$ process and their related uncertainties, for a SM Higgs boson of mass $m_H = 125$ GeV at $\sqrt{s} = 13$ TeV.	177
7.2	Branching ratio and their uncertainties for the Higgs boson of mass $m_H = 125$ GeV.	180
7.3	The 76 dimension-6 baryon number conserving operators in the so-called Warsaw basis.	190
7.4	Combined 68 % CL and 95 % CL intervals from the $H \rightarrow \tau^+\tau^-$ and $H \rightarrow 4\ell$ results on effective couplings using all the Run 1 plus partial Run 2 data collected by the CMS experiment.	201
8	Experimental aspects of the Higgs to Four Leptons decay	204
8.1	Summary of the generators and settings used for the generation of Standard Model Higgs boson signal samples.	207
8.2	Modelling of each contribution entering the $\ell\ell + \mu\mu$ background estimation. . . .	214
8.3	Results of the MC closure test for the $\ell\ell + \mu\mu$ background modelling.	215
8.4	Final $\ell\ell + \mu\mu$ background estimates in the “relaxed” validation region in the full $m_{4\ell}$ range for each background component.	217
8.5	Comparison of the $\ell\ell + \mu\mu$ background estimation in the signal region obtained from the three methods.	218
8.6	Efficiency of isolation and impact parameter selections for background muons selected in $Z+\mu$ events.	218
8.7	Isolation and impact-parameter efficiencies for background muons in the light-flavour enriched and heavy-flavour enriched $Z+\mu$ samples.	219
8.8	Fit result for the various background component yields in the $3\ell+X$ control region using the full Run 2 dataset.	223
8.9	Data/MC efficiency scale factors for the f and γ component of the background evaluated on the X electron candidate of the $Z+X$ control region.	224
8.10	Typical impact of the dominant systematic uncertainties on the measured inclusive fiducial cross-section and on the STXS Stage-0 production mode cross-sections for the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ Run 2 analysis.	226
8.11	Input variables used to train the MLP and the two rNNs for each reconstruction category of the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ Run 2 analysis.	235
8.12	Number of expected and observed events with final-states breakdown in the signal region for the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ Run 2 analysis.	235
8.13	Expected and observed Higgs boson signal and background yields in the different reconstruction categories of the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ Run 2 analysis in the signal region for an integrated luminosity of 139 fb^{-1} at $\sqrt{s} = 13$ TeV and assuming the SM Higgs boson signal with a mass $m_H = 125$ GeV.	237

9	Measurement of CP-odd Higgs-boson couplings	239
9.1	Comparison of limits on the CP -odd Wilson coefficients obtained using truth-level pure VBF samples and from cross-section measurements.	257
9.2	Truth-level sensitivity scans along the $\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$ and $c_{H\tilde{W}}$ axes, without background and assuming 100 % reconstruction efficiency and with backgrounds (with and without BDT-binned analysis) assuming 45 % reconstruction efficiency.	260
A	Layout of distorted geometries	271
B	STXS Stage-0 categorisation yields	275
B.1	Expected number of Higgs boson events for each STXS stage-0 production mode in the different reconstruction categories of the ATLAS $H \rightarrow ZZ^{(*)} \rightarrow 4\ell$ Run 2 analysis in the signal region for an integrated luminosity of 139 fb^{-1} at $\sqrt{s} = 13 \text{ TeV}$ and assuming the SM Higgs boson signal with a mass $m_H = 125 \text{ GeV}$	275
C	Synthèse en français	277
C.1	Comparaison des limites sur différents coefficients de Wilson impairs par CP ($\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$, $\tilde{c}_{\gamma\gamma}$, $c_{H\tilde{W}}$, $c_{H\tilde{B}}$, $c_{H\tilde{W}B}$) obtenus par un scan de vraisemblance avec un échantillon VBF pur, et des valeurs de couplage donnant une section efficace multipliée VBF par le rapport d'embranchement $H \rightarrow 4\ell$ 25 % supérieure à la valeur du Modèle Standard.	295
C.2	Résultats de sensibilité aux effets des coefficients de Wilson $\tilde{c}_{zz}$, $\tilde{c}_{z\gamma}$, $c_{H\tilde{W}}$ dérivés des scans de vraisemblance utilisant $\mathcal{OO}_{1,jj}$, avec ou sans bruits de fond et effets du détecteur, et en séparant l'analyse en cinq régions de BDT, pour les niveaux de confiance 68 % et 95 %.	296

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¹↑ Typical day at LAL, 11:53am: "Antoine you coming? Yes, two minutes. Sure? Hum, actually... no, just go without me, I'll join soon(ish).".

²↑ Et oui Konie, je sais, tu n'es plus une stagiaire maintenant. Mais fais attention aux coin-coins quand même.

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Enfin « je verrais bien une dernière catégorie, où s'inscrirait, unique, ... »³ toi qui m'a accompagné ces dernières années et dont la seule présence suffisait à me redonner le sourire. Pour ça, et pour tout le reste, merci.

A.

