

MEASUREMENT OF ASSOCIATED PRODUCTION OF HIGGS BOSONS DECAYING TO PAIRS OF W BOSONS
WITH THE ATLAS DETECTOR AT THE LARGE HADRON COLLIDER

by

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Graduate Department of Physics
University of Toronto

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Abstract

This thesis describes the measurement of Higgs boson production in association with a vector boson, WH/ZH , using the $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$ and $H \rightarrow WW^* \rightarrow \ell\nu q\bar{q}$ decay channels, where ℓ is an electron or a muon. The analysis uses proton-proton collision data at centre-of-mass energy 13 TeV delivered by the Large Hadron Collider and recorded by the ATLAS detector between 2015 and 2018, corresponding to an integrated luminosity of 139 fb^{-1} . The analysis is performed in several channels across 2-, 3-, and 4-lepton final states and utilizes a diverse set of multivariate techniques for signal extraction. The WH and ZH production cross sections times the $H \rightarrow WW^*$ branching ratio are measured to be $0.13^{+0.08}_{-0.07}(\text{stat.})^{+0.05}_{-0.04}(\text{sys.}) \text{ pb}$ and $0.31^{+0.09}_{-0.08}(\text{stat.}) \pm 0.03(\text{sys.}) \text{ pb}$, respectively, in agreement with the Standard Model expectations.

To the people I love, who make my life as colourful as quantum chromodynamics.

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Contributions

My thesis describes the measurement of associated production of Higgs bosons (VH) in the WW^* decay channel ($H \rightarrow WW^*$) using the full LHC Run 2 proton-proton dataset collected by the ATLAS detector [1]. A particle physics analysis is a collaborative endeavour; accordingly, the results presented in this thesis are the culmination of the efforts of many individuals. I describe all aspects of the VH , $H \rightarrow WW^*$ measurement for completeness. In terms of my specific contributions, I served as the primary analyzer for the 2ℓ different-flavour, opposite-sign (DFOS) channel, produced the data and Monte Carlo (MC) inputs for all channels, developed the machinery for applying the misidentified (“fake”) lepton estimation for the 2ℓ and 3ℓ channels, estimated the theoretical systematic uncertainties for the 2ℓ DFOS and 3ℓ channels, contributed to the development of the inclusive and Simplified Template Cross Section (STXS) statistical fits, and provided software support in various capacities.

While this measurement has been the primary focus of my doctoral studies, I have additionally contributed to research and development towards ATLAS’s Inner Tracker (ITk) upgrade as well as future lepton colliders. Papers which I have served as one of the contacts for – which includes conceiving the study; obtaining, analyzing, and visualizing the results; and writing and editing the corresponding paper – are:

- “A Starry Byte — proton beam measurements of single event upsets and other radiation effects in ABCStar ASIC Versions 0 and 1 for the ITk strip tracker” (published in *JINST*) [2];
- “Strange quark as a probe for new physics in the Higgs sector” (published in the proceedings of the 2022 Snowmass Summer Study) [3];
- “Asymptotic formulae for estimating statistical significance in particle physics analyses” (published in *J. Phys. G: Nucl. Part. Phys.*) [4].

Papers which I have contributed to – in the form of ideas and/or data – and/or edited include:

- “Analysis of humidity sensitivity of silicon strip sensors for ATLAS upgrade tracker, pre- and post-irradiation” (submitted to *JINST*) [5];
- “Pre-Production Results from ATLAS ITk Strip Sensors Quality Assurance Testchip” (published in *JINST*) [6];
- “Electrical characterization of surface properties of the ATLAS17LS sensors after neutron, proton and gamma irradiation” (published in *Nucl. Instrum. Meth. A*) [7];
- “Charge collection study with the ATLAS ITk prototype silicon strip sensors ATLAS17LS” (published in *Nucl. Instrum. Meth. A*) [8];
- “The ABC130 barrel module prototyping programme for the ATLAS strip tracker” (published in *JINST*) [9];
- “Testbeam studies of barrel and end-cap modules for the ATLAS ITk strip detector before and after irradiation” (published in *Nucl. Instrum. Meth. A*) [10];

Beyond the VH , $H \rightarrow WW^*$ analysis, I have made significant contributions to the measurement of gluon-gluon fusion (ggF) and vector boson fusion (VBF) production of Higgs bosons in the WW^* decay channel using the full LHC Run 2 dataset collected by the ATLAS detector. These contributions include the production

of data and MC inputs, optimization of the lepton selections, the estimation of the fake lepton background, and the development of machinery for applying the fake lepton estimation. The analysis has resulted in two conference papers [11, 12] as well as a paper [13] submitted to *Phys. Rev. D*.

I have contributed to combined performance recommendations in ATLAS as an active member of the vertexing subgroup of the Inner Tracking performance group. I studied and developed the LHC Run 3 track-to-vertex association (TTVA) working points for a new vertexing algorithm. The TTVA working points describe the quality of the assignment of tracks to vertices for the purposes of jet building, tau building, and quantifying lepton isolation, among other aspects of reconstruction. The working points have been presented in ATLAS Weekly [14].

I have also had the opportunity to present my research at several national and international conferences, covering my ATLAS physics analysis, ITk hardware, and future lepton collider studies. These presentations include:

- Couplings and STXS measurements of Higgs boson decays to bosons by ATLAS and CMS, presented at the 10th Annual Large Hadron Collider Physics Conference [15];
- Prospects for strange tagging and measurements of Standard Model $H \rightarrow s\bar{s}$ at future lepton colliders, presented at the XXIX International Workshop on Deep-Inelastic Scattering and Related Subjects [16], the European Committee for Future Accelerators 1st Workshop of the Higgs/Top/EW Group [17], Higgs 2021 [18], and the International Workshop on Future Linear Colliders 2021 [19];
- Measurements of radiation effects in the ITk's front-end readout electronics from testbeams, presented at the 148th LHC Committee Meeting [20] and the Canadian Association of Physicists Congress 2021 [21];
- Per-channel and combination measurements of Higgs boson production by ATLAS and CMS, presented at Higgs 2020 [22];
- The use of machine learning for improved measurements of $VH, H \rightarrow WW^*$ in the 2ℓ DFOS channel with ATLAS, presented at the Canadian Association of Physicists Congress 2020 [23].

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Chapter 1

The Standard Model of Particle Physics and the Higgs Boson

High-energy particle physics studies the elementary particles which comprise the universe’s matter and the interactions of those particles with one another. It achieves this through collider experiments, where the byproducts of particle-particle collisions are measured, or observational apparatus, where particles originating from cosmic or terrestrial sources are measured. By understanding the universe at its smallest scales, high-energy particle physics provides insight into the universe at its largest scales including its origin and early conditions.

The present-day theoretical understanding of high-energy particle physics has its beginnings in the 20th century with the advent of quantum mechanics and special relativity. A holistic treatment of both fields, quantum field theory, led to the development of the Standard Model (SM) of particle physics, the theoretical framework describing the elementary particles of the universe and the forces by which they interact. The SM, in its current form, is the culmination of decades of theoretical and experimental advances and collaboration in the fields of quantum, nuclear, accelerator, and instrumentation physics.

In 2012, the existence of the Higgs boson, the particle which was long-postulated to give mass to elementary particles, was experimentally discovered by the ATLAS [24] and CMS [25] experiments at the Large Hadron Collider (LHC). In 2013, the Nobel Prize in Physics was awarded to Higgs and Englert for the theoretical foundation of the Higgs boson. The discovery of the Higgs boson represents a triumph of the SM and of the close collaboration of the theoretical and experimental high-energy physics communities.

Since its discovery, there has been a global effort to decipher the nature of the observed Higgs boson. This thesis is a contribution to that effort. The following sections provide an introduction to the physics of the SM (Section 1.1), the Higgs boson (Section 1.2), and the measurement of associated production of Higgs bosons using the ATLAS detector (Section 1.3). The measurement of associated production of Higgs bosons is the focus of this thesis.

1.1 The Standard Model of Particle Physics

The “particle zoo” of the SM consists of fermions and bosons. Fermions are half-integer-spin particles which comprise matter and include leptons, neutrinos, and quarks. They are further subdivided into three generations or “flavours” based on their mass. For each fermion, there is a corresponding “antifermion” or, more broadly,

antiparticle which has the same mass but opposite quantum numbers. Bosons, in contrast, are integer-spin particles which mediate the interactions between fermions and include photons, gluons, vector bosons, and Higgs bosons. The fundamental fermions and bosons of the SM are listed in Tables 1.1 and 1.2, respectively.

The SM Lagrangian, which encodes the possible interactions between fermions and bosons, is based on a product of local gauge symmetries, with each symmetry corresponding to set of conserved charges and forces. Quantum electrodynamics describes the interactions of particles carrying electric charge, which includes charged leptons, quarks, and charged vector bosons. Its corresponding force, the electromagnetic force, is mediated by the massless photon. Analogously, quantum chromodynamics describes the interactions of particles carrying colour charge, which includes quarks and gluons. Its corresponding force, the strong force, is mediated by the massless gluon. There are eight types of gluons carrying different combinations of colour-anticolour charge. The weak interaction describes the interactions of particles carrying weak isospin, the charge of the weak interaction’s underlying symmetry, or a combination of weak isospin and electric charge. Its corresponding force, the weak force, is mediated by the massive vector bosons, $W^\pm$ and Z . Whether a particle carries weak isospin depends on its chirality or “handedness”: *only* particles with left-handed chirality and antiparticles with right-handed chirality carry nonzero weak isospin. It follows that the strength of the weak interaction also depends on a particle’s handedness.

The Higgs boson, a massive excitation of a spin-0 field, arises due to spontaneous electroweak symmetry breaking and couples to massive fermions and bosons. The Higgs boson introduces mass terms for fermions and the vector bosons into the SM Lagrangian in way which preserves gauge invariance, completing the modern-day picture of particle physics.

Table 1.1: Fermions included in the Standard Model. All of the fermions have a spin of $1/2$. Quoted masses are taken from the Particle Data Group [26]. The weak isospins in the table only correspond to fermions with left-handed chirality and antifermions with right-handed chirality – all other fermions have weak isospins of 0.

Generation	Particle	Mass [MeV]	Electric charge	Weak isospin
First	Electron e	0.511	-1	$-1/2$
	Electron neutrino ν_e	< 0.00046	0	$+1/2$
	Down quark d	4.67	$-1/3$	$-1/2$
	Up quark u	2.16	$+2/3$	$+1/2$
Second	Muon μ	106	-1	$-1/2$
	Muon neutrino ν_μ	< 0.19	0	$+1/2$
	Strange quark s	93	$-1/3$	$-1/2$
	Charm quark c	1270	$+2/3$	$+1/2$
Third	Tau τ	1777	-1	$-1/2$
	Tau neutrino ν_τ	< 18.2	0	$+1/2$
	Bottom quark b	4180	$-1/3$	$-1/2$
	Top quark t	172760	$+2/3$	$+1/2$

The SM has been highly successful in its physics predictions. In addition to predicting the existence of the Higgs boson, the SM accommodates mixing between all three generations of quarks via the weak interaction, and the strength of a particular interaction – within or between generations – is dictated by the Cabibbo-Kobayashi-Maskawa (CKM) matrix [27, 28]. The CKM matrix is a 3×3 unitary matrix whose free parameters [29, 30] have since been constrained to $\lesssim \mathcal{O}(1\text{--}10\%)$ precision via independent and complementary measurements of nuclear and flavour physics [26], the latter corresponding to the study of the interactions of different flavours of particles.

Table 1.2: Bosons included in the Standard Model. Quoted masses are taken from the Particle Data Group [26].

Particle	Mass [GeV]	Electric charge	Spin	Weak isospin
Photon γ	0	0	1	0
Charged vector boson $W^\pm$	80.379	± 1	1	± 1
Neutral vector boson Z	91.1876	0	1	0
(8 $\times$) Gluon g	0	0	1	0
Higgs boson H	125.25	0	0	0

The SM also predicted the variation in the strength of the strong force as a function of the four-momentum transfer of the interaction, q^2 . Gross and Wilczek [31] and Politzer [32] demonstrated quantum chromodynamics to be “asymptotically free” whereby the strength of the strong force *decreases* with increasing q^2 , a prediction in excellent agreement with experiment [33]. Asymptotic freedom is what permits the treatment of quarks as quasi-free particles in high- q^2 interactions, an assumption which is paramount to physics predictions at the LHC.

Despite all of its successes, the SM is known to be an incomplete description of the universe. One unexplained phenomenon is the origin of neutrino masses: the SM assumes neutrinos are *massless*. However, the observation of neutrino oscillations, a variation in the probability of measuring a particular flavour of neutrino as it travels, by the Super-Kamiokande [34] and SNO [35, 36] experiments has demonstrated that neutrinos must be *massive*.

Another phenomenon unaccounted for by the SM is the universe’s matter-antimatter asymmetry. In the early universe, it was expected that matter and antimatter should have been created equally by the process of baryogenesis [37, 38]. However, there is a clear abundance of matter over antimatter in the universe today, and so there is broad interest in understanding the ways in which matter and antimatter differ. The ALPHA experiment [39, 40] at CERN is one such experiment probing the matter-antimatter differences and has so far demonstrated consistency between matter and antimatter for various quantum mechanical transitions [41–43], as predicted by charge-parity-time symmetry.

Galactic observations, including translational velocities of clusters of galaxies, rotational velocities of individual galaxies, and gravitational lensing, indicate the presence of additional mass or “dark matter” in the universe [44, 45]. However, there is no conclusive theory explaining the existence of dark matter or how it interacts with the particles of the SM.

Lastly, it is theoretically appealing to construct a model of the universe where the gravitational force is united with the existing forces of the SM. A massless spin-2 boson, the “graviton”, is postulated to be the mediator of the gravitational force [46]. However, there is no experimental evidence supporting the existence of the graviton or, as for dark matter, how it interacts with the particles of the SM.

1.1.1 Quantum Field Theory

Quantum field theory (QFT) is the framework used to describe the physics of multiparticle, relativistic quantum mechanics. The usual Hilbert space occupied by the single-particle basis states is elevated to a Fock space, which includes 2-, 3-, $\dots$, and n -particle basis states as well.

The quantum field $\phi \equiv \phi(x)$ is a local operator which acts upon the vacuum at spacetime coordinate x , creating and annihilating particles in a manner analogous to the creation and annihilation operators of the simple harmonic oscillator in non-relativistic quantum mechanics. And analogous to classical mechanics, the

behaviour of the system is that which minimizes the action:

$$S = \int \mathcal{L}(\phi, \partial_\mu \phi) d^4x, \quad (1.1)$$

where $\mathcal{L}$ is the Lagrange density, integrated over a 4-dimensional spacetime volume, and the derivative ∂_μ is defined as:

$$\partial_\mu \equiv (\partial/\partial x^0, \partial/\partial x^1, \partial/\partial x^2, \partial/\partial x^3) = (\partial_t, \vec{\nabla}). \quad (1.2)$$

By varying the action and requiring $\delta S = 0$ with fixed end points, one finds the Euler-Lagrange equation governing the behaviour of the field:

$$\partial_\mu \left(\frac{\partial \mathcal{L}}{\partial(\partial_\mu \phi)} \right) - \frac{\partial \mathcal{L}}{\partial \phi} = \partial_\mu \Pi^\mu - \frac{\partial \mathcal{L}}{\partial \phi} = 0. \quad (1.3)$$

Here, $\Pi \equiv \Pi^0$ is referred to as the “momentum conjugate to the field” and is similarly promoted to a local operator as part of the canonical quantization of the theory. Analogous to quantization in non-relativistic quantum mechanics, ϕ and Π are required to obey the equal-time commutation relations:

$$\begin{aligned} [\phi(t, \vec{x}), \phi(t, \vec{y})] &= [\Pi(t, \vec{x}), \Pi(t, \vec{y})] = 0, \\ [\phi(t, \vec{x}), \Pi(t, \vec{y})] &= i\delta^3(\vec{x} - \vec{y}), \end{aligned} \quad (1.4)$$

where δ^3 is the Dirac delta function in 3-dimensions. Bosonic fields, which have bosons as their quanta, obey a set of relations similar to Eq. 1.4. In contrast, fermionic fields, which have fermions as their quanta, obey equal-time *anti*-commutation relations.

Symmetry considerations are paramount to constructing QFTs. Noether’s theorem states that symmetries of a system lead to conserved quantities [47], and this is no different for quantum fields. In particular, one can consider a transformation $\phi(x) \rightarrow \phi(x, \lambda)$ with $\phi(x, 0) \equiv \phi(x)$, where λ is some parameter of the transformation. Such a transformation is a symmetry if and only if:

$$D\mathcal{L}(\phi, \partial_\mu \phi) = \partial_\mu F(\phi, \partial_t \phi, t) \text{ where } D \equiv \left. \frac{\partial}{\partial \lambda} \right|_{\lambda=0}, \quad (1.5)$$

where F is some function. When imposed on the Lagrangian, Eq. 1.5 leads to a conserved current:

$$J^\mu = \Pi^\mu D\phi - F^\mu. \quad (1.6)$$

Examples of symmetries include spacetime translations or Lorentz transformations, which collectively define the Poincaré group. As will be discussed in the following sections, imposing symmetries on the SM Lagrangian results in conserved quantities like electric charge, weak isospin, and colour charge.

In QFT, it is common practice to assume natural units:

$$\hbar = c = 1, \quad (1.7)$$

where $\hbar$ is the reduced Planck’s constant and c is the speed of light. This convention will be adopted in all subsequent sections and chapters. The choice of natural units implies energy, momentum, and mass to have units of electron-volts (eV), length and time to have units of eV^{-1} , and angular momentum to be dimensionless.

1.1.2 Quantum Electrodynamics

Quantum electrodynamics (QED) describes the interactions of photons with electrically-charged matter. A photon is a spin-1 massless boson with a corresponding vector field A_μ , whose free-space behaviour is governed by the Maxwell Lagrangian:

$$\mathcal{L}_{\text{EM}} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu}, \quad (1.8)$$

which consists of only a kinematic term, where $F_{\mu\nu}$ is the electromagnetic field tensor, $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$. Electrically-charged matter is described by a spin-1/2 massive fermion field ψ , whose free-space behaviour is governed by the Dirac Lagrangian [48]:

$$\mathcal{L}_D = \bar{\psi}(i\partial - m)\psi, \quad (1.9)$$

which consists of the kinematic term and the mass term, where m is the mass of the fermion. Additionally, $\bar{\psi} \equiv \psi^\dagger \gamma^0$ is the Dirac adjoint and $\partial \equiv \partial_\mu \gamma^\mu$. The famous “gamma” matrices, γ^μ , obey the anticommutation relation $\{\gamma^\mu, \gamma^\nu\} = 2g^{\mu\nu}$, where $g^{\mu\nu}$ is the Minkowski metric.

It is worth noting that Eq. 1.9 has both positive and negative energy solutions moving forward in time. The latter can also be understood as positive energy solutions moving backwards in time [49]. These negative energy solutions correspond to antiparticles.

The QED Lagrangian is the sum of Eqs. 1.8 and 1.9. An interaction term between the fields is introduced via the “minimal coupling” prescription: the derivative ∂_μ in Eq. 1.9 is replaced with the gauge covariant derivative $D_\mu = \partial_\mu + ieQA_\mu$, where e is the coupling strength of the interaction and Q is the electric charge of the fermion. Such a replacement yields:

$$\begin{aligned} \mathcal{L}_{\text{QED}} &= \bar{\psi}(i\partial - m)\psi - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} \\ &= \bar{\psi}(i\partial - m)\psi - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} - e\bar{\psi}\not{A}\psi, \end{aligned} \quad (1.10)$$

which is invariant under a U(1) local gauge symmetry:

$$\begin{aligned} \psi(x) &\rightarrow \exp(i\lambda(x))\psi(x), \\ A_\mu(x) &\rightarrow A_\mu(x) - \frac{1}{e}\partial_\mu\lambda(x), \end{aligned} \quad (1.11)$$

where λ is a function of spacetime position x . If the Lagrangian for the vector field, Eq. 1.8, included a mass term, $+\frac{\mu^2}{2}A_\mu A^\mu$, then Eq. 1.10 would *not* be invariant under a U(1) local gauge symmetry. It is this symmetry of the QED Lagrangian which leads to the conservation of electric charge, as per Noether’s theorem. The coupling strength in Eq. 1.10 is therefore identified as the fundamental unit of electric charge, $e = 1.602 \times 10^{-16}$ C.

Perturbation Theory: QED is regarded as an “interacting” theory, whereby the minimal coupling procedure introduces an interaction term connecting the free field theories for ψ and A_μ . With interacting theories, it is typically impossible to exactly calculate observables of interest. Approximation methods must be used, the most popular being perturbation theory. The underlying assumption of the use of perturbation theory is that

the coupling strength of the interaction is “small”, allowing calculations to be expanded to some fixed order in the coupling strength. In QED, for example, calculations are expanded in powers of the fine structure constant:

$$\alpha(q^2) \equiv \frac{e^2(q^2)}{4\pi} \stackrel{q^2 \approx 0}{\approx} \frac{1}{137} \ll 1. \quad (1.12)$$

The power of 2 on the coupling strength e naturally arises when calculating the matrix element $\mathcal{M}$ for particular observable O : $|\mathcal{M}|^2 \equiv |\langle f|O|i \rangle|^2$, where i and f are the initial and final states, respectively.

1.1.3 Quantum Chromodynamics

Quantum chromodynamics (QCD) describes the interactions of coloured particles, quarks and gluons. Quarks, coined by Gell-Mann [50], carry one unit of colour charge – red (r), green (g), or blue (b) – and antiquarks carry one unit of anticolour charge. As such, quarks are represented by colour triplets:

$$q_j = \begin{pmatrix} q_j^r \\ q_j^g \\ q_j^b \end{pmatrix} \forall j = 1, \dots, 6, \quad (1.13)$$

where the index j runs over the six flavours of quarks. Gluons, in contrast, carry one unit of colour charge and one unit of anticolour charge and are represented by colour octets. The colour confinement property of QCD dictates that coloured particles are only found in composite particles called hadrons. Hadrons correspond to colour-singlet states and can consist of two, in the case of mesons, or three, in the case of baryons, quarks. Recently, measurements of five-quark resonances known as pentaquarks have been reported [51, 52].

QCD differs from QED in that its underlying symmetry is an $SU(3)$ local gauge symmetry, as first proposed by Greenberg [53] and Han and Nambu [54]. The $SU(3)$ group is a *non*-Abelian group: its generators, $\vec{T}$, do *not* commute. As such, the quark and gluon fields transform as:

$$\begin{aligned} q_j(x) &\rightarrow \exp(i\vec{T} \cdot \vec{\theta}(x)) q_j(x) \quad \forall j = 1, \dots, 6, \\ G_\mu^a(x) &\rightarrow G_\mu^a(x) - \frac{1}{g_S} \partial_\mu \theta^a(x) - f_{abc} \theta^b(x) G_\mu^c(x) \quad \forall a = 1, \dots, 8. \end{aligned} \quad (1.14)$$

Here, G_μ^a are the gluon fields, $\vec{\theta}$ are functions of spacetime position x , g_S is the coupling strength of the $q_j - G_\mu^a$ interaction, and f_{abc} are the structure constants of $SU(3)$.¹ There are eight gluon fields as there are eight generators of $SU(3)$. These generators are typically written using the Gell-Mann matrices $\vec{\lambda}$: $\vec{T} = \vec{\lambda}/2$. The conserved quantities of the symmetry are the colour (C) charges, and so QCD is referred to as an $SU(3)_C$ local gauge symmetry.

Writing the Lagrangian for QCD follows very nearly the same procedure as for QED. Quarks are governed by the Dirac Lagrangian, Eq. 1.9, and gluons are governed by the Yang-Mills Lagrangian [55], leading to:

$$\mathcal{L}_{\text{QCD}} = \sum_{j=1}^6 \bar{q}_j (i\not{D} - m_j) q_j - \frac{1}{4} \sum_{a=1}^8 G_{\mu\nu}^a G^{a\mu\nu}, \quad (1.15)$$

where the $q_j - G_\mu^a$ interaction term is included via the minimal coupling prescription, $D_\mu = \partial_\mu + ig_S \sum_{a=1}^8 T^a G_\mu^a$, and $G_{\mu\nu}^a$ are the field tensors for the gluons, $G_{\mu\nu}^a = \partial_\mu G_\nu^a - \partial_\nu G_\mu^a - g_S f_{abc} G_\mu^b G_\nu^c$. The final term in the field

¹The structure constants f_{abc} for a group are defined by $[T^a, T^b] = if_{abc} T^c$, where T^a are the generators of the group.

tensor, $-g_s f_{abc} G_\mu^b G_\nu^c$, is due to the non-Abelian nature of QCD. One can see this term gives rise to 3- and 4-point gluon interaction terms not present in QED. As with QED, including a mass term for the gluons would break the local gauge symmetry, and so such a mass term is excluded.

Analogously to QED, predictions in QCD rely on perturbation theory, expanded in terms of the strong coupling constant:

$$\alpha_S(q^2) \equiv \frac{g_S^2(q^2)}{4\pi}. \quad (1.16)$$

The coupling constants for QED, α , and QCD, α_S , depend on the four-momentum transferred during interactions. For typical LHC proton-proton collisions, $\alpha_S \sim O(0.1)$.

1.1.4 The Weak Force and Electroweak Unification

The weak force describes the interactions of the spin-1 massive vector bosons, $W^\pm$ and Z , with the spin-1/2 fermions, leptons, neutrinos, and quarks. The boson-fermion interaction terms added to the SM Lagrangian take the form of:

$$\mathcal{L}_I \supset i\bar{\psi}\gamma^\mu(g_V - g_A\gamma_5)\psi V_\mu \text{ where } g_V \neq 0, g_A \neq 0, \quad (1.17)$$

where g_V and g_A are the “vector” and “axial” coupling strengths, respectively, and $V_\mu \in [W_\mu^\pm, Z_\mu]$. Additionally, $\gamma^5 \equiv -\frac{i}{4!}\epsilon^{\mu\nu\alpha\beta}\gamma_\mu\gamma_\nu\gamma_\alpha\gamma_\beta$ and $\epsilon^{\mu\nu\alpha\beta}$ is a fully antisymmetric tensor with $\epsilon^{0123} = 1$.

A key difference between Eq. 1.17 and the interaction term in Eq. 1.10 is the presence of nonzero axial couplings, which indicates the weak interaction to be parity-violating. In particular, interactions involving Z and $W^\pm$ bosons have $|g_V| < |g_A|$ and $g_V = g_A$, respectively. The latter corresponds to an interaction containing a left-handed chiral projection operator, $2ig_V\bar{\psi}\gamma^\mu P_L\psi W_\mu^\pm$, indicating that $W^\pm$ bosons couple only to left-handed (LH) chiral particles and right-handed (RH) chiral antiparticles. In other words, the weak charged current is *maximally* parity-violating.

In the 1960s, Glashow [56, 57], Weinberg [58], and Salam [59] postulated the existence of four vector fields, B_μ and $W_\mu^a \forall a \in [1, 2, 3]$, which interact with the fermions of the SM and linearly combine to form the observed bosons of QED and the weak force, a theory known as electroweak (EW) unification. The symmetry group of the Lagrangian governing the vector field B_μ is a $U(1)$ local gauge symmetry, where the fields transform in the usual way:

$$\begin{aligned} \psi(x) &\rightarrow \exp\left(i\frac{Y}{2}\lambda(x)\right)\psi(x), \\ B_\mu(x) &\rightarrow B_\mu(x) - \frac{1}{g'}\partial_\mu\lambda(x), \end{aligned} \quad (1.18)$$

where g' is the coupling strength of $\psi - B_\mu$ interaction and Y is the weak hypercharge, the conserved quantity of the symmetry. As such, the symmetry is labelled as $U(1)_Y$.

The symmetry group of the Lagrangian governing of the vector fields $W_{\mu\nu}^a$ is an $SU(2)$ local gauge symmetry, where the fields transform as:

$$\begin{aligned}\psi(x) &\rightarrow \exp(i\vec{T} \cdot \vec{\theta}(x))\psi(x), \\ W_\mu^a(x) &\rightarrow W_\mu^a(x) - \frac{1}{g_W} \partial_\mu \theta^a(x) - f_{abc} \theta^b(x) W_\mu^c(x) \quad \forall a = 1, \dots, 3\end{aligned}\quad (1.19)$$

where g_W is the coupling strength of the $\psi - W_\mu^a$ interaction, $\vec{T}$ are the generators of SU(2), and $\vec{\theta}$ are functions of spacetime position x . There are three components of the vector field as there are three generators of SU(2). These generators are typically written using the Pauli spin matrices $\vec{\sigma}$: $\vec{T} = \vec{\sigma}/2$. Each generator has an associated charge or weak isospin, T^1 , T^2 , and T^3 ; however, only the third component of the weak isospin T^3 and the total weak isospin T are conserved, the former referred to simply as the weak isospin. With respect to this last point: *only* LH chiral particles and RH chiral antiparticles transform under this symmetry, which are grouped according to generation into weak isospin doublets Ψ_L :

$$\Psi_L \in \left[\begin{pmatrix} \nu_e \\ e \end{pmatrix}_L, \begin{pmatrix} \nu_\mu \\ \mu \end{pmatrix}_L, \begin{pmatrix} \nu_\tau \\ \tau \end{pmatrix}_L, \begin{pmatrix} u \\ d \end{pmatrix}_L, \begin{pmatrix} c \\ s \end{pmatrix}_L, \begin{pmatrix} t \\ b \end{pmatrix}_L \right], \quad (1.20)$$

of total weak isospin $T = 1/2$ with a third component of weak isospin $T^3(\nu_e) = T^3(u) = \dots = +1/2$ and $T^3(e) = T^3(d) = \dots = -1/2$. As such, the symmetry is labelled as SU(2)_L. RH chiral particles, alternatively, are placed into weak isospin singlets ψ_R :

$$\psi_R \in [e_R, \mu_R, \tau_R, u_R, d_R, c_R, s_R, t_R, b_R]. \quad (1.21)$$

To decouple these states from the SU(2) symmetry, they are assigned $T(\psi_R) = T^3(\psi_R) = 0$.

Following the minimal coupling prescription, the Lagrangian governing the electroweak theory for a spin-1/2 fermion ψ is:

$$\mathcal{L}_{EW} = (\bar{\Psi}_L + \bar{\psi}_R)(i\not{D} - m)(\Psi_L + \psi_R) - \frac{1}{4} B_{\mu\nu} B^{\mu\nu} - \frac{1}{4} \sum_{a=1}^3 W_{\mu\nu}^a W^{a,\mu\nu}, \quad (1.22)$$

and:

$$D_\mu = \partial_\mu + ig' \frac{Y}{2} B_\mu + ig_W \sum_{a=1}^3 \frac{\sigma^a}{2} W_\mu^a, \quad (1.23)$$

is the covariant derivative. The vector fields B_μ^a and W_μ^a are *massless* in order to preserve gauge invariance. The terms for the $\Psi_L - W_\mu^a$ interaction take the form:

$$\mathcal{L}_{EW} \supset \frac{g_W}{2} \sum_{a=1}^3 \bar{\Psi}_L \sigma^a W^a \Psi_L = \frac{g_W}{\sqrt{2}} \bar{\Psi}_L \left(\sigma^+ W^+ + \sigma^- W^- + \frac{1}{\sqrt{2}} \sigma^3 W^3 \right) \Psi_L, \quad (1.24)$$

where the following definitions have been used:

$$\begin{aligned}W_\mu^\pm &\equiv \frac{1}{\sqrt{2}} (W_\mu^1 \mp iW_\mu^2), \\ \sigma^\pm &\equiv \frac{1}{2} (\sigma^1 \pm i\sigma^2).\end{aligned}\quad (1.25)$$

The first two terms of Eq. 1.24 lead to interaction terms proportional to $\bar{\nu}_{e,L} W^+ e_L$ and $\bar{e}_L W^- \nu_{e,L}$, respectively, using electrons and electron neutrinos as an example. Accordingly, $W_\mu^\pm$ are identified as fields for the SM $W^\pm$ bosons, albeit *without* mass.

In an analogous way, one can define the vector fields A_μ and Z_μ , which are linear combinations of the vector fields B_μ and W_μ^3 :

$$\begin{aligned} A_\mu &\equiv +B_\mu \cos \theta_W + W_\mu^3 \sin \theta_W, \\ Z_\mu &\equiv -B_\mu \sin \theta_W + W_\mu^3 \cos \theta_W, \end{aligned} \quad (1.26)$$

where θ_W is a parameter known as the weak mixing angle. These new fields have nonzero couplings to both LH and RH chiral particles. The vector field A_μ is identified as the field for the SM photon. By fixing the couplings of this field to those expected from QED, one finds the following relationship between weak hypercharge, electric charge, and weak isospin: $Y = 2(Q - T^3)$. One also finds following relationship between the coupling strengths of the various interactions: $e = g_W \sin \theta_W = g' \cos \theta_W$. Accordingly, Z_μ is identified as the field for the SM Z boson, albeit *without* mass. The coupling strength of the $\psi - Z_\mu$ interaction, $g_Z = g_W / \cos \theta_W$, is entirely determined by e and g_W , and the electroweak theory is unified.

It is worth noting that the quarks' mass eigenstates, the eigenstates of the free-field Lagrangian, are different from the flavour eigenstates, the eigenstates of the interaction Lagrangian and what are observed in nature. In particular, the flavour eigenstates are obtained via a unitary transformation of the mass eigenstates. The matrix representing this unitary transformation is referred to as the Cabibbo-Kobayashi-Maskawa (CKM) matrix [27, 28]. This permits interactions between quarks of different generations via the weak force, albeit suppressed by the magnitude of the off-diagonal elements of the CKM matrix. For simplicity, there will be no further distinction between mass and flavour eigenstates in subsequent sections and chapters.

1.2 The Standard Model Higgs Boson

In Section 1.1.4, the electroweak Lagrangian, Eq. 1.17, has a fermion mass term $-m(\bar{\Psi}_L \psi_R + \bar{\psi}_R \Psi_L)$. However, Ψ_L transforms under the $SU(2)_L$ gauge symmetry and ψ_R does not, and so the mass term overall is *not* invariant under the $SU(2)_L$ gauge symmetry and necessitates $m = 0$. It was for the same reason that mass terms were not included for the gauge bosons of the weak force, $W^\pm$ and Z . However, such a conclusion is clearly incorrect, as both fermions and weak gauge bosons are experimentally known to have nonzero masses.

The Higgs mechanism, the spontaneous breaking of the local gauge symmetry of the electroweak sector induced by the Higgs field, provides a theoretical resolution to this inconsistency by generating mass terms for fermions and weak gauge bosons. It also generates interaction terms coupling the Higgs field to massive particles known as Yukawa couplings.

1.2.1 Spontaneous Symmetry Breaking

Spontaneous symmetry breaking refers the process by which a Lagrangian's symmetry becomes "broken" by way of the system choosing to occupy a particular vacuum state out of two or more degenerate vacuum states which are not invariant under the original symmetry.

As an example, consider a simple Lagrangian for a scalar field ϕ :

$$\mathcal{L} = \frac{1}{2}(\partial_\mu \phi)(\partial^\mu \phi) - V(\phi) \text{ where } V(\phi) = \frac{1}{2}\mu^2 \phi^2 + \frac{1}{4}\lambda \phi^4, \quad (1.27)$$

where the mass of the field is $\sqrt{\mu^2}$ and $\lambda > 0$ corresponds to the strength of the quartic self-interaction. The Lagrangian is invariant under the global $\mathbb{Z}_2$ symmetry $\phi \rightarrow -\phi$.

The potential from Eq. 1.27 is plotted in Fig. 1.1 for both $\mu^2 > 0$ and $\mu^2 < 0$. In the former case, the vacuum state occurs at $\phi = 0$. In the latter case, the vacuum states occur at a nonzero value of the field, $\phi = \pm v = \pm \sqrt{-\mu^2/\lambda}$.

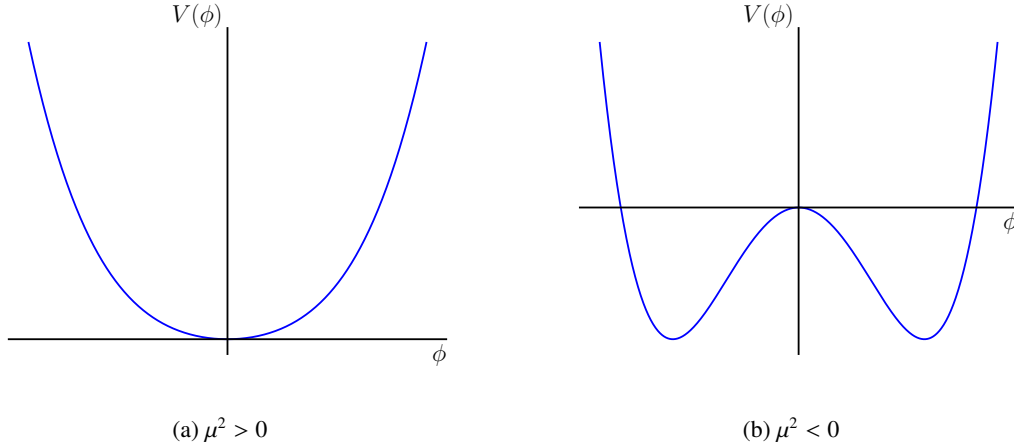


Figure 1.1: Plots of the potential from Eq. 1.27 for different signs of μ^2 .

Nature requires that the system is in a vacuum state. For the case of $\mu^2 < 0$, an arbitrary choice must be made between $+v$ and $-v$. Choosing $+v$ for definiteness, the field may be expanded about its vacuum state: $\phi = \phi' + v$. Making this replacement in Eq. 1.27 and dropping constant terms recovers the original Lagrangian but with an additional term proportional to ϕ'^3 . This triplet interaction term is *not* invariant under the global $\mathbb{Z}_2$ symmetry of the original Lagrangian, and so the symmetry has been broken.

1.2.2 The Higgs Mechanism

The underlying theory of the Higgs mechanism was developed by Englert and Brout [60]; Guralnik, Hagen, and Kibble [61]; and Higgs [62].² Embedded in the $SU(2)_L \otimes U(1)_Y$ local gauge symmetry of the electroweak sector, the model introduces two complex scalar fields which are placed into a weak isopin doublet Φ :

$$\Phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix} = \frac{1}{\sqrt{2}} \begin{pmatrix} \phi_1 + i\phi_2 \\ \phi_3 + i\phi_4 \end{pmatrix}, \quad (1.28)$$

and as such transforms under the $SU(2)_L$ local gauge symmetry, $\Phi(x) \rightarrow \exp(i\vec{T} \cdot \vec{\theta}(x))\Phi(x)$. Here, ϕ^+ is electrically-charged and ϕ^0 is electrically-neutral. The corresponding Lagrangian is:

$$\mathcal{L}_H = (D_\mu \Phi)^\dagger (D^\mu \Phi) - V(\Phi) \text{ where } V(\Phi) = \mu^2 \Phi^\dagger \Phi + \lambda (\Phi^\dagger \Phi)^2, \quad (1.29)$$

where the vector fields B_μ and W_μ^a are minimally coupled to scalar fields via the covariant derivative D_μ , given

²The discussion in this section closely follows that of Chapter 17 in Thomson [63].

by Eq. 1.23. For $\mu^2 < 0$, there are infinitely many vacuum states, $\Phi^\dagger \Phi = v^2/2 = -\mu^2/(2\lambda)$. The fields may be expanded about one of these vacuum states as, breaking the local gauge symmetry:

$$\Phi = \frac{1}{\sqrt{2}} \begin{pmatrix} \phi_1 + i\phi_2 \\ (h + v) + i\phi_4 \end{pmatrix}, \quad (1.30)$$

where $\phi_3(x) = h(x) + v$ and $h(x)$ is a real scalar field and v is a real constant. When Eq. 1.30 is substituted into Eq. 1.29, the fields ϕ_1 , ϕ_2 , and ϕ_4 are now found to be massless. The bosons corresponding to these fields are referred to as Goldstone bosons [64, 65]. The same substitution also demonstrates that the Goldstone bosons have unphysical couplings to the gauge bosons of electroweak sector; however, the Goldstone bosons may be removed by an $SU(2)_L$ local gauge transformation with a judicious choice of the gauge-fixing functions $\tilde{\theta}(x)$. This particular choice is known as the unitary gauge, and in the unitary gauge, Eq. 1.30 becomes:

$$\Phi = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ h + v \end{pmatrix}. \quad (1.31)$$

Boson Masses and Couplings: The minimally-coupled kinematic term of the scalar field's Lagrangian, Eq. 1.29, may be expanded using the scalar field in the unitary gauge:

$$\begin{aligned} (D_\mu \Phi)(D^\mu \Phi) &= \frac{1}{2}(\partial_\mu h)(\partial^\mu h) + \frac{1}{8}g_W^2(W_\mu^1 + iW_\mu^2)(W^{1,\mu} - iW^{2,\mu})(h + v)^2 \\ &\quad + \frac{1}{8}(g_W W_\mu^3 - g' B_\mu)(g_W W^{3,\mu} - g' B^\mu)(h + v)^2. \end{aligned} \quad (1.32)$$

The terms $\frac{1}{8}v^2 g_W^2(W_\mu^1 W^{1,\mu} + W_\mu^2 W^{2,\mu})$ are identified as mass terms for $W^{1/2}$, and therefore the $W^\pm$ bosons, with mass $m_W = \frac{1}{2}g_W v$. The mass terms for the W_μ^3 and B_μ fields are:

$$(D_\mu \Phi)(D^\mu \Phi) \supset \frac{1}{2} \begin{pmatrix} W_\mu^3 & B_\mu \end{pmatrix} \mathbf{M} \begin{pmatrix} W_\mu^3 \\ B_\mu \end{pmatrix} \text{ where } \mathbf{M} = \frac{1}{4}v^2 \begin{pmatrix} g_W^2 & -g_W g' \\ -g_W g' & g'^2 \end{pmatrix}, \quad (1.33)$$

where $\mathbf{M}$ is the mass matrix. This matrix can be readily diagonalized to yield the mass eigenstates A_μ and Z_μ with eigenvalues $m_A^2 = 0$ and $m_Z^2 = \frac{1}{4}v^2(g_W^2 + g'^2)$, respectively. If the ratio of the $U(1)_Y$ -to- $SU(2)_L$ fermion-boson couplings, g'/g_W , is required to equal $\tan \theta_W$, then the mass eigenstates correspond exactly to those for the photon and Z boson fields, Eq. 1.26, with masses of m_A and m_Z , respectively. Thus, all of the gauge boson masses in the SM are accommodated by the Higgs mechanism.

Still working in the basis where the $W_\mu^3 - B_\mu$ mass matrix is diagonalized and using the definition of $W_\mu^\pm$, Eq. 1.25, the Higgs boson couples to massive gauge bosons with interaction terms:

$$(D_\mu \Phi)(D^\mu \Phi) \supset \frac{1}{2}g_V^2 v V_\mu V^\mu h + \frac{1}{4}g_V^2 V_\mu V^\mu h h, \quad (1.34)$$

where $V_\mu \in [W_\mu^\pm, Z_\mu]$. It can be shown that coupling strength of the Higgs boson to massive gauge bosons is proportional to the mass of the boson: $\frac{1}{2}g_V^2 v = g_V m_V$. There are no direct interaction terms coupling the Higgs boson to pairs of photons or gluons as photons and gluons are massless. Instead, the $h - A_\mu / h - G_\mu^a$ couplings are indirectly generated via a massive fermion loop, typically a top quark, or a boson loop.

Fermion Masses and Couplings: The fermion masses in the SM are accommodated by coupling the Higgs doublet to weak isospin doublets. The combination, $\bar{\Psi}_L \Phi$, is invariant under an $SU(2)_L$ local gauge symmetry.

Invariance under a $U(1)_Y$ local gauge symmetry is introduced by coupling this combination to a weak isospin singlet. As the Hermitian conjugate of the total combination is also invariant, the Higgs-fermion interaction terms must be of the form:

$$\mathcal{L}_{\text{Yukawa}} = -g_\psi(\bar{\Psi}_L\Phi\psi_R + \bar{\psi}_R\Phi^\dagger\Psi_L), \quad (1.35)$$

where g_ψ is referred to as the *Yukawa* coupling for ψ . Following spontaneous symmetry breaking, the Higgs doublet Φ may be written in the unitary gauge and Eq. 1.35 becomes:

$$\begin{aligned} \mathcal{L}_{\text{Yukawa}} &= -\frac{g_\psi}{\sqrt{2}}v(\bar{\psi}_L\psi_R + \bar{\psi}_R\psi_L) - \frac{g_\psi}{\sqrt{2}}h(\bar{\psi}_L\psi_R + \bar{\psi}_R\psi_L) \\ &= -m_\psi(\bar{\psi}_L\psi_R + \bar{\psi}_R\psi_L) - \frac{m_\psi}{v}h(\bar{\psi}_L\psi_R + \bar{\psi}_R\psi_L), \end{aligned} \quad (1.36)$$

where ψ_L is the *lower* spin-1/2 field of the corresponding weak isospin doublet and $m_\psi = g_\psi v / \sqrt{2}$ has been identified as the mass of the fermion. In Eq. 1.36, the second term is referred to as a Yukawa interaction. The corresponding coupling strength referred to as the Yukawa coupling.

To construct the mass terms for the upper spin-1/2 fields of the weak isospin doublets, the fermion fields are coupled to the charge conjugate of Φ :

$$\Phi_c = -i\sigma_2\Phi^* = \begin{pmatrix} -\phi^{0*} \\ \phi^- \end{pmatrix} = \frac{1}{\sqrt{2}}\begin{pmatrix} -\phi_3 + i\phi_4 \\ \phi_1 - i\phi_2 \end{pmatrix}, \quad (1.37)$$

via interaction terms of the form:

$$\mathcal{L}_{\text{Yukawa}} = g_\psi(\bar{\Psi}_L\Phi_c\psi_R + \bar{\psi}_R\Phi_c^\dagger\Psi_L). \quad (1.38)$$

Following spontaneous symmetry breaking and after writing Φ_c in the unitary gauge, mass terms and Yukawa interactions identical to Eq. 1.36 are obtained for the up, charm, and top quarks. Thus, through the Higgs mechanism, fermion masses are accounted for in the SM.

1.2.3 Higgs Observables and Properties

Using the complete dataset obtained by the Tevatron at the end of its operation in late 2011, the mass of a particle consistent the SM Higgs boson was limited to the range $115 < m_H < 140 \text{ GeV}$ [66], with a maximum local significance of 3σ at 120 GeV. In late 2012, a particle consistent with the SM Higgs boson was experimentally confirmed by the ATLAS [24] and CMS [25] experiments at the LHC, where a local significance of 5σ at 125–126 GeV was observed. Since that time, there has been a strong program of precision measurements of the Higgs boson's properties as well as searches for di-Higgs production [67–73] and studies of extended Higgs sector models [74–84].

Two of the most common observables reported by measurements in particle physics are production cross sections and branching ratios for decay – this is no different for measurements characterizing the Higgs boson. Both observables correspond to transitions between quantum mechanical states and are therefore understood using the relativistic formulation of Fermi's Golden rule [85, 86]. That is, the rate of transition is proportional to the Lorentz-invariant matrix elements integrated over the Lorentz-invariant phase space.

A production cross section describes the rate of two particles scattering to produce one or more daughter particles and is given by [63]:

$$\text{Cross section} \equiv \sigma = \frac{\text{number of interactions per unit time per target particle}}{\text{incident flux}}. \quad (1.39)$$

Cross sections are typically given in units of barns (b) where $1 \text{ b} \equiv 10^{-28} \text{ m}^2$.

A branching ratio describes the relative rate of a particle decaying into two or more daughter particles. The total rate of decay for a particle depends on its total transition rate or “width”, Γ , which is given by the sum of the partial widths Γ_i for all possible sets of final state particles or “channels”:

$$\frac{dN}{dt} = -N \sum_{i \in \text{channels}} \Gamma_i = -N \times \Gamma \text{ where } \Gamma \equiv \sum_{i \in \text{channels}} \Gamma_i. \quad (1.40)$$

For a particle species X , the branching ratio for the $X \rightarrow n$ channel, $\mathcal{B}_{X \rightarrow n}$, is defined as the partial width for that channel over the total width of X :

$$\text{Branching ratio for } X \rightarrow n \equiv \mathcal{B}_{X \rightarrow n} = \frac{\Gamma_{X \rightarrow n}}{\Gamma_X}. \quad (1.41)$$

Observables for the Higgs boson, such as its production cross sections and branching ratios, are functions of the properties of the Higgs boson, such as its mass and width. Measurements of the former may be used to constrain the latter and vice versa.

Mass: The most precise measurements of the mass of the Higgs boson come from the ATLAS³ [88] and CMS [89] experiments. The world average is found by the PDG to be $m_H = 125.25 \pm 0.17 \text{ GeV}$ [26], with the corresponding vacuum expectation value of the Higgs field $v \approx 246 \text{ GeV}$. However, most physics predictions which use m_H as an input parameter, including measurements of the production cross sections and branching ratios, use the value obtained from the LHC Run 1 combination measurement by ATLAS and CMS: $125.09 \pm 0.21(\text{stat.}) \pm 0.11(\text{syst.}) \text{ GeV}$ [90].

Width: The most precise measurements of the width of the Higgs boson come from the CMS experiment, which has reported $\Gamma_H = 3.2^{+2.4}_{-1.7} \text{ MeV}$ using measurements of on- and off-shell Higgs couplings in the $H \rightarrow ZZ^* \rightarrow 4\ell/2\ell 2\nu$ channels [91], in agreement with the SM expectation of 4.1 MeV . The ATLAS experiment previously performed a similar measurement using only the 2015–16 LHC Run 2 dataset [92], reporting a 95% upper limit of 14.4 MeV .

Cross Sections and Couplings: The tree-level Feynman diagrams for the main production modes of the Higgs boson in pp collisions at centre-of-mass energy $\sqrt{s} = 13 \text{ TeV}$ are shown in Fig. 1.2. In order of highest-to-lowest cross section, the main production modes are:

1. Gluon-gluon fusion (ggF or ggH), which corresponds to the coupling of gluons to a Higgs boson via a heavy fermion loop;
2. Vector boson fusion (VBF or qqH), which occurs when scattering quarks radiate vector bosons which then couple to a Higgs boson;

³A more precise measurement by ATLAS has since been performed using the $H \rightarrow ZZ^* \rightarrow 4\ell$ channel [87].

3. Production in association with a vector boson (VH), which occurs when a vector boson radiates a Higgs boson or couples to one via a heavy fermion loop;
4. Production in association with a top ($t\bar{t}H$) or bottom ($b\bar{b}H$) quark pair.

As the coupling strength of the Higgs boson to a particle is a monotonically increasing function of the particle's mass, the Higgs boson will preferentially decay into heavier particles. The leading decays include $H \rightarrow b\bar{b}$, $H \rightarrow WW^*$, $H \rightarrow \tau\tau$. However, measurements in these channels must contend with large hadronic backgrounds or final states which cannot be fully reconstructed. The $H \rightarrow ZZ^*$ and $H \rightarrow \gamma\gamma$ decays have considerably smaller branching ratios but final states with clean experimental signatures.

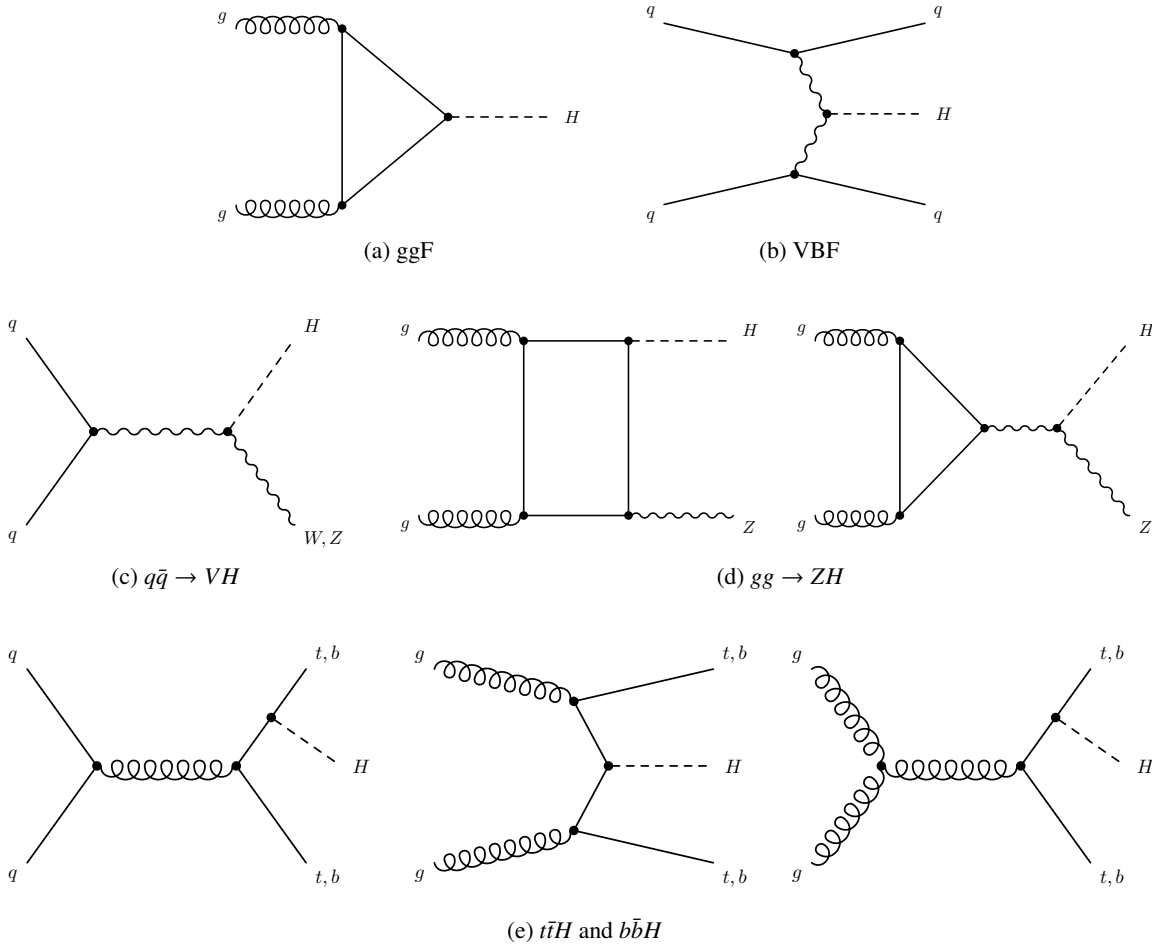


Figure 1.2: Tree-level Feynman diagrams for the main production modes of the Higgs boson in pp collisions at $\sqrt{s} = 13$ TeV [93].

State-of-the-art calculations are available for the production cross sections of the Higgs and its branching ratios [94–97]. Figs. 1.3a and 1.3b show the predicted Higgs boson production cross sections and branching ratios, respectively. Tables 1.3 and 1.4 list the cross sections and branching ratios, respectively, assuming fixed values of $m_H = 125.09$ GeV and $\sqrt{s} = 13$ TeV.

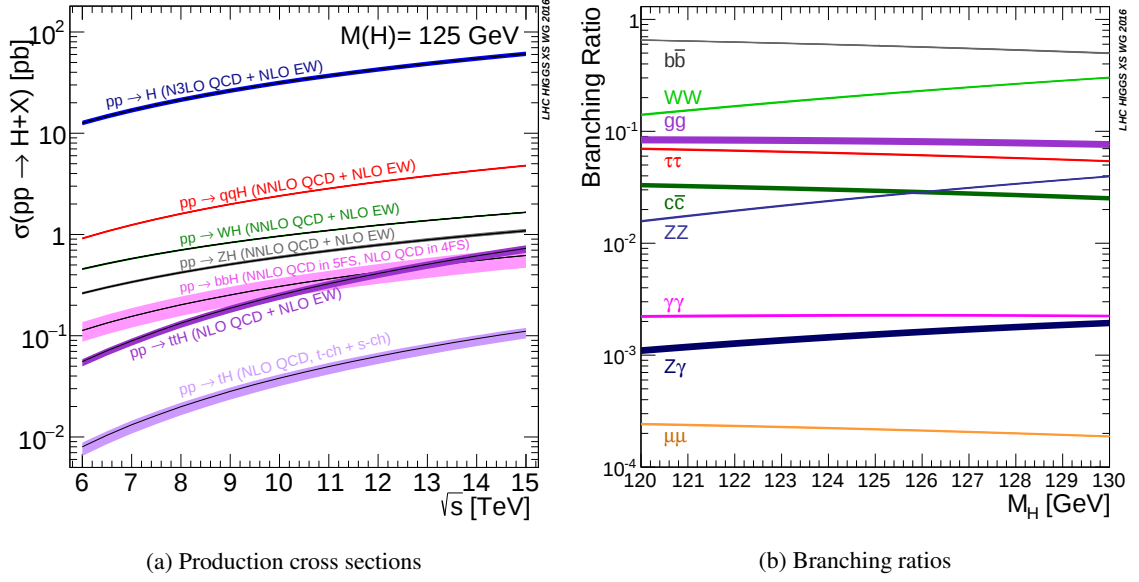


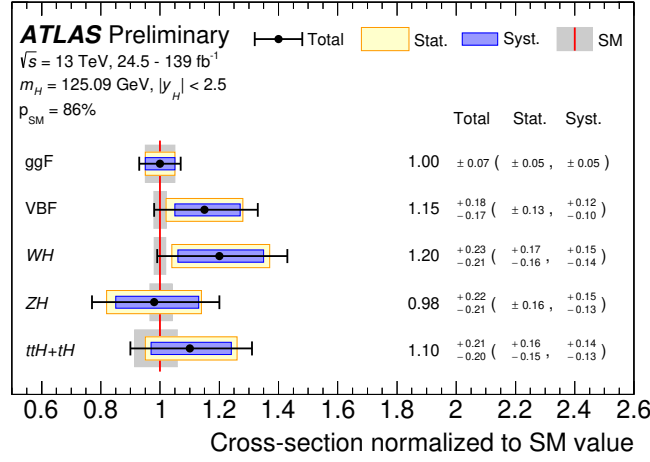
Figure 1.3: Higgs boson production cross sections in pp collisions as a function of centre-of-mass energy $\sqrt{s}$ and (b) Higgs boson branching ratios as a function of the mass of the Higgs boson m_H [97]. In both plots, the uncertainty bands correspond to a combination of theoretical uncertainties as well as uncertainties on the parameters entering the calculations.

Table 1.3: SM Higgs boson production cross sections in pp collisions assuming a Higgs boson mass $m_H = 125.09$ GeV and a centre-of-mass energy $\sqrt{s} = 13$ TeV [97]. The uncertainties correspond to theoretical uncertainties.

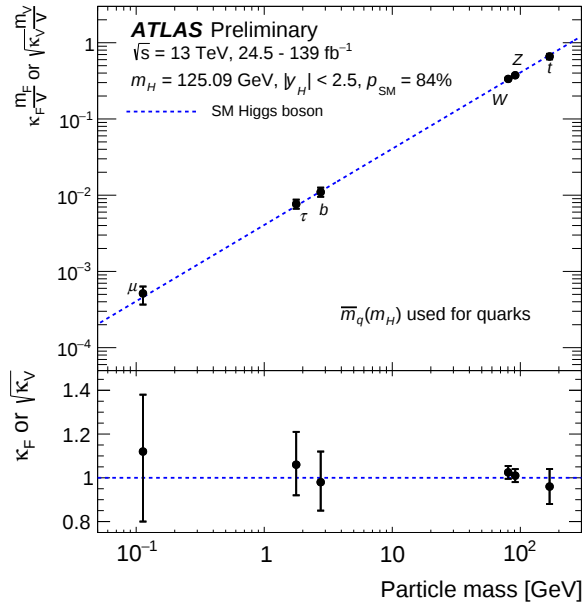
Production mode	Cross section [pb]	Uncertainty [%]
ggF	48.5	± 3.9
VBF	3.78	± 2.1
WH	1.37	± 2.0
ZH	0.882	+4.1 / -3.5
$t\bar{t}H$	0.507	+6.9 / -9.9
$b\bar{b}H$	0.486	+20.1 / -23.9

Cross sections measured by ATLAS in a recent combination measurement are shown in Fig. 1.4a, where excellent consistency with the SM is observed. Fig. 1.4b shows the Higgs-boson and Higgs-fermion coupling strengths as a function of the mass of the boson or fermion. For vector bosons as well as third generation fermions, there is excellent consistency with the SM.

With the Yukawa couplings to third generation fermions well established, ATLAS and CMS have begun studying the Yukawa couplings to second generation fermions, including $H \rightarrow \mu\bar{\mu}$ [98, 99] and $H \rightarrow c\bar{c}$ [100, 101]. The results obtained thus far are consistent with the SM but do not reach the same level of precision as measurements of $H \rightarrow \tau\bar{\tau}$ and $H \rightarrow b\bar{b}$. As the LHC heads into Run 3 and beyond, improving the sensitivity of these channels will be critical for testing the flavour universality of the Yukawa couplings – that is, that there are no additional corrections to the Yukawa couplings beyond their mass dependence.



(a) Production cross sections



(b) Yukawa couplings

Figure 1.4: (a) Measured cross sections for the production modes of the Higgs boson in pp collisions at $\sqrt{s} = 13 \text{ TeV}$ normalized to their SM expectations [102]. The yellow boxes include only the statistical uncertainty, the blue boxes include only the systematic uncertainty, and the black error bars include both statistical and systematic uncertainties (b) Measured coupling strengths for fermions, $\kappa_F(m_F/v) \forall F \in [t, b, \tau, \mu]$, and bosons, $\sqrt{\kappa_V}(m_V/v) \forall V \in [W, Z]$, as a function of their mass [102]. The black error bars indicate the 68% confidence intervals for the measured parameters.

Table 1.4: SM Higgs boson branching ratios assuming a Higgs boson mass of $m_H = 125.09$ GeV [97]. The uncertainties correspond to a combination of theoretical uncertainties as well as uncertainties on the parameters entering the calculations.

Decay channel	Branching ratio	Uncertainty [%]
$H \rightarrow b\bar{b}$	5.81×10^{-1}	± 1.7
$H \rightarrow WW^*$	2.15×10^{-1}	$+2.2 / -2.1$
$H \rightarrow gg$	8.18×10^{-2}	$+7.3 / -7.2$
$H \rightarrow \tau\bar{\tau}$	6.26×10^{-2}	± 2.3
$H \rightarrow c\bar{c}$	2.88×10^{-2}	$+6.6 / -2.8$
$H \rightarrow ZZ^*$	2.64×10^{-2}	$+2.2 / -2.1$
$H \rightarrow \gamma\gamma$	2.27×10^{-3}	$+2.9 / -2.8$
$H \rightarrow Z\gamma$	1.54×10^{-3}	$+6.8 / -6.9$
$H \rightarrow \mu\bar{\mu}$	2.17×10^{-4}	± 2.4

1.3 Associated Production of Higgs Bosons in the $H \rightarrow WW^*$ Decay Channel

Higgs boson production in association with a W or Z boson, which are respectively denoted by WH and ZH and collectively referred to as associated production, provides direct access to the Higgs boson’s couplings to vector bosons given knowledge of the other branching fractions of the Higgs. The values of these couplings are predicted by measurements of the electroweak coupling strengths and the vacuum expectation value of the Higgs field, and so measurements of associated production provide stringent tests of the SM.

In this thesis, a measurement of associated production of Higgs bosons in the $H \rightarrow WW^*$ decay channel is described. The measurement utilizes the full LHC Run 2 dataset, corresponding to an integrated luminosity of 139 fb^{-1} recorded by the ATLAS detector using proton-proton collisions at centre-of-mass energy $\sqrt{s} = 13 \text{ TeV}$.

1.3.1 Channel Overview

The $VH, H \rightarrow WW^*$ measurement is performed in several channels based on the decays of the W bosons from the Higgs and the associated vector boson V . The channels are split according to the number leptons in the final state as well as the charge configuration of those leptons. The channels are defined such that they are “orthogonal” to one another – events found in one channel will not be found in any other channel, ensuring independent measurements. The tree-level Feynman diagrams for the signal topologies of interest are shown in Fig. 1.5.

Within each channel, one or more signal regions enriched in $VH, H \rightarrow WW^*$ events are constructed. $W \rightarrow \tau\nu$ and $Z \rightarrow \tau\bar{\tau}$ decays with leptonically-decaying tau leptons for any of the vector bosons in the $VH, H \rightarrow WW^*$ topology are also included as signal; however, $VH, H \rightarrow \tau\bar{\tau}$ is treated as background. For some channels, control regions normalizing one or more of the dominant backgrounds to data are also defined.

To achieve the highest precision measurement possible, all channels use *multivariate discriminants* to extract the $VH, H \rightarrow WW^*$ signal within their respective signal regions. Multivariate discriminants are capable of efficiently utilizing correlations between input variables to isolate target processes, making them advantageous over kinematic selections alone. The multivariate discriminants in the signal regions as well as the event yields in the control regions are fitted in the combined measurement, and the WH/ZH production

cross sections times the $H \rightarrow WW^*$ branching ratio are extracted as the final result.

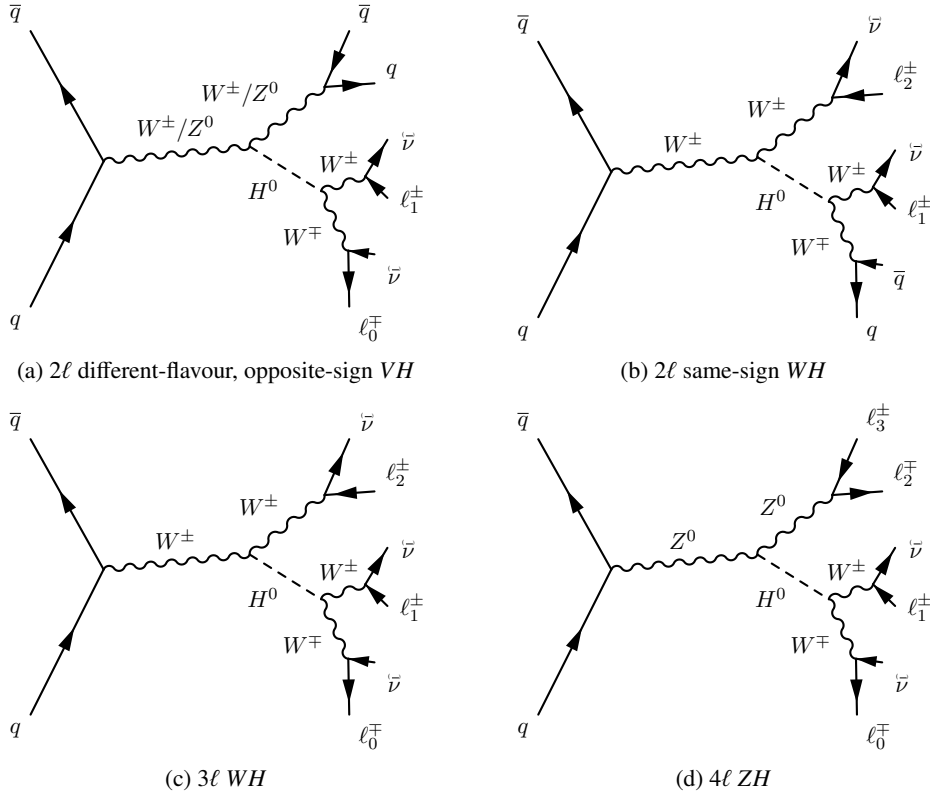


Figure 1.5: Tree-level Feynman diagrams for the signal topologies considered by the VH , $H \rightarrow WW^*$ analysis [103].

For final states with two leptons (2ℓ), there are two channels. The different-flavour, opposite-sign (DFOS) VH channel targets two DFOS leptons, $e\mu/\mu e$, from the decay of the Higgs and two or more jets from the decay of the associated vector boson. As there is no leptonic decay required of the vector boson, both WH and ZH are included as a signal. The channel contends with large dilepton backgrounds, including top-antitop pair production ($t\bar{t}$), Z boson production in association with jets (Z +jets), and diboson production ($q\bar{q} \rightarrow WW$). Control regions are used to normalize the top, Z +jets, and $q\bar{q} \rightarrow WW$ backgrounds to data.

The 2ℓ same-sign WH channel targets two same-sign leptons and one or more jets. The 2ℓ same-sign channel is divided into three subchannels based on the flavours of the leading leptons: $\mu\mu$, ee , and $e\mu/\mu e$. The largest backgrounds include trileptonic WZ and same-sign WW production. This channel also has considerable contamination from the misidentified lepton background. This background is estimated using a data-driven procedure common to the $2\ell/3\ell$ channels. For the ee and $e\mu/\mu e$ subchannels, the background attributed to electrons with misidentified charge is estimated using a data-driven procedure specific to the 2ℓ same-sign WH channel.

The 3ℓ channel is subdivided according to the number of same-flavour, opposite-sign (SFOS) lepton pairs in the final state: the Z -depleted channel requires zero SFOS lepton pairs and the Z -dominated channel requires one or more SFOS lepton pairs. Both channels are inclusive in the number of jets. There is greater background contamination in the Z -dominated channel than in the Z -depleted channel. In the case of Z -dominated, the background consists primarily of trileptonic WZ with much smaller contributions from ZZ and misidentified leptons. In the Z -depleted channel, trileptonic WWW , trileptonic WZ , and misidentified lepton backgrounds

contribute all roughly equally. Control regions normalizing the WZ background to data are shared between the 3ℓ channels.

The 4ℓ channel, following a similar approach to the 3ℓ channel, is subdivided according to the number of SFOS lepton pairs in the final state. The 1 SFOS (1SFOS) channel contends with backgrounds from four-lepton ZZ , $t\bar{t}Z$, and misidentified lepton processes, while the 2 SFOS (2SFOS) channel is dominated by ZZ production. Both channels are inclusive in the number of jets. A data-driven procedure is used to estimate the misidentified lepton background. This procedure is different from that used for the $2\ell/3\ell$ channels as the composition of the misidentified leptons is different. A control region normalizing the ZZ background to data is shared between the 4ℓ channels.

1.3.2 Existing Results

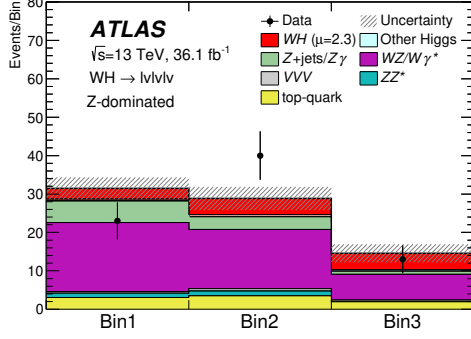
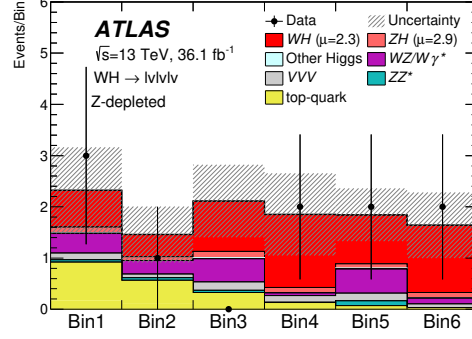
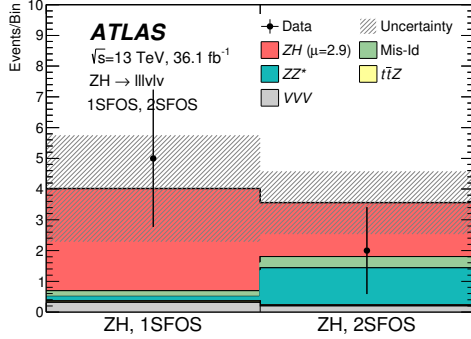
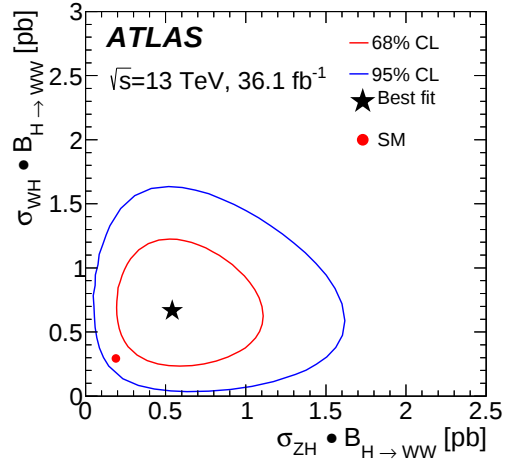
Previous iterations of the analysis were performed by ATLAS and CMS using the LHC’s Run 1 dataset at $\sqrt{s} = 7$ and 8 TeV [103, 104] as well as the 2015–16 Run 2 dataset at $\sqrt{s} = 13$ TeV [105, 106]. CMS has also published results using the full Run 2 dataset [107, 108].

In its 2015–16 analysis [105], ATLAS measured the WH/ZH production cross sections, $\sigma_{WH/ZH}$, times the $H \rightarrow WW^*$ branching ratio, $\mathcal{B}_{H \rightarrow WW^*}$, to be:

$$\begin{aligned}\sigma_{WH} \times \mathcal{B}_{H \rightarrow WW^*} &= 0.67^{+0.31}_{-0.27}(\text{stat.})^{+0.18}_{-0.14}(\text{sys.}) \text{ pb}, \\ \sigma_{ZH} \times \mathcal{B}_{H \rightarrow WW^*} &= 0.54^{+0.31}_{-0.24}(\text{stat.})^{+0.15}_{-0.07}(\text{sys.}) \text{ pb},\end{aligned}\tag{1.42}$$

However, this measurement only included the $3\ell/4\ell$ channels and utilized a much smaller fraction of the Run 2 integrated luminosity, 36.1 fb^{-1} . As these results are limited by the statistical uncertainty in data, the inclusion of the full Run 2 dataset will improve the precision significantly with an expected reduction in the statistical uncertainty by a factor of 2. The signal region distributions following the statistical fit (or “post-fit”) and the 2D likelihood contours for the measured cross sections from the 2015–16 analysis are shown in Fig. 1.6.

Given that VH production has been observed in the $H \rightarrow b\bar{b}$ decay channel [109–111] and $H \rightarrow WW^*$ has been observed from a combination measurement of ggF , VBF , and VH production [112], it is expected that VH , $H \rightarrow WW^*$ will be observed with sufficiently high statistics. However, there are still Beyond the SM (BSM) effects which could lead to a modification of its cross section and couplings to vector bosons. BSM particles – for example, additional Higgs bosons introduced via two-Higgs-doublet models [113, 114] or new heavy vector bosons [115, 116] – sharing the same signature as VH can lead to enhancements to its cross section [117, 118]. Pseudoscalar Higgs bosons can also mix with the SM Higgs boson and modify the charge-parity (CP) structure of its couplings to vector bosons [119–122]; however, CP-odd contributions are already significantly constrained by ATLAS [123, 124] and CMS [125, 126].

(a) 3ℓ Z-dominated WH BDT(b) 3ℓ Z-depleted WH BDT(c) 4ℓ 1SFOS and 2SFOS ZH yields

(d) 2D likelihood contours

Figure 1.6: The post-fit boosted decision trees (BDTs) for the 3ℓ (a) Z-dominated and (b) Z-depleted WH channels and (c) the post-fit yields for the 4ℓ 1SFOS and 2SFOS ZH channels from the VH , $H \rightarrow WW^*$ measurement utilizing the 2015–16 Run 2 dataset measured by ATLAS [105]. The error bands include both statistical and systematic uncertainties. (d) The 2D likelihood contours for the measured WH/ZH production cross sections times the $H \rightarrow WW^*$ branching ratio from the same analysis.

Chapter 2

The ATLAS Detector at the Large Hadron Collider

The ATLAS (A Toroidal LHC ApparatuS) detector [127] is one of four large detectors operating at CERN’s Large Hadron Collider (LHC) [128], the others being the CMS (Compact Muon Solenoid) [129], LHCb (LHC beauty) [130], and ALICE (A Large Ion Collider Experiment) [131] detectors. ATLAS and CMS are both general purpose detectors designed for high-luminosity proton-proton collisions and measurements of a diverse range of phenomena. LHCb studies flavour physics and specifically the interactions of b -hadrons, providing insight into the asymmetries of nature. ALICE uses heavy ion collisions to study quark-gluon plasma, whose conditions are similar to those in the earliest moments of the universe.

In this chapter, the design and operation of the LHC (Section 2.1) as well as the ATLAS detector (Section 2.2) – the experimental apparatus of the $VH, H \rightarrow WW^*$ measurement – are described.

2.1 The Large Hadron Collider

The Large Hadron Collider (LHC) [128, 132–134] is a particle-particle collider located beneath the French-Swiss border at the European Organization for Nuclear Research (CERN). The LHC is roughly circular in design, with a circumference of 26.7 km which is divided into eight arcs, each with a radius of curvature of 2.8 km, and eight straight sections, each with a length of 528 m. Of the eight straight sections, four of which are equipped with interaction points for particle-particle collisions and are instrumented by one of the LHC’s large detectors: ATLAS, CMS, LHCb, or ALICE. Two of these interaction points, located diametrically opposite one another along the LHC’s ring, support high-luminosity collisions and are instrumented by ATLAS and CMS. A diagram of the LHC accelerator complex, including the positions of the LHC’s major experiments, is shown in Fig. 2.1.

