

Search for Supersymmetry and Large Extra Dimensions with the ATLAS Experiment

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Abstract

The Large Hadron Collider is the most powerful particle accelerator built to date. It is a proton-proton and heavy ion collider which in 2015 and 2016 operated at an unprecedented center of mass energy of $\sqrt{s} = 13$ TeV. The Tile Calorimeter is the ATLAS hadronic calorimeter covering the central region of the detector. It is designed to measure hadrons, jets, tau particles and missing energy. In order to accurately be able to properly reconstruct these physical objects a careful description of the electronic noise is required. This thesis presents the work done in updating, monitoring and studying the noise calibration constants used in the processing and identification of hadronic jet in the 2011 data.

Moreover the results of the searches for compressed supersymmetric squark-neutralino and large extra dimensions models are also presented in this thesis. The present work uses an experimental signature with a high energy hadronic jet and large missing transverse energy later often referred to as *monojet* signature. The search for supersymmetry is carried out using an integrated luminosity of 3.2 fb^{-1} recorded by the ATLAS experiment in 2015.

The search for large extra dimensions presented in this work uses the full 2015 + 2016 dataset of 36.1 fb^{-1} . No significant excess compared to the Standard Model prediction has been observed on the production of squark pairs with the subsequent decay of the squark in a quark and a neutrino. Exclusion limits are set on squark production as a function of the neutralino mass. Squark masses up to 608 GeV are excluded for a mass difference between the squark and the neutralino of 5 GeV. In the second search for the presence of large extra spatial dimensions in the Arkani-Hamed, Dimopoulos and Dvali model scenario a good agreement between data and Standard Model prediction is observed and exclusion limits are set on the effective Planck scale M_D of 7.7 and 4.8 TeV for two and six hypothesized large extra dimensions respectively significantly improving earlier results.

Keywords: *ATLAS, CERN, large extra dimensions, supersymmetry, compressed, squark, neutralino, monojet, ADD, standard model.*

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questo lavoro non sarebbe
stato possibile.

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*Si sta,
come d'autunno,
sugli alberi,
le foglie.*

[G. UNGARETTI, Soldati]

Acronyms

E_T transverse energy. [36](#)

η pseudorapidity. [36](#)

E_T^{miss} missing transverse momentum. [68](#)

p_T transverse momentum. [36](#)

pp proton-proton. [33](#)

ADC Analog to Digital Converter. [47](#)

ADD Arkani-Hamed, Dimopoulos, Dvali. [xiv](#), [1–3](#), [30](#), [31](#), [41](#), [73](#), [97–103](#), [109](#), [117–120](#), [123](#), [125](#), [128](#)

ALFA Absolute Luminosity For ATLAS. [40](#)

ATLAS A Toroidal LHC apparatus. [1–3](#), [22](#), [28](#), [31](#), [34](#), [36](#), [67](#), [71](#), [72](#), [75](#), [86](#), [88](#), [95](#), [97](#), [98](#), [102](#), [110](#), [115](#), [119](#), [122](#), [127](#), [128](#)

BG background. [81](#)

BIB Beam Induced Background. [83](#)

BSM Beyond Standard Model. [25](#), [73](#), [77](#), [82](#), [95](#), [99](#), [122](#), [127](#)

CB Combined. [62](#)

CERN European Organization for Nuclear Research. [1](#), [33](#)

CIS Charge Injection System. [47](#)

CL Confidence Level. [92](#), [120](#), [122](#), [127](#), [128](#)

CMS Compact Muon Solenoid. [1](#), [28](#)

COOL ATLAS Condition Database. [48](#)

CR Control Region. [78](#), [81](#), [82](#), [106](#), [109](#)

CSC Cathode Strip Chambers. [39](#)

CST Calorimeter Soft Term. [68](#)
CT Calorimeter Tagged. [62](#)
EB Extended Barrel. [45](#)
EFT Effective Field Theory. [30](#), [100](#), [102](#)
EM electromagnetic. [38](#), [65](#), [66](#)
EW Electroweak. [107](#), [108](#), [113](#), [114](#)
FCal LAr Forward Calorimeter. [39](#)
FEB Front-End Board. [61](#)
FSR Final State Radiation. [116](#), [119](#)
GCW Global Cell Weighting. [66](#)
GR General Relativity. [29](#)
GS Global Sequential. [66](#)
HEC Hadronic End-cap Calorimeter. [39](#)
HFN High Frequency Noise. [48](#)
HG High Gain. [47](#)
HGHG High Gain - High Gain. [49](#)
HGLG Hig Gain - Low Gain. [49](#)
HLT High Level Trigger. [40](#), [41](#)
HS Hard Scatter. [59](#), [64](#), [75](#)
HV High Voltage. [61](#)
IBL Insertable B-Layer. [38](#)
ID Inner Detector. [36](#), [61](#), [69](#)
IOV Interval Of Validity. [52](#)
IP interaction point. [34](#)
ISR Initial State Radiation. [93](#), [116](#), [119](#)
ITC Intermediate Tile Calorimeter. [46](#)
JES Jet Energy Scale. [65](#), [66](#), [110](#)
JVF Jet Vertex Fraction. [63](#)

JVT Jet Vertex Tagger. [64](#), [75–77](#)

KK Kaluza-Klein. [29](#), [97](#), [98](#)

L1 Level One. [40](#)

LAr Liquid Argon. [39](#)

LB Long Barrel. [45](#)

LCW Local Cluster Weighting. [41](#), [66](#)

LED Large Extra Dimension. [1](#), [97](#), [102](#), [123](#)

LFN Low Frequency Noise. [48](#)

LG Low Gain. [47](#)

LGHG Low Gain - High Gain. [49](#)

LGLG Low Gain - Low Gain. [49](#)

LHC Large Hadron Collider. [xiv](#), [1–3](#), [28](#), [33](#), [36](#), [38](#), [88–90](#), [92](#), [97](#), [98](#), [100](#)

LO Leading Order. [78](#), [98](#), [107](#), [108](#)

LSP Lightest Supersymmetric Particle. [27](#)

LUCID LUMinosity measurement using Cerenkov Integrating Detector. [40](#)

LVPS Low Voltage Power Supply. [49](#)

MB Minimum Bias. [88](#)

MC Monte Carlo. [48](#), [65](#), [78–81](#), [84](#), [85](#), [88](#), [99](#), [103](#), [104](#), [106–108](#), [110](#), [112](#), [114–116](#), [118](#), [123](#)

MDT Monitored Drift Tubes. [39](#)

MHT Missing Hadronic Trigger. [41](#)

MS Muon Spectrometer. [39](#)

MSSM Minimal Supersymmetric Standard Model. [27](#)

NCB Non Collision Background. [77](#), [103–107](#)

NLL Next to Leading Logarithm. [107](#), [108](#)

NLO Next to Leading Order. [78](#), [107](#), [108](#), [113](#), [115](#), [116](#)

NNLO Next to Next to Leading Order. [107](#), [108](#), [113](#)

OF Optimal Filtering. [47](#)

OR Overlap Removal. [67](#)

PDF Parton Distribution Function. [78](#), [85–88](#), [91](#), [100](#), [101](#), [103](#), [113](#), [114](#), [117](#), [118](#)

PMT PhotoMultiplier Tube. [45](#)

PS Proton Synchrotron. [33](#)

PSB Proton Synchrotron Booster. [33](#)

PU Pile Up. [64](#)

PV Primary Vertex. [59](#)

QCD Quantum Chromo Dynamics. [5](#), [77](#), [87](#), [88](#), [105–108](#), [113](#), [114](#)

RMS Root Mean Square. [48](#)

ROD ReadOut Driver. [47](#)

RoIs Region of Interest. [40](#)

RPC Resistive Plate Chambers. [39](#)

SA Standalone. [62](#)

SCT SemiConductor Tracker. [37](#)

SM Standard Model. [1](#), [6](#), [90](#), [120](#), [123](#), [124](#), [127](#)

SPS Super Proton Synchrotron. [33](#)

SR Signal Region. [74](#), [78](#), [81](#), [109](#)

ST Segment Tagged. [62](#)

SUSY Supersymmetry. [1](#), [3](#), [27](#), [28](#), [71–73](#), [75](#), [76](#), [85–89](#), [92](#), [93](#), [95](#), [96](#), [103](#), [108](#), [117–119](#), [127](#), [128](#)

TGC Thin Gap Chamber. [39](#)

TileCal Tile Calorimeter. [39](#)

TNF Tile Noise Filter. [55](#)

TRT Transition Radiation Tracker. [37](#)

TST Track Soft Term. [68](#), [69](#), [73](#)

TUCS TileCal Universal Calibration Software. [52](#)

WIMP Weakly Interacting Massive Particle. [26](#), [73](#)

WSF Wavelength Shifting Fibers. [45](#), [46](#)

ZDC Zero-Degree Calorimeter. [40](#)

Sammanfattning

Den stora hadron kollideraren (Large Hadron Collider eller LHC) vid CERN är världens kraftfullaste partikelacceleratoranläggning. Den krockar protoner med protoner och även tunga joner. Sedan 2015 har LHCs proton-proton kollisionerna en mass-centrum-energi som uppgår till 13 TeV. ATLAS är en av fyra stora experiment som registrerar LHCs kollisioner.

Tile Calorimeter är den centrala delen av ATLAS hadron kalorimetern. Den mäter upp enstaka hadroner, hadron kvastar, tau partiklar och den totala rörelsemängden. För att noggrant kunna mäta dessa partiklar krävs en exakt model av det elektroniska bruset i kalorimetern. Arbetet för att mäta och studera bruset samt beräkna relaterade kalibreringskonstanter som används för precisionsfysik i ATLAS presenteras.

Avhandlingen presenterar resultat från två olika sökningar efter nya fenomen m.h.a ATLAS experimentets data. Det första arbetet är en sökning efter produktion av supersymmetriska kvarkar, så kallade skvarkar, som sönderfaller till vanliga kvarkar och den supersymmetriska partikeln $\tilde{\chi}_1^0$. $\tilde{\chi}_1^0$ skulle kunna utgöra universums mörk materia. Det andra arbetet är en sökning efter stora extra rumsdimensioner i Arkani-Hamed, Dimopoulos, Dvali (ADD) modellen. Båda arbeten använder kollisioner som uppvisar stor odetekterad rörelsemängd och en högenergetisk hadron kvast i motsatt riktning.

Sökningen efter skvarkar använder data från 2015, sammanlagt 3.2 fb^{-1} , och sökningen efter extra rumsdimensioner använder data från 2015 och 2016, sammanlagt 36.1 fb^{-1} . I båda fallen stämmer det experimentella utfallet överens med bakgrundsberäkningarna. Därmed utesluts supersymmetriska modeller med skvarkar lättare än 608 GeV (för en masskillnad mot $\tilde{\chi}_1^0$ ner till 5 GeV). För modeller med extra rumsdimensioner, nya gränser sätts för den effektiva Planck massan M_D i ADD modeller up till 7.7 TeV för två extra rumsdimensioner och upp till 4.8 TeV för modeller med sex extra rumsdimensioner.