³↑ Pierre Desproges, Chroniques de la Haine Ordinaire, « l'Humanité » (10 mars 1986).

Titre : Inter-calibration des couches du calorimètre électromagnétique d'ATLAS et mesure de couplages CP -impairs du boson de Higgs dans son canal de désintégration en quatre leptons avec les données du Run 2 au LHC

Mots clés : ATLAS, Calibration en énergie, Higgs, 4 leptons, Couplages CP -impair BSM, Théorie Effective

Résumé : Après la découverte du boson de Higgs en 2012 au LHC, l'intérêt s'est porté sur l'étude de ses propriétés pour vérifier le Modèle Standard et pour sonder la nouvelle physique. L'une de ses propriétés fondamentales est sa spin-parité (CP), dont le Modèle Standard prédit la valeur 0^+ . Les analyses menées sur les données récoltées au Run 1 du LHC ont rejeté toutes les hypothèses d'état pur de spin-parité autre que cette valeur. Cependant des états mixtes de CP sont toujours possibles, ce qui indiquerait une violation de symétrie CP dans le secteur du Higgs.

La première partie de cette thèse se concentre sur la calibration du calorimètre électromagnétique d'ATLAS permettant d'atteindre une résolution de l'ordre du pour mille sur l'énergie des électrons et photons, primordiaux dans les analyses du boson de Higgs. Une des étapes est l'inter-calibration des couches du calorimètre électromagnétique, corrigeant des effets résiduels de calibration électronique et de diaphonie (cross-talk). La méthode établie au Run 1 a montré ses limites devant les niveaux d'empilement mesurés au Run 2, et une nouvelle analyse a été alors développée, assurant le contrôle précis des incertitudes systématiques.

La deuxième partie de cette thèse porte sur la mesure des couplages CP -impairs du boson de Higgs aux bosons vecteurs, étudié dans le canal de désintégration du boson de Higgs en quatre leptons. Malgré une faible statistique, ce canal offre une signature propre et un rapport signal sur bruit de plus de deux, permettant l'analyse précise des propriétés du boson de Higgs. Le mode de production par fusion de bosons vecteurs offre la meilleure sensibilité aux effets de CP grâce à la présence de deux jets dans l'état final. La pollution venant du mode de production par fusion de gluon avec des jets additionnels est réduite grâce à l'utilisation de réseaux neuronaux. Pour distinguer de manière univoque les effets CP -impair d'éventuels effets CP -pair encore inconnus, une nouvelle variable est construite dont l'asymétrie de forme dépend uniquement d'effets CP -impairs. Composée d'éléments de matrice, cette variable utilise les informations cinématiques du boson de Higgs et des jets de manière maximale. Les résultats sont interprétés en termes de théorie effective, et la sensibilité statistique à 68% de confiance sur le coefficient de Wilson $\tilde{c}_{zz}$ est estimée à $[-0.80, 0.80]$.

Title: Layer Intercalibration of the ATLAS Electromagnetic Calorimeter and CP -odd Higgs Boson Couplings Measurements in the Four-Lepton Decay Channel with Run 2 Data of the LHC

Keywords: ATLAS, Energy calibration, Higgs boson, 4 leptons, BSM CP -odd couplings, Effective Field Theory

Abstract: After the Higgs boson discovery at the LHC in 2012, interest turned to Higgs boson property measurements to refine the tests of the Standard Model and probe for new physics. One of its key properties is its spin-parity (CP), predicted to be 0^+ in the Standard Model. Analyses of data collected during the Run 1 of the LHC rejected all pure spin-parity state other than 0^+ . However mixed CP states are still possible, and would indicate CP violation in the Higgs sector.

The first part of this thesis focuses on the ATLAS electromagnetic calorimeter calibration, needed to reach a permil level on electron and photon energy resolution which are of prime importance for Higgs boson studies. One step of the calibration sequence consists of the layer intercalibration of the electromagnetic calorimeter, needed to correct residual electronics miscalibration and cross-talk effects. The Run 1 method has proven to be unreliable for the pileup levels in Run 2 and a new method was developed, ensuring a precise control on the systematic uncertainties.

The second part of this thesis puts emphasis on the Higgs boson to vector boson CP -odd couplings, with the Higgs boson decaying to four leptons. This channel, despite low statistics, provides a clean signature and a signal-to-noise ratio over two, allowing for a precise determination of the Higgs boson properties. The vector-boson fusion production mode offers the best sensitivity to CP effects thanks to its two characteristic tagging jets in the final state. The contamination from the gluon fusion production mode with additional jets is reduced using neural networks. To unambiguously distinguish yet unknown CP -even from possible CP -odd effects, a variable whose shape asymmetry only depends on CP -odd effects is built. This observable is based on the matrix element computation, maximally using the kinematic information available from Higgs boson and associated jets. Results are interpreted in a context of effective field theory, and the statistical precision on the $\tilde{c}_{zz}$ Wilson coefficient is estimated to $[-0.80, 0.80]$ at the 68% confidence level.