2.1.1 Design and Operation

The LHC was designed for proton-proton (pp) collisions with a beam energy of 7 TeV and a luminosity of $O(10^{34}) \text{ cm}^{-2}\text{s}^{-1}$ and heavy ion collisions with a beam energy of 2.8 TeV per nucleon and a peak luminosity $O(10^{27}) \text{ cm}^{-2}\text{s}^{-1}$. As a particle-particle collider, the LHC consists of two independent rings with counter-rotating beams. Each ring has its own magnetic field and vacuum chamber, with common sections only at the

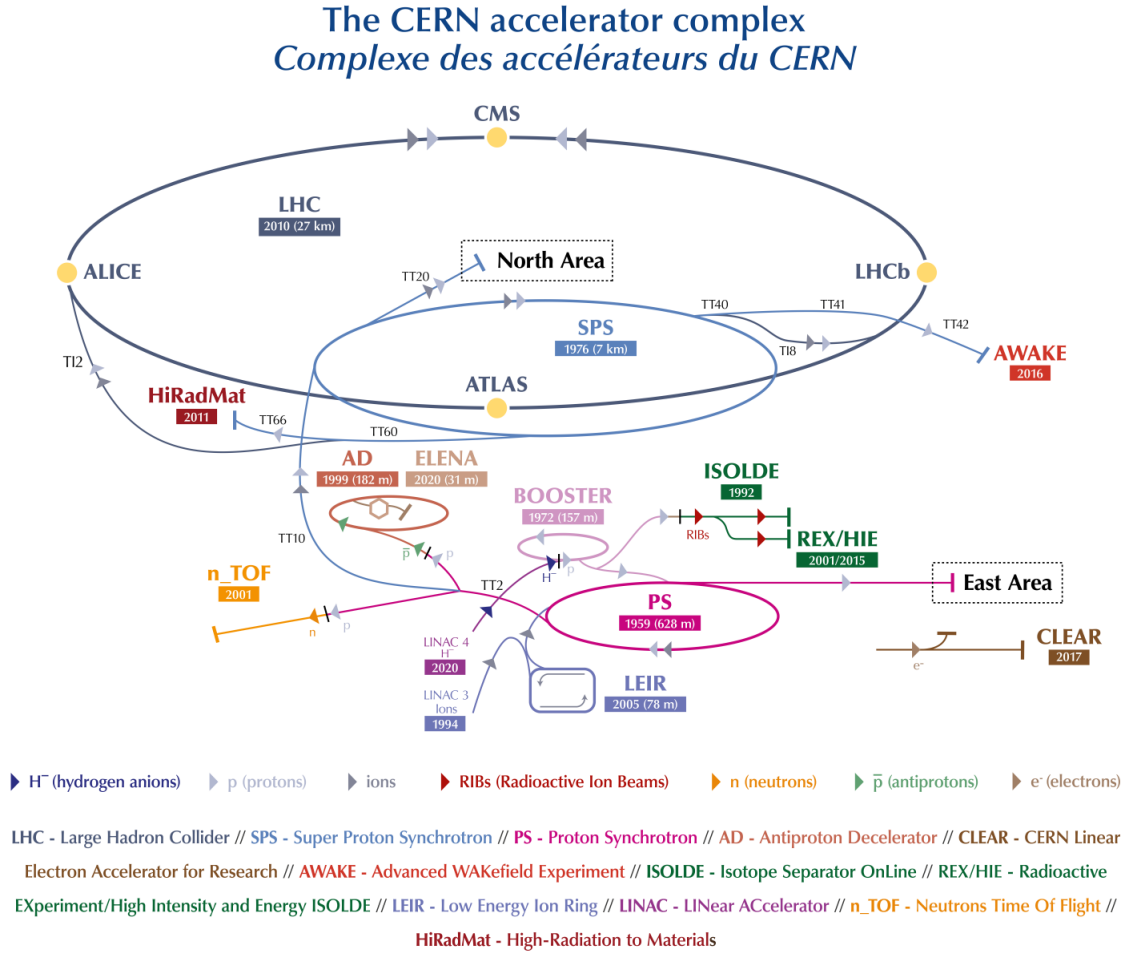


Figure 2.1: The LHC accelerator complex at CERN [135].

interaction points. Because there is insufficient room in the tunnel for two separate rings, the LHC utilizes twin-bore magnets [136]. Each magnet contains two sets of coils and beam channels, both sharing the same mechanical and cooling infrastructure. A cross section of a typical LHC magnet is shown in Fig. 2.2.

The LHC accelerator chain begins with hydrogen gas. Protons are stripped away from the gas and accelerated to 50 MeV via Linac2 [138]. Coming from Linac2, protons are bunched into synchronous groups or “bunches” by the Proton Synchrotron Booster (PSB) followed by the Proton Synchrotron (PS) [139]. The bunches are spaced 25 ns apart and accelerated to 25 GeV. Coming from the PS, protons are accelerated to 450 GeV by the Super Proton Synchrotron (SPS) [140] and injected into the LHC. The bunches have an intensity of 1.15×10^{11} protons/bunch at the injection, and the LHC is capable of supporting 2,808 bunches per ring during collisions.

Once in the LHC, the beam is accelerated using a 400 MHz (RF) superconducting cavity system, with eight multi-cavity cryomodules per beam [128, 141] and an energy gain of 485 keV per turn. Each cryomodule consists of four single-cell cavities operating at a field strength of 5.5 MV/m and sharing a common cryostat and klystron [142]. The cavities themselves are made of copper sputtered with superconducting niobium, which has the benefit of being less susceptible to quenching than solid niobium. The cryostats maintain the cavities

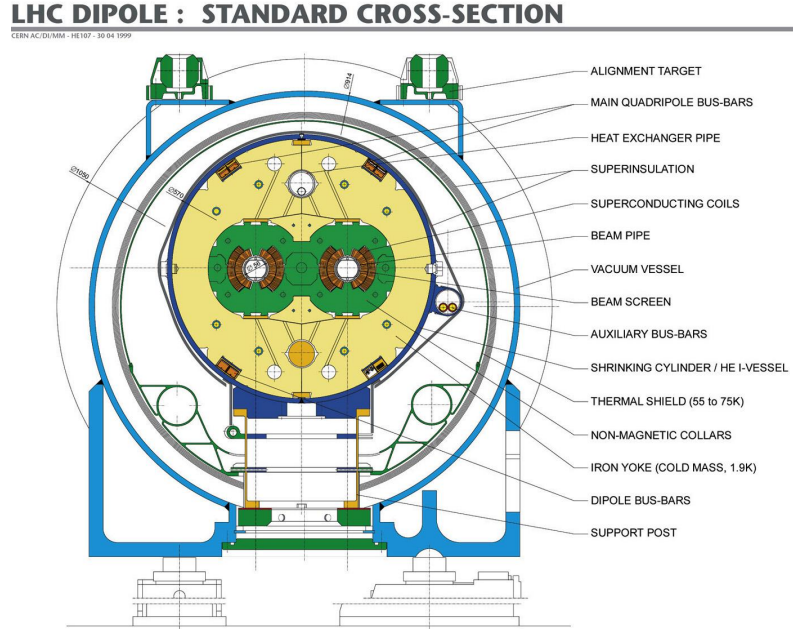


Figure 2.2: Cross section of an LHC dipole magnet [137].

at 1.9 K using superfluid helium. Additionally, niobium couplers are used to dampen the higher-order modes (HOMs) of the RF cavities, preventing beam instabilities [143]. A cross section of an LHC RF cryomodule is shown in Fig. 2.3.

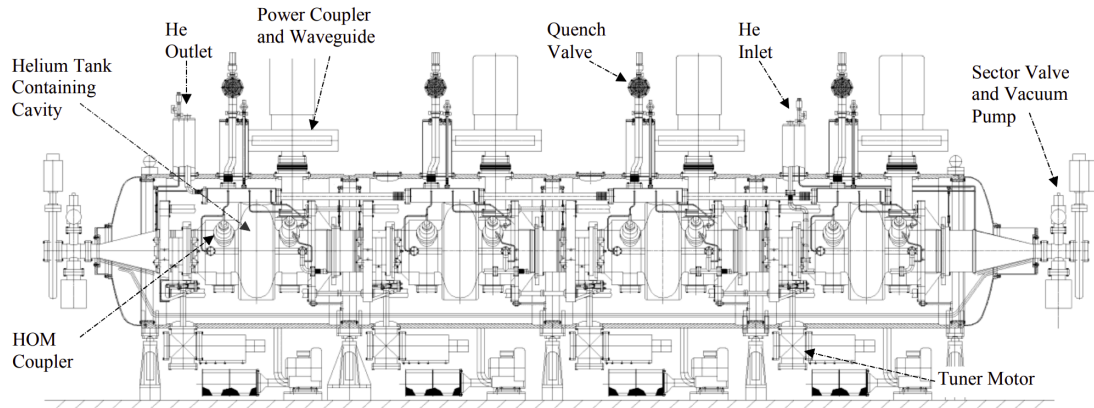


Figure 2.3: Cross section of an LHC RF cryomodule [132].

Dipole magnets are responsible for bending the beams along the arcs of the LHC. For a proton beam energy of 7 TeV, the required field strength is 8.33 T. This necessitates the use of superconducting magnets – the LHC uses magnets with niobium-titanium windings and cooled to 1.9 K with superfluid helium. Complementing the dipoles are the main quadrupole magnets, which are responsible for focusing each beam in the transverse plane. The main quadrupoles have a nominal field gradient of 223 T/m [128]. In total, there are 1,232 dipoles and 392 main quadrupoles along the length of the LHC [144]. Additional magnets perform further refinement of each beam's optics, including modifying their chromaticity, matching different sections of the accelerator, and ultimately colliding the beams in the interaction regions [145]. There are approximately 6,000 of these

additional magnets for corrections, matching, and collisions.

During collisions, there is $O(1)$ GJ of energy stored in the circulating beams and the magnet system. The LHC’s beam dump system is used to safely dispose of the beam at the end of a run or during a malfunction or an emergency. The beam dump system consists of kicker magnets [146] which instrument a dedicated interaction region, deflecting each beam above and away from the LHC cryostat with a rise time of $3\text{ }\mu\text{s}$ when a beam dump is triggered. Each beam travels towards a beam dump cavern, $\sim 750\text{ m}$ away from the kicker magnets, where they enter a 7.7 m long carbon absorber. The carbon absorber is water-cooled and surrounded by ~ 900 tonnes of radiation shielding [128].

During the LHC’s Run 1 (2009–2013), the beam energy was 3.5 or 4 TeV. During the LHC’s Run 2 (2015–2018), the beam energy was increased to 6.5 TeV, corresponding to a centre-of-mass energy $\sqrt{s} = 13\text{ TeV}$ [147].

2.1.2 Luminosity

Exploration of the energy frontier of particle-particle interactions necessitates a sufficiently large set of measured events. In particular, the event rate dN/dt for a particular physics process is given by:

$$\frac{dN}{dt} = L\sigma, \quad (2.1)$$

where L is the instantaneous luminosity and σ is the cross section for the physics process of interest. The instantaneous luminosity is given in units of flux (e.g., $\text{cm}^{-2}\text{s}^{-1}$). To indicate the amount of collision data measured by a detector, the (time) integrated luminosity, $\mathcal{L}$, is typically reported:

$$\mathcal{L} \equiv \int L dt, \quad (2.2)$$

and so the total event yield may be written as $N = \mathcal{L}\sigma$ – that is, the event yield is *proportional* to the integrated luminosity. The integrated luminosity is given in units of inverse area and typically inverse femtobarns (fb^{-1}) where $\text{fb}^{-1} \equiv 10^{43}\text{ m}^{-2}$. Measurements of precision physics and searches for rare processes therefore benefit greatly from the LHC’s high instantaneous luminosity and long data-taking periods.

Assuming Gaussian beam distribution, the instantaneous luminosity at the LHC is given by:

$$L = \frac{N_b^2 n_b f_{\text{rev}} \gamma_r}{4\pi \varepsilon_n \beta^*} F, \quad (2.3)$$

where N_b is the number of particles per bunch, n_b is the number of bunches per beam, f_{rev} is the revolution frequency, γ_r is the relativistic gamma factor, ε_n is the normalized transverse beam emittance, β^* is the beta function at the interaction point, and F is the geometric reduction factor due to the crossing angle of the beams at the interaction point, $150\text{--}200\text{ }\mu\text{rad}$. The revolution frequency and relativistic gamma factor are fixed by the beam energy of 6.5 TeV, $f_{\text{rev}} = 11.245\text{ kHz}$ and $\gamma_r = 7460.5$. The Run 2 values for all other parameters in Eq. 2.3 are given in Table 2.1.

Due to loss of beam during collisions, multiple Coulomb scattering within bunches or “intrabeam” scattering, and emittance growth, the luminosity decays exponentially with a lifetime of 14.9 hours. In order to maximize the luminosity achievable, the LHC is operated in runs of length 5.5 or 12 hours with turnaround times of 1.15 and 7 hours, respectively [128]. Gas ionization chambers monitor the luminosity at the LHC [151].

The ATLAS detector has its own dedicated subsystems for luminosity measurements, the primary detector being LUCID-2 [152]. LUCID-2 consists of two modules placed $\pm 17\text{ m}$ longitudinally from the interaction point, and each module is composed of sixteen photomultiplier tubes using quartz windows as Cherenkov

Table 2.1: Design [128] and actual [148–150] LHC beam parameters for pp collisions at $\sqrt{s} = 13$ TeV during Run 2. The bunch spacing was 25 ns for all years. The variables are described in Section 2.1.2.

Parameter	Design	2015	2016	2017	2018
N_b [$\times 10^{11}$ protons/bunch]	1.15	1.1	1.1	1.1	1.1
n_b [bunches/beam]	2,808	2,232	2,208	2,544	2,544
ε_n [μm]	3.75	3.5	2.5	2.0	2.0
β^* [m]	0.55	0.8	0.4	0.3	0.3

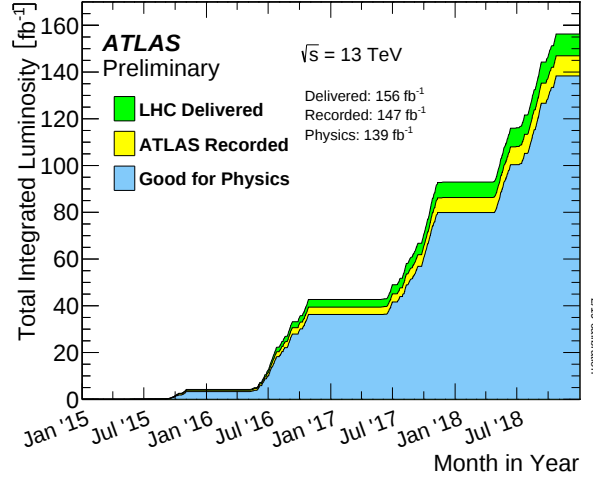


Figure 2.4: The cumulative integrated luminosity provided by the LHC (“LHC Delivered”) and measured by ATLAS (“ATLAS Recorded”) using pp collisions at $\sqrt{s} = 13$ TeV for the full Run 2 data-taking period [155]. The “Good for Physics” fraction requires the reconstructed physics objects to be of good quality.

media. LUCID-2 measures the per-bunch visible interaction rate, μ_{vis} , by counting hits. The visible inelastic cross section, σ_{vis} , represents an absolute calibration and is measured from special LHC runs where the colliding beams are scanned in the transverse plane, also known as van der Meer scans [153, 154]. Using μ_{vis} and σ_{vis} , the per-bunch instantaneous luminosity, L_b , may be calculated:

$$L_b = \frac{\mu_{\text{vis}} f_{\text{rev}}}{\sigma_{\text{vis}}} . \quad (2.4)$$

The total instantaneous luminosity is obtained by summing Eq. 2.4 over all colliding bunch pairs, and the total integrated luminosity is obtained by integrating the total instantaneous luminosity over the data taking period:

$$\mathcal{L} = \int_{\text{period}} dt \sum_{\text{bunches}} L_b . \quad (2.5)$$

For the LHC’s Run 2, ATLAS measured a total integrated luminosity of $\mathcal{L} = 139 \text{ fb}^{-1}$ with a relative uncertainty of 1.7% [148]. This number represents the fraction of the total measured luminosity passing data quality monitoring and corresponds to a data quality efficiency of 95.6% [155]. Fig. 2.4 shows the integrated luminosity delivered by the LHC and measured by ATLAS as a function of time over Run 2.

2.2 The ATLAS Detector

The ATLAS detector [127] is a cylindrically-symmetric detector with a near 4π solid angle coverage. It is situated at Interaction Point 1 at the CERN LHC, 100 m below ground. ATLAS reconstructs collisions with high efficiency and precision using a diverse array of subsystems including:

- An inner tracking system (Section 2.2.1) for reconstructing charged particle tracks, which bend in the field of a central solenoid, as well as primary and secondary vertices;
- A calorimetry system (Section 2.2.2) for reconstructing electrons, photons, jets, and missing transverse momentum;
- A muon spectrometer (Section 2.2.3) for reconstructing muons deflecting in the field of a toroidal magnet system.

A trigger and data acquisition system (Section 2.2.4) uses a combination of hardware-based and software-based triggers to reduce the event rate from 40 MHz to ~ 1.2 kHz. A diagram of the ATLAS detector is shown in Fig. 2.5. In terms of scale, ATLAS is 44 m in length, 25 m in diameter, and weighs 7,000 tonnes.

ATLAS uses a right-handed coordinate system with its origin at the nominal interaction point in the centre of the detector. The z -axis points along the beam pipe, the x -axis points from the interaction point to the centre of the LHC ring, and the y -axis points upwards. Cylindrical coordinates (r, ϕ) are used in the transverse plane, where r is the radial extent and ϕ is the azimuthal angle around the z -axis. The polar angle θ is defined relative to the z -axis, and the pseudorapidity is defined in terms of the polar angle as $\eta = -\ln \tan(\theta/2)$. The pseudorapidity is approximately equal to the rapidity at LHC beam energies. Angular distance is measured in units of $\Delta R = \sqrt{(\Delta\eta)^2 + (\Delta\phi)^2}$. Additional details on the four-momentum parameterization adopted at hadron colliders can be found in Appendix A.

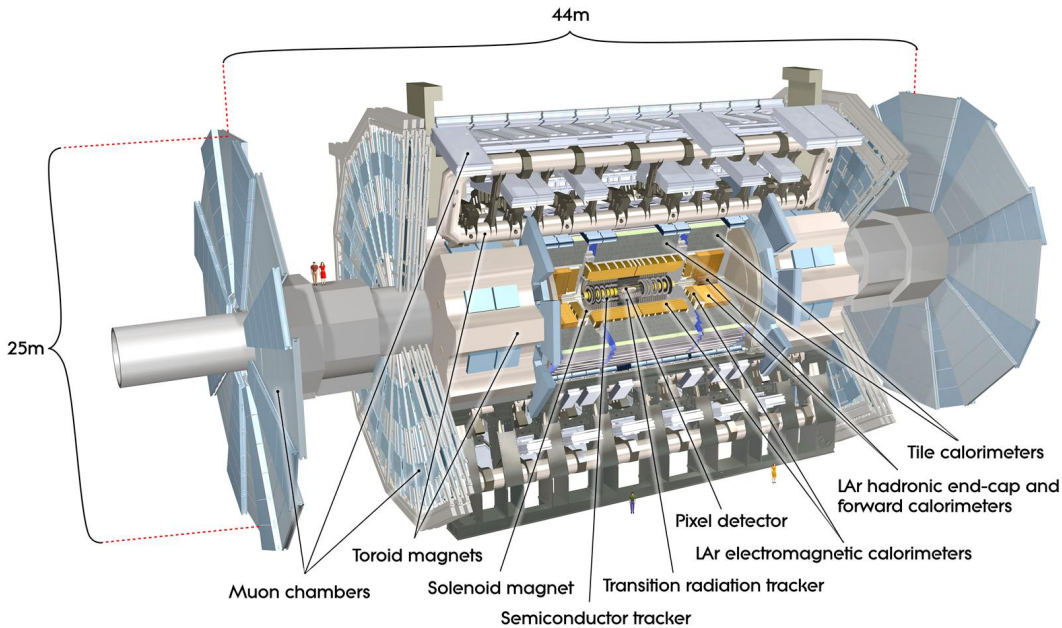


Figure 2.5: The ATLAS detector [156].

2.2.1 Inner Tracking System

The inner tracking system [157, 158], also known as the inner tracker or inner detector, performs the reconstruction of charged particle tracks and vertices, providing momentum measurements, lepton identification, and flavour tagging capabilities in the region $|\eta| < 2.5$ and for track momenta as low as 500 MeV. The design momentum resolution of the inner tracker is:

$$\frac{\sigma_{p_T}}{p_T} = 0.05\% \times p_T \oplus 1\% \text{ for } |\eta| < 2.5 \text{ [127]}. \quad (2.6)$$

Situated closest to the interaction point, the inner tracker utilizes a highly segmented design with radiation-tolerant components to handle the high occupancy of hits, $O(10^3)$ particles every 25 ns, and the high fluence of radiation, up to an integrated fluence of $O(10^{14}) n_{eq}/\text{cm}^2$ at end-of-life, delivered during pp collisions. The inner tracker’s low material budget helps minimize particle scattering and ensures optimal performance of the detector [127].

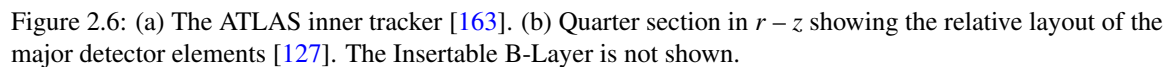
Three subdetectors comprise the inner tracker: the pixel detector, the SemiConductor Tracker (SCT), and the Transition Radiation Tracker (TRT), in order of increasing radial distance from the beam axis. The pixel detector and the SCT are collectively known as the silicon tracker. A diagram of the inner tracker is shown in Fig. 2.6a and its relative layout is shown in Fig. 2.6b.

The inner tracker is enveloped by the central solenoid [159, 160], which provides a 2 T magnetic field [161] and enables measurements of a track’s momentum and the sign of its charge. The central solenoid consists of a single niobium-titanium superconducting coil and is cooled to 4.5 K using liquid helium [162].

Silicon Tracker: For the pixel detector and the SCT, the operation relies on reverse-biased silicon sensors: the sensor’s readout implants (i.e., the anode) are biased with -150 V relative to the sensor’s backplane (i.e., the cathode) at the start of operation and up to -600 V at end-of-life [127]. Ionizing particles traversing the depletion layer in the silicon’s bulk generate electron-hole pairs, $O(100)$ pairs/ μm [164]. The electrons and holes drift in the electric field within the silicon and are measured as a current by readout electronics. If the total deposited charge exceeds a set threshold, it is regarded as a “hit”. The thickness of the sensors is typically a compromise between the operating voltage, signal strength, and design complexity, while the segmentation of the sensors is chosen for the desired precision, particle occupancy, and noise performance [127].

The pixel detector [165] is closest to the interaction point in radius and has the highest granularity of the tracking subdetectors. The pixel detector consists of three concentric cylinders or “barrels” surrounding the beam axis and two “endcaps”, each consisting of three disks perpendicular to the beam axis and longitudinally placed at opposite ends of the barrel. The barrel layers and endcaps are composed of units called modules [166], and each module performs the readout of a $256 \mu\text{m}$ thick n^+ -in- n -type silicon sensor [167]. Each sensor is segmented as a 328×144 array of pixels of size $50 \times 400 \mu\text{m}^2$ in $r - \phi \times z$ with intrinsic accuracies of $10 \mu\text{m}$ in $r - \phi$ and $115 \mu\text{m}$ in r/z . The readout electronics [168] provide time-of-arrival information such that each hit can be uniquely associated to a specific bunch crossing. A schematic view of a barrel pixel module is shown in Fig. 2.8a.

During the long shutdown between the LHC’s Run 1 and 2, the pixel detector was expanded by the installation of the Insertable B-Layer (IBL) [169–171] at a radius of 33.5 mm. The IBL consists of a single barrel layer of modules built from $200 \mu\text{m}$ thick n^+ -in- n -type silicon sensors. Each sensor is segmented as a 336×80 array of pixels of size $50 \times 250 \mu\text{m}^2$ in $r - \phi \times z$. As the IBL’s end-of-life expected fluence is 20 times larger than that of the innermost pixel detector layer, the readout electronics of the IBL were redesigned



with an improved radiation hardness with respect to the pixel detector [172]. A 3D rendering of the IBL is shown in Fig. 2.7. Collectively, the IBL and pixel detector provide typically four hits per track [127].

Following the pixel detector is the SCT [173, 174]. The SCT consists four barrel layers and two endcaps, each consisting of nine disks. Analogous to the pixel detector, the barrel layers and endcaps are composed of modules [175, 176] which perform the readout of 285 μm thick p -in- n -type silicon sensors [177]. Each sensor is segmented longitudinally into strips with a mean pitch of 80 μm and intrinsic accuracies of 17 μm in $r - \phi$ and 580 μm in r/z . A single module, shown schematically in Fig. 2.8b, consists of two sensors mounted back-to-back with a 40 mrad stereo angle rotation for simultaneous measurements of ϕ and r/z . A binary readout architecture is used for all modules [178, 179]. The SCT provides typically eight hits per track [127].

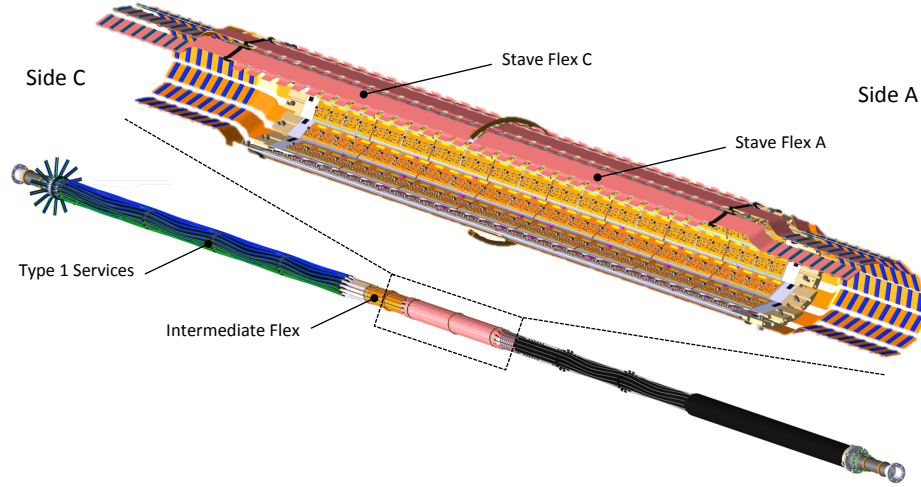


Figure 2.7: The ATLAS Insertable B-Layer [171].

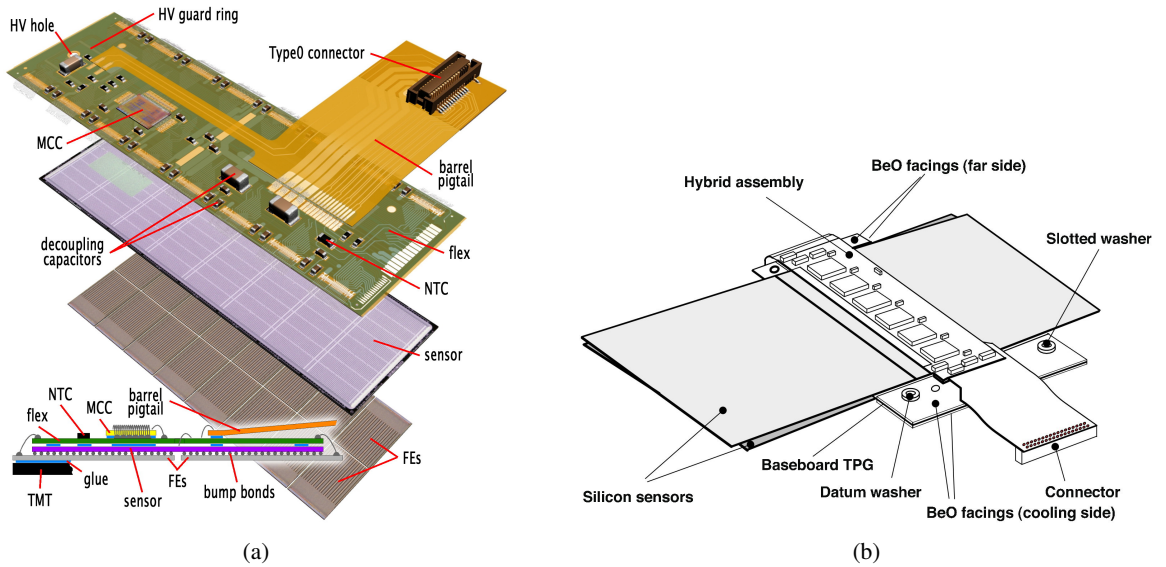


Figure 2.8: (a) Schematic view of a barrel pixel module [166]. The front-end chips (FES), module control chip (MCC) [180], temperature monitoring (NTC), and sensor elements are indicated. (b) Schematic view of a barrel SCT module [175].

Transition Radiation Tracker: The outermost subdetector of the inner tracker is the TRT [181, 182], whose basic detection units are 4 mm diameter gaseous drift (“straw”) tubes. Each tube contains a gold-plated tungsten anode wire, which runs along the tube’s axis of symmetry, and an associated aluminum cathode layer, which is embedded in the tube’s polyimide wall and biased at -1.53 kV relative to the anode. The volume of each tube is filled with a 70% Xe / 27% CO₂ / 3% O₂ gas mixture [183, 184]. When a charged particle passes through, the gas ionizes and generates a cluster of charge. The cluster avalanches with a gain of 2.5×10^4 as it drifts in the electric field towards the anode, where its arrival time is measured by readout electronics [185]. Using knowledge of the drift velocity of electrons in the gas mixture, the track’s position may be determined.

The TRT consists of a barrel detector [186], with straw tubes arranged parallel to the beam axis, and two endcap detectors [187], with straw tubes arranged radially into layers. The barrel detector is divided into three concentric rings, with the readout of each ring instrumented by 32 modules. Each endcap detector is divided into twenty wheels, each constituting a single module. Interspersed between the barrel straw tubes and the endcap straw layers are polypropylene wires and foils, respectively, which act as transition radiation materials. The transition radiation emitted by charged particles can result in further ionization and additional hits. The TRT provides typically 36 hits per track [127].

Environmental Control: Evaporative cooling is used to maintain the temperature of the inner tracker [127, 188]. In order to minimize the impact of radiation damage on the silicon sensors, the pixel detector and the SCT are both operated at -7 °C. The coolant, C₃F₈, is circulated through capillaries located on each detector. Additionally, the volume enclosing each detector is flushed with dry nitrogen gas to prevent condensation on the cold surfaces. The TRT, in contrast, is operated at room temperature, 20 °C. Carbon dioxide is circulated in the volume between the straw tubes, acting as a thermal conduit to the edges of each module where capillaries carrying C₆F₁₄ remove excess heat.

2.2.2 Calorimetry System

Surrounding the inner tracker and the central solenoid is the calorimetry system [189, 190], which performs the reconstruction of electrons, photons, jets, and missing transverse momentum as well as the identification of electrons and photons in the region $|\eta| < 4.9$.

The electromagnetic calorimeters measure the electromagnetic showers created by electrons, which lose energy primarily through bremsstrahlung, and photons, which interact primarily through pair production [26, 191]. The electromagnetic calorimeters have sufficient depth, >22 and >24 radiation lengths in the barrel and endcaps, respectively, to contain the showers [127]. The energy resolution of the electromagnetic calorimeters is:

$$\frac{\sigma_E}{E} = \begin{cases} \frac{10\%}{\sqrt{E}} \oplus 0.5\% & |\eta| < 3.2 \text{ [189]} \\ \frac{30\%}{\sqrt{E}} \oplus 3.5\% & 3.1 < |\eta| < 4.9 \text{ [192]} \end{cases}. \quad (2.7)$$

In contrast, the hadronic calorimeters measure the hadronic showers created by charged or neutral hadrons interacting via the strong or electromagnetic interactions. The hadronic calorimeters have sufficient depth, 9.7 and 10 interaction lengths in the barrel and endcaps, respectively, to contain the hadronic showers [127], shielding the muon system from punch-through and providing good measurements of missing transverse momentum. The energy resolution of the hadronic calorimeters is:

$$\frac{\sigma_E}{E} = \begin{cases} \frac{50\%}{\sqrt{E}} \oplus 3\% & |\eta| < 3.2 \text{ [190]} \\ \frac{95\%}{\sqrt{E}} \oplus 7.5\% & 3.1 < |\eta| < 4.9 \text{ [192]} \end{cases}. \quad (2.8)$$

In general, the response of a hadronic calorimeter for the electromagnetic (e) and non-electromagnetic (h) components of a shower are *not* equal. The ratio, e/h , is a function of the fraction of the electromagnetic component [193], which in turn is a function of the energy of the incident particle [194, 195], as well as depends on the calorimeter's material.

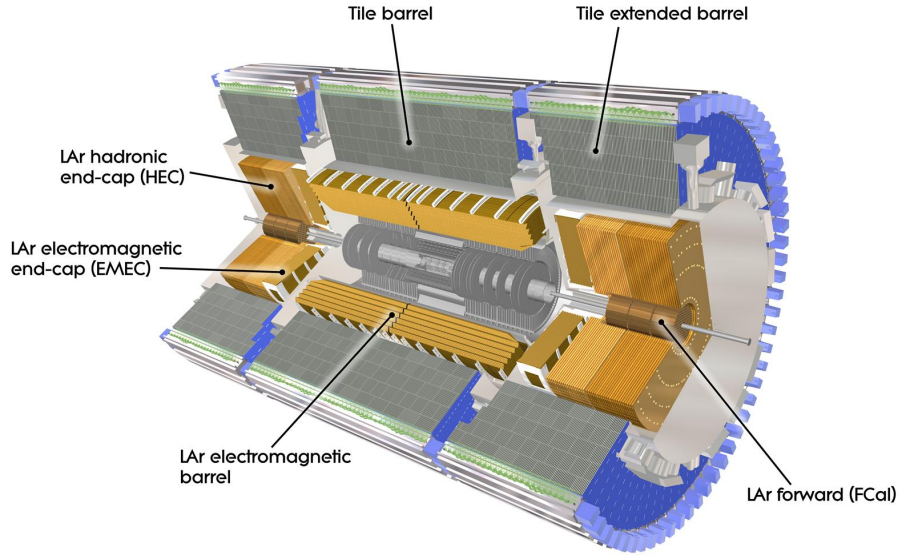
All of the calorimeters in ATLAS are *sampling* calorimeters, consisting of absorber layers, which reduce the energy of the incident particle, interleaved with active layers, which produce a measurable signal dependent on the energy of the incident particle [191]. In contrast, *homogeneous* calorimeters – for example, the CMS electromagnetic calorimeter [196] – are composed of a single material which acts as both the absorber and the active material. The intrinsic energy resolution of homogeneous calorimeters is typically better than that of sampling calorimeters. In sampling calorimeters, the energy deposited in the active layers fluctuates event-by-event due to the presence of independent absorber layers, and so $\sigma_E/E \propto \sqrt{t/E}$ where t is the thickness of the absorber layers [191].

The calorimetry system is divided into three subsystems: the electromagnetic calorimeters, the hadronic calorimeters, and the forward calorimeters. The central solenoid is enveloped by the electromagnetic barrel calorimeter, which in turn is enveloped by the hadronic barrel calorimeter. The forward calorimeters instrument the region close to the beam axis. A diagram of the calorimetry system is shown in Fig. 2.9a and its relative layout is shown in Fig. 2.9b.

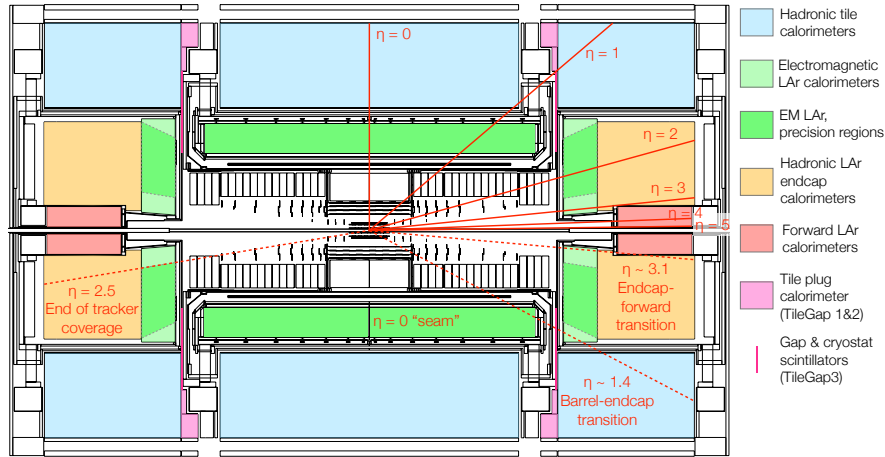
Electromagnetic Calorimetry: The electromagnetic calorimetry subsystem is composed of a barrel calorimeter [199], which instruments $|\eta| < 1.475$, and two endcap calorimeters [200], which instrument $1.375 < |\eta| < 3.2$. It utilizes liquid argon (LAr) as an active material, which possesses good linearity, stability, and radiation tolerance, between accordion-shaped layers of lead absorbers, which offer high granularity and good azimuthal coverage. The thickness of the absorbers varies from 1.13 to 1.53 mm in the barrel and from 1.7 to 2.2 mm in the endcaps, and the spacing between the absorbers is 4.2 mm in the barrel and varies from 0.9 to 3.1 mm in the endcaps. Both the thickness and the spacing vary with location in $|\eta|$ and r . Copper-polyimide electrodes are placed halfway between subsequent absorber layers [201]. The electrodes consist of two high-voltage copper layers capacitively coupled to an inner copper layer at ground potential, which performs signal readout. The high-voltage layers are biased with +2 kV in the barrel; in the endcaps, the potential varies from +1 to +2.5 kV as the spacing between the absorbers also varies. All of ATLAS's liquid argon calorimeters utilize a common set of readout electronics [202, 203].

The barrel calorimeter is divided into two halves at $z = 0$. Each half-barrel consists of sixteen identical sectors or modules [204] and each endcap calorimeter consists of eight identical modules [205]. Modules are further divided into three sampling layers of varying $\Delta\eta \times \Delta\phi$ segmentation. The first sampling layer is segmented finely as $0.025/8 \times 0.1$. The second sampling layer collects the majority of an electromagnetic shower's energy and is segmented as finely as 0.025×0.025 . The third sampling layer collects the tail of a shower and has uniform segmentation, 0.050×0.025 [127]. A sketch of an electromagnetic calorimeter module is shown in Fig. 2.10a.

A presampler, placed immediately in front of the electromagnetic calorimeter and instrumenting $|\eta| < 1.8$, is used to correct for the energy lost by incident particles upstream of the calorimeter [206]. It consists of two half-barrels and two endcaps, each composed of 32 modules with uniform $\Delta\eta \times \Delta\phi$ segmentation, 0.025×0.1 .



(a)



(b)

Figure 2.9: (a) The ATLAS calorimetry system [197]. (b) Cross section in $r-z$ showing the relative layout of the major detector elements [198].

Each module utilizes a thin active layer of liquid argon enveloped by readout electrodes.

Hadronic Calorimetry: The hadronic calorimetry subsystem is composed of a barrel calorimeter [207–209], which instruments $|\eta| < 1.7$, and two endcap calorimeters [210], which instrument $1.5 < |\eta| < 3.2$. Unlike the electromagnetic calorimetry system, the hadronic calorimetry system uses different active and absorber layers in the barrel versus the endcaps.

The barrel calorimeter, known as the tile calorimeter, utilizes steel as the absorber and scintillating tiles as the active layers. It is divided into a central barrel and two extended barrels, each composed of 64 modules. Special modules instrument the gap regions between the barrel and each extended barrel [127]. The scintillating tiles are made of polystyrene [211] and are arranged radially and normal to the beam axis. The tiles are built into the steel bulk of each module in a staggered fashion to ensure good azimuthal coverage, resulting in a 4.7:1 steel-scintillator volume ratio. Wavelength-shifting fibres are used to collect the scintillation light produced by each tile, and groups of fibres are bundled together and read out by photomultiplier tubes [212, 213]. The granularity of bundling is used to define three radial sampling depths of varying $\Delta\eta \times \Delta\phi$ segmentation, 0.1×0.1 for the inner and middle layers and 0.2×0.1 for outer layer [127]. A sketch of an tile calorimeter module is shown in Fig. 2.10b.

The endcap calorimeters utilize liquid argon an active material between plates of copper absorbers. Each endcap is divided into a front wheel and a rear wheel, each composed of 32 identical modules. The absorber plates have a typical thickness of 25 and 50 mm in the front and rear wheels, respectively, leading to finer and coarser sampling of hadronic showers [210]. The liquid argon gaps between the plates have a uniform thickness of 8.5 mm, and each gap is segmented into four drift zones of equal width by three copper-polyimide electrodes. The outer electrodes are biased with +2 kV relative to the central electrode, which performs signal readout [210]. By etching these electrodes, readout cells of varying $\Delta\eta \times \Delta\phi$ segmentation are defined, 0.1×0.1 for $|\eta| < 2.5$ and 0.2×0.2 for $|\eta| > 2.5$ [127]. The calorimeter's signal-to-noise ratio is optimized by its integrating front-end electronics into the cold mass of the endcaps [214].

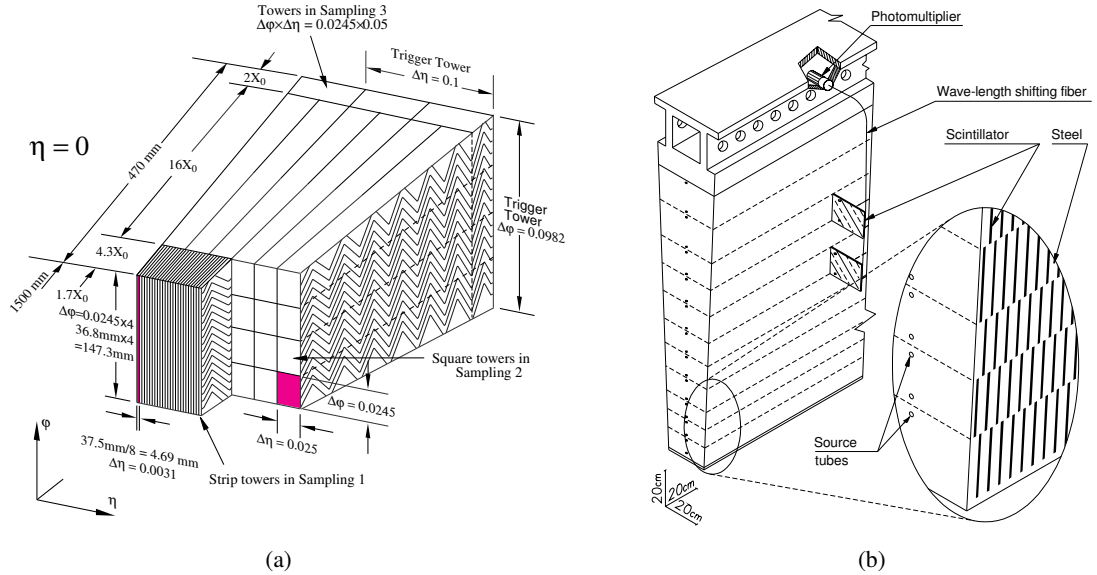


Figure 2.10: (a) Sketch of an electromagnetic calorimeter barrel module [189]. The three sampling layers of varying $\Delta\eta \times \Delta\phi$ segmentation are indicated. (b) Sketch of a tile calorimeter module [127].

Forward Calorimetry: The forward calorimetry subsystem [192] instruments the region close to the beam axis, $3.1 < |\eta| < 4.9$. It is composed of two endcaps calorimeters, each of which is composed of an electromagnetic module and two hadronic modules. Each module consists of a 45 cm long cylindrical absorber with liquid argon as the active material. For the electromagnetic module, the absorber is built from a stack of copper plates as copper provides the optimal resolution and thermal control [192]; for the hadronic modules, the absorber is built from a matrix of tungsten slugs as tungsten provides good containment of showers and minimizes their lateral spread [127]. A 45 cm long cylindrical brass plug is placed directly behind the third module of each endcap to shield the muon system from radiation damage. The relative layout of the forward calorimetry subsystem is shown in Fig. 2.11a.

The modules are segmented parallel to the beam axis by cylindrical bores housing the readout electrodes [127]. Each readout electrode consists of a grounded copper tube enclosing a high-voltage coaxial rod. The rod is composed of copper, in the case of electromagnetic modules, or tungsten, in the case of hadronic modules, and biased such that a 1 kV/mm electric field is generated. The gap between the tube and the rod constitutes the active liquid argon volume. The gap thickness, which varies from 0.27 to 0.51 mm, is nearly an order of magnitude smaller than that of the other liquid argon subsystems to avoid problems due to positive ion build-up [215]. Each module is segmented in $\Delta x \times \Delta y$ by ganging together varying numbers of electrodes [192]. The segmentation can be as fine as 2 and 3.5 cm² in the electromagnetic and hadronic modules, respectively [127]. A schematic of the electrode structure is shown in Fig. 2.11b.

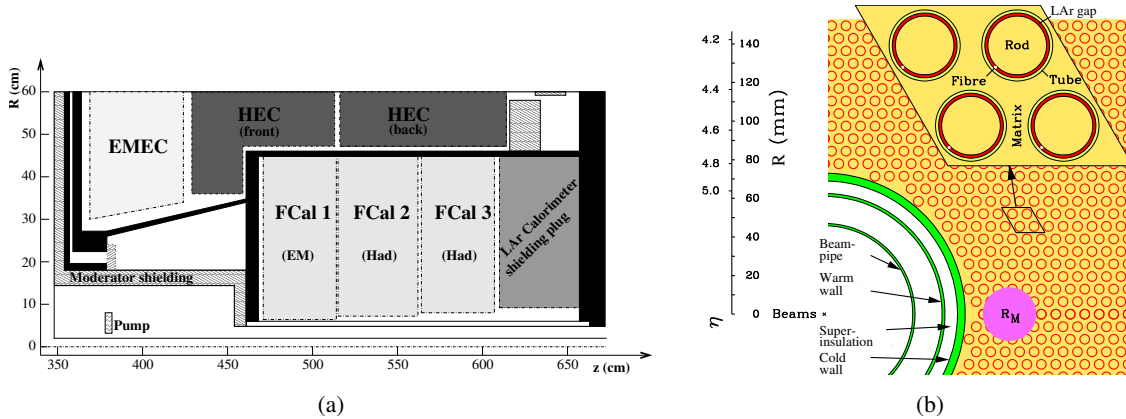


Figure 2.11: (a) Quarter section of the ATLAS endcap calorimeters showing the relative layout of the major detector elements [127]. The electromagnetic endcap (EMEC), the hadronic endcap (HEC), and the electromagnetic (FCal 1) and hadronic (FCal 2 and FCal 3) forward calorimeters are indicated. (b) Schematic of the forward calorimeter electrode structure [127]. The relative size of the Molière radius (R_M) is also indicated.

Environmental Control: All of ATLAS's liquid argon calorimeters are housed in cryostats which maintain the liquid argon at a cold and stable temperature, 87 K. The central solenoid and electromagnetic barrel calorimeter share a common cryostat [216], where an inner cavity houses the solenoid and an outer cavity houses the calorimeter. Each set of endcap calorimeters, including the electromagnetic endcap, hadronic endcap, and forward calorimeters, share a common cryostat [217].

2.2.3 Muon Spectrometer:

The muon spectrometer [218] is the outermost system of the ATLAS detector. It performs the reconstruction of charged particle tracks exiting the barrel and endcap calorimeters, providing momentum measurements and lepton identification in the region $|\eta| < 2.7$ and for track momenta as low as 3 GeV and as high as 3 TeV. The design momentum resolution of the muon spectrometer is:

$$\frac{\sigma_{p_T}}{p_T} = 0.01\% \times p_T \text{ for } p_T > 300 \text{ GeV [218]}. \quad (2.9)$$

Additionally, it is designed to trigger on tracks in the region $|\eta| < 2.4$ [127].

The muon spectrometer is divided into a barrel, which consists of three concentric cylindrical shells, as well as two endcaps, each of which consists of four wheels. The operation of the muon spectrometer is realized through a suite of detector technologies embedded within large air-core toroid magnets sharing the ϕ symmetry of the spectrometer. A diagram of the muon spectrometer is shown in Fig. 2.12a, and its relative layout in $r - z$ and $x - y$ is shown in Figs. 2.12b and 2.12c, respectively.

Precision tracking measurements are performed by the Monitored Drift Tubes (MDTs). The MDTs measure a track's coordinates in the bending ($r - z$) plane of the toroid, adopting the coordinate in the non-bending ($x - y$) plane from a matched trigger chamber [127]. In the more forward regions of the spectrometer, the MDTs are complemented by the Cathode Strip Chambers (CSCs). The CSCs are segmented both longitudinally and transversely relative to the bending plane and therefore provide coordinates in the bending and non-bending planes simultaneously. An optical alignment system [219–221] is used to monitor the positions and internal deformations of the precision tracking chambers, ensuring the locations of the detector elements are known to better than 40 μm .

The Resistive Plate Chambers (RPCs) and the Thin Gap Chambers (TGCs) perform fast triggering on tracks in the barrel and endcaps, respectively. Both the RPCs and the TGCs are capable of delivering information on a track within a few tens of nanoseconds following its passage [127]. Their spatial resolution enables excellent discrimination of muon p_T both centrally and forwardly, and their timing resolution enables bunch crossing identification with an accuracy of 99%. Owing to their longitudinal and transverse segmentation relative to the bending plane of a track, the coordinates measured by the trigger chambers complement those measured by the precision tracking chambers.

The air-core toroid magnet system consists of a barrel [222] and two endcap [223] toroids, each divided in eight octants and providing magnetic fields of 0.5 and 1 T, respectively. The toroid magnets are built from niobium-titanium-copper superconducting coils [224]. They are contained within individual stainless-steel cryostats [225–227] and cooled to 4.6 K using liquid helium [228]. The magnetic field of the toroid magnet system is continuously monitored by 1,700 3D Hall sensors distributed throughout the muon spectrometer [229]. The geometry of the magnet system is shown in Fig. 2.13.

Monitored Drift Tubes: The MDTs [231, 232] instrument precision tracking in the region $|\eta| < 2.7$. A single MDT consists of a 3 cm diameter aluminum tube filled with a 93% Ar / 7% CO₂ / $\leq 0.1\%$ H₂O gas mixture [233]. A gold-plated, tungsten-rhenium wire runs along the tube's axis of symmetry [234]. The wire, which acts as the anode, is biased at +3.08 kV relative to the tube, which acts as the cathode. As ionizing particles traverse the volume of the tube, their deposited charge avalanches with a gain of 2×10^4 as it drifts in the electric field towards the anode. Readout electronics measure the leading edge of the resulting signal, which yields the distance of the track from the wire with an accuracy of 80 μm [235]. A cross section of a

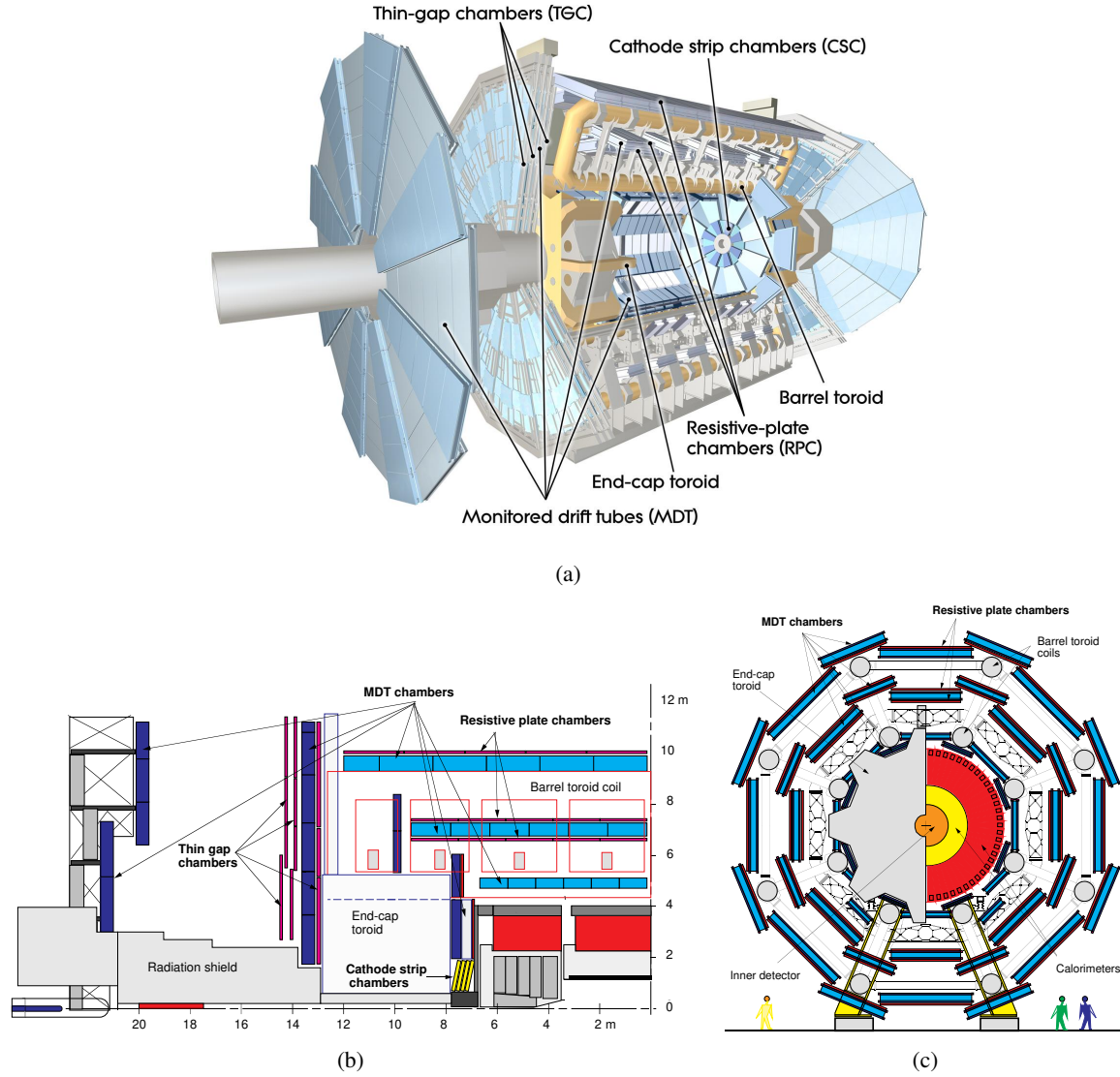


Figure 2.12: (a) The ATLAS muon spectrometer [230]. (b) Quarter section in $r-z$ and (c) cross section in $x-y$ showing the relative layout of the major detector elements [218].

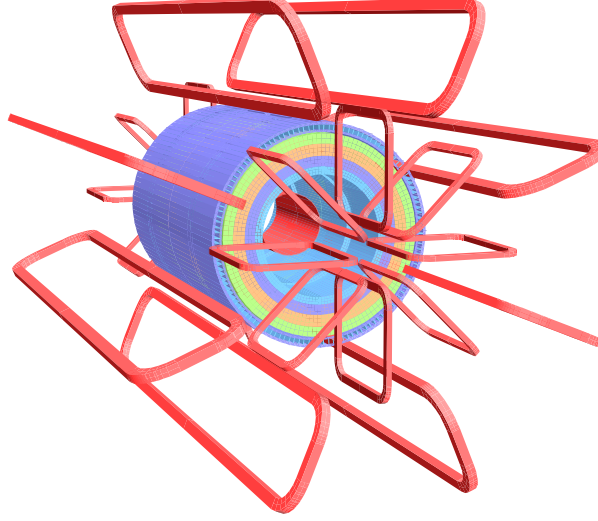


Figure 2.13: Geometry of the ATLAS magnet system [127], showing the windings of the central solenoid (red, innermost), the steel of the tile calorimeter and the return yoke (blue, purple, orange, and green), and the windings of the barrel and endcap toroids (red, outermost). To give a sense of the scale, the central solenoid's outer diameter is 2.56 m and its axial length is 5.8 m, while the barrel toroid's outer diameter is 20.1 m and its axial length is 25.3 m [127].

MDT is shown in Fig. 2.14a.

Groups of MDTs are assembled into rectangular (in the barrel) or trapezoidal (in the endcaps) multi-layer chambers [236], each with a combined accuracy of $35\ \mu\text{m}$ in the bending plane. The shapes and dimensions are chosen such the solid angle coverage is maximized [237]. The multi-layer chambers, in turn, are assembled into three barrel layers and two endcaps, each consisting of four disks. In both the barrel and the endcaps, the axis of symmetry of each MDT points along the ϕ direction [127]. The mechanical structure of a MDT chamber is shown in Fig. 2.14b. The MDTs provide 20 coordinate measurements per track in both the barrel and the endcaps [127].

Cathode Strip Chambers: In the region $2 < |\eta| < 2.7$ for the innermost wheel of each endcap, the counting rates exceed the limit for safe operation of the MDTs; instead, the CSCs [238] instrument this region due to their higher rate capability and timing resolution [127]. A CSC is a multi-wire proportional chamber with gold-plated, tungsten-rhenium anode wires arranged radially from the beam axis and planes of cathode strips arranged perpendicular and parallel to the anode wires, enabling simultaneous measurements of η and ϕ . The anode wires are biased with +1.9 kV relative to the cathode strips, and the volume of each CSC is filled with an 80% Ar / 20% CO₂ gas mixture. Ionizing particles traversing the gas generate clusters of charge which avalanche with a gain of 6×10^4 . The clusters are collected at the anode, and the induced charge distribution on the cathode strips is read out by radiation-hard electronics [239]. A cross section of a CSC cell is shown in Fig. 2.15a.

Four planes of CSCs are assembled into chambers of large and small volume designs [127]. To maximize the precision of the charge distribution measurement while minimizing the number of readout channels, every two neighbouring strips which are read out have two floating strips between them [240, 241], resulting in an average readout strip pitch of 5.4 mm. The capacitive coupling of the floating strips to the readout strips reduces the non-linearity of the position measurement [127]. Each CSC additionally uses the relative charge

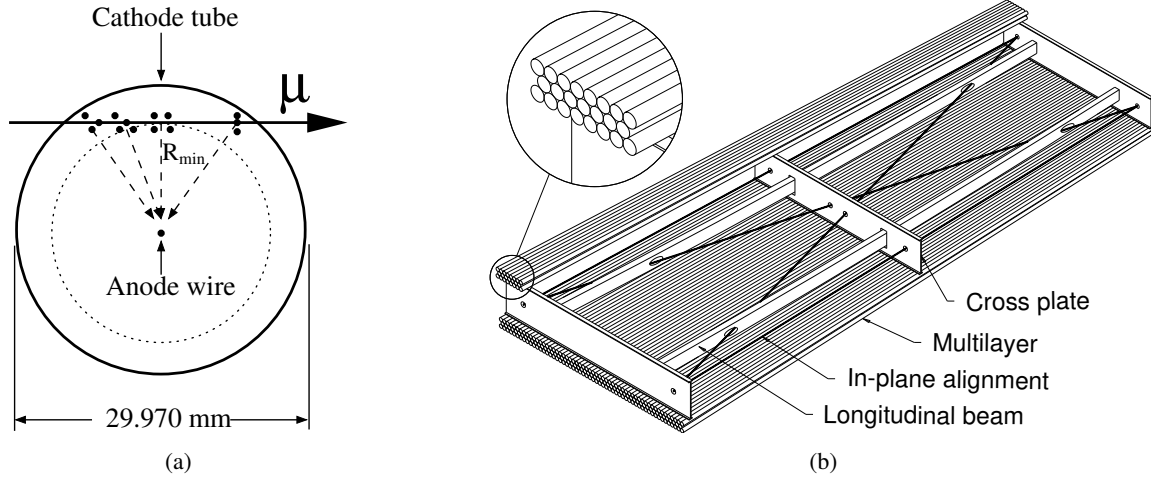


Figure 2.14: (a) Cross section of a MDT [127]. (b) Mechanical structure of a MDT chamber [218]. The aluminum frame carries two multi-layers, each with three or four layers of MDTs. The width of the frame varies from 1 to 2 m, and the length of the frame varies from 1 to 6 m [127].

distribution across the cathode strips to obtain the precision coordinate, making the measurement immune to fluctuations in the gas. Altogether, the spatial resolution in the bending plane is $40\mu\text{m}$. The cathodes strips parallel the anode wires have a coarser resolution of 5 mm in the non-bending plane [127]. The time of arrival of the clusters is obtained using the logical OR of the signals in the corresponding cathodes strips of each plane, resulting in a timing resolution of 7 ns and ensuring reliable bunch crossing identification [127].

The CSC system consists of two endcap disks, each composed of eight large chambers and eight small chambers. A 3D view of an endcap disk is shown in Fig. 2.15b. The CSCs provide four coordinate measurements per track [127].

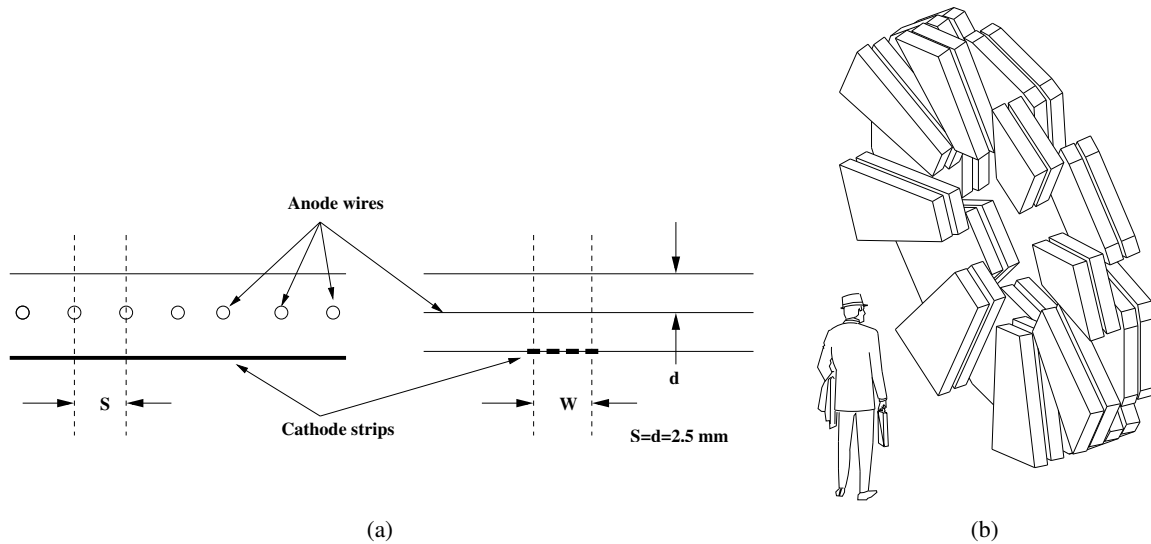


Figure 2.15: (a) Left: the structure of a CSC cell, looking down the wires; right: the structure of a CSC cell in the bending plane, looking down the strips [127]. The wire pitch s is equal to the anode-cathode spacing $d = 2.5\text{ mm}$. (b) A 3D view of the CSC chamber layout [218].

Resistive Plate Chambers: The RPCs [242–244] comprise the trigger system in the central region of the muon spectrometer, $|\eta| < 1.05$. A RPC consists of two parallel resistive plates, each composed of phenolic-melaminic plastic laminate with a thin graphite coating and spaced 2 mm apart [245]. The gap between the plates is filled with a 94.7% $\text{C}_2\text{H}_2\text{F}_4$ / 5% iso- C_4H_{10} / 0.3% SF_6 gas mixture [246], and the plates are biased such that a uniform 4.9 kV/mm electric field is generated. When an ionizing particle traverses the gaseous volume, the resulting avalanche current induces a signal on planes of copper pick-up strips flanking the outermost edges of each plate. The strips have a 25–35 mm pitch, with one plane segmented along the η direction and the other plane segmented along the ϕ direction. The layout of a readout strip plane is shown in Fig. 2.16a. Gallium-arsenide readout electronics, chosen for their reliability, are mounted directly onto the RPCs [247, 248]. The RPCs have a good spatial resolution, 10 mm in both z and ϕ , and an excellent timing resolution, 1.5 ns [127]. Additionally, they are capable of operating at high counting rates, up to a maximum of 700 Hz/cm² [249].

Two independent detector layers, each with a gaseous volume and two planes of readout strips, are embedded in a paper honeycomb matrix and constitute a unit [127], as shown in Fig. 2.16b. Contiguous units are joined to form sectors, and three barrel layers are built from sixteen sectors each.¹ The large transverse separation between the innermost and outermost barrel layers, 2.4–3 m, enables triggering on high- p_T tracks, 9–35 GeV. In contrast, the small transverse separation between the innermost and middle layers, ~ 0.5 m, enables triggering on low- p_T tracks, 6–9 GeV. The RPCs provide six coordinate measurements per track [127].

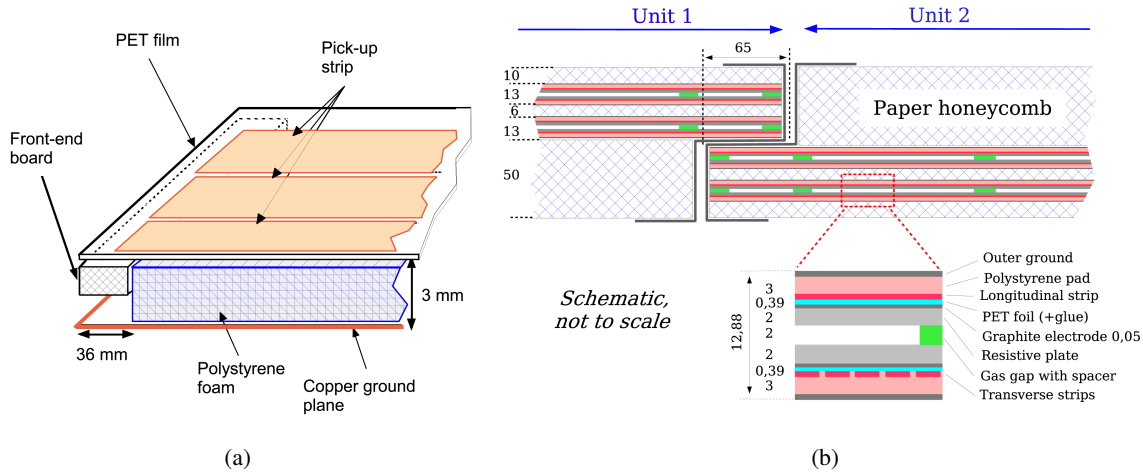


Figure 2.16: (a) Layout of a RPC pick-up strip plane [127]. (b) Cross section of a RPC sector showing units [127]. The dimensions are given in millimeters.

Thin Gap Chambers: The TGCs [242, 251] comprise the trigger system in the forward regions of the muon spectrometer, $1.05 < |\eta| < 2.4$. A TGC is a multi-wire proportional chamber with gold-plated, tungsten anode wires arranged azimuthally to the beam axis [218]. The wires, which are spaced 1.8 mm apart, are sandwiched halfway between two cathode planes, which are spaced 2.8 mm apart. Each cathode plane is composed of a fiberglass laminate (G-10) plate with a thin graphite coating. One of the planes is etched with copper pick-up strips extending radially from the beam axis, segmented with an azimuthal granularity of 2–3 mrad [127]. The volume of the chamber is filled with a 55% CO_2 / 45% n- C_5H_{12} gas mixture [252]. The gas confers an

¹During the LHC's shutdown between Run 1 and Run 2, a fourth layer of RPCs were installed in the detector feet region, $-2.16 < \phi < -1.77$ and $-1.37 < \phi < -0.98$, to cover holes in the geometrical acceptance caused by support structures [250].

avalanche gain of 3×10^5 when the anode wires are biased with +2.9 kV relative to the cathode strips. Anode wires are ganged together in groups of 6–31 wires and measure the coordinate in the bending plane with a resolution of 2–6 mm, while the cathode strips measure the coordinate in the non-bending plane with a resolution of 3–7 mm [127]. Owing to the high electric field between cathode planes and the small anode wire pitch, the TGCs have an excellent timing resolution, 3.7 ns [253]. A cross section of a TGC is shown in Fig. 2.17a.

The TGCs are embedded in a paper honeycomb matrix and assembled into doublet and triplet modules with two and three TGCs, respectively, as shown in Fig. 2.17b. The modules are then assembled into an inner endcap instrumenting $1.05 < |\eta| < 1.92$ and an outer endcap instrumenting $1.92 < |\eta| < 2.4$. The inner endcap is composed of a single layer of doublet modules, while the outer endcap is composed of one layer of triplet modules followed by two layers of doublet modules. The TGCs provide nine coordinate measurements per track [127].

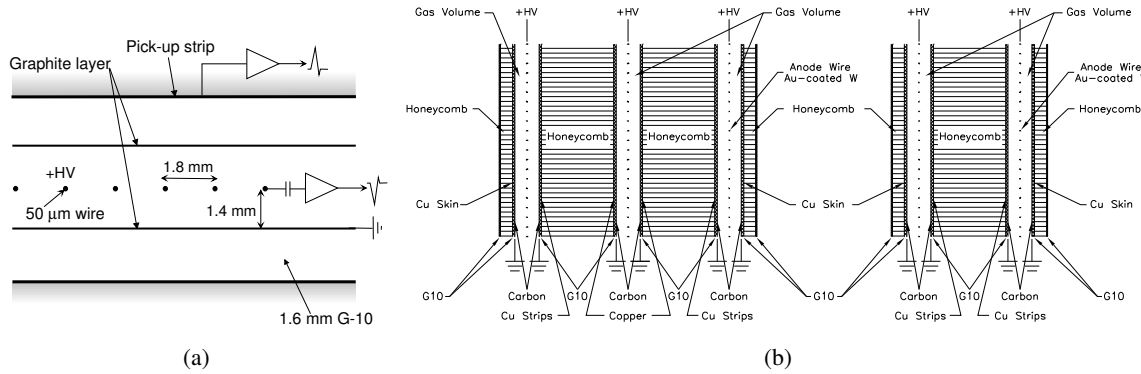


Figure 2.17: (a) Cross section of a TGC [127]. (b) Cross sections of TGC triplet (left) and doublet (right) modules [127]. The relative width of the gas gaps is increased for clarity.

2.2.4 Trigger and Data Acquisition System

The trigger and data acquisition (TDAQ) system [254–257] performs the online processing, selecting, and storing of collision events of interest for offline analysis. The TDAQ system utilizes a two-stage trigger, consisting of a Level-1 (L1) trigger and High-Level Trigger (HLT), and is responsible for reducing the data rate from the LHC bunch crossing frequency, 40 MHz, to a level acceptable by bandwidth and storage constraints. A diagram of the TDAQ system during Run 2 is shown in Fig. 2.18.

The L1 trigger [254], the first stage, is a hardware-based trigger and uses reduced-granularity information from the calorimeters and muon detectors [257]. In particular, the RPCs and TGCs are used to trigger on muons, and the calorimeters are used to trigger on electrons/photons, taus, jets, missing transverse momentum, and total transverse energy [127]. The L1 trigger accepts events with a maximum read-out rate of 100 kHz and a latency of 2.5 μs [257].

The HLT [255] is a software-based trigger. The HLT uses fast trigger algorithms for early event rejection followed by regional reconstruction using information seeded by the L1 trigger to reduce the read-out rate to ~ 1.2 kHz [257]. The latency of the HLT is $O(200)$ ms due to the CPU-intensive algorithms used [258].

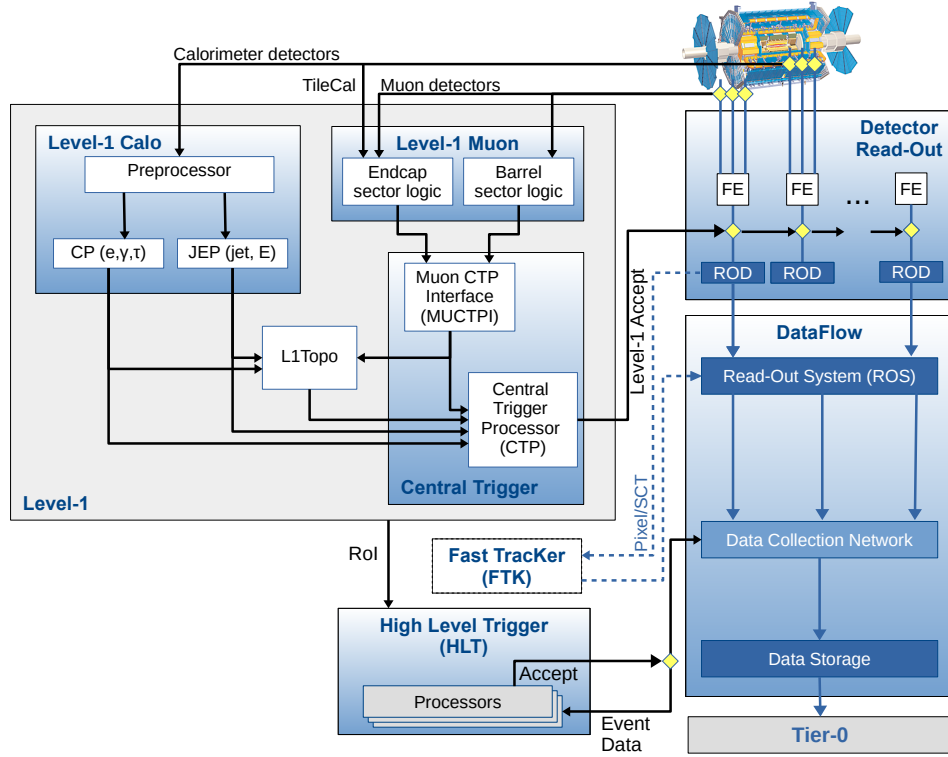


Figure 2.18: The ATLAS TDAQ system during the Run 2 [257]. The Fast Tracker (FTK) system [259] was not used.

Level-1 Trigger: The L1 calorimeter (L1Calo) trigger [260] is a pipelined digital system using custom electronics. As input, the L1Calo trigger works with $\sim 7,200$ analog trigger towers, primarily of granularity $\Delta\eta \times \Delta\phi = 0.1 \times 0.1$, from the electromagnetic and hadronic calorimeters. A pre-processor [261, 262] digitizes the analog signals, associates them with specific LHC bunch crossings using a digital filter, and converts the digital signals into measurements of transverse energy [260]. These measurements are sent in parallel to the Cluster Processor (CP) and the Jet/Energy-sum Processor (JEP), both of which use sliding window algorithms to search for trigger features [263].

The CP system implements the electron/photon and tau trigger algorithms, identifying 2×2 trigger tower clusters as Regions-of-Interest (RoIs) by requiring the sum of the transverse energy for at least one of the four possible nearest neighbour towers, 1×2 or 2×1 , to pass a programmable threshold [258, 260]. As an isolation-veto, the twelve towers surrounding the RoI are summed and required to be less than a programmable threshold. Both the RoI and the isolation-veto clusters are shown schematically in Fig. 2.19a. In order to resolve the ambiguity due to overlapping trigger tower windows, the total sum of the 2×2 core towers, calculated for both the electromagnetic and hadronic calorimeters and added together, must be a local maximum with respect to its eight nearest neighbours [260]. The local maximum test is shown pictorially in Fig. 2.19b.

The JEP system implements the jet trigger algorithm and produces global sums of the missing transverse momentum and the total transverse momentum. Jet elements, the sum of 2×2 trigger towers in the electromagnetic calorimeter added to the sum of 2×2 trigger towers in the hadronic calorimeter, are the basic units of the trigger algorithm. The transverse energy is summed within a window of 2×2 , 3×3 , or 4×4 jet

elements, depending on the desired jet multiplicity, and the total sum is required to be above a programmable threshold [260]. Within each window, a 2×2 region is required to be a local maximum in order to resolve ambiguous hits – this region is identified as the RoI [260]. The different window sizes are shown in Fig. 2.19c.

The L1 muon (L1Muon) trigger [250], like the L1Calo trigger, is a pipelined digital system using custom electronics. The momentum of a muon, estimated using the degree of deviation from the hit pattern expected for a muon with infinite momentum, is required to pass a programmable threshold to yield an accept decision [250]. The region of the detector associated with the trigger is identified as the RoI, typically of granularity $\Delta\eta \times \Delta\phi = 0.1 \times 0.1$ and 0.3×0.3 in the barrel and endcaps, respectively [250].

In the barrel region, $|\eta| < 1.05$, the accept decision is based on the coincidence of hits in two or three RPC barrels for the low- and high- p_T thresholds, respectively [264]. The high- p_T thresholds are used for single-muon signatures, while the low- p_T thresholds are used in conjunction with other triggers for multi-object signatures [250]. A schematic of the barrel trigger is shown in Fig. 2.20a.

In the endcap region, $1.05 < |\eta| < 2.5$, the accept decision is based on the coincidence of hits in the inner and outer TGC endcaps, which reduces the trigger background due to low- p_T charged particles emerging from the endcaps [250]. A coincidence between the TGCs and the tile calorimeter is used to further reject fake muon triggers in the region $1.05 < |\eta| < 1.3$ [250]. A chamber-by-chamber optimization of the coincidence windows is performed to account for the alignment of the detector [250]. Like the barrel trigger, the L1Muon endcap trigger instruments both low- and high- p_T thresholds. A schematic of the endcap trigger is shown in Fig. 2.20b.