Abstract

The Large Hadron Collider (LHC) is the most powerful particle accelerator built to date. It is a proton-proton and heavy ion collider which in 2015 and 2016 operated at an unprecedented center of mass energy of $\sqrt{s} = 13$ TeV. The Tile Calorimeter is the ATLAS hadronic calorimeter covering the central region of the detector. It is designed to measure hadrons, jets, tau particles and missing energy. In order to properly reconstruct these physical objects a careful description of the electronic noise is required. This thesis presents the work done in updating, monitoring and studying the noise calibration constants used in the processing and identification of hadronic jet in the 2011 data.

The LHC proton-proton data collected by the ATLAS experiment at 13 TeV center of mass is used in this thesis to perform two searches for new phenomena: a search for supersymmetric particles in the compressed supersymmetric squark-neutralino model and a search for heavy gravitons in the large extra dimension Arkani-Hamed, Dimopoulos, Dvali (ADD) model. This work exploits an experimental signature with a high energy hadronic jet and large missing transverse energy, often referred to as monojet signature.

The search for supersymmetry presented in this thesis is based on the data recorded by the ATLAS experiment in 2015 which represents an integrated luminosity of 3.2 fb^{-1} . No significant excess compared to the Standard Model prediction is observed and exclusion limits are set on squark pair production with the subsequent squark decay $\tilde{q} \rightarrow q + \tilde{\chi}_1^0$ ($q = u, d, c, s$). The limits are set as function of the neutralino mass and focus on the least experimentally constrained situation with a small mass difference between the squark and the neutralino. Squark masses up to 608 GeV are excluded for $m_{\tilde{q}} - m_{\tilde{\chi}_1^0} = 5$ GeV. The second work presented in this thesis is a search for heavy gravitons in the large extra dimension ADD model using the full 2015 + 2016 dataset, representing an integrated luminosity of 36.1 fb^{-1} . A good agreement between data and Standard Model prediction is observed and exclusion limits are set on the effective Planck scale M_D . Values of M_D lower than 7.7 and 4.8 TeV for two and six hypothesized large extra dimensions are excluded. Both searches significantly improve earlier results.

Chapter 1

Introduction

The Standard Model (SM) of particle physics is the theory used to describe the elementary constituents of matter and their interactions and through the years it has been tested by many experiments [1]. Despite its success it cannot explain for instance the so called hierarchy and dark matter problems (see Chapter 3 for more details).

Supersymmetry [2] is an extension of the Standard Model that solves the hierarchy problem introducing a supersymmetric partner to each SM particle canceling in this way the contribution to the quantum correction to the Higgs mass of the Standard Model particles that push it to the Planck scale. When the lightest *neutralino* (denoted $\tilde{\chi}_1^0$) is also the lightest Supersymmetry (SUSY) particle and is stable, then it is a good candidate for Dark Matter.

Large Extra Dimension (LED) in the ADD model scenario [3] sets the Planck scale at the order of the weak scale formally eliminating the distinction between the two energy scales and introduces large extra spatial dimensions in which only the *graviton* (the hypothesized particle that mediates the gravitational interaction) is allowed to propagate in order to recover the strength of the gravitational interaction.

The LHC is a proton-proton and heavy ion collider designed to operate at a center of mass energy of $\sqrt{s} = 13$ TeV located at European Organization for Nuclear Research (CERN). In the context of a minimal supersymmetric model the neutralino could be produced in squark ($\tilde{q}$) pair production in proton-proton collisions with $\tilde{q} \rightarrow q + \tilde{\chi}_1^0$ ($q = u, d, c, s$). Similarly gravitons could be produced in association with jets. If the LHC produces such kind of particles in proton-proton collisions, the two main general purpose detectors A Toroidal LHC apparatus (ATLAS) and Compact Muon Solenoid (CMS) should be able to infer their presence. This thesis presents two searches: the search for compressed supersymmetric squark-neutralino signal with the ATLAS detector in the 3.2 fb^{-1} delivered in 2015 and for large extra dimensions in the ADD model scenario using 36.1 fb^{-1} collected by ATLAS in 2015 and 2016 in an experimental signature with jets and large missing transverse momentum in the final state.

1.1 Content Overview

An overview of the theoretical framework of the Standard Model is given in Chapter 2, it describes how the different elementary particles discovered so far are organized and the main ideas that lead to the current formulation of the Standard Model. Chapter 3 presents the hierarchy and dark matter problems and introduces supersymmetry and large extra dimensions. The LHC and the ATLAS experiment are introduced in Chapter 4 where the different sub-systems that make up the ATLAS detector are described. Studies on the stability of the electronic noise in the hadronic calorimeter and work to derive noise thresholds for the reprocessing and hadronic jet reconstruction in 2011 data are described in Chapter 5.

In order to search for physics beyond the standard model it is necessary to define criteria based on measured quantities to identify for example electrons, muons, jets, missing energy and displaced vertices from b -quark decays in an ATLAS collision. The criteria for identifying these particles and objects are described in some detail in Chapter 6 expanding in particular on some aspects that are specific to the searches presented in this thesis. The methodology and results for the search for supersymmetry in a compressed scenario with an energetic jet and large missing energy with the 2015 data are presented in Chapter 7. The results of the search for large extra dimensions in the ADD model scenario with the combined 2015 + 2016 data are presented in Chapter 8. Finally Chapter 9 presents an overall conclusion of this thesis.

1.2 Author's Contribution

My contribution to the ATLAS experiment started in early 2013 by studying the electronic noise in the hadronic calorimeter as part of my work to become a signing author of the ATLAS collaboration. The Tile Calorimeter is designed to measure jets, tau particles, missing momentum and for the energy reconstruction of hadrons, thus having an up-to-date description of the noise in the detector is important for most physics analysis in ATLAS. For the reprocessing of the 2011 data I developed a set of python scripts to study the noise constants variations over several calibration runs. Some of the software I produced is still in use today. Based on these scripts I carried out new studies of noise evolution some of which are described in Chapter 5 and produced noise calibration constants used in the reprocessing of the ATLAS data which is later used in precision measurements performed with the full ATLAS run 1 (2010 - 2012) dataset. I presented the performance of the ATLAS Tile calorimeter on behalf of the ATLAS collaboration in a poster session at the “XXVII International Symposium on Lepton Photon Interactions at High Energies (2015)”, this work appears in the proceedings of the conference in Ref. [4].

I later started analyzing the ATLAS data joining the monojet team in the effort of trying to answer relevant Standard Model open questions. During run 1 my contribution was limited to the study of the parton distribution function systematic uncertainties associated to the cross section and acceptance for extra dimension models.

For LHC run 2 (2015 - 2018) the ATLAS software framework was radically changed along with the formats for data analysis thus the Stockholm analysis code had to be largely rewritten. I am the main author of this software that applies calibrations to electrons, muons and jets, calculates systematic uncertainties and selects the events required for data analysis. This software also produces end-user last analysis stage ROOT [5] files on the LHC computing grid that are small enough to be imported to the department computers for final stage analysis and statistical interpretation.

I performed for the first time studies that demonstrated the sensitivity of the monojet analysis to the compressed supersymmetric squark-neutralino model. In order to ensure orthogonality with other SUSY searches it was decided to apply an upper cut on the number of jets. I performed optimization studies of such cut along with signal region definition using a multivariate method based on the TMVA [6] package in order to retain sensitivity to the SUSY compressed. I contributed to the development of the limit setting framework and to its optimization that allowed to reduce the turn around time from many days down to few hours allowing to quickly check and debug the analysis results. Finally I have performed the analysis of the 2015 data and derived myself the exclusion limits for the compressed squark-neutralino model and studied the theoretical systematics for this model.

The second analysis presented in this thesis is based on the combined 2015 + 2016 data collected by the ATLAS detector and is used to derive new experimental constraints on the large extra dimension in the ADD model scenario. In this context I calculated the theoretical systematic uncertainties affecting the cross section and acceptance for this signal and implemented them in the fit. I then calculated the exclusion limits on the effective Planck scale for large extra spatial dimensions. The results presented in this thesis led to the following publications:

- A - *Search for new phenomena in final states with an energetic jet and large missing transverse momentum in pp collisions at $\sqrt{s} = 13$ TeV using the ATLAS detector* given in Ref. [7].
- B - *Search for dark matter and other new phenomena in events with an energetic jet and large missing transverse momentum using the ATLAS detector* given in Ref. [8]. This conference note is being converted to a paper and is in the final stages of ATLAS approval.

As mentioned earlier I also contributed to the calculations of the systematic uncertainties related to the parton distribution functions for the ADD large extra dimension model that was part of the run 1 monojet paper:

- C - *Search for new phenomena with the monojet and missing transverse momentum signature using the ATLAS detector in proton-proton collisions* given in Ref. [9].

Finally, as stated in the beginning of this section, I am the author of the ATLAS Tile calorimeter performance poster for the 2015 Lepton Photon conference that resulted in the following conference proceeding:

D - *Performance of the ATLAS Tile Calorimeter* given in Ref. [4].

The text of this thesis relies in part on my licentiate thesis with the title: *Search for Supersymmetry in Monojet Final States with the ATLAS Experiment* [10] and for this reason some of the chapters in this thesis are taken or adapted from there. In particular Chapter 2 has been largely revisited, Chapter 3 has been extended, Chapter 4 is taken to a large extent from the licentiate and a section has been added to Chapter 5. Chapter 6 has been revised and extended to include aspects relevant to the analysis of the 2015 and 2016 data for the search for large extra dimensions. Chapter 7 has been revisited and expanded while Chapter 8 is entirely new.

Chapter 2

Theoretical Overview of the Standard Model

2.1 Introduction

It is possible to set the start of elementary particle physics in 1897 with the discovery of the electron by Thompson [11] but it was not until the beginning of the 1900 when the idea of the photon came along that the first big classification of particles could be made. In modern particle physics we classify particles in elementary *fermions*, the matter constituents, and *bosons*, the force carriers, and describe interactions using the principles of *gauge theories*.

In 1961 Glashow [12] introduced a model for the weak and electromagnetic interactions based on the $SU(2) \times U(1)$ symmetry group. Requiring gauge invariance under phase transformations imply that the gauge bosons are massless but it was known that they should be about hundred times heavier than a proton in order to account for the short range of the weak interaction. In 1967 and 1968 Weinberg [13] and Salam [14] independently published a paper where they used the $SU(2) \times U(1)$ symmetry group but by introducing the spontaneous breaking of the gauge symmetry recovered the masses of the weak force gauge bosons while leaving the photon massless (see Section 2.6). In order for the breaking of the symmetry to occur, the *Higgs* (H) field must be introduced. The Higgs field is a massive scalar particle (spin zero) that, unlike the other gauge bosons, is not the mediator of any interaction but is responsible of giving mass to the particles coupling to it. The Higgs boson was discovered in 2012 [15, 16]. At high energies the theory developed by Glashow, Weinberg and Salam unify the weak and the electromagnetic interactions in the *electroweak* one and is known as the Standard Model of electroweak interactions.