To cope with the increased instantaneous luminosity and therefore the increased background rate during the LHC's Run 2, the L1 trigger was updated to include topological information on each trigger object's direction in space in addition to multiplicity [265]. This updated system, the L1 topological (L1Topo) trigger [265, 266], receives energy/momentum and topological information for each trigger object, enabling accept decisions which select on high-level variables such as the invariant mass or angular distance [267].

The L1-Accept (L1A) trigger decision is formed by the Central Trigger Processor (CTP) [268, 269]. The CTP receives inputs from the L1Calo, L1Muon, and L1Topo trigger systems as well as other detector subsystems including the LUCID-2 Cherenkov counters [152], the Minimum Bias Trigger Scintillators² [270], and the Zero-Degree Calorimeters³ [271]. For each L1-accepted event, the front-end data is read out for all subdetectors, transferred to the associated ReadOut Drivers (RODs) for initial processing, and finally buffered in the ReadOut System (ROS) [257]. In parallel, the RoIs identified by the L1 trigger are transferred to the HLT for regional reconstruction. The CTP additionally applies a preventative dead time between subsequent L1As to avoid overlapping readout windows and overflowing front-end buffers [257].

High-Level Trigger: The HLT selects events using trigger chains, each of which consists of a L1 trigger and a sequence of HLT algorithms which reconstruct physics objects, apply kinematic selections, and form the HLT-accept decision [257]. Most HLT algorithms use fast reconstruction for initial event rejection followed by slower precision reconstruction for the final selection [258]. As input, the HLT receives the event data

²The Minimum Bias Trigger Scintillators, which instrument the region $2.08 < |\eta| < 3.75$, provide an efficient and unbiased trigger during low-luminosity runs. They consist of 2 cm thick polystyrene scintillator disks mounted in front of the electromagnetic endcap calorimeters at $z \approx \pm 3.6$ m. The scintillation light is guided via optical fibers to photomultiplier tubes, which measure the arrival time of the signal as well as the total charge collected [270].

³The Zero-Degree Calorimeters, located at $z = \pm 140$ m and $\theta \approx 0^\circ$ ($|\eta| > 8.3$), enhance ATLAS's acceptance for diffractive processes and provide an minimum bias trigger during pp collisions [127]. Each calorimeter consists of one electromagnetic module and three hadronic modules, and each module consists of tungsten absorber plates perforated by an array of quartz rods. Incident particles emit Cherenkov light in the quartz, and the light is subsequently measured by photomultiplier tubes to yield energy and transverse position information with a 100 ps timing resolution [271].

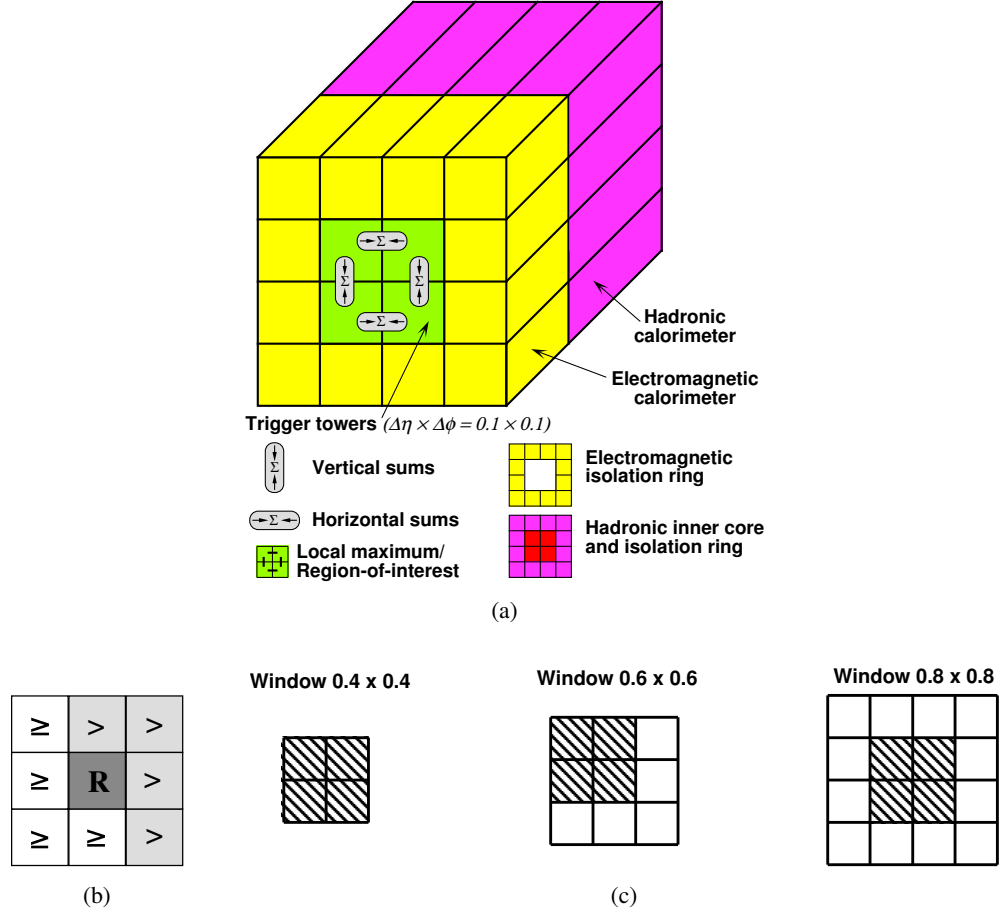


Figure 2.19: (a) A schematic view of the trigger towers used as input to the L1Calo electron/photon and tau trigger algorithms [127]. (b) The local maximum test used for ambiguity solving in the L1Calo electron/photon and tau trigger algorithms [127]. The possibility of equal digital values are accounted for by the “greater-than-or-equal-to” operators. The horizontal and vertical axes correspond to η and ϕ axes, respectively. Each square represents a 2×2 cluster of trigger towers and the square labelled “R” indicates the RoI. (c) Window sizes used by L1Calo jet trigger algorithm [127]. Each square corresponds a jet element and the shaded area indicates the RoI.

measured within the RoIs identified by the L1 trigger [257]. For some trigger chains – for example, those relying on missing transverse momentum – the event data from the full detector is used [267]. Events accepted by the HLT are transferred to the Tier-0 facility at CERN for offline reconstruction [272].

The available trigger chains cover diverse range of physics signatures, including electrons [273], photons [273], muons [250] taus [274], jets [275], b -jets [276], missing transverse momentum [277], total transverse energy [267], and B -mesons [267]. The list of trigger chains used during data-taking constitutes the trigger menu [257]. The main goal of the trigger menu during Run 2 was to maintain single-electron and single-muon triggers with p_T thresholds around 25 GeV to collect the majority of events with leptonic W/Z decays [257]. In order to accommodate increases in the LHC’s instantaneous luminosity, the trigger menu was adjusted yearly over the duration of Run 2 [278–281].

For the $VH, H \rightarrow WW^*$ measurement described in this thesis, the data is selected using only single-electron and single-muon triggers. The precision reconstruction of electron and muon candidates in the HLT closely

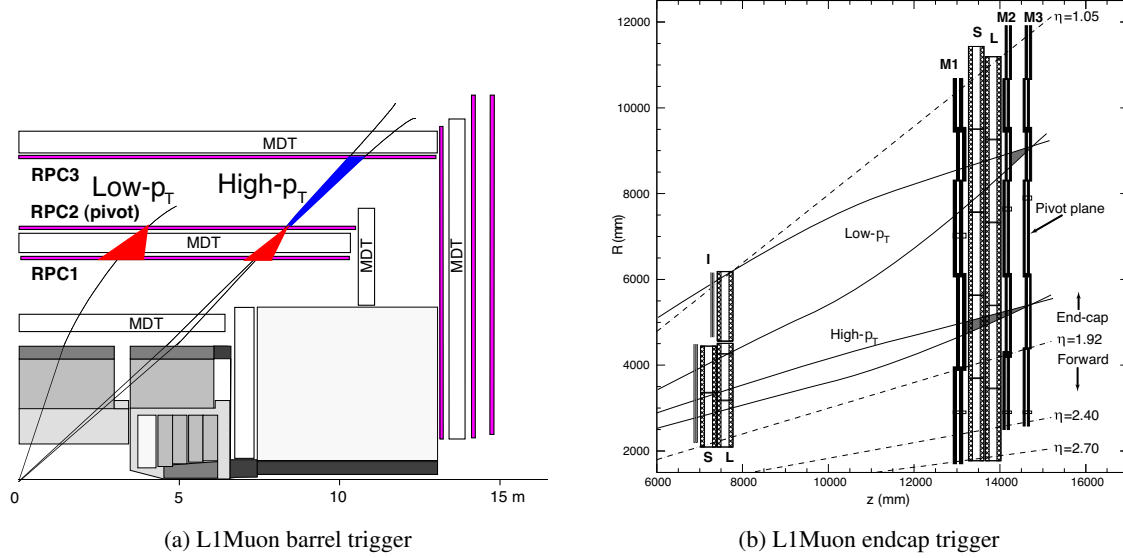


Figure 2.20: A schematic of the L1Muon trigger in the $r-z$ plane for (a) the barrel and (b) the endcaps [127]. Both low- and high- p_T muons as well as their associated coincidence windows are shown. In (a), “RPC1”, “RPC2”, and “RPC3” refer to the three RPC barrel layers. In (b), “I” indicates the TGC doublet of the inner endcap and “M1”, “M2”, and “M3” indicate the single TGC triplet and two TGC doublets of the outer endcap.

follows that used in offline reconstruction. The trigger efficiencies for the single-electron and single-muon triggers used by the $VH, H \rightarrow WW^*$ measurement are shown in Fig. 2.21.

For the single-electron triggers, the fast electron reconstruction proceeds via a cut-based algorithm ($E_T < 15$ GeV) or a neural-network-based (“Ringer”) algorithm ($E_T > 15$ GeV) [273]. The cut-based algorithm builds clusters using calorimeter cells from the electromagnetic calorimeter within the RoI. Each cluster is centered on the most energetic cell of the second sampling layer and is required to be a local maximum in its transverse energy sum. Additionally, each cluster must pass electron identification based on its transverse energy and shower shape parameters [273]. The Ringer algorithm [282] relies on the fact that electromagnetic showers develop laterally in a conic fashion, receiving as input the transverse energy sums for concentric rings centered on the most energetic cell of each sampling layer [273]. The Ringer algorithm ensures optimal fake rejection while maintaining a high trigger efficiency. For both algorithms, electron candidates must be matched to tracks reconstructed by the inner detector within the RoI [283].

The precision electron reconstruction proceeds via a sliding window algorithm operating on a grid of projective towers within a given sampling layer of the electromagnetic calorimeter [284]. Window sizes of $\Delta\eta \times \Delta\phi = 0.075 \times 0.175$ and 0.125×0.125 are used in the barrel and endcaps, respectively [258]. The algorithm scans the grid until the transverse energy contained within the window is found to be a local maximum. If the local maximum is above a set threshold, the window is identified as an electromagnetic cluster. This process is repeated for each sampling layer [258]. A multivariate technique is then used to calibrate the energy of each cluster to the electron energy scale [285], and a likelihood-based discriminant is used to identify electrons [286]. Tracks reconstructed by the inner detector are assigned to the clusters using distance matching, $|\Delta\eta(\text{track}, \text{cluster})| < 0.05$ and $|\Delta\phi(\text{track}, \text{cluster})| < 0.05$ [273]. An isolation criterion is optionally applied by requiring the p_T -sum of the inner detector tracks within a cone of size $\Delta R = 0.2$ centered on and excluding the electron candidate to be less than 10% of the electron candidate’s p_T [273].

For the single-muon triggers, the fast muon reconstruction proceeds by fitting standalone muon candidates within the RoI, using a parameterized function to reconstruct each candidate's transverse momentum [250]. The MDT hits, excluding noise hits, are provided as input to the fit and are complemented by hit information from the endcap MDTs and the CSCs in the regions $1.05 < |\eta| < 1.35$ and $2.0 < |\eta| < 2.4$, respectively [250]. The standalone muon tracks are back-extrapolated to the interaction point and combined with inner detector tracks to produce combined muon candidates with refined track parameter resolutions [250].

The precision muon reconstruction proceeds by building combined muon candidates in a manner analogous to fast reconstruction; however, if no combined muons are found, they are searched for by extrapolating tracks from the inner detector to the muon spectrometer [250]. The transverse momentum of the muon candidate is required to be above a set threshold. Additionally, an isolation criterion is applied by requiring the p_T -sum of the inner detector tracks within a longitudinal distance of 6 or 2 mm of the primary vertex to be less than 6 or 7% of the muon candidate's p_T [250]. The muon track is excluded from the p_T -sum, and the particular choice of distance and fractional p_T cut varies between data-taking years. To recover missing muons due to L1 trigger inefficiencies, the muon candidates are searched for in all muon detectors and provided to the precision reconstruction algorithm as additional RoIs [250].

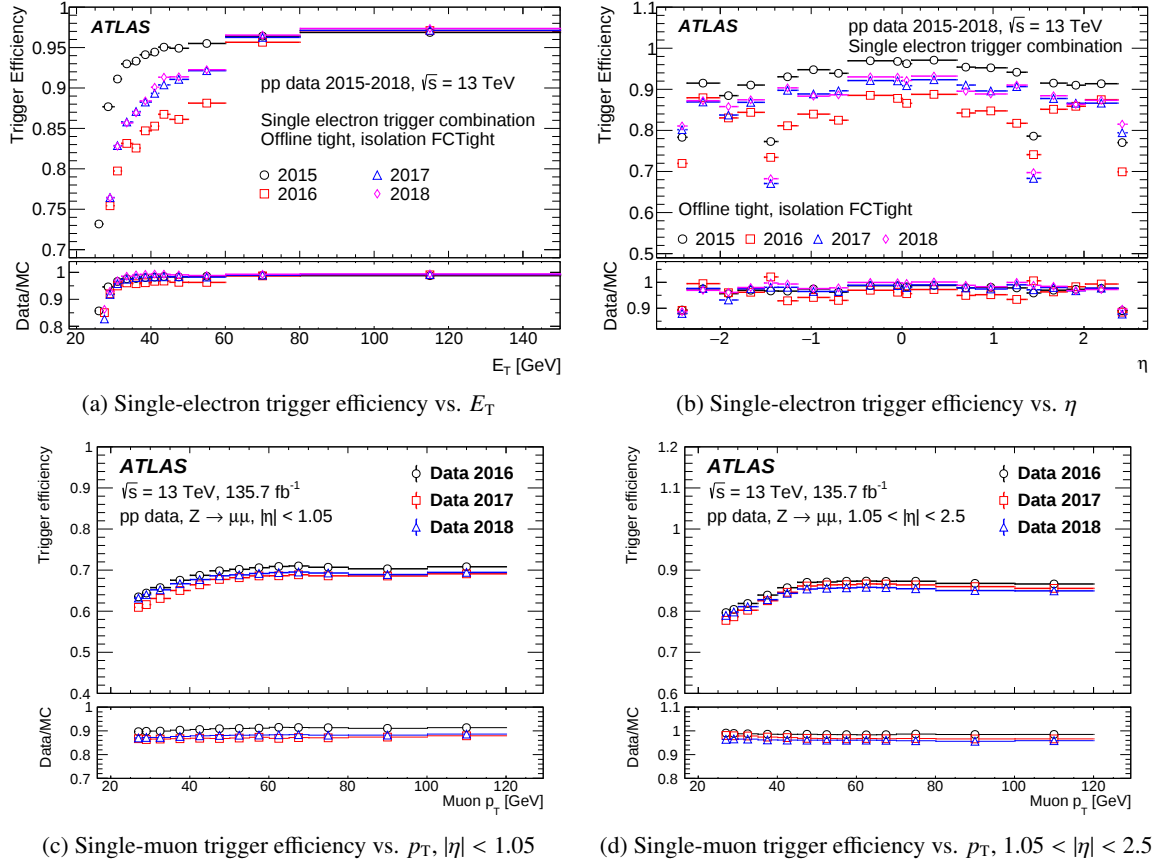


Figure 2.21: Single-electron [273] and single-muon [250] trigger efficiencies in Run 2 data and simulation. In (a) and (b), the efficiencies correspond to the logical OR of all single-electron triggers and are obtained using $Z \rightarrow e\bar{e}$ events. The efficiencies are plotted as a function of offline electron E_T and η . The errors bars include both statistical and systematic uncertainties. In (c) and (d), the efficiencies correspond to the logical OR of the single-muon triggers with p_T thresholds of 26 and 50 GeV and are obtained using $Z \rightarrow \mu\bar{\mu}$ events. The efficiencies are plotted as a function of offline muon p_T . The error bars indicate the statistical uncertainty.

Chapter 3

Monte Carlo Simulation

In high-energy scattering experiments, Monte Carlo (MC) simulation is used to make predictions for the rates and kinematics of the physics processes involved, incorporating theoretical and experimental knowledge of the phenomenology of particle physics with computational tools. The programs used to make these predictions are known as MC event generators. Examples of commonly used MC event generators include POWHEG Box [287–289], PYTHIA [290], SHERPA [291], MADGRAPH5_aMC@NLO [292], and HERWIG [293, 294].

MC event generation corresponds to a random sampling of the underlying probability density function governing the physics. Each sampling is realized as an exclusive event. To allow for direct comparisons with data, the effect of each event on the ATLAS detector is simulated and digitized, allowing physics objects to be reconstructed in the same way as collision data. The simulation and digitization is typically performed using the GEANT4 toolkit [295–297].

This chapter provides an overview of the MC simulation used to model proton-proton (pp) collisions at $\sqrt{s} = 13$ TeV (Section 3.1) as well as describes the simulated samples used as inputs to the $VH, H \rightarrow WW^*$ analysis (Section 3.2).

3.1 Event Generation and Detector Simulation

MC simulation is a stochastic simulation of fully exclusive collision events. It seeks to factorize the short-range (high- q^2) and long-range (low- q^2) physics into a series of discrete steps. These steps correspond to the matrix element calculation, parton showering, hadronization, and finally detector simulation.

3.1.1 Matrix Element Calculation

Physics predictions at hadron colliders are complicated by the fact that hadrons are composite particles – any of the quarks or gluons inside the hadrons may interact with one another during collisions, each carrying an indeterminate amount of the parent hadron’s momentum. It is necessary, then, to integrate over the possible interacting constituent particles as well as over the momentum fractions carried by those constituents. These constituents are referred to as partons [298, 299], and the momentum fraction x of a parent hadron h carried by a constituent parton p is described by a probability density function known as a parton distribution function (PDF), $f_{p/h}$. PDFs are not known *a priori* and must be measured experimentally. Examples of commonly used PDFs include the NNPDF3.0 [300] and PDF4LHC15 [301] PDF sets, and proton PDFs for various flavours of partons are shown in Fig. 3.1.

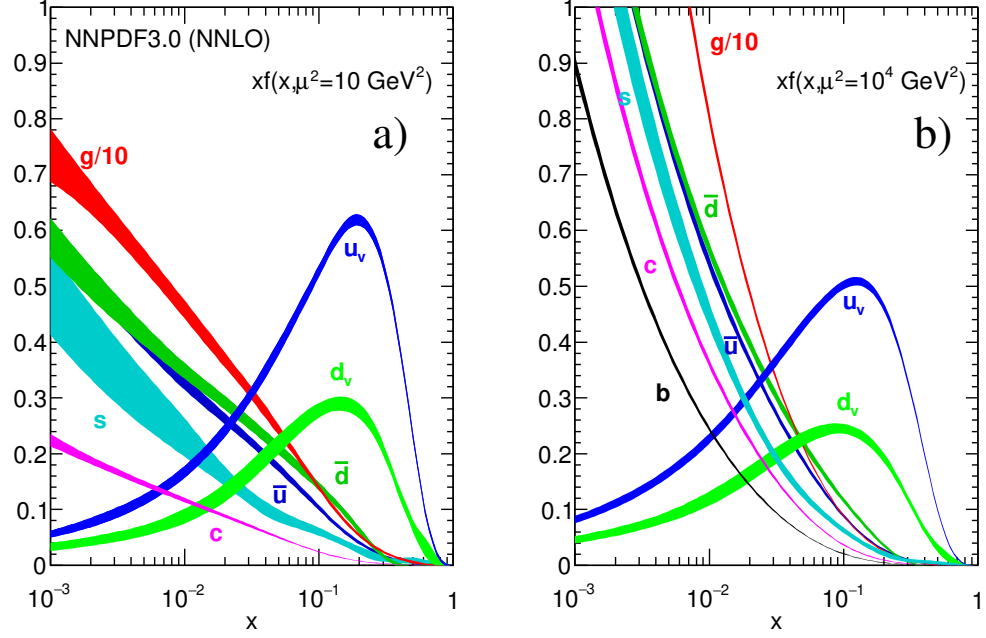


Figure 3.1: Proton PDFs $xf(x, q^2)$ for various flavours of partons from the NNPDF3.0 PDF set [300]. The PDFs are plotted as a function of the momentum fraction x for two different values of the four-momentum transfer q^2 , 10 GeV^2 (left) and 10^4 GeV^2 (right). Figure obtained from the Particle Data Group [26].

The pp collision with the highest four-momentum transfer, the hard-scatter collision, consists two initial state partons scattering to produce n final state particles. The cross section for the hard-scatter collision, $\sigma_{pp \rightarrow n}$, is given by [302, 303]:

$$\sigma_{pp \rightarrow n}(s, \mu_R^2, \mu_F^2) = \underbrace{\sum_i \sum_j}_{(a)} \underbrace{\int_0^1 \int_0^1 dx_i dx_j f_i(x_i, \mu_F^2) f_j(x_j, \mu_F^2) \hat{\sigma}_{ij \rightarrow n}(x_i, x_j, s, \mu_R^2, \mu_F^2)}_{(b)} \quad (3.1)$$

where the partonic cross section $\hat{\sigma}_{ij \rightarrow n}$ is given by:

$$\hat{\sigma}_{ij \rightarrow n}(x_i, x_j, s, \mu_R^2, \mu_F^2) = \underbrace{\frac{1}{2x_i x_j s}}_{(c)} \underbrace{\prod_{k=1}^n \left(\int \frac{d^4 p_k}{(2\pi)^4} \right)}_{(d)} \underbrace{|\mathcal{M}_{ij \rightarrow n}|^2(\phi_1, \dots, \phi_n, \mu_R^2, \mu_F^2)}_{(e)}, \quad (3.2)$$

Eqs. 3.1 and 3.2 may be broken down into several logical blocks:

- (a) The summation over the initial state parton flavours, i and j , for each proton;
- (b) The integration over the proton PDFs for parton flavours i and j , $f_i(x_i, \mu_F^2)$ and $f_j(x_j, \mu_F^2)$, respectively. Here, μ_F corresponds to the factorization scale, which is introduced to cure infrared divergences due to collinear and soft parton emissions;
- (c) The parton flux, where x_i and x_j are the momentum fractions carried by partons i and j , respectively, and s is the centre-of-mass energy squared;
- (d) The product of the Lorentz-invariant phase space integrals for all outgoing particles, $1, \dots, n$;

- (e) The Lorentz-invariant matrix element squared for the scattering of initial state partons i and j , which is a function of the final-state phase space $\phi_1, \dots, \phi_n$, the renormalization scale μ_R , and the factorization scale μ_F . The renormalization scale is introduced to cure ultraviolet divergences due to loop integrals.

The partonic cross section, Eq. 3.2, is evaluated through a perturbative expansion in the strong coupling constant α_S :

$$\hat{\sigma} \approx \sum_{i=0}^n \alpha_S^{i+1} \hat{\sigma}_i, \quad (3.3)$$

where $n = 0$ denotes the leading-order calculation, $n = 1$ denotes the next-to-leading-order calculation, and so on. For four-momentum transfers typical of the hard-scatter collision, $|q| \sim \mathcal{O}(0.1-1)$ TeV, the strong coupling constant is small, $\alpha_S \sim \mathcal{O}(0.1)$, and so the perturbative expansion is valid. However, during the evaluation of the contributions or “diagrams” to the partonic cross section at fixed-order, one will encounter multidimensional phase space integrals which are impossible to solve exactly – instead, MC integration algorithms are used for their fast convergence in many dimensions, ability to integrate arbitrary regions, and straightforward error estimates [304]. The integration proceeds by generating sets of final state particles whose kinematics are drawn from a uniform distribution. Sets of final state particles accepted by the algorithm correspond to exclusive events which are kinematically consistent with the physics process being simulated.

A problematic aspect of the perturbative expansion is the appearance of ultraviolet divergences in the matrix element calculation due to loop integrals. These divergences are renormalized using dimensional regularization [305, 306] to yield well defined poles and minimal subtraction [307, 308] to cancel the poles using counter terms added to the Lagrangian. In using dimensional regularization, a renormalization scale μ_R is introduced; in using minimal subtraction, a dependence of the strong coupling constant on the four-momentum scale q^2 is introduced such that:

$$\alpha_S(q^2) = \alpha_S(\mu_R^2) \times \left(1 + b_0 \alpha_S(\mu_R^2) \log \left(\frac{q^2}{\mu_R^2} \right) \right)^{-1}, \quad (3.4)$$

where b_0 is just a number. By including counter terms for all 1-loop diagrams in the perturbative expansion, it can be shown that b_0 is *greater* than 0 and therefore α_S *decreases* with increasing q^2 . In other words, QCD is *asymptotically free* [31, 32]. The so-called running of α_S is shown in Fig. 3.2. Asymptotic freedom justifies the treatment of quarks as quasi-free particles in high- q^2 inelastic scattering, as would be the case for pp collisions at $\sqrt{s} = 13$ TeV.

Normalization: Each exclusive event produced by an MC generator has an associated event weight, w_i , which corresponds to the differential cross section of the point in phase space represented by the event [309]. The sum-of-weights per physics process is normalized to the predicted number of events, $N_{\text{predicted}}$, for that same process:

$$w_i \rightarrow w'_i = \frac{N_{\text{predicted}}}{\sum_i w_i} w_i = \frac{\mathcal{L}\sigma}{\sum_i w_i} w_i \quad \forall i, \quad (3.5)$$

where $\mathcal{L}$ is the integrated luminosity and σ is the predicted cross section for the process of interest. The number of generated MC events is typically very large, $\mathcal{O}(10^5-10^7)$, in order to yield statistically significant results. In particular, the statistical uncertainty on the predicted number of events, $\sigma_{N_{\text{predicted}}}^{\text{stat}}$, follows weighted Poisson counting statistics:

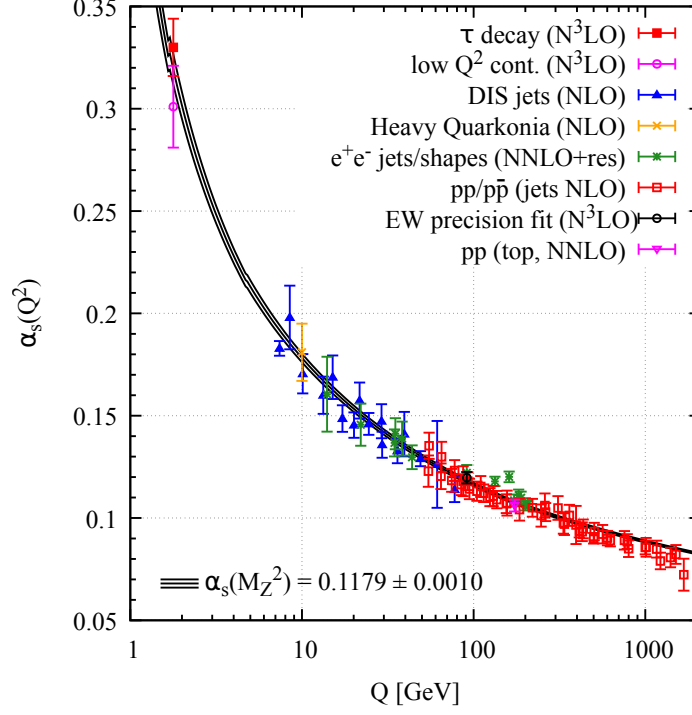


Figure 3.2: Measurements of the strong coupling constant α_s plotted as a function of the four-momentum scale $|q|$. The accuracy in QCD used for the extraction of α_s for each measurement is indicated in the brackets. Figure obtained from the Particle Data Group [26].

$$\sigma_{N_{\text{predicted}}}^{\text{stat}} = \sqrt{\sum_i (w'_i)^2} = \frac{\mathcal{L}\sigma}{\sum_i w_i} \sqrt{\sum_i w_i^2}, \quad (3.6)$$

calculated per bin for differential distributions.

3.1.2 Parton Showering and Jets

Point-like quarks and gluons at high four-momentum scales are what remain at the end of the matrix element calculation. Parton showering describes the evolution of these partons from the high scales, $|q| \sim \mathcal{O}(0.1\text{--}1)$ TeV, to the low scales, $|q| \sim \mathcal{O}(1)$ GeV. It proceeds by the repeated emission of gluon radiation from the incoming and outgoing partons [302, 303], resulting in collimated sprays of hadronic particles known as jets. A key assumption of parton showering is that it is *independent* of the initial hard-scattering process.

At the crux of parton showering is probability of a parton emission occurring. In particular, the differential probability of a parton i splitting into partons j and k , $d\mathcal{P}_{i \rightarrow jk}$, is given by:

$$d\mathcal{P}_{i \rightarrow jk}(q^2, z) = \frac{\alpha_s}{2\pi} \frac{dq^2}{q^2} P_{i \rightarrow jk}(z),^1 \quad (3.7)$$

where q^2 is the four-momentum transfer of parton i , z is the momentum fraction of parton i carried by parton j , and $P_{i \rightarrow jk}$ is a universal splitting kernel for the $i \rightarrow jk$ emission [303]. The functional behaviour of the splitting kernel is dependent on the flavour of the emission: $q \rightarrow qg$, $q \rightarrow gq$, $g \rightarrow qq$, or $g \rightarrow gg$, where q and g correspond to quarks and gluons, respectively. Eq. 3.7 is infrared divergent in two distinct ways: the $q^2 \rightarrow 0$

¹N.B. in Eqs. 3.7 and 3.8, the flavour of parton k is entirely determined by the flavours of partons i and j .

or “collinear emission” limit, which corresponds to a parton emission at an angle $\theta \rightarrow 0$ with respect to the parent parton, and the $z \rightarrow 0$ or “soft gluon emission” limit, which is implicit in the $P_{g \rightarrow gg}$ and $P_{q \rightarrow gq}$ splitting functions due to factors of $1/z$ [310].

Parton showering only concerns *resolved* emissions as physical measurements cannot distinguish between two collinear partons and a single parton with the same total momentum [303]. As a result, the collinear and soft divergences are cured by introducing a cutoff scale, $Q_0 \sim O(1) \text{ GeV}$. Two partons are resolved if their relative p_T is greater than Q_0 and unresolved otherwise. Equipped with this cutoff scale, parton showering proceeds by calculating the Sudakov form factor [311], $\mathcal{P}_i(q^2, Q_0^2)$, which corresponds to the probability of a parton with four-momentum transfer q^2 producing no resolved emissions in the range $Q_0^2 < k^2 < q^2$ [303]:

$$\mathcal{P}_i(q^2, Q_0^2) = \exp \left(-\frac{\alpha_S}{2\pi} \int_{Q_0^2}^{q^2} \frac{dk^2}{k^2} \int_{Q_0^2/k^2}^{1-Q_0^2/k^2} dz \sum_{j \in [q, g]} P_{i \rightarrow jk}(z) \right). \quad (3.8)$$

The second integral’s limits ensure the momentum fraction of the emitted parton z respects the first integral’s limits. A random number ρ may be drawn from a uniform distribution between 0 and 1 to decide if an emission occurs. If $\rho > \mathcal{P}_i(q^2, Q_0^2)$, an emission occurs and the momentum fraction of the emitted parton z may be chosen according to the splitting kernels $P_{i \rightarrow jk}(z)$. The process repeats by calculating the Sudakov form factors for each daughter parton. If $\rho < \mathcal{P}_i(q^2, Q_0^2)$, the shower terminates for the current parton [303].

The same collinear and soft singularities arise when studying additional radiation emitted by initial state partons. The solution is to regularize the associated phase space integral by introducing a factorization scale μ_F [302]. The divergent behaviour is subsequently absorbed into the PDFs, which become “running” functions of μ_F whose evolution is governed by the Dokshitzer-Gribov-Lipatov-Altarelli-Parisi (DGLAP) equations [312–315]:

$$\frac{df_i(x, \mu_F^2)}{d \log \mu_F^2} = \frac{\alpha_S}{2\pi} \int_x^1 \frac{dy}{y} \sum_{j \in [q, g]} P_{j \rightarrow ik}(x/y) f_j(y, \mu_F^2) \quad \forall i \in [q, g]. \quad (3.9)$$

The summation includes all partons j which can split into parton i , and the integration is such that only partons with momentum fraction $y > x$ may split into a parton with momentum fraction x . The flow of the proton PDFs in $x - q^2$ space, so-called “Bjorken scaling”, is shown in Fig. 3.3.

The fixed-order matrix elements and parton showers complement one another: the matrix elements best simulate well-separated and hard partons, while the parton showers best simulate soft or collinear partons [302, 303]. However, matrix elements are *inclusive* in the number of partons at fixed-order in α_S and parton showers are *exclusive* in the number of partons to all orders in α_S . To prevent undercounting or overcounting the number of partons in different regions of phase space, jet matching [316] and merging [317] procedures are used. With matching, the parton shower expression is computed at fixed-order and subtracted from the higher-order matrix element, allowing the subtracted result to be processed by the parton shower with no double counting [318]. Commonly used matching methods include aMC@NLO [319] and PowHEG [287, 288]. With merging, a separate tree-level calculation is performed for each parton multiplicity of interest, using cutoff scales to regulate the soft and collinear singularities. The parton shower is then combined with these calculations, using vetoes on shower branchings to prevent double counting [318]. Commonly used matching methods include the Catani-Kuhn-Krauss-Webber (CKKW) [317] and MEPS@NLO [320, 321] methods.

As a final note, the integration of the soft and collinear divergences can lead to potentially large logarithmic enhancements to the partonic cross section, whose perturbative expansion to all orders now takes the form:

$$\hat{\sigma} \simeq \sum_{n=0}^{\infty} \sum_{m=0}^{2n} \alpha_S^n \ln^m \left(\frac{|q|}{Q_0} \right) \hat{\sigma}_{nm}. \quad (3.10)$$

If $\alpha_S \ln^2(|q|/Q_0) \gg 1$, a resummation of the dominant logarithmically-enhanced terms to all orders in α_S is required [322], where summing the terms with $m = 2n$ corresponds to leading-logarithm accuracy, summing the terms with $m = 2n - 1$ corresponds to next-to-leading-logarithm accuracy, and so on [26].

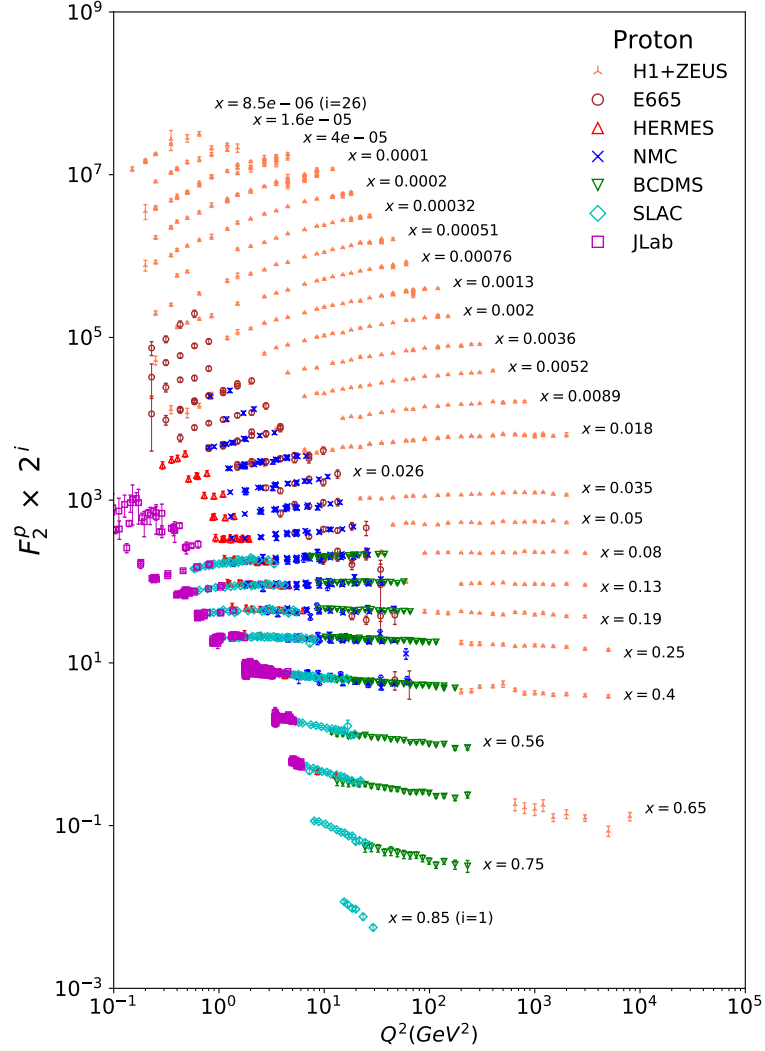


Figure 3.3: Measurements of the proton structure function F_2^p plotted as a function of the momentum fraction x and the four-momentum transfer q^2 . The measurements come from a variety of collision and fixed-target experiments [323–333]. The structure function is defined in terms of the convolution integral of the proton PDFs for each flavour of parton and summed over all parton flavours. For clarity, F_2^p is multiplied by a factor of 2^{i_x} , where i_x is the index of the bin in momentum fraction x . Figure obtained from the Particle Data Group [26].

3.1.3 Hadronization and Decays

Hadronization or jet fragmentation describes the process of producing colour-singlet final states using the quarks and gluons left-over from the hard-scatter collision and the parton shower [334]. These partons exist at the scale of non-perturbative QCD, $|q| \sim \mathcal{O}(1)$ GeV, and so the description of the physics relies on phenomenological models. The two models in use today are the string and cluster models.

The (Lund) string model [335–337] is based on the observation that the quark-antiquark potential is linear, $V(r) \approx (1 \text{ GeV/fm}) \times r$ [334]. A Lorentz-covariant and causal treatment of energy flow due to this potential is achieved by treating the longitudinal axis of the quark-antiquark system as a massless relativistic string [338], with each quark corresponding to an endpoint of the string. As the quark and the antiquark move apart at the speed of light, the stored potential energy can lead to the creation of a new quark-antiquark pair (or diquark-antidiquark pair, in the case of baryon production) by introducing a “break” in the string [334]. This process may be repeated until no further breaks are possible. In the string model, gluons are treated as “kinks” on the string and assigned a sum of one colour charge and one anticolour charge [303, 339]. A benefit of the string model is that the hadronization proceeds unaffected by the emission of collinear or soft gluons [340]. PYTHIA [290] is the only MC event generator which applies the string model to hadronization.

The cluster model [341] is based on the preconfinement property of QCD: colour-singlet combinations of partons or “clusters” may be formed at any point in the parton shower and the invariant mass distribution of the clusters is universal [342]. That is, the invariant mass distribution only depends on the cutoff scale of the parton shower, $Q_0 \sim \mathcal{O}(1)$ GeV, and the scale at which QCD becomes non-perturbative, $\Lambda_{\text{QCD}} \approx 300 \text{ MeV} \sim 1 \text{ fm}^{-1}$ [303, 334]. Following parton showering, gluons undergo non-perturbative splitting into quark-antiquark (or diquark-antidiquark) pairs [343, 344]. Clusters with an invariant mass less than 4 GeV or “light” clusters, formed from colour-connected pairs, transform into hadrons through two-body decays. In contrast, heavy clusters must first undergo non-perturbative splitting into two light clusters [303, 334]. This process repeats until there are no remaining clusters. SHERPA [291] and HERWIG [293, 294] both apply the cluster model to hadronization.

The string and cluster models both offer a good description of hadronization. The cluster model has fewer free parameters than the string model, but the string model has better modelling of the physics [303]. The parameters of a hadronization model are typically not known *a priori* and must be measured using data. A particular choice for the parameters is known as a tune. Examples of tunes used in ATLAS include the AZNLO [345], A14 [346], and A3 [347] tunes.

3.1.4 Detector Simulation

The final step in the simulation chain is simulating the physics interactions of the final state particles with the ATLAS detector and simulating the response of the ATLAS detector to those interactions [348, 349]. The latter step is known as digitization. To be considered by the detector simulation, the final state particles must be stable, $c\tau > 10 \text{ mm}$, or decayed into stable secondaries, where c is the speed of light and τ is the proper lifetime of the particle [26, 349].

To simulate the physics interactions of the final state particles with the ATLAS detector, the GEANT4 toolkit is used [295–297]. A comprehensive model of the ATLAS detector is provided as input. This model details the geometry and the granularity of the ATLAS detector as well as encodes the materials which comprise the sensitive detector elements, the shielding, and the support structures [349]. Such a description is necessary for accurately modelling the detector’s response and its efficiencies [349] – for example, its reconstruction

and identification efficiencies for each type of physics object. GEANT4 uses a fourth-order Runge-Kutta method to solve the equations of motion governing the trajectory of each particle in ATLAS’s magnetic field, which is provided to GEANT4 as a 3D map [295, 349]. As a particle is transported through the detector’s geometry, GEANT4 simulates the relevant physical interactions in each material through which the particle passes including [295]:

- Decay processes, and specifically those of pions and short-lived hadrons;
- Electromagnetic processes, including Compton scattering, pair production, the photoelectric effect, bremsstrahlung, ionization, and electron-positron annihilation;
- Hadronic processes, including inelastic scattering, capture of neutral particles, induced fission, and elastic scattering;
- Optical processes, including Cherenkov radiation, scintillation, absorption and Rayleigh scattering, reflection and refraction, and transition radiation.

The physical interaction of a particle with a sensitive detector element – for example, its energy deposition in that element – is tracked by GEANT4 as a “hit” [295]. The conversion of a hit into a “digit” or digitization – for example, the voltage or current over threshold induced by the hit – is performed by ATLAS’s offline software framework, Athena [350]. The digitization is optimized such that only the necessary information for reproducing the performance of each of ATLAS’s subsystems is stored [349]. At the end of digitization, the simulated digits are equivalent to those measured during actual operations. A simulated event display is shown in Fig. 3.4.

At this point, it is useful to make a distinction between the “particle-level” event and the “reconstruction-level” event. The particle-level event corresponds to the state of a simulated event after hadronization but before detector simulation. In the particle-level event, all of the final state particles are required to be stable. The reconstruction-level event corresponds to the state of a real or simulated event after detector digitization and reconstruction. It is assumed that a quantity or distribution is given at the reconstruction-level unless specified otherwise.

3.2 Simulated Samples

The MC generators used to model the signal and background processes for the $VH, H \rightarrow WW^*$ measurement are listed in Table 3.1. In general, POWHEG is used to generate Higgs and top samples, while SHERPA is used to generate Z +jets and multiboson samples. PYTHIA 8 and HERWIG 7 are only used to model the parton shower and non-perturbative physics.

As a result of the LHC’s high instantaneous luminosity, typically more than one pp collision occurs per bunch crossing: there is a single hard-scatter collision accompanied by any number of low- p_T collisions known as pile-up. Pile-up, which is part of the larger class of soft processes known as the underlying event, affects the reconstruction of the hard-scatter collision and therefore must be included in the event modelling [352]. Fig. 3.5 shows the mean number of interactions per bunch cross measured by ATLAS during the LHC’s Run 2. There are ~ 30 pile-up collisions per bunch crossing, on average.

For all MC samples, the effect of pile-up is modelled by overlaying simulated inelastic pp events onto the hard-scatter collision. The inelastic pp events are generated with PYTHIA 8.186 [353] using the

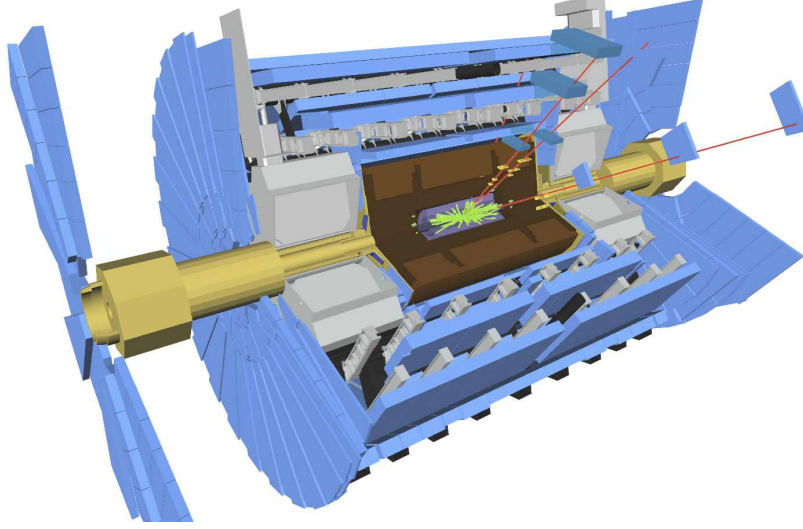


Figure 3.4: A simulated VP1 [351] event display of a $H \rightarrow ZZ^* \rightarrow 4\mu$ decay [349]. The muons are shown in red, the inner detector tracks are shown in green, and the calorimeter energy deposits are shown in yellow.

NNPDF2.3Lo [300] PDF set and the A3 tune [347]. A reweighting procedure is applied to the pile-up distribution in MC to better match that observed in data. Additionally, a scale factor of 1.03^{-1} is applied to data prior to reweighting. This is done because the MC simulation for a given pile-up best models data with a higher pile-up.

Table 3.1: Overview of simulation tools used to generate the signal and background processes in the VH , $H \rightarrow WW^*$ measurement. Also summarized are the corresponding PDF sets as well as the models used for the underlying event/parton shower (UE/PS) simulation. When higher-orders are specified, QCD is implied if not explicitly stated.

Process	Generator	PDF set	UE/PS model	Accuracy
$q\bar{q}/qg \rightarrow VH$	PowHEG Box v2 [287–289, 354]	PDF4LHC15NLO [301]	PYTHIA 8.212 [290]	NNLO QCD + NLO EW [355–359]
$gg \rightarrow ZH$	PowHEG Box v2 [287–289, 354]	PDF4LHC15NLO	PYTHIA 8.212	NLO + NLL [360, 361]
ggF	PowHEG Box v2 [287–289, 362, 363]	PDF4LHC15NNLO [301]	PYTHIA 8.212	N ³ LO QCD + NLO EW [97, 364–373]
VBF	PowHEG Box v2 [287–289, 374]	PDF4LHC15NNLO	PYTHIA 8.230 [290]	NNLO QCD + NLO EW [375–377]
$t\bar{t}H$	PowHEG Box v2 [287–289, 378, 379]	NNPDF3.0NLO [300]	PYTHIA 8.230	NLO QCD + NLO EW [97]
$q\bar{q} \rightarrow VV$	SHERPA 2.2.2 [291]	NNPDF3.0NNLO [300]	SHERPA 2.2.2 [317, 321, 334, 339, 380, 381]	NLO [382–384]
$gg \rightarrow VV$	SHERPA 2.2.2	NNPDF3.0NNLO	SHERPA 2.2.2	NLO [385, 386]
VVV	SHERPA 2.2.2	NNPDF3.0NNLO	SHERPA 2.2.2	NLO [382–384]
$V+\gamma$	SHERPA 2.2.8 [291]	NNPDF3.0NNLO	SHERPA 2.2.8 [317, 321, 334, 339, 380, 381]	NLO [382, 383, 387]
$q\bar{q} \rightarrow WWq\bar{q}$	MadGraph5_aMC@NLO [292]	NNPDF3.0NLO	PYTHIA 8.244 [290]	LO
$V+\text{jets}$	SHERPA 2.2.1 [291]	NNPDF3.0NNLO	SHERPA 2.2.1 [317, 321, 334, 339, 380, 381]	NNLO [388]
$t\bar{t}$	PowHEG Box v2 [287–289, 378]	NNPDF3.0NLO	PYTHIA 8.230	NNLO + NNLL [389–395]
Wt	PowHEG Box v2 [287–289, 396]	NNPDF3.0NLO	PYTHIA 8.230	NLO [397, 398]
$t\bar{t}V$	MadGraph5_aMC@NLO 2.3.3 [292]	NNPDF3.0NLO	PYTHIA 8.212	NLO QCD + NLO EW [97]
tZ	MadGraph5_aMC@NLO 2.3.3	NNPDF3.0NLO	PYTHIA 8.212	NLO
tWZ	MadGraph5_aMC@NLO 2.3.3	NNPDF3.0NLO	PYTHIA 8.212	NLO

3.2.1 Higgs Samples

Higgs boson production in association with a vector boson (VH) was simulated using PowHEG Box v2 [287–289, 354]. The PowHEG prediction is accurate to next-to-leading-order (NLO) in QCD for VH plus one-jet production. The loop-induced $gg \rightarrow ZH$ process was generated separately at leading-order (LO) accuracy. The MC prediction was normalized to a cross section calculated at next-to-next-to-leading-order (NNLO) accuracy in QCD with NLO electroweak corrections for $q\bar{q}/qg \rightarrow VH$ [355–359] and at NLO and next-to-leading-logarithm (NLL) accuracy in QCD for $gg \rightarrow ZH$ [360, 361].

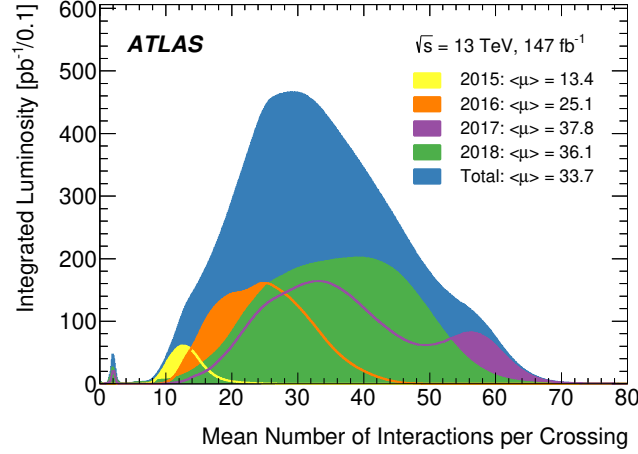


Figure 3.5: Luminosity-weighted distribution of the mean number of interactions per bunch crossing, μ , measured by ATLAS using pp collisions at $\sqrt{s} = 13$ TeV [155]. Here, μ corresponds to the mean of the Poisson distribution of the number of interactions per bunch crossing and calculated for each bunch. Also shown is the average μ , $\langle\mu\rangle$, for each year.

Higgs boson production via gluon-gluon fusion (ggF) was simulated at NNLO accuracy in QCD using PowHEG Box v2 [287–289, 362, 363]. The ggF prediction from the MC samples was normalized to a cross section calculated at next-to-next-to-next-to-leading-order (N^3 LO) accuracy in QCD with NLO electroweak corrections [97, 364–373].

Higgs boson production via vector boson fusion (VBF) was simulated using PowHEG Box v2 [287–289, 374]. The PowHEG prediction is accurate to NLO in QCD. It was tuned to match calculations with effects due to finite heavy-quark masses and soft-gluon resummation up to next-to-next-to-leading-logarithm (NNLL) accuracy. The MC prediction was normalized to an approximate NNLO-accurate QCD cross section with NLO electroweak corrections [375–377].

For VH , ggF, and VBF production, the PDF4LHC15_{NLO} (VH) or PDF4LHC15_{NNLO} (ggF and VBF) [301] PDF sets were used. The generators were interfaced with PYTHIA 8.212 (VH and ggF) or 8.230 (VBF) [290] for parton showering and hadronization, with parameters set according to the AZNLO tune [345]. Additionally, the normalization for each sample accounts for the branching ratio for decay calculated with HDECAY [399–401] and PROPHECY4F [402–404].

Higgs boson production in association with a top quark pair ($t\bar{t}H$) was simulated using PowHEG Box v2 [287–289, 378, 379] at NLO accuracy in QCD and the NNPDF3.0_{NLO} [300] PDF set. The generator was interfaced with PYTHIA 8.230 using the A14 tune [346] to model the parton shower, hadronization, and underlying event. The decays of bottom and charm hadrons were simulated using EVTGEN 1.6.0 [405]. The MC prediction was normalized to a cross section calculated at NLO accuracy in QCD with NLO electroweak corrections [97].

All of the Higgs boson cross sections and branching ratios are calculated assuming a Higgs boson mass of 125.09 GeV [90].

3.2.2 Multiboson Samples

The production of diboson (VV) final states was simulated with SHERPA 2.2.2. Fully leptonic final states were generated using matrix elements at NLO accuracy in QCD for up to one additional parton and at LO accuracy for up to three additional parton emissions. Samples for the loop-induced processes ($gg \rightarrow VV$) were generated

using matrix elements at LO accuracy for up to one additional parton emission.

The production of on-shell triboson (VVV) final states was simulated with SHERPA 2.2.2 using factorized gauge-boson decays. The matrix elements are accurate to NLO in QCD for the inclusive process and to LO for up to two additional parton emissions.

The production of $V+\gamma$ final states was simulated with SHERPA 2.2.8 [291]. The matrix elements are accurate to NLO in QCD for up to one additional parton and to LO for up to three additional parton emissions.

The matrix element calculations were matched and merged with the SHERPA parton shower [380, 381] using the MEPS@NLO prescription [317, 321, 334, 339] and the set of tuned parameters developed by the SHERPA authors. The virtual QCD corrections for matrix elements at NLO accuracy were provided by the OPENLOOPS (VV and VVV) or OPENLOOPS 2 ($V+\gamma$) libraries [382–384, 387]. All of the samples were generated using the NNPDF3.0NNLO [300] PDF set.

Electroweak WW production in association with two jets ($WWq\bar{q}$) was generated by MADGRAPH5_aMC@NLO [292] with LO-accurate matrix elements [385] and using the NNPDF3.0NNLO [300] PDF set. The events were interfaced to PYTHIA 8.244 [290] using the A14 tune to model the parton shower, hadronization, and underlying event.

3.2.3 V +jets Samples

The production of V +jets was simulated with SHERPA 2.2.1 [291] using matrix elements at NLO accuracy in QCD for up to two partons and at LO accuracy for up to four partons, calculated with the COMIX [381] and OPENLOOPS 2 libraries [382, 383, 387]. They were matched and merged with the SHERPA parton shower [380, 381] using the MEPS@NLO prescription [317, 321, 334, 339] and the set of tuned parameters developed by the SHERPA authors. The samples were generated using the NNPDF3.0NNLO PDF set, and they were normalized to a cross section calculated at NNLO accuracy in QCD [388].

3.2.4 Top Samples

The production of $t\bar{t}$ dilepton events was modelled using POWHEG Box v2 [287–289, 378] at NLO accuracy in QCD with the h_{damp} parameter² set to $1.5 \times m_{\text{top}}$ [406]. In order to correct for a known mismodelling of the leading lepton p_T due to missing higher-order corrections, an NNLO reweighting procedure is performed [407]. The MC prediction was normalized to a cross section calculated at NNLO accuracy in QCD including the resummation of NNLL soft gluon terms calculated using TOP++ 2.0 [389–395].

The associated production of top quarks with W bosons (Wt) was modelled using POWHEG Box v2 [287–289, 396] at NLO accuracy in QCD using the five-flavour scheme and the NNPDF3.0NNLO PDF set.

The production of top quark pairs with vector bosons ($t\bar{t}V$), the production single top quarks with Z bosons (tZ), and the production of single top quarks with WZ (tWZ) were modelled using the MADGRAPH5_aMC@NLO 2.3.3 [292] at NLO accuracy in QCD. For $t\bar{t}V$, the MC prediction was normalized to a cross section calculated at NLO accuracy in QCD with NLO electroweak corrections [97].

For all top processes, the NNPDF3.0NNLO PDF set was used. Additionally, the events were interfaced to PYTHIA 8.230 ($t\bar{t}$ and Wt) or 8.212 ($t\bar{t}V$, tZ , and tWZ) using the A14 tune to model the parton shower, hadronization, and underlying event. The decays of bottom and charm hadrons were simulated using EVTGEN 1.6.0 ($t\bar{t}$ and Wt) or 1.2.0 ($t\bar{t}V$, tZ , and tWZ) [405].

²The h_{damp} parameter is a resummation damping factor and one of the parameters that controls the matching of POWHEG matrix elements to the parton shower, effectively regulating the high- p_T radiation against which the $t\bar{t}$ system recoils.

Chapter 4

Reconstruction and Selection of Physics Objects

Reconstruction describes the process of building physics objects using the digitized measurements obtained from detector operation or simulation – it is the final step before a physics measurement is performed. The reconstructed objects are subject to calibrations as well as quality selections. The calibrations better match the kinematics of the reconstructed objects to the kinematics of the underlying physics processes as well as improve the modelling between data and simulation. The quality selections are typically a balance between selecting desired objects and rejecting undesired objects. ATLAS’s offline software framework, Athena [350], contains a suite of sophisticated algorithms and tools for reconstructing, calibrating, and selecting all physics objects of interest.

This chapter describes the reconstruction of physics objects pertinent to the $VH, H \rightarrow WW^*$ analysis (Section 4.1) as well as the selections placed on those physics objects (Section 4.2).

4.1 Object Reconstruction

The physics objects entering the $VH, H \rightarrow WW^*$ measurement include tracks, vertices, electrons, muons, jets, and missing transverse momentum. Fig. 4.1 sketches how each type of object interacts with ATLAS’s subdetectors.

4.1.1 Tracks

A track corresponds to the trajectory of a charged particle as it passes through the ATLAS detector. Tracks are used in the reconstruction of nearly every other physics object, and so reconstructing tracks with high efficiency and precision is an imperative.

A track is described by a set of five parameters at its perigee, the point of closest approach to the interaction point. The first parameter is the transverse impact parameter (d_0), given by the displacement of the track in the transverse plane relative to the beam axis at the track’s perigee. The second parameter is the longitudinal impact parameter (z_0), given by the displacement of the track in the longitudinal plane relative to the interaction point at the track’s perigee. The other three parameters are the polar (θ) and azimuthal (ϕ) angles as well as the charge-to-momentum ratio (q/p) of the track, each given at the track’s perigee [409].

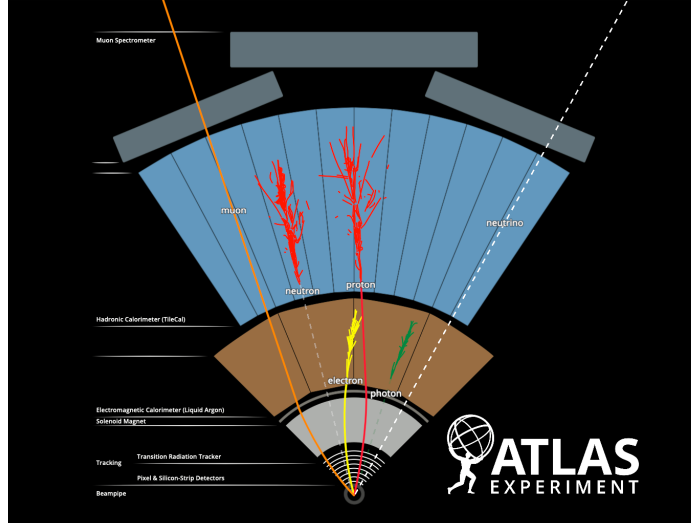


Figure 4.1: Slice of the ATLAS detector showing the interactions of different particles with the various subdetectors [408]. Protons and neutrons correspond to jets and neutrinos correspond to missing transverse momentum.

Charged particles traversing the sensitive detector elements of the inner tracker deposit energy, resulting in the separation (“deposition”) of electric charge. Charge deposition above a threshold set in the readout electronics yields a detector “hit”, the fundamental unit for track reconstruction, for the associated channel. In ATLAS, track reconstruction [410, 411] begins by assembling clusters of hits in the pixel detector and the SCT. A connected-component analysis¹ [412] groups hits with common edges or corners into clusters, and the clusters are used to create 3D measurements known as space-points [410, 411]. Each space-point corresponds to the coordinates where a charged particle traversed the detector. In the pixel detector, each cluster corresponds to a single space-point; in the SCT, clusters from each sensor of a double-sided module are combined to form a single space-point [410, 411].

In the pixel detector, incident particles typically deposit charge across multiple adjacent pixel sensors. In a sufficiently dense environment, it can happen that charge deposits associated with multiple particles are reconstructed as a single cluster known as a merged cluster [410]. In contrast, clusters whose charge is attributed to a single particle are known as single-particle clusters. Both types are shown pictorially in Fig. 4.2. A third type, a shared cluster, is shared among multiple reconstructed tracks but is not sufficiently compatible with the properties of a merged cluster to be identified as such [410].

Seeds are formed using sets of three space-points on different detector layers [410, 411]. An estimate of each track’s momentum and impact parameters is obtained by assuming a perfectly helical trajectory originating from the interaction point. The seeds are further subject to selections on their momentum, impact parameters, and multiplicity of additional compatible space-points to improve the quality of the resulting tracks. A combinatorial Kalman filter² [413] is then used to build track candidates from the seeds by incorporating additional space-points compatible with the preliminary trajectory. A seed may have multiple track candidates [410, 411].

¹A connected-component analysis corresponds to a sequence of local operations performed on a binary image which divide the image into a set of “connected” regions [412].

²A Kalman filter is a progressive method for track finding and fitting using time-ordered hits, propagating both the track’s trajectory and its covariance [413]. A Kalman filter combines forward filtering: estimating each “present” hit using all measured “past” hits; backward smoothing: estimating each past hit using all measured hits between it and the present; and outlier rejection [411, 413].

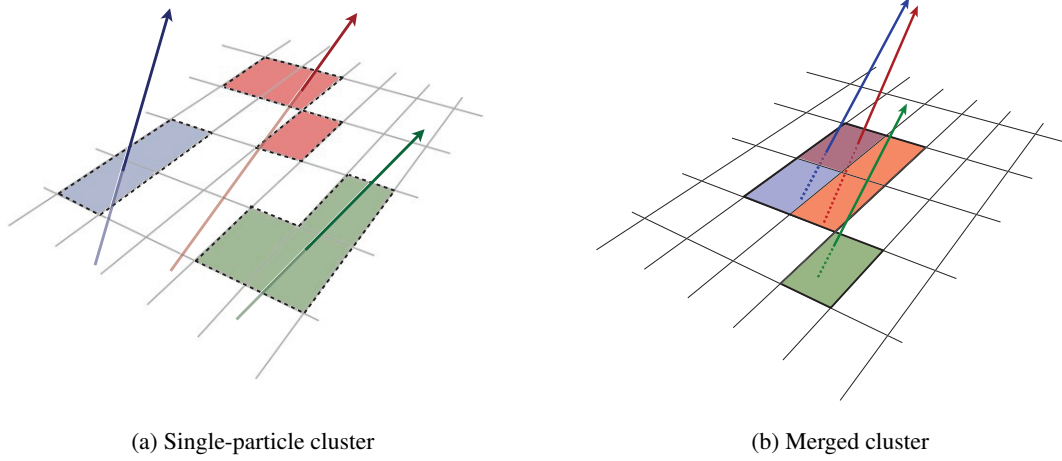


Figure 4.2: Illustration of (a) a single-particle cluster and (b) a merged cluster on a pixel module [410]. The charged particle trajectories are shown as arrows and the different colours correspond to charge deposited by different particles.

Each track candidate is assigned a track “score” based on quality metrics including the number of assigned space-points, the number of holes,³ the χ^2 of the fit, and the logarithm of the track’s p_T [410]. The scored tracks are ranked and provided to the ambiguity solver [414], which immediately rejects tracks with low scores; the accepted tracks are subsequently fitted. If a cluster is assigned to multiple tracks, a neural network [415] is used to identify the cluster as merged or shared. Clusters may only be shared with tracks previously accepted by the ambiguity solver. Additionally, clusters can be shared with no more than two tracks and a track can have no more than two shared clusters – a shared cluster is removed from a track candidate if it causes the candidate to fail these requirements [410]. The updated candidates are subsequently re-scored. Lastly, the ambiguity solver applies basic quality criteria to all tracks before final acceptance including $p_T > 500 \text{ MeV}$, $|\eta| < 2.5$, and selections on the minimum number of silicon hits, maximum number of silicon holes, and maximum number of shared clusters [410]. The logic of the ambiguity solver is shown pictorially in Fig. 4.3.

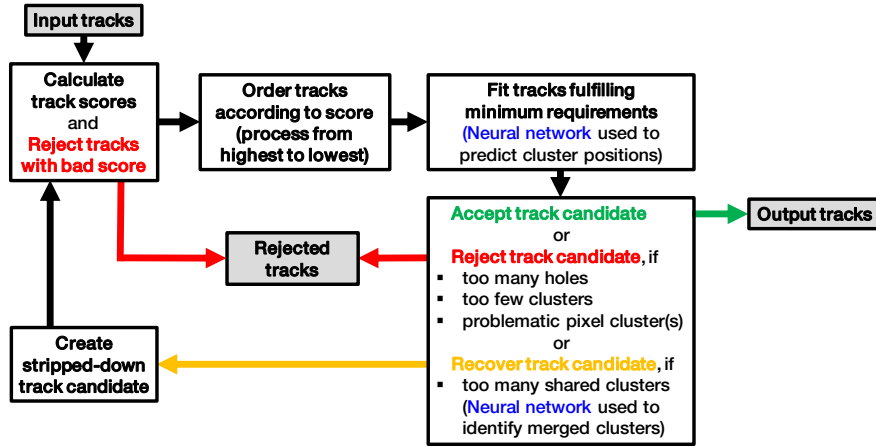


Figure 4.3: The flow of tracks through the ambiguity solver [410].

³A “hole” is defined as the absence of an expected detector hit. If the intersection of a track’s trajectory with a sensitive detector element does not contain a matching cluster within uncertainty, it is identified as a hole [410].

A high-resolution fit is performed for each track candidate accepted by the ambiguity solver. For these high-resolution fits, the position and uncertainty on the intersection of each track with each of its assigned clusters are obtained using additional neural networks [415]. TRT measurements are included in track reconstruction by extrapolating each track’s fit into the TRT and adding hits compatible with its trajectory. A sequence of outside-in track reconstruction is also performed by identifying TRT segments using a Hough transform⁴ [416] and backtracking the segments into the silicon tracker, improving efficiency by recouping silicon seeds missed during inside-out track reconstruction [411]. The average number of tracks reconstructed by ATLAS in 2017 data is shown in Fig. 4.4.

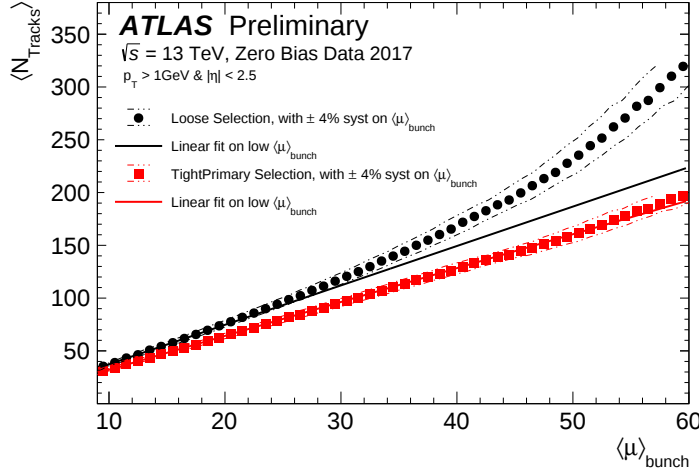


Figure 4.4: Average number of tracks reconstructed by the ATLAS detector in 2017 data as a function of the average number of interactions per bunch crossing [417]. The “Loose” and “TightPrimary” selections correspond to different sets of quality track criteria [418].

4.1.2 Vertices

Vertices are defined as the points in space where proton-proton interactions have occurred [419]. Per event, there are two types of vertices: the hard-scatter or *primary* vertex, which corresponds to the hard-scatter collision, and the pile-up vertices, which correspond to the pile-up collisions. The reconstruction of the primary vertex and the correct assignment of its tracks is an essential aspect of event reconstruction at the LHC.

Vertex reconstruction [419–421], also known as vertex finding and fitting, starts with the track collection passing the ambiguity solver’s quality selections. This track collection constitutes the track pool available. A vertex seed is identified (i.e., the *finding*) as the position of maximum track density along the longitudinal axis of the beam spot, the luminous region where collisions occur, using an iterative mode-finding algorithm [422]. Tracks are associated to the seed using a distance metric based on impact parameters.

The seed is then fitted (i.e., the *fitting*) using an adaptive Kalman filter [423]. Using the seed’s position as the starting point and its associated tracks as the input measurements, an iterative χ^2 minimization is performed to find the optimal vertex position [419]. The vertex is fitted, and each input track is assigned a weight reflecting its χ^2 compatibility with the vertex’s position estimate. The vertex is refitted using the weighted tracks, with low weights having less of an impact on the fit than high weights [419]. Each track is

⁴A Hough transform is a method for detecting patterns in binary image data by identifying local maxima in a parameter space, which is constructed using a back-projection equation relating the curve parameters and the image points [416].

subsequently re-weighted, and the algorithm is repeated. The tracks are re-scored once more following the final iteration, and tracks incompatible with the vertex by more than seven standard deviations, $\chi^2 > 49$, are removed from vertex candidate and returned to the track pool [419]. The tracks used in vertex's fit are not returned to the pool.

The procedure of vertex finding and fitting repeats until the track pool is empty or no additional vertices may be reconstructed. All vertices are required to have at least two associated tracks. The size of the beam spot may also be used to constrain each fit, resulting in improved transverse position resolution [419]. The vertexing algorithm described is known as the Iterative Vertex Finder (IVF) [424], whose logic is shown pictorially in Fig. 4.5. The average number of vertices reconstructed by ATLAS in 2018 data, along with a distribution of longitudinal position, is shown in Fig. 4.6.

Finally, the primary vertex is identified as the vertex whose associated tracks result in the largest p_T^2 -sum. Additional track-to-vertex association (TTVA) criteria based on impact parameters are used during the reconstruction of electrons, muons, and jets for assessing compatibility of tracks with the primary vertex. These criteria include selections on a track's transverse impact parameter, $|d_0|$, transverse impact parameter significance, $|d_0|/\sigma_{d_0}$, and absolute displacement from the primary vertex in the $y-z$ plane, $|\Delta z \sin \theta|$. Here, σ_{d_0} is the measured uncertainty on d_0 and Δz is the longitudinal displacement between a track's perigee, z_0 , and the primary vertex, z_{PV} : $\Delta z \equiv z_0 - z_{PV}$.

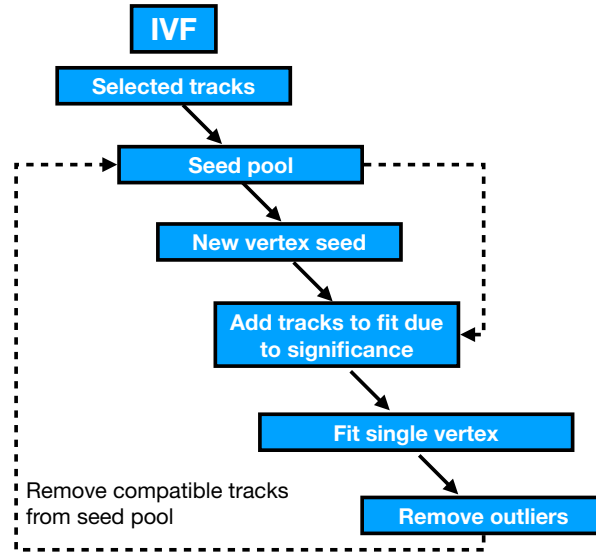


Figure 4.5: The flow of tracks through the Iterative Vertex Finder [424].

4.1.3 Electrons

An electron consists of a cluster of energy deposits in the electromagnetic calorimeter known as a super-cluster, which corresponds to the electromagnetic shower resulting from the electron's bremsstrahlung, and a matched track in the inner detector, which corresponds to the electron's trajectory through the inner detector [286, 426]. A diagram of the electron reconstruction algorithm is shown in Fig. 4.7.

Reconstruction: Reconstruction of electrons [286, 426] begins by forming electromagnetic topo-clusters [284, 427] by seeding and iteratively merging calorimeters cell based on their signal-to-noise ratio, considering only

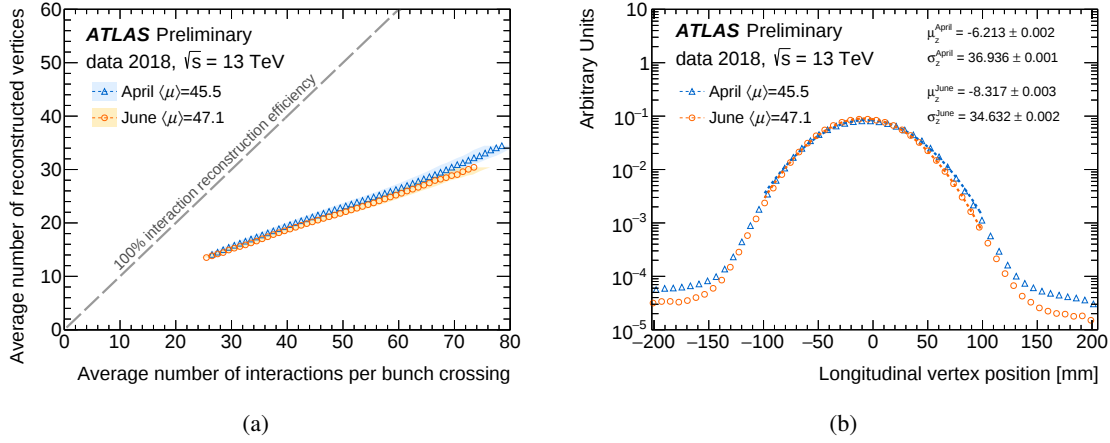


Figure 4.6: (a) Average number of vertices reconstructed by the ATLAS detector in 2018 data as a function of the average number of interactions per bunch crossing [425]. Two different average pile-up fills are shown. The errors bands are statistical with a $\pm 4\%$ systematic uncertainty included. (b) Normalized distribution of the longitudinal position of vertices reconstructed by the ATLAS detector in 2018 data [425]. Two different average pile-up fills are shown, and each is fitted with a Gaussian function.

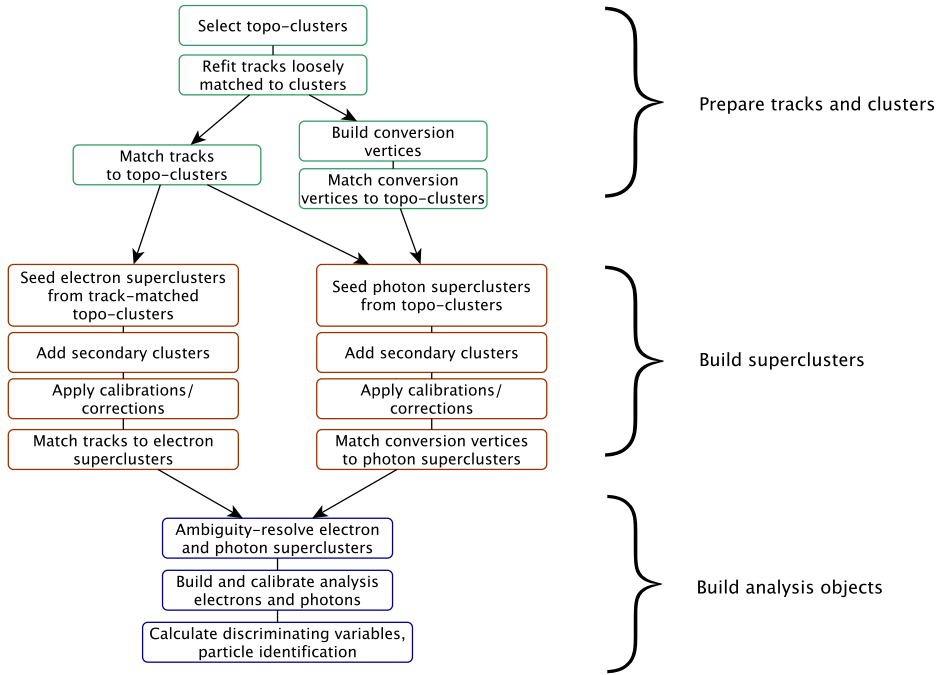


Figure 4.7: Diagram of the algorithm used for electron and photon reconstruction [426].

the second and third sampling layers of the electromagnetic calorimeter. Pattern recognition [411] is used to form track candidates in the inner detector, which are seeded by and subsequently matched to fixed-size electromagnetic clusters. If standard pattern recognition fails, a modified algorithm allowing for up to 30% energy loss at each material intersection is used instead [286, 426]. Track candidates are then fitted using a global χ^2 fitter [428], while tracks loosely matched to fixed-size clusters are refitted using a Gaussian sum filter [429], a generalization of the Kalman filter which accounts for bremsstrahlung [286]. Each track is extrapolated from its perigee to the second sampling layer of the electromagnetic calorimeter and matched to a topo-cluster by requiring $|\Delta\eta(\text{track}, \text{cluster})| < 0.05$ and $-0.10 < q \times \Delta\phi(\text{track}, \text{cluster}) < 0.05$, where q is the reconstructed charge of the track [426]. If a significant amount of energy is lost due to bremsstrahlung, a rescaling of the track momentum to the cluster energy is performed. If multiple tracks are matched to the same cluster, priority is given to the track with the highest number of hits in the innermost tracking layers and closest to the topo-cluster in ΔR [426].

Super-clusters [430] are reconstructed using the topo-clusters and their matched tracks. The topo-clusters are first sorted according to E_T , and clusters with $E_T > 1$ GeV and a matched track with at least four silicon hits are identified as seeds. For each seed cluster, satellite clusters within a window of size $\Delta\eta \times \Delta\phi = 0.075 \times 0.125$ (or $\Delta\eta \times \Delta\phi = 0.125 \times 0.300$, if the clusters share a matched track) around the seed barycenter are added to the seed cluster. The seed clusters and their associated satellite clusters constitute super-clusters. Tracks are then matched to super-clusters following the same procedure as for the topo-clusters. Finally, an ambiguity solver is used to determine the “author” of particles reconstructed as both electrons and photons (i.e., sharing a seed cluster): the author can be an electron, a photon, or ambiguous [426]. The electron reconstruction efficiency in simulation is shown in Fig. 4.8.

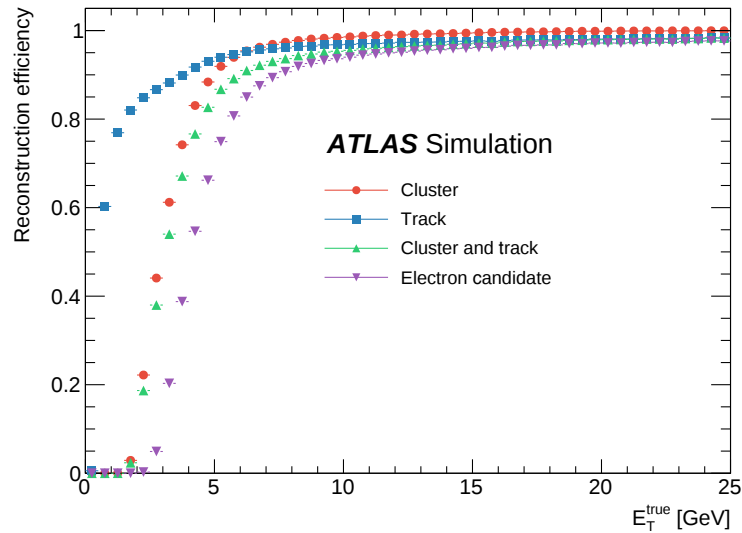
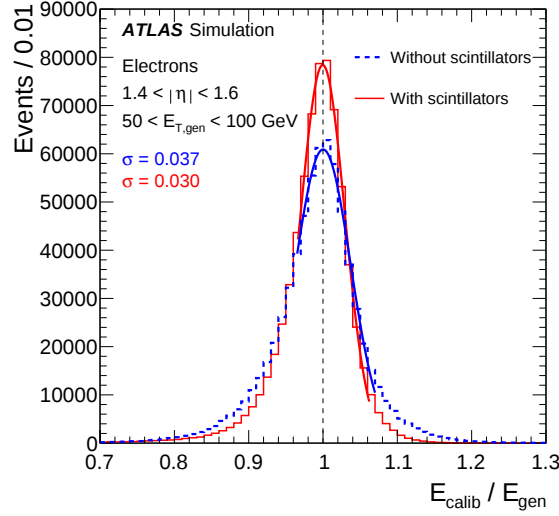


Figure 4.8: Reconstruction efficiencies in simulation for super-clusters, tracks, super-clusters and tracks, and electrons as a function of the particle-level electron transverse energy, E_T^{true} [426].

Calibration: The electron energy scale calibration, which restores the reconstructed electron energies to that of particle-level electrons, is performed using a boosted decision tree (BDT). Trained using $Z \rightarrow e\bar{e}$ events and utilizing shower energy and topological inputs, the BDT corrects for energy lost upstream/downstream of the

electromagnetic calorimeter or not contained within the clusters [285]. Distributions of the calibrated electron energy, E_{calib} , divided by the particle-level electron energy, E_{gen} , are shown in Fig. 4.9a. The energy resolution is extracted as the interquartile range of the distribution of $E_{\text{calib}}/E_{\text{gen}}$. As shown in Figs. 4.9b and 4.9c, the use of super-clusters over fixed-size clusters leads to an improvement in the energy resolution as a super-cluster collects more of an electron's deposited energy [426]. Additional η -dependent corrections, measured in $Z \rightarrow e\bar{e}/\mu\mu$ events, are applied to match the energy scale and resolution in data and simulation [285]. In data, these corrections account for the intercalibration of different sampling layers and the energy shifts induced by pile-up as well as improve the uniformity of the energy response; in simulation, these corrections constitute an offset to the energy resolution [426].



(a) Energy scale

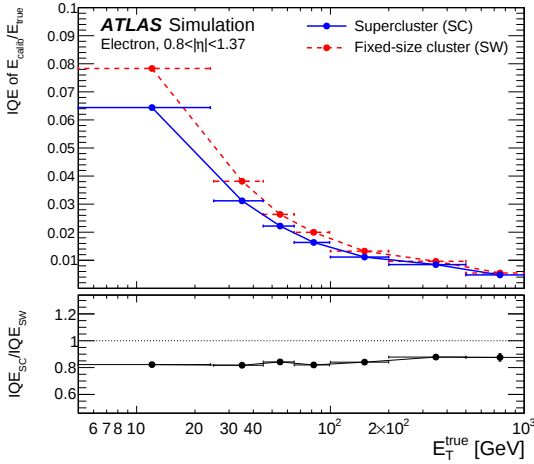
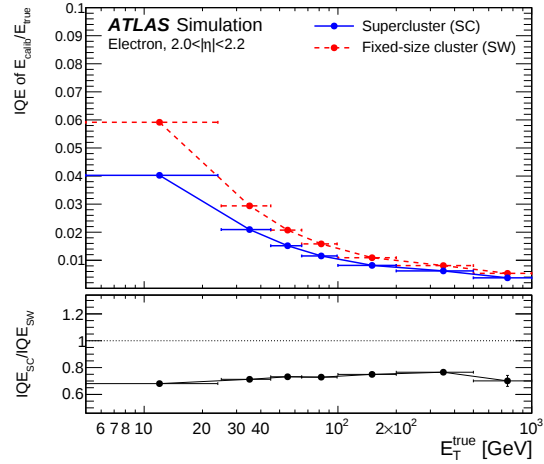
(b) Energy resolution, $0.8 < |\eta| < 1.37$ (c) Energy resolution, $2.0 < |\eta| < 2.2$

Figure 4.9: (a) Distributions of the calibrated electron energy, E_{calib} , divided by the particle-level electron energy, E_{gen} [285]. The (E_4) scintillators [127] instrument the transition region between the electromagnetic barrel and endcap calorimeters, leading to an improved energy resolution. (b, c) Calibrated electron energy resolution as a function of the particle-level electron transverse energy, E_T^{true} [426]. The energy resolution is given by the interquartile range (“IQE”) of the distribution of $E_{\text{calib}}/E_{\text{true}}$.

Identification: Identification criteria are used to improve the purity of selected electrons [286, 426]. A likelihood-based discriminant is constructed using $J/\psi \rightarrow e\bar{e}$ and $Z \rightarrow e\bar{e}$ candidates. The discriminant utilizes inputs relating to the primary electron track, the lateral and longitudinal development of the electromagnetic shower, and the spatial compatibility of the primary electron track and the super-cluster. Selections are placed on the likelihood discriminant in bins of $|\eta| - E_T$ such that three working points are defined: “Loose”, “Medium”, and “Tight”, corresponding to average identification efficiencies of 93%, 88%, and 80%, respectively, for electrons from electroweak processes [426]. The identification efficiencies in $Z \rightarrow e\bar{e}$ events for all three working points are shown in Fig. 4.10a.

Isolation: To quantify the activity near electrons, two isolation variables are defined. The calorimeter-based isolation variable E_T^{coneXX} is defined as the sum of the transverse energy of positive-energy topological clusters whose barycenters are within a cone of radius $\Delta R = \text{XX}/100$ centered on the electron cluster barycenter [426]. The transverse energy of the electron cluster’s core is subtracted from the sum, and corrections accounting for energy leakage and pile-up/underlying-event contributions [431] are applied. The track isolation variable $p_T^{\text{varconeXX}}$ is defined as the sum of the transverse momentum of selected tracks whose trajectories are within a variable-size cone of max radius $\Delta R_{\text{max}} = \text{XX}/100$ centered on and excluding the electron track [426]. The tracks must be of a good quality as well as satisfying $p_T > 1 \text{ GeV}$, $|\eta| < 2.5$, and $|\Delta z \sin \theta| < 3 \text{ mm}$. The radius of the cone is given by $\Delta R = \min((10 \text{ GeV})/p_T, \Delta R_{\text{max}})$, where p_T is the transverse momentum of electron track. A cone radius inversely proportional to p_T is optimal for topologies where other objects are in close proximity to an energetic lepton [286, 432]. The track- and calorimeter-based isolation variables complement one another, the former having a good resolution and low pile-up dependence and the latter accounting for neutral particles and particles below the track p_T threshold [432].

Isolation criteria are used to select prompt leptons, which result from the “prompt” decays of W/Z bosons originating from the hard-scatter vertex, while rejecting non-prompt leptons, which result from the “non-prompt” decays of b/c -hadrons or photons or due to the misidentification of a jet. Non-prompt leptons are typically characterized by a higher level of nearby activity. Two electron isolation working points, “Loose” and “Tight”, are defined using fixed cuts on the E_T^{cone20} and $p_T^{\text{varcone20}}$ isolation variables [426]. The Loose working point corresponds to $E_T^{\text{cone20}}/p_T < 0.20$ and $p_T^{\text{varcone20}}/p_T < 0.15$, while the Tight working point corresponds to $E_T^{\text{cone20}}/p_T < 0.06$ and $p_T^{\text{varcone20}}/p_T < 0.06$. The isolation efficiencies in $Z \rightarrow e\bar{e}$ events for the Loose and Tight working points are shown in Fig. 4.10b.

As an alternative to the fixed-cut working points, lepton isolation is also applied using a multivariate discriminant, the Prompt Lepton Improved Veto (PLIV). The PLIV utilizes inputs from the lepton, including its p_T , isolation variables, calorimeter cluster energy, and lifetime variables. The lifetime variables capture features of b -hadron decays. The PLIV also utilizes inputs from the closest track jet⁵ in ΔR , including the lepton p_T relative to the jet p_T , the lepton p_T projected onto the jet p_T , the lepton-jet ΔR , and the number of tracks contained within the jet. BDTs are trained separately for barrel ($|\eta| < 1.37$) electrons, endcap ($|\eta| > 1.52$) electrons, and muons. Selections are placed on each BDT discriminant as a function of lepton p_T , chosen such that the signal sensitivity is maximized when applied to the 2ℓ same-sign WH , $t\bar{t}(H/W)$ multilepton [435], and same-sign WW [436] analyses. This procedure has resulted in two working points, “PLImprovedTight” and “PLImprovedVeryTight”, for both electrons and muons. The isolation efficiencies in simulated $t\bar{t}$ events for the PLIV working points are shown in Fig. 4.11. In the ee same-sign WH channel, the use of PLImprovedTight

⁵Track jets are reconstructed from inner detector hits using the anti- k_t algorithm [433] with radius parameter $R = 0.4$. The tracks must be of a good quality as well as satisfying $p_T > 500 \text{ MeV}$, $|d_0| < 0.5 \text{ mm}$, and $|\Delta z \sin \theta| < 1.5 \text{ mm}$ [434].

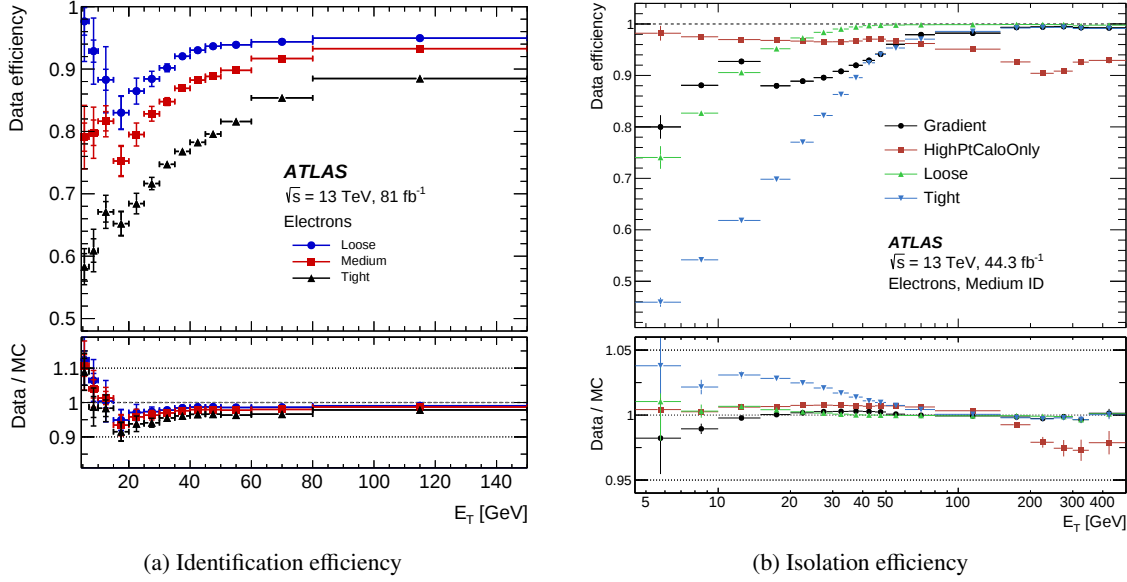


Figure 4.10: (a) Electron identification efficiencies in $Z \rightarrow e\bar{e}$ events as a function of E_T for the Loose, Medium, and Tight working points [426]. The inner error bars include only the statistical uncertainty and the outer error bars include both statistical and systematic uncertainties. (b) Electron isolation efficiencies in $Z \rightarrow e\bar{e}$ events as a function of E_T for the Loose, Tight, Gradient, and HighPtCaloOnly working points [426]. The electrons are required to pass Medium identification. The errors bars include both statistical and systematic uncertainties.

over the Tight working point improves the non-prompt rejection by a factor of 1.4 while rejecting only 9.2% more signal events.

Charge Misidentification: The electric charge of a reconstructed electron is determined using the curvature of its associated track in the inner detector [286, 426]. Interactions of the electron with the detector can lead to distortions of the primary electron track and an incorrect determination of the sign of the track’s curvature, and the presence secondary particles can pollute the region near the primary electron track and lead to an incorrectly matched track – either case may result in incorrect charge reconstruction [286, 426]. In order to suppress charge misidentification, a BDT is trained to separate $Z \rightarrow e\bar{e}$ data events with opposite- and same-sign electrons, utilizing the electron’s E_T , η , charge, and impact parameters as well as the charge, momentum, and topological information for its matched track. A selection on the output of this BDT, known as electron charge identification selection (ECIDS), leads to a 98% efficiency for selecting events with correctly reconstructed charge while rejecting 90% of events with incorrectly reconstructed charge [426]. The ECIDS efficiency in $Z \rightarrow e\bar{e}$ events is shown in Fig. 4.12.

Efficiency Scale Factors: Efficiency scale factors (SFs) are used to correct a particular selection’s efficiency in simulation, ε_{MC} , to that observed in data, ε_{data} :

$$SF(p^\mu) = \frac{\varepsilon_{data}(p^\mu)}{\varepsilon_{MC}(p^\mu)}, \quad (4.1)$$

where the efficiencies, and therefore the efficiency scale factors, are functions of the kinematics of the physics object under study, p^μ .

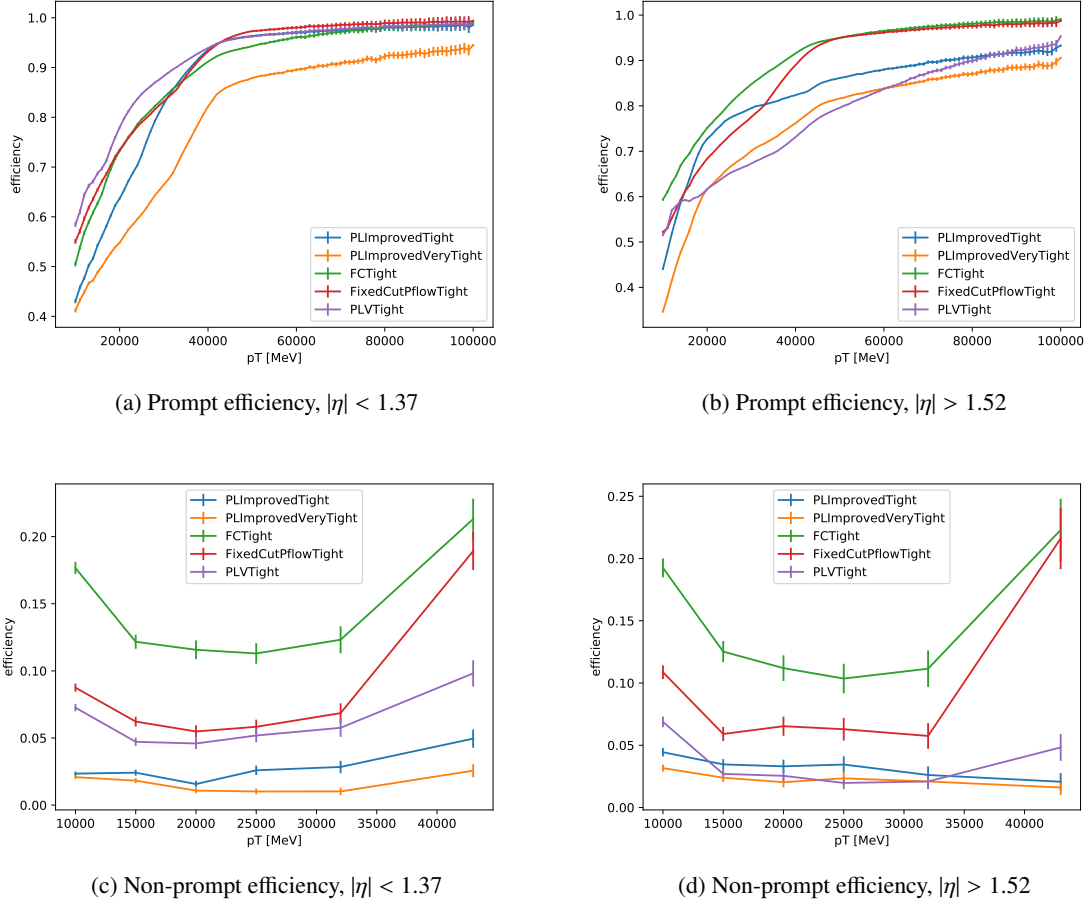


Figure 4.11: (a, b) Prompt and (c, d) non-prompt electron isolation efficiencies in simulated $t\bar{t}$ events as a function of p_T for the PLImprovedTight, PLImprovedVeryTight, Tight (“FCTight”), FixedCutPflowTight, and PLVTight working points. The electrons are required to pass Loose identification. The error bars indicate the statistical uncertainty.

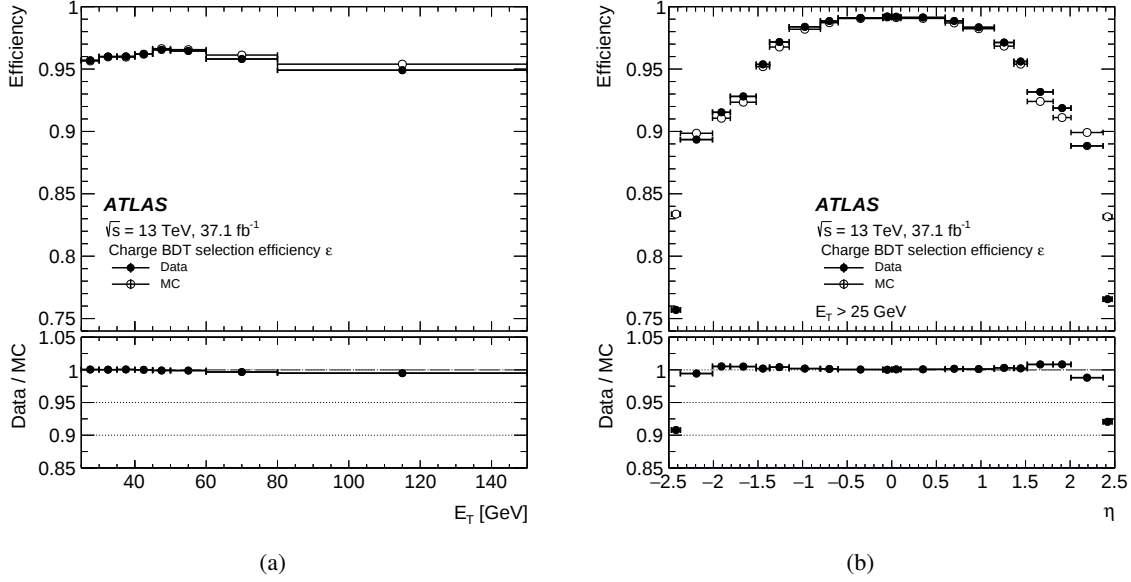


Figure 4.12: Election ECIDS efficiency in $Z \rightarrow e\bar{e}$ events as a function of (a) E_T and (b) η [286]. The electrons are required to pass Medium identification. The errors bars are smaller than the markers and therefore not visible.

Efficiency scale factors are estimated using the tag-and-probe method, which directly measures a particular selection’s efficiency in data and simulation. A tightly-selected physics object known as the “tag” is used to select only the events of interest, and a loosely-selected physics object known as the “probe” is used to measure the efficiencies, which are given as the number of probe objects passing the selection under study over the total number of tag (and therefore probe) objects [286].

Simulation-to-data scale factors for the electron trigger, reconstruction, identification, and isolation efficiencies are applied to the simulated yields containing reconstructed electrons. The scale factors and their associated uncertainties are estimated from $J/\psi \rightarrow e\bar{e}$ and $Z \rightarrow e\bar{e}$ events using the tag-and-probe method.

4.1.4 Muons

The $VH, H \rightarrow WW^*$ analysis uses only *combined* muons – that is, muons reconstructed using a combination of hits in the inner detector and muon spectrometer [432, 437]. Muons can also be reconstructed using only the muon spectrometer (“muon-spectrometer-extrapolated”) or only the inner detector and electromagnetic calorimeter (“calorimeter-tagged”). However, they are not considered in this thesis.⁶

Reconstruction: Reconstruction of combined muons [432, 437] begins by identifying linear track segments within the muon spectrometer by means of a Hough transform [416]. Track segments are loosely combined on the basis of a parabolic trajectory originating from the interaction point, which provides better matching between segments than a straight-line trajectory [432]. Muon candidates are built by assigning coordinates to the track segments using hit measurements in the bending and non-bending planes obtained from the precision tracking and trigger chambers, respectively. A global χ^2 fit [428] of the muon trajectory in the magnetic field is

⁶This is not entirely true – muon types other than combined muons are used during E_T^{miss} building (Section 4.1.6) and overlap removal (Section 4.2). However, only combined muons enter the $VH, H \rightarrow WW^*$ measurement.

performed, accounting material effects and detector alignment [438, 439]. The fit is performed a second time, after removing outliers and adding missing hits, as well as a third time, after removing low quality tracks. The third fit additionally accounts for energy loss in the calorimeters and loosely constrains the muon trajectory to originate from the interaction point. Finally, the muon candidates are back-extrapolated and matched to inner detector tracks and all hits are iteratively refitted. The resulting p_T of the muon candidate is expressed at the interaction point [432].

Calibration: Combined muons in simulation are calibrated to the muon momentum scale and resolution observed in data using fits to the dimuon invariant mass distributions for $J/\psi \rightarrow \mu\bar{\mu}$ and $Z \rightarrow \mu\bar{\mu}$ candidates [437], as shown in Fig. 4.13. The scale and resolution corrections are derived in regions of $\eta - \phi$ and independently for the transverse momentum measured using only the inner detector and the transverse momentum measured using only the muon spectrometer. In the latter case, the momentum imbalance between the inner detector and the muon spectrometer measurements is also fitted. The corrected momentum of the combined muon is assumed to be a weighted average of the two independent measurements [437].

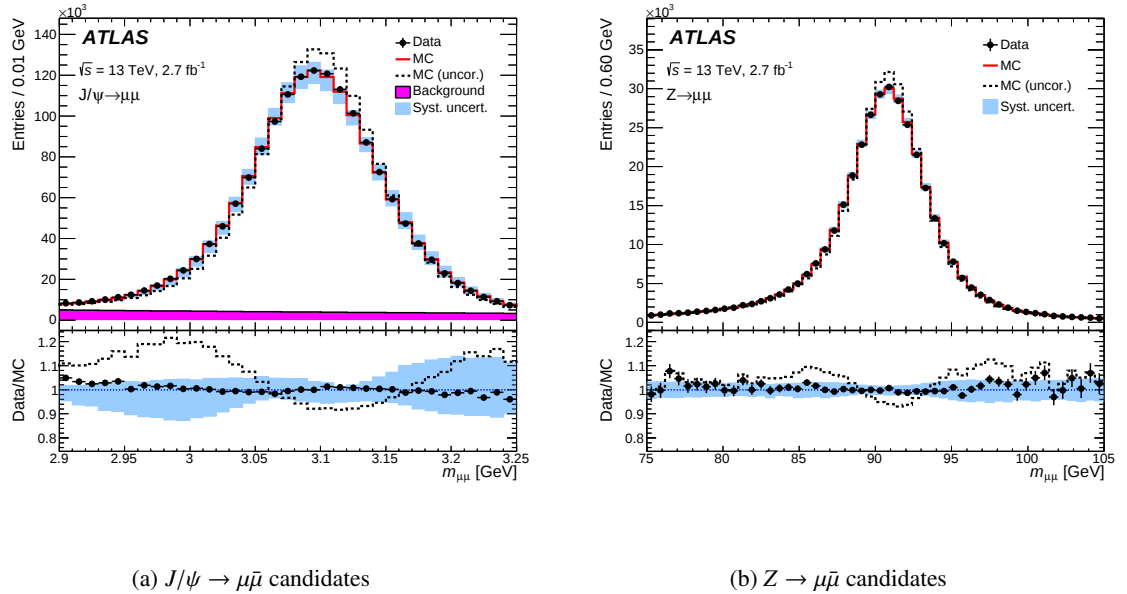


Figure 4.13: Dimuon invariant mass distributions reconstructed using combined muons from (a) $J/\psi \rightarrow \mu\bar{\mu}$ and (b) $Z \rightarrow \mu\bar{\mu}$ candidates [437]. Both the uncorrected and corrected simulated mass distributions are shown. The error band indicates the systematic uncertainty on the MC momentum corrections.

Identification: Identification proceeds by applying quality requirements to muon candidates such that prompt muons are selected with high efficiency while suppressing backgrounds from in-flight decays of pions and kaons [432, 437]. The quality requirements are based on the number of hits/holes in precision tracking detectors of the muon spectrometer, the charge-over-momentum compatibility between the momentum measurements in the inner detector and the muon spectrometer, the absolute difference between the momentum measurements in the inner detector and the muon spectrometer divided by the momentum of the combined muon, and the normalized χ^2 of the combined track's fit. Three working points, each associated with a particular set of selections, are defined: “Loose”, “Medium”, and “Tight”, corresponding identification efficiencies of 99%,

97%, and 93%, respectively, for muons with $p_T > 20$ GeV [432]. The identification efficiency in $J/\psi \rightarrow \mu\bar{\mu}$ and $Z \rightarrow \mu\bar{\mu}$ events for the Medium quality working point is shown in Fig. 4.14a.

Track-To-Vertex Association: TTVA requirements are applied to further reject muons originating from in-flight decays of hadrons or pile-up interactions. In particular, the muon tracks are required to pass association to the primary vertex: $|d_0|/\sigma_{d_0} < 3$ and $|\Delta z \sin \theta| < 0.5$ mm [432].

Isolation: For muon isolation, two working points, “Loose” and “Tight”, are defined using fixed cuts on the E_T^{cone20} and $p_T^{\text{varcone30}}$ isolation variables [432]. The Loose working point corresponds to $E_T^{\text{cone20}}/p_T < 0.30$ and $p_T^{\text{varcone30}}/p_T < 0.15$, while the Tight working point corresponds to $E_T^{\text{cone20}}/p_T < 0.15$ and $p_T^{\text{varcone30}}/p_T < 0.04$. Both working points prioritize high efficiency for prompt muons over low efficiency for non-prompt muons [432]. The isolation efficiency in $Z \rightarrow \mu\bar{\mu}$ events for the Loose working point is shown in Fig. 4.14b, and the isolation efficiencies in simulated $t\bar{t}$ events for the PLIV working points are shown in Fig. 4.15. In the $\mu\mu$ same-sign WH channel, the use of PLImprovedTight over the Tight working point improves the non-prompt rejection by a factor of 2.1 while rejecting only 6.5% more signal events.

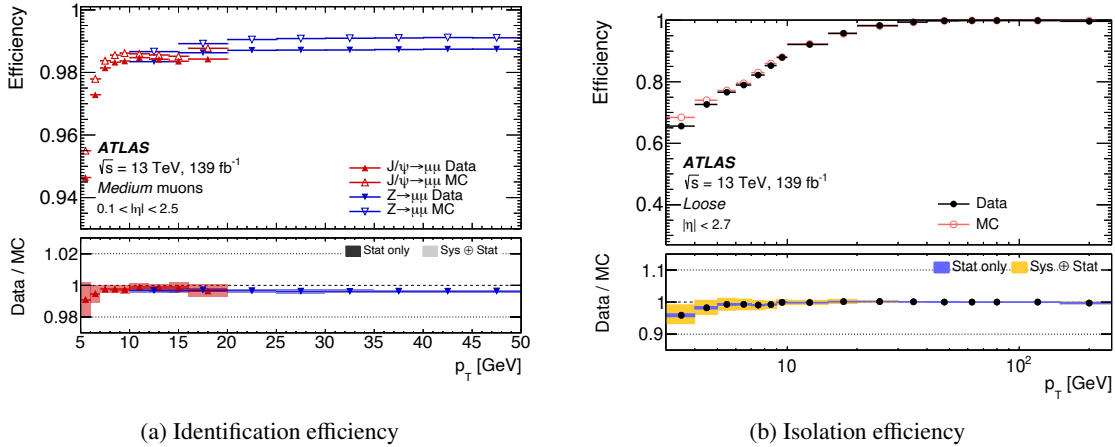


Figure 4.14: (a) Muon reconstruction and identification efficiencies in $J/\psi \rightarrow \mu\bar{\mu}$ and $Z \rightarrow \mu\bar{\mu}$ events as a function of p_T for the Medium working point [432]. The error band includes both statistical and systematic uncertainties – the statistical component is not visible. (b) Muon isolation efficiency in $Z \rightarrow \mu\bar{\mu}$ data events as a function of p_T for the Loose working point [432]. The blue error band includes only the statistical uncertainty and the yellow error band includes both statistical and systematic uncertainties.

Efficiency Scale Factors: Simulation-to-data scale factors for the muon trigger, reconstruction/identification, TTVA, and isolation efficiencies are applied to the simulated yields containing reconstructed muons. The scale factors and their associated uncertainties are estimated from $J/\psi \rightarrow \mu\bar{\mu}$ and $Z \rightarrow \mu\bar{\mu}$ events using the tag-and-probe method.

4.1.5 Jets

ATLAS reconstructs jets [440] using the particle-flow algorithm, which incorporates measurements from the inner detector and the calorimeter in a complementary manner. Low- p_T charged particles are best measured

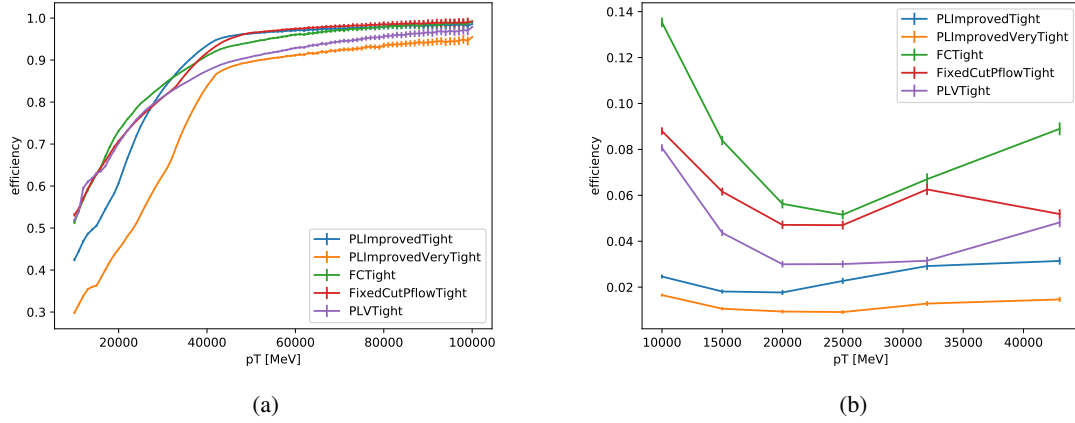


Figure 4.15: (a) Prompt and (b) non-prompt muon isolation efficiencies in simulated $t\bar{t}$ events as a function of p_T for the PLImprovedTight, PLImprovedVeryTight, Tight (“FCTight”), FixedCutPflowTight, and PLVTight working points. The muons are required to pass Medium identification. The error bars indicate the statistical uncertainty.

by the inner detector, while charged and neutral particles at high energies are best measured by the calorimeter. The integration of tracking has additional advantages, including improved clustering of low- p_T charged particles as well as improved rejection of pile-up jets by identifying jets associated to the primary vertex [440].

The ALEPH experiment at LEP pioneered the use of particle-flow algorithms for object reconstruction [441]. The ATLAS [442] and CMS [443–445] experiments have used particle-flow algorithms to reconstruct hadronic taus and, in the case of CMS, jets since Run 1. However, jet reconstruction via particle-flow in ATLAS was a development introduced during Run 2, with most analyses previously utilizing jets built from topological clusters of calorimeter cells [427].

Reconstruction: A flow chart for the particle-flow algorithm is shown in Fig. 4.16. Analogous to the reconstruction of electrons, the particle flow algorithm starts by forming topo-clusters [284, 427] by seeding and iteratively merging calorimeter cells in a given sampling layer based on their signal-to-noise ratio. Both electromagnetic and hadronic calorimeter cells are used, and the resulting topo-clusters are calibrated to the electromagnetic energy scale. Tracks reconstructed by the inner detector are matched to topo-clusters such that the ΔR -based significance, calculated using the track’s position and the topo-cluster’s barycenter in the second sampling layer of the electromagnetic calorimeter, is minimized. The tracks must be of good quality as well as satisfying $0.5 < p_T < 100$ GeV and $|\eta| < 2.5$. Additionally, tracks matched to candidate electrons or muons are excluded [198, 440]. Only topo-clusters satisfying $E_{\text{cluster}}/p_{\text{track}} > 0.1$ are considered for matching, where E_{cluster} is the topo-cluster’s energy and p_{track} is the track’s momentum. Tracks with a minimum significance of 1.64 or greater are not matched to any topo-cluster [440].

For tracks matched to topo-clusters, a combination of the inner detector and the calorimeter measurements is used. The energy deposited by each track in the calorimeter must therefore be subtracted to avoid double-counting [440]. For each track matched to a topo-cluster, the average energy deposited in the calorimeter, $\langle E_{\text{dep}} \rangle$, and its associated spread, $\sigma(E_{\text{dep}})$, are calculated using calibrations measured in single-particle samples without pile-up. The significance of the difference between the measured and the expected energy of the matched topo-cluster, $(E_{\text{cluster}} - \langle E_{\text{dep}} \rangle) / \sigma(E_{\text{dep}})$, is required to be greater than -1 . Otherwise, the track’s

energy is likely split across multiple topo-clusters, and a split-shower recovery procedure is performed by matching topo-clusters to the track if they fall within a cone of radius $\Delta R = 0.2$ centered on the track's position in the second sampling layer of the electromagnetic calorimeter [440].

For each track, the subtraction step is *only* performed if the calorimeter energy in a cone of radius $\Delta R = 0.15$ around the extrapolated track, E_{cluster} , satisfies $(E_{\text{cluster}} - \langle E_{\text{dep}} \rangle) / \sigma(E_{\text{dep}}) < 33.2 \times \log_{10}((40 \text{ GeV}) / p_{\text{T}}^{\text{track}})$ [198]. This condition avoids subtraction in very dense environments, where the accurate removal of the track's energy is difficult. The subtraction step is executed in two parts: the first step is cell-by-cell subtraction [440]. If $\sigma(E_{\text{dep}})$ exceeds the total energy of the matched topo-clusters, the topo-clusters are removed. Otherwise, rings in $\eta - \phi$ are formed around the extrapolated track such that they contain at least one cell, and the rings are then ranked in descending order according to their contained energy density, E_{ring} . Starting with the ring with the highest energy density, rings are successively removed until $\sum_{\text{removed}} E_{\text{ring}} \geq \langle E_{\text{dep}} \rangle$, where the energy of the final ring is scaled down instead of removed if $\sum_{\text{removed}} E_{\text{ring}} > \langle E_{\text{dep}} \rangle$ [440]. In situations where no subtraction has been performed, the track momentum is scaled down to avoid double-counting. Further scaling is applied to all particle-flow objects to smooth the transition between the low- and high-energy scales [198]. The second step is remnant removal [440]. If the energy of the cells surviving cell-by-cell subtraction is less than $1.5\sigma(E_{\text{dep}})$, the energy is assumed to originate from shower fluctuations of a single particle and the cells are removed. Otherwise, the energy is likely attributed to multiple particles and the cells are retained [440].

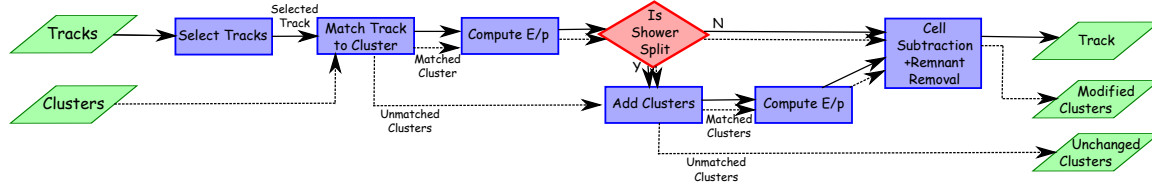


Figure 4.16: A flow chart of ATLAS's particle-flow algorithm for jet reconstruction [440].

The particle-flow algorithm results in a set of selected tracks, modified topo-clusters, and unmodified topo-clusters. The tracks are required to be matched to the primary vertex, $|\Delta z \sin \theta| < 2 \text{ mm}$, in order to reject tracks from pile-up interactions [440]. Particle-flow jets are then reconstructed from this ensemble of objects using the anti- k_t jet clustering algorithm [433] with radius parameter $R = 0.4$ using FASTJET [446]. A key aspect of the anti- k_t algorithm is its infrared and collinear safety, resulting in jets whose boundaries are unaffected by the addition of soft radiation. The anti- k_t algorithm produces hard jets with approximately circular shapes of radius R in $\eta - \phi$ and softer jets with complex shapes [433]. Examples of jet properties in $Z \rightarrow \mu\bar{\mu}$ events are shown in Fig. 4.17.

Calibration: The jet energy scale (JES) calibration, shown diagrammatically in Fig. 4.18, restores the reconstructed jet energy to that of particle-level jets⁷ [198] – that is, the “true” jet energy. Pile-up corrections remove excess energy due to pile-up interactions and are applied to the jet p_{T} in two steps [198]. The first step is a correction based on jet area (in $\eta - \phi$) and p_{T} density in the central region of the detector; the second step is a residual correction to account for pile-up sensitivity in the forward regions of the detector and in the core of high- p_{T} jets. The residual correction, which is derived in different regions of $|\eta|$, is function of the mean number of interactions per bunch crossing as well as the number of reconstructed vertices in the event.

⁷Particle-level jets are built from stable, particle-level objects originating from the hard-scatter vertex.

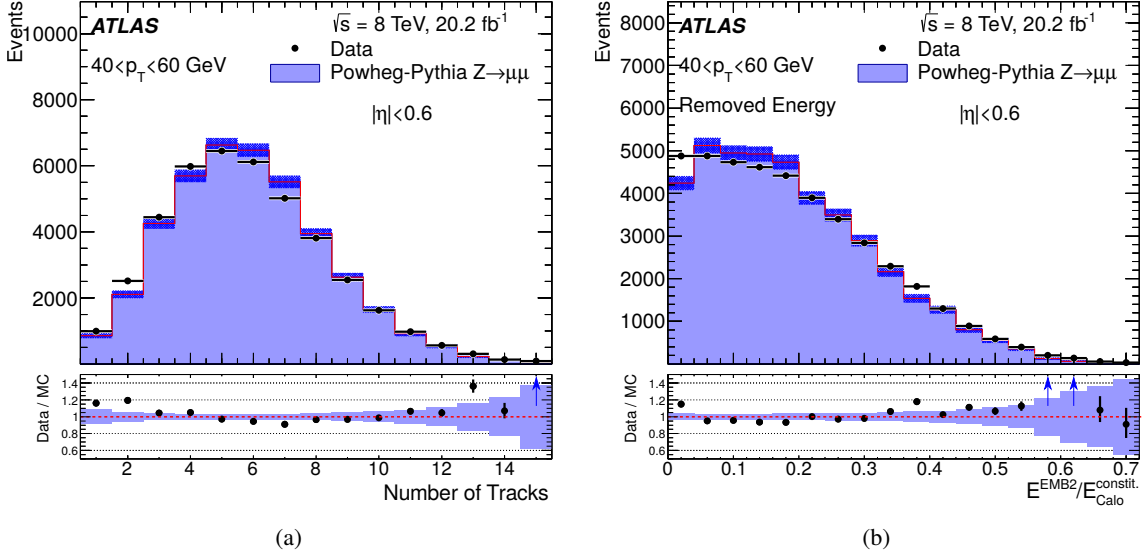


Figure 4.17: Examples of jet properties in $Z \rightarrow \mu\mu$ events, including (a) the number of tracks in the jet which originate from the primary vertex and (b) the energy removed by the second sampling layer of the electromagnetic barrel calorimeter (“EMB2”) relative to the total energy of the constituents of the matched calorimeter jet [440]. The error bands indicate the statistical uncertainty.

Following pile-up corrections, an absolute JES and η calibration is applied to account for the non-compensating calorimeter response, energy losses in passive materials, out-of-cone effects, and biases in jet η reconstruction due to transitions in calorimeter granularity or technology [198]. Both calibrations are derived using simulated dijet events. A global sequential calibration, also derived using simulated dijet events, consists of a series of multiplicative corrections to reduce the dependence of the jet p_T response on observables reconstructed using the tracking, calorimetry, and muon subsystems [198]. In doing so, the global sequential calibration improves the jet resolution without changing the average jet response. The pile-up correction, absolute JES/ η calibration, and global sequential calibration are applied to both data and simulation. A final *in situ* calibration is applied to the jet response in data to correct for the remaining differences between data and simulation due to imperfect event generation and detector simulation [198]. The *in situ* calibration is derived using well-measured reference objects, including photons, Z bosons, and calibrated jets. The ratio of the jet response in data to that in simulation obtained from the *in situ* measurements is shown in Fig. 4.19a.

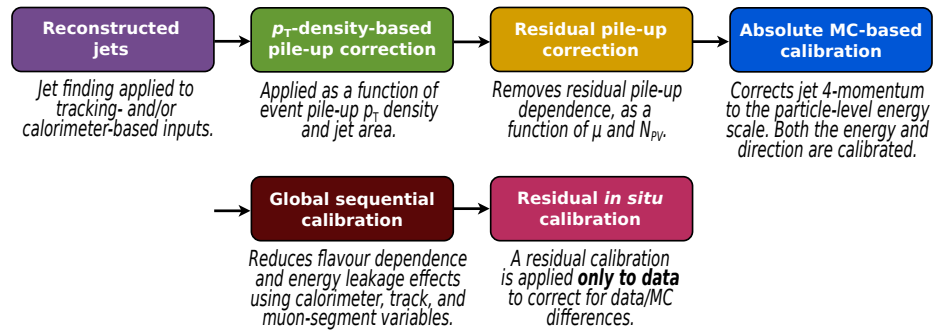


Figure 4.18: The flow of stages through the jet energy scale calibration [198].

The jet energy resolution (JER) is measured using the dijet balance method, which relies on the approximate scalar balance between the transverse momenta of the two leading jets in dijet events [198]. Measurements of the p_T asymmetry between a probe jet and a well-measured reference jet are iteratively fitted, accounting for event-by-event physics effects such as additional radiation or non-perturbative processes, to extract the resolution in regions of p_T and η . The noise contribution to the JER is constrained using measurements of the fluctuations in the energy deposited due to pile-up [198]. Using data events collected with random unbiased triggers, the calorimeter energy deposits contained within two cones at random ϕ and opposite η are computed. The noise resolution is extracted as the width of the distribution of the random cone balance, the energy sum difference between the two cones, as a function of η and the average number of interactions per bunch crossing. The random cones measurement is applied as a constraint to the dijet balance measurement and the JER, shown in Fig. 4.19b, is obtained from the combined statistical fit [198].

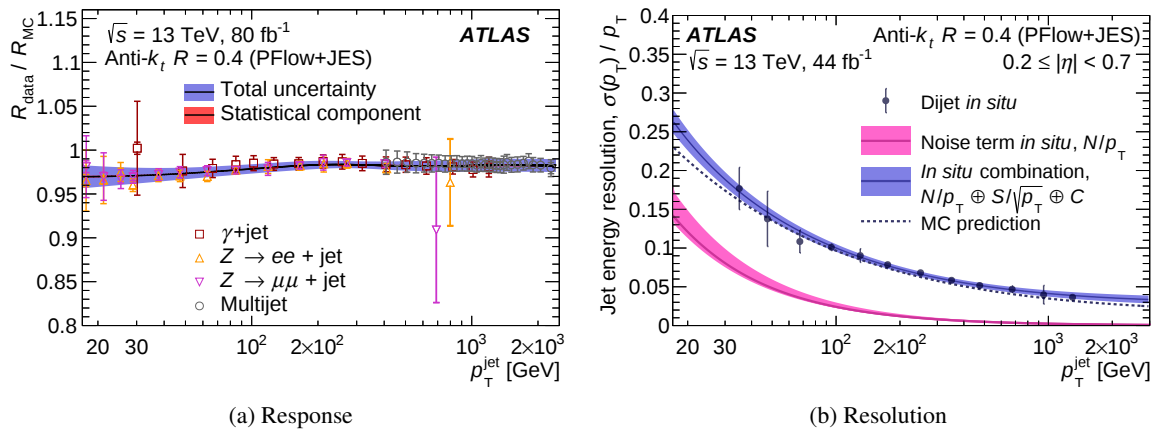


Figure 4.19: (a) Ratio of the jet response in data, R_{data} , to that in simulation, R_{MC} , as a function of jet p_T for the combination of the *in situ* measurements [198]. For the *in situ* measurements (i.e., data points), the inner error bars include only the statistical uncertainty and the outer error bars include both statistical and systematic uncertainties. For the final correction (i.e., solid line), the error band includes both statistical and systematic uncertainties – the statistical component is not visible. (b) Relative jet energy resolution as a function of jet p_T for the combination of the dijet balance and random cones measurements [198]. The error bars/bands include both statistical and systematic uncertainties.

Pile-Up Jet Tagging: A likelihood-based discriminant, the Jet Vertex Tagger (JVT) [447], is used to suppress jets originating from pile-up in the central region of the detector. The JVT is constructed using simulated dijet events and describes the relative probability of a jet originating from the hard-scatter vertex. For each input jet, the p_T -sum of the tracks contained within the jet and associated to the hard-scatter vertex is calculated. The resulting quantity is expressed relative to the p_T -sum of the tracks contained within the jet and associated to any vertex (the jet vertex fraction or JVF) as well as relative to the jet's p_T (R_{p_T}) [447]. The JVT is evaluated at this point in the $\text{JVF} - R_{p_T}$ plane and returned as the output score. For jets satisfying $20 < p_T < 60$ GeV and $|\eta| < 2.4$, the JVT output score is required to be greater than 0.5, corresponding to the “Tight” working point. The JVT selection efficiency in $Z \rightarrow \mu\mu$ events for the Tight working point is shown in Fig. 4.20a.

A second discriminant, the Forward JVT (fJVT) [448, 449], is used to suppress jets originating due to pile-up in the forward regions of the detector. The fJVT constructs a discriminant based on the missing transverse momentum per vertex, utilizing the JVT to remove central hard-scatter jets with high efficiency.

For each forward jet, the missing transverse momentum per vertex is projected onto the jet p_T , and the maximum value of the normalized projection across all vertices corresponds to the fJVT output score [449]. The normalized projection is close to 0 for a forward hard-scatter jet as the jet’s momentum is not expected to be balanced by the missing transverse momentum. For jets satisfying $20 < p_T < 60$ GeV and $2.5 < |\eta| < 4.5$, the fJVT output score is required to be less than 0.5, corresponding to the “Loose” working point. The fJVT selection efficiency in simulated $Z \rightarrow \mu\bar{\mu}$ events is shown in Fig. 4.20b.

Simulation-to-data scale factors for the JVT and fJVT selection efficiencies are applied to the simulated yields containing reconstructed jets. The scale factors and their associated uncertainties are estimated from $Z \rightarrow \mu\bar{\mu}$ events using the tag-and-probe method, where the Z boson and the leading jet are used as the tag and the probe, respectively [447, 448].

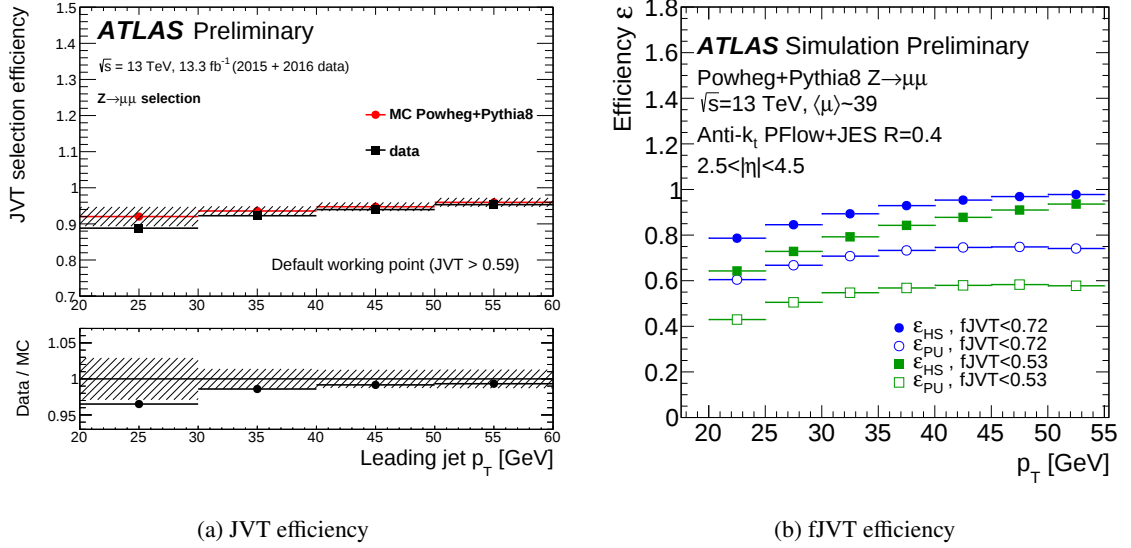


Figure 4.20: (a) JVT selection efficiency in $Z \rightarrow \mu\bar{\mu}$ events as a function of jet p_T for the Tight working point [450]. The jets used in this plot are *calorimeter* jets, reconstructed from topological clusters [284, 427] using a local cluster weighting algorithm [451]. Accordingly, the selection on the JVT output score is different from that used for particle-flow jets. The error band includes both statistical and systematic uncertainties. (b) Hard-scatter and pile-up jet selection efficiencies using the fJVT in simulated $Z \rightarrow \mu\bar{\mu}$ events as function of jet p_T [449].

Flavour Tagging: The identification of jets containing b -hadrons (“ b -jets”) is critical in identifying top quark processes, which primarily decay via $t \rightarrow Wb$ [452, 453]. Algorithms for doing so – b -tagging or, more generally, flavour tagging algorithms – exploit the long lifetime ($\langle c\tau \rangle \approx 458 \mu\text{m}$), high mass ($m_b \approx 4.2$ GeV), and high decay multiplicity of b -hadrons as well as the properties of b -quark fragmentation [454]. In ATLAS, low-level algorithms are first used to reconstruct properties of b -jets, and these properties are subsequently provided to a high-level algorithm consisting of a multivariate classifier. The output of the classifier corresponds to the relative probability of a jet containing a b -hadron [454].

In this thesis, b -jets are identified using the DL1r tagger [455, 456], a deep neural network with a multidimensional output for concurrently identifying b -jets, jets containing c -hadrons but no b -hadrons (“ c -jets”), and jets containing neither b - or c -hadrons (“light-jets”). Per jet, the DL1r tagger receives information relating to the jet’s p_T and $|\eta|$, the impact parameters and their correlations for the tracks contained within the jet,

and any displaced vertices [456]. A displaced vertex has a significant transverse displacement relative to the longitudinal axis of the beam spot due to the in-flight decay of a long-lived particle, such as a b -hadron. Displaced vertices are reconstructed using the tracks contained within the jet [457, 458]. Properties of displaced vertices utilized by the DL1r tagger include displacement from the primary vertex, invariant mass, energy fraction, track multiplicity, and track p_T and η [454]. The DL1r tagger is trained using b -, c -, and light-jets in simulated $t\bar{t}$ events.

Working points are defined using cuts on the output of the DL1r tagger such that specific b -tagging efficiencies are achieved for the b -jets in simulated $t\bar{t}$ events [454]. In this thesis, the 85% (b -tagging) efficiency working point is used. The 85% efficiency working point has rejection factors of 2.6 and 29 for c - and light-jets, respectively [454]. The corresponding b -tagging efficiencies are measured in data using dileptonic $t\bar{t}$ events and assuming c - and light-jet mistag rates obtained from independent measurements [455, 459]. A maximum-likelihood fit is performed to extract the efficiencies as function of jet p_T . Simulation-to-data scale factors are derived to correct the tagging efficiency in simulation to data [454, 455, 459]. Additionally, as the b -tagging efficiency varies between event generators due to usage of different fragmentation models, simulation-to-simulation scale factors are derived to correct the tagging efficiency in simulation when the event generator is different from that used to derive the simulation-to-data scale factors [454]. These scale factors are applied to the simulated yields when a requirement is placed on the number of b -jets.

The b -tagging efficiency for the DL1 tagger [456] at the 85% efficiency working point is shown in Fig. 4.21a. The corresponding simulation-to-data scale factors are shown in Fig. 4.21b. Strictly speaking, the DL1 tagger is an earlier iteration of the DL1r tagger; however, the former's performance is representative of the latter's.

4.1.6 Missing Transverse Momentum

Missing transverse momentum is an experimental proxy for the momentum carried by undetected particles [460]. In the context of the $VH, H \rightarrow WW^*$ measurement, these particles are neutrinos. The missing transverse momentum is calculated as the *negative* vectorial sum of the transverse momenta for all hard objects in the event (i.e., the “hard” term) as well as for tracks associated to the primary vertex but not to any of the hard objects (i.e., the “soft” term):

$$\vec{E}_T^{\text{miss}} = - \underbrace{\sum_{\text{selected electrons}} \vec{p}_T^e - \sum_{\text{selected muons}} \vec{p}_T^\mu - \sum_{\text{accepted taus}} \vec{p}_T^{\tau_{\text{had}}} - \sum_{\text{accepted photons}} \vec{p}_T^\gamma - \sum_{\text{accepted jets}} \vec{p}_T^{\text{jet}}}_{\text{hard term}} - \underbrace{\sum_{\text{unused tracks}} \vec{p}_T^{\text{track}}}_{\text{soft term}}. \quad (4.2)$$

The x and y components of the vectorial missing transverse momentum are denoted by E_x^{miss} and E_y^{miss} , respectively, leading to the definition of the magnitude of the missing transverse momentum:

$$E_T^{\text{miss}} \equiv |\vec{E}_T^{\text{miss}}| = \sqrt{(E_x^{\text{miss}})^2 + (E_y^{\text{miss}})^2}, \quad (4.3)$$

and its azimuthal angle:

$$\phi^{\text{miss}} \equiv \arctan(E_y^{\text{miss}}/E_x^{\text{miss}}). \quad (4.4)$$

In the remainder of this thesis, the phrase “missing transverse momentum” refers exclusively to its magnitude. Where applicable, the phrase “vectorial missing transverse momentum” is used to refer to Eq. 4.2.

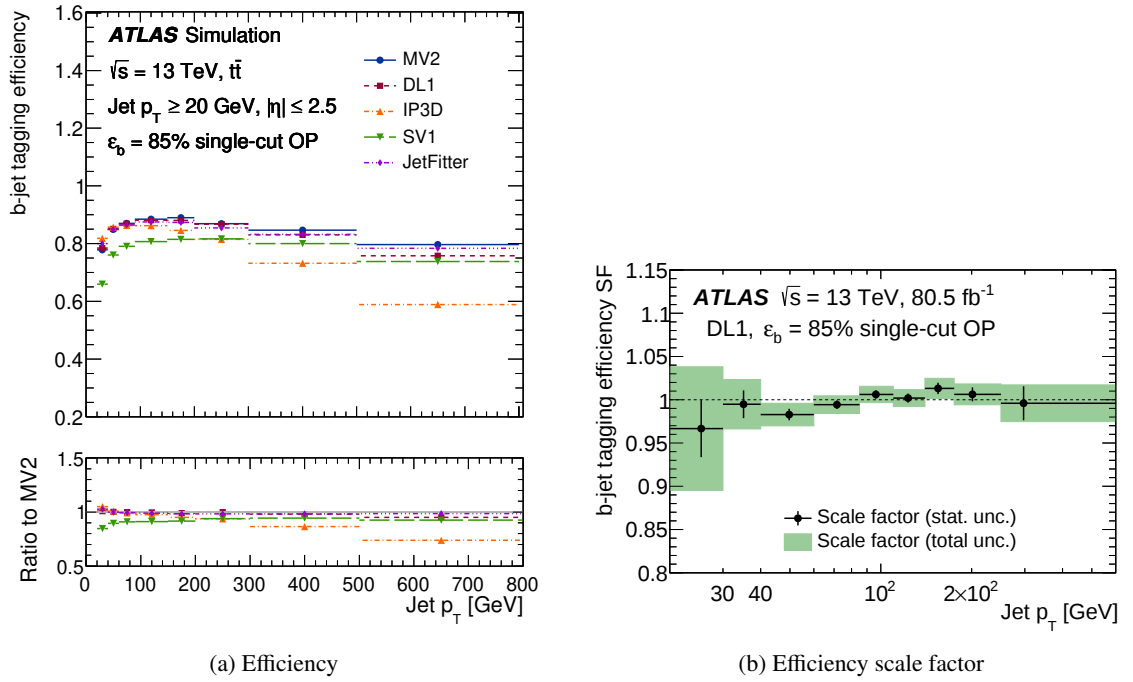


Figure 4.21: (a) The b -tagging efficiency in simulated $t\bar{t}$ events as a function of jet p_T [454]. The 85% efficiency working point is shown for the MV2 [457], DL1 [457], IP3D [457], SV1 [457], and JETFITTER [458] algorithms. (b) The simulation-to-data efficiency scale factors for the DL1 tagger at the 85% working point [454]. The scale factors are shown as a function of jet p_T . The error bars include only the statistical uncertainty and the error band includes both statistical and systematic uncertainties.

The E_T^{miss} definition, Eq. 4.2, explicitly depends on the selections applied to the objects entering the sum [460]. The objects must be reconstructed using mutually exclusive detector signals and so a priority is assigned to their reconstruction: electrons, photons [286, 426], hadronically-decaying taus [461, 462], muons, jets, and tracks, from highest-to-lowest priority. Lower-priority objects are typically removed when overlapping with higher-priority objects. Muons have little overlap with the other objects, the exception being non-isolated muons. Tracks associated to the primary vertex and matched to a jet are added to the soft term if the jet is removed [460].

The usage of fully-calibrated hard objects and tracks originating from only the primary vertex ensures that E_T^{miss} reconstruction is insensitive to pile-up effects [460]. Neutral particle signals from the calorimeter are significantly affected by pile-up and are therefore excluded from the soft term. Due to the limitations of detector acceptance and resolution, not all physics objects originating from the primary vertex are reconstructed, resulting in an observation bias towards nonzero E_T^{miss} values [460].

The E_T^{miss} energy response and resolution are measured in $Z \rightarrow \mu\bar{\mu}$ events, where the particle-level missing transverse momentum, $E_T^{\text{miss,true}}$, is zero, as well as in $W \rightarrow e\nu/\mu\nu$ and semileptonic $t\bar{t}$ events, where $E_T^{\text{miss,true}}$ is nonzero [460]. In latter case, the relative deviation of the reconstructed E_T^{miss} from its particle-level value is largest at low values of $E_T^{\text{miss,true}}$, where the observation bias leads to a nonzero offset in the reconstructed E_T^{miss} as shown in Fig. 4.22a. The corresponding resolution, given as the width of the $(E_T^{\text{miss}} - E_T^{\text{miss,true}})$ distribution, is shown in Fig. 4.22b. Generally, the resolution is poorer in $t\bar{t}$ final states than in $W \rightarrow e\nu/\mu\nu$ final states [460].

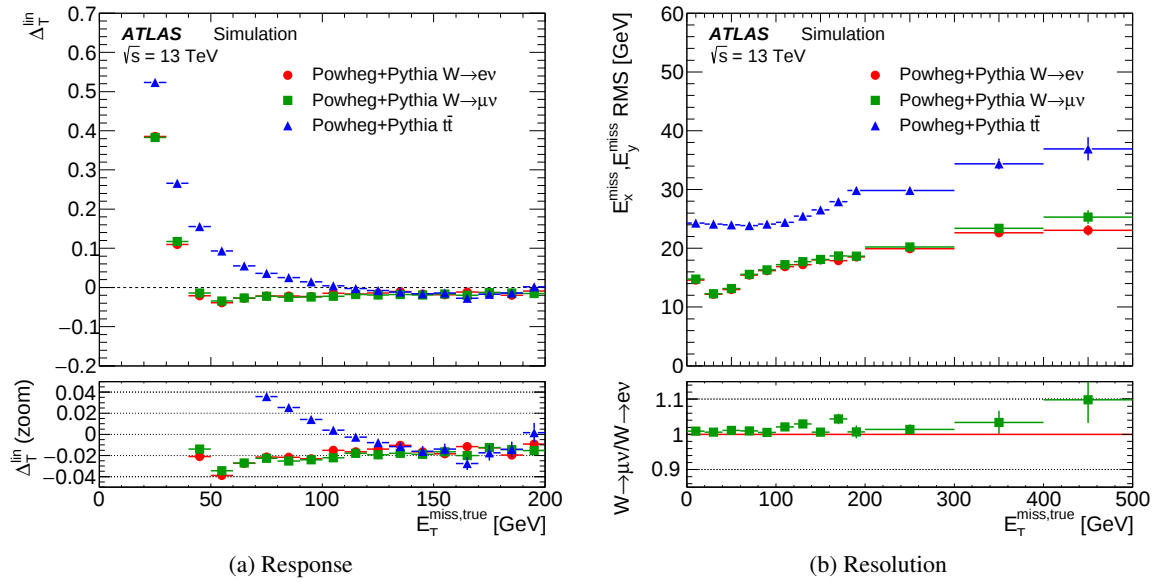


Figure 4.22: (a) Deviation of the E_T^{miss} response from linearity, $(E_T^{\text{miss}} - E_T^{\text{miss,true}})/E_T^{\text{miss,true}}$, and (b) the E_T^{miss} resolution, extracted as the width (“RMS”) of the $(E_T^{\text{miss}} - E_T^{\text{miss,true}})$ distribution, in simulated $W \rightarrow e\nu/\mu\nu$ and $t\bar{t}$ events [460]. Both the response and resolution are shown as functions of the particle-level missing transverse momentum, $E_T^{\text{miss,true}}$. The error bars indicate the statistical uncertainty.

An object-based E_T^{miss} significance [463] is used to reject events where the E_T^{miss} arises due to the mis-reconstruction of the physics objects entering the calculation and *not* due to the presence of neutrinos. The p_T and ϕ resolutions for the input objects, shown in Fig. 4.23a, are used to derive E_T^{miss} resolutions in the planes transverse and longitudinal to its direction [463]. The transverse and longitudinal resolutions are used as inputs to a hypothesis test asserting that the momentum carried by invisible particles is zero, and the associated

significance is calculated [463]. The object-based E_T^{miss} significance in $Z \rightarrow e\bar{e}$ events is shown in Fig. 4.23b.

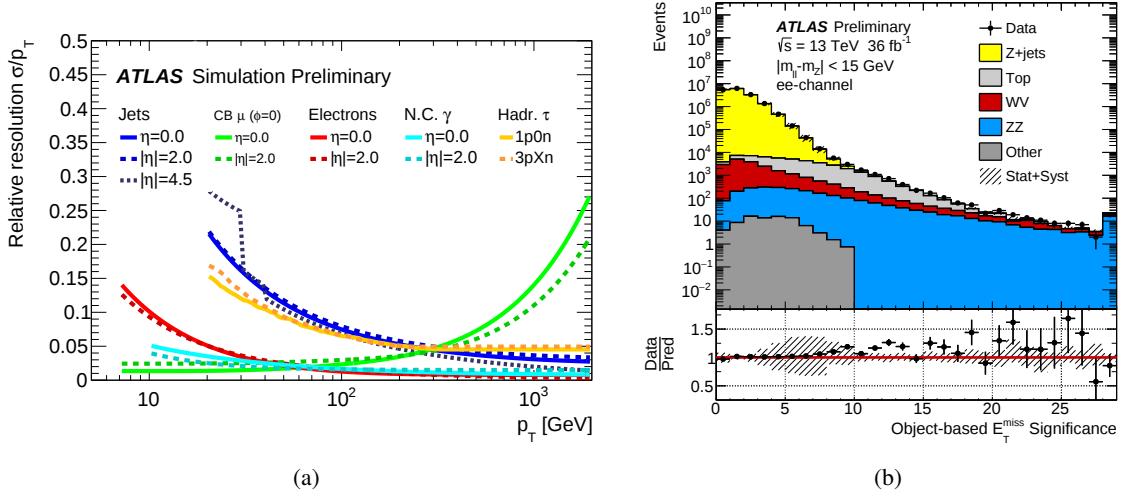


Figure 4.23: (a) Relative momentum resolutions for the physics objects entering E_T^{miss} reconstruction, including jets, combined muons (“CB μ ”), electrons, unconverted photons (“N.C. γ ”), and hadronically-decaying taus (Hadr. τ) [463]. (b) Distribution of the object-based E_T^{miss} significance in $Z \rightarrow e\bar{e}$ events [463]. The error band includes both statistical and systematic uncertainties.

4.2 Object Selection

The selections applied to the physics objects entering the $VH, H \rightarrow WW^*$ measurement are shown in Table 4.1.

Electrons are collected using the logical OR of single electron triggers. They are required to have a minimum p_T of 10 GeV and to be reconstructed in region $|\eta| < 2.47$ – electrons reconstructed in the transition region between the electromagnetic barrel and endcap calorimeters, $1.37 < |\eta| < 1.52$, are excluded. The primary electron tracks are required to pass association to the primary vertex: $|\Delta z \sin \theta| < 0.5 \text{ mm}$ and $|d_0|/\sigma_{d_0} < 5$. An author cut of 1 ensures electrons are reconstructed exclusively as electrons, removing possible electron-photon ambiguity due to photon conversions. The choice of identification and isolation working points is channel-dependent, chosen such that the expected sensitivity of the measurement is maximized. In the $2\ell/3\ell$ channels, electrons are required to pass Tight identification and PLImprovedTight isolation; in the 4ℓ channels, electrons are required to pass Medium identification and Loose isolation. In the 2ℓ same-sign WH and 3ℓ Z-depleted WH channels, electrons are additionally required to pass ECIDS in order to reject electrons with misidentified charge.

Muons are collected using the logical OR of single muon triggers. They are required to have a minimum p_T of 10 GeV and to be reconstructed in the region $|\eta| < 2.5$. The muon tracks are required to pass association to the primary vertex: $|\Delta z \sin \theta| < 0.5 \text{ mm}$ and $|d_0|/\sigma_{d_0} < 3$. Muons are required to pass Medium identification as well as PLImprovedTight and Loose isolation in the $2\ell/3\ell$ and 4ℓ channels, respectively.

Jets are required to have a minimum p_T of 20 GeV and to be reconstructed in the region $|\eta| < 4.5$. Jets with $20 < p_T < 60 \text{ GeV}$ in the region $|\eta| < 2.4$ are required to pass the Tight JVT working point, and jets with $20 < p_T < 60 \text{ GeV}$ in the region $2.5 < |\eta| < 4.5$ are required to pass the Loose fJVT working point. For the purposes of jet counting, jets are required to have a minimum p_T of 30 GeV. In the region $|\eta| < 2.5$, b -jets

are identified using the DL1r tagger at the 85% b -tagging efficiency working point. For the purposes of b -jet counting, jets are required to have a minimum p_T of 20 GeV.

Prior to the “final” selections described above, physics objects are required to initially pass a set of looser selections⁸ followed by an overlap removal procedure [464]. The overlap removal proceeds as follows:

- Electron-electron: if two electrons have overlapping clusters in the second sampling layer of the electromagnetic calorimeter or share a matched track, the electron with the lower p_T is removed;
- Electron-muon: if a combined muon shares an inner detector track with an electron, the electron is removed; if a calorimeter-tagged muon [432] shares an inner detector track with an electron, the muon is removed;
- Electron-jet: if $\Delta R(\text{jet}, \text{electron}) < 0.2$, the jet is removed. For any surviving jets, the electron is removed if $\Delta R(\text{jet}, \text{electron}) < \min(0.4, 0.04 + (10 \text{ GeV})/p_T^{\text{electron}})$;
- Muon-jet: if $\Delta R(\text{jet}, \text{muon}) < 0.2$, the jet is removed if it has less than three associated tracks or if both of the following conditions are met: $p_T^{\text{muon}}/p_T^{\text{jet}} > 0.5$ and $p_T^{\text{muon}}/\sum_{\text{jet}} p_T^{\text{track}} > 0.7$, where $\sum_{\text{jet}} p_T^{\text{track}}$ is the p_T -sum of the tracks contained within the jet with $p_T > 500 \text{ MeV}$. For any surviving jets, the muon is removed if $\Delta R(\text{jet}, \text{muon}) < \min(0.4, 0.04 + (10 \text{ GeV})/p_T^{\text{muon}})$.

If a jet is b -tagged, it is always saved from removal.

Table 4.1: The selections applied to all of the physics objects utilized by the $VH, H \rightarrow WW^*$ analysis. The details for each selection are described in the text. Additionally, all objects are required to pass overlap removal.

Object	Selection	2 ℓ /3 ℓ channels	4 ℓ channels
Electrons	Transverse momentum	$p_T > 10 \text{ GeV}$	
	Pseudorapidity	$(\eta < 1.37) \parallel (1.52 < \eta < 2.47)$	
	Impact parameters	$ \Delta z \sin \theta < 0.5 \text{ mm} \ \&\& \ d_0 /\sigma_{d_0} < 5$	
	Author	Author = 1	
	Identification	Tight	Medium
	Isolation	PLImprovedTight	Loose
Muons	Transverse momentum	$p_T > 10 \text{ GeV}$	
	Pseudorapidity	$ \eta < 2.5$	
	Impact parameters	$ \Delta z \sin \theta < 0.5 \text{ mm} \ \&\& \ d_0 /\sigma_{d_0} < 3$	
	Identification	Medium	
	Isolation	PLImprovedTight	Loose
Jets	Transverse momentum	$p_T > 20 \text{ GeV}$ ($p_T > 30 \text{ GeV}$ for counting)	
	Pseudorapidity	$ \eta < 4.5$	
	JVT	Tight	
	fJVT	Loose	
b -jets	Transverse momentum	$p_T > 20 \text{ GeV}$	
	Pseudorapidity	$ \eta < 2.5$	
	JVT	Tight	
	Flavour tagger	DL1r @ 85% b -tagging efficiency	

⁸The only differences are for leptons: electrons are required to satisfy $p_T > 10 \text{ GeV}$, $|\eta| < 2.47$, VeryLoose identification [286], $|\Delta z \sin \theta| < 0.5 \text{ mm}$, and $|d_0|/\sigma_{d_0} < 5$; muons are required to satisfy $p_T > 10 \text{ GeV}$, $|\eta| < 2.7$, Loose identification, $|\Delta z \sin \theta| < 0.5 \text{ mm}$, and $|d_0|/\sigma_{d_0} < 15$.

Chapter 5

Estimation of Data-Driven Backgrounds

The $VH, H \rightarrow WW^*$ analysis utilizes a data-driven estimation of the misidentified lepton background and the electron background with misidentified charge. Both backgrounds are attributed to incorrect object reconstruction, a phenomenon which is typically difficult to simulate. The usage of data-driven estimation, where the misidentification rates are directly measured in data, is a more reliable approach than simulation.

Misidentified leptons, also known as non-prompt or “fake” leptons, come from various sources, including jets incorrectly reconstructed as leptons and leptons resulting from decays which are not directly related to hard-scatter collision, including leptonically-decaying hadrons or γ -conversions. Fake leptons are often characterized by a high level of track- or calorimeter-based activity around the lepton, and so identification and isolation criteria are powerful tools for rejecting fake leptons.

Electrons with incorrectly measured or “misidentified” charge, also known as “charge-flip” electrons, are a result of an incorrect determination of the sign of an electron track’s curvature or the usage of an incorrectly matched electron track with curvature opposite to that of the “true” track. In either case, the reconstructed charge of an electron is “flipped”. Charge-flip electrons are rejected using electron charge identification selection (ECIDS).

In this chapter, the estimation of the data-driven backgrounds is described. The fake factor method (Section 5.1) is used in $2\ell/3\ell$ channels (Section 5.2) for estimating the fake muon background. The fake electron background is estimated using a scale factor method multiplying MC predictions. For the 4ℓ channel (Section 5.3), the fake factor method is used for the estimating both flavours of the fake lepton background. In the 2ℓ same-sign WH channel, the electron charge-flip rates in data are measured using Z +jets events (Section 5.4). The muon charge-flip rates are negligible in the p_T range relevant to this analysis.

5.1 The Fake Factor Method

The fake factor method is a widely-used technique for estimating fake lepton backgrounds. It begins with the construction of a control region targeting the source of the fake leptons. Within the fake lepton control region, two sub-regions are defined: the identified (ID) control region, where all of the leptons in the event are required to pass the nominal lepton selections, and the anti-identified (anti-ID) control region, where at least one of the leptons is required to pass a looser set of lepton selections. The usage of a looser set of lepton selections enriches the anti-ID control region in fake leptons.

The number of fake events in the ID control region divided by the number of fake events in the anti-ID

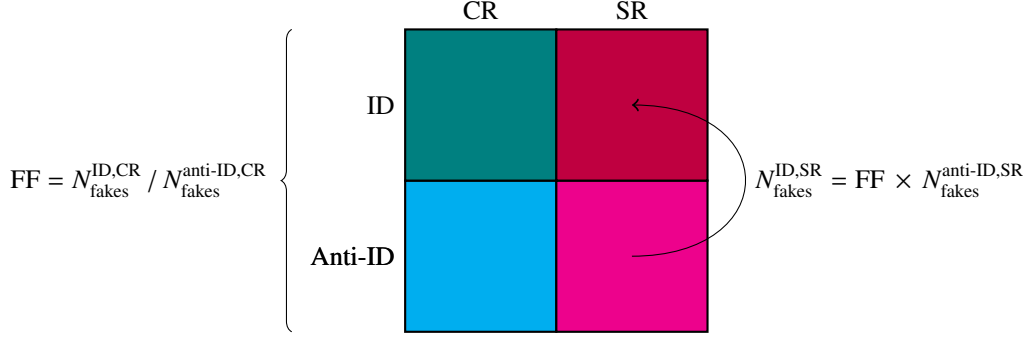


Figure 5.1: A pictorial representation of the fake factor method: fake lepton control regions (“CR”) with anti-ID (“anti-ID”) or ID (“ID”) selections applied to one lepton are used to calculate the fake factor (“FF”). The fake factor is then applied to the anti-ID fake yields in the signal region (“SR”) to estimate the ID fake yields. Event yields are denoted by N .

control region corresponds to an anti-ID-to-ID extrapolation factor known as the fake factor. In particular, the data-driven fake factor (FF), measured using the fake lepton control region, is given by:

$$\text{FF} = \frac{N_{\text{fakes}}^{\text{ID,CR}}}{N_{\text{fakes}}^{\text{anti-ID,CR}}} = \frac{N_{\text{data}}^{\text{ID,CR}} - N_{\text{EW}}^{\text{ID,CR}}}{N_{\text{data}}^{\text{anti-ID,CR}} - N_{\text{EW}}^{\text{anti-ID,CR}}}, \quad (5.1)$$

where N_{fakes} is the data-driven estimate of the number of fake events: $N_{\text{fakes}} = N_{\text{data}} - N_{\text{EW}}$, where N_{data} is the measured number of data events and N_{EW} is the simulated number of events with *only* electroweak (EW), or prompt, leptons. Whether a given event yield corresponds to the ID or anti-ID control region is indicated by the superscript.