Currently the Standard Model of particle physics also includes the *strong* interaction and is the theoretical model that describes the properties of all the known particles and their interactions. It is based on the $SU(3)_C \times SU(2)_L \times U(1)_Y$ symmetry group where $SU(3)_C$ is the symmetry group of the Quantum Chromo Dynamics (QCD) and the “C” stands for the color charge carried by quarks and gluons, the $SU(2)_L$ is the weak isospin group acting on *left-handed*

doublet of fermions while the $U(1)_Y$ group is the *hypercharge* symmetry group. The Standard Model is a very successful and well tested theory that manage to explain most of the observed phenomena in particle physics. Some of the still open questions of the SM are addressed in Chapter 3.

2.2 Particle Classification

Particles are classified in the Standard Model according to their properties, the first one that can be used is the spin. Fermions have *half integer* while bosons have *integer* spin. Fermions are further divided in *leptons* and *quarks* which are currently assumed to be truly elementary particles meaning that they lack substructure and can be treated as point-like. Leptons, unlike quarks, do not have a *color charge* and therefore cannot experience the *strong* interaction. There are six types or *flavors* of leptons divided in *three generations* according to their *lepton number*. Each generation consists of an electrically charged and a neutral lepton, the *electron* (e^-) and the *electron neutrino* (ν_e) constitute the first generation, they both carry *electron number* ($L_e = 1$). The *muon* (μ^-) and the *muon neutrino* (ν_μ) have a *muon number* ($L_\mu = 1$) and make up the second generation. The *tau* (τ^-) and the *tau neutrino* (ν_τ) in the third generation carry *tau number* ($L_\tau = 1$).

There are also six flavors of quarks divided into three generations each of which consists of one $+2/3$ and one $-1/3$ electric charge quark. The first generation groups together the *up* (u) and *down* (d) quarks, the second the *charm* (c) and *strange* (s) while the third the *top* (t) and the *bottom* (b). As mentioned earlier, each quark also has color charge with three possibilities: *Red*, *Green* and *Blue* (R, G, B) and *baryonic number* of $1/3$. All the visible (detectable and stable) matter in the universe is made of fermions belonging to the first generation while the other two quickly decay to the first one.

Elementary bosons of the Standard Model have *spin one* and mediate the three fundamental interactions: the *electromagnetic*, the *weak* and the *strong*. The *photon* (γ) mediates the electromagnetic interaction, the $W^\pm$ (with electric charge) and the Z^0 (neutral) bosons are the weak interaction carriers and finally there are eight *gluons* that are responsible of carrying the strong force. The Standard Model also contains a spin zero boson, the *Higgs* boson, discussed in Section 2.6. The Standard Model does not describe gravitation at the microscopic level but it is believed to be mediated by a massless gauge boson with spin two called the *graviton* which so far (as of 2017) has not been detected by experiments. The photon and the gluons are massless thus the interaction they mediate has infinite range. This is true for the photon but since gluons carry the color charge themselves they can self-interact thus shortening the range of the strong interaction.

Finally the *hadron* category groups together *mesons* and *baryons*. Mesons are bosons that are a bound state of one quark and one anti-quark while *baryons* are fermions made of three quarks.

2.3 Electroweak Symmetry Group

In 1930 Pauli suggested the introduction of a new particle in order to explain the continuous spectrum of β decays [17]. Since the neutron was not discovered yet Pauli named his hypothesized particle that way as it had to be electrically neutral in order to conserve charge. It was not until Fermi proposed his theory for β decay [18, 19] that the new particle was accepted and named *neutrino*. The Lagrangian in Fermi's theory is of the current-current type and can be written as:

$$\mathcal{L}_{\text{weak}} = \frac{G_F}{\sqrt{2}} J_\mu^\dagger J^\mu \quad (2.3.1)$$

where G_F is the Fermi constant and the J_μ current is given by:

$$J_\mu = \bar{\psi} \gamma^\mu \psi, \quad (2.3.2)$$

here ψ , $\bar{\psi}$ are the spinors of the fermions involved in the interaction and γ_μ ($\mu = 0, 1, 2, 3$) are the Dirac matrices. In 1949 Powell [20] discovered in cosmic rays two particles that he called the tau and the theta mesons. The tau meson could decay by weak interaction into three pions ($\tau^+ \rightarrow \pi^+ \pi^+ \pi^-$) and the theta decays into two pions ($\theta^+ \rightarrow \pi^+ \pi^0$). The measured lifetime and mass of the tau and the theta mesons turned out to be compatible within experimental uncertainties hinting that they could be the same particle. If that was the case parity¹ conservation would be violated. This is known as the $\tau - \theta$ puzzle. In 1950 Lee and Yang proposed that parity could be violated in weak interactions [21] which was experimentally confirmed in 1956 by Wu [22]. The tau and theta mesons were identified as the *kaon* which decays are both parity conserving ($K^+ \rightarrow \pi^+ \pi^+ \pi^-$) and parity violating ($K^+ \rightarrow \pi^+ \pi^0$). Wu's experiment involved the β decay of Cobalt 60 atoms to Nickel 60. The Cobalt atoms were immersed in a magnetic field in order to align the spin of the nucleus with it. The observable in the experiment is the product between the spin of the electron and its momentum ($\vec{J} \cdot \vec{p}$). Reverting the magnetic field is the same as applying a parity transformation and since the momentum transforms as a vector and changes sign under space reflection ($\vec{p} \rightarrow -\vec{p}$) but the spin transforms as an axial vector and remains unchanged ($\vec{J} \rightarrow \vec{J}$) the product $\vec{J} \cdot \vec{p}$ is not parity conserving and an asymmetry in the distributions of the experiment and its reflected version were observed.

In 1958 Feynman and Gell-Mann proposed a parity violating extension of Fermi's theory [23] called the V-A (vector minus axial vector) theory which introduces the *chiral* operators:

$$P_L = \frac{1}{2}(1 - \gamma^5); \quad P_R = \frac{1}{2}(1 + \gamma^5) \quad (2.3.3)$$

¹Parity is a space-time symmetry in which $\vec{r} \rightarrow \vec{r}' = -\vec{r}$ and $t \rightarrow t' = t$ where $\vec{r}$ and t are the position vector and time respectively thus $\vec{p} \rightarrow \vec{p}' = -\vec{p}$, $\vec{L} = \vec{r} \times \vec{p} \rightarrow \vec{L}' = \vec{L} \Rightarrow \vec{J} \rightarrow \vec{J}' = \vec{J}$ where $\vec{p}$, $\vec{L}$ and $\vec{J}$ are the momentum, the angular momentum and the total angular momentum respectively. For the helicity we have also $\lambda = \vec{J} \cdot \vec{p}/|\vec{p}| \rightarrow \lambda' = -\lambda$.

where $\gamma_5 = i\gamma_0\gamma_1\gamma_2\gamma_3$. The chiral operators projects the left and right chirality when applied on a spinor ψ . The weak current can be re-written as:

$$\begin{aligned} J_\mu &\equiv J_\mu^{(+)} = \bar{\psi}_{\nu_e} \gamma_\mu \frac{1}{2} (1 - \gamma_5) \psi_e \\ J_\mu^\dagger &\equiv J_\mu^{(-)} = \bar{\psi}_e \gamma_\mu \frac{1}{2} (1 - \gamma_5) \psi_{\nu_e}, \end{aligned} \quad (2.3.4)$$

here $J_\mu^{(+)}$ and $J_\mu^{(-)}$ are the charged weak currents and are a combination of vector and axial vector currents. The $\frac{1}{2}(1 - \gamma_5)$ operator only selects left-handed fermions (or right-handed anti-fermions) thus only these components participate in the weak interaction. Indicating for simplicity $\psi_e(x)$ as $e(x)$ and $\psi_{\nu_e}(x)$ as $\nu_e(x)$, Eq. (2.3.4) can be written as:

$$\begin{aligned} J_\mu^{(+)} &= \bar{\nu}_{eL} \gamma_\mu e_L \\ J_\mu^{(-)} &= \bar{e}_L \gamma_\mu \nu_{eL} \end{aligned} \quad (2.3.5)$$

where the subscript L indicates that the $\frac{1}{2}(1 - \gamma_5)$ have been applied and only the left-handed fermions are selected.

One problem of the V-A theory is that the total cross section for the inelastic neutrino scattering process $\nu_\mu + e^- \rightarrow \nu_e + \mu^-$ given by [24]:

$$\sigma = \frac{G_F^2}{\pi} s \quad (2.3.6)$$

where $s = 4E^2$ is the center of mass energy of the $\nu_\mu - e$ system grows indefinitely with energy. In his attempt to elaborate a theory for beta decay, Fermi assumed it to be point-like [19]. As we now know the short range of the weak force is due to the large masses of the mediators making it only appear point-like in relatively low energy phenomena such as the beta decay. Let us consider the following expansion in partial waves:

$$\frac{d\sigma}{d\Omega} = \left| \frac{1}{2ik} \sum_{\ell=0}^{\infty} (2\ell+1) A_\ell P_\ell(\cos \theta) \right|^2 \quad (2.3.7)$$

where A_ℓ is the partial wave amplitude for the orbital angular momentum ℓ and P_ℓ are the Legendre polynomials. For point-like scattering (also referred to as s -wave scattering) the $\ell = 0$ mode dominates [24] and there is no angular dependence ($P_0(\cos \theta) = 1$) thus we have:

$$\frac{d\sigma}{d\Omega} = \frac{1}{4E^2} |A_0|^2. \quad (2.3.8)$$

Requiring the unitarity ($|A_\ell| \leq 1$) for each partial wave, we have that:

$$\frac{d\sigma}{d\Omega} \leq \frac{1}{4E^2} \quad (2.3.9)$$

and for the total cross section:

$$\sigma = \int \frac{d\sigma}{d\Omega} d\Omega = 4\pi \frac{d\sigma}{d\Omega} \leq \frac{\pi}{E^2} \quad (2.3.10)$$

and we see that the cross section is subject to a *unitarity bound*. If we use the result of Eq. (2.3.6) we see that for energies greater than:

$$E \geq \sqrt{\frac{\pi}{2G_F}} \approx 367 \text{ GeV} \quad (2.3.11)$$

unitarity is violated. The V-A theory is thus only valid for energies comparable to the Fermi scale ($\approx 367 \text{ GeV}$) while for higher energies the point-like nature of the theory violates the unitarity of the scattering matrix. Feynman and Gell-Mann [25] proposed that the weak interaction could happen through the exchange of spin one bosons, this prompted Schwinger [26] to introduce a gauge theory of the weak interactions mediated by the $W^\pm$ bosons where he also tried to include the photon. A consistent theory for the weak interaction at high energies thus required the exchange of mediator particles and could be obtained by analogy to the electromagnetic case by formulating it as a gauge theory.