A key assumption of the fake factor method is that the fake lepton efficiency only depends on the properties of the lepton and not on the remainder of the event, as lepton reconstruction uses information from a localized area of the detector. The implication of this assumption is that the fake factor derived in the fake lepton control region is equally applicable in a channel’s signal region which, in this context, refers to any region where the fake lepton background estimate will be applied. One may analogously define ID and anti-ID sub-regions of the signal region, naturally leading to an expression for the number of fake events in the ID signal region:

$$N_{\text{fakes}}^{\text{ID,SR}} = \text{FF} \times N_{\text{fakes}}^{\text{anti-ID,SR}} = \text{FF} \times (N_{\text{data}}^{\text{anti-ID,SR}} - N_{\text{EW}}^{\text{anti-ID,SR}}). \quad (5.2)$$

Whether a given event yield corresponds to the ID or anti-ID signal region is indicated by the superscript. A diagram summarizing the fake factor method is shown in Fig. 5.1.

5.2 Misidentified Lepton Estimation for the $2\ell/3\ell$ Channels

For the 2ℓ same-sign WH and 3ℓ WH channels, the primary sources of fakes are V +jets, $t\bar{t}$, and V + γ ; for the 2ℓ DFOS VH channel, the primary sources of fakes are W +jets and W + γ .

The ID and anti-ID selections for the $2\ell/3\ell$ channels are detailed in Table 5.1. Anti-ID muons have a looser cut on the significance of the transverse impact parameter and an absent isolation working point selection with respect to ID muons; anti-ID electrons have looser identification and isolation working point selections as well as no author cut with respect to ID electrons. The absence of the author cut increases the fraction of

electrons resulting from photon conversions. Additionally, all lepton pairs are required to pass a minimum angular separation cut, $\Delta R_{\ell\ell} > 0.1$. This is because fake events were sometimes found to have overlapping leptons, $\Delta R_{\ell\ell} < 0.1$. Overlapping leptons in prompt lepton events, on the other hand, were found to be rare.

Table 5.1: The ID and anti-ID criteria applied to leptons for the fake lepton estimation in $2\ell/3\ell$ channels. For electrons, the crack region between the barrel and endcap electromagnetic calorimeters, $1.37 < |\eta| < 1.52$, is excluded. An author cut of 1 ensures electrons are reconstructed exclusively as electrons. ECIDS is *only* applied to electrons in the 2ℓ same-sign WH and 3ℓ Z-depleted WH channels. “Veto ID” indicates that a lepton is not classified as anti-ID if it passes the ID criteria.

Criteria	Muons		Electrons	
	ID	Anti-ID	ID	Anti-ID
p_T [GeV]		> 15		> 15
$ \eta $		< 2.5		< 2.47 , excluding crack region
$ \Delta z \sin \theta $ [mm]		< 0.5		< 0.5
$ d_0 /\sigma_{d_0}$	< 3			< 5
Author		–	Author = 1	–
Identification		Medium	Tight	Loose
Isolation	PLImprovedTight	–	PLImprovedTight	Loose && PLVLoose
ECIDS	–	–	No (Yes)	–
Veto ID	–	Yes	–	Yes

5.2.1 Misidentified Muons

The fake muon background is estimated using the fake factor method for the $2\ell/3\ell$ channels. A 3ℓ control region enriched in $t\bar{t}$ events is constructed for the estimation of the fake factors. The selections for the $t\bar{t}$ control region are shown in Table 5.2. Two same-sign muons, one electron, and one b -jet are selected, making the control region orthogonal to the 3ℓ Z-depleted WH signal region.

The fake factors are measured in two fake muon p_T bins: $15 < p_T < 30$ GeV and $p_T > 30$ GeV. In the ID control region, the subleading muon is selected as the fake muon, as the muon from the decay of the b quark is expected to be softer than that from the decay of the W boson. In the anti-ID control region, the anti-ID muon is selected as the fake muon.

Table 5.2: Event selections for the $t\bar{t}$ control regions used to estimate the muon fake factors in $2\ell/3\ell$ channels.

$t\bar{t}$ control region	
ID control region	Anti-ID control region
Exactly 3 leptons with at least 1 ID lepton matched to a single lepton trigger	
Lepton flavour/charge: $\mu^\pm e^\mp \mu^\pm$	
$N_{b\text{-jets}} = 1$	
3 ID leptons	2 ID leptons and 1 anti-ID muon

Following the procedure used in the ATLAS same-sign WW measurement [436], the p_T of the anti-ID muon, $p_T^{\ell_{\text{anti-ID}}}$, is corrected using its $p_T^{\text{varcone30}}$ isolation variable, typically 5 GeV or 20% of the anti-ID muon’s p_T :

$$p_T^{\ell_{\text{anti-ID}}, \text{corr}} = p_T^{\ell_{\text{anti-ID}}} + p_T^{\text{varcone30}}(\ell_{\text{anti-ID}}). \quad (5.3)$$

This is done to reduce the p_T spectrum differences between the anti-ID muon and the underlying jet which has

been misidentified as the anti-ID muon, as the muon carries only a fraction of the jet's total momentum. Doing so also reduces the flavour dependence on the fake lepton source used to estimate the fake factors: $t\bar{t}$ versus Z +jets. For the purposes of the $p_T > 15$ GeV selection and the binning of the fake factors, it is the *corrected* anti-ID muon p_T which is used.

Fig. 5.2 shows distributions of the subleading muon p_T and the corrected anti-ID muon p_T in the ID and anti-ID $t\bar{t}$ control regions, respectively. The fake factors are extracted from these distributions by fitting a likelihood function¹ parameterized in terms of the fake factors, $\vec{\text{FF}}$, the nuisance parameters on the electroweak backgrounds, $\vec{\theta}$, and the event yields in the anti-ID control regions, $\vec{v}$.² The likelihood function is written as a product of Poisson probability density functions (PDFs) for each p_T bin of the ID and anti-ID control regions and Gaussian constraints for the theoretical systematics on the largest electroweak backgrounds:

$$\mathcal{L}(\vec{\text{FF}}, \vec{\theta}, \vec{v}) = \left[\prod_i^{p_T \text{ bins}} P(N_i^{\text{data,ID}} | N_i^{\text{predicted,ID}}(\vec{\text{FF}}_i, \vec{\theta}, v_i)) \times P(N_i^{\text{data,anti-ID}} | v_i) \right] \times \left[\prod_k^{\text{uncertainties}} G(\tilde{\theta}_k | \theta_k) \right], \quad (5.4)$$

with:

$$N_i^{\text{predicted,ID}}(\vec{\text{FF}}_i, \vec{\theta}, v_i) = N_i^{\text{EW,ID}}(\vec{\theta}) + \text{FF}_i \times (v_i - N_i^{\text{EW,anti-ID}}(\vec{\theta})), \quad (5.5)$$

where:

- P is a Poisson PDF and G is a Gaussian constraint;
- $N_i^{\text{data,ID}}$ and $N_i^{\text{data,anti-ID}}$ are the data yields in p_T bin i of the ID and anti-ID control regions, respectively;
- $N_i^{\text{predicted,ID}}(\vec{\text{FF}}_i, \vec{\theta}, v)$ is the predicted yield in p_T bin i of the ID control region;
- v_i is the (floating) event yield in p_T bin i of the anti-ID control region, constrained by $N_i^{\text{data,anti-ID}}$;
- $N_i^{\text{EW,ID}}(\vec{\theta})$ and $N_i^{\text{EW,anti-ID}}(\vec{\theta})$ are the MC electroweak background yields in p_T bin i of the ID and anti-ID control regions, respectively;
- FF_i is the fake factor in p_T bin i .

The Gaussian constraints allow the yields for the largest components of the electroweak background, $t\bar{t}V$ and $t\bar{t}H$, to float in the fit. The size of the fluctuations are dictated by the uncertainties on their respective cross sections, $\Delta\sigma_{t\bar{t}V} = 13\%$ [465] and $\Delta\sigma_{t\bar{t}H} = 13\%$ [465]. Other components of the electroweak background include WZ , ZZ , VVV , VH , and tZ ; however, these components are constant in the fit.

The fitted fake factors are shown in Table 5.3. Two sources of uncertainty on the fake factors are quantified: the statistical and cross section uncertainty and the non-closure uncertainty.

The statistical and cross section uncertainty, $\Delta\text{FF}_{\text{stat+XS}}$, is a result of the limited statistics in the fake muon control regions as well as the uncertainties on the $t\bar{t}V$ and $t\bar{t}H$ cross sections. It is given by the 68% confidence interval on best-fit values of the fake factors.

The non-closure uncertainty, $\Delta\text{FF}_{\text{non-closure}}$, is calculated by comparing the MC prediction for the fake muon background in the ID 3ℓ Z -dominated WH signal region to the estimate obtained by scaling the MC

¹The likelihood function describes the “likelihood” of measuring a particular set of parameters defining the measurement model given the observed data. More details on likelihood functions are given in Chapter 8.

²Throughout this chapter as well as the following chapters, an overhead arrow (e.g., $\vec{\text{FF}}$ or $\vec{\theta}$) is used to denote a vector.

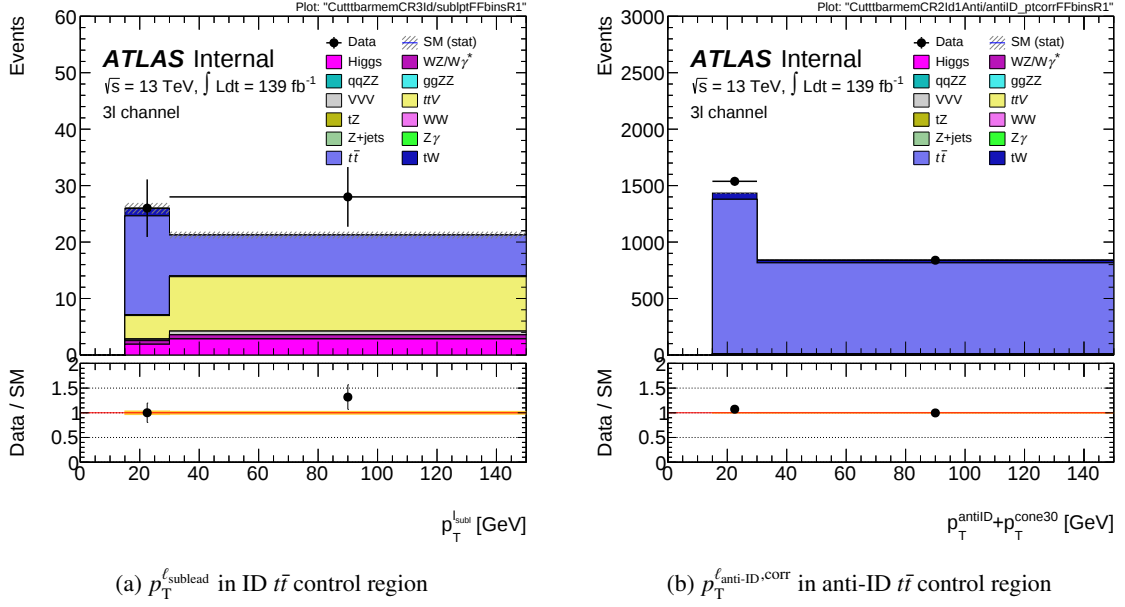


Figure 5.2: Distributions of (a) the subleading muon p_T and (b) the corrected anti-ID muon p_T in the $t\bar{t}$ control regions used for the fake muon background estimation in the $2\ell/3\ell$ channels. The error bands indicate the statistical uncertainty. Here, $t\bar{t}$ corresponds to the MC prediction for the fake lepton background.

prediction in the anti-ID signal region by MC-driven fake factors. Per p_T bin, the relative difference in the fake yield is symmetrized, summed in quadrature with the statistical uncertainty on the MC prediction in the ID signal region, and assigned as the uncertainty.

Table 5.3: Data-driven fake factors (FFs) for the fake muon estimation in the $2\ell/3\ell$ channels. Also shown are the statistical and theoretical cross section ($\Delta\text{FF}_{\text{stat+XS}}$), the non-closure ($\Delta\text{FF}_{\text{non-closure}}$), and the total ($\Delta\text{FF}_{\text{total}}$) uncertainties for each fake factor. The total uncertainty is given by the quadrature sum of $\Delta\text{FF}_{\text{stat+XS}}$ and $\Delta\text{FF}_{\text{non-closure}}$.

p_T bin	FF	$\Delta\text{FF}_{\text{stat+XS}}$ [%]	$\Delta\text{FF}_{\text{non-closure}}$ [%]	$\Delta\text{FF}_{\text{total}}$ [%]
$15 < p_T < 30 \text{ GeV}$	0.0124	19.2	72.7	75.2
$30 \text{ GeV} < p_T$	0.0170	27.3	23.0	35.7

5.2.2 Misidentified Electrons

In the $2\ell/3\ell$ channels, control regions independently measuring the fake factors for each source of fake electrons could not be constructed – instead, the scale factor method is used. Control regions enriched in fake electrons are split into sub-regions with ID or anti-ID criteria applied to one of the electrons. The normalization of the MC prediction for each source of fake electrons is measured by fitting the MC predictions in the anti-ID control regions to data in the ID control regions. The normalization factor for a particular source of fake leptons is referred to as its scale factor (SF). By multiplying the scale factor for each source onto the corresponding MC prediction in the anti-ID signal region, an estimate of fake electron background in the ID signal region is obtained. For example, if there is only one source of fake electrons, then the estimated number of fake electron events in the ID signal region is given by:

$$N_{\text{fakes}}^{\text{ID,SR}} = \text{SF} \times N_{\text{fakes}}^{\text{anti-ID,SR}} = \frac{N_{\text{data}}^{\text{ID,CR}} - N_{\text{EW}}^{\text{ID,CR}}}{N_{\text{fakes}}^{\text{anti-ID,CR}}} \times N_{\text{fakes}}^{\text{anti-ID,SR}}, \quad (5.6)$$

where $N_{\text{fakes}}^{\text{anti-ID,SR}}$ and $N_{\text{fakes}}^{\text{anti-ID,CR}}$ are the MC predictions for the number of fake electron events in the anti-ID signal and control regions, respectively.

Three 3ℓ control regions are constructed, one for each source of fake electrons: γ -conversion, light-flavour, and heavy-flavour fakes. The selections for each control region are shown in Table 5.4.

The $Z+\gamma$ control region targets γ -conversion fakes from photons in final state radiation, and the $Z+\text{jets}$ control region targets light-flavour fakes from misidentified jets. Both select two opposite-sign muons and one electron. An upper cut on the missing transverse momentum, $E_{\text{T}}^{\text{miss}} < 30 \text{ GeV}$, ensures the control regions are orthogonal to the 3ℓ Z -dominated WH signal region. Requiring zero b -jets reduces each control region's $t\bar{t}$ contamination, and a lower cut on the minimum opposite-sign dilepton invariant mass, $m_{\ell\ell,\text{OS}}^{\text{min}} > 12 \text{ GeV}$, is used to remove low-mass resonances. Cuts on the transverse mass of the W boson:

$$m_{\text{T}}^W = \sqrt{2p_{\text{T}}^{\ell_{\text{fake}}} E_{\text{T}}^{\text{miss}} (1 - \cos \Delta\phi_{\ell_{\text{fake}},\text{miss}})}, \quad (5.7)$$

and the trilepton invariant mass, $m_{\ell\ell\ell}$, are used to further improve each control region's purity. Finally, cuts on the dimuon invariant mass, $m_{\mu\mu}$, are used to select softer muons, in the case of non-resonant Z production via $Z+\gamma$, or harder muons, in the case of resonant Z production via $Z+\text{jets}$.

The $t\bar{t}$ control region targets heavy-flavour fakes from $t\bar{t} \rightarrow W^+bW^-\bar{b}$ decays. Two same-sign electrons and one opposite sign muon are selected, which ensures that the muon is prompt and one of the electrons is non-prompt. Requiring one b -jet improves the control region's $t\bar{t}$ purity. Additionally, the lepton flavour/charge and b -jet selections ensure the region is orthogonal to the $t\bar{t}$ control region used to extract the muon fake factors and the 3ℓ Z -depleted WH signal region.

Within each control region, the MC events are split into γ -conversion, light-flavour, and heavy-flavour sources using a truth classifier. Per event, the fake electron is associated to a truth-level particle using a matching procedure based on minimum distance in ΔR , and the type of the truth-level particle, or the type of its progenitor, is used to assign the electron's fake source. In the ID control regions, the subleading electron is selected as the fake electron; in the anti-ID control regions, the anti-ID electron is selected as the fake electron.

Independent sets of scale factors are derived for electron ID with and without ECIDS. Additionally, each scale factor is measured in three fake electron p_{T} bins: $15 < p_{\text{T}} < 20 \text{ GeV}$, $20 < p_{\text{T}} < 30 \text{ GeV}$, and $p_{\text{T}} > 30 \text{ GeV}$.

Table 5.4: Event selections for the $Z+\gamma$, $Z+\text{jets}$, and $t\bar{t}$ control regions (CRs) used to estimate the electron scale factors in $2\ell/3\ell$ channels.

$Z+\gamma$ CR		$Z+\text{jets}$ CR		$t\bar{t}$ CR	
ID CR	Anti-ID CR	ID CR	Anti-ID CR	ID CR	Anti-ID CR
Exactly 3 leptons with at least 1 ID lepton matched to a single lepton trigger					
Lepton flavour/charge: $\mu^\pm\mu^\mp e$			Lepton flavour/charge: $e^\pm\mu^\mp e^\pm$		
$N_{b\text{-jets}} = 0$			$N_{b\text{-jets}} = 1$		
$E_{\text{T}}^{\text{miss}} < 30 \text{ GeV}$					
$m_{\ell\ell,\text{OS}}^{\text{min}} > 12 \text{ GeV}$					
$m_{\text{T}}^W < 40 \text{ GeV}$					
$ m_{\ell\ell\ell} - m_Z < 10 \text{ GeV}$		$m_{\ell\ell\ell} < 200 \text{ GeV}$			
$40 < m_{\mu\mu} < 80 \text{ GeV}$		$80 < m_{\mu\mu} < 100 \text{ GeV}$			
3 ID leptons	2 ID leptons and 1 anti-ID electron	3 ID leptons	2 ID leptons and 1 anti-ID electron	3 ID leptons	2 ID leptons and 1 anti-ID electron

Fig. 5.3 shows distributions of the subleading electron p_T and anti-ID electron p_T in the ID and anti-ID control regions, respectively, for each source of fake electrons. The scale factors are extracted from these distributions by fitting a likelihood function parameterized in terms of the scale factors, $\vec{SF}$, and the nuisance parameters on the electroweak backgrounds, $\vec{\theta}$. The likelihood function is written as a product of Poisson PDFs for each p_T bin of each control region and Gaussian constraints for the theoretical systematics on the electroweak backgrounds:

$$\mathcal{L}(\vec{SF}, \vec{\theta}) = \left[\prod_i^{p_T \text{ bins}} \prod_j^{\text{control regions}} P(N_{ij}^{\text{data,ID}} | N_{ij}^{\text{predicted,ID}}(\vec{SF}, \vec{\theta})) \right] \times \left[\prod_k^{\text{uncertainties}} G(\tilde{\theta}_k | \theta_k) \right], \quad (5.8)$$

with:

$$N_{ij}^{\text{predicted,ID}}(\vec{SF}, \vec{\theta}) = N_{ij}^{\text{EW,ID}}(\vec{\theta}) + \sum_{\ell}^{\text{fake sources}} (SF_{i\ell} \times N_{ij}^{\text{fake } \ell, \text{anti-ID}}), \quad (5.9)$$

where:

- P is a Poisson PDF and G is a Gaussian constraint;
- $N_{ij}^{\text{data,ID}}$ is the data yield in p_T bin i of the ID control region j ;
- $N_{ij}^{\text{predicted,ID}}(\vec{SF}, \vec{\theta})$ is the predicted yield in p_T bin i of the ID control region j ;
- $N_{ij}^{\text{EW,ID}}(\vec{\theta})$ is the MC electroweak background yield in p_T bin i of the ID control region j ;
- $N_{ij}^{\text{fake } \ell, \text{anti-ID}}$ is the MC yield for fake source ℓ in p_T bin i of anti-ID control region j ;
- $SF_{i\ell}$ is the scale factor for fake source ℓ in p_T bin i .

The Gaussian constraints allow the yields for all components of the electroweak background to float in the fit. The size of the fluctuations are dictated by the uncertainties on their respective cross sections, $\Delta\sigma_{\bar{t}tV} = 13\%$ [465], $\Delta\sigma_{\bar{t}tH} = 13\%$ [465], $\Delta\sigma_{WZ} = 5\%$ [466], $\Delta\sigma_{ZZ} = 6\%$ [467], $\Delta\sigma_{VVV} = 35\%$ [468].

The fitted scale factors are shown in Table 5.5. Two sources of uncertainty on the scale factors are quantified: the statistical and cross section uncertainty and the non-closure uncertainty.

The statistical and cross section uncertainty, $\Delta SF_{\text{stat+XS}}$, is a result of the limited statistics in the fake electron control regions as well as the uncertainties on the cross sections for each component of the electroweak background. It is given by the 68% confidence interval on best-fit values of the scale factors. If the fit converges to 0 for a particular scale factor, the 68% confidence interval is instead used as the central value of the scale factor and a 100% uncertainty is assigned.

The non-closure uncertainty, $\Delta SF_{\text{non-closure}}$, has two contributions. The first contribution is calculated by comparing the MC prediction for the fake electron background in the ID Z +jets control region to the estimate obtained by scaling the MC prediction in the anti-ID control region by MC-driven scale factors. Per p_T bin, the relative difference in the fake yield is symmetrized and summed in quadrature with the statistical uncertainty on the MC prediction in the ID control region. Each source of fake electrons shares the same uncertainty in a given p_T bin. The second contribution is calculated by comparing the MC prediction for the fake electron background in the ID μe same-sign WH signal region to the estimate obtained by scaling the MC prediction for the fake electron background in the anti-ID signal region by MC-driven scale factors. Per p_T bin, the relative difference in the fake yield is symmetrized and summed in quadrature with the statistical uncertainty

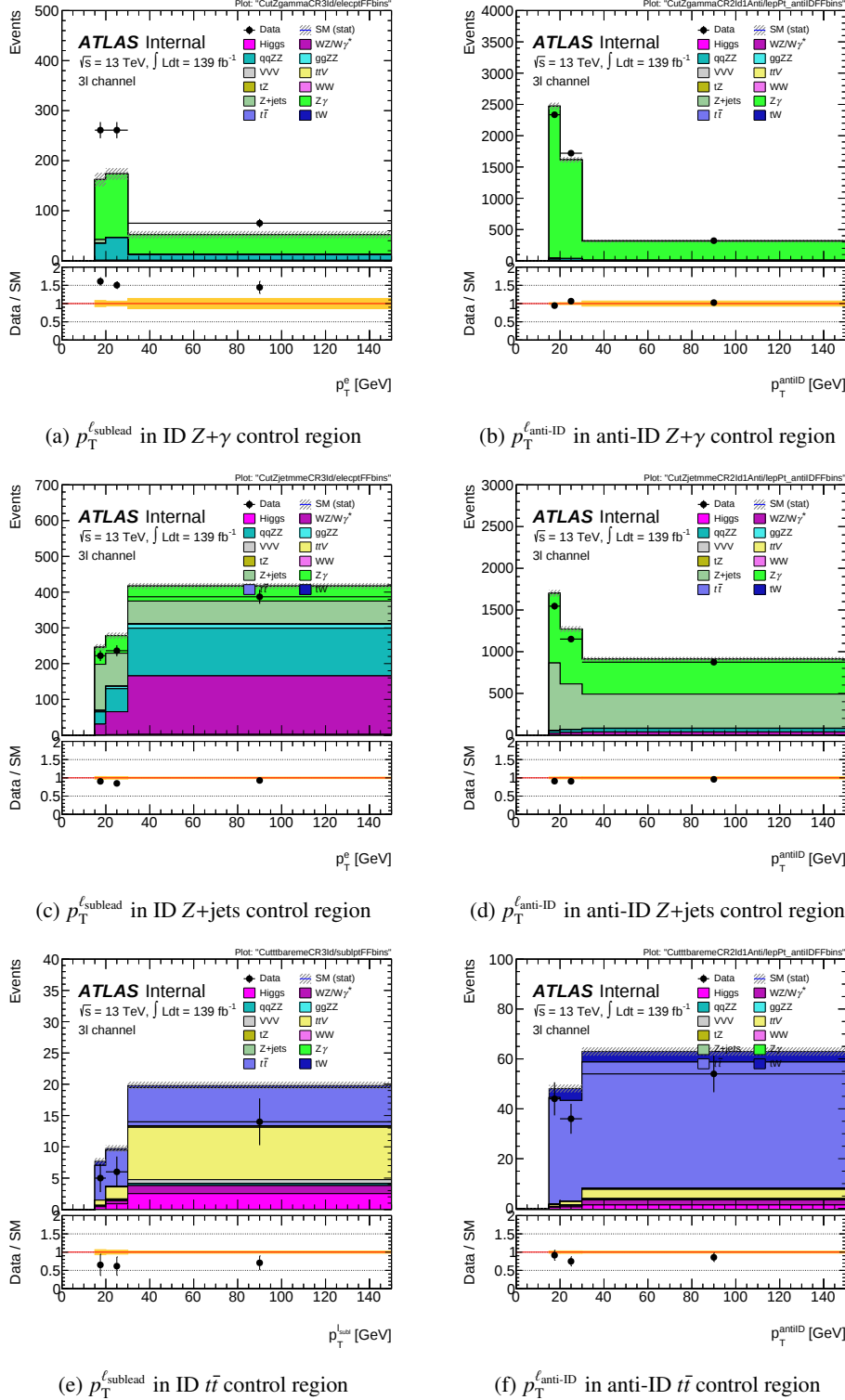


Figure 5.3: Distributions of (a, c, e) subleading electron p_T and (b, d, f) anti-ID electron p_T in the $Z+\gamma$, $Z+\text{jets}$, and $t\bar{t}$ control regions used for the fake electron background estimation in the $2\ell/3\ell$ channels. The error bands indicate the statistical uncertainty. These distributions correspond to electron ID without ECIDS. In the ID control regions, $Z+\gamma$, $Z+\text{jets}$, and $t\bar{t}$ correspond to the MC prediction for the fake lepton background, explaining the imperfect data-MC agreement.

on the MC prediction in the ID signal region. Each source of fake electrons has its own set of uncertainties, as the ID signal region is split into sub-regions of γ -conversion, light-flavour, and heavy-flavour fakes based on the source of the subleading electron. Both contributions are summed in quadrature to yield the non-closure uncertainty.

Table 5.5: Data-driven scale factors (SFs) for the fake electron estimation in the $2\ell/3\ell$ channels. Also shown are the statistical and cross section ($\Delta\text{SF}_{\text{stat+XS}}$), the non-closure ($\Delta\text{SF}_{\text{non-closure}}$), and the total ($\Delta\text{SF}_{\text{total}}$) uncertainties for each scale factor. The total uncertainty is given by the quadrature sum of $\Delta\text{SF}_{\text{stat+XS}}$ and $\Delta\text{SF}_{\text{non-closure}}$. These scale factors correspond to electron ID without ECIDS – the scale factors for electron ID with and without ECIDS agree within uncertainty, and so only the latter are shown for brevity.

Fake flavour	p_T bin	SF	$\Delta\text{SF}_{\text{stat+XS}}$ [%]	$\Delta\text{SF}_{\text{non-closure}}$ [%]	$\Delta\text{SF}_{\text{total}}$ [%]
γ -conversion	$15 < p_T < 20 \text{ GeV}$	0.091	6.5	79.0	79.3
	$20 < p_T < 30 \text{ GeV}$	0.144	6.0	25.1	25.8
	$30 \text{ GeV} < p_T$	0.207	9.1	27.6	29.1
Heavy	$15 < p_T < 20 \text{ GeV}$	0.064	100.0	61.3	117.3
	$20 < p_T < 30 \text{ GeV}$	0.079	100.0	49.9	111.8
	$30 \text{ GeV} < p_T$	0.049	100.0	57.4	115.3
Light	$15 < p_T < 20 \text{ GeV}$	0.113	25.2	43.9	50.6
	$20 < p_T < 30 \text{ GeV}$	0.025	100.0	49.4	111.5
	$30 \text{ GeV} < p_T$	0.025	100.0	53.8	113.6

5.2.3 Validation

The data-driven fake factors and scale factors are applied to same-sign validation regions in order to assess the data-MC fidelity of the fake lepton background estimate. The selections for each validation region are shown in Table 5.6.

The $\mu^\pm\mu^\pm$ validation region requires two same-sign muons. It is further divided into a b -veto sub-region by requiring zero b -jets and zero jets, where the jet veto is applied to ensure orthogonality with the $\mu\mu$ same-sign WH signal region, and b -tag sub-region by requiring one b -jet. The dominant fake lepton backgrounds in the b -veto and b -tag sub-regions are W +jets and $t\bar{t}$, respectively.

The $e^\pm\mu^\pm/\mu^\pm e^\pm$ validation region requires two same-sign, different-flavour leptons. It is further divided into a b -veto sub-region by requiring zero b -jets and zero jets, where the jet veto is applied to ensure orthogonality with the $e\mu/\mu e$ same-sign WH signal region, and b -tag sub-region by requiring one b -jet. The dominant fake lepton backgrounds in the b -veto and b -tag sub-regions are V +jets/ $V+\gamma$ and $t\bar{t}$, respectively.

Table 5.6: Event selections for the same-sign validation regions (VRs) used to validate the fake lepton background estimate in the $2\ell/3\ell$ channels.

$\mu^\pm\mu^\pm$ VR		$e^\pm\mu^\pm/\mu^\pm e^\pm$ VR	
b -veto VR	b -tag VR	b -veto VR	b -tag VR
Exactly ID 2 leptons with at least 1 lepton matched to a single lepton trigger			
Lepton flavour/charge: $\mu^\pm\mu^\pm$		Lepton flavour/charge: $e^\pm\mu^\pm/\mu^\pm e^\pm$	
$N_{b\text{-jets}} = 0$ && $N_{\text{jets}} = 0$	$N_{b\text{-jets}} = 1$	$N_{b\text{-jets}} = 0$ && $N_{\text{jets}} = 0$	$N_{b\text{-jets}} = 1$

The subleading lepton p_T is sensitive to the modelling of the fake lepton background, as the subleading lepton is expected to be the fake lepton. Fig. 5.4 shows distributions of the subleading lepton p_T in each of the

same-sign validation regions, where the fake lepton background is obtained using the data-driven estimation. Good data-MC agreement is observed in all distributions.

5.3 Misidentified Lepton Estimation for the 4ℓ ZH Channel

For the 4ℓ ZH channel, the primary sources of fakes are WZ and $t\bar{t}W/tZ$, each with one misidentified lepton in the final state, and Z +jets and $t\bar{t}/Wt$, each with two misidentified leptons in the final state. The latter are referred to as *double* fakes.

The ID and anti-ID selections for the 4ℓ channel are detailed in Table 5.7. Anti-ID muons have a looser cut on the transverse impact parameter significance, a looser identification working point selection, and an absent isolation working point selection with respect to ID muons; anti-ID electrons have a looser identification working point selection, an absent isolation selection, and no author cut with respect to ID electrons.

Table 5.7: The ID and anti-ID criteria applied to leptons for the fake lepton estimation in the 4ℓ ZH channel. For electrons, the crack region between the barrel and endcap electromagnetic calorimeters, $1.37 < |\eta| < 1.52$, is excluded. An author cut of 1 ensures electrons are reconstructed exclusively as electrons. “Veto ID” indicates that a lepton is not classified as anti-ID if it passes the ID criteria.

Criteria	Muons		Electrons	
	ID	Anti-ID	ID	Anti-ID
p_T [GeV]		> 10		> 10
$ \eta $		< 2.5		< 2.47 , excluding crack region
$ \Delta z \sin \theta $ [mm]		< 0.5		< 0.5
$ d_0 /\sigma_{d_0}$	< 3	< 15		< 5
Author		–	Author = 1	–
Identification	Medium	Loose	Medium	VeryLoose
Isolation	Loose	–	Loose	–
Veto ID	–	Yes	–	Yes

Both the fake muon and fake electron backgrounds are estimated using the fake factor method, proceeding in a manner similar to the estimation of fake muon background for the $2\ell/3\ell$ channels. The fake factors are not binned in the kinematics of fake lepton due to limited data statistics. With the inclusion of double fakes, Eq. 5.2 is modified as:

$$N_{\text{fakes}}^{4\ell, \text{SR}} = \text{FF} \times N_{\text{fakes}}^{3\ell+\ell, \text{SR}} - \text{FF}^2 \times N_{\text{fakes}}^{2\ell+2\ell, \text{SR}}, \quad (5.10)$$

where $N_{\text{fakes}}^{n\ell+m\ell, \text{SR}}$ denotes the fake yield in the signal region with n ID leptons and m anti-ID leptons. Accounting for double fakes constitutes a *correction* to the total fake yield. Regardless of the number of anti-ID leptons selected in a given region, the data-driven estimation of the fake yield proceeds in the same manner as before: $N_{\text{fakes}}^{n\ell+m\ell, \text{SR}} = N_{\text{data}}^{n\ell+m\ell, \text{SR}} - N_{\text{EW}}^{n\ell+m\ell, \text{SR}}$.

A 3ℓ control region enriched in Z +jets events is constructed for the estimation of the fake factors. The selections for the Z +jets control region are shown in Table 5.8. At least one same-flavour, opposite-sign (SFOS) lepton pair is required, selecting the pair, ℓ_0 and ℓ_1 , with an invariant mass closest to that of the Z boson as the Z boson candidate. The leptons belonging to this pair, referred to as the Z -tagged pair, must pass ID selections and at least one of them must be matched to a single lepton trigger. The invariant mass of the Z candidate is required to be within a window around the mass of the Z boson: $70 < m_{\ell_0\ell_1} < 110$ GeV, in the case of $Z \rightarrow \mu\bar{\mu}$, or $80 < m_{\ell_0\ell_1} < 110$ GeV, in the case of $Z \rightarrow e\bar{e}$. The Z -tagged leptons are also required

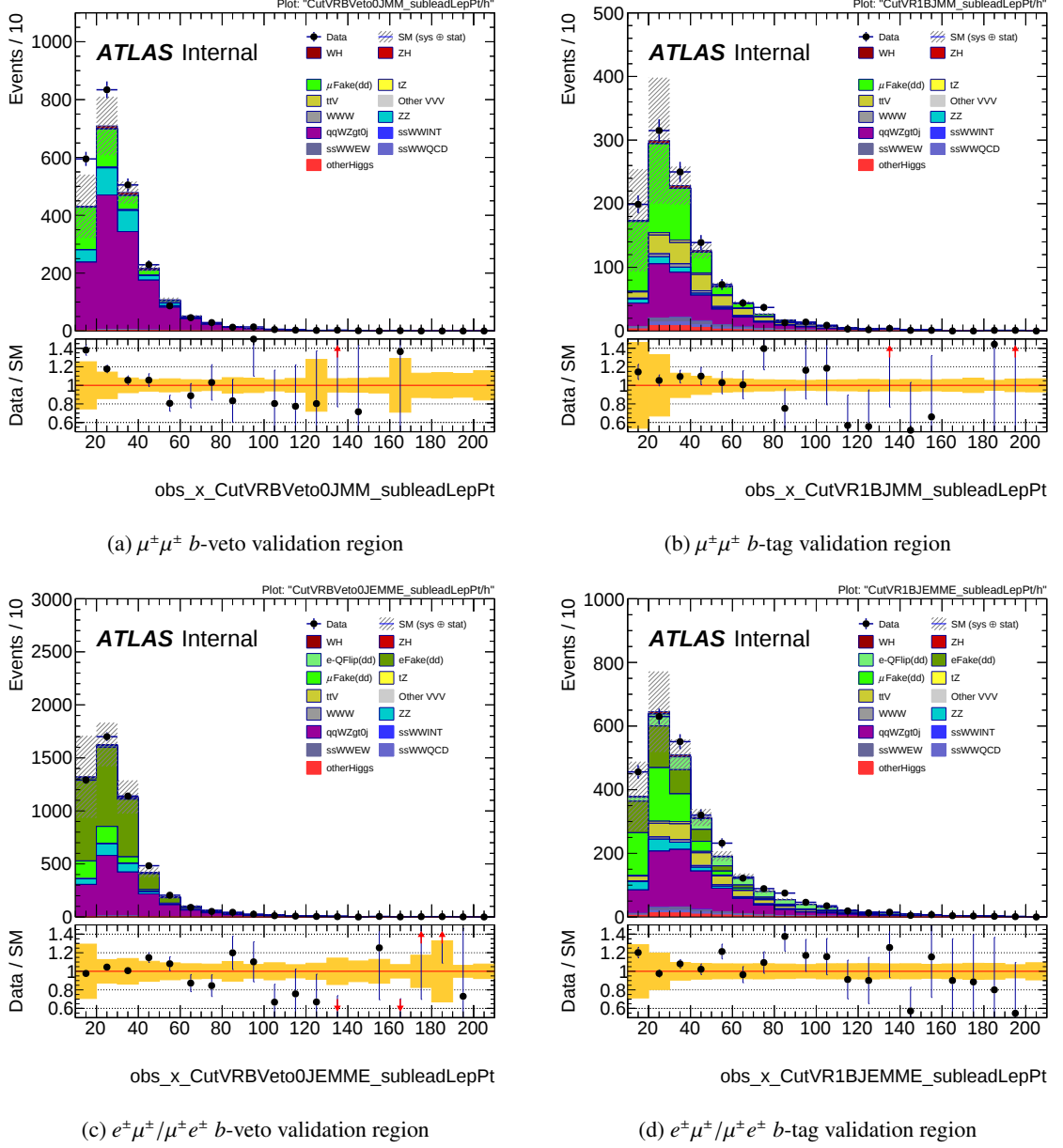


Figure 5.4: Distributions of the subleading lepton p_T in each of the same-sign validation regions used to validate the estimation of the data-driven fake lepton background in the $2\ell/3\ell$ channels. The error bands include the statistical uncertainty as well as the systematic uncertainties relating to the estimation of fake lepton background.

to pass tighter p_T cuts, $p_T^{\ell_0} > 25$ GeV and $p_T^{\ell_1} > 25$ GeV. Finally, an upper cut on the transverse mass of the W boson, $m_T^W < 50$ GeV, is used to reject WZ , the largest electroweak background in the control region. Distributions of $m_{\ell_0\ell_1}$ and m_T^W are shown in Fig. 5.5. The data-MC agreement is imperfect due to the use of Z +jets as the MC prediction for the fake lepton background.

The flavour of the fake lepton is used to further split the ID and anti-ID control regions such that muon and electron fake factors may be independently measured. In the ID control region, the fake lepton is selected as the lepton not assigned to the Z candidate; in the anti-ID control region, the fake lepton is selected as the anti-ID lepton.

The WZ component of the Z +jets control region is normalized to data using a WZ control region, which is constructed by inverting the m_T^W cut of the Z +jets control region. The normalization factor is found to be $1.032 \pm 0.010(\text{stat.})$.

Table 5.8: Event selections for the Z +jets and WZ control regions (CRs) used to estimate the fake factors in 4ℓ channel.

Z+jets CR		WZ CR
ID CR	Anti-ID CR	
Exactly 3 leptons with 1 Z-tagged pair, ℓ_0 and ℓ_1		
70 (80) < $m_{\ell_0\ell_1}$ < 110 GeV if ℓ_0, ℓ_1 are muons (electrons)		
$p_T^{\ell_0} > 25$ GeV && $p_T^{\ell_1} > 25$ GeV		
$m_T^W < 50$ GeV		$m_T^W > 50$ GeV
3 ID leptons	2 ID leptons and 1 anti-ID lepton	

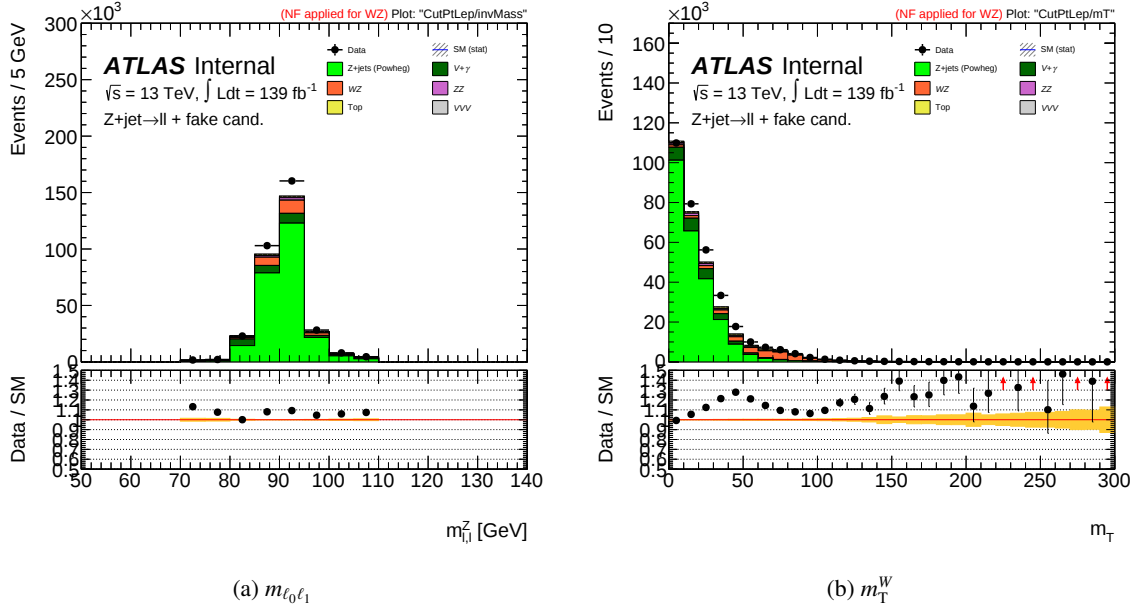


Figure 5.5: Distributions of (a) the invariant mass of the Z candidate and (b) the transverse mass of the W boson following the $p_T^{\ell_0} > 25$ GeV && $p_T^{\ell_1} > 25$ GeV selection from Table 5.8. A normalization factor of 1.032 is applied to the WZ event yield. The error bands indicate the statistical uncertainty. Here, Z +jets corresponds to the MC prediction for the fake lepton background.

The fake factors, shown in Table 5.9, are inclusive in the kinematics of the fake lepton and therefore are

calculated from Eq. 5.1 without a maximum-likelihood fit. Three different sources of uncertainty on the fake factors are quantified: the statistical uncertainty, the electroweak subtraction uncertainty, and the non-closure uncertainty. The uncertainties are calculated separately for each flavour of fake lepton.

The statistical uncertainty, $\Delta\text{SF}_{\text{stat}}$, is a result of the limited statistics in the fake lepton control regions. It is calculated from the definition of the fake factor, Eq. 5.1, by propagating the statistical uncertainties in data and MC.

Attributed to the experimental and theoretical uncertainties on the components of the electroweak background, the electroweak subtraction uncertainty, $\Delta\text{FF}_{\text{EW subtraction}}$, is calculated by varying the total electroweak background yield up and down by 30% and recalculating the fake factor in each case. The larger difference relative to the nominal fake factor is symmetrized and assigned as the uncertainty.

The non-closure uncertainty, $\Delta\text{FF}_{\text{non-closure}}$, is calculated by comparing the MC prediction for WZ , the dominant source of fake leptons in the 4ℓ ZH channel, at the b -jet veto selection in the ID 4ℓ 1SFOS ZH signal region to the estimate obtained by scaling the MC prediction for WZ in the anti-ID signal region by MC-driven fake factors. The relative difference in the fake yield is symmetrized and assigned as the uncertainty.

Table 5.9: Data-driven fake factors (FFs) for the fake lepton estimation in the 4ℓ ZH channel. Also shown are the statistical ($\Delta\text{FF}_{\text{stat}}$), the electroweak subtraction ($\Delta\text{FF}_{\text{EW subtraction}}$), the non-closure ($\Delta\text{FF}_{\text{non-closure}}$), and the total ($\Delta\text{FF}_{\text{total}}$) uncertainties for each flavour of fake factor. The total uncertainty is given by the quadrature sum of $\Delta\text{SF}_{\text{stat}}$, $\Delta\text{FF}_{\text{EW subtraction}}$, and $\Delta\text{SF}_{\text{non-closure}}$.

Fake flavour	FF	$\Delta\text{FF}_{\text{stat}}$ [%]	$\Delta\text{FF}_{\text{EW subtraction}}$ [%]	$\Delta\text{FF}_{\text{non-closure}}$ [%]	$\Delta\text{FF}_{\text{total}}$ [%]
Muons	0.227	1.5	14	49	51
Electrons	0.079	1.2	8	34	35

The fake lepton estimate and its associated uncertainties are of minor importance to the 4ℓ ZH channel. In some bins of the 4ℓ 1SFOS ZH signal region, however, the number of fake events is *negative* due to limited data statistics. In these cases, the number of fake events is set to a small positive number, 1×10^{-12} , to avoid pathological behaviour in the statistical fit. The uncertainty on the number of data events in the anti-ID signal region, and therefore on the number of fake events in the ID signal region, is inflated to account for this. Specifically, the uncertainty on the number of data events in the anti-ID signal region is taken as the square root of the 68% upper confidence bound on the number of data events, assuming Poisson statistics.

5.4 Charge-Flip Electron Estimation for the 2ℓ Channel

A measurement of the electron charge-flip rate is performed using samples of same- and opposite-sign data. It begins with the construction of two dilepton control regions, where at least one of the leptons must be an electron. The control regions have *identical* selections with the *exception* of their charge requirement: one control region requires same-sign leptons and the other requires opposite-sign leptons. In the same-sign control region, the residual between the number of data events, $N_{\text{data}}^{\text{SS,CR}}$, and the number of expected events, $N_{\text{expected}}^{\text{SS,CR}}$, is attributed to charge-flip electrons. The size of this residual is assumed to be proportional to the number of data events in the opposite-sign control region, $N_{\text{data}}^{\text{OS,CR}}$, leading to the definition of the electron charge-flip rate:

$$\epsilon = \frac{N_{\text{data}}^{\text{SS,CR}} - N_{\text{expected}}^{\text{SS,CR}}}{N_{\text{data}}^{\text{OS,CR}}} . \quad (5.11)$$

The electron charge-flip rate is subsequently applied to a same-sign dilepton signal region, where at least one of the leptons must be an electron. An opposite-sign dilepton signal region, identical to the same-sign signal region with the exception of its charge requirement, is constructed. The number of data events in the opposite-sign signal region, $N_{\text{data}}^{\text{OS,SR}}$, scaled by the electron charge-flip rate provides the data-driven estimate of the number of charge-flip electron events in the same-sign signal region:

$$N_{Q\text{-flip}}^{\text{SS,SR}} = \begin{cases} \epsilon \times N_{\text{data}}^{\text{OS,SR}} & \text{if } e\mu/\mu e \\ (\epsilon + \epsilon' - 2\epsilon\epsilon') \times N_{\text{data}}^{\text{OS,SR}} & \text{if } ee \end{cases}. \quad (5.12)$$

For same-flavour signal regions, either electron could have misidentified charge. Accordingly, the charge-flip rate is modified as $(\epsilon + \epsilon' - 2\epsilon\epsilon')$ to account for this, where ϵ' indicates that the charge-flip rates could differ for each electron and $-2\epsilon\epsilon'$ constitutes a double-counting correction.

For the 2ℓ same-sign WH channel, the primary sources of charge-flip electrons are Z +jets and $t\bar{t}$. A Z +jets control region, whose selections are shown in Table 5.10, is constructed for the estimation of the charge-flip rates. Two same-sign electrons are required, and the leading and subleading electrons are required to pass minimum p_T cuts of 22 and 15 GeV, respectively. To reject $t\bar{t}$, zero b -jets are required; to select Z +jets, the dilepton invariant mass is required to be within 20 GeV of the mass of the Z boson. The dilepton invariant mass cut also ensures the region is orthogonal to the ee same-sign WH signal region.

Table 5.10: Event selections for the Z +jets control region used to estimate the electron charge-flip rates in the 2ℓ same-sign WH channel.

Z +jets control region
Exactly 2 same-sign (or opposite-sign) ID electrons
At least 1 ID electron matched to a single lepton trigger
$p_T^{\ell_{\text{lead}}} > 22 \text{ GeV} \ \&\& \ p_T^{\ell_{\text{sublead}}} > 15 \text{ GeV}$
$N_{b\text{-jets}} = 0$
$ m_{\ell\ell} - m_Z < 20 \text{ GeV}$

The charge-flip rates are measured in bins of electron p_T and $|\eta|$. There are 4 bins in p_T between 15 and 1000 GeV as well as 6 bins in $|\eta|$ between 0.6 and 2.47, excluding the transition region $1.37 < |\eta| < 1.52$. In the following, ϵ_i denotes the rate in the i -th cell of $p_T - |\eta|$ space.

The rates are extracted by fitting a likelihood function parameterized in terms of the rates, $\vec{\epsilon}$, and the nuisance parameters on the data-driven fake lepton and electroweak backgrounds, $\vec{\theta}$. The likelihood function is written as a product of Poisson PDFs for each $p_T - |\eta|$ bin of the control region per electron and Gaussian constraints for the systematics on the backgrounds:

$$\mathcal{L}(\vec{\epsilon}, \vec{\theta}) = \left[\prod_{i,j}^{p_T - |\eta| \text{ bins}} P(N_{ij}^{\text{data,SS}} | N_{ij}^{\text{expected,SS}}(\epsilon_i, \epsilon_j, \vec{\theta})) \right] \times \left[\prod_k^{\text{uncertainties}} G(\tilde{\theta}_k | \theta_k) \right], \quad (5.13)$$

with:

$$N_{ij}^{\text{expected,SS}}(\epsilon_i, \epsilon_j, \vec{\theta}) = (\epsilon_i + \epsilon_j - 2\epsilon_i\epsilon_j) \times N_{ij}^{\text{data,OS}} + N_{ij}^{\text{EW,SS}}(\vec{\theta}) + N_{ij}^{\text{fakes,SS}}(\vec{\theta}). \quad (5.14)$$

An event yield denoted by N_{ij} corresponds to the yield in the sub-region of the control region where the leading and subleading electrons are in the i -th and j -th cells of $p_T - |\eta|$ space, respectively. Additionally:

- P is a Poisson PDF and G is a Gaussian constraint;

- $N_{ij}^{\text{expected,SS}}(\epsilon_i, \epsilon_j, \vec{\theta})$ is the expected same-sign yield;
- $N_{ij}^{\text{data,SS}}$ and $N_{ij}^{\text{data,OS}}$ are the same- and opposite-sign data yields, respectively;
- $N_i^{\text{EW,SS}}(\vec{\theta})$ is the same-sign MC electroweak background yield;
- $N_i^{\text{fakes,SS}}(\vec{\theta})$ is the same-sign data-driven fake lepton background yield;

The Gaussian constraints allow the yields for the WZ component of the electroweak background and the fake lepton background to float in the fit. For WZ , the size of the fluctuations is dictated by the cross section uncertainty of $\Delta\sigma_{WZ} = 13\%$ [465]. For the fake lepton background, the size of the fluctuations is dictated by the systematic uncertainties on the fake factors and scale factors, which are described in Section 5.2.

The fitted charge-flip rates are shown in Table 5.11 and their corresponding uncertainties are shown in Fig. 5.6. The rate is an increasing function of electron p_T and $|\eta|$ – this is expected because a harder electron has less curvature and because pattern recognition becomes more complicated in the forward regions of the detector. The rates themselves are very low but result in a sizable background in the 2ℓ same-sign WH channel. Three sources of uncertainty are quantified: the statistical, the background estimation, and the mass window uncertainties.

The statistical uncertainty is a result of the limited statistics in the Z +jets control region. It is given by the 68% confidence interval on best-fit values of the charge-flip rates with uncertainties on the WZ cross section and the fake lepton background fixed. In the barrel region, the uncertainty can be as large as 22%.

The background estimation uncertainty is a result of the uncertainties on the WZ cross section and the fake lepton background. It is given by the 68% confidence interval on best-fit values of the charge-flip rates with the statistical uncertainties fixed.

The mass window uncertainty is attributed to the particular choice of mass window applied to the Z candidate, 20 GeV. It is evaluated by varying the choice of mass window and refitting the charge-flip rates. The relative difference between the nominal and varied rates is assigned as the uncertainty.

Table 5.11: Data-driven charge-flip rates (in %) for the charge-flip electron estimation in the 2ℓ same-sign WH channel. The uncertainties correspond to the quadrature sum of the statistical, the background estimation, and the mass window uncertainties.

		p_T bin			
		$15 < p_T < 60 \text{ GeV}$	$60 < p_T < 90 \text{ GeV}$	$90 < p_T < 130 \text{ GeV}$	$130 \text{ GeV} < p_T$
$ \eta $ bin	$0 < \eta < 0.6$	0.0013 ± 0.0003	0.0048 ± 0.0024	0.0278 ± 0.0070	0.1147 ± 0.0220
	$0.6 < \eta < 1.1$	0.0038 ± 0.0006	0.0228 ± 0.0030	0.0644 ± 0.0119	0.2384 ± 0.0365
	$1.1 < \eta < 1.37$	0.0087 ± 0.0010	0.0534 ± 0.0070	0.1867 ± 0.0240	0.5473 ± 0.0640
	$1.52 < \eta < 1.7$	0.0183 ± 0.0019	0.0793 ± 0.0118	0.2664 ± 0.0395	0.9444 ± 0.1157
	$1.7 < \eta < 2.3$	0.0381 ± 0.0020	0.1291 ± 0.0094	0.4584 ± 0.0332	1.5050 ± 0.1245
	$2.3 < \eta $	0.0816 ± 0.0052	0.4586 ± 0.0470	1.4437 ± 0.1444	4.0149 ± 0.4790

5.4.1 Validation

The data-driven charge-flip rates are applied to a Z +jets validation region in order to assess the data-MC fidelity of the charge-flip electron background estimate. The validation region, whose selections are shown Table 5.6, is a subset of the Z +jets control region used to measure the rates with the additional requirement of one or more jets.

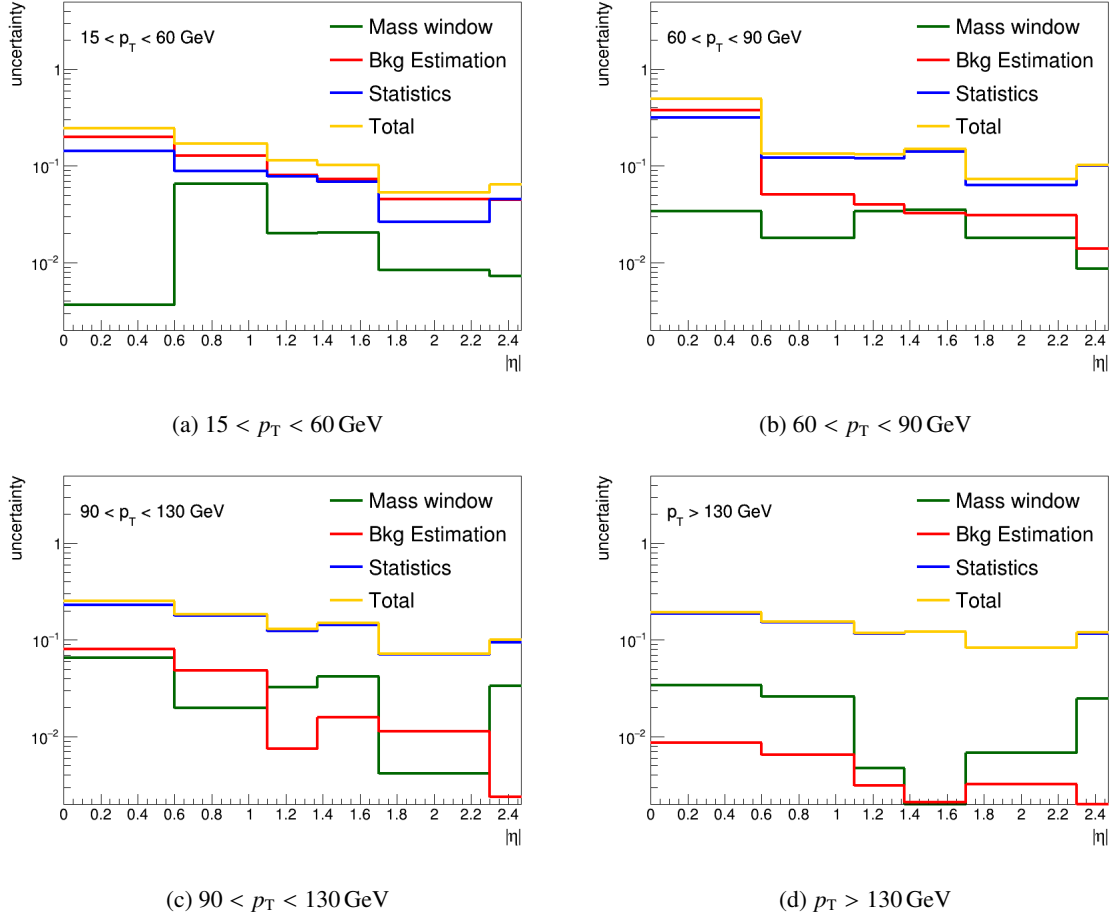


Figure 5.6: Relative uncertainties on the data-driven charge-flip rates for the electron charge-flip estimation in the 2ℓ same-sign WH channel. For each region of electron $p_T - |\eta|$, the contributions to the total uncertainty as well as the total uncertainty, calculated as the quadrature sum of the individual contributions, are shown.

Table 5.12: Event selections for the Z+jets validation region used to validate the charge-flip electron background estimate in the 2ℓ same-sign WH channel.

Z+jets validation region
Exactly 2 same-sign ID electrons
At least 1 ID electron matched to a single lepton trigger
$p_T^{\ell_{\text{lead}}} > 22 \text{ GeV} \ \&\& \ p_T^{\ell_{\text{sublead}}} > 15 \text{ GeV}$
$N_{b\text{-jets}} = 0$
$ m_{\ell\ell} - m_Z < 20 \text{ GeV}$
$N_{\text{jets}} \geq 1$

Distributions of a few interesting variables in the Z+jets validation region are shown in Fig. 5.7. There is shape mismodelling in the lepton p_T and dilepton invariant mass distributions; however, the effect of this mismodelling is expected to be limited for two reasons. Firstly, all other variables including the recurrent neural network output – the signal discriminant for the 2ℓ same-sign WH channel (Section 6.2) – are well modelled. Secondly, the 2ℓ same-sign WH signal region vetoes the Z boson mass window targeted by the validation region.

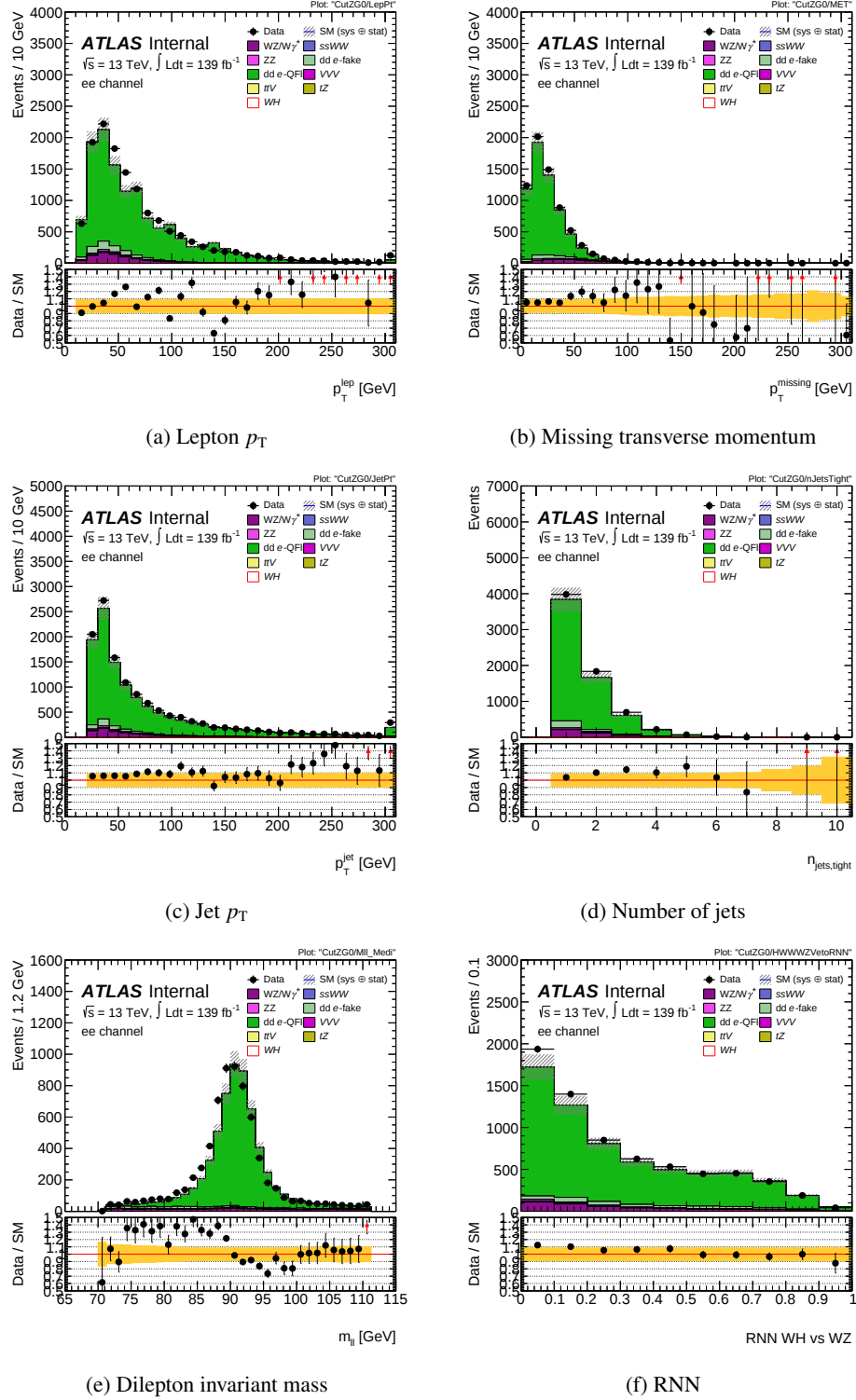


Figure 5.7: Distributions of the interesting variables in the Z+jets validation region used to validate the estimation of the data-driven charge-flip electron background in the 2ℓ same-sign WH channel. The error bands include the statistical uncertainty as well as the systematic uncertainties relating to the estimation of charge-flip electron and fake lepton backgrounds. The recurrent neural network (RNN) is described in Section 6.2.

Chapter 6

Measurement Channels

As introduced in Section 1.3, the $VH, H \rightarrow WW^*$ measurement is performed in several channels across 2-, 3-, and 4-lepton final states, which are subdivided further according to the charge and flavour configuration of the leptons. The high-level event selections for each channel are shown in Table 6.1. Each channel uses a combination of kinematic selections and multivariate discriminants to construct its signal and control regions. The multivariate discriminants targeting the $VH, H \rightarrow WW^*$ process are optimally binned within each signal region such that their expected sensitivity is maximized, and the optimally-binned signal regions as well as the control regions are provided as input to the combined fit. Both the optimization of the binning and the combined fit are described in Chapter 8.

The remainder of this chapter describes the construction of the signal, control, and validation regions used by the 2ℓ DFOS VH (Section 6.1), 2ℓ same-sign WH (Section 6.2), 3ℓ Z-dominated WH (Section 6.3), 3ℓ Z-depleted WH (Section 6.4), and 4ℓ ZH (Section 6.5) channels.

Table 6.1: High-level event selections for each channel of the $VH, H \rightarrow WW^*$ measurement. The dilepton angular separation is denoted by $\Delta R_{\ell\ell}$, and the minimum same-flavour, opposite-sign (SFOS) dilepton invariant mass and minimum different-flavour, opposite-sign (DFOS) dilepton invariant mass are denoted by $m_{\ell\ell,\text{SFOS}}^{\min}$ and $m_{\ell\ell,\text{DFOS}}^{\min}$, respectively. In the 3ℓ channels, the $\Delta R_{\ell\ell}$ selection is applied to all combinations of leptons; in the 4ℓ channels, it is applied only to the leptons originating from the $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$ decay.

Selection	Channel					
	2ℓ		3ℓ		4ℓ	
	DFOS	Same-sign	Z-depleted	Z-dominated	1SFOS	2SFOS
Lepton p_T [GeV]	>15	>15	>15	>15	>10	>10
Number of leptons	2	2	3	3	4	4
Total lepton charge	0	± 2	± 1	± 1	0	0
Number of SFOS lepton pairs	0	0	0	≥ 1	1	2
$\Delta R_{\ell\ell}$	>0.1	>0.4	>0.1	>0.1	>0.2	>0.2
$m_{\ell\ell,\text{SFOS}}^{\min}$ [GeV]	–	–	–	>12	>12	>12
$m_{\ell\ell,\text{DFOS}}^{\min}$ [GeV]	>10	–	–	–	>10	>10
Number of jets	≥ 2	≥ 1	–	–	–	–
Number of b -jets	0	0	0	0	0	0

Common Observables: A common set of $H \rightarrow WW^*$ decay observables is exploited among the channels. The $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$ topology is characterized by a low azimuthal separation, $\Delta\phi_{\ell\ell}$, and a low angular

separation, $\Delta R_{\ell\ell}$, between the leptons. This is because the spin-0 initial state and the vector-axial structure of the weak interaction tends to result in collimated products. The dilepton invariant mass, $m_{\ell\ell}$, is expected to be low for the same reason.

The mass of the Higgs boson cannot be perfectly reconstructed due to the transverse momentum carried away by the neutrinos. Instead, the transverse mass of the dilepton system, m_T , is used as a proxy:

$$m_T = \sqrt{(E_T^{\ell\ell} + E_T^{\text{miss}})^2 - |\vec{p}_T^{\ell\ell} + \vec{E}_T^{\text{miss}}|^2} \quad \text{where} \quad E_T^{\ell\ell} = \sqrt{|\vec{p}_T^{\ell\ell}|^2 + m_{\ell\ell}^2}, \quad (6.1)$$

where $E_T^{\ell\ell}$ and $\vec{p}_T^{\ell\ell}$ are the transverse energy and the vectorial transverse momentum of the dilepton system, respectively. For $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$ decays, m_T is expected to peak between 50 GeV and $m_H \approx 125$ GeV. For $H \rightarrow WW^* \rightarrow \ell\nu q\bar{q}$ decays, a similar quantity may be constructed using the lepton-jet invariant mass, $m_{\ell j}$.

Background rejection is similarly achieved through a common set of observables. Low-mass resonances – for example, $J/\psi \rightarrow e\bar{e}/\mu\bar{\mu}$ – are rejected by requiring the minimum dilepton invariant mass to exceed 10 or 12 GeV, depending on the flavour configuration of the leptons. A selection on the dilepton angular separation, $\Delta R_{\ell\ell} > 0.1$, is applied across channels to remove events with overlapping leptons, a common feature of fake lepton events. Rejection of top quark processes is achieved using a b -jet veto, $N_{b\text{-jets}} = 0$.

The $Z \rightarrow e\bar{e}/\mu\bar{\mu}$ background is rejected by vetoing events where the dilepton invariant mass is consistent with the mass of the Z boson or where the missing transverse momentum, E_T^{miss} , is low, as there is no “true” E_T^{miss} expected. Similarly, the $Z \rightarrow \tau\bar{\tau}$ background with leptonically-decaying tau leptons is rejected by vetoing events where the invariant mass of the tau lepton pair, $m_{\tau\tau}$, is consistent with the mass of Z boson. The individual neutrinos are expected to be collimated with the tau leptons as per the collinear approximation [469]. This allows for the computation of the energy fractions of the neutrinos, the only source of E_T^{miss} in the event, and therefore $m_{\tau\tau}$.

6.1 2ℓ Different-Flavour, Opposite-Sign VH Channel

The 2ℓ different-flavour, opposite-sign (DFOS) VH channel is characterized the presence of two DFOS leptons from the $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$ decay and two or more jets from the hadronic decay of the associated vector boson. In this channel, the leading and subleading leptons in p_T are denoted by ℓ_0 and ℓ_1 , respectively, and the leading and subleading jets in p_T are denoted by j_0 and j_1 , respectively.

The leading and subleading leptons are required to exceed transverse momentum thresholds of 22 and 15 GeV, respectively. A selection on the dilepton angular separation, $\Delta R_{\ell\ell} > 0.1$, is used to remove overlapping leptons, and a selection on the dilepton invariant mass, $m_{\ell\ell} > 10$ GeV, is used to remove low-mass resonances. A b -jet veto is applied to reject $t\bar{t}/Wt$ production. The hadronic decay of the associated vector boson is characterized by a dijet mass consistent with that of a W/Z boson and relatively collinear jets. To ensure the selected events are orthogonal to the ≥ 2 -jets phase space targeted by the ggF+VBF, $H \rightarrow WW^*$ measurement [13], selections are applied to the dijet invariant mass, $|m_{jj} - 85| < 15$ GeV, and dijet rapidity separation, $\Delta Y_{jj} < 1.2$.

The 2ℓ DFOS VH signal region is defined using a multivariate discriminant known as an artificial neural network (ANN). Following the dijet selections, a cut on the ANN discriminant targeting the VH , $H \rightarrow WW^*$ process, $\text{ANN}_{VH} \geq 0.2$, is used to remove the majority of background events and improve the signal-over-background ratio. This final cut constitutes the signal region. Within the signal region, ANN_{VH} is optimally binned in its partial range, $[0.2, 1]$, and provided to the combined fit.

The selections for each region of the 2ℓ DFOS VH channel are shown in Table 6.2. Control regions are used to normalize the top, Z +jets, and $q\bar{q} \rightarrow WW$ backgrounds to data. In the following sections, $q\bar{q} \rightarrow WW$ is referred to as simply WW . The fake lepton background is estimated using the data-driven approach described in Chapter 5. The shape and normalization of all other backgrounds – which include non- WW diboson, triboson, and ggF production – are taken from simulation. A top validation region based on a b -jet requirement is used to assess the modelling of the top quark background, the largest background in the 2ℓ DFOS VH channel.

Table 6.2: Event selections for the signal (SR), control (CR), and validation (VR) regions of the 2ℓ DFOS VH channel. The artificial neural network (ANN) is described in Section 6.1.1.

2ℓ DFOS VH SR	Top VR	Top CR	Z +jets CR	WW CR
Exactly 2 different-flavour, opposite-sign ID leptons				
At least 1 ID lepton matched to a single lepton trigger				
$p_T^{\ell_{\text{lead}}} > 22 \text{ GeV} \ \&\& \ p_T^{\ell_{\text{sublead}}} > 15 \text{ GeV}$				
$m_{\ell\ell} > 10 \text{ GeV}$				
$\Delta R_{\ell\ell} > 0.1$				
$N_{\text{jets}} \geq 2$				
$N_{b\text{-jets}} = 0$	$N_{b\text{-jets}} = 1$	$N_{b\text{-jets}} = 0$		
$ m_{jj} - 85 < 15 \text{ GeV} \ \&\& \ \Delta Y_{jj} < 1.2$		–		
$\text{ANN}_{VH} \geq 0.2$		$\text{ANN}_{VH} < 0.2$		
		$\text{ANN}_{\text{top}} \geq 0.1$	$\text{ANN}_{\text{top}} < 0.1$	$\text{ANN}_{\text{top}} < 0.5$
		$\text{ANN}_{Z+\text{jets}} < 0.6$	$\text{ANN}_{Z+\text{jets}} \geq 0.3$	$\text{ANN}_{Z+\text{jets}} < 0.1$
		$\text{ANN}_{WW} < 0.7$	$\text{ANN}_{WW} < 0.2$	$\text{ANN}_{WW} \geq 0.7$

6.1.1 Multivariate Analysis

The multivariate discriminant used by the 2ℓ DFOS VH channel is an artificial neural network (ANN), which assigns input events into two or more output classes corresponding to different physics processes. The ANN achieves this through a series of linear transformations or “layers” applied sequentially to an input event, which is represented as a vector, with a nonlinear “activation” function applied element-wise to the result of each linear transformation.

Given a particular truth class for an input event, the goal of the ANN is to correctly reproduce the truth class at its output. The predicted class’s output “score” is constrained between 0 and 1 and may be interpreted as the “probability” of the input event belonging to the truth class. A *binary* classifier provides discrimination between two classes using a single output score. In contrast, a *multi*-classifier provides discrimination between three or more classes. There is one output score per class, and the output score for a given class corresponds to the relative probability of the event belonging that class. Per event, the output scores sum to unity.

The free parameters of the linear transformation for each layer of the ANN constitute its total set of free parameters. These free parameters are optimally chosen by minimizing a “loss” function quantifying how disparate the truth and the predicted classes are. This minimization is referred to as “training” the ANN and typically proceeds via backpropagation [470], a type of stochastic gradient descent, using a dedicated set of training events.

The 2ℓ DFOS VH channel utilizes a multiclassifier ANN for discriminating between VH , ggF, VBF, $t\bar{t}/Wt$ (denoted “top”), Z +jets, and WW processes. The use of a multiclassifier provides greater freedom for categorizing events and simultaneously provides a way for constructing control regions enriched in the top,

Z +jets, and WW backgrounds. For each event, the multiclassifier receives basic kinematic information relating to the leading and subleading leptons, the leading and subleading jets, and the E_T^{miss} as well as higher-level derived quantities. The input variables are listed in Table 6.3, and the inputs themselves are produced at the level of the b -jet veto in the 2ℓ DFOS VH signal region, given in Table 6.2. Input variables with varying shapes for one or more classes are generally chosen. The 2ℓ DFOS VH and 3ℓ WH channels additionally train their respective multivariate discriminants many times with different input variables, selecting the set which maximizes the separation of the signal from the background. Prior to training, each variable has its median subtracted and is scaled to its interquartile range using the Python package `scikit-learn` [471]. This has the effect of roughly centring each variable on 0 with a range that extends to $O(\pm 1)$. Shape comparisons among classes for each variable are shown in Figs. 6.1 and 6.2.

The multiclassifier is trained using Keras [472] with the TensorFlow [473] backend. As the simulated samples used to train the multiclassifier are the *same* samples that the multiclassifier will be deployed on, a 5-way k -folding prescription is used to avoid possible bias. The network is trained 5 times, $i = 0, \dots, 5$, where i -th network is trained using events which satisfy $(\text{event number}) \bmod 5 \neq i$. The i -th network is said to have been trained using k -fold i , and it is subsequently deployed on events which satisfy $(\text{event number}) \bmod 5 = i$. In this way, a network is never deployed on the same events it is trained with.

Table 6.3: Input variables to the multiclassifier used by the 2ℓ DFOS VH channel. Variables are grouped based on their purpose; however, some variables may serve more than one purpose.

Purpose	Variable(s)	Description
$H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$ discrimination	$p_T^{\ell_0}, p_T^{\ell_1}$	Leading and subleading lepton transverse momentum
	$m_{\ell\ell}$	Dilepton invariant mass
	$\Delta\phi_{\ell\ell}$	Dilepton azimuthal separation
	$\Delta Y_{\ell\ell}$	Dilepton rapidity separation
	m_T	Transverse mass (Eq. 6.1)
$V(\rightarrow q\bar{q})H$ / VBF discrimination	$p_T^{j_0}, p_T^{j_1}$	Leading and subleading jet transverse momentum
	m_{jj}	Dijet invariant mass
	$\Delta\phi_{jj}$	Dijet azimuthal separation
	ΔY_{jj}	Dijet rapidity separation
	$m_{\ell_0 j_0}, m_{\ell_0 j_1}, m_{\ell_1 j_0}, m_{\ell_1 j_1}$	All lepton-jet invariant mass combinations
Z +jets separation	$m_{\tau\tau}$	Invariant mass of the tau lepton pair
	E_T^{miss}	Missing transverse momentum
	$\mathcal{S}_{\text{miss}}$	Object-based E_T^{miss} significance
Top / WW separation	H_T	Transverse momentum sum of all hard objects (see Eq. 4.2)

Overtraining occurs when the number of free parameters of the network is comparable to the number of events used during training. In this situation, it is possible for the network to perfectly reproduce the truth class for each training event while simultaneously failing to generalize to new events. To assess whether the multiclassifier is overtrained, it is deployed on a dedicated set of testing events which are independent of the training events but share the same set of truth classes. The multiclassifier is expected to achieve similar performance in both the training and testing datasets. Shape comparisons among the training and testing datasets for each output score of the network trained using k -fold 0 are shown in Fig. 6.3. Shape comparisons among the training and testing datasets for the networks trained using k -folds 1 through 4 are shown in Figs. B.1 through B.4 of Appendix B. Good agreement is observed between the shapes, and so it may be concluded that no overtraining is present.