In order to determine the symmetry group from the weak interactions, let us simplify the notation by writing:

$$\chi_L = \begin{pmatrix} \nu_{eL} \\ e_L^- \end{pmatrix} \equiv \begin{pmatrix} \nu_e \\ e^- \end{pmatrix} \quad (2.3.12)$$

and using the Pauli matrices

$$\tau_\pm = \frac{1}{2}(\tau_1 \pm i\tau_2) \quad (2.3.13)$$

we have for the two charged currents of Eq. (2.3.4):

$$\begin{aligned} J_\mu^{(+)} &= \bar{\chi}_L \gamma_\mu \tau_+ \chi_L \\ J_\mu^{(-)} &= \bar{\chi}_L \gamma_\mu \tau_- \chi_L \end{aligned} \quad (2.3.14)$$

we can complete a possible $SU(2)$ symmetry group by introducing a *neutral* current:

$$J_\mu^{(3)} = \bar{\chi}_L \gamma_\mu \frac{\tau_3}{2} \chi_L = \frac{1}{2} \bar{\nu}_L \gamma_\mu \nu_L - \frac{1}{2} \bar{e}_L \gamma_\mu e_L \quad (2.3.15)$$

and we thus have an *isospin* triplet of weak currents:

$$J_\mu^i = \bar{\chi}_L \gamma_\mu \frac{\tau_i}{2} \chi_L \quad (2.3.16)$$

whose *generators* $T_i = \tau_i/2$ give an $SU(2)_L$ algebra. The subscript “L” indicates that only the left-handed chiral components of the fermions couples to the isospin triplet of weak currents. Global transformation of the $SU(2)_L$ group are of the form:

$$\chi_L(x) \rightarrow \chi'_L(x) = e^{i\vec{\varepsilon} \cdot \vec{T}} \chi_L(x) = e^{i\vec{\varepsilon} \cdot \frac{\vec{\tau}}{2}} \chi_L(x), \quad (2.3.17)$$

where χ_L is the *fundamental representation* of the group, the right handed fermions are singlet for the $SU(2)_L$, thus:

$$e_R \rightarrow e'_R = e_R. \quad (2.3.18)$$

Since we are considering the global transformations, we have no interaction, so the Lagrangian can be written as:

$$\mathcal{L} = \bar{e}i\gamma^\mu\partial_\mu e + \bar{\nu}i\gamma^\mu\partial_\mu\nu \equiv \bar{\chi}_L i\gamma\partial\chi_L + \bar{e}_R i\gamma\partial e_R; \quad (2.3.19)$$

where we have set $m_e = 0$ because the mass term couples right and left components of the fermions and it is not $SU(2)_L$ invariant. In 1973, the Gargamelle bubble chamber experiment [27] detected events of the type:

$$\bar{\nu}_\mu e^- \rightarrow \bar{\nu}_\mu e^- \quad (2.3.20)$$

$$\begin{cases} \nu_\mu N \rightarrow \nu_\mu X \\ \bar{\nu}_\mu N \rightarrow \bar{\nu}_\mu X \end{cases} \quad (2.3.21)$$

which are evidence of a neutral current. It was natural to try to identify this neutral current with the one of Eq. (2.3.15) but it was seen that the experimentally observed current has both left-handed and right-handed components thus the $J_\mu^{(3)}(x)$ current introduced above can not be used as it involves only left-handed fermions and needs to be combined with some other current with also right-handed components. The electromagnetic current is neutral and mixes left-handed and right-handed components:

$$J_\mu^{\text{em}}(x) = \bar{\psi}(x)\gamma_\mu Q\psi(x) \quad (2.3.22)$$

where Q is the electric charge operator (with eigenvalue $Q = -1$ for the electron) and is also the generator of the $U(1)_{\text{em}}$ symmetry group of the electromagnetism. For example in the case of the electron we have:

$$j_\mu^{\text{em}} = -\bar{e}\gamma_\mu e = -\bar{e}_L\gamma_\mu e_L - \bar{e}_R\gamma_\mu e_R. \quad (2.3.23)$$

The idea is to find the minimal symmetry group that contains the $SU(2)_L$ and the $U(1)_{\text{em}}$ generators and a current that completes the weak isospin triplet but is invariant for $SU(2)_L$ transformations. To this end we can define the *hypercharge* operator:

$$Y = 2(Q - T_3) \Rightarrow Q = T_3 + \frac{Y}{2}, \quad (2.3.24)$$

for the current we can write:

$$J_\mu^{\text{em}} = J_\mu^{(3)} + \frac{1}{2}J_\mu^Y \quad (2.3.25)$$

where:

$$J_\mu^Y = \bar{\chi}_L\gamma_\mu\chi_L - 2\bar{e}_R\gamma_\mu e_R. \quad (2.3.26)$$

The hypercharge Y generates a $U(1)_Y$ symmetry and since it is a $SU(2)_L$ singlet, leaves the Lagrangian of Eq. (2.3.19) invariant under the transformations:

$$\begin{aligned} \chi_L(x) &\rightarrow \chi'_L(x) = e^{i\beta Y}\chi_L(x) \equiv e^{i\beta y_L}\chi_L \\ e_R(x) &\rightarrow e'_R(x) = e^{i\beta Y}e_R(x) \equiv e^{i\beta y_R}e_R. \end{aligned} \quad (2.3.27)$$

We thus have unified the electromagnetic and weak interactions extending the gauge group to $SU(2)_L \times U(1)_Y$ and instead of having a single symmetry group we have a direct product of groups each with his own coupling constant thus in addition to the electric charge e we will have another coupling to be found. As mentioned in Section 2.2 the idea of an extended symmetry group that could unify the weak interaction and electromagnetism was first explored by Glashow and later by Weinberg and Salam. Since we have a direct product of symmetry groups, the generators of $SU(2)_L$ and those of $U(1)_Y$ commute. The commutation relations are give by:

$$[T_+, T_-] = 2T_3; \quad [T_3, T_\pm] = \pm T_\pm; \quad [Y, T_\pm] = [Y, T_3] = 0. \quad (2.3.28)$$

Members of the same isospin triplet, have same hypercharge eigenvalue. The relevant quantum numbers are summarized in Table 2.1.

Lepton	T	$T^{(3)}$	Q	Y	Quark	T	$T^{(3)}$	Q	Y
ν_e	$\frac{1}{2}$	$\frac{1}{2}$	0	-1	u_L	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{2}{3}$	$\frac{1}{3}$
e_L^-	$\frac{1}{2}$	$-\frac{1}{2}$	-1	-1	d_L	$\frac{1}{2}$	$-\frac{1}{2}$	$-\frac{1}{2}$	$\frac{1}{3}$
					u_R	0	0	$\frac{2}{3}$	$\frac{4}{3}$
e_R^+	0	0	-1	-2	d_R	0	0	$-\frac{1}{3}$	$-\frac{2}{3}$

Table 2.1: Weak isospin and hypercharge quantum numbers of leptons and quarks.

2.4 Electroweak Interactions

As seen in Section 2.3 the electroweak interaction is mediated by gauge bosons. In order to generate them, we consider *local* gauge transformations of the $SU(2)_L \times U(1)_Y$ group:

$$\begin{aligned} \chi_L &\rightarrow \chi'_L = e^{i\vec{\epsilon}(x) \cdot \vec{T} + i\beta(x)Y} \chi_L \\ \psi_R &\rightarrow \psi'_R = e^{i\beta(x)Y} \psi_R, \end{aligned} \quad (2.4.1)$$

introducing four gauge bosons, $\vec{W} \equiv (W_\mu^{(1)}, W_\mu^{(2)}, W_\mu^{(3)})$ and B_μ (same as the number of generators) and the *covariant derivative*:

$$\begin{aligned} D_\mu \chi_L &= \left(\partial_\mu + ig \frac{\vec{\tau}}{2} \cdot \vec{W}_\mu(x) + i \frac{g'}{2} y_L B_\mu(x) \right) \chi_L \\ &= \left(\partial_\mu + ig \frac{\vec{\tau}}{2} \cdot \vec{W}_\mu(x) - i \frac{g'}{2} B_\mu(x) \right) \chi_L \\ D_\mu \psi_R &= \left(\partial_\mu + i \frac{g'}{2} y_R B_\mu(x) \right) \psi_R \\ &= \left(\partial_\mu - i \frac{g'}{2} B_\mu(x) \right) e_R \end{aligned} \quad (2.4.2)$$

the Lagrangian of Eq. (2.3.19) can be written as:

$$\begin{aligned} \mathcal{L} = & \bar{\chi}_L i \gamma \partial \chi_L + \bar{e}_R i \gamma \partial e_R - g \bar{\chi}_L \gamma^\mu \frac{\vec{\tau}}{2} \chi_L \vec{W}_\mu \\ & + \frac{g'}{2} (\bar{\chi}_L \gamma^\mu \chi_L + 2 \bar{e}_R \gamma^\mu e_R) B_\mu - \frac{1}{4} \vec{W}_{\mu\nu} \vec{W}^{\mu\nu} - \frac{1}{4} B_{\mu\nu} B^{\mu\nu} \end{aligned} \quad (2.4.3)$$

where:

$$\begin{aligned} \vec{W}_{\mu\nu} &= \partial_\mu \vec{W}_\nu - \partial_\nu \vec{W}_\mu - g \vec{W}_\mu \times \vec{W}_\nu \\ B_{\mu\nu} &= \partial_\mu B_\nu - \partial_\nu B_\mu \end{aligned} \quad (2.4.4)$$

are the kinetic plus non abelian interaction term for the $SU(2)_L$ symmetry (first equation) and the kinetic term for the abelian symmetry group $U(1)_Y$ (second equation). We can now split the Lagrangian terms to find out the field of the vector bosons coupled to the charged currents and to the neutral current.

2.4.1 Charged Currents Interaction

Let us consider the term:

$$\mathcal{L}_{int}^{ew} = -g \bar{\chi}_L \gamma_\mu \frac{\vec{\tau}}{2} \chi_L \vec{W}_\mu + \frac{g'}{2} \bar{\chi}_L \gamma_\mu \chi_L B^\mu + g' \bar{e}_R \gamma_\mu e_R B^\mu \quad (2.4.5)$$

defining:

$$W_\mu^\pm = \frac{1}{\sqrt{2}} \left(W^{(1)}_\mu \mp i W^{(2)}_\mu \right) \quad (2.4.6)$$

we can write:

$$\mathcal{L}^{CC} = -\frac{g}{\sqrt{2}} \left(J_\mu^{(+)} W^{-\mu} + J_\mu^{(-)} W^{+\mu} \right) \quad (2.4.7)$$

and recognize two charged vector bosons, the W^+ and the W^- , with coupling given by “ g ”.

2.4.2 Neutral Current Interaction

The relevant term left to consider for what concerns the electroweak currents is:

$$\mathcal{L}^{NC} = -g J_\mu^{(3)} W^{(3)\mu} - \frac{g'}{2} J_\mu^Y B^\mu, \quad (2.4.8)$$

the electromagnetic interaction, $-ie J^{(em)\mu} A_\mu$, is embedded in this expression as will become clear considering the *spontaneously broken symmetry* phenomena, for now, is sufficient to define:

$$\begin{aligned} W_\mu^{(3)} &= \cos \theta_w Z_\mu + \sin \theta_w A_\mu \\ B_\mu &= -\sin \theta_w Z_\mu + \cos \theta_w A_\mu \end{aligned} \quad (2.4.9)$$

and invert to get:

$$\begin{aligned} A_\mu &= \sin \theta_w W_\mu^{(3)} + \cos \theta_w B_\mu \\ Z_\mu &= \cos \theta_w W_\mu^{(3)} - \sin \theta_w B_\mu \end{aligned} \quad (2.4.10)$$

where θ_w is the electroweak *mixing angle*. Plugging this into Eq. (2.4.8) and rearranging the terms we have:

$$\begin{aligned} \mathcal{L}^{NC} = - & \left[\left(g \sin \theta_w J_\mu^{(3)} + \frac{g'}{2} \cos \theta_w J_\mu^Y \right) A^\mu \right. \\ & \left. + \left(g \cos \theta_w J_\mu^{(3)} - \frac{g'}{2} \sin \theta_w J_\mu^Y \right) Z^\mu \right] \end{aligned} \quad (2.4.11)$$

since A^μ is the photon field, the first parenthesis must be identified with the electromagnetic current, thus:

$$e J_\mu^{\text{em}} A^\mu \equiv e \left(J_\mu^{(3)} + \frac{J_\mu^Y}{2} \right) A^\mu = \left(g \sin \theta_w J_\mu^{(3)} + \frac{g'}{2} \cos \theta_w J_\mu^Y \right) A^\mu \quad (2.4.12)$$

from which we get the relation:

$$g \sin \theta_w = g' \cos \theta_w = e \quad (2.4.13)$$

and we can rewrite Eq. (2.4.8) as:

$$\mathcal{L}^{NC} = - \frac{g}{\cos \theta_w} \left[J_\mu^{(3)} - \sin^2 \theta_w J_\mu^{(em)} \right] Z^\mu \quad (2.4.14)$$

so that Z^μ can be identified with the field for the weak neutral vector boson.

2.5 Strong Interaction

The dynamic of the quarks can be described by the Lagrangian:

$$\mathcal{L} = \bar{q}_i (i \gamma^\mu \partial_\mu - m) q_i + \mathcal{L}_{\text{int}} \quad (2.5.1)$$

where q_i ($i = 1, 2, 3$) are the color fields and m is the mass of the quark and it only depends on the quark flavor and not on the color charge. Let us consider only one of the three possible color fields, thus $q_i(x) \equiv q(x)$. To generate the interaction Lagrangian, we can impose on the free Lagrangian:

$$\mathcal{L}_0 = \bar{q} (i \gamma^\mu \partial_\mu - m) q \quad (2.5.2)$$

gauge invariance under local $\text{SU}(3)_C$ transformations of the type:

$$q(x) \rightarrow q'(x) = e^{i \sum_{k=1}^8 \epsilon^k(x) T^k} q(x) \equiv \Omega(\epsilon(x)) q(x) \quad (2.5.3)$$

where $T^k = (T^k)^\dagger \equiv \lambda^k/2$ ($k = 1, \dots, 8$) are the generators of $\text{SU}(3)_C$, $\epsilon^k(x)$ are arbitrary space-time functions and a common choice for the $\lambda^k/2$ are the Gell-Mann matrices. The commutation relation of the group generators is:

$$[T^k, T^l] = i f^{klm} T^m \quad (2.5.4)$$

where f^{klm} are the structure constants of the group and it can be seen that not all the generators commute thus the symmetry group is non-Abelian (i.e.

$\Omega(\epsilon_1)\Omega(\epsilon_2) \neq \Omega(\epsilon_2)\Omega(\epsilon_1)$). If we consider the derivative of the transformed color field:

$$\begin{aligned} \partial_\mu q &\rightarrow \partial_\mu q' = \Omega(\epsilon)\partial_\mu q + (\partial_\mu u \Omega(\epsilon)) q \\ &\equiv \Omega \left[\partial_\mu + \Omega^{-1} \partial_\mu \Omega \right] q \neq \Omega \partial_\mu q \end{aligned} \quad (2.5.5)$$

and we see that the invariance under the $\Omega(\epsilon(x))$ transformation is broken. We need to define a derivative such that:

$$D_\mu q \rightarrow (D_\mu q)' = \Omega(\epsilon(x)) D_\mu q \equiv (\Omega D_\mu \Omega^{-1}) \Omega q \equiv D'_\mu q', \quad (2.5.6)$$

to this end we introduce eight vector fields $G_\mu^k(x)$ ($k = 1, \dots, 8$) such that:

$$D_\mu q(x) = (\partial_\mu + ig T^k G_\mu^k) q(x) \quad (2.5.7)$$

where we use the conventions on repeated indexes $T^k G_\mu^k \equiv \sum_k T^k G_\mu^k(x)$. We want to find how the $G_\mu^k(x)$ fields transform under the wanted transformation property of Eq. (2.5.6), to this end:

$$(D_\mu q)' = \left(\Omega \partial_\mu + \partial_\mu \Omega + ig T^k G_\mu^k \Omega \right) q \equiv \Omega D_\mu q = \Omega (\partial_\mu + ig T^k G_\mu^k) q \quad (2.5.8)$$

from which:

$$T^k G_\mu^k = \Omega T^k G_\mu^k \Omega^{-1} + \frac{i}{g} (\partial_\mu \Omega) \Omega^{-1} = \Omega T^k G_\mu^k \Omega^{-1} - \frac{i}{g} \Omega \partial_\mu \Omega^{-1}. \quad (2.5.9)$$

Considering the infinitesimal transformations of the type:

$$\Omega(\epsilon) \sim 1 + i\epsilon^k T^k; \quad \Omega^{-1}(\epsilon) \sim 1 - i\epsilon^k T^k, \quad (2.5.10)$$

from Eq. (2.5.9), we obtain:

$$\begin{aligned} T^k G_\mu^k &= (1 + i\epsilon^l T^l) T^m G_\mu^m (1 - i\epsilon^l T^l) \\ &\quad - \frac{i}{g} (1 + i\epsilon^l T^l) (-i\partial_\mu \epsilon^k) T^k \\ &\sim T^k G_\mu^k + i\epsilon^l \underbrace{[T^l, T^m]}_{if^{lmk} T^k} G_\mu^m - \frac{1}{g} (\partial_\mu \epsilon^k) T^k \\ &\equiv T^k \left[G_\mu^k - f^{klm} \epsilon^l G_\mu^m - \frac{i}{g} \partial_\mu \epsilon^k \right]. \end{aligned} \quad (2.5.11)$$

And the looked for transformations are:

$$G_\mu^k(x) \rightarrow G_\mu^k(x) = G_\mu^k(x) - f^{klm} \epsilon^l(x) G_\mu^m(x) - \frac{i}{g} \partial_\mu \epsilon^k(x) \quad (2.5.12)$$

and we are sure that the Lagrangian:

$$\mathcal{L}(q, D_\mu q) = \bar{q} (i\gamma^\mu D_\mu - m) q \quad (2.5.13)$$

is invariant under $SU(3)_C$ local gauge transformations. The $G_\mu^k(x)$ ($k = 1, \dots, 8$) fields are called *gluons* and are the carriers of the strong force between quarks as a consequence of the color charge.

To be able to write the full Lagrangian we need to find the equivalent of $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ of the Abelian case also for the gluons. This is needed since up to now the gluons were not treated as a dynamic degree of freedom and we need to add their kinetic term in the Lagrangian. It can be seen that in the Abelian case:

$$[D_\mu, D_\nu] = -ieF_{\mu\nu} \quad (2.5.14)$$

thus, by analogy, we can calculate:

$$\begin{aligned} [D_\mu, D_\nu] &= [\partial_\mu + igT^k G_\mu^k, \partial_\nu + igT^k G_\nu^k] \\ &= [\partial_\mu, \partial_\nu] + ig [T^k G_\mu^k, \partial_\nu] + ig [\partial_\mu, T^k G_\nu^k] - g^2 [T^k G_\mu^k, T^k G_\nu^k] \\ &= ig (\partial_\mu T^k G_\nu^k - \partial_\nu T^k G_\mu^k) - g^2 \underbrace{[T^k, T^l]}_{ifklm T^m} G_\mu^k G_\nu^l \\ &= ig T^m (\partial_\mu G_\nu^m - \partial_\nu G_\mu^m - g f^{mkl} G_\mu^k G_\nu^l) \\ &\equiv ig T^m G_{\mu\nu} \equiv ig G_{\mu\nu} \end{aligned} \quad (2.5.15)$$

where $G_{\mu\nu}^k = \partial_\mu G_\nu^k - \partial_\nu G_\mu^k - g f^{klm} G_\mu^l G_\nu^m$ and $G_{\mu\nu} = T^k G_{\mu\nu}^k$ have been defined. Finally indicating with $G_\mu \equiv T^k G_\mu^k$ we have:

$$G_{\mu\nu} = \partial_\mu G_\nu - \partial_\nu G_\mu + ig [G_\mu, G_\nu]. \quad (2.5.16)$$

The field strength tensor of Eq. (2.5.16) transforms according to:

$$G'_{\mu\nu} = \Omega G_{\mu\nu} \Omega^{-1} \quad (2.5.17)$$

and for any product of these tensors we have:

$$G'_{\mu\nu} G'_{\rho\sigma} \dots G'_{\lambda\tau} = U (G_{\mu\nu} G_{\rho\sigma} \dots G_{\lambda\tau}) U^{-1} \quad (2.5.18)$$

and the trace of these products is gauge invariant:

$$\text{Tr}(G'_{\mu\nu} \dots G'_{\lambda\tau}) = \text{Tr}(U G_{\mu\nu} \dots G_{\lambda\tau} U^{-1}) = \text{Tr}(G_{\mu\nu} \dots G_{\lambda\tau}). \quad (2.5.19)$$

Thus in order to retain local gauge invariance for the gluon kinetic term the simplest Lagrangian can be written as:

$$\mathcal{L}_G = \text{const Tr}(G_{\mu\nu} G^{\mu\nu}) \equiv \text{const Tr}(T^k T^l) G_{\mu\nu}^k G^{\mu\nu l} \quad (2.5.20)$$

using the known property of the Gell-Mann matrices $\text{Tr}(T^k T^l) = \text{Tr}(\lambda^k \lambda^l)/2$ and choosing the constant term we have:

$$\mathcal{L}_G = -\frac{1}{2} \text{Tr}(G_{\mu\nu} G^{\mu\nu}) = -\frac{1}{4} \sum_{k=1}^8 G_{\mu\nu}^k G^{\mu\nu k}. \quad (2.5.21)$$

The full Lagrangian can finally be written as:

$$\mathcal{L} = \bar{q} (i\gamma^\mu D_\mu - m) q - g (\bar{q} \gamma^\mu T^k q) G_\mu^k - \frac{1}{4} G_{\mu\nu}^k G^{\mu\nu k} \quad (2.5.22)$$

where the first term is the free Lagrangian, the second one describes the interaction between quarks and gluons and the final ones are the term for free gluon motion and the gluon self interaction.