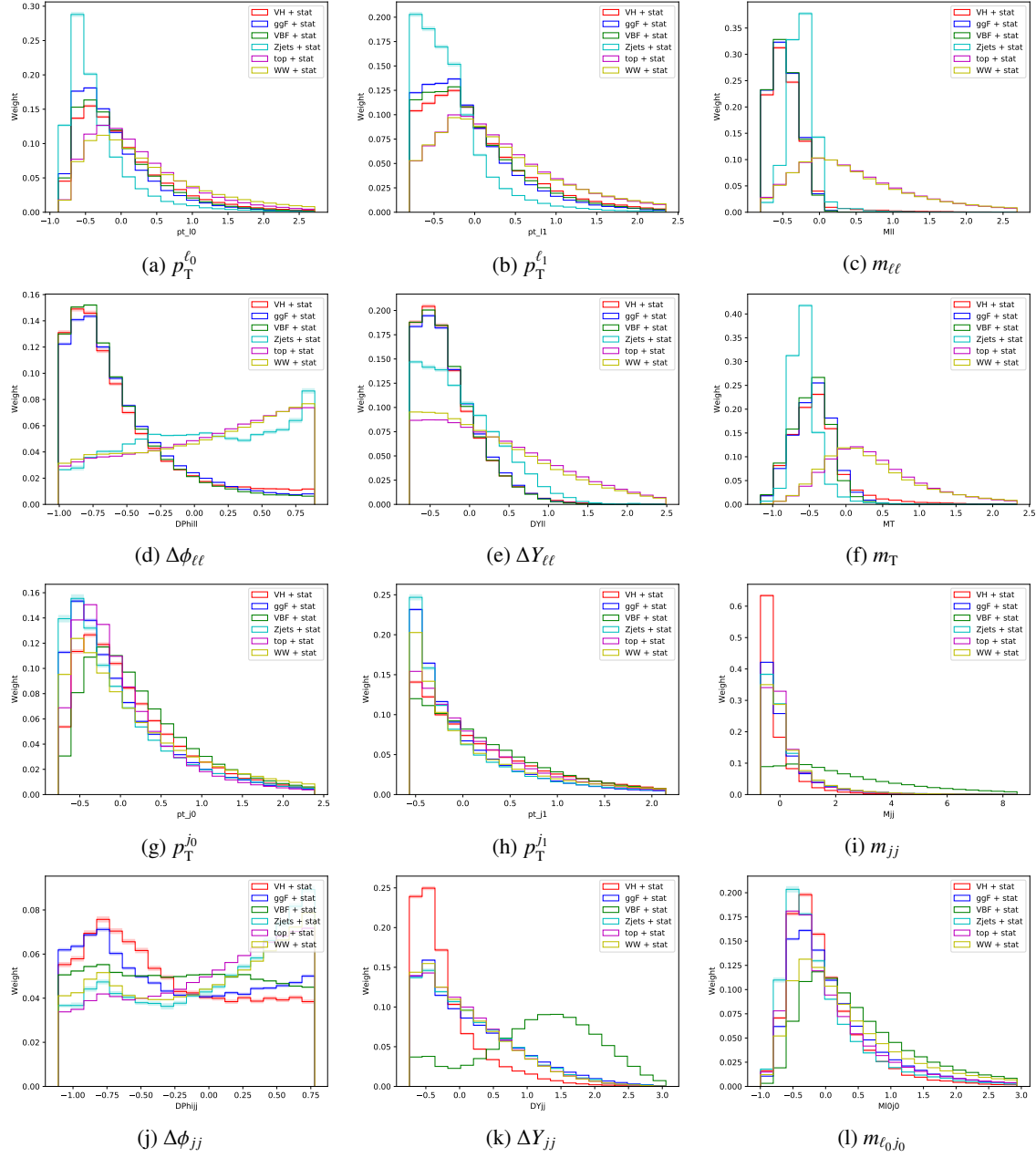


Figure 6.1: Shape comparisons among classes for each input variable of the multiclassifier used by the 2ℓ DFOS VH channel. The sum-of-weights for each class is normalized to unity. Each variable has its median subtracted and is scaled to its interquartile range. The error bands indicate the statistical uncertainty.

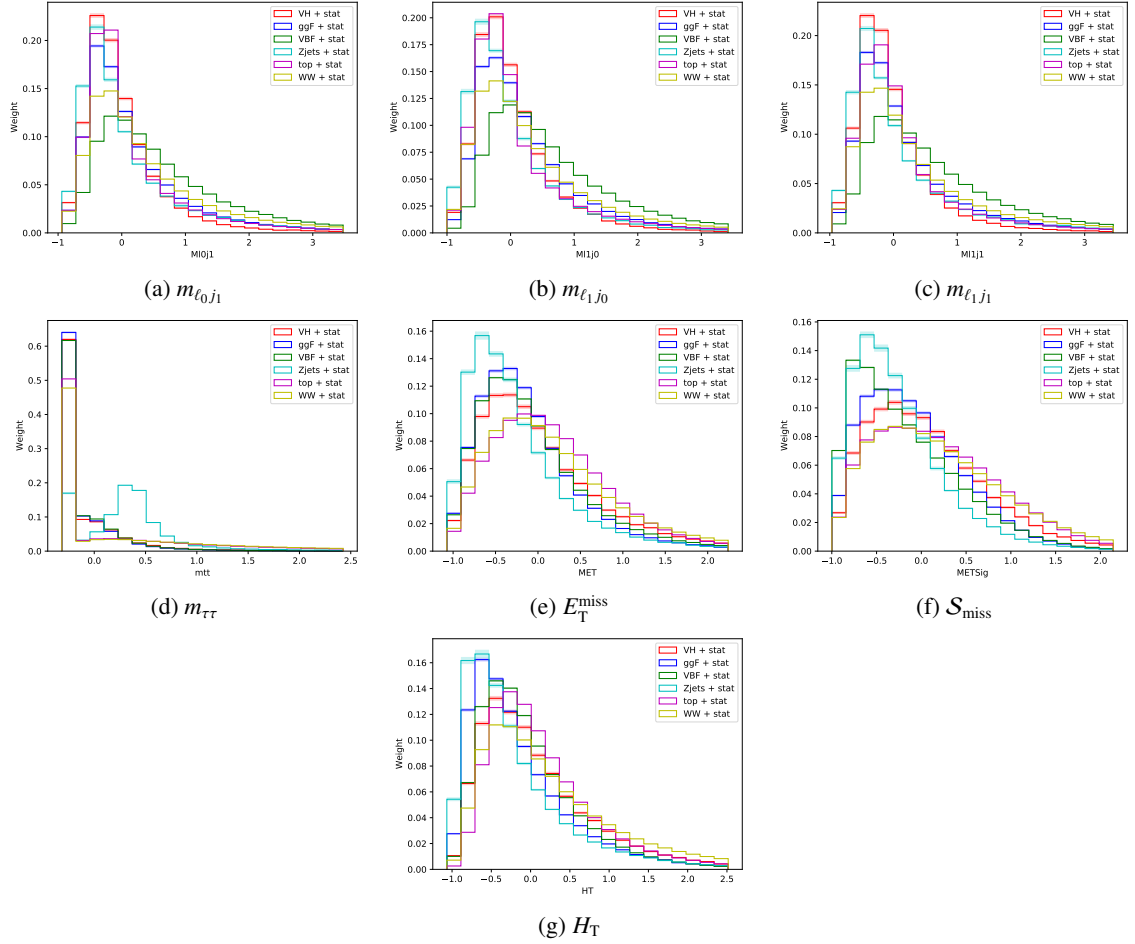


Figure 6.2: Shape comparisons among classes for each input variable of the multiclassifier used by the 2ℓ DFOS VH channel. The sum-of-weights for each class is normalized to unity. Each variable has its median subtracted and is scaled to its interquartile range. The error bands indicate the statistical uncertainty.

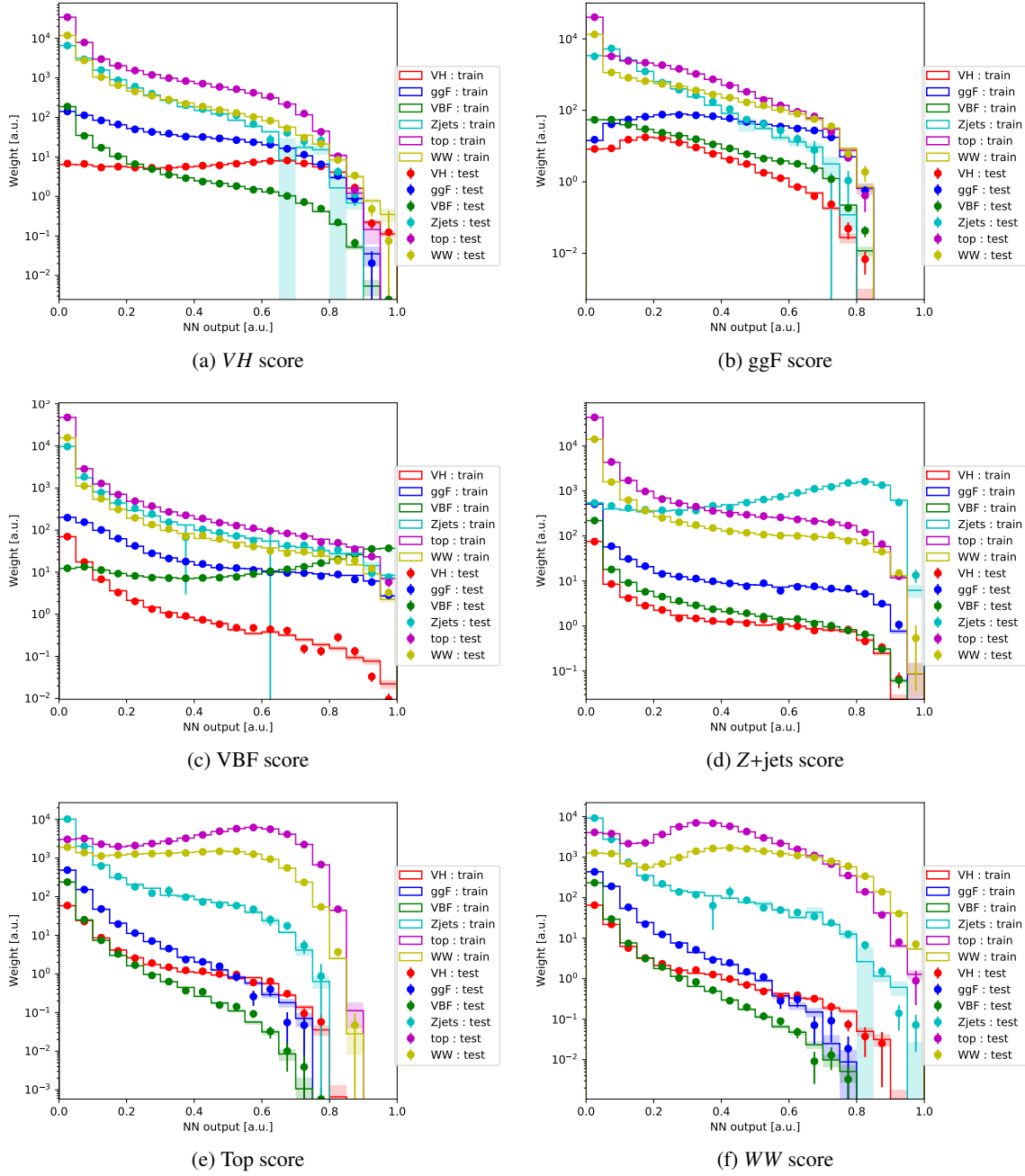


Figure 6.3: Output scores for the multiclassifier used by the 2ℓ DFOS VH channel. The multiclassifier is trained using events satisfying (event number) $\bmod 5 \neq 0$ (i.e., “ k -fold 0”), and it is deployed on both the training and testing datasets to assess the level of agreement. The error bands/bars indicate the statistical uncertainty.

The distribution of the multiclassifier's VH output score in the 2ℓ DFOS VH signal region is shown in Fig. 6.4. Good agreement is observed between data and simulation. In the rightmost bin, there is an excess in data: $18.4 \pm 2.2(\text{stat.})$ events were expected whereas 25 events were measured.

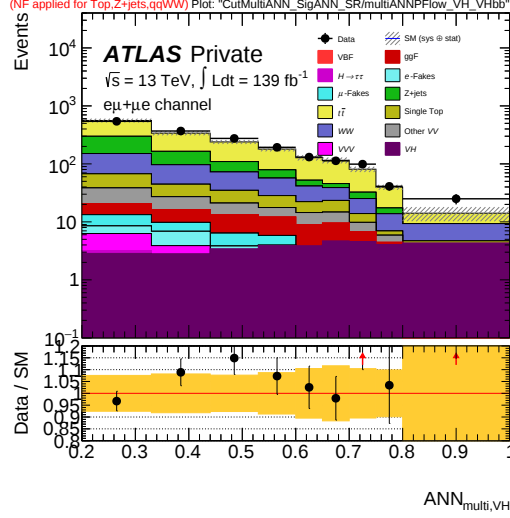


Figure 6.4: Distribution of the multiclassifier's VH output score in the signal region of the 2ℓ DFOS VH channel. Normalization factors are applied to the top, Z +jets, and WW event yields. The error band includes the statistical uncertainty as well as the dominant experimental systematic uncertainties.

6.1.2 Control and Validation Regions

The region orthogonal to the signal region, $\text{ANN}_{VH} < 0.2$, is dominated by background processes and contributes little sensitivity to the measurement. Instead, the top, Z +jets, and WW control regions are defined within this orthogonal region using cuts on the multiclassifier's top (ANN_{top}), Z +jets ($\text{ANN}_{Z+\text{jets}}$), and WW (ANN_{WW}) output scores. A 3D grid search is performed to select the cuts, constrained between 0 and 1, which minimize the statistical uncertainties on the associated normalization factors while maintaining orthogonality between the control regions. The results of the grid search are shown in Table 6.4. Due to similarities in the shapes of the input variables, it is difficult to separate top and WW production and therefore to produce highly pure regions in each. Normalization factors of $0.98 \pm 0.01(\text{stat.})$, $0.89 \pm 0.01(\text{stat.})$, and $0.89 \pm 0.04(\text{stat.})$ are obtained for the top, Z +jets, and WW processes, respectively. Distributions of the dilepton invariant mass and azimuthal separation in the control regions are shown in Figs. 6.5a through 6.5c. With the application of the normalization factors, good agreement between data and simulation is observed.

Table 6.4: The definitions of the top, Z +jets, and WW control regions used by the 2ℓ DFOS VH channel. The number of signal events, N_{signal} , corresponds to the event yield for the process targeted by a given control region; the number of background events, $N_{\text{background}}$, corresponds to the event yield for all other processes. Purity is given by $N_{\text{signal}}/(N_{\text{signal}} + N_{\text{background}})$.

Control region	ANN_{top}	$\text{ANN}_{Z+\text{jets}}$	ANN_{WW}	N_{signal}	$N_{\text{background}}$	Purity [%]
Top	≥ 0.1	< 0.6	< 0.7	40,141	15,699	72
Z +jets	< 0.1	≥ 0.3	< 0.2	9,380	2,862	77
WW	< 0.5	< 0.1	≥ 0.7	1,758	1,269	58

Top quark processes comprise the dominant background of the 2ℓ DFOS VH channel; accordingly, it is important to assess the quality of the top quark modelling. To do so, a top validation region is constructed by inverting the b -jet veto of the signal region and requiring one b -jet, generating a high-statistics region with a top purity of 97%. The distribution of the multiclassifier's VH output score in the top validation region is shown in Fig. 6.5d. Simulation disagrees with data by an overall normalization factor, a feature which is caused by applying a top normalization factor obtained using a b -veto region to a b -tag region. Discrepancies between data and simulation of 5–10% are expected to be covered by top theory uncertainties and b -tagging efficiency uncertainties. It may be concluded that top quark processes are adequately modelled.

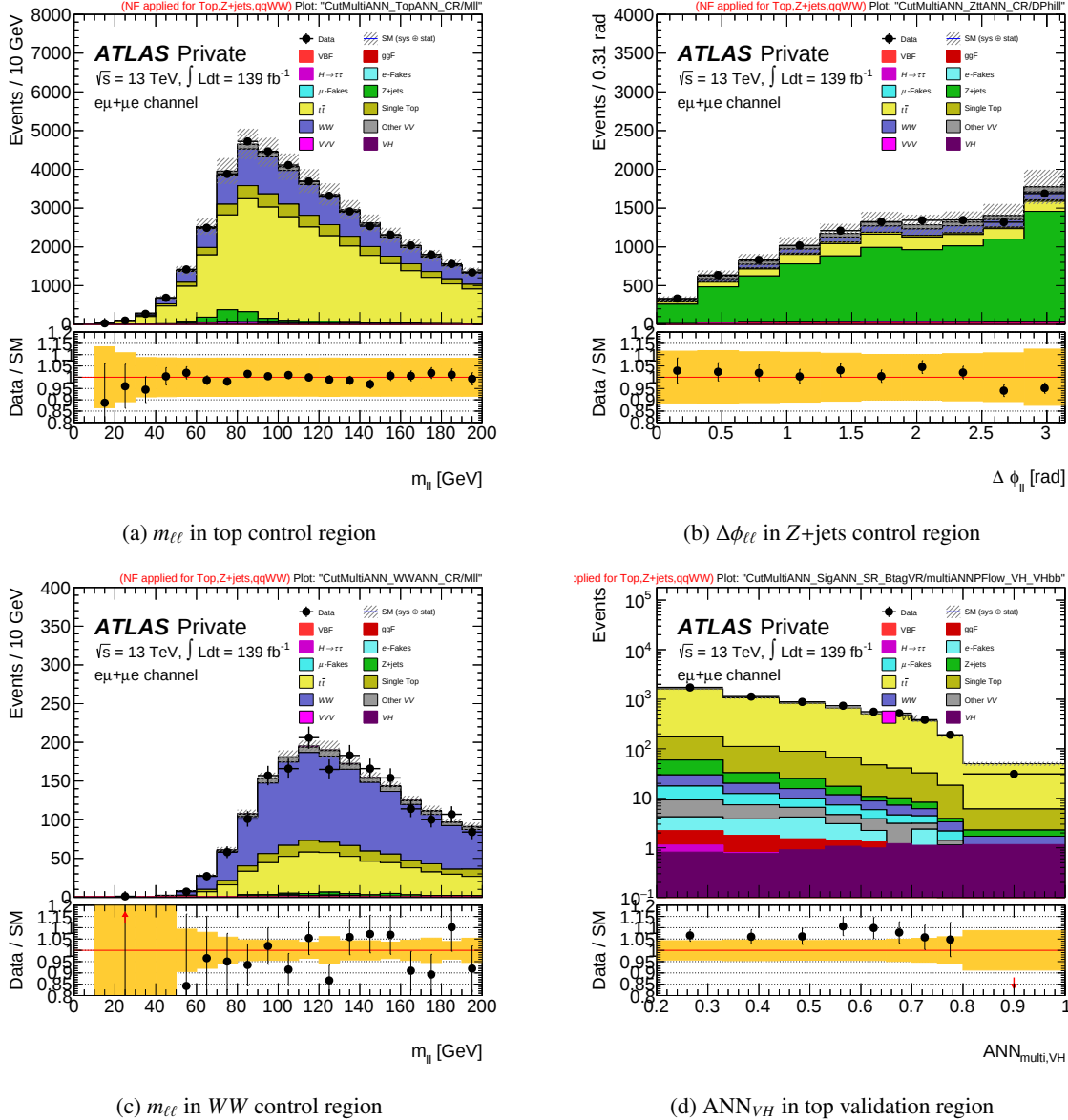


Figure 6.5: Distributions of (a) the dilepton invariant mass in the top control region, (b) the dilepton azimuthal separation in the Z+jets control region, (c) the dilepton invariant mass in the WW control region, and (d) the multiclassifier's VH output score in the top validation region of the 2ℓ DFOS VH channel. Normalization factors are applied to the top, Z+jets, and WW event yields. The error bands include the statistical uncertainty as well as the dominant experimental systematic uncertainties.

6.2 2ℓ Same-Sign WH Channel

The 2ℓ same-sign WH channel is characterized by the presence of two same-sign leptons and one or more jets. To achieve a final state with two same-sign leptons, one lepton must originate from the leptonic decay of the associated $W^\pm$ boson and the other lepton must originate from the leptonic decay of the $W^\pm$ boson produced in the $H \rightarrow WW^*$ decay. The 2ℓ same-sign WH channel is further divided into three subchannels based on the flavour of the leptons: $\mu\mu$, ee , and $e\mu/\mu e$.

The leading and subleading electrons are required to exceed transverse momentum thresholds of 22 and 15 GeV. A selection on the dilepton angular separation, $\Delta R_{\ell\ell} > 0.4$, is used to suppress $W+\gamma$, the largest source of fake electrons, and a b -jet veto is applied to reject $t\bar{t}$, a source of fake leptons and charge-flip electrons. For events with two or more jets, a selection on the dijet invariant mass, $m_{jj} < 500$ GeV, is used to suppress electroweak and strong production of same-sign WW . For the ee channel, a window cut on the dilepton invariant mass, $|m_{\ell\ell} - m_Z| > 20$ GeV, is used to reject Z +jets, a source of charge-flip electrons.

The 2ℓ same-sign WH channel uses a recurrent neural network (RNN) to extract the WH , $H \rightarrow WW^*$ process. Within each of the $\mu\mu$, ee , and $e\mu/\mu e$ signal regions, the RNN output score is optimally binned in its full range, $[0, 1]$, and provided to the combined fit. The same binning is used for all signal regions.

The selections for each region of the 2ℓ same-sign WH channel are shown in Table 6.5. While there are no dedicated control regions, the normalization of the (≥ 1 -jets) WZ background is left floating in the fit and is constrained by the background-like bins of the RNN. The fake lepton and charge-flip electron backgrounds are estimated using the data-driven approach described in Chapter 5. The shape and normalization of all other backgrounds – which include same-sign WW , ZZ , triboson, and $t\bar{t}V/tZ$ production – are taken from simulation. To validate the modelling of the RNN, two validation regions are defined based on the requirement of a third lepton. The selections applied to the third lepton are shown in Table 6.6.

Table 6.5: Event selections for the signal (SR) and validation (VR) regions of the 2ℓ same-sign WH channel. The selections applied to the third lepton in the validation regions are shown in Table 6.6.

$\mu\mu$ SR	2ℓ same-sign WH SR		Loose RNN VR	Tight RNN VR
	ee SR	$e\mu/\mu e$ SR		
Exactly 2 same-sign ID leptons with electrons passing ECIDS				
At least 1 ID lepton matched to a single lepton trigger				
$p_T^{\ell_{\text{lead}}} > 22 \text{ GeV} \ \&\& \ p_T^{\ell_{\text{sublead}}} > 15 \text{ GeV}$				
$\Delta R_{\ell\ell} > 0.4$				
$N_{\text{jets}} \geq 1$				
$N_{b\text{-jets}} = 0$				
$m_{jj} < 500 \text{ GeV}$ if $N_{\text{jets}} \geq 2$				
Lepton flavour: $\mu\mu$	Lepton flavour: ee $ m_{\ell\ell} - m_Z > 20 \text{ GeV}$	Lepton flavour: $e\mu/\mu e$	Third lepton	Third lepton, fails ID

6.2.1 Multivariate Analysis

The multivariate discriminant used by the 2ℓ same-sign WH channel is a recurrent neural network (RNN). A RNN is a type of ANN that operates on *variable*-size input sequences. In the case of the 2ℓ same-sign WH channel, these input sequences consist of the leptons, jets, and missing transverse momentum belonging to a given event.

The RNN is trained to discriminate between WH and WZ production, the latter corresponding to the dominant electroweak background. The events used during training are required to have exactly two same-sign

Table 6.6: The selections applied to the third lepton in the loose and tight RNN validation regions of the 2ℓ same-sign WH channel. Additionally, the lepton is required to pass overlap removal.

Object	Selection	
Electrons	Transverse momentum	$p_T > 15 \text{ GeV}$
	Pseudorapidity	$(\eta < 1.37) \parallel (1.52 < \eta < 2.47)$
	Impact parameters	$ \Delta z \sin \theta < 0.5 \text{ mm} \ \&\& \ d_0 /\sigma_{d_0} < 5$
	Identification	Medium
	Isolation	Loose
Muons	Transverse momentum	$p_T > 15 \text{ GeV}$
	Pseudorapidity	$ \eta < 2.5$
	Impact parameters	$ \Delta z \sin \theta < 0.5 \text{ mm} \ \&\& \ d_0 /\sigma_{d_0} < 3$
	Identification	Medium
	Isolation	Loose

leptons and one or more jets – there is no splitting based on the flavour of the leptons. To increase the training statistics, the leptons and jets are required to pass a modified set of object selections, which are shown in Table 6.7. The selections for leptons and jets are looser and tighter, respectively, than those given in Table 4.1.

Table 6.7: The selections applied to the leptons and jets which are used to train the RNN of the 2ℓ same-sign WH channel. All objects are required to pass overlap removal.

Object	Selection	
Electrons	Transverse momentum	$p_T > 10 \text{ GeV}$
	Pseudorapidity	$(\eta < 1.37) \parallel (1.52 < \eta < 2.47)$
	Author	Author = 1
	Identification	Loose
	Isolation	Loose
Muons	Transverse momentum	$p_T > 10 \text{ GeV}$
	Pseudorapidity	$ \eta < 2.5$
	Identification	Medium
	Isolation	Loose
Jets	Transverse momentum	$\begin{cases} p_T > 25 \text{ GeV} & \eta < 2.5 \\ p_T > 30 \text{ GeV} & 2.5 < \eta < 4.5 \end{cases}$
	Pseudorapidity	$ \eta < 4.5$
	JVT	Tight
	fJVT	Loose

It has been shown that neural networks are capable of reproducing nonlinear features of high-energy particle-particle collisions using basic kinematic quantities [474]. For example, a neural network may “learn” to calculate the invariant mass of a lepton-jet system using only the individual objects’ four-momenta. Inspired by this study, the RNN receives only the p_T , η , and ϕ for each lepton and jet as well as the magnitude and azimuthal angle of the missing transverse momentum. Additionally, the RNN receives an integer variable indicating the type of each object: lepton, jet, or missing transverse momentum. The input variables are listed in Table 6.8.

The geometry of the ATLAS has a rotation symmetry in the ϕ direction and a reflection symmetry in the η direction. Therefore, the RNN should yield the same result if the system of input objects is rotated by an arbitrary ϕ or reflected in the η direction. These symmetries are encoded by preprocessing the input objects

Table 6.8: Input variables to the RNN used by the 2ℓ same-sign WH channel.

Variable	Description
p_T^ℓ	Lepton transverse momentum
η_ℓ	Lepton pseudorapidity
ϕ_ℓ	Lepton azimuthal angle
Object type	1 for leptons
E_T^{miss}	Missing transverse momentum
η_{miss}	Pseudorapidity of E_T^{miss} (always 0)
ϕ_{miss}	Azimuthal angle of E_T^{miss}
Object type	0 for E_T^{miss}
p_T^j	Jet transverse momentum
η_j	Jet pseudorapidity
ϕ_j	Jet azimuthal angle
Object type	-1 for jets

prior to training such that the leading lepton is always at $\phi = 0$ and positive η . Shape comparisons among the WH and WZ classes for each variable following preprocessing are shown in Figs. 6.6.

The input objects are provided to the RNN as a sequence of $(p_T, \eta, \phi, \text{type})$ vectors. As the recurrent layers [475] used by the RNN are sensitive to the ordering of the input objects, the input objects for each event are always provided to the RNN in the following order:

1. Leading lepton in p_T ;
2. Subleading lepton in p_T ;
3. Missing transverse momentum;
4. Up to five jets, ordered from highest-to-lowest in p_T .

The RNN is trained using Keras [472] with the TensorFlow [473] backend. To assess whether the RNN is overtrained, it is deployed on a dedicated set of testing events. The shape comparison among the training and testing datasets in the RNN output score is shown in Fig. 6.7. Good agreement is observed, and so it may be concluded that no overtraining is present.

Distributions of the RNN’s output score in the 2ℓ same-sign WH signal regions are shown in Fig. 6.8. Good agreement is observed between data and simulation except in the rightmost bin of the $\mu\mu$ signal region, where there is a deficit in data: $44.3 \pm 4.1(\text{stat.})$ events were expected whereas 28 events were measured.

6.2.2 Validation Regions

Validation regions are used to assess the quality of the modelling of the RNN output score in data. They are constructed using the $\mu\mu$, ee , and $e\mu/\mu e$ signal region selections with the additional requirement of a third lepton. This lepton is required to pass a looser set of selections, given in Table 6.6, in the “loose” RNN validation region. The “tight” RNN validation region, which is a subset of the loose RNN validation region, additionally requires the third lepton to fail ID selections, given in Table 4.1. While the loose RNN validation region has larger statistics, the phase space of the tight RNN validation region is closer to that of the signal region.

Distributions of the RNN’s output score in all six RNN validation regions are shown in Fig. 6.9. Adequate agreement between data and simulation is observed in all distributions. A moderate deficit is observed in data

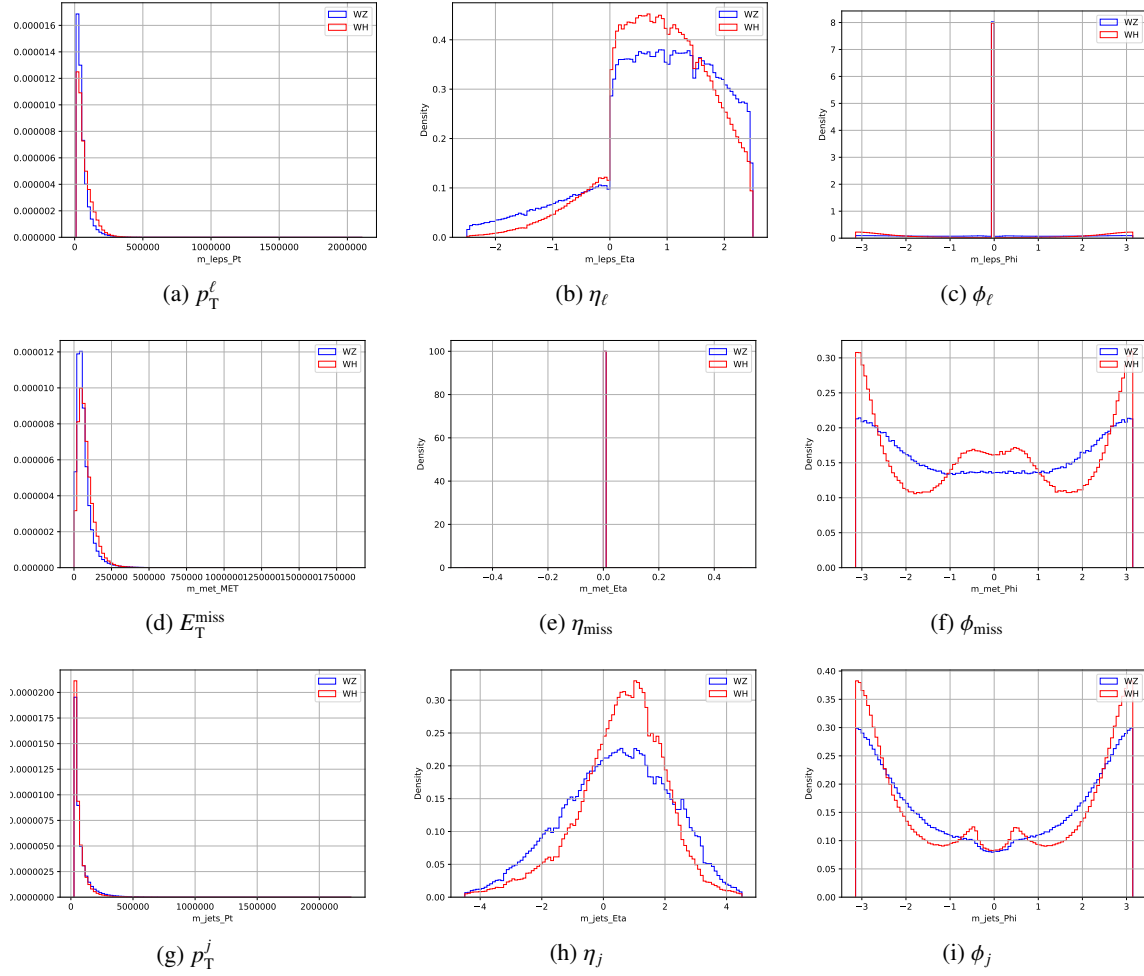


Figure 6.6: Shape comparisons among the WH and WZ classes for each input variable of the RNN used by the 2ℓ same-sign WH channel. The η and ϕ variables are subject to the reflections and rotations described in the text. The sum-of-weights for each class is normalized to unity.

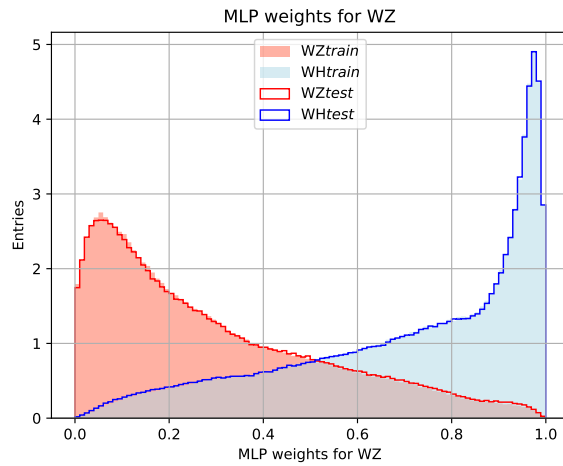


Figure 6.7: Output score for the RNN used by the 2ℓ same-sign WH channel. The RNN is deployed on both the training and testing datasets to assess the level of agreement. The sum-of-weights for each class is normalized to unity.

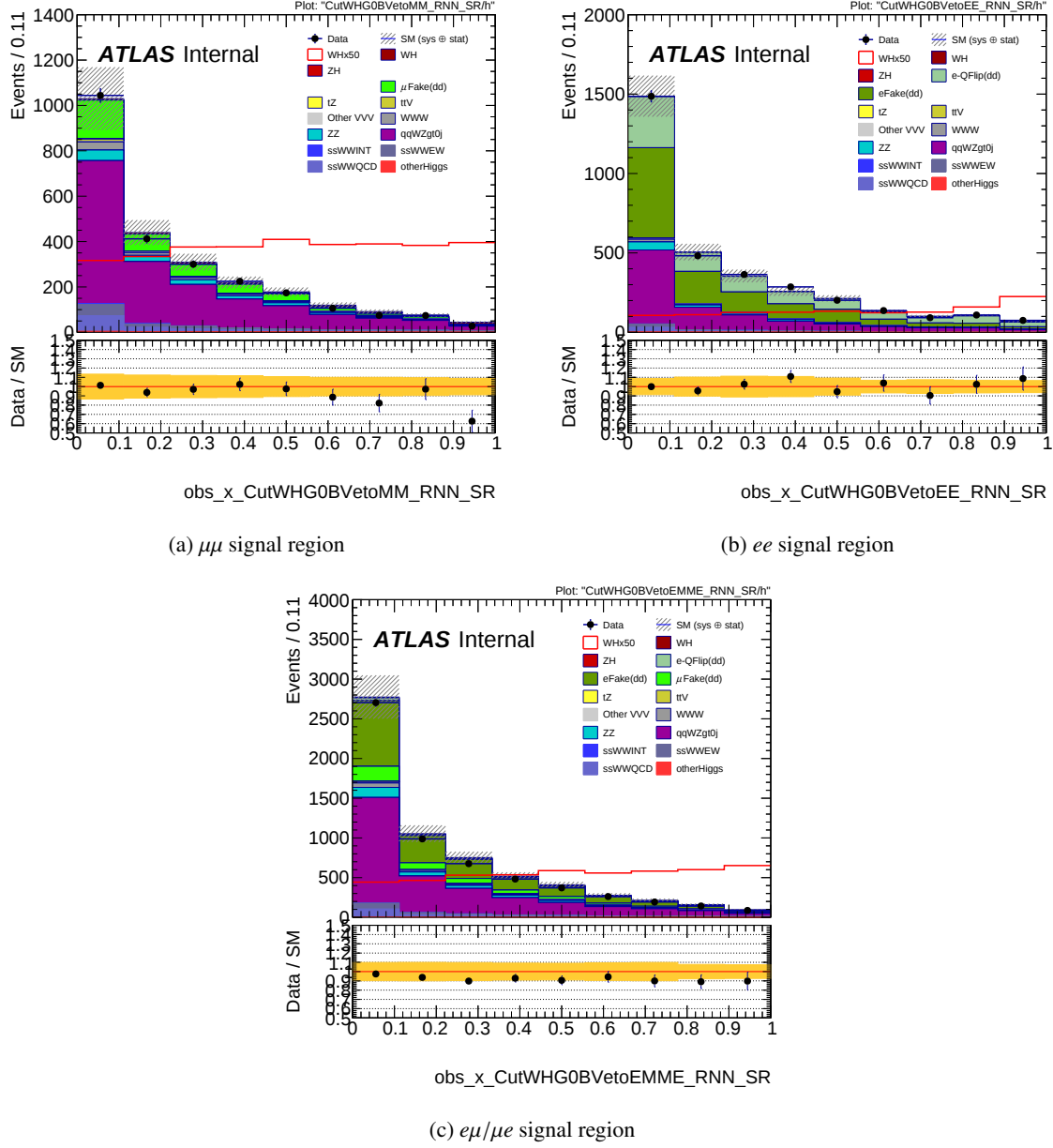


Figure 6.8: Distributions of the RNN's output score in the (a) $\mu\mu$, (b) ee , and (c) $e\mu/\mu e$ signal regions of the 2ℓ same-sign WH channel. The distributions are remapped to bins of equal width. The error bands include the statistical uncertainty as well as the dominant experimental and theoretical systematic uncertainties. The WH process is displayed both as stacked as well as unstacked and scaled up by a factor of 50 for visibility.

in the rightmost bins of $\mu\mu$ and $e\mu/\mu e$ channels. Further investigation revealed a significant mismodelling in events with exactly two jets in the loose RNN validation regions, as shown in Fig. 6.10. Notably, there is good modelling of events with a number of jets not equal to two. As the validation regions are highly pure in WZ events, the mismodelling is attributed to the WZ background. An alternative multivariate technique – a boosted decision tree – as well as alternative lepton isolation working points¹ were both tested, but neither were found to improve the modelling.

As the source of the mismodelling could not be identified, a RNN shape systematic is assigned to the WZ background instead. For each of the $\mu\mu$, ee , and $e\mu/\mu e$ RNN validation regions, the WZ event yield is normalized to the data yield ($N_{WZ}^{\text{normalized,VR}}$), and the per-bin ratio of the data yield ($N_{\text{data}}^{\text{VR}}$) less the non- WZ event yield ($N_{\text{non-}WZ}^{\text{VR}}$) to the normalized WZ event yield is multiplied onto the WZ event yield in the corresponding signal region bin ($N_{WZ}^{\text{nominal,SR}}$) to yield the varied WZ event in that same signal region bin ($N_{WZ}^{\text{varied,SR}}$):

$$N_{WZ}^{\text{varied,SR}} = \frac{N_{\text{data}}^{\text{VR}} - N_{\text{non-}WZ}^{\text{VR}}}{N_{WZ}^{\text{normalized,VR}}} \times N_{WZ}^{\text{nominal,SR}}. \quad (6.2)$$

The absolute difference between the varied and the nominal WZ event yields in each signal region bin is symmetrized and assigned as the corresponding uncertainty. The size of this uncertainty varies from 4% to 18%.

¹In the $\mu\mu$ and $e\mu/\mu e$ channels, the mismodelling appears strongly at low values of ΔR between a muon and the nearest jet. The isolation efficiency for muons often depends on this quantity, hence making the choice of lepton isolation working point a possible candidate for the mismodelling.

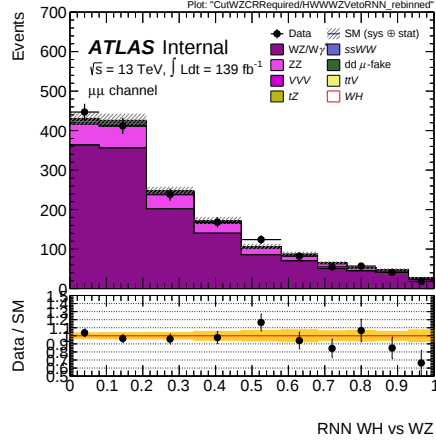
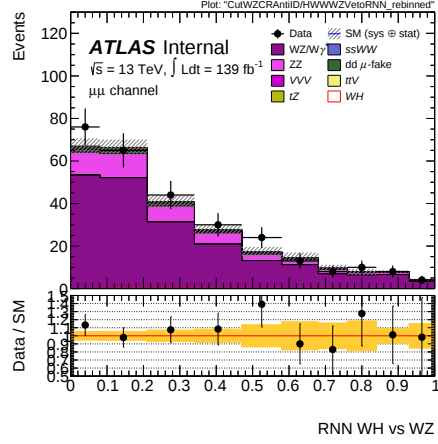
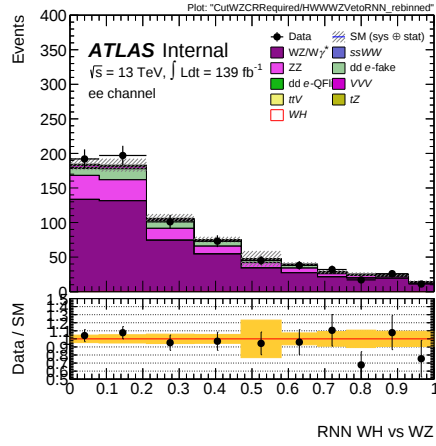
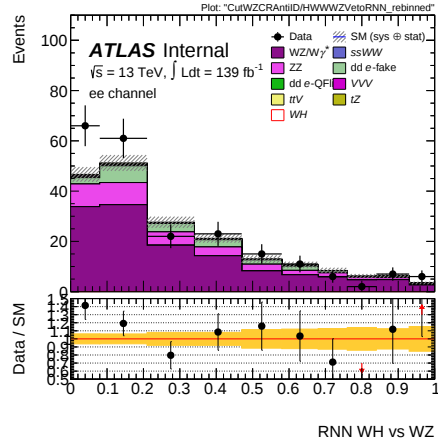
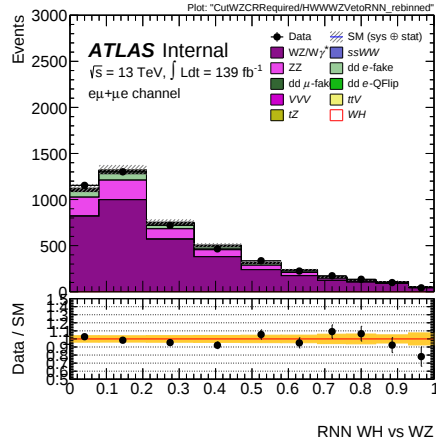
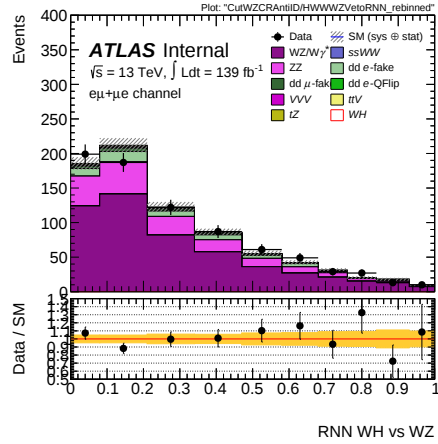
(a) $\mu\mu$ loose RNN validation region(b) $\mu\mu$ tight RNN validation region(c) ee loose RNN validation region(d) ee tight RNN validation region(e) $e\mu/\mu e$ loose RNN validation region(f) $e\mu/\mu e$ tight RNN validation region

Figure 6.9: Distributions of the RNN's output score in the (a, b) $\mu\mu$, (c, d) ee , and (e, f) $e\mu/\mu e$ RNN validation regions of the 2ℓ same-sign WH channel. Both the loose (left) and tight (right) RNN validation regions are shown. The error bands include the statistical uncertainty as well as the dominant experimental systematic uncertainties.

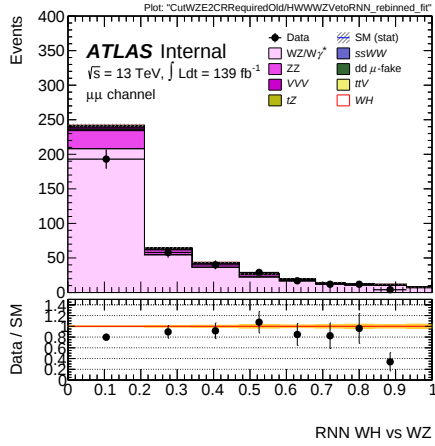
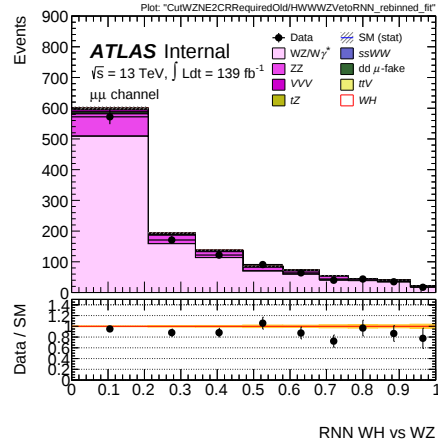
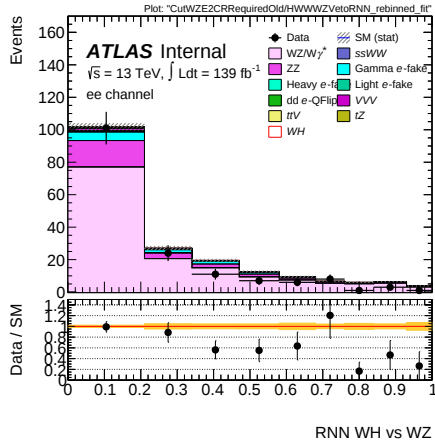
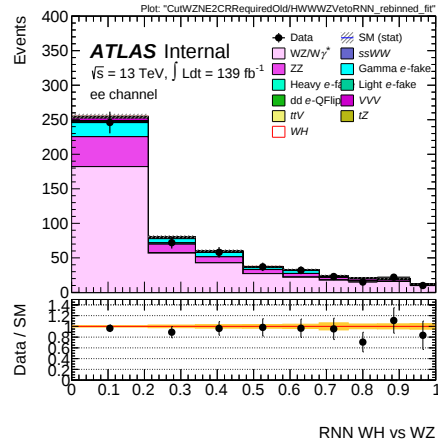
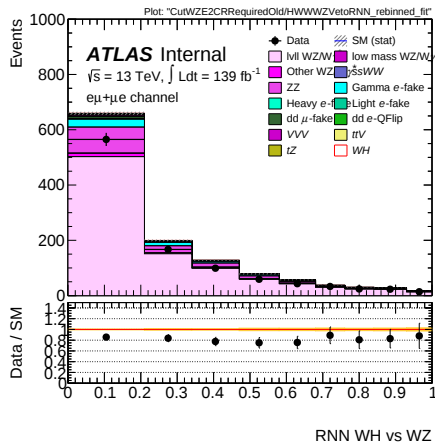
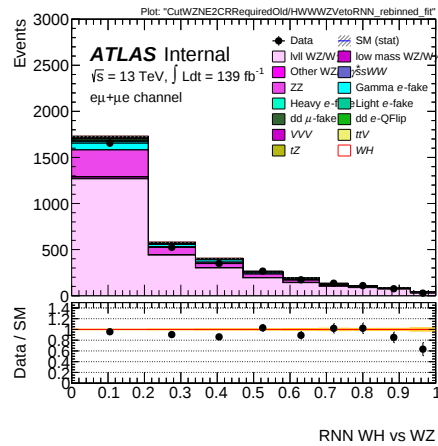
(a) $\mu\mu$ channel, $N_{\text{jets}} = 2$ (b) $\mu\mu$ channel, $N_{\text{jets}} \neq 2$ (c) ee channel, $N_{\text{jets}} = 2$ (d) ee channel, $N_{\text{jets}} \neq 2$ (e) $e\mu/\mu e$ channel, $N_{\text{jets}} = 2$ (f) $e\mu/\mu e$ channel, $N_{\text{jets}} \neq 2$

Figure 6.10: Distributions of the RNN's output score in the (a, b) $\mu\mu$, (c, d) ee , and (e, f) $e\mu/\mu e$ loose RNN validation regions of the same-sign WH channel. Within each validation region, events are split according to $N_{\text{jets}} = 2$ (left) and $N_{\text{jets}} \neq 2$ (right). The error bands indicate the statistical uncertainty.

6.3 3ℓ Z-Dominated WH Channel

The 3ℓ WH channel is characterized by the presence of three leptons of total charge ± 1 . One lepton originates from the leptonic decay of the associated W boson and the other two leptons originate from the $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$ decay. The small opening angle between the leptons from $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$ decay naturally leads to a topology-based lepton naming scheme – the lepton with the unique charge must come from the $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$ decay, as shown in Fig. 1.5c, and is denoted as ℓ_0 . The lepton closest to ℓ_0 in ΔR is denoted as ℓ_1 , and the remaining lepton is denoted as ℓ_2 . In 99% of cases, ℓ_0 originates from the decay of the Higgs boson; in 85% of cases, ℓ_1 originates from the decay of the Higgs boson.

Between the three leptons, the Z-dominated channel requires one or more same-flavour, opposite-sign (SFOS) lepton pairs, making it “dominated” by Z decays. In contrast, the Z-depleted channel requires zero SFOS lepton pairs, making it “depleted” in Z decays. While the Z-depleted channel shares a common signature with the Z-dominated channel, its analysis strategy is sufficiently different such that it is given a dedicated section, Section 6.4. The remainder of this section focuses exclusively on the Z-dominated channel.

All three leptons are required to exceed a transverse momentum threshold of 15 GeV. A selection on the dilepton angular separation, $\Delta R_{\ell\ell} > 0.1$, is used to remove overlapping leptons and is applied to all combinations of leptons, and a selection on the minimum dilepton invariant mass of all SFOS lepton pairs, $m_{\ell\ell,\text{SFOS}}^{\min} > 12$ GeV, is used to remove low-mass resonances. A cut on the missing transverse momentum, $E_T^{\text{miss}} > 30$ GeV, is used to select final states with neutrinos, and a b -jet veto is applied to reject $t\bar{t}$ production, a source of fake leptons. A window cut on the invariant mass of all SFOS lepton pairs, $|m_{\ell\ell,\text{SFOS}} - m_Z| > 25$ GeV, is used to veto Z boson production. This includes WZ and ZZ production as well as Z +jets production, a source of fake leptons.

The 3ℓ Z-dominated WH channel uses an ANN to extract the WH , $H \rightarrow WW^*$ process. Within the signal region, the ANN output score is optimally binned in its full range, $[0, 1]$, and provided to the combined fit.

The selections for each region of the 3ℓ Z-dominated WH channel are shown in Table 6.9. Control regions are used to separately normalize the components the WZ background with zero reconstructed jets (“0-jet WZ ”) and one or more reconstructed jets (“ ≥ 1 -jets WZ ”) to data. The rationale for doing so is that the theoretical uncertainties on WZ production affect each component very differently, and so decorrelating the two components results in a more stable statistical fit. The fake lepton background is estimated using the data-driven approach described in Chapter 5. The shape and normalization of all other backgrounds – which include ZZ , triboson, and $t\bar{t}V/tZ$ production – are taken from simulation. A top validation region based on b -jet requirement is used to assess the modelling of the top quark background.

6.3.1 Multivariate Analysis

The multivariate discriminant utilized by the 3ℓ Z-dominated VH channel is an ANN. The ANN is trained to discriminate between WH and WZ production, the latter corresponding to the dominant background. For the purposes of training, dedicated WH and WZ training samples are used, ensuring there is no bias introduced by deploying the ANN on the same events which were used to train it. The alternative WH sample, like the nominal sample, is generated with PowHEG Box v2 [287–289, 374] and showered with PYTHIA 8 [290]; the alternative WZ sample, like the nominal sample, is generated with SHERPA 2.2.2 [291].

For each event, the ANN receives kinematic and topological input variables related to the trilepton system and the missing transverse momentum. Also included is the transverse mass of the W boson:

Table 6.9: Event selections for the signal (SR), control (CR), and validation (VR) regions of the 3ℓ Z-dominated WH channel.

3ℓ Z-dominated WH SR	0-jet WZ CR	≥ 1 -jets WZ CR	Top VR
Exactly 3 ID leptons with total charge ± 1 At least 1 ID lepton matched to a single lepton trigger All lepton $p_T > 15$ GeV All lepton pairs $\Delta R_{\ell\ell} > 0.1$ ≥ 1 same-flavour, opposite-sign (SFOS) lepton pairs $m_{\ell\ell, \text{SFOS}}^{\min} > 12$ GeV			
$E_T^{\text{miss}} > 30$ GeV $N_{b\text{-jets}} = 0$ 0 SFOS pairs with $ m_{\ell\ell, \text{SFOS}} - m_Z < 25$ GeV	$N_{\text{jets}} = 0$	$E_T^{\text{miss}} > 30$ GeV $N_{b\text{-jets}} = 0$ ≥ 1 SFOS pairs with $ m_{\ell\ell, \text{SFOS}} - m_Z < 25$ GeV	$E_T^{\text{miss}} > 50$ GeV $N_{\text{jets}} \geq 1$ $N_{b\text{-jets}} = 1$ 0 SFOS pairs with $ m_{\ell\ell, \text{SFOS}} - m_Z < 25$ GeV

$$m_T^W = \sqrt{2p_T^{\ell_W} E_T^{\text{miss}} (1 - \cos \Delta\phi_{\ell_W, \text{miss}})}, \quad (6.3)$$

where ℓ_W is the lepton, either ℓ_1 or ℓ_2 , *not* belonging to the SFOS lepton pair with an invariant mass closest to the mass of the Z boson. The input variables are listed in Table 6.10, and the inputs themselves are produced within the 3ℓ Z-dominated WH signal region. Distributions of the input variables in the signal region are shown in Figs. 6.11 and 6.12. Prior to training, each variable has its mean subtracted and is divided by its standard deviation. This has the effect of roughly centring each variable on 0 with a range that extends to $O(\pm 1)$.

The ANN is trained using Keras [472] with the TensorFlow [473] backend. To assess whether the ANN is overtrained, it is deployed on a dedicated set of testing events. The shape comparison among the training and testing datasets in the ANN output score is shown in Fig. 6.13. Good agreement is observed, and so it may be concluded that no overtraining is present.

Table 6.10: Input variables to the ANN used by the 3ℓ Z-dominated WH channel.

Variable(s)	Description
$p_T^{\ell_0}$	Transverse momentum of ℓ_0
$ \sum_{i=0}^2 \vec{p}_T^{\ell_i} $	Magnitude of the vectorial sum of the lepton transverse momenta
$\Delta\eta_{\ell_0\ell_1}, \Delta\eta_{\ell_1\ell_2}$	Pseudorapidity separation between ℓ_0 and ℓ_1 as well as ℓ_1 and ℓ_2
$\Delta\phi_{\ell_0\ell_2}$	Azimuthal separation between ℓ_0 and ℓ_2
$\Delta R_{\ell_0\ell_1}, \Delta R_{\ell_0\ell_2}$	Angular separation between ℓ_0 and ℓ_1 as well as ℓ_0 and ℓ_2
$m_{\ell_0\ell_1}, m_{\ell_0\ell_2}, m_{\ell_1\ell_2}$	Dilepton invariant mass for each combination of leptons
E_T^{miss}	Missing transverse momentum
$\Delta\phi_{\ell_0, \text{miss}}, \Delta\phi_{\ell_1, \text{miss}}, \Delta\phi_{\ell_2, \text{miss}}$	Azimuthal separation between each lepton and the E_T^{miss}
m_T^W	Transverse mass of the W boson (Eq. 6.3)

The distribution of the ANN's output score in the 3ℓ Z-dominated WH signal region is shown in Fig. 6.14a. Good agreement is observed between data and simulation.

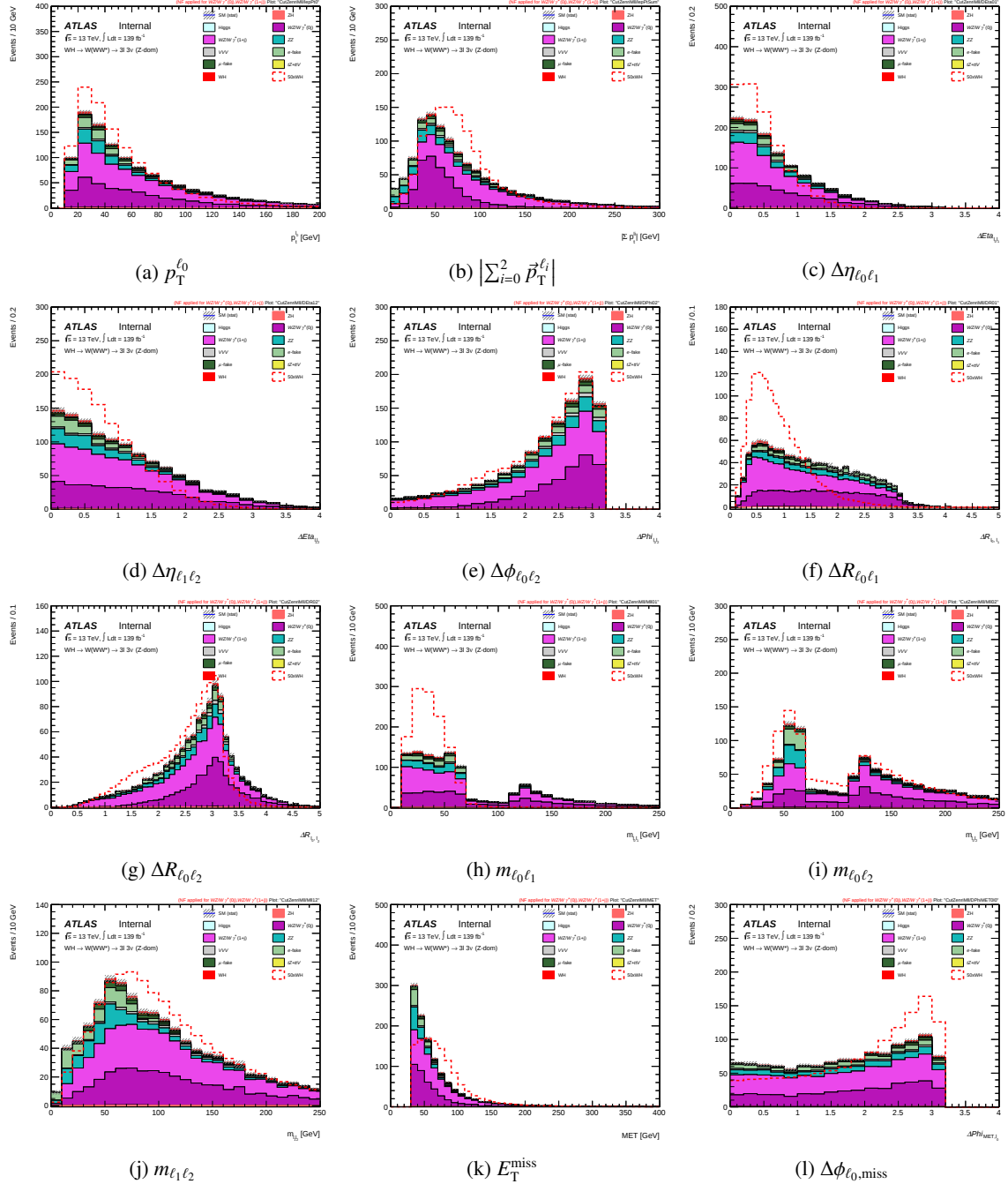


Figure 6.11: Distributions of the input variables to the ANN in the 3ℓ Z-dominated WH signal region. Normalization factors are applied to the 0-jet and ≥ 1 -jets WZ event yields. The error bands indicate the statistical uncertainty. The WH process is unstacked and scaled up by a factor of 50 for visibility.

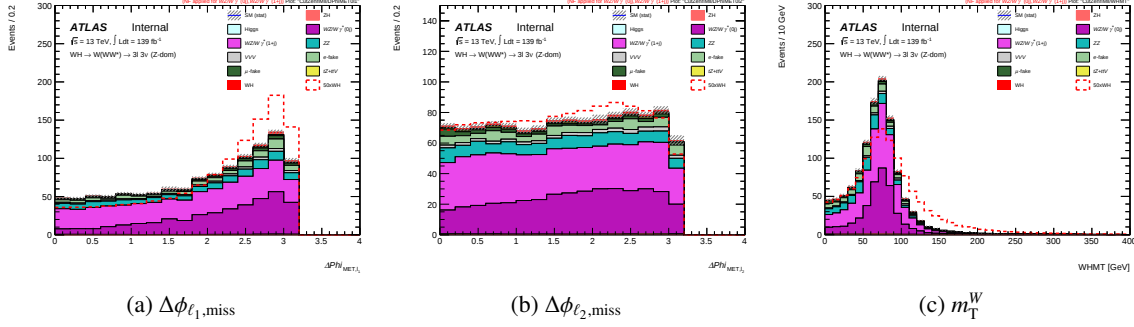


Figure 6.12: Distributions of the input variables to the ANN in the 3ℓ Z-dominated WH signal region. Normalization factors are applied to the 0-jet and ≥ 1 -jets WZ event yields. The error bands indicate the statistical uncertainty. The WH process is unstacked and scaled up by a factor of 50 for visibility.

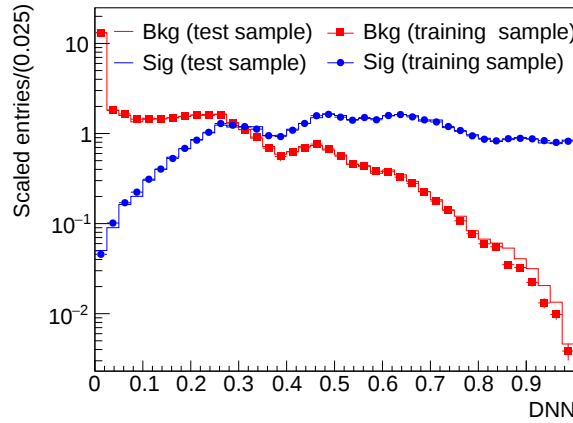


Figure 6.13: Output score for the ANN used by the 3ℓ Z-dominated WH channel. “Sig” corresponds to WH production and “Bkg” corresponds to WZ production. The ANN is deployed on both the training and testing datasets to assess the level of agreement. The error bars indicate the statistical uncertainty.

6.3.2 Control and Validation Regions

The 0-jet and ≥ 1 -jets WZ control regions require zero and one or more jets, respectively. They are constructed by inverting the Z boson veto of the signal region and requiring one or more SFOS lepton pairs with an invariant mass within 25 GeV of the mass of the Z boson. The 0-jet WZ control region is 92% pure in 0-jet WZ events with a corresponding normalization factor of 1.04, while the ≥ 1 -jets WZ control region is 88% pure in ≥ 1 -jets WZ events with a corresponding normalization factor of 0.90. Distributions of the ANN output score in both control regions are shown in Figs. 6.14b and 6.14c. With the application of the normalization factors, good agreement between data and simulation is observed.

The top validation region is used to assess the top quark modelling, which includes electroweak production of $t\bar{t}V/tZ$ as well as $t\bar{t}$ production with a misidentified lepton. This region is constructed by inverting the b -jet veto of the signal region and requiring one or more jets and one b -jet. Additionally, the missing transverse momentum cut is tightened with respect to the signal region: $E_T^{\text{miss}} > 50$ GeV. The distribution of the ANN’s output score in the top validation region is shown in Fig. 6.14d. Good agreement between data and simulation is observed, and so it may be concluded that top quark processes are well modelled.

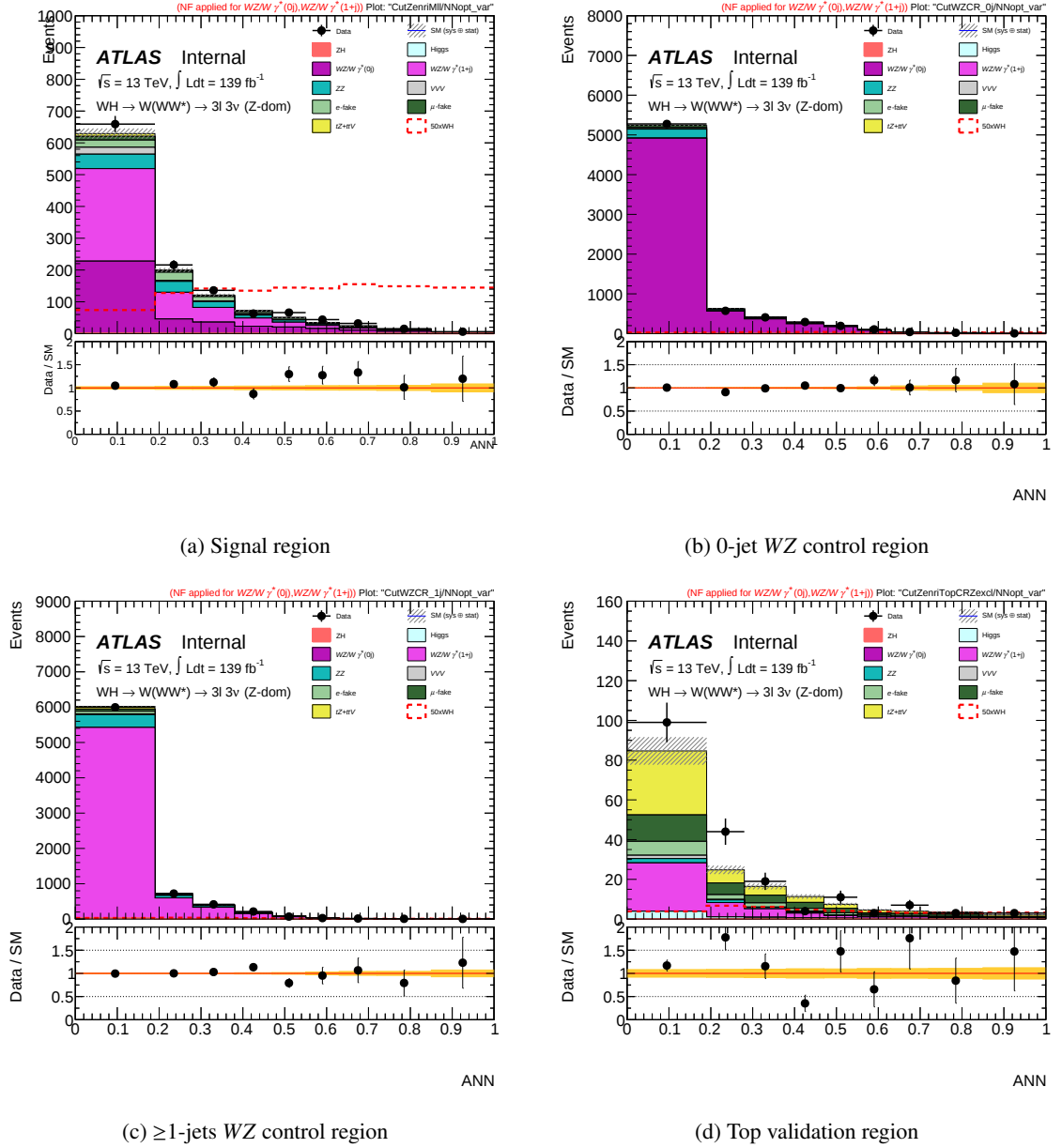


Figure 6.14: Distributions of the ANN output score in the (a) signal region, (b) 0-jet WZ control region, (c) ≥ 1 -jets WZ control region, and (d) top validation region of the 3ℓ Z-dominated WH channel. The signal region binning is used in all distributions. Normalization factors are applied to the 0-jet and ≥ 1 -jets WZ event yields. The error bands include the statistical uncertainty as well as the dominant experimental systematic uncertainties. The WH process is unstacked and scaled up by a factor of 50 for visibility.

6.4 3ℓ Z-Depleted WH Channel

The 3ℓ Z-depleted WH channel, first introduced in Section 6.3, is characterized by the presence of three leptons of total charge ± 1 . Between the three leptons, the Z-depleted channel requires zero SFOS lepton pairs. All three leptons are required to exceed a transverse momentum threshold of 15 GeV. A selection on the dilepton angular separation, $\Delta R_{\ell\ell} > 0.1$, is used to remove overlapping leptons and is applied to all combinations of leptons, and a b -jet veto is applied to reject $t\bar{t}$ production, a source of fake leptons.

The 3ℓ Z-depleted WH signal region is defined using a multiclassifier ANN discriminating between WH , $t\bar{t}$, WZ , and WWW production. Following the b -jet veto, a cut on the multiclassifier's $t\bar{t}$ output score, $\text{ANN}_{t\bar{t}} < 0.25$, is used to reject $t\bar{t}$ production, one of the largest backgrounds in the Z-depleted channel. This final cut constitutes the signal region. Within the signal region, the WH , $H \rightarrow WW^*$ process is extracted using a discriminant calculated as the difference between the multiclassifier's WH output score and its WZ and WWW output scores:

$$\text{ANN}_{\Delta} \equiv \text{ANN}_{WH} - \text{ANN}_{WZ} - \text{ANN}_{WWW} . \quad (6.4)$$

ANN_{Δ} is optimally binned in its full range, $[-1, +1]$, and provided to the combined fit.

The selections for each region of the 3ℓ Z-depleted WH channel are shown in Table 6.11. The Z-depleted channel shares the 0-jet and ≥ 1 -jets WZ control regions with the Z-dominated channel. Additionally, the normalization of the WWW background is left floating in the fit and is ultimately constrained by the background-like bins of the ANN_{Δ} distribution. The fake lepton background is estimated using the data-driven approach described in Chapter 5. The shape and normalization of all other backgrounds – which include ZZ , non- WWW triboson, and $t\bar{t}V/tZ$ production – are taken from simulation. A top validation region based on b -jet requirement is used to assess the modelling of the top quark background. This validation region is different from that used by the Z-dominated channel.

Table 6.11: Event selections for the signal (SR) and validation (VR) regions of the 3ℓ Z-depleted WH channel. The artificial neural network (ANN) is described in Section 6.4.1. The Z-depleted channel shares the 0-jet and ≥ 1 -jets WZ control regions, given in Table 6.9, with the Z-dominated channel.

3ℓ Z-depleted WH SR	WWW VR	Top VR
Exactly 3 ID leptons with total charge ± 1 with electrons passing ECIDS		
At least 1 ID lepton matched to a single lepton trigger		
All lepton $p_T > 15$ GeV		
All lepton pairs $\Delta R_{\ell\ell} > 0.1$		
0 same-flavour, opposite-sign (SFOS) lepton pairs		
–		$N_{\text{jets}} \geq 1$
$N_{b\text{-jets}} = 0$		$N_{b\text{-jets}} = 1$
$\text{ANN}_{t\bar{t}} < 0.25$		No ECIDS requirement
–	$\text{ANN}_{WWW} > 0.5$	

6.4.1 Multivariate Analysis

The multiclassifier is trained using the nominal WH and WWW samples as well as alternative WZ and $t\bar{t}$ samples. The alternative WZ sample is identical to that used to train the ANN of the Z-dominated channel. The $t\bar{t}$ sample is generated with POWHEG Box v2 [287–289, 378] and showered with PYTHIA 8 [290].

To avoid possible bias by training the multiclassifier with the same events it will be deployed on, the input variables entering the training are produced in different regions of phase space. For the WH and WWW samples, the inputs are produced in the Z-dominated channel at the requirement of one or more SFOS lepton pairs, given in Table 6.9. Both WH and WWW production do not include a Z boson and therefore should not depend on the number of SFOS lepton pairs. For the $t\bar{t}$ and WZ samples, the inputs are produced in the Z-depleted channel at the requirement of zero SFOS lepton pairs, given in Table 6.11. Additionally, the isolation criteria applied to all leptons is loosened from the PLImprovedTight working point to the Tight working point to enhance statistics.

For each event, the multiclassifier receives kinematic and topological input variables related to the trilepton system, the reconstructed jets, and the missing transverse momentum. To separate $WZ \rightarrow W\tau\tau \rightarrow \ell\nu\ell\nu\ell\nu$ decays, the compatibility of the event with the WZ hypothesis, F_α , is also provided. Given the reconstructed lepton momenta and the vectorial missing transverse momentum, the event kinematics can be calculated under the $WZ \rightarrow W\tau\tau \rightarrow \ell\nu\ell\nu\ell\nu$ (“WZ”) hypothesis assuming collinear tau decays. This yields a single unknown: the ratio of the energy of a single tau lepton to the energy of the lepton originating from its decay. This unknown is varied and the number of physical solutions is taken as a measure of the compatibility with the WZ hypothesis. The input variables are listed in Table 6.12, and distributions of the input variables in the signal region are shown in Figs. 6.15 through 6.17. Prior to training, each variable has its median subtracted and is scaled to its interquartile range using the Python package `scikit-learn` [471].

The multiclassifier is trained using Keras [472] with the TensorFlow [473] backend. To assess whether the multiclassifier is overtrained, it is deployed on a dedicated set of testing events. For the purposes of testing, all inputs are produced in the Z-depleted channel at the requirement of zero SFOS lepton pairs – the nominal WZ and $t\bar{t}$ samples are also used. Shape comparison among the training and testing datasets in the multiclassifier’s output scores are shown in Fig. 6.18. Good agreement is observed, and so it may be concluded that no overtraining is present.

Table 6.12: Input variables to the multiclassifier used by the 3ℓ Z-depleted WH channel.

Variable(s)	Description
$p_T^{\ell_0}, p_T^{\ell_1}, p_T^{\ell_2}$	Transverse momentum for each lepton
$\Delta\eta_{\ell_0\ell_1}, \Delta\eta_{\ell_0\ell_2}, \Delta\eta_{\ell_1\ell_2}$	Dilepton pseudorapidity separation for each combination of leptons
$\Delta R_{\ell_0\ell_1}, \Delta R_{\ell_0\ell_2}, \Delta R_{\ell_1\ell_2}$	Dilepton angular separation for each combination of leptons
$m_{\ell_0\ell_1}, m_{\ell_0\ell_2}, m_{\ell_1\ell_2}$	Dilepton invariant mass for each combination of leptons
$m_T^{\ell_0\ell_1}, m_T^{\ell_0\ell_2}, m_T^{\ell_1\ell_2}$	Dilepton transverse mass for each combination of leptons
$ \sum_{i=0}^2 \vec{p}_T^{\ell_i} $	Magnitude of the vectorial sum of the lepton transverse momenta
$m_{\ell\ell\ell}$	Trilepton invariant mass
N_{jets}	Number of jets
$N_{b\text{-jets}}$	Number of b -jets
$p_T^{j_0}$	Transverse momentum of the leading jet (0 if $N_{\text{jets}} = 0$)
E_T^{miss}	Missing transverse momentum
$\Delta\phi_{\ell_0, \text{miss}}, \Delta\phi_{\ell_1, \text{miss}}, \Delta\phi_{\ell_2, \text{miss}}$	Azimuthal separation between each lepton and the E_T^{miss}
$S_{\text{miss}}/E_T^{\text{miss}}$	Object-based E_T^{miss} significance divided by the E_T^{miss}
F_α	Compatibility of the event with the WZ hypothesis

The distribution of ANN_Δ in the 3ℓ Z-depleted signal region is shown in Fig. 6.19a. Good agreement is observed between data and simulation except in the leftmost bin, where an excess in data is observed: ~ 24 events were expected whereas 36 events were measured.

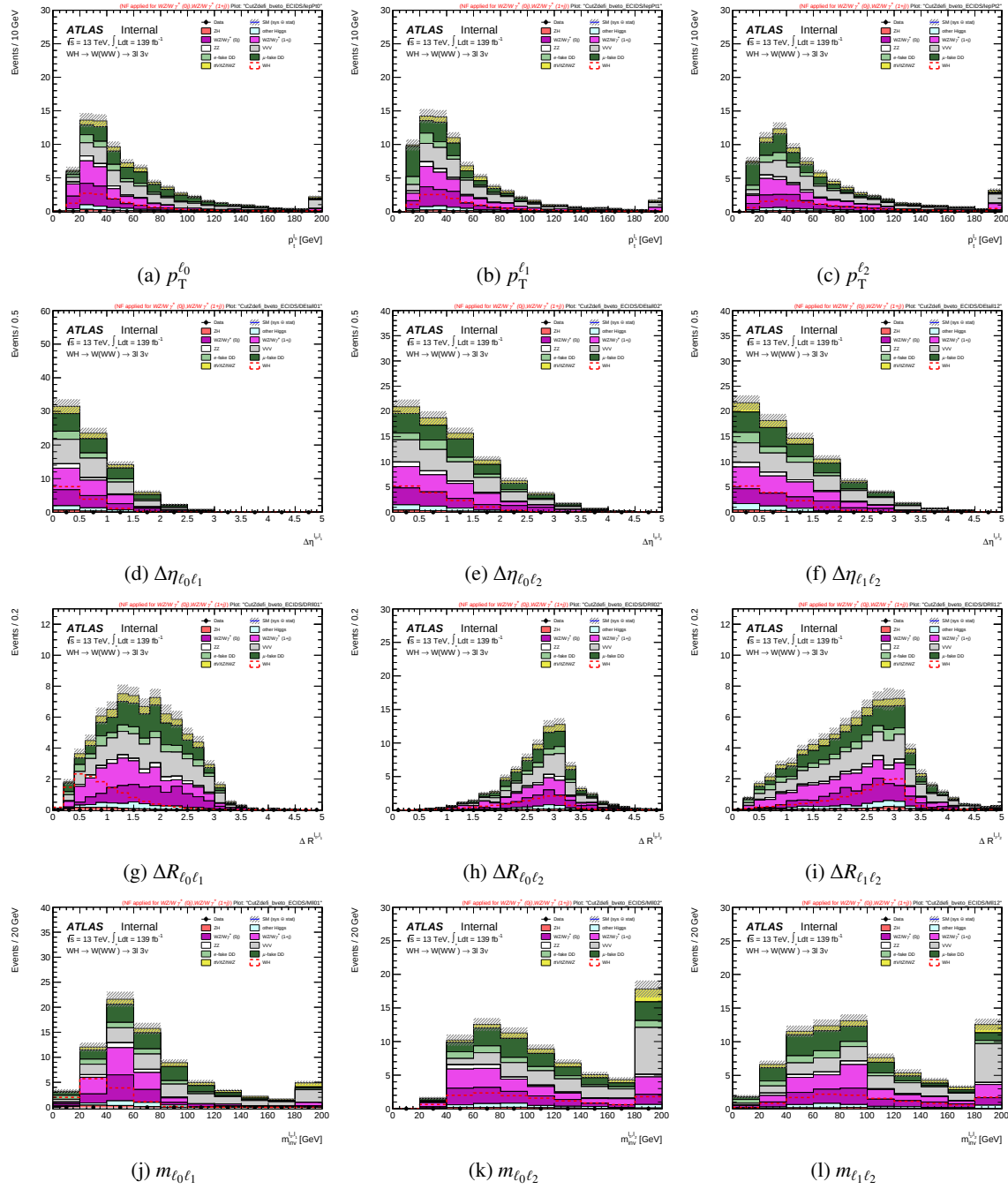


Figure 6.15: Distributions of the input variables to the multiclassifier in the 3ℓ Z-depleted WH signal region. Normalization factors are applied to the 0-jet and ≥ 1 -jets WZ event yields. The error bands include the statistical uncertainty as well as the dominant experimental systematic uncertainties. The WH process is unstacked for visibility.