2.6 The Mass Generation Mechanism

In Section 2.4 we have generated massless gauge bosons for the electroweak interaction, in fact no term such as $M^2 B_\mu B^\mu / 2$ appear in the Lagrangian of Eq. (2.4.3). These kind of terms are not gauge invariant and lead to divergences that would prevent using the theory to make any prediction. Nevertheless the weak gauge bosons $W^\pm$ and Z^0 must be massive in order to account for the short range of the weak interaction. A gauge invariant way to recover both the fermion and boson masses, is to spontaneously break the local $SU(2)_L \times U(1)_Y$ electroweak symmetry. Spontaneous breaking of a symmetry occurs in different phenomena in nature, for example when a bowl of water freezes the ice crystals form in a particular direction which cannot be predicted and that breaks the full rotational symmetry that the system originally had. In the case of a ferromagnet at high temperature the spin orientations are not well defined and globally (at the macroscopic level) this disorder gives a rotational symmetry since, for example, the value of the magnetization vanish. If it is cooled down below a critical temperature it may acquire a magnetization in two possible directions but there is no way to tell in advance which one breaking in this way the rotational invariance of the ferromagnet for global rotations. In Section 2.6.1 we will see the case of the spontaneous breaking of a global gauge symmetry. For a more detailed explanation of the mechanism see for example [28].

2.6.1 Global Symmetry Spontaneous Breaking

Let us consider the Lagrangian for a complex scalar field ϕ :

$$\mathcal{L} = (\partial_\mu \phi^*)(\partial_\mu \phi) - \underbrace{\mu^2 \phi^* \phi - \lambda(\phi^* \phi)^2}_{V(\phi^* \phi)} \quad (2.6.1)$$

where:

$$\begin{aligned} \phi &= \frac{\phi_1 + i\phi_2}{\sqrt{2}} \\ \phi^* &= \frac{\phi_1 - i\phi_2}{\sqrt{2}} \end{aligned} \quad (2.6.2)$$

and for the Lagrangian we thus get:

$$\mathcal{L} = \frac{1}{2}(\partial_\mu \phi_1)^2 + \frac{1}{2}(\partial_\mu \phi_2)^2 - \frac{\mu^2}{2}(\phi_1^2 + \phi_2^2) - \frac{\lambda}{4}(\phi_1^2 + \phi_2^2)^2. \quad (2.6.3)$$

This Lagrangian is invariant for global transformations ($U(1)$ symmetry group) of the form:

$$\begin{cases} \phi(x) \rightarrow \phi'(x) = e^{i\epsilon} \phi(x) \\ \phi^*(x) \rightarrow \phi'^*(x) = e^{-i\epsilon} \phi^*(x). \end{cases} \quad (2.6.4)$$

There are two possible choices for the potential:

A - $\mu^2 > 0$, which gives a stable configuration around $|\phi| = 0$ with a unique minimum.

B - $\mu^2 < 0$, which gives a local maximum for $\phi = 0$ and a circle of degenerate minima such that $\phi_1^2 + \phi_2^2 = v^2$, with $v^2 = -\mu^2/\lambda$.

In case A the potential has the shape of a parabola and if we perform a power expansion around $\phi = \phi^* = 0$ we obtain a massive particle ($m = \mu^2$) with ± 1 charge with an interaction given by the ϕ^4 term and the Lagrangian is invariant for global U(1) phase rotations. In the B case the potential has the Mexican hat shape where all the different states corresponding to the same minimum energy have different orientations in the complex plane thus breaking the invariance under U(1) rotations, in fact:

$$\phi_0 = \langle 0|\phi|0\rangle \rightarrow \frac{v}{\sqrt{2}}e^{i\alpha} \quad \text{if} \quad \phi \rightarrow e^{i\alpha}\phi. \quad (2.6.5)$$

In order to explore the physics content of this theory we make a perturbative expansion around one particular vacuum state, we chose one for example $\alpha = 0$ for which $\phi_1 = v$ and $\phi_2 = 0$ and introduce the two perturbations $\eta(x)$ and $\xi(x)$ so that:

$$\phi(x) = \frac{1}{\sqrt{2}} \left[\overbrace{v + \xi(x)}^{\phi_1} + i \overbrace{\eta(x)}^{\phi_2} \right] \quad \eta, \xi \text{ real} \quad (2.6.6)$$

and plug them in the Lagrangian of Eq. (2.6.3) we obtain that:

$$\begin{aligned} \mathcal{L}'(\xi, \eta) = & \frac{1}{2}(\partial_\mu \xi)^2 + \frac{1}{2}(\partial_\mu \eta)^2 - \frac{1}{2}(-2\mu^2)\eta^2 \\ & - \lambda v(\eta^2 + \xi^2)\eta - \frac{1}{4}(\eta^2 + \xi^2)^4 + \dots \end{aligned} \quad (2.6.7)$$

and as we can see, the third term looks like a mass term so that the field η has mass $m_\eta^2 = -2\mu^2$ while we have no mass term for the field ξ . Thus by choosing a particular vacuum state among the infinite degenerate ones we lost the original U(1) invariance (we thus *broke the symmetry*) and recovered the mass for the boson η but generated an additional massless field ξ . The generation of a massless scalar (spin zero) boson, known as *Goldstone boson*, for each generator of the symmetry is a general property formulated in the *Goldstone theorem* which was proved by Goldstone, Salam and Weinberg in 1962 [29]. Since no massless scalar particle were observed a new mechanism was introduced in 1964 to generate the mass for the gauge bosons without the presence of the Goldstone bosons. This is briefly treated in 2.6.2.

2.6.2 The Higgs Mechanism

In 1964 Englert and Brout [30], Guralnik, Hagen and Kibble [31] and Higgs [32] independently published a paper where they introduced a way to give mass to the gauge bosons and to get rid of the Goldstone fields. It involves the spontaneous breaking of *local* gauge symmetries. In order to see how the mechanism works, let us consider the Lagrangian:

$$\mathcal{L} = (\partial_\mu \phi)^\dagger (\partial^\mu \phi) - \mu^2 \phi^\dagger \phi - \lambda (\phi^\dagger \phi)^2, \quad (2.6.8)$$

where ϕ is a doublet of complex scalar fields with four degrees of freedom:

$$\phi \equiv \begin{pmatrix} \phi_i \\ \phi_j \end{pmatrix} = \frac{1}{\sqrt{2}} \begin{pmatrix} \phi_1 + i\phi_2 \\ \phi_3 + i\phi_4 \end{pmatrix}. \quad (2.6.9)$$

The Lagrangian is invariant under global SU(2) phase transformations:

$$\phi(x) \rightarrow \phi'(x) = e^{i \sum_{k=1}^3 \epsilon^k T^k} \phi(x), \quad (2.6.10)$$

where $T^k = \tau^k/2$ are the generators and $[T^i, T^j] = i\epsilon^{ijk}T^k$ with $i, j, k = 1, 2, 3$ are the commutation relations. To achieve SU(2) invariance of the Lagrangian for local transformations, that is when the ϵ_k parameter becomes a function of the space-time coordinates ($\epsilon_k \equiv \epsilon_k(x^\mu)$), we need to introduce the covariant derivative:

$$D_\mu = \partial_\mu + ig \frac{\vec{\tau}}{2} \cdot \vec{W}_\mu(x) \quad (2.6.11)$$

where $\vec{W}_\mu(x) = (W^0, W^1, W^2)$ are three gauge fields. For simplicity we can consider infinitesimal transformations of the fields:

$$\phi(x) \rightarrow \phi'(x) \simeq \left(1 + i\vec{\epsilon}(x) \cdot \frac{\vec{\tau}}{2}\right) \phi(x), \quad (2.6.12)$$

for the gauge bosons transformations we have:

$$\vec{W}_\mu(x) \rightarrow \vec{W}_\mu(x) - \frac{1}{g} \partial_\mu \vec{\epsilon}(x) - \vec{\epsilon}(x) \times \vec{W}_\mu(x). \quad (2.6.13)$$

The gauge invariant expression of the Lagrangian is then:

$$\mathcal{L} = \left(\partial_\mu \phi + ig \frac{\vec{\tau}}{2} \cdot \vec{W}_\mu \phi \right)^\dagger \left(\partial_\mu \phi + ig \frac{\vec{\tau}}{2} \cdot \vec{W}_\mu \right) - V(\phi) - \frac{1}{4} \vec{W}_{\mu\nu} \cdot \vec{W}^{\mu\nu} \quad (2.6.14)$$

where the potential is given by:

$$V(\phi) = \mu^2 \phi^\dagger \phi + \lambda (\phi^\dagger \phi)^2 \quad (2.6.15)$$

and in the kinetic term $\vec{W}_{\mu\nu} \cdot \vec{W}^{\mu\nu}/4$ we have introduced:

$$\vec{W}_{\mu\nu} = \partial_\mu \vec{W}_\nu - \partial_\nu \vec{W}_\mu - g \vec{W}_\mu \times \vec{W}_\nu. \quad (2.6.16)$$

Also in this case there are two possibilities for the potential in Eq. (2.6.15), if $\mu^2 > 0$ the Lagrangian of Eq. (2.6.14) describes four scalar particles, the ϕ_i with $i = 1, \dots, 4$ of Eq. (2.6.9), with mass μ that interacts with the three massless gauge bosons $\vec{W}$. We are interested in the second case where $\mu^2 < 0$ and $\lambda > 0$, the minimum of the potential is at:

$$\phi^\dagger \phi = \frac{1}{2} (\phi_1^2 + \phi_2^2 + \phi_3^2 + \phi_4^2) = -\frac{\mu^2}{2\lambda} \quad (2.6.17)$$

and to study the particle spectrum we have to expand the field $\phi(x)$ around a minimum, we can choose one of them, for example:

$$\phi_0 \equiv \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v \end{pmatrix} \quad (2.6.18)$$

that is $\phi_1 = \phi_2 = \phi_4 = 0$ and $\phi_3^2 = -\mu^2/\lambda \equiv v^2$. Also in this case any other vacuum choice is related to ϕ_0 by a global phase transformation and thus ϕ_0 is not invariant under $SU(2)$ transformations and the symmetry is spontaneously broken. Making a small perturbation around this particular minimum we have:

$$\phi(x) = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + h(x) \end{pmatrix} \quad (2.6.19)$$

where $h(x)$ is a scalar field called the *Higgs field*. The fact that only one field out of the initial four remains can be better understood by parameterizing the fluctuations around the vacuum ϕ_0 using four real fields θ_i with $i = 1, 2, 3$ and h :

$$\phi(x) = \frac{1}{\sqrt{2}} e^{i \frac{\vec{\tau} \cdot \vec{\theta}}{v}} \begin{pmatrix} 0 \\ v + h(x) \end{pmatrix} \quad (2.6.20)$$

and expanding at the lowest order in $1/v$:

$$\begin{aligned} \phi(x) &= \frac{1}{\sqrt{2}} \begin{pmatrix} 1 + i\theta_3/v & i(\theta_1 - i\theta_2)/v \\ i(\theta_1 + i\theta_2)/v & 1 - \theta_3/v \end{pmatrix} \begin{pmatrix} 0 \\ v + h(x) \end{pmatrix} \\ &\sim \frac{1}{\sqrt{2}} \begin{pmatrix} \theta_2 + i\theta_1 \\ v + h - i\theta_3 \end{pmatrix}. \end{aligned} \quad (2.6.21)$$

We see that these four fields are independent and fully parametrize the perturbation of the vacuum ϕ_0 , θ_i with $i = 1, 2, 3$ are the Goldstone bosons that can be absorbed with a local $SU(2)$ gauge transformation of the form given in Eq. (2.6.10) where the functions $\epsilon^k \equiv \epsilon^k(x)$ are chosen accordingly, the θ_1 , θ_2 and θ_3 fields disappear from the Lagrangian and the parameterization of Eq. (2.6.19) is found.

To generate the masses for the three gauge bosons W^1 , W^2 and W^3 , we substitute the vacuum expectation value ϕ_0 in the Lagrangian of Eq. (2.6.14). For the mass term we have:

$$\begin{aligned} \left| ig \frac{\vec{\tau}}{2} \cdot \vec{W}_\mu \phi_0 \right|^2 &= \\ &= \frac{g^2}{2} \left| \begin{pmatrix} W^3 & W^1 - iW^2 \\ W^1 + iW^2 & -W^3 \end{pmatrix} \begin{pmatrix} 0 \\ v \end{pmatrix} \right|^2 = \\ &= \frac{g^2 v^2}{8} (W^1 + iW^2, -W^3) \begin{pmatrix} W^1 - iW^2 \\ -W^3 \end{pmatrix} = \\ &= \frac{g^2 v^2}{8} \left[(W_\mu^1)^2 + (W_\mu^2)^2 + (W_\mu^3)^2 \right]. \end{aligned} \quad (2.6.22)$$

By comparing this result with a general mass term for a boson $M^2 W_\mu^2/2$ we see that the three vector gauge bosons gain a mass $M = \frac{gv}{2}$. The Lagrangian thus describes three massive vector bosons where the three spin degrees of freedom are a consequence of the gauge fixing used to absorb the Goldstone bosons and a massive scalar Higgs field with mass $m_H = 2v^2\lambda$. This mechanism can be used to generate the masses of the electroweak gauge bosons preserving the renormalizability of the theory.

2.6.3 Masses for the $W^\pm$ and Z^0 Gauge Bosons

In order to apply the Higgs mechanism introduced in Section 2.6.2 to generate the masses of the $W^\pm$ and Z^0 weak bosons keeping the photon massless and preserving the renormalizability of the electroweak theory, we add to the Lagrangian of Eq. (2.4.3) the Higgs Lagrangian:

$$\mathcal{L}_\phi = \left| \left(\partial_\mu + ig \vec{T} \cdot \vec{W}_\mu + i \frac{g'}{2} Y B_\mu \right) \phi \right|^2 - V(\phi) \quad (2.6.23)$$

with $V(\phi) = \mu^2 \phi^\dagger \phi + \lambda (\phi^\dagger \phi)^2$ and $\lambda > 0$. To preserve $SU(2)_L \times U(1)_Y$ gauge invariance for this Lagrangian we choose four fields to be arranged in an isospin doublet with hypercharge $Y = 1$:

$$\phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix} = \frac{1}{\sqrt{2}} \begin{pmatrix} \phi_1 + i\phi_2 \\ \phi_3 + i\phi_4 \end{pmatrix}. \quad (2.6.24)$$

To generate the boson masses we consider the case where in the potential $V(\phi)$ we have $\mu^2 < 0$ and choose a particular vacuum expectation value of ϕ :

$$\phi_0 \equiv \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v \end{pmatrix}. \quad (2.6.25)$$

With this choice the $SU(2)_L \times U(1)_Y$ symmetry is broken, in fact the vacuum is no more invariant under transformations of this group, for example:

$$T_3 |\phi_0\rangle = \frac{1}{2} \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix} \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v \end{pmatrix} = -\frac{1}{2\sqrt{2}} \begin{pmatrix} 0 \\ v \end{pmatrix} \quad (2.6.26)$$

and we get that:

$$\phi_0 \rightarrow \phi'_0 = e^{i\epsilon^3 T_3} \phi_0 \simeq (1 + \epsilon^3 T_3) \phi_0 \neq \phi_0 \quad (2.6.27)$$

and thus the three weak gauge bosons will acquire a mass through the Higgs mechanism. We can also see that:

$$\begin{aligned} Q |\phi_0\rangle &= \left(T_3 + \frac{Y}{2} \right) |\phi_0\rangle = \left[\begin{pmatrix} 1/2 & 0 \\ 0 & -1/2 \end{pmatrix} + \begin{pmatrix} 1/2 & 0 \\ 0 & 1/2 \end{pmatrix} \right] \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v \end{pmatrix} \\ &= \begin{pmatrix} 1 & 0 \\ 0 & 0 \end{pmatrix} \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v \end{pmatrix} = 0 \end{aligned} \quad (2.6.28)$$

and the vacuum is invariant for $U(1)_{\text{em}}$ transformations:

$$\phi_0 \rightarrow \phi'_0 = e^{i\epsilon(x)Q} \phi_0 = \phi_0 \quad \forall \epsilon(x) \quad (2.6.29)$$

and this invariance ensures that the photon will not gain a mass through the Higgs mechanism ($m_\gamma = 0$). To summarize, the choice of a particular vacuum expectation value ϕ_0 breaks the $SU(2)_L \times U(1)_Y$ symmetry and by applying the Higgs mechanism the masses of the weak gauge bosons are generated but since ϕ_0 is invariant under transformations of the $U(1)_{\text{em}}$ subgroup the photon remains massless.

As seen in Section 2.6.1 the gauge bosons masses are obtained simply substituting the vacuum expectation value ϕ_0 in the Lagrangian:

$$\phi(x) \rightarrow \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + h(x) \end{pmatrix}, \quad (2.6.30)$$

for the mass term we have:

$$\begin{aligned} \left| \left(g \frac{\vec{\tau}}{2} \cdot \vec{W}_\mu + \frac{g'}{2} B_\mu \right) \phi \right|^2 &= \\ &= \frac{1}{8} \left| \begin{pmatrix} gW_\mu^3 + g'B_\mu & g(W_\mu^1 - iW_\mu^2) \\ g(W_\mu^1 + iW_\mu^2) & -gW_\mu^3 + g'B_\mu \end{pmatrix} \begin{pmatrix} 0 \\ v \end{pmatrix} \right|^2 \\ &= \frac{1}{8} v^2 g^2 \left[(W_\mu^1)^2 + (W_\mu^2)^2 \right] + \frac{1}{8} v^2 (g'B_\mu - gW_\mu^3) (g'B_\mu - gW_\mu^3) \\ &= \left(\frac{1}{2} gv \right)^2 W_\mu^+ W^{-\mu} + \frac{1}{8} v^2 \begin{pmatrix} W_\mu^3 & B_\mu \end{pmatrix} \begin{pmatrix} g^2 & -gg' \\ -gg' & g'^2 \end{pmatrix} \begin{pmatrix} W_\mu^3 \\ B_\mu \end{pmatrix} \end{aligned} \quad (2.6.31)$$

where we used $W^\pm = (W^1 \mp iW^2)/\sqrt{2}$. The first term of the equation is similar to the expected mass term for a charged boson $M^2 W_\mu^+ W^{-\mu}$ and by comparison we conclude that:

$$M_W = \frac{1}{2} gv. \quad (2.6.32)$$

The last term of Eq. (2.6.31) is off diagonal in the (W_μ^3, B_μ) basis but can be diagonalized solving the eigenvalue equation in terms of the eigenstates Z_μ and A_μ :

$$\begin{aligned} A^\mu &= \frac{g'W_\mu^3 + gB_\mu}{\sqrt{g^2 + g'^2}} \\ Z^\mu &= \frac{gW_\mu^3 - g'B_\mu}{\sqrt{g^2 + g'^2}}, \end{aligned} \quad (2.6.33)$$

and we can write:

$$\begin{aligned} \frac{1}{8} v^2 \left[g^2 (W_\mu^3)^2 - 2gg'W_\mu^3 B_\mu + g'^2 B_\mu^2 \right] &= \\ &= \frac{v^2}{8} [gW_\mu^3 - gB_\mu]^2 + 0 [g'W_\mu^3 - g'B_\mu]^2 = \\ &= \frac{v^2}{8} (g^2 + g'^2) Z_\mu Z^\mu + 0 (g^2 + g'^2) A_\mu A^\mu. \end{aligned} \quad (2.6.34)$$

The general mass term for neutral fields can be written as:

$$\frac{1}{2} M_Z^2 Z_\mu Z^\mu + \frac{1}{2} M_A^2 A_\mu A^\mu \quad (2.6.35)$$

and comparing this with Eq. (2.6.34) we have that:

$$\begin{aligned} M_A &= 0 \\ M_Z &= \frac{1}{2} v \sqrt{g^2 + g'^2}. \end{aligned} \quad (2.6.36)$$

which are the mass term for the photon and for the neutral weak vector boson fields.

After resolving the mass term we can explore the form of the Lagrangian in Eq. (2.6.23) after the substitution of Eq. (2.6.30):

$$\mathcal{L}_\phi = \frac{1}{2} (\partial_\mu h)^2 + \frac{g^2}{4} (v+h)^2 W_\mu^+ W^{-\mu} + \frac{1}{8} (v+h)^2 Z_\mu Z^\mu - V(\phi), \quad (2.6.37)$$

the terms:

$$\frac{g^2 v}{2} W_\mu^+ W^{-\mu} h \quad \text{and} \quad \frac{g^2 v}{4} Z_\mu Z^\mu h \quad (2.6.38)$$

show that the coupling between the Higgs and the weak gauge fields is proportional to their mass while the Higgs quadratic terms (h^2) represents the interaction of two Higgs fields with two vector bosons. For the potential $V(\phi)$ we have:

$$V(\phi) = \frac{1}{2} (-2\mu^2) h^2 + \lambda v h^3 + \frac{\lambda}{4} h^4 \quad (2.6.39)$$

and the first term describes the mass of the Higgs boson:

$$m_h = \sqrt{2\mu^2} \quad (2.6.40)$$

Since the value of μ^2 is not known the mass of the Higgs boson cannot be predicted and must be inferred from experiments. In 2012 the ATLAS [15] and CMS [16] experiments at CERN discovered the Higgs boson needed to complete the Standard Model of weak interactions with a mass of $m_h = 125.36$ GeV.

2.7 Experimental Confirmation of the Standard Model

The Standard Model is a very well tested theory capable of giving precise predictions on particles and their interactions. Figure 2.1 gives the comparison between experimental measurements performed by the ATLAS collaboration and theoretical predictions for several Standard Model production cross sections, in general the predictions agree with the measured data. Figure 2.2 reports the measured signal strength for a Higgs boson mass of 125.36 GeV normalized to the Standard Model expectations in $H \rightarrow \gamma\gamma$, $H \rightarrow ZZ^* \rightarrow \ell\ell\ell\ell$, $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$, $H \rightarrow \tau\tau$ and $H \rightarrow b\bar{b}$ final states, as can be seen the ATLAS collaboration has, within uncertainties, sensitivity in most of these searches.