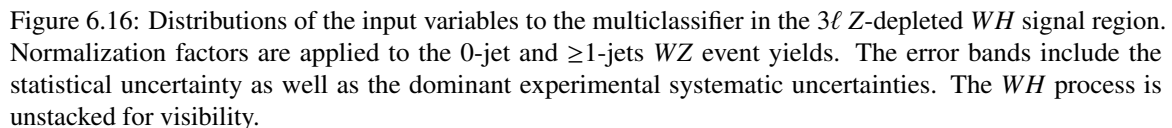


Figure 6.16: Distributions of the input variables to the multiclassifier in the 3ℓ Z-depleted WH signal region. Normalization factors are applied to the 0-jet and ≥ 1 -jets WZ event yields. The error bands include the statistical uncertainty as well as the dominant experimental systematic uncertainties. The WH process is unstacked for visibility.

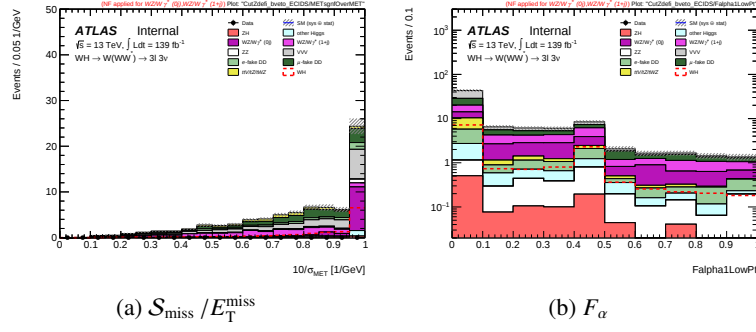


Figure 6.17: Distributions of the input variables to the multiclassifier in the 3ℓ Z-depleted WH signal region. Normalization factors are applied to the 0-jet and ≥ 1 -jets WZ event yields. The error bands include the statistical uncertainty as well as the dominant experimental systematic uncertainties. The WH process is unstacked for visibility.

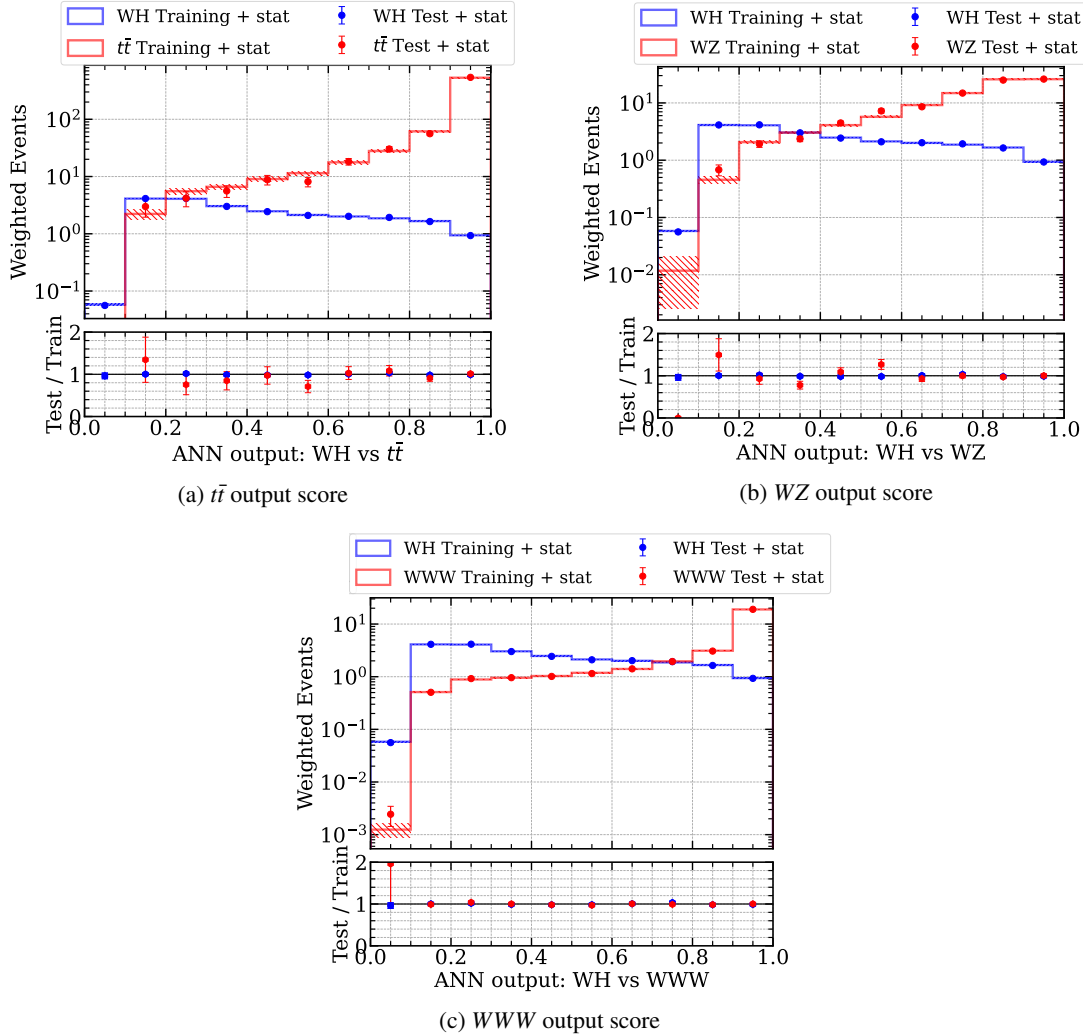


Figure 6.18: Output scores for the multiclassifier used by the 3ℓ Z-depleted WH channel. In each subplot, only the class targeted by the associated output score and the WH class are shown. The multiclassifier is deployed on both the training and testing datasets to assess the level of agreement. The error bands/bars indicate the statistical uncertainty.

6.4.2 Control and Validation Regions

The Z-depleted channel shares the 0-jet and ≥ 1 -jets WZ control regions with the Z-dominated channel. Distributions of ANN_Δ in both control regions are shown in Figs. 6.19b and 6.19c. With the application of the normalization factors, good agreement between data and simulation is observed.

The top validation region is used to assess the top quark modelling, which includes electroweak production of $t\bar{t}V/tZ$ as well as $t\bar{t}$ production with a misidentified lepton. This region is constructed by inverting the b -jet veto of the signal region and requiring one or more jets and one b -jet. Additionally, the top validation region does not include cuts on the multiclassifier's $t\bar{t}$ output score and does not require electrons to pass ECIDS. The distribution of ANN_Δ in the top validation region is shown in Fig. 6.19d. Good agreement between data and simulation is observed, and so it may be concluded that top quark processes are well modelled.

On-shell WWW production is expected to be one of the largest backgrounds in the Z-depleted channel. A recent measurement of WWW production by ATLAS [476] measured a $WWW + WH$ signal strength² of $1.61 \pm 0.25(\text{stat.}+\text{sys.})$, corresponding to a 2.5σ excess above the SM expectation. Motivated by this result, a WWW validation region is used to assess the modelling of WWW production. Within the signal region, the WWW validation region is constructed by requiring $\text{ANN}_{WWW} > 0.5$, a cut which suppresses WH production almost completely. The distribution of the multiclassifier's WWW output score in the WWW validation region is shown in Fig. 6.20a, where three bins of equal width are plotted. An excess in data is observed in the middle bin, corresponding to a p -value of 0.04%. This excess is primarily attributed to events with $N_{\text{jets}} = 1$, as shown in Fig. 6.20b, and events with $N_{b\text{-jets}} = 0$, as shown in Fig. 6.20c. However, no other features of the excess could be reliably identified due to low statistics. Scaling the WWW event yield by 1.61 increases the p -value to 1.3%. To account for these observations, the WWW background's normalization is left floating in the combined fit.

²The signal strength floats the event yield for the signal process in the statistical fit. It constitutes the parameter-of-interest of the measurement. More details are given in Chapter 8.

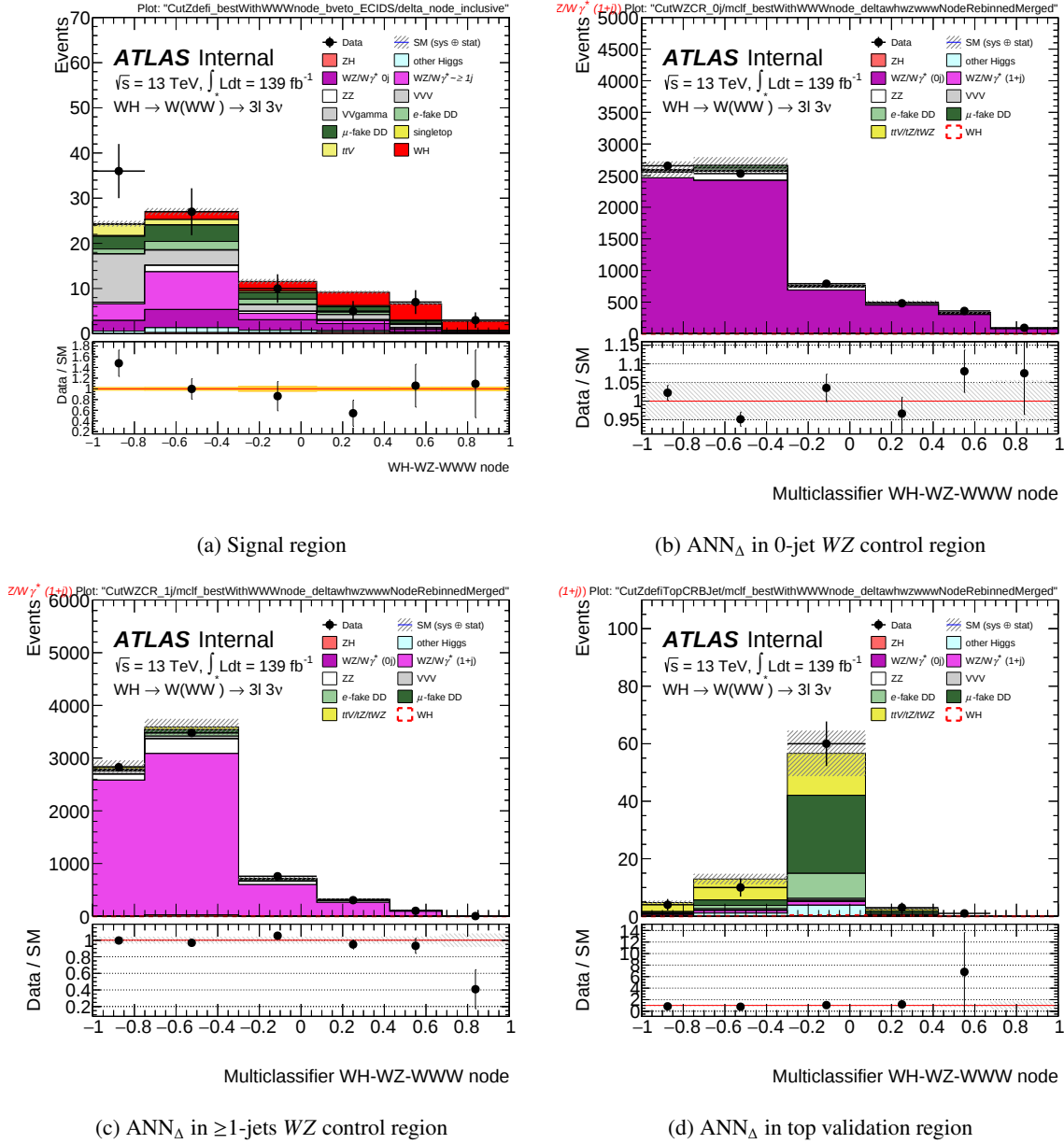


Figure 6.19: Distributions of ANN_{Δ} , Eq. 6.4, in the (a) signal region, (b) 0-jet WZ control region, (c) ≥ 1 -jets WZ control region, and (d) top validation region of the 3ℓ Z-depleted WH channel. The signal region binning is used in all distributions. Normalization factors are applied to the 0-jet and ≥ 1 -jets WZ event yields. The error bands include the statistical uncertainty as well as the dominant experimental systematic uncertainties. In (b), (c), and (d), the WH process is unstacked.

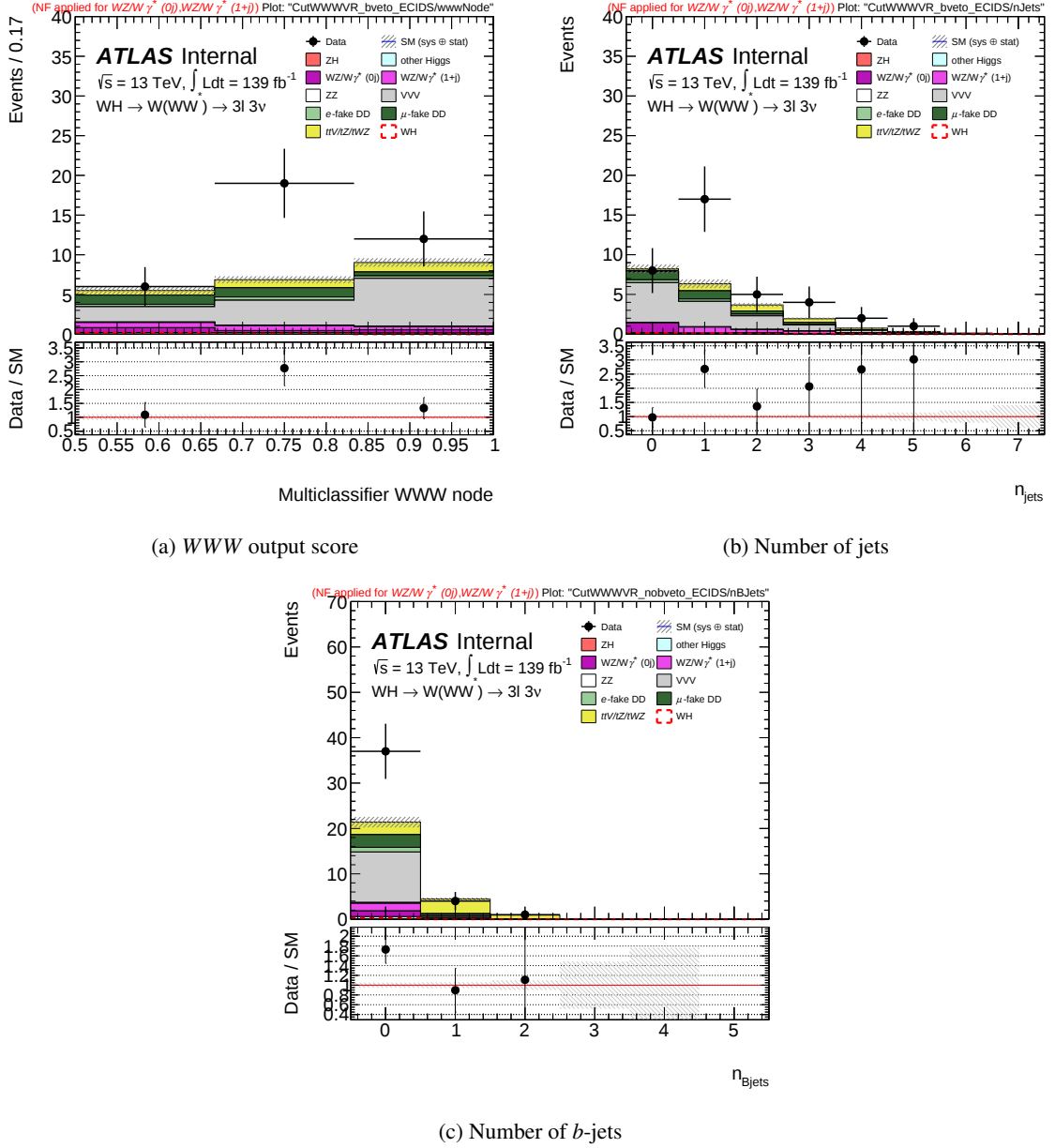


Figure 6.20: Distributions of (a) the multiclassifier's WWW output score, (b) the number of jets, and (c) the number of b -jets in the WWW validation region of the 3ℓ Z-depleted WH channel. Normalization factors are applied to the 0-jet and ≥ 1 -jets WZ event yields. The error bands include the statistical uncertainty as well as the dominant experimental systematic uncertainties. The WH process is unstacked.

6.5 4ℓ ZH Channel

The 4ℓ ZH channel is characterized by the presence of four leptons with a total null charge. Two leptons originate from the leptonic decay of the associated Z boson – the lepton pair with an invariant mass closest to the mass of the Z boson are denoted by ℓ_2 and ℓ_3 , in order of decreasing p_T . The remaining two leptons originate from the $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$ decay and are denoted by ℓ_0 and ℓ_1 , in order of decreasing p_T . The 4ℓ ZH channel is further divided into two subchannels based on the number of SFOS lepton pairs: the 1SFOS channel requires one SFOS lepton pair and the 2SFOS channel requires two SFOS lepton pairs.

All four leptons are required to exceed a transverse momentum threshold of 15 GeV and the leading lepton is required to exceed a transverse momentum threshold of 22, 25, or 27 GeV, depending the data-taking period. Selections on the minimum dilepton invariant mass of all SFOS lepton pairs, $m_{\ell\ell,\text{SFOS}}^{\min} > 12$ GeV, and the minimum dilepton invariant mass of all DFOS lepton pairs, $m_{\ell\ell,\text{DFOS}}^{\min} > 10$ GeV, are used to remove low-mass resonances, and a selection on the minimum angular separation between leptons 0 and 1, $\Delta R_{\ell_0\ell_1} > 0.2$, is used to remove overlapping leptons. A b -jet veto, $N_{b\text{-jets}} = 0$, is used to reject top quark production, which includes electroweak production of $t\bar{t}Z/tWZ$ as well as $t\bar{t}W/tZ$ production with a misidentified lepton. An upper cut on the invariant mass of the tau lepton pair calculated using leptons 0 and 1, $m_{\tau\tau} < 5$ TeV, is used to remove events with unphysically-high $m_{\tau\tau}$ values, which were found to adversely affect this channel’s multivariate analysis. For the 2SFOS channel only, a cut on the invariant mass of leptons 0 and 1, $m_{\ell_0\ell_1} < 50$ GeV, is used to suppress ZZ production. Additionally, a window cut on the invariant mass of the 4-lepton system, $|m_{4\ell} - 122.5| > 7.5$ GeV, is used to ensure orthogonality with the ATLAS $H \rightarrow ZZ^* \rightarrow 4\ell$ measurement [477].

The 4ℓ ZH channel use a multivariate discriminant known as a boosted decision tree (BDT) to extract the ZH , $H \rightarrow WW^*$ process. A BDT is trained for each of the 1SFOS and 2SFOS channels, and each BDT’s output score is optimally binned in its full range, $[-1, +1]$, within its respective signal region. The optimally-binned BDTs are provided to the combined fit.

The selections for each region of the 4ℓ ZH channel are shown in Table 6.13. A control region is used to normalize the ZZ background to data. The fake lepton background is estimated using the data-driven approach described in Chapter 5. The shape and normalization of all other backgrounds – which include $ZZq\bar{q}$ (vector boson scattering or VBS), triboson, and $t\bar{t}Z/tWZ$ production – are taken from simulation. A ZZ validation region is used to assess the modelling of ZZ production with $Z \rightarrow \tau\bar{\tau}$ in the 1SFOS channel, and a $t\bar{t}Z$ validation region based on b -jet requirement is used to assess the modelling of the $t\bar{t}Z$ background.

Table 6.13: Event selections for the signal (SR), control (CR), and validation (VR) regions of the 4ℓ ZH channel. The selections in the $t\bar{t}Z$ validation region indicated with double asterisks (**) are only applied if there are two SFOS lepton pairs.

4 ℓ ZH SR		ZZ CR	ZZ VR	$t\bar{t}$ Z VR		
1SFOS SR	2SFOS SR					
Exactly 4 ID leptons with total charge 0						
At least 1 ID lepton matched to a single lepton trigger						
All lepton $p_T > 10$ GeV with $p_T^{\ell_{\text{lead}}} > 22/25/27$ GeV						
$m_{\ell\ell,\text{SFOS}}^{\text{min}} > 12$ GeV && $m_{\ell\ell,\text{DFOS}}^{\text{min}} > 10$ GeV						
$\Delta R_{\ell_0\ell_1} > 0.2$						
$N_{b\text{-jets}} = 0$		$N_{b\text{-jets}} = 0$		$N_{\text{jets}} \geq 2$ && $N_{b\text{-jets}} \geq 1$		
$m_{\tau\tau} < 5$ TeV		–		–		
1 SFOS lepton pair	2 SFOS lepton pairs	2 SFOS lepton pairs	1 SFOS lepton pair	≥ 1 SFOS lepton pairs		
	$10 < m_{\ell_0\ell_1} < 50$ GeV	$m_{\ell_0\ell_1} > 50$ GeV	$m_{\tau\tau} > 60$ GeV	$ m_{\ell_2\ell_3} - m_Z < 10$ GeV		
	$ m_{4\ell} - 122.5 > 7.5$ GeV	$ m_{\ell_2\ell_3} - m_Z < 10$ GeV		$ m_{\ell_0\ell_1} - m_Z > 10$ GeV**		
				$m_{\ell_0\ell_1} > 40$ GeV**		

6.5.1 Multivariate Analysis

The multivariate discriminant used by the 4ℓ ZH channel is a boosted decision tree (BDT). A decision tree is trained to discriminate between two classes by recursively splitting the input events using rectangular cuts. Each split or “node” corresponds to a pass/fail decision. The decision tree terminates at a (terminal) node or “leaf”, which represents the event’s predicted class. Analogous to a neural network, a decision tree’s cuts – its free parameters – are obtained by minimizing a loss function using a dedicated set of training events. A leaf’s predicted class is assigned based on the majority of truth classes for events reaching that particular leaf.

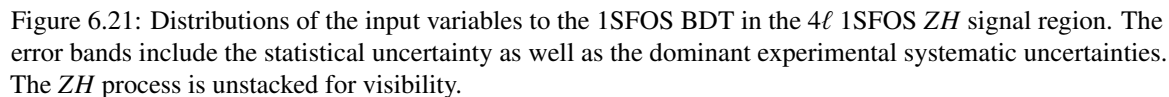
A *boosted* decision tree seeks to improve the performance of a decision tree by training a collection or “forest” of decision trees using adaptive boosting [478]. With adaptive boosting, the decision trees are trained in a sequence, and the input events for each tree are reweighted according to the misclassification rate of the previous tree such that previously misclassified events are assigned a higher weight and therefore importance [479]. The output score of the BDT is given by the weighted average of the output scores for all trees in the forest, where the weight per tree is based on its misclassification rate. Adaptive boosting improves both the stability of the classifier and typically results in improved performance over the use of a single tree [479].

For the 4ℓ ZH channel, independent BDTs, each discriminating between ZH and ZZ production, are trained for the 1SFOS and 2SFOS channels. The BDTs share a common set of input variables inspired by the signal region selections in the previous iteration of the analysis [105]. These inputs comprise kinematic and topological variables related to the 4-lepton system and the missing transverse momentum. Also included is the jet multiplicity and the azimuthal separation between leptons 0 and 1 in the Higgs candidate’s frame, $\Delta\phi_{\ell_0\ell_1}^{\text{boost}}$. This frame cannot be reconstructed due to the presence of neutrinos but is approximated using the transverse momentum of the reconstructed Z boson, $-\vec{p}_T^Z$. In the presence of at least one reconstructed jet, this frame is approximated as $-\vec{p}_T^Z - \sum \vec{p}_T^{\text{jet}}$. The input variables are listed in Table 6.14, and the inputs themselves are produced within each BDT’s associated signal region. Distributions of the input variables in the 1SFOS and 2SFOS signal regions are shown in Figs. 6.21 through 6.24.

The BDTs are trained using TMVA [479]. As the simulated samples used to train the BDTs are the same samples that the BDTs will be deployed on, a k -folding prescription is used. In general, n -way k -folding utilizes a larger training dataset per BDT for larger values of n . As there are 50% fewer ZZ training events for the 1SFOS BDT than for the 2SFOS BDT, 5-way k -folding is used to train the 1SFOS BDT and 3-way k -folding is used to train the 2SFOS BDT. Shape comparisons among the training and testing datasets for each k -fold of the 1SFOS and 2SFOS BDTs are shown in Fig. 6.25. Good agreement is observed between the shapes, and so it may be concluded that no overtraining is present.

Table 6.14: Input variables to the BDTs used by the 4ℓ ZH channel.

Variable(s)	Description
$p_T^{\ell_0}, p_T^{\ell_1}, p_T^{\ell_2}, p_T^{\ell_3}$	Transverse momentum for each lepton
$m_{\ell_0\ell_1}, m_{\ell_2\ell_3}$	Dilepton invariant masses for ℓ_0 and ℓ_1 as well as ℓ_2 and ℓ_3
$p_T^{4\ell}$	Transverse momentum of the 4-lepton system
$m_{4\ell}$	Invariant mass of the 4-lepton system
$m_{\tau\tau}$	Invariant mass of the tau lepton pair
N_{jets}	Number of jets
E_T^{miss}	Missing transverse momentum
$\Delta\phi_{\ell_0\ell_1,\text{miss}}$	Azimuthal separation between the dilepton system composed of ℓ_0 and ℓ_1 and the E_T^{miss}
$\Delta\phi_{\ell_0\ell_1}^{\text{boost}}$	Azimuthal separation between ℓ_0 and ℓ_1 in the Higgs candidate’s frame



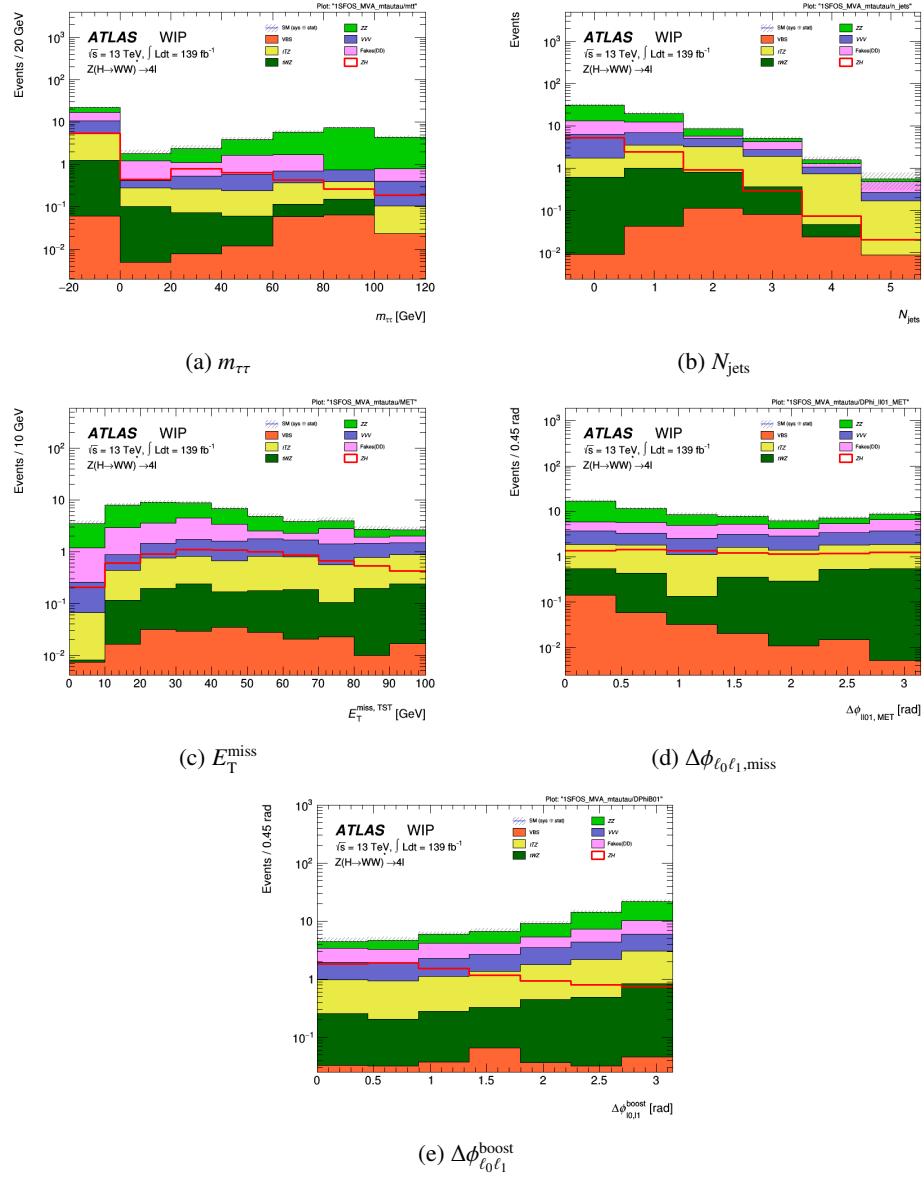


Figure 6.22: Distributions of the input variables to the 1SFOS BDT in the 4ℓ 1SFOS ZH signal region. The error bands include the statistical uncertainty as well as the dominant experimental systematic uncertainties. The ZH process is unstacked for visibility.

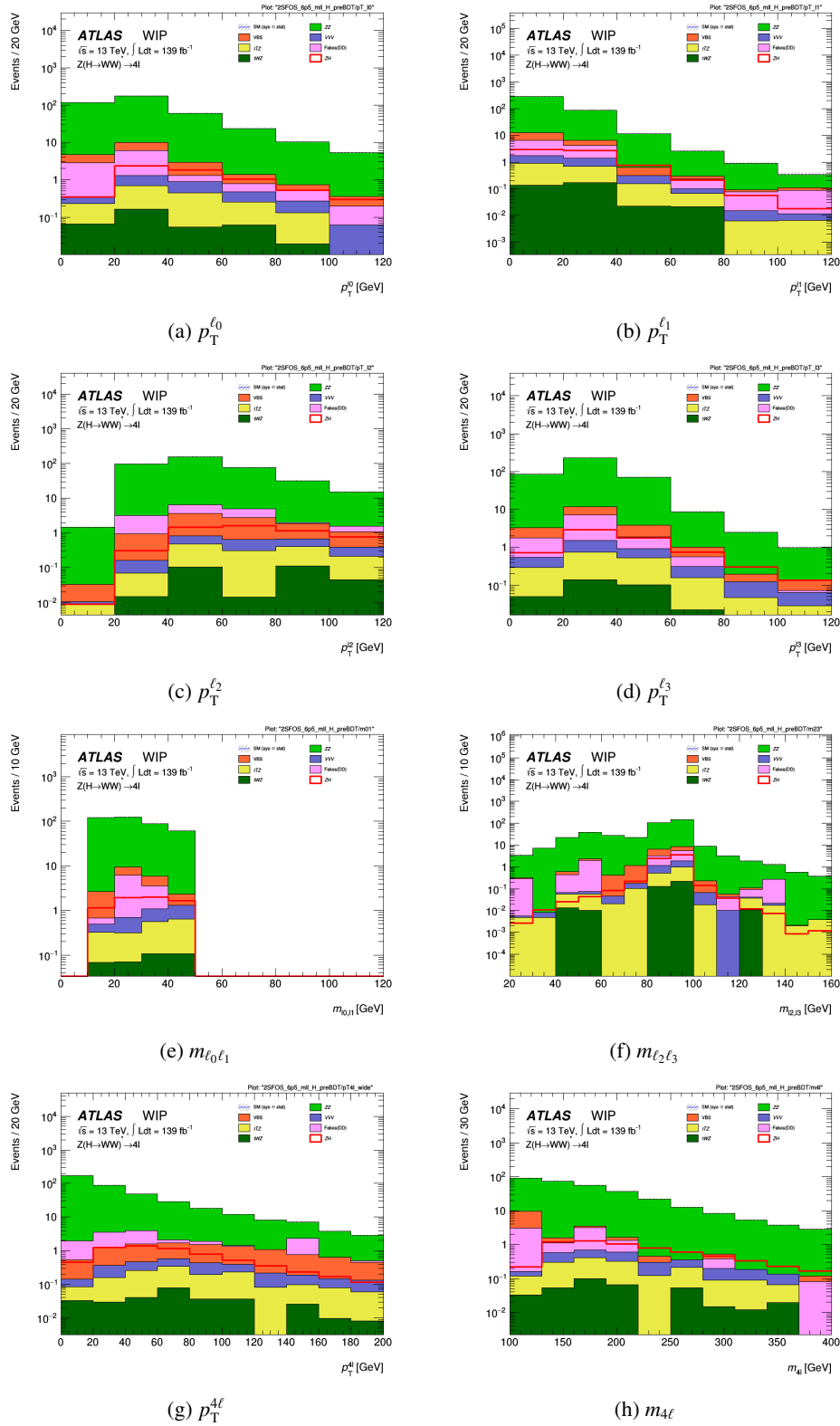


Figure 6.23: Distributions of the input variables to the 2SFOS BDT in the 4ℓ 2SFOS ZH signal region. The error bands include the statistical uncertainty as well as the dominant experimental systematic uncertainties. The ZH process is unstacked for visibility.

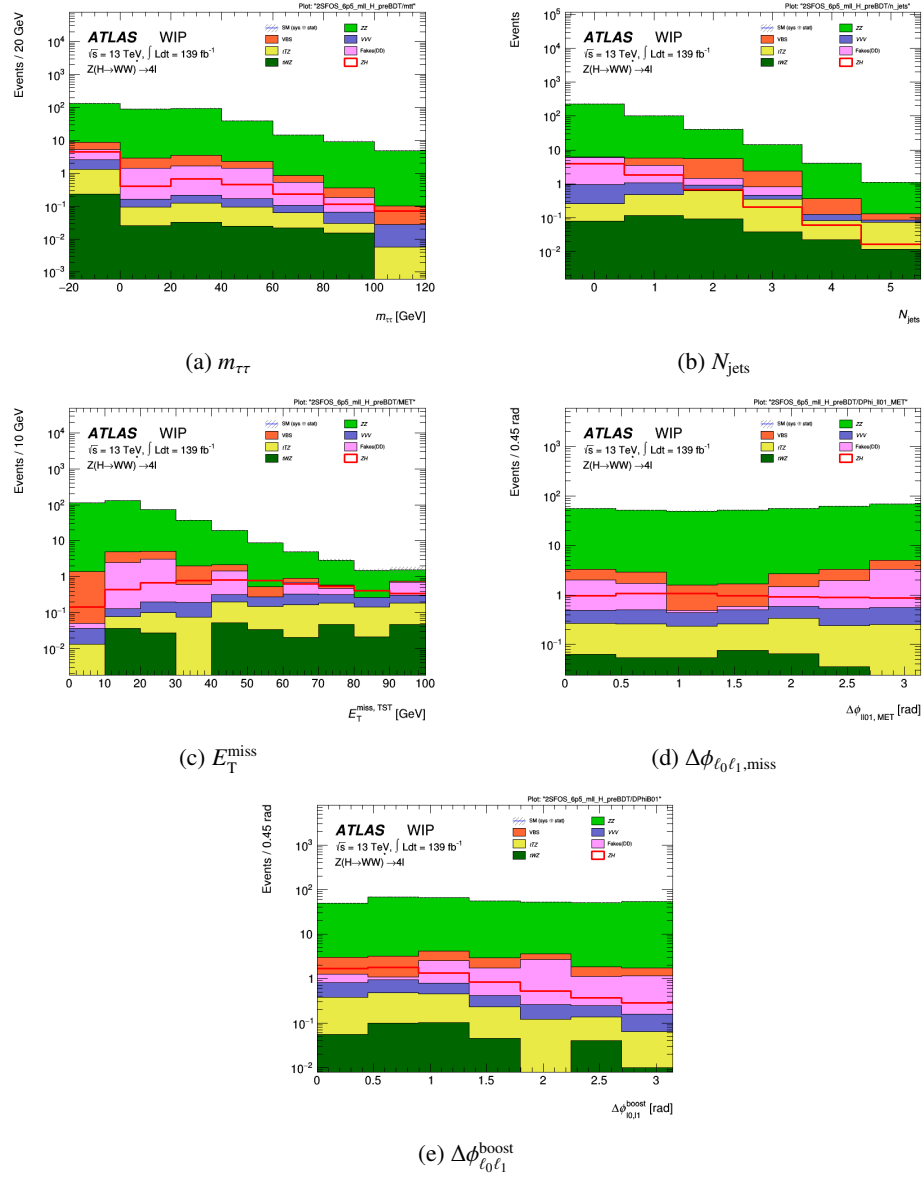


Figure 6.24: Distributions of the input variables to the 2SFOS BDT in the 4ℓ 2SFOS ZH signal region. The error bands include the statistical uncertainty as well as the dominant experimental systematic uncertainties. The ZH process is unstacked for visibility.

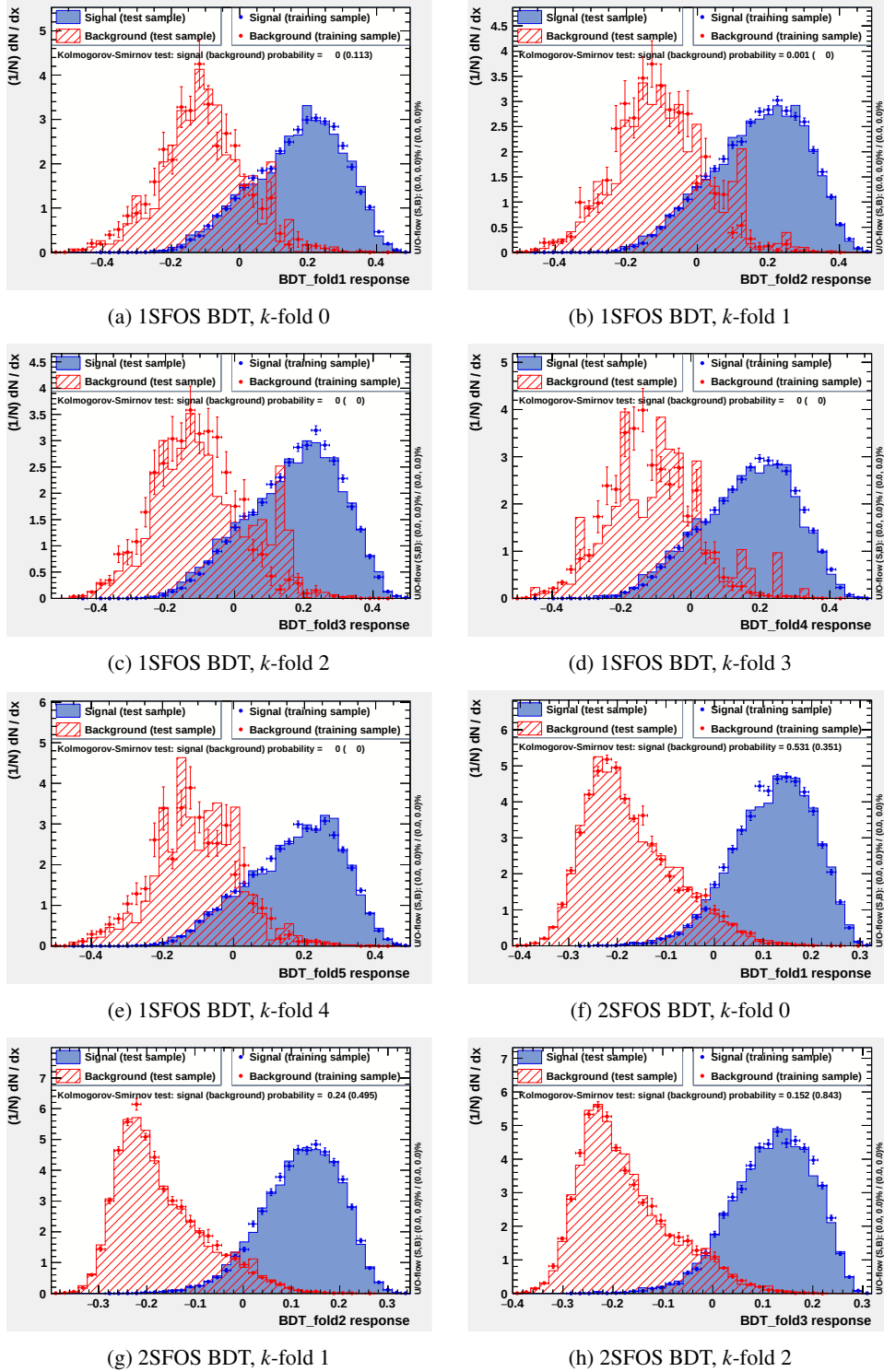


Figure 6.25: Output scores for the (a–e) 1SFOS and (f–h) 2SFOS BDTs used by the 4ℓ ZH channel. Each subplot shows a BDT trained using a different k -fold, where k -fold i corresponds to the set of events satisfying $(\text{event number}) \bmod n \neq i$ and $n = 5$ and 3 for the 1SFOS and 2SFOS BDTs, respectively. Each BDT is deployed on both the training and testing datasets to assess the level of agreement. The sum-of-weights for each class is normalized to unity. The error bars indicate the statistical uncertainty.

Distributions of the 1SFOS BDT output score in the 1SFOS signal region and the 2SFOS BDT output score in the 2SFOS signal region are shown in Fig 6.26a and 6.26b, respectively. Good agreement is observed between data and simulation.

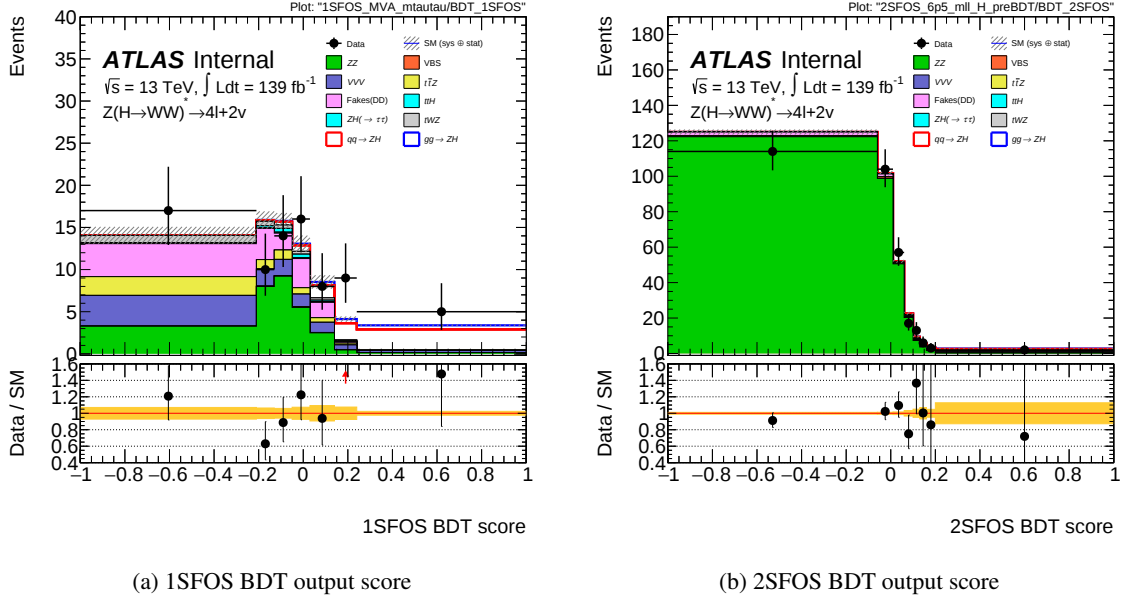
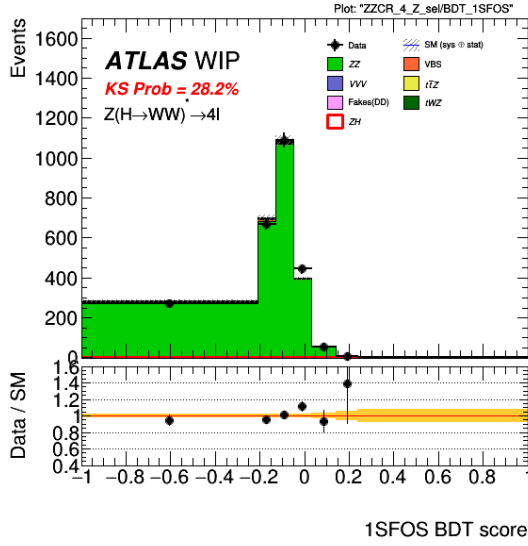


Figure 6.26: Distributions of (a) the 1SFOS BDT output score in the 1SFOS signal region and (b) the 2SFOS BDT output score in the 2SFOS signal region of the 4ℓ ZH channel. A normalization factor is applied to the ZZ event yield. The error bands include the statistical uncertainty as well as the dominant experimental systematic uncertainties.

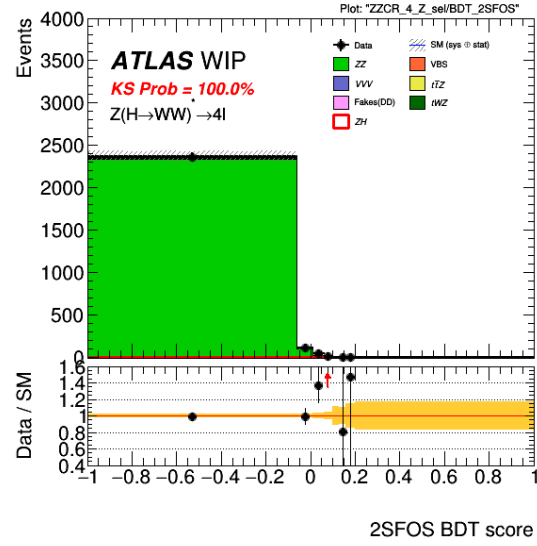
6.5.2 Control and Validation Regions

The ZZ control region is constructed by inverting the $m_{\ell_0\ell_1}$ cut of the 2SFOS signal region and requiring $m_{\ell_0\ell_1} > 50 \text{ GeV}$, which suppresses ZH production. Additionally, a window cut is applied to the mass of the Z boson candidate, $|m_{\ell_2\ell_3} - m_Z| < 10 \text{ GeV}$, to enrich the region in ZZ events. The resulting region is 98% pure in ZZ events with a corresponding normalization factor of $1.01 \pm 0.02(\text{stat.})$, which is applied to the event yields for both $q\bar{q} \rightarrow ZZ$ and $g\bar{g} \rightarrow ZZ$ production. Distributions of the 1SFOS and 2SFOS BDT output scores in the ZZ control region are shown in Figs. 6.27a and 6.27b, respectively. Good agreement between data and simulation is observed except in the 2SFOS BDT bins with edges $[0.01, 0.06]$ and $[0.06, 0.10]$. The discrepancy is largest in the latter bin: $15.1 \pm 0.7(\text{stat.})$ event were expected whereas 28 events were observed. The leading theoretical uncertainties on ZZ production in each of the 1SFOS and 2SFOS BDT bins are shown in Figs. 6.27c and 6.27d, respectively.³ In the $[0.06, 0.10]$ bin, theoretical and statistical sources in simulation result in an approximate 10% uncertainty on the expected yield. Combined with the statistical uncertainty in data, there is a total uncertainty of ± 6 events, which corresponds to an approximate 2σ excess in data over the SM expectation. Such an excess is consistent with a statistical fluctuation and therefore is not of practical concern.

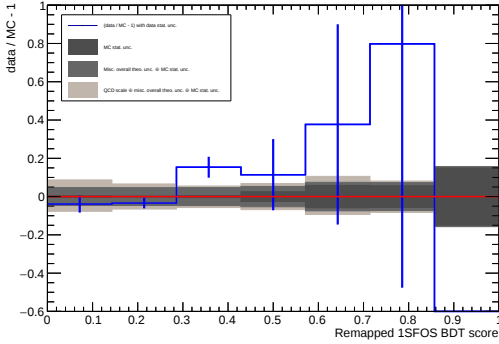
³These uncertainties include the uncertainty due to missing higher-order corrections (“QCD scales”); the uncertainties on the central values of the parton distribution function set and the strong coupling constant; and the uncertainties relating to the choice of soft gluon resummation scale, the jet merging scale, and the parton shower recoil scheme. The estimation of these uncertainties is described in Chapter 7.



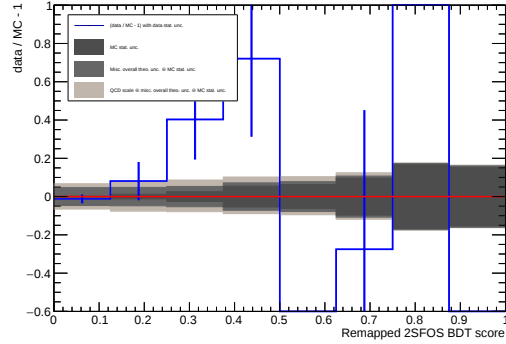
(a) 1SFOS BDT output score



(b) 2SFOS BDT output score



(c) Per-bin uncertainties for the 1SFOS BDT output score

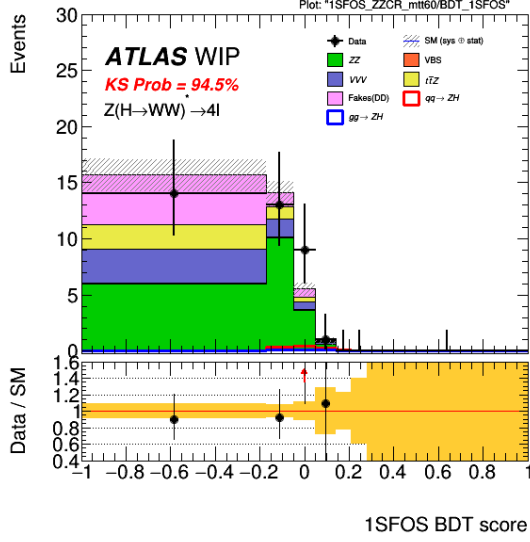


(d) Per-bin uncertainties for the 2SFOS BDT output score

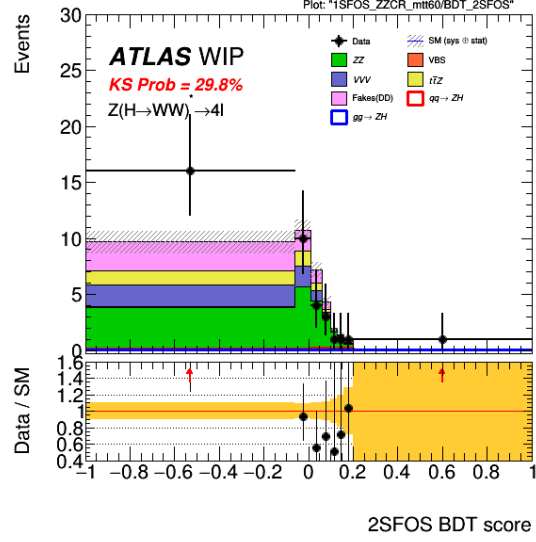
Figure 6.27: (a, b) Distributions of the (a) 1SFOS and (b) 2SFOS BDT output scores in the ZZ control region. A normalization factor is applied to the ZZ event yield. The error bands include the statistical uncertainty as well as the dominant experimental systematic uncertainties. (c, d) Per-bin uncertainties on the data-to-simulation ratio for the (c) 1SFOS and (d) 2SFOS BDT output scores in the ZZ control region. The distributions are remapped to bins of equal width. The error bars indicate the statistical uncertainty in data and the error bands indicate the combination of the statistical uncertainty in simulation and the leading theoretical uncertainties on ZZ production.

In the 1SFOS channel, ZZ production is characterized by the presence of a $Z \rightarrow \tau\bar{\tau}$ decay resulting in a DFOS lepton pair. To assess the modelling of ZZ production in the 1SFOS channel, a ZZ validation region is constructed using the same selections as the 1SFOS signal region but with an additional selection on the invariant mass of the tau lepton pair, $m_{\tau\tau} > 60$ GeV. This selection ensures the ZZ validation region is depleted in ZH events but does not make it orthogonal to the signal region; however, the ZH yield is covered by the total uncertainty on the background yield. The resulting region is 60% pure in ZZ events with a corresponding normalization factor of $1.01 \pm 0.31(\text{stat.})$ – this normalization factor is not included in the combined fit due to its large statistical uncertainty. Distributions of the 1SFOS and 2SFOS BDT output scores in the ZZ validation region are shown in Figs. 6.28a and 6.28b, respectively. Good agreement between data and simulation is observed, and so it may be concluded that ZZ production with $Z \rightarrow \tau\bar{\tau}$ is well modelled.

The $t\bar{t}Z$ validation region is used to assess the modelling of $t\bar{t}Z$ production. This region is constructed by inverting the b -jet veto of the signal region and requiring one or more jets and one or more b -jets. To select events with Z bosons, one or more SFOS lepton pairs are required and a window cut is applied to the mass of the Z boson candidate, $|m_{\ell_2\ell_3} - m_Z| < 10$ GeV. If there are two SFOS lepton pairs, a Z boson veto is applied to the mass of the Higgs boson candidate, $|m_{\ell_0\ell_1} - m_Z| > 10$ GeV, to select top quark pairs. The mass of Higgs boson candidate is additionally required to exceed 40 GeV to reject ZH production. Distributions of the 1SFOS and 2SFOS BDT output scores in the $t\bar{t}Z$ validation region are shown in Figs. 6.28c and 6.28d, respectively. Good agreement between data and simulation is observed, and so it may be concluded that $t\bar{t}Z$ production is well modelled.



(a) 1SFOS BDT output score in ZZ validation region



(b) 2SFOS BDT output score in ZZ validation region

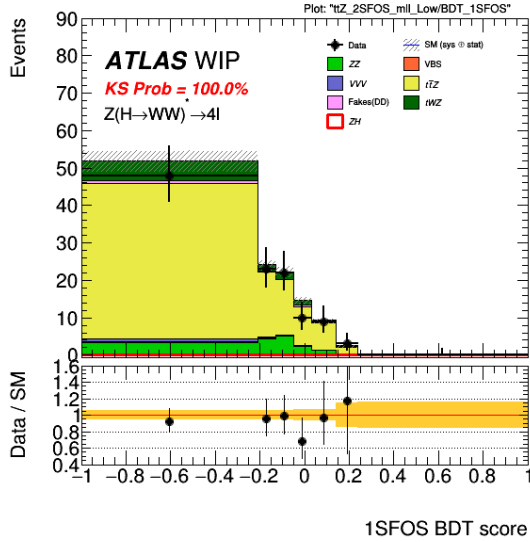
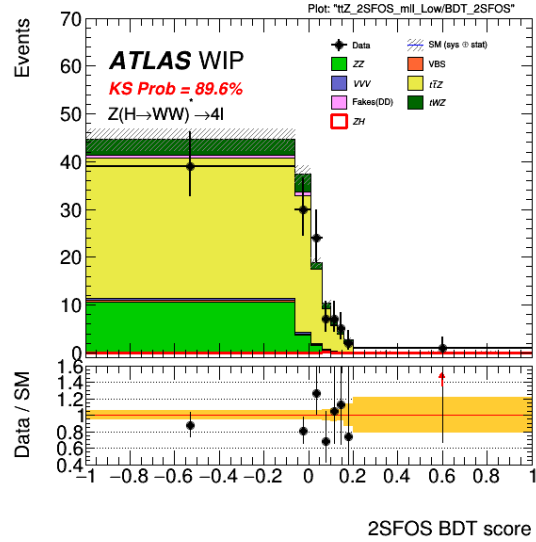
(c) 1SFOS BDT output score in $t\bar{t}Z$ validation region(d) 2SFOS BDT output score in $t\bar{t}Z$ validation region

Figure 6.28: Distributions of the 1SFOS (left) and 2SFOS (right) BDT output scores in the (a, b) ZZ and (c, d) $t\bar{t}Z$ validation regions of the 4ℓ ZH channel. A normalization factor is applied to the ZZ event yield. The error bands include the statistical uncertainty as well as the dominant experimental systematic uncertainties.

Chapter 7

Systematic Uncertainties

As with all measurements, the $VH, H \rightarrow WW^*$ measurement is affected by different sources of systematic uncertainties on the event yields in data and simulation. These uncertainties have effect of modifying the normalization and shape of each channel’s multivariate discriminant, providing additional degrees of freedom for achieving compatibility between data and simulation and reducing the sensitivity of the measurement. Therefore, it is important to enumerate and quantify the different sources of systematic uncertainty in order to provide a realistic estimate of the measurement’s precision. This chapter describes both the experimental (Section 7.1) and theoretical (Section 7.2) sources of these systematics.

7.1 Experimental Systematics

The majority of the experimental systematic uncertainties considered by the $VH, H \rightarrow WW^*$ measurement are “object-level” systematic uncertainties, which relate to the performance of specific physics objects. These include uncertainties on the simulation-to-data efficiency scale factors as well as uncertainties on the scale and resolution of the energy or momentum.

Additionally, two “event-level” systematic uncertainties are considered. The first corresponds to the uncertainty on the integrated luminosity for the Run 2 dataset, 1.7% [148], which affects the total number of predicted events. The second corresponds to an uncertainty on the central value of the data scale factor (1.03^{-1}) used during pile-up reweighting. The systematic uncertainty is assigned by varying the scale factor up (0.99^{-1}) and down (1.07^{-1}) and recomputing the pile-up reweighting for each case.

Systematic uncertainties affecting the data-driven backgrounds belong in this category as well. The estimation of these uncertainties is described in Chapter 5.

7.1.1 Efficiency Scale Factor Systematics

The systematic uncertainty on a simulation-to-data efficiency scale factor is evaluated by varying parameters of the simulated events or features of the associated tag-and-probe analysis and calculating their relative effect on the scale factor. Table 7.1 shows the typical size of the uncertainty for each scale factor applied by the $VH, H \rightarrow WW^*$ measurement.

A systematic uncertainty on a scale factor is translated into a systematic uncertainty on the event yield by computing the difference between the event yield with the nominal scale factors applied and the event

Table 7.1: Typical uncertainties on the simulation-to-data efficiency scale factors applied by the $VH, H \rightarrow WW^*$ measurement.

Efficiency scale factor	Typical uncertainty
Electron trigger	0.1% [273]
Electron reconstruction	0.4% for $E_T \approx 20$ GeV and decreasing with E_T , reaching better than 0.1% for $E_T > 30$ GeV [480]
Electron identification	7% for $E_T = 4.5$ GeV and decreasing with E_T , reaching better than 1% for $E_T > 30$ GeV [426]
Electron isolation	5% for $E_T < 7$ GeV and decreasing with E_T , reaching better than 0.5% for high E_T [426]
Muon trigger	$O(1\%)$ [250]
Muon identification	0.1% for $p_T \approx 20$ GeV and decreasing with p_T [432]
Muon TTVA	Better than 0.1% [432]
Muon isolation	$O(1\%)$ for $p_T < 20$ GeV and decreasing with p_T , reaching better than 0.1% for $p_T > 30$ GeV [432]
JVT selection	3% for jet $p_T = 25$ GeV and decreasing with p_T , reaching 1% for $p_T > 50$ GeV [450]
fJVT selection	Better than 1% [448]
b -jet tag	7% for jet $p_T < 30$ GeV and decreasing with p_T , reaching 1% for $p_T > 100$ GeV [454]
c -jet mistag	3–17% [455]
Light-jet mistag	10–50% [459]

yield with the scale factor of interest varied by its systematic uncertainty. The difference is assigned as the uncertainty.

7.1.2 Scale and Resolution Systematics

The systematic uncertainties on the scale and resolution of the energy or momentum are evaluated by varying parameters of the simulated events or features of the associated calibration analysis. For each variation, the relative effect on the energy or momentum is calculated. Systematic uncertainties on the scale and resolution of the energy or momentum are considered for electrons, muons, jets, and missing transverse momentum.

The electron energy scale uncertainty is shown in Fig. 7.1a. For $E_T < 40$ GeV, the uncertainty is dominated by the contributions due to the intercalibration of the first and second sampling layers of the electromagnetic calorimeter (indicated in the referenced figure by the “ $\alpha_{1/2}$ ” label) as well as the material between the inner detector and the presampler (“Material ID to PS”). For $E_T \approx 40$ GeV, the uncertainty is dominated by the contribution due to the η -dependent corrections applied to data (“ $Z \rightarrow ee$ calib.”). For $E_T > 40$ GeV, the uncertainty is dominated by the contribution due to the non-linearity of the calorimeter cell energy measurement (“MG/HG gain”). For the electron E_T range applicable to the $VH, H \rightarrow WW^*$ measurement, $10 < E_T \lesssim 200$ GeV, the energy scale uncertainty does not exceed 0.3% [285].

The electron energy resolution uncertainty is shown in Figs. 7.1b and 7.1c. For $10 < E_T \lesssim 200$ GeV, the uncertainty is dominated by the contribution due to the detector material in front of the calorimeter (indicated in the referenced figure by the “Material unc.” label). In this range, the energy resolution uncertainty varies from 5% at $E_T = 10$ GeV to as high as 20% at $E_T = 200$ GeV [285].

The muon momentum scale and resolution uncertainties are extracted from fits to the dimuon invariant mass in $J/\psi \rightarrow \mu\bar{\mu}$ and $Z \rightarrow \mu\bar{\mu}$ candidates [437], whose mean invariant mass and mass resolution are shown in

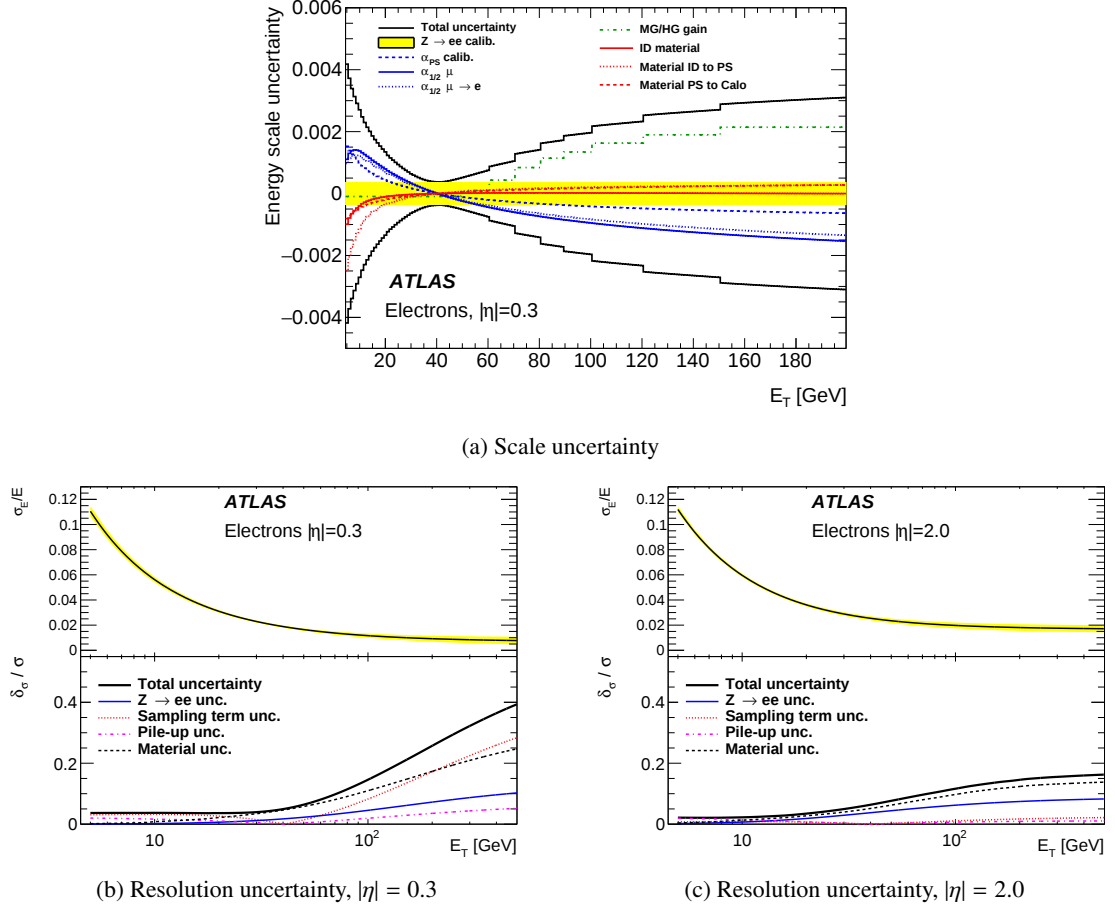


Figure 7.1: (a) Fractional energy scale calibration uncertainty for electrons as a function of E_T at $|\eta| = 0.3$ [285]. (b, c) Relative energy resolution (top) and relative energy resolution uncertainty (bottom) for electrons as a function of E_T at (b) $|\eta| = 0.3$ and (c) 2.0 [285]. In all plots, the total uncertainty as well as the individual contributions, given as one-sided variations, are shown.

Fig. 7.2. The momentum scale and resolution uncertainties are affected by contributions related to the analysis selections, the background parameterization, the momentum scale of the inner detector, and the alignment of the muon spectrometer chambers. In particular, the uncertainty on the momentum scale varies from 0.05% for $|\eta| < 1$ to a maximum of 0.3% for $|\eta| \sim 2.5$, and the uncertainty on the momentum resolution varies from 5% to 10% [437].

The jet energy scale uncertainty is shown in Fig. 7.3a. For $p_T < 30$ GeV, the uncertainty is dominated by the contribution related to the pile-up correction (indicated in the referenced figure by the “Pile-up” label) and specifically the uncertainty in the pile-up modelling in simulation. For $E_T > 30$ GeV, the uncertainty is dominated by the contribution due to the response to gluon-initiated jets (“Flav. response”). For the jet p_T range applicable to the $VH, H \rightarrow WW^*$ measurement, $20 < p_T \lesssim 300$ GeV, the energy scale uncertainty varies from 5% at $p_T = 20$ GeV to 1% at $p_T = 300$ GeV [198].

The jet energy resolution uncertainty is shown in Figs. 7.3b. For $10 < p_T \lesssim 300$ GeV, the uncertainty is dominated by the contributions related to the jet energy scale, the MC generator modelling, and the non-closure between the resolutions measured at reconstruction-level and at particle-level. In this range, the energy resolution uncertainty varies from 1.5% at $p_T = 20$ GeV to 0.5% at $p_T = 300$ GeV [198].

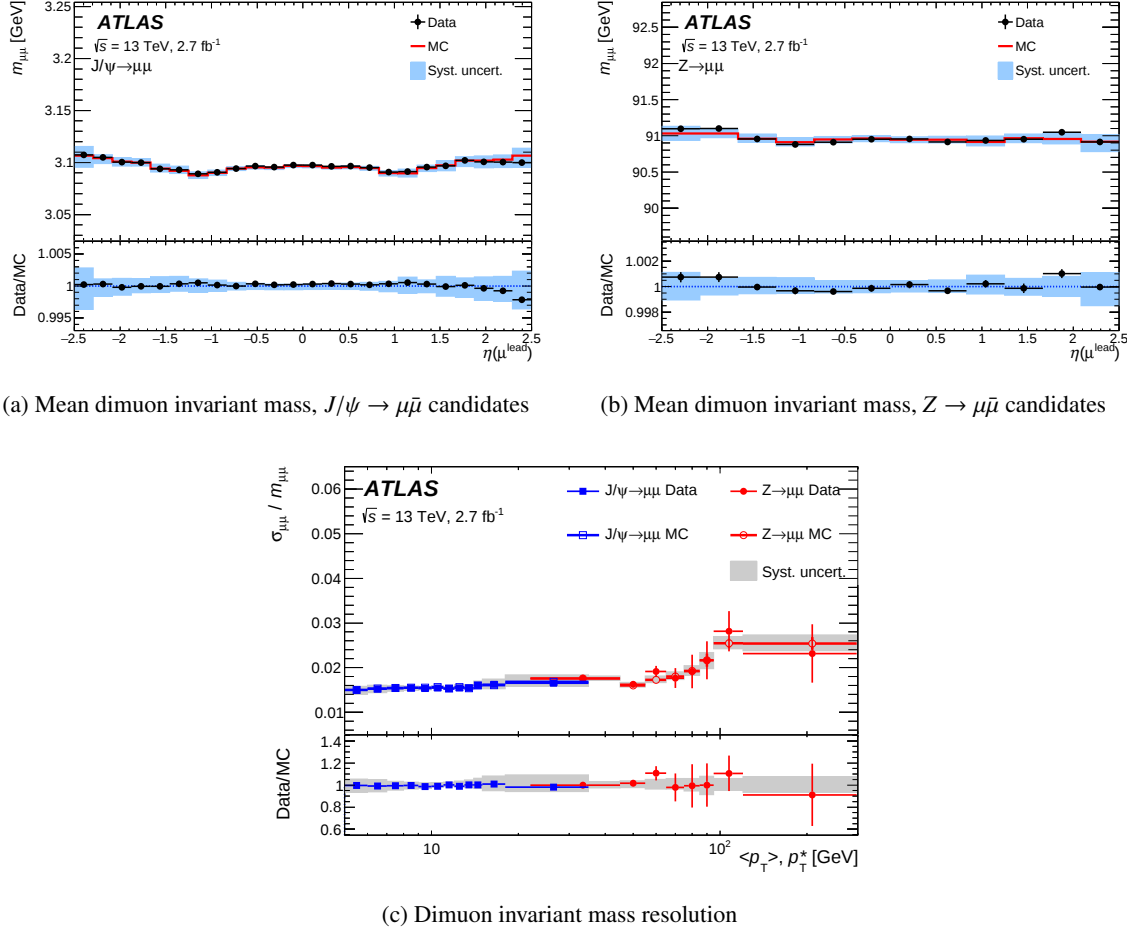


Figure 7.2: (a, b) Mean dimuon invariant mass reconstructed using combined muons from (a) $J/\psi \rightarrow \mu\bar{\mu}$ and (b) $Z \rightarrow \mu\bar{\mu}$ candidates as a function of the leading muon's η [437]. (c) Relative dimuon invariant mass resolution for $J/\psi \rightarrow \mu\bar{\mu}$ and $Z \rightarrow \mu\bar{\mu}$ candidates as a function of the dimuon p_T proxy [437]. In all plots, data and corrected simulation, along with the systematic uncertainty on the correction, are shown.

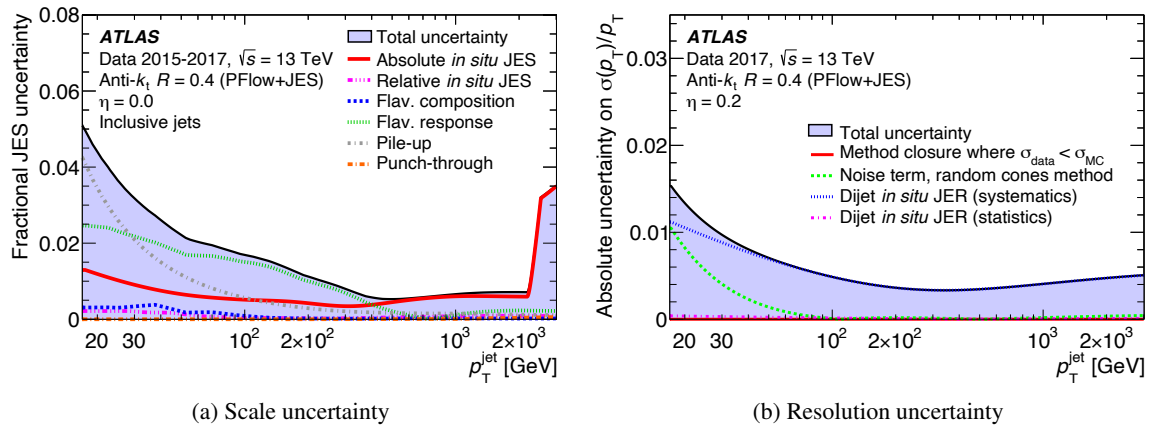


Figure 7.3: (a) Fractional energy scale calibration uncertainty for particle-flow jets as a function of jet p_T at $|\eta| = 0.0$ [198]. (b) Absolute uncertainty on the relative jet energy resolution as a function of jet p_T at $|\eta| = 0.2$ [198]. In both plots, the total uncertainty as well as the individual contributions, given as one-sided variations, are shown.

Uncertainties on the response and resolution of the missing transverse momentum are measured using 0-jet $Z \rightarrow \mu\bar{\mu}$ events, where the hard term transverse momentum, $\vec{p}_T^{\text{hard}}$, is expected to be balanced by the soft term transverse momentum, $\vec{p}_T^{\text{soft}}$ [460]. In particular, the average parallel and perpendicular projections of $\vec{p}_T^{\text{soft}}$ onto $\vec{p}_T^{\text{hard}}$, denoted by $\langle \mathcal{P}_{\parallel} \rangle$ and $\langle \mathcal{P}_{\perp} \rangle$, respectively, are computed. The energy response uncertainty is then obtained from data-to-simulation comparisons of $\langle \mathcal{P}_{\parallel} \rangle$, as shown in Fig. 7.4a, while the energy resolution uncertainty is obtained from data-to-simulation comparisons of the variances in $\mathcal{P}_{\parallel}$ and $\mathcal{P}_{\perp}$, as shown in Figs. 7.4b and 7.4c, respectively.

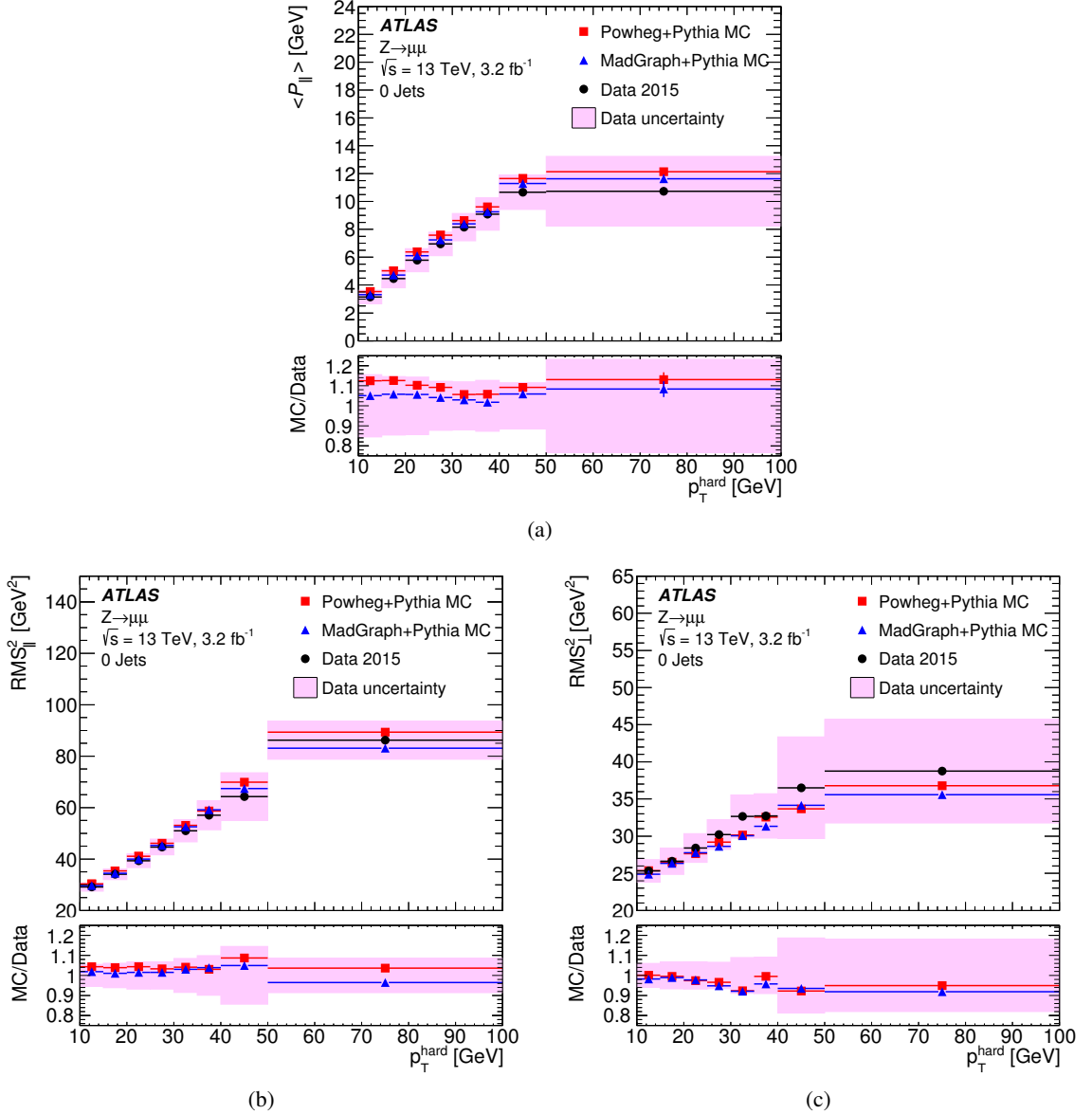


Figure 7.4: (a) Average value of the parallel projection of $\vec{p}_T^{\text{soft}}$ onto $\vec{p}_T^{\text{hard}}$ in 0-jet $Z \rightarrow \mu\bar{\mu}$ events [460]. (b, c) Variances of the (b) parallel and (c) perpendicular projections of $\vec{p}_T^{\text{soft}}$ onto $\vec{p}_T^{\text{hard}}$ in 0-jet $Z \rightarrow \mu\bar{\mu}$ events [460]. In all plots, the error bands indicate the systematic uncertainty on the E_T^{miss} energy response or resolution.

A systematic uncertainty on an energy or momentum scale is translated into a systematic uncertainty on the event yield by computing the difference between the event yield with the nominal scales and the event yield

with the scale of interest varied by its systematic uncertainty. The difference is assigned as the uncertainty. A systematic uncertainty on an energy or momentum resolution is translated into a systematic uncertainty on the event yield by computing the quadrature difference between the nominal and systematically-varied values of the resolution of interest. This difference is used to define the width of a zero-centered Gaussian, which is used to apply noise to the response in simulation, a procedure known as smearing.

7.2 Theoretical Systematics

Theoretical systematic uncertainties correspond to uncertainties on the simulation of signal and background processes. These include uncertainties due to missing higher-order corrections, uncertainties on the central values of the parton distribution function set and on the strong coupling constant, and uncertainties on the parton shower modelling.

The total theoretical uncertainty for each process and region entering the $VH, H \rightarrow WW^*$ measurement is given in Table 7.2. The following paragraphs review the prescriptions general to all processes, while the subsequent sections discuss process-specific uncertainties. Unless specified otherwise, the estimation of the uncertainties is performed using reconstruction-level samples and a full simulation of the ATLAS detector.

Missing Higher-Order Corrections: For the purposes of event generation, the partonic cross section, Eq. 3.2, is expanded in perturbation theory as:

$$\hat{\sigma} \approx \sum_{i=0}^n \alpha_S^{i+1} \hat{\sigma}_i, \quad (7.1)$$

and is truncated at a fixed order n in QCD. The effect of the missing higher-order corrections, $\mathcal{O}(\alpha_S^{n+2})$, on the partonic cross section is governed by the renormalization and factorization scales. As choice of these scales is somewhat arbitrary, the uncertainty due to missing higher-order corrections is evaluated by varying each scale up and down by a factor of 2. The following combinations are considered:

$$(\mu_R^{\text{var}}, \mu_F^{\text{var}}) = (N \times \mu_R^{\text{nom}}, M \times \mu_F^{\text{nom}}) \forall N, M \in [0.5, 1, 2], \quad (7.2)$$

where μ^{var} and μ^{nom} are the varied and nominal scales, respectively. The combinations $(N, M) = (0.5, 2)$ and $(2, 0.5)$ are excluded in order to avoid large logarithms in the cross section calculation.

The process's cross section is recomputed for each combination of scales, defining a 7-point envelope. The uncertainty due to missing higher-order corrections, referred to simply as the scale uncertainty, is given by the maximum variation of this envelope.

Parton Distribution Function: Parton distribution functions (PDFs) are affected by experimental uncertainties on the datasets which enter the PDF fits as well as on the functional forms of the fitted models. These uncertainties are provided through PDF error sets, each with a modified parameterization with respect to the nominal PDF set. When estimating the PDF uncertainty, a process's cross section must be recomputed for each error set, as the PDF affects the cross section in a nonlinear way.

For Higgs predictions, the recommended PDF set is PDF4LHC15 [301]. The associated error sets are treated as independent nuisance parameters in the measurement, with the total uncertainty given by:

$$\Delta\sigma_{\text{PDF}} = \pm \sqrt{\sum_{i=1}^{N_{\text{err}}} (\sigma_i - \sigma_0)^2}, \quad (7.3)$$

where σ_0 is the cross section evaluated using the nominal PDF set and σ_i is the cross section evaluated using the i -th error set. A total of $N_{\text{err}} = 30$ error sets exist.

For all of the background processes, including VV , VVV , Z +jets, and $t\bar{t}$, the recommended PDF set is NNPDF3.0 [300]. The total uncertainty is calculated as the standard deviation of the associated error sets:

$$\Delta\sigma_{\text{PDF}} = \pm \sqrt{\frac{1}{N_{\text{err}}} \sum_{i=1}^{N_{\text{err}}} (\sigma_i - \langle\sigma\rangle)^2}, \quad (7.4)$$

where σ_i is the cross section evaluated using the i -th error set and $\langle\sigma\rangle \equiv \sigma_0$ is the average cross section across all error sets, $\langle\sigma\rangle = N_{\text{err}}^{-1} \sum_{i=1}^{N_{\text{err}}} \sigma_i$. A total of $N_{\text{err}} = 100$ error sets exist.

Strong Coupling Constant: The strong coupling constant, α_S , is affected by experimental uncertainties on the datasets which enter the α_S fit as well as missing higher-order corrections to the renormalization group equations governing its behaviour. The effect of the uncertainty on α_S on a process's cross section σ is evaluated by recomputing the cross section for up and down variations in α_S and symmetrizing the midpoint:

$$\Delta\sigma_{\alpha_S} = \pm \frac{|\sigma(\alpha_S(m_Z^2) + \Delta\alpha_S) - \sigma(\alpha_S(m_Z^2) - \Delta\alpha_S)|}{2}, \quad (7.5)$$

where α_S is quoted at the scale of the mass of the Z boson and $\Delta\alpha_S$ is its associated uncertainty. Depending on the PDF set used, the uncertainty on α_S may vary: for PDF4LHC15, $\alpha_S(m_Z^2) = 0.1180 \pm 0.0015$ [301]; for NNPDF3.0, $\alpha_S(m_Z^2) = 0.118 \pm 0.001$ [300].

7.2.1 Higgs Boson Production

For both the VH and ggF production modes, a parton shower modelling uncertainty is evaluated by computing the relative difference in the event yield between the nominal sample, which was showered with PYTHIA 8 [290], and an alternative sample, which was showered with HERWIG 7 [293, 294]. For VH production, the difference is computed at particle-level, while for ggF production, the difference is computed at reconstruction-level. The difference is symmetrized and assigned as the uncertainty.

A $H \rightarrow WW^*$ branching ratio uncertainty of $+2.2\%$ / -2.1% is assigned to the event yields for all samples which correspond to the production of Higgs bosons decaying to pairs of W bosons [97].

As VBF production is only relevant for the 2ℓ DFOS VH channel and because it is strongly rejected by that channel's multivariate discriminant, there are no dedicated theoretical systematics on VBF production included in the measurement.

7.2.2 Multiboson Production

The SHERPA 2.2.2 [291] VV samples contain matrix elements for 0 or 1 jets at next-to-leading-order accuracy in QCD and 3 or 4 jets at leading-order in QCD, with the merging of matrix elements for different jet multiplicities performed using the MEPS@NLO prescription [317, 321, 334, 339]. A jet merging scale uncertainty is evaluated by varying the choice of the scale Q used to separate the phase space covered by the matrix element and the phase space covered by the parton shower. The scale of nominal sample, $Q = 20 \text{ GeV}$, is varied up to

30 GeV and down to 15 GeV, resulting in two alternative samples. The relative difference in the event yield between the nominal sample and each alternative sample is computed, and the larger difference is symmetrized and assigned as the uncertainty.

A resummation scale uncertainty is evaluated by varying the soft gluon resummation scale of the nominal sample up and down by a factor of 2, resulting in two alternative samples [481]. The relative difference in the event yield between the nominal sample and each alternative sample is computed, and the larger difference is symmetrized and assigned as the uncertainty.

The uncertainty due to the choice of parton shower recoil scheme, which controls the evaluation of the momentum recoil during parton showering, is evaluated by comparing the nominal sample to an alternative sample, each of which applies a different parton shower recoil scheme. In the nominal sample [380], the parton shower proceeds via the ordered splitting of “dipoles”, each consisting of a parton which actually splits (i.e., the “splitter”) and a parton which spectates (i.e., the “spectator”). The transverse momentum of the parton emitted by the splitter is compensated by the spectator such that the virtuality and four-momentum of the splitter-spectator system is invariant [482]. This strategy is known as the Catani-Seymour dipole formalism [483, 484]. In the alternative sample [482], the transverse momentum of the emitted parton is compensated by the set of all final state particles. The relative difference in the event yield between the nominal sample and the alternative sample is computed, symmetrized, and assigned as the uncertainty.

For the uncertainties covering the jet merging scale, the resummation scale, and the parton shower recoil scheme, the differences are computed at particle-level.

For SHERPA 2.2.2 VVV production, a modelling uncertainty is evaluated by computing the relative difference in the event yield between the nominal sample and an alternative sample, which was generated with MADGRAPH5_aMC@NLO [292].¹ The difference is symmetrized and assigned as the uncertainty.

For MADGRAPH5_aMC@NLO $WWqq$ production, only scale, PDF, and α_s uncertainties are considered. For $gg \rightarrow WW$ production, a conservative uncertainty of +100% / –50% is assigned to the cross section to account for the leading-order accuracy of the sample. An uncertainty of $\pm 15\%$ is similarly assigned to the $gg \rightarrow ZZ$ production cross section [386].

7.2.3 Z+jets Production

For Z+jets production, a generator modelling uncertainty is evaluated by computing the relative difference in the event yield between the nominal sample, which was generated with SHERPA 2.2.1 [291], and an alternative sample, which was generated with MADGRAPH5_aMC@NLO 2.2.2 [292].² The difference is symmetrized and assigned as the uncertainty.

7.2.4 Top Production

For $t\bar{t}$ production, a fragmentation and hadronization modelling uncertainty is evaluated by computing the relative difference in the event yield between the nominal sample, which was showered with PYTHIA 8, and an alternative sample, which was showered with HERWIG 7.04 [293, 294] using the H7UE tune [294]. The difference is symmetrized and assigned as the uncertainty.

¹The MADGRAPH5_aMC@NLO events were interfaced to PYTHIA 8 [290] using the A14 tune [346] to model the parton shower, hadronization, and underlying event.

²The MADGRAPH5_aMC@NLO events were simulated at leading-order accuracy in QCD for up to four final-state partons. The events were interfaced to PYTHIA 8.186 [353] using the A14 tune [346] to model the parton shower, hadronization, and underlying event. A next-to-next-leading-order cross section [388] was used to normalize the MC prediction.

A generation and matching modelling uncertainty is evaluated by computing the relative difference in the event yield between the nominal sample, which was generated with POWHEG Box v2 [287–289, 374], and an alternative sample, which was generated with MADGRAPH5_aMC@NLO 2.6.0 [292]. The difference is symmetrized and assigned as the uncertainty.

For the calculation of the fragmentation/hadronization and generation/matching modelling uncertainties, ATLFast-II [485], a fast simulation framework for the ATLAS detector, is used to simulate both the nominal and alternative samples.

The uncertainty due to additional radiation is evaluated by varying internal weights and computing their effect on the production cross section. For initial-state radiation (ISR), the A14 Var3c eigentune [346] is varied up and down, corresponding to a variation in the renormalization scale for the ISR shower up and down by roughly a factor of 2. For final-state radiation (FSR), the renormalization scale for the FSR shower is varied up and down by a factor of 2. In both cases, the variations have the effect of modifying the value of $\alpha_S(m_Z^2)$ for their respective showers.

For $t\bar{t}Z$ production, an uncertainty of +10% / –12% is assigned to the cross section [465].

Table 7.2: The total theoretical uncertainty per process in each signal region (SR) and control region (CR) entering the $VH, H \rightarrow WW^*$ measurement. For each process and region, the total uncertainty is obtained by summing the independent contributions in quadrature and is given as a relative variation on the event yield. An empty cell indicates that no uncertainty is given for the corresponding process and region. Electroweak or strong production of same-sign (SS) WW is denoted by the subscripts “EW” and “QCD”, respectively, and “ $N_j WZ$ ” denotes WZ production with N reconstructed jets. The total uncertainties on Higgs boson production include the $H \rightarrow WW^*$ branching ratio uncertainty.

Region	Total theoretical uncertainty per process [%]															
	WH	ZH	ggF	WW	$gg \rightarrow WW$	$SS WW_{EW}$	$SS WW_{QCD}$	$0j WZ$	$\geq 1j WZ$	ZZ	$gg \rightarrow ZZ$	WWW	WWZ	$Z+jets$	$t\bar{t}$	$t\bar{t}Z$
2ℓ DFOS VH SR	+6.0 / -6.4	+9.2 / -7.5	+27 / -20	+29 / -22	+100 / -50	–	–	–	–	–	–	–	–	+20 / -15	± 23	–
Top CR	–	–	–	+28 / -21	–	–	–	–	–	–	–	–	–	+13 / -10	± 24	–
$Z+jets$ CR	–	–	–	+27 / -21	–	–	–	–	–	–	–	–	–	+13 / -10	+21 / -20	–
WW CR	–	–	–	+28 / -21	–	–	–	–	–	–	–	–	–	+11 / -12	+25 / -19	–
$2e$ SS WH SR	± 5.2	–	–	–	–	+11 / -8	+32 / -22	–	+17 / -13	+10 / -9	–	+14 / -11	–	–	–	–
2μ SS WH SR	± 5.3	–	–	–	–	+10 / -8	+32 / -22	–	+19 / -15	+14 / -11	–	+13 / -10	–	–	–	–
$e\mu/\mu e$ SS WH SR	+5.3 / -5.2	–	–	–	–	+10 / -8	+32 / -22	–	+17 / -14	+11 / -9	–	+13 / -10	–	–	–	–
3ℓ Z-dominated WH SR	+5.7 / -5.8	–	–	–	–	–	–	+5.3 / -5.6	+20 / -16	+11 / -9	–	+16 / -13	–	–	–	–
3ℓ Z-depleted WH SR	+5.9 / -6.0	–	–	–	–	–	–	+9 / -11	+24 / -20	± 11	–	+16 / -13	–	–	–	–
$0j WZ$ CR	–	–	–	–	–	–	–	+3.4 / -3.9	–	–	–	–	–	–	–	–
$\geq 1j WZ$ CR	–	–	–	–	–	–	–	–	+20 / -15	–	–	–	–	–	–	–
4ℓ 1SFOS ZH SR	–	+8.6 / -7.3	–	–	–	–	–	–	–	+6.8 / -6.5	± 15	–	± 2.1	–	–	+10 / -12
4ℓ 2SFOS ZH SR	–	+8.2 / -6.4	–	–	–	–	–	–	–	+8.6 / -7.7	± 15	–	± 0.9	–	–	+10 / -12
ZZ CR	–	–	–	–	–	–	–	–	–	+7.4 / -7.1	± 15	–	–	–	–	–

Chapter 8

Statistical Analysis and Results

The $VH, H \rightarrow WW^*$ measurement culminates with the statistical analysis, which is roughly divided into three steps. The first step is the construction of a model describing the results. In the context of the $VH, H \rightarrow WW^*$ measurement, this model is a likelihood function. The second step is fitting the model to the measured data. The inputs to this step are the event yields in the signal and control regions, which are described in Chapter 6, as well as the systematic uncertainties affecting those yields, which are described in Chapter 7. The third and final step is the interpretation of the results. In this step, the significance of the measured results over the background-only hypothesis is calculated.

This chapter describes the statistical analysis of the $VH, H \rightarrow WW^*$ measurement. Section 8.1 provides a description of the likelihood model and significance calculation, and Section 8.2 presents the measured results.

8.1 Statistical Analysis

The statistical analysis proceeds by a frequentist approach with the construction of a likelihood function, $\mathcal{L}$, which describes the “likelihood” of measuring a particular set of parameters defining the measurement model, $\vec{\Theta}$, given the observed data. For counting experiments typical of particle physics measurements – that is, where events occur with a constant average rate – the likelihood function is written as the product of Poisson probability density functions (PDFs), one for each region or bin:

$$\mathcal{L}(\vec{\Theta} | \vec{n}) = \prod_i^{\text{regions}} P(n_i | \nu_i(\vec{\Theta})) \quad \text{where} \quad P(n | \nu) = \frac{\nu^n}{n!} \exp(-\nu). \quad (8.1)$$

Here, n_i and ν_i are the observed and expected numbers of events in region i , respectively. N.B. the expected number of events is a function of the model’s parameters. The most probable set of parameters are those which *maximize* the likelihood function given the observed data, known as the maximum likelihood estimators (MLEs). The MLE of the parameter set is implicitly defined by:

$$\left. \frac{\partial \mathcal{L}}{\partial \theta} \right|_{\vec{\Theta}=\vec{\hat{\Theta}}} = 0 \quad \forall \theta \in \vec{\Theta}. \quad (8.2)$$

As is common practice, the MLE of a particular parameter θ (or set of parameters $\vec{\Theta}$) is denoted by a hat: $\hat{\theta}$ ($\vec{\hat{\Theta}}$).

The model’s parameter set is divided into two subsets: the parameters one is interested in measuring and the parameters which are measured but may not be physically interesting. The former set of parameters are

known as “parameters-of-interest” (POIs) and correspond to a set of scale factors or “signal strengths”, each of which multiplies the event yield for an individual signal process. The latter set of parameters are known as “nuisance parameters”. In the following sections, the POIs are denoted by $\vec{\mu}$ and the nuisance parameters are denoted by $\vec{\theta}$ such that $\vec{\Theta} = \vec{\mu} \cup \vec{\theta}$.

Examples of nuisance parameters include normalization factors obtained using control regions, systematic uncertainties, and statistical uncertainties. They are encoded in the likelihood function using constraint terms, one for each nuisance parameter. A constraint term represents an auxiliary measurement which “constrains” the effect of the nuisance parameter on the event yields. With the inclusion of constraints, Eq. 8.1 becomes:

$$\mathcal{L}(\vec{\mu}, \vec{\theta} | \vec{n}) = \left[\prod_i^{\text{regions}} P(\vec{n}_i | \nu_i(\vec{\mu}, \vec{\theta})) \right] \times \left[\prod_j^{\text{nuisance parameters}} C_j(\vec{\theta}_j | \theta_j) \right], \quad (8.3)$$

where C_j is the constraint term for nuisance parameter j , θ_j , whose central value is given by $\tilde{\theta}_j$. Each constraint term is a function of one nuisance parameter only.

8.1.1 Likelihood Function for the Combined Fit

The likelihood function for the $VH, H \rightarrow WW^*$ combined fit is given by the total product of Poisson PDFs for each bin of each channel’s signal region distribution, Poisson PDFs for each control region, and constraint terms for each nuisance parameter:

$$\begin{aligned} \mathcal{L}(\vec{\mu}, \vec{\theta} | \vec{n}, \vec{m}) = & \left[\prod_{i=1}^{N_{\text{channels}}} \prod_{j=1}^{N_{\text{bins in channel } i}} P\left(n_{ij} \mid \sum_{k=1}^{N_{\text{signals}}} \mu_k s_{ijk}(\vec{\theta}) + \sum_{\ell=1}^{N_{\text{backgrounds}}} b_{ij\ell}(\vec{\theta})\right) \right] \\ & \times \left[\prod_{i'=1}^{N_{\text{control regions}}} P\left(m_{i'} \mid \sum_{k=1}^{N_{\text{signals}}} \mu_k s_{i'k}(\vec{\theta}) + \sum_{\ell=1}^{N_{\text{backgrounds}}} b_{i'\ell}(\vec{\theta})\right) \right] , \quad (8.4) \\ & \times \left[\prod_{j'=1}^{N_{\text{nuisance parameters}}} C_{j'}(\vec{\theta}_{j'} | \theta_{j'}) \right] \end{aligned}$$

where N denotes “number of”, $\vec{n}$ are the observed event yields in the signal regions, $\vec{m}$ are the observed event yields in the control regions, $s(\vec{\theta})$ corresponds to the expected signal yield in a particular region, and $b(\vec{\theta})$ corresponds to the expected background yield in a particular region. The subscripts on an event yield uniquely map it to a particular bin of a signal region distribution or control region. For the inclusive measurement of $VH, H \rightarrow WW^*$ production, there is one signal process ($N_{\text{signals}} = 1$) and one POI, μ_{VH} . For the resolved measurement of $WH, H \rightarrow WW^*$ and $ZH, H \rightarrow WW^*$ production, there are two signal processes ($N_{\text{signal}} = 2$) and two POIs, μ_{WH} and μ_{ZH} .

The expected signal and background yields are functions of the nuisance parameters and are parameterized such that their response to each nuisance parameter is factorized from their nominal rates: $s(\vec{\theta}) = s_0 \times \prod_j \nu_j(\theta_j)$. The choice of the response function, $\nu(\theta)$, and the associated constraint term depends on the type of nuisance parameter [486]. For systematic uncertainties, the constraint takes the form of a unit Gaussian; for statistical uncertainties and normalization factors obtained using control regions, the constraint takes the form of a Poisson PDF. As an intuitive example, the response function for a normalization factor obtained using a control region is the normalization factor itself, applied only to its associated background.

8.1.2 Definition of the Test Statistic

To assess the level of agreement of the observed data with the statistical model, Eq. 8.4, the profile likelihood ratio is constructed [487, 488]:

$$\lambda(\tilde{\mu}) = \frac{\mathcal{L}(\tilde{\mu}, \vec{\hat{\theta}})}{\mathcal{L}(\hat{\mu}, \vec{\hat{\theta}})}, \quad (8.5)$$

where $\tilde{\mu}$ is the *hypothesized* value of the signal strength, $\hat{\mu}$ and $\vec{\hat{\theta}}$ are the *unconditional* MLEs, and $\vec{\hat{\theta}}$ is the *conditional* MLE when μ is fixed to $\tilde{\mu}$. For simplicity, only a single POI is assumed. The value of profile likelihood ratio varies between 0 and 1, where $\lambda \sim 1$ indicates good agreement between the hypothesized value of μ and its unconditional MLE while $\lambda \ll 1$ indicates disagreement.

A test statistic is a quantity for hypothesis testing, encapsulating both the statistical model and the observed data in order to assess their compatibility. By the Neyman-Pearson lemma [489], the profile likelihood ratio corresponds to the most powerful¹ test statistic for hypothesis testing. As is common practice in particle physics, the test statistic for the $VH, H \rightarrow WW^*$ measurement is the negative log-likelihood ratio:

$$q_{\tilde{\mu}} = -2 \ln \lambda(\tilde{\mu}), \quad (8.6)$$

where $q_{\tilde{\mu}} \sim 0$ indicates good agreement between the hypothesized value of μ and its MLE and increasing $q_{\tilde{\mu}}$ indicates increasing disagreement [487, 488].

The test statistic is used to define the p -value of the measurement, which corresponds to the probability of observing equal or greater disagreement with the hypothesized value of μ :

$$p_{\tilde{\mu}} = \int_{q_{\tilde{\mu},\text{obs}}}^{\infty} f(q_{\tilde{\mu}} | \tilde{\mu}) dq_{\tilde{\mu}}, \quad (8.7)$$

where $q_{\tilde{\mu},\text{obs}}$ is the observed value of the test statistic and $f(q_{\tilde{\mu}} | \tilde{\mu})$ is the PDF of the test statistic assuming a hypothesized $\mu = \tilde{\mu}$. The p -value is typically translated into the number of standard deviations from the mean of a standard Gaussian whose integrated, one-sided tail equals such a probability:

$$Z_{\tilde{\mu}} = \Phi^{-1}(1 - p_{\tilde{\mu}}), \quad (8.8)$$

where Φ^{-1} is the inverse cumulative function of a standard Gaussian. This quantity is referred to as the *significance* or the *sensitivity* of the measurement, and a significance of $Z = n$ is referred to as an “ $n\sigma$ significance”.

The $VH, H \rightarrow WW^*$ analysis corresponds to a *discovery* measurement, where one looks to measure the presence of signal among background. The corresponding null hypothesis is that no signal is present, $\tilde{\mu} = 0$, and the alternative hypothesis is that signal is present in some nonzero amount. The analysis utilizes the test statistic defined in Eq. 8.6 with $\tilde{\mu} = 0$ and the additional constraint that $q_0 = 0$ if $\hat{\mu} < 0$. The PDF of the test statistic is obtained by sampling the PDFs governing the likelihood function many times and calculating the test statistic for each iteration (“throwing toys”) or via asymptotic formulae [488, 490], which approximate the PDF of the test statistic in the limit that the number of observations is large [491, 492].

¹The Type-II error rate, β , corresponds to the probability of accepting the null hypothesis when the alternative hypothesis is true. The power of the test is given by $1 - \beta$ [487].

8.2 Results

The statistical model and combined fit are constructed using the software packages `HistFactory` [486], `Roofit` [493], and `RooStats` [494], with minimization performed using the program `MINUIT` [495].

The expected sensitivity for a given fit is obtained using the “Asimov” dataset [488], which is constructed by fixing the data yields in each region and bin to their corresponding simulated yields. The uncertainties on a given fit’s POIs and nuisance parameters are obtained using the program `MINOS` [495]: for each parameter θ of the unconditional maximum-likelihood fit, the negative log-likelihood function is scanned in the vicinity of the parameter’s MLE, $\hat{\theta}$, and therefore the negative log-likelihood function’s minimum, $\mathcal{L}(\hat{\theta})$, with all other parameters fixed to their MLEs. The parameter’s 68% upper and lower confidence bounds, θ_{upper} and θ_{lower} , are identified as the roots of $-2(\ln \mathcal{L}(\theta) - \ln \mathcal{L}(\hat{\theta})) = 1$ subject to $\theta > \hat{\theta}$ and $\theta < \hat{\theta}$, respectively. As the negative log-likelihood function is asymmetric near its minimum, the resulting confidence bounds are also asymmetric.

8.2.1 Optimization of the Binning for Fitted Distributions

The signal region distributions provided to the $VH, H \rightarrow WW^*$ combined fit are divided into several bins with varying signal-to-background ratios. The bin edges are chosen in order to maximize the statistical sensitivity of the fit using an optimization algorithm closely following that used by the Run 1 $VH, H \rightarrow b\bar{b}$ analysis [496]. The optimization algorithm is seeded with an input distribution’s full range, $[x_{\min}, x_{\max}]$. An objective function is defined:

$$Z([x_i, x_j]) = z_s \frac{n_s([x_i, x_j])}{N_s} + z_b \frac{n_b([x_i, x_j])}{N_b}, \quad (8.9)$$

where:

- $[x_i, x_j]$ is the interval between bin edges x_i and x_j ;
- $n_s([x_i, x_j])$ and $n_b([x_i, x_j])$ are the number of signal and background events in the interval $[x_i, x_j]$, respectively;
- $N_s \equiv n_s([x_{\min}, x_{\max}])$ and $N_b \equiv n_b([x_{\min}, x_{\max}])$ are the number of signal and background events in the full range, respectively;
- z_s and z_b are parameters which tune the minimum number of signal and background events required per bin, respectively.

The optimization algorithm proceeds as follows:

1. Starting at the rightmost edge, $x_{\max}$, the range of the interval is increased from right-to-left in steps of 0.01: $[\tilde{x}, x_{\max}]$, where $(x_{\max} - \tilde{x})$ is an integer multiple of 0.01 and $\tilde{x}$ satisfies $x_{\min} \leq \tilde{x} < x_{\max}$;
2. At each step, $Z([\tilde{x}, x_{\max}])$ is calculated;
3. If $Z([x_0, x_{\max}]) > 1$ at the step $\tilde{x} = x_0$, the interval $[x_0, x_{\max}]$ is assigned as a bin;
4. Steps 1–3 are repeated, starting from leftmost edge of the previous bin, x_0 , and increasing the interval right-to-left, same as before;
5. If the leftmost edge, $x_{\min}$, is reached at any point in Steps 1–4, the interval $[x_{\min}, x_n]$ is assigned as the last bin and the algorithm terminates.

The resulting bin edges are given by:

$$(x_{\min}, x_n, x_{n-1}, \dots, x_1, x_0, x_{\max}),$$

corresponding to $n + 2$ bins. The results of the binning optimization are shown in Table 8.1.

Table 8.1: Results of the binning optimization for each channel’s signal region distribution. For each signal region, the values of the parameters z_s and z_b as well as the expected sensitivity, Z_0 , for the single-channel fit are shown. The single-channel fits use the optimized bin edges and only include statistical uncertainties. N.B. a finer step size, 0.005, was used for the binning optimization of the 3ℓ Z-depleted WH signal region distribution.

Signal region	(z_s, z_b)	Z_0	Bin edges
2ℓ DFOS VH	(8, 2)	1.29	(0.20, 0.33, 0.44, 0.53, 0.60, 0.65, 0.70, 0.75, 0.80, 1.00)
$\mu\mu$ same-sign WH	(8, 2)	1.63	(0.00, 0.21, 0.34, 0.47, 0.58, 0.68, 0.76, 0.84, 0.93, 1.00)
ee same-sign WH		0.61	
$e\mu/\mu e$ same-sign WH		1.71	
3ℓ Z-dominated WH	(8, 2)	1.62	(0.00, 0.19, 0.28, 0.38, 0.47, 0.55, 0.63, 0.72, 0.85, 1.00)
3ℓ Z-depleted WH	(5, 3)	2.92	(−1.000, −0.750, −0.300, 0.075, 0.425, 0.675, 1.000)
4ℓ 1SFOS ZH	(3, 4)	2.90	(−1.00, −0.21, −0.13, −0.05, 0.03, 0.14, 0.24, 1.00)
4ℓ 2SFOS ZH	(7, 2)	1.36	(−1.00, −0.06, 0.01, 0.06, 0.10, 0.13, 0.16, 0.20, 1.00)

8.2.2 Combined Fit

Table 8.2 shows the regions which enter the $VH, H \rightarrow WW^*$ combined fit. The signal region distributions constitute histograms with ≥ 2 bins, while the control regions constitute histograms with 1 bin or “counters”. In the combined fit, a parameter is fully correlated across all of the regions it affects with the exception of the ≥ 1 -jets WZ normalization factor, which is decorrelated between the 2ℓ and 3ℓ channels.

Table 8.2: The regions entering the $VH, H \rightarrow WW^*$ combined fit. The pre-fit distributions within these regions are referenced. Whether an input distribution constitutes a histogram or a counter is denoted by an “H” or a “C”, respectively. The POIs, μ_{WH} and μ_{ZH} , and the normalization factors (NFs) affecting the event yields are also listed. A checkmark (✓) indicates that a parameter affects the event yields in a given region. The subscript on a normalization factor indicates the process it is applied to.

Region	Figure	Type	Parameter									
			μ_{WH}	μ_{ZH}	NF _{top}	NF _{Z+jets}	NF _{WW}	NF _{2ℓ $\geq 1j$ WZ}	NF _{3ℓ $0j$ WZ}	NF _{3ℓ $\geq 1j$ WZ}	NF _{WWW}	NF _{ZZ}
2ℓ DFOS VH SR	6.4	H	✓	✓	✓	✓	✓	–	–	–	✓	–
Top CR	6.5a	C	✓	✓	✓	✓	✓	–	–	–	✓	–
Z+jets CR	6.5b	C	✓	✓	✓	✓	✓	–	–	–	✓	–
WW CR	6.5c	C	✓	✓	✓	✓	✓	–	–	–	✓	–
$\mu\mu$ same-sign WH SR	6.8a	H	✓	✓	–	–	–	✓	–	–	✓	–
ee same-sign WH SR	6.8b	H	✓	✓	–	–	–	✓	–	–	✓	–
$e\mu/\mu e$ same-sign WH SR	6.8c	H	✓	✓	–	–	–	✓	–	–	✓	–
3ℓ Z-dominated WH SR	6.14a	H	✓	✓	–	–	–	–	✓	✓	✓	–
3ℓ Z-depleted WH SR	6.19a	H	✓	✓	–	–	–	–	✓	✓	✓	–
$0j$ WZ CR	6.14b	C	✓	✓	–	–	–	–	✓	✓	✓	–
$\geq 1j$ WZ CR	6.14c	C	✓	✓	–	–	–	–	✓	✓	✓	–
4ℓ 1SFOS ZH SR	6.26a	H	–	✓	–	–	–	–	–	–	–	✓
4ℓ 2SFOS ZH SR	6.26b	H	–	✓	–	–	–	–	–	–	–	✓
ZZ CR	6.27	C	–	✓	–	–	–	–	–	–	–	✓

The post-fit expected and observed yields in each signal region are shown in Table 8.3, and post-fit distributions of the signal-sensitive discriminants in the 2ℓ and $3\ell/4\ell$ channels are shown in Figs. 8.1 and 8.2, respectively. The rates and shapes of the observed and expected event yields agree in all signal regions. The best-values of the POIs and normalization factors are shown in Table 8.4, and 2D likelihood contours of the measured cross sections are shown in Fig. 8.3.

For the combined 1-POI fit, the VH signal strength is measured to be:

$$\mu_{VH} = 0.92^{+0.21}_{-0.20}(\text{stat.})^{+0.14}_{-0.12}(\text{sys.}),$$

corresponding to a 4.6σ significance over the background-only hypothesis. The VH cross section times the $H \rightarrow WW^*$ branching ratio is measured to be:

$$\sigma_{VH} \times \mathcal{B}_{H \rightarrow WW^*} = 0.44 \pm 0.10(\text{stat.})^{+0.06}_{-0.05}(\text{sys.}) \text{ pb},$$

in agreement with the SM expectation of 0.48 ± 0.01 pb.

For the combined 2-POI fit, the WH and ZH signal strengths are measured to be:

$$\begin{aligned} \mu_{WH} &= 0.45^{+0.27}_{-0.25}(\text{stat.})^{+0.18}_{-0.15}(\text{sys.}), \\ \mu_{ZH} &= 1.64^{+0.50}_{-0.44}(\text{stat.})^{+0.23}_{-0.17}(\text{sys.}), \end{aligned} \tag{8.10}$$

corresponding to significance levels of 1.5σ and 4.6σ , respectively, over the background-only hypothesis. For each production mode, the significance is calculated from a scenario where the production mode of interest is considered absent while the other is treated as a background normalized from data. The WH/ZH cross sections times the $H \rightarrow WW^*$ branching ratio are measured to be:

$$\begin{aligned} \sigma_{WH} \times \mathcal{B}_{H \rightarrow WW^*} &= 0.13^{+0.08}_{-0.07}(\text{stat.})^{+0.05}_{-0.04}(\text{sys.}) \text{ pb}, \\ \sigma_{ZH} \times \mathcal{B}_{H \rightarrow WW^*} &= 0.31^{+0.09}_{-0.08}(\text{stat.}) \pm 0.03(\text{sys.}) \text{ pb}. \end{aligned} \tag{8.11}$$

in agreement with the SM expectations of 0.294 ± 0.009 pb and $0.190^{+0.009}_{-0.008}$ pb, respectively. The measurements of WH and ZH production agree with the SM at the 1.8σ and 1.2σ levels, respectively. The observed WH significance is smaller than the expected significance of 3.3σ due to lower-than-expected signal yields in the 2ℓ same-sign WH and 3ℓ Z-depleted WH channels. In the 2ℓ same-sign WH channel, this is attributed to the mismodelling of the WZ background, enhancing the expected event yields in the signal-sensitive bins; in the 3ℓ Z-depleted WH channel, this is attributed to a higher-than-expected WWW event yield, resulting in a WWW normalization factor of ~ 2 and reducing the WH signal strength.

The contributions to the total uncertainties on the observed values of the signal strengths are listed in Table 8.5. For all POIs, the total uncertainties are dominated by statistical uncertainties in data. For the measurements of μ_{VH} and μ_{WH} , the RNN shape uncertainty for WZ and statistical uncertainties in simulation are the leading sources of systematic uncertainty. For the measurement of μ_{WH} production, theoretical uncertainties relating to WWW and 0-jet WZ production also contribute substantially to the total uncertainty. Specifically, the modelling uncertainty for WWW production and the uncertainty due to the choice of parton shower recoil scheme for 0-jet WZ production are dominant. For the measurement of μ_{ZH} , theoretical uncertainties relating to ZH production and experimental uncertainties relating to muons are the leading

sources of systematic uncertainty. Specifically, the uncertainty due to missing higher-order corrections to the prediction for ZH production and the uncertainties on the muon isolation efficiency scale factors are dominant.

Table 8.3: Post-fit expected and observed yields in each signal region as measured by the combined 2-POI fit. The top and Z +jets processes for a given channel correspond to those with only prompt leptons. The uncertainties correspond to the total of all statistical and systematic sources. The quadrature sum of the individual sources may differ from the total uncertainty due to correlations.

Process	2ℓ DFOS VH			2ℓ same-sign WH					
	$\mu\mu$			ee			$e\mu/\mu e$		
WH	10.1	$\pm$	6.9	31	$\pm$	21	11.2	$\pm$	7.6
ZH	20.4	$\pm$	6.1	7.3	$\pm$	2.3	4.1	$\pm$	1.3
Other Higgs	58.3	$\pm$	7.3	19.0	$\pm$	2.2	9.3	$\pm$	1.1
WW	265	$\pm$	47	252	$\pm$	35	106	$\pm$	14
WZ	41.4	$\pm$	3.1	1410	$\pm$	150	841	$\pm$	94
ZZ	7.7	$\pm$	0.8	143	$\pm$	18	100.1	$\pm$	8.2
$V+\gamma$	10.3	$\pm$	6.9	—			—		
VVV	18.6	$\pm$	6.3	166	$\pm$	47	75	$\pm$	21
Z +jets	301	$\pm$	27	—			—		
Top	1016	$\pm$	64	32.9	$\pm$	2.7	16.1	$\pm$	1.4
Misidentified leptons	26.5	$\pm$	6.3	360	$\pm$	140	1110	$\pm$	100
Charge-flip electrons	—			—			921	$\pm$	67
Total	1776	$\pm$	39	2420	$\pm$	47	3198	$\pm$	51
Observed	1788	$\pm$	42	2438	$\pm$	49	3233	$\pm$	57

Process	3ℓ WH			4ℓ ZH					
	Z-dominated			Z-depleted			1SFOS		
WH	11.0	$\pm$	7.5	5.7	$\pm$	3.9	—		
ZH	4.5	$\pm$	1.4	1.73	$\pm$	0.53	14.8	$\pm$	4.4
Other Higgs	1.80	$\pm$	0.25	2.81	$\pm$	0.31	1.87	$\pm$	0.22
WZ	876	$\pm$	20	24.6	$\pm$	1.6	—		
ZZ	130	$\pm$	17	2.63	$\pm$	0.28	29.2	$\pm$	1.5
WWW	63	$\pm$	16	34.5	$\pm$	8.9	—		
WWZ	3.78	$\pm$	0.11	1.63	$\pm$	0.06	10.81	$\pm$	0.97
WZZ	0.31	$\pm$	0.01	< 0.1			0.44	$\pm$	0.03
Top	14.4	$\pm$	1.3	4.80	$\pm$	0.43	8.71	$\pm$	0.95
Misidentified leptons	102	$\pm$	14	13.3	$\pm$	3.0	14.5	$\pm$	3.2
Total	1208	$\pm$	23	91.8	$\pm$	8.0	80.3	$\pm$	5.5
Observed	1237	$\pm$	35	88	$\pm$	9	79	$\pm$	9

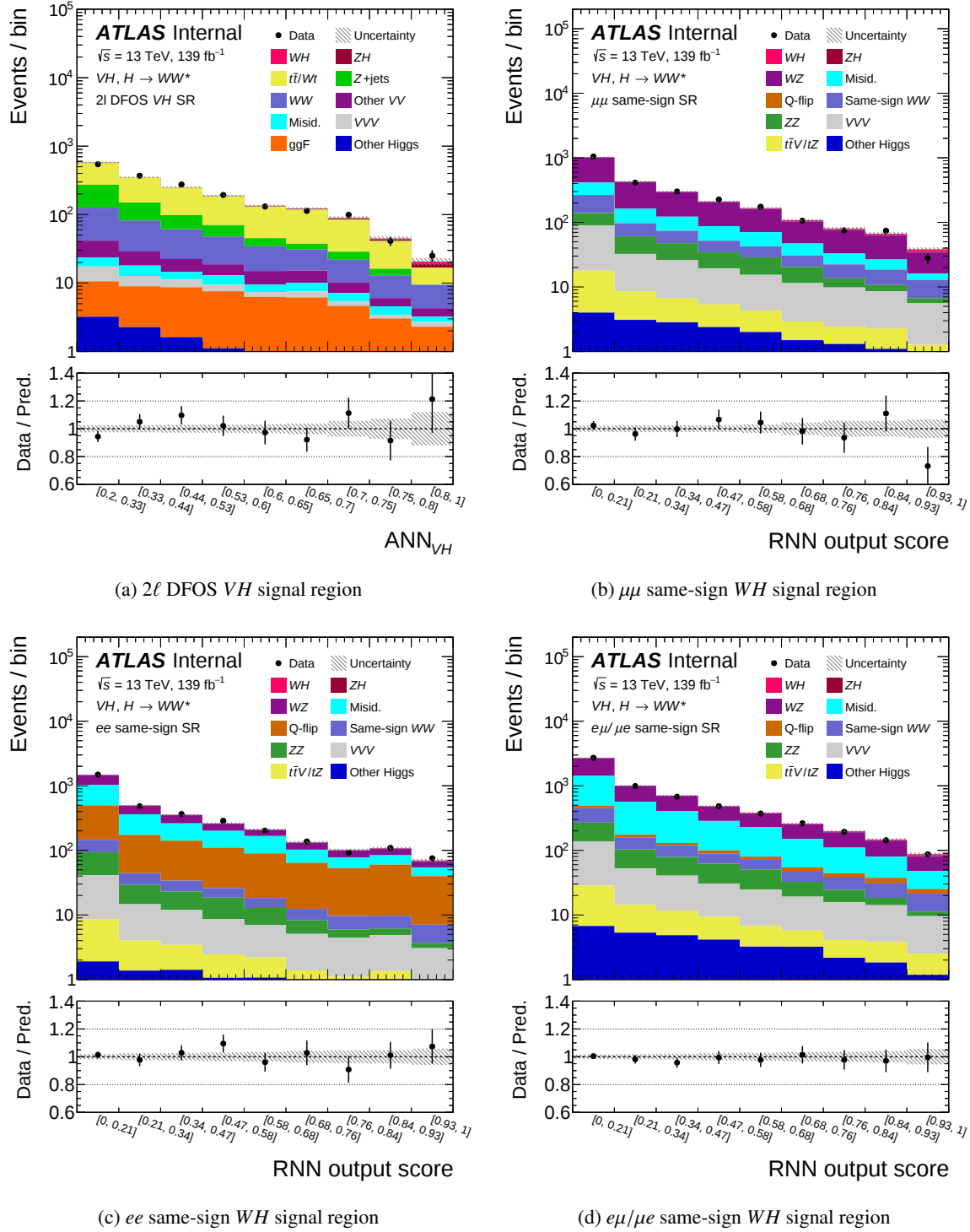


Figure 8.1: Post-fit distributions of the signal-sensitive multivariate discriminants in the (a) 2ℓ DFOS VH , (b) $\mu\mu$ same-sign WH , (c) ee same-sign WH , and (d) $e\mu/\mu e$ same-sign WH signal regions. The post-fit results are obtained from the combined 2-POI fit. The error bands include the statistical uncertainty as well as all sources of systematic uncertainties.

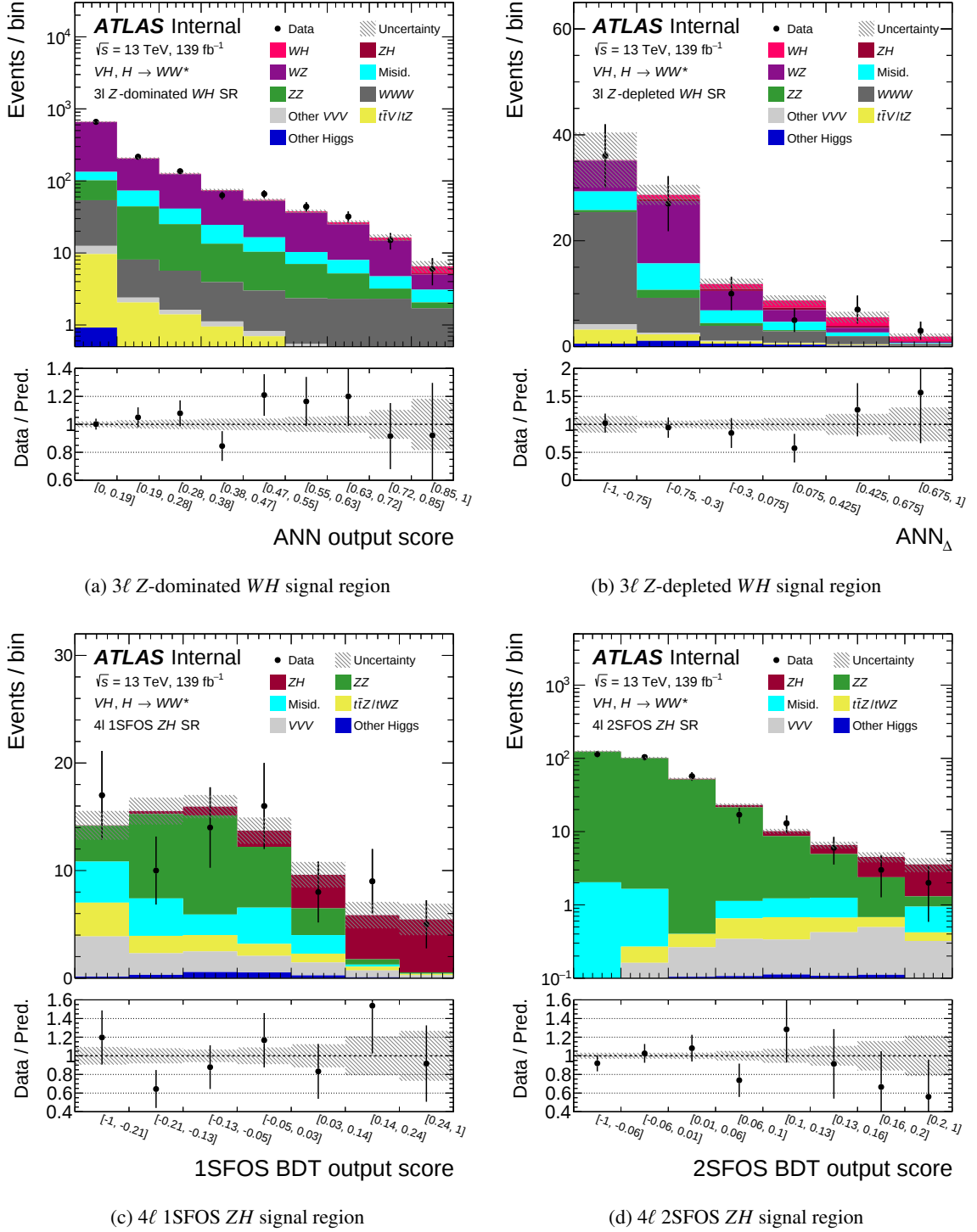


Figure 8.2: Post-fit distributions of the signal-sensitive multivariate discriminants in the (a) 3ℓ Z-dominated WH, (b) 3ℓ Z-depleted WH, (c) 4ℓ 1SFOS ZH, and (d) 4ℓ 2SFOS ZH signal regions. The post-fit results are obtained from the combined 2-POI fit. The error bands include the statistical uncertainty as well as all sources of systematic uncertainties.

Table 8.4: The results of the $VH, H \rightarrow WW^*$ combined fit. The uncertainties correspond to the total of all statistical and systematic sources. (a) The observed values of the signal strengths and the corresponding sensitivities, Z_0 , for the single-channel fits as well as the combined 1- and 2-POI fits. For each fit, both the expected and observed results are shown. (b) The observed best-fit values of the normalization factors from the 2-POI combined fit. The subscript on a normalization factor indicates the process it is applied to.

(a) Parameters-of-interest				(b) Normalization factors	
Channel	POI / Z_0	Expected	Observed		
2ℓ DFOS VH	μ_{VH}	$1.00^{+1.02}_{-0.98}$	$1.94^{+1.07}_{-1.02}$	NF_{top}	$0.99^{+0.31}_{-0.22}$
	Z_0	1.0	1.9	NF_{Z+jets}	$0.86^{+0.15}_{-0.14}$
2ℓ same-sign WH	μ_{WH}	$1.00^{+0.61}_{-0.60}$	-0.08 ± 0.58	NF_{WW}	$0.89^{+0.27}_{-0.24}$
	Z_0	1.6	0.0	$NF_{\geq 1j WZ}^{2\ell}$	$0.90^{+0.17}_{-0.16}$
3ℓ WH	μ_{WH}	$1.00^{+0.44}_{-0.40}$	$0.64^{+0.42}_{-0.37}$	$NF_{0j WZ}^{3\ell}$	1.03 ± 0.06
	Z_0	2.8	1.8	$NF_{\geq 1j WZ}^{3\ell}$	$0.86^{+0.16}_{-0.15}$
4ℓ ZH	μ_{ZH}	$1.00^{+0.47}_{-0.39}$	$1.59^{+0.54}_{-0.47}$	NF_{WWW}	$2.17^{+0.73}_{-0.60}$
	Z_0	3.1	4.5	NF_{ZZ}	$0.99^{+0.08}_{-0.07}$
Combined 1-POI	μ_{VH}	$1.00^{+0.27}_{-0.25}$	$0.92^{+0.25}_{-0.23}$		
	Z_0	4.7	4.6		
Combined 2-POI	μ_{WH}	$1.00^{+0.35}_{-0.33}$	$0.45^{+0.32}_{-0.30}$		
	μ_{ZH}	$1.00^{+0.47}_{-0.39}$	$1.64^{+0.55}_{-0.47}$		
	Z_0^{WH}	3.3	1.5		
	Z_0^{ZH}	3.1	4.6		

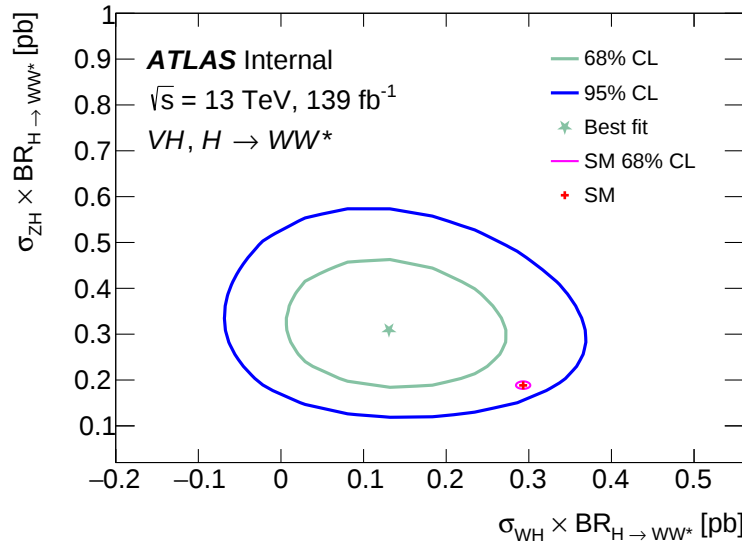


Figure 8.3: Two-dimensional likelihood contours of the observed values of $\sigma_{ZH} \times \mathcal{B}_{H \rightarrow WW^*}$ versus $\sigma_{WH} \times \mathcal{B}_{H \rightarrow WW^*}$ for the 68% and 95% confidence levels (CLs). The 68% confidence level on the SM expectations is also indicated.

Table 8.5: Breakdown of the average contributions to the total uncertainties on the observed values of the signal strengths. Indentation is used to denote subcategories. “Floating normalization” refers to the uncertainties on the normalization factors obtained using control regions. N.B. the quadrature sum of the individual sources may differ from the total uncertainty due to correlations.

Source	$\Delta(\mu_{VH})/\mu_{VH}$ [%]	$\Delta(\mu_{WH})/\mu_{WH}$ [%]	$\Delta(\mu_{ZH})/\mu_{ZH}$ [%]
Statistical uncertainties in data	22.2	57.8	28.4
Systematic uncertainties	13.9	36.9	12.0
Statistical uncertainties in simulation	6.4	14.3	5.9
Experimental systematic uncertainties	5.2	9.5	6.0
Electrons	1.2	1.8	1.6
Muons	2.5	2.8	4.1
Jet energy scale	0.6	2.1	0.4
Jet energy resolution	0.4	1.4	0.4
Flavour tagging	0.9	1.4	0.7
Missing transverse momentum	0.6	0.2	0.9
Pile-up	1.1	1.5	0.8
Luminosity	2.3	2.4	2.1
Misidentified leptons	2.9	7.1	2.7
Charge-flip electrons	1.5	4.5	< 0.1
Theoretical uncertainties	7.2	19.5	8.2
WH	3.7	5.0	0.1
ZH	1.8	1.2	7.3
Top	2.9	5.8	2.4
Z+jets	1.7	3.3	1.5
WW	1.0	3.3	0.2
0-jet WZ	3.2	11.3	0.2
≥ 1 -jets WZ	0.2	0.8	0.3
ZZ	0.8	1.4	0.6
VVV	2.4	12.6	0.3
RNN shape uncertainty for WZ	8.7	27.3	0.3
Floating normalizations	< 0.1	0.1	< 0.1
Total	26.2	68.6	30.9

Chapter 9

Conclusion

It has been ten years since a particle consistent with the SM Higgs boson was discovered by the ATLAS and CMS experiments at the LHC. Since that time, a comprehensive experimental program has measured many properties of the observed Higgs boson to accuracies better than 10% [497, 498]. This thesis is a contribution to that experimental program, reporting the measurement of associated production of Higgs bosons decaying to pairs of W bosons using the LHC Run 2 pp dataset collected by the ATLAS detector. Using 2-, 3-, and 4-lepton final states and a suite of multivariate analysis techniques, the WH and ZH production cross sections times the $H \rightarrow WW^*$ branching ratio are measured to be $0.13^{+0.08}_{-0.07}(\text{stat.})^{+0.05}_{-0.04}(\text{sys.})$ pb and $0.31^{+0.09}_{-0.08}(\text{stat.}) \pm 0.03(\text{sys.})$ pb, respectively, in agreement with the Standard Model expectations.

Looking towards the future, the High-Luminosity upgrade of the LHC (HL-LHC) [499] is expected to be operational in 2029. This upgrade will enable pp collisions with ten times the instantaneous luminosity of the LHC, resulting in ten times more data over the HL-LHC's lifetime. Along with a reduction in systematic uncertainties [500], the larger dataset will allow the ATLAS and CMS experiments to establish di-Higgs production and rare Higgs decays including $H \rightarrow \mu\bar{\mu}$, $H \rightarrow c\bar{c}$, and $H \rightarrow Z\gamma$ [498]. Even further into the future is the prospect of a lepton collider such as the International Linear Collider [501] or the Future Circular Collider [502]. With known initial states and a clean collision environment, a lepton collider will enable measurements of many Higgs properties to better than 1% and more tightly constrain potential BSM models [503]. Lepton colliders also offer an interesting opportunity for studying flavour tagging beyond bottom or charm quark tagging – with improvements to particle identification, strange quark tagging may be possible [3, 504, 505].

At start of this thesis, it was noted that the Standard Model is an incomplete picture of the universe – measurements of the Higgs boson offer a portal for exploring some of its shortcomings. The remarkable success of the Standard Model is also its curse, but it is a curse that the high-energy particle physics community is intent on breaking.

Appendix A

Four-Momentum Parameterization at Hadron Colliders

While the four-momentum of a particle is most intuitively represented by (E, p_x, p_y, p_z) , the common parameterization adopted at hadron colliders is (E, p_T, η, ϕ) , where:

Energy E ,

Transverse momentum $p_T = \sqrt{p_x^2 + p_y^2}$,

Pseudorapidity $\eta = \frac{1}{2} \ln \left(\frac{|\vec{p}| + p_z}{|\vec{p}| - p_z} \right) = -\ln \left(\tan \left(\frac{\theta}{2} \right) \right)$, (A.1)

Azimuthal angle $\phi = \arctan \left(\frac{p_y}{p_x} \right)$.

The choice of p_T is sensible as the longitudinal momentum p_z is not known *a priori* at hadron colliders and so only the transverse momentum may be constrained: $\sum p_T = 0$. Additionally, pseudorapidity has the nice property of being approximately equal to the rapidity y for $E \gg m$ (i.e., $E \approx |\vec{p}|$):

$$\eta \equiv \frac{1}{2} \ln \left(\frac{|\vec{p}| + p_z}{|\vec{p}| - p_z} \right) \stackrel{E \gg m}{\approx} \frac{1}{2} \ln \left(\frac{E + p_z}{E - p_z} \right) \equiv y. \quad (\text{A.2})$$

As differences in rapidity are Lorentz invariant, differences in pseudorapidity are approximately Lorentz invariant. For completeness, one has the reverse transformations:

$$\begin{aligned} p_x &= p_T \cos \phi, \\ p_y &= p_T \sin \phi, \\ p_z &= p_T \sinh \eta, \\ |\vec{p}| &= p_T \cosh \eta. \end{aligned} \quad (\text{A.3})$$

Appendix B

Additional Distributions from Chapter 6

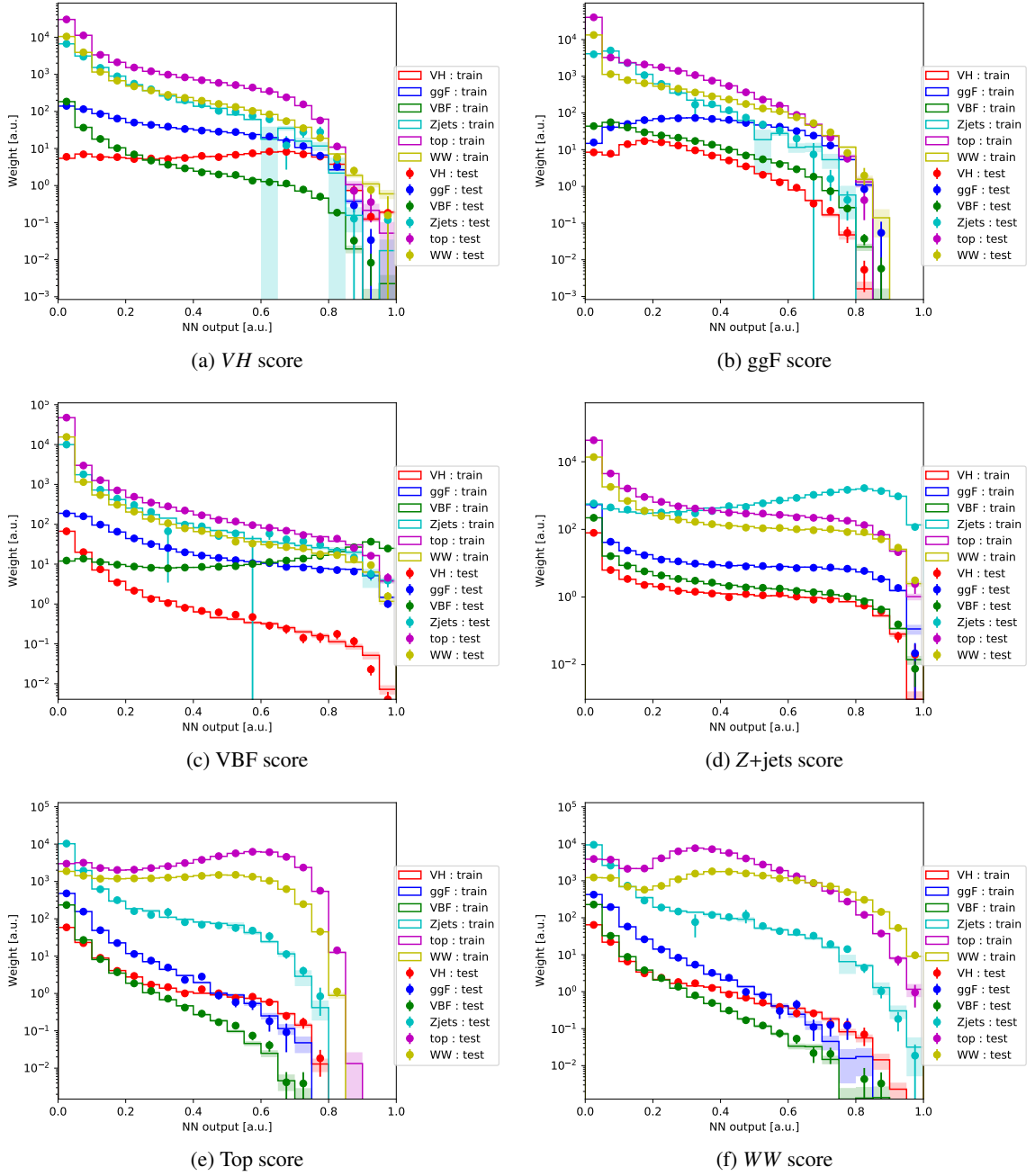


Figure B.1: Output scores for the multiclassifier used by the 2ℓ DFOS *VH* channel. The multiclassifier is trained using events satisfying $(\text{event number}) \bmod 5 \neq 1$ (i.e., “ k -fold 1”), and it is deployed on both the training and testing datasets to assess the level of agreement. The error bands/bars indicate the statistical uncertainty.

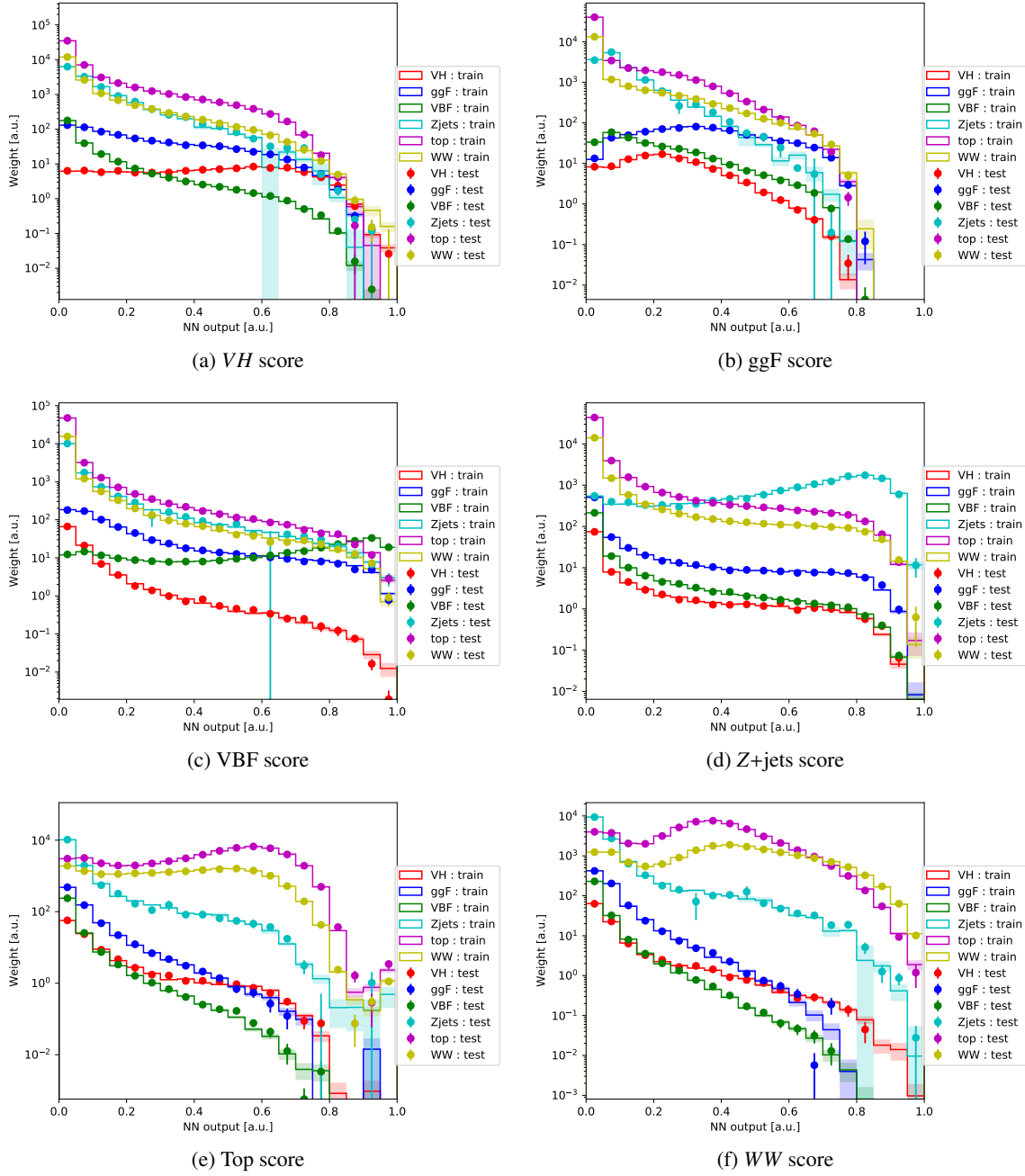


Figure B.2: Output scores for the multiclassifier used by the 2ℓ DFOS VH channel. The multiclassifier is trained using events satisfying $(\text{event number}) \bmod 5 \neq 2$ (i.e., “ k -fold 2”), and it is deployed on both the training and testing datasets to assess the level of agreement. The error bands/bars indicate the statistical uncertainty.

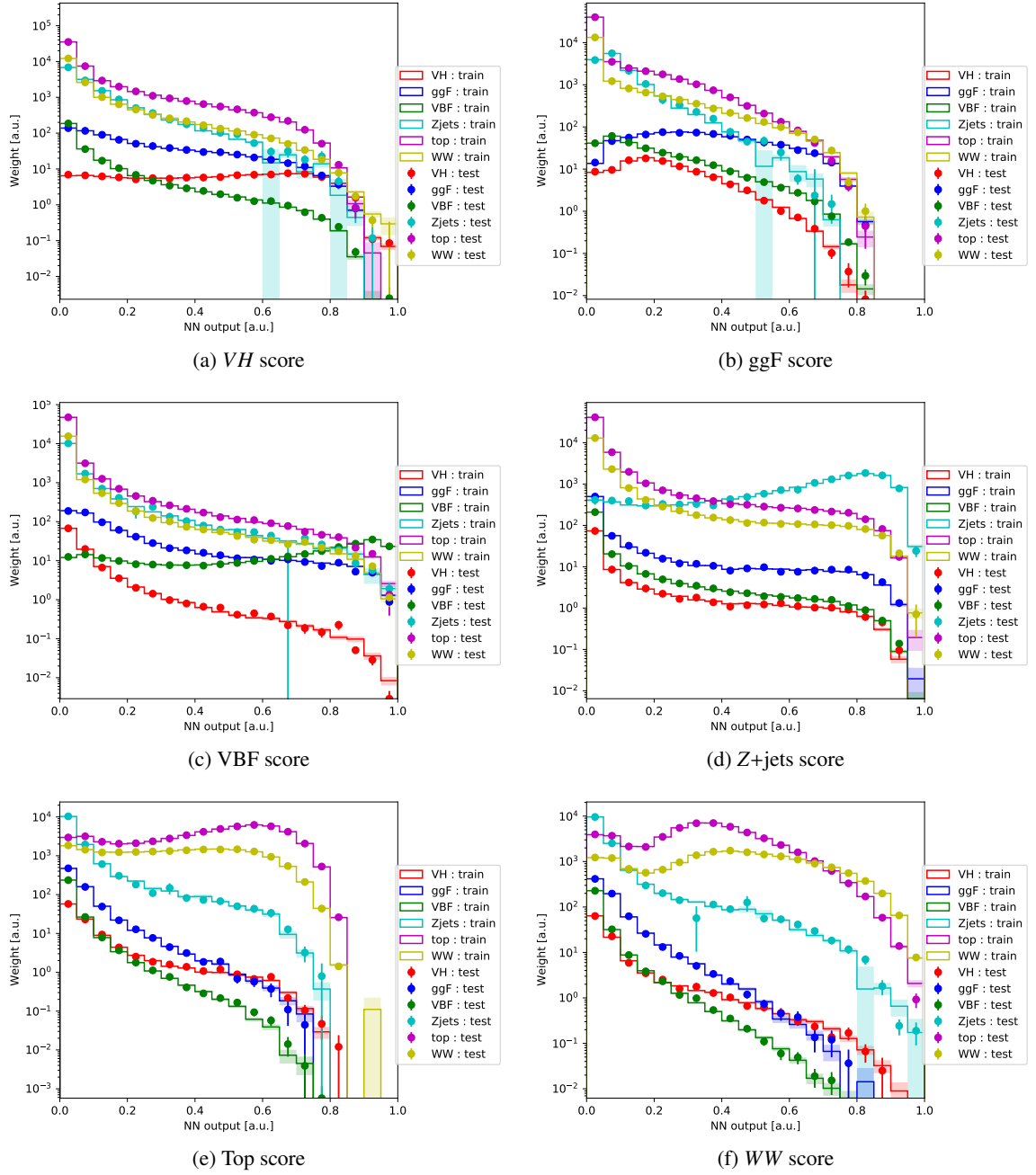


Figure B.3: Output scores for the multiclassifier used by the 2ℓ DFOS VH channel. The multiclassifier is trained using events satisfying $(\text{event number}) \bmod 5 \neq 3$ (i.e., “ k -fold 3”), and it is deployed on both the training and testing datasets to assess the level of agreement. The error bands/bars indicate the statistical uncertainty.

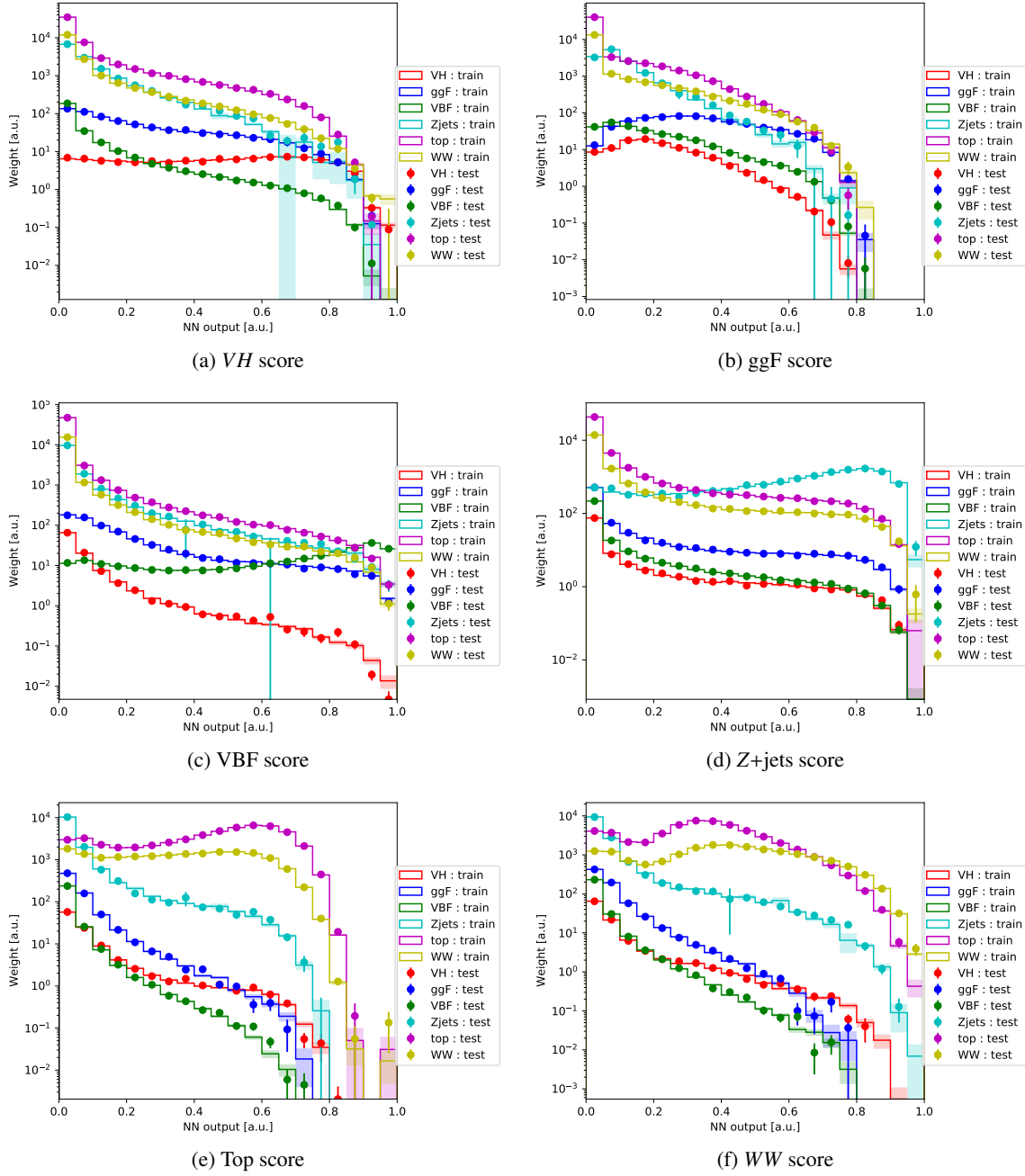


Figure B.4: Output scores for the multiclassifier used by the 2ℓ DFOS *VH* channel. The multiclassifier is trained using events satisfying (event number) mod 5 $\neq$ 4 (i.e., “*k*-fold 4”), and it is deployed on both the training and testing datasets to assess the level of agreement. The error bands/bars indicate the statistical uncertainty.

Glossary

ALEPH Apparatus for LEP Physics.

ALICE A Large Ion Collider Experiment.

ALPHA Antihydrogen Laser Physics Apparatus.

ANN Artificial neural network.

ATLAS A Toroidal LHC Apparatus.

BDT Boosted decision tree.

BSM Beyond the Standard Model.

CERN European Organization for Nuclear Research.

CKKW Catani-Kuhn-Krauss-Webber (method).

CKM Cabibbo-Kobayashi-Maskawa (matrix).

CL Confidence level.

CMS Compact Muon Solenoid.

CP Charge-parity (symmetry/violation); Cluster Processor.

CR Control region.

CSC Cathode Strip Chamber.

CTP Central Trigger Processor.

DFOS Different-flavour, opposite-sign (lepton pair).

DGLAP Dokshitzer-Gribov-Lipatov-Altarelli-Parisi (equations).

ECIDS Electron charge identification selection.

EMEC Electromagnetic endcap (calorimeter).

EW Electroweak.

FCal Forward calorimeter.

FE Front-end (chip).

FF Fake factor.

fJVT Forward Jet Vertex Tagger.

FSR Final-state radiation.

FTK Fast TracKer.

ggF Gluon-gluon fusion.

HEC Hadronic endcap (calorimeter).

HL-LHC High-Luminosity LHC.

HLT High-Level Trigger.

HOM Higher-order mode.

IBL Insertable B-Layer.

ID Identified or identification.

ISR Initial-state radiation.

ITk Inner Tracker.

IVF Iterative Vertex Finder.

JEP Jet/Energy-sum Processor.

JER Jet energy resolution.

JES Jet energy scale.

JVF Jet Vertex Fraction.

JVT Jet Vertex Tagger.

L1 Level-1 (trigger).

L1A L1-Accept.

L1Calo L1 calorimeter (trigger).

L1Muon L1 muon (trigger).

L1Topo L1 topological (trigger).

LAr Liquid argon.

LEP Large Electron-Positron (Collider).

LH Left-handed.

LHC Large Hadron Collider.

LHCb LHC beauty.

LO Leading-order (accuracy).

MC Monte Carlo.

MCC Module control chip.

MDT Monitored Drift Tube.

MLE Maximum likelihood estimator.

N³LO Next-to-next-to-next-to-leading-order (accuracy).

NF Normalization factor.

NLL Next-to-leading-logarithm (accuracy).

NLO Next-to-leading-order (accuracy).

NNLL Next-to-next-to-leading-logarithm (accuracy).

NNLO Next-to-next-to-leading-order (accuracy).

NTC Negative temperature coefficient (thermistor).

OS Opposite-sign (lepton pair).

PDF Parton distribution function; probability density function.

PLIV Prompt Lepton Improved Veto.

POI Parameter-of-interest.

PS Proton Synchrotron; parton shower.

PSB Proton Synchrotron Booster.

QCD Quantum chromodynamics.

QED Quantum electrodynamics.

QFT Quantum field theory.

RF Radio frequency.

RH Right-handed.

RNN Recurrent neural network.

ROD ReadOut Driver.

RoI Region-of-Interest.

ROS ReadOut System.

RPC Resistive Plate Chamber.

SCT SemiConductor Tracker.

SF Scale factor.

SFOS Same-flavour, opposite-sign (lepton pair).

SM Standard Model.

SNO Sudbury Neutrino Observatory.

SPS Super Proton Synchrotron.

SR Signal region.

SS Same-sign (lepton pair).

STXS Simplified Template Cross Section.

TDAQ Trigger and data acquisition.

TGC Thin Gap Chamber.

TRT Transition Radiation Tracker.

TTVA Track-to-vertex association.

UE Underlying event.

VBF Vector boson fusion.

VBS Vector boson scattering.

VR Validation region.

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