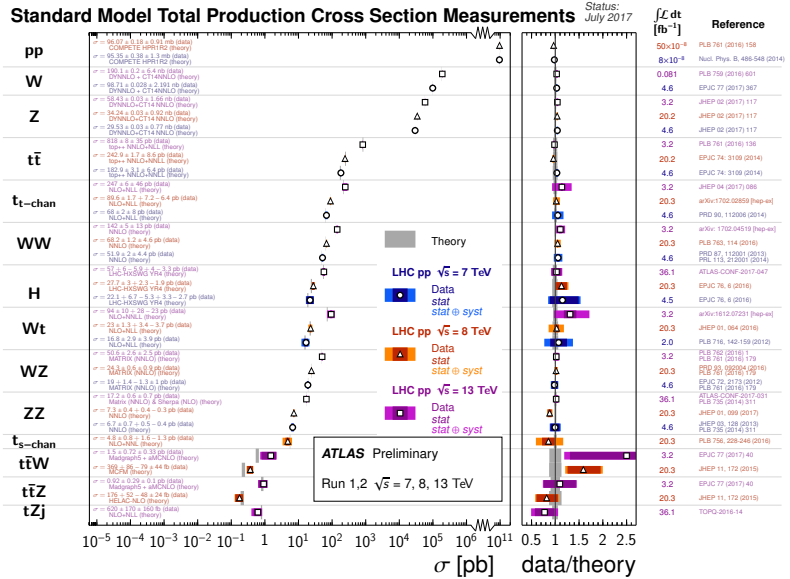


Figure 2.1: Comparison between experimental measurements and theoretical predictions for several Standard Model production cross sections corrected for leptonic branching fractions. The statistical uncertainty is reported in the plot with a dark-colored error bar. The light-colored error band represents the full uncertainty including statistics, systematics and luminosity uncertainties. The reference for each measurement as long as the ratio between data and theoretical prediction are also reported in the plot [33].

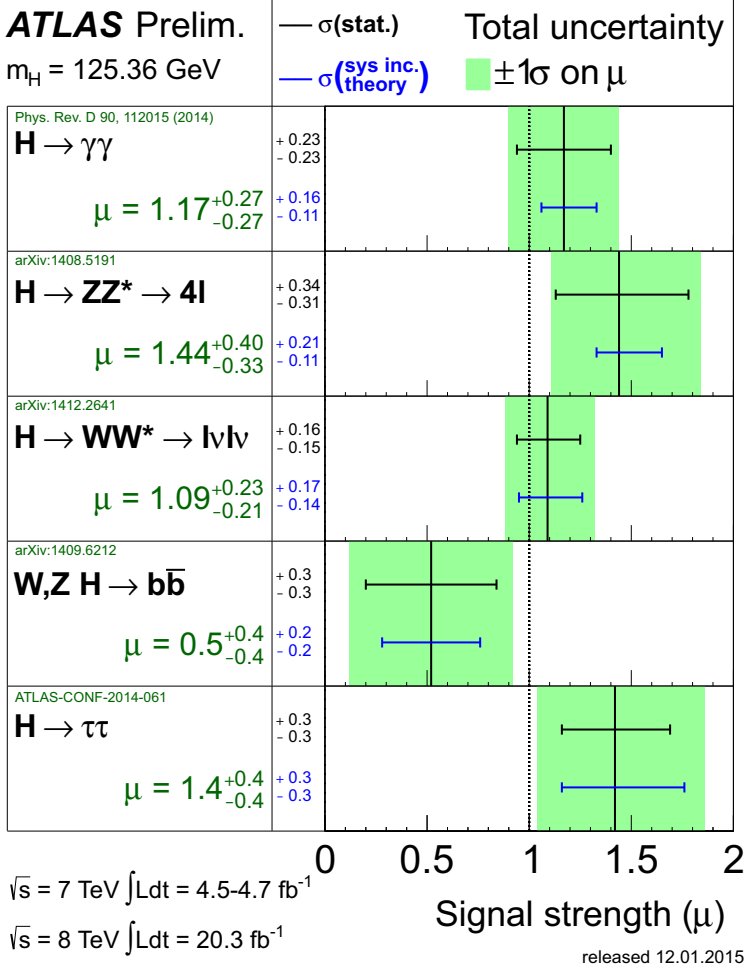


Figure 2.2: Measured signal strength for a Higgs boson mass of 125.36 GeV normalized to the Standard Model expectations in $H \rightarrow \gamma\gamma$, $H \rightarrow ZZ^* \rightarrow llll$, $H \rightarrow WW^* \rightarrow l\nu l\nu$, $H \rightarrow \tau\tau$ and $H \rightarrow b\bar{b}$ final states. The solid vertical lines represent the best fitted value, the shaded band is the total $\pm 1\sigma$ uncertainty. The statistical and total (experimental and theoretical) individual uncertainties are also reported [34].

Chapter 3

Beyond the Standard Model

3.1 Open Questions of the Standard Model

Despite its great success, the Standard Model only explains three of the four fundamental interactions, failing to incorporate gravity. There are also questions that are not answered by the theory, two of these are briefly explained in the sections below, namely: the *hierarchy problem* and *dark matter*. These questions motivate the introduction of Beyond Standard Model (BSM) theories.

3.1.1 The Hierarchy Problem and Naturalness

The *naturalness criterion* states that: “one such [dimensionless and measured in units of the cut-off (Λ)] parameter is allowed to be much smaller than unity only if setting it to zero increases the symmetry of the theory. If this does not happen, the theory is unnatural” [35].

One important concept in physics that enter in the formulation of the naturalness principle is that of symmetries. *Symmetries* are closely connected to conservation laws through the Noether’s theorem, moreover theory parameters that are protected by a symmetry, if smaller than the unit, are not problematic according to the naturalness criterion.

Let us consider the strength of the gravitational force, characterized by the Newton’s constant, G_N and the weak force, characterized by the Fermi’s constant G_F , if we take the ratio of these we get:

$$\frac{G_F \hbar^2}{G_N c^2} = 1.738 \times 10^{33}. \quad (3.1.1)$$

The reason why this number is worth some attention is that theory parameters close to the order of unit in the SM, may be calculated in a more fundamental theory, if any, using constants like π or e while numbers that deviates from

one, may not have such a simple mathematical expression and thus may lead to uncover new properties of the fundamental theory.

This number becomes even more interesting if we consider quantum effects. *Virtual particles* are off-shell ($E \neq m^2 + p^2$) and according to the *uncertainty principle*, $\Delta t \Delta E \geq \hbar/2$, can appear out of the vacuum for a short time that depends on the energy of the virtual particle; according to quantum field theory, the vacuum is populated with virtual particles. The Higgs field has the property to couple with other SM particles with a strength proportional to their mass. All these virtual particles have a mass determined by the cut-off energy Λ and when the Higgs field travels through space, couples with these virtual particles and its mass squared gets a contribution proportional to Λ (see [36]):

$$\delta m_H^2 = k\Lambda^2, \text{ with } k = \frac{3G_F}{4\sqrt{2}\pi^2}(4m_t^2 - 2m_W^2 - m_Z^2 - m_H^2). \quad (3.1.2)$$

Since $k \approx 10^{-2}$ [36], the value of Higgs mass $m_H (\sim G_F^{-1/2})$, should be close to the maximum energy scale Λ and if we assume this to be the Planck scale $M_{Pl} = G_N^{-1/2}$, the ratio G_F/G_N , should be close to the unity which contradicts Eq. (3.1.1), this goes by the name of *hierarchy problem*.

The large quantum corrections in Eq. (3.1.2) are mainly due to the fact that in the SM, there is no symmetry protecting the mass of the Higgs field.

3.1.2 Dark Matter

Observations on galaxies rotation suggests that their mass is not enough to generate the gravitational force needed to counteract the effect of the centripetal force and to prevent them to tore apart [37]. Since observation also suggest that stars and planets in galaxies are indeed kept together by the gravitational attraction, we conclude that there must be extra mass that we cannot see that generate the gravity needed to hold them together. This unknown matter goes by the name of *dark matter* and is thought to be made of Weakly Interacting Massive Particles (WIMPs). These particles do not interact or interact weakly with the electromagnetic force, as a consequence, dark matter does not emit, reflect or absorb light making it hard to detect. Cosmological measurements on the content of matter of the universe have been performed [38] that seems to suggest that the visible matter (the one we can see and measure), only accounts for 5% of the total energy of the universe and that dark matter constitutes roughly the 27%, the remaining 68% is called *dark energy* and is believed to be responsible for the accelerated expansion of the universe. Some theories [39] predict that dark matter particles should be light enough to be produced at hadron colliders but due to their “dark” nature, they would escape detection leading to an energy imbalance in the detector that could be used as an hint of their existence.

3.2 Supersymmetry

One possible solution to the dark matter and the hierarchy problem is achieved introducing a symmetry, called SUSY [2], that relates fermions and bosons. Supersymmetry is capable of solving the hierarchy problem by canceling out the quantum corrections that bring m_H close to Λ thus restoring the naturalness of the SM. SUSY introduces a set of new particles that are the *superpartners* of the SM ones. These supersymmetric particles or *sparticles* have the same quantum numbers and couplings of their SM counterpart but turn with a supersymmetry transformation a spin 1/2 fermionic SM particle into a bosonic spin 0 sparticle and a spin 1 SM boson in a spin 1/2 fermion while a scalar particle becomes a spin 1/2 sparticle. The superpartners of the SM leptons are called *sleptons*, adding the “s” prefix (that stands for “scalar”) to their SM fermionic partners and are denoted with a “~” on top of their symbol, so for example the SM electrons are called *selectrons* and denoted with $\tilde{e}_L$ and $\tilde{e}_R$ where the R and L does not refer to their helicity (being 0-spin bosons) but to that of their SM superpartners. Similarly for *smuons* and *staus*: $\tilde{\mu}_L, \tilde{\mu}_R, \tilde{\tau}_L, \tilde{\tau}_R$ and for the *squarks*: $\tilde{q}_L, \tilde{q}_R$ (with $q = u, d, s, b, c, t$). The supersymmetric partners of the gauge bosons are called *gauginos*, adding the “ino” suffix to their SM partners, for instance, the W and Z bosons have the *wino* ($\tilde{W}$) and *zino* ($\tilde{Z}$) superpartners, the photon has the *photino* ($\tilde{\gamma}$) and the Higgs field has the *Higgsino* ($\tilde{H}$). In addition to these new particles, the minimal supersymmetric extension of the Standard Model, called Minimal Supersymmetric Standard Model (MSSM) [40], introduces three neutral and two charged Higgs bosons leading to five physical Higgs states. The neutral superpartners of these Higgs states ($\tilde{H}^0$), combine with the bino and the wino ($\tilde{B}, \tilde{W}^0$) in four new electrically neutral mass eigenstates called *neutralinos* ($\tilde{\chi}_i^0$ with $i = 1, 2, 3, 4$). The charged higgsinos ($\tilde{H}^+, \tilde{H}^-$) and the winos ($\tilde{W}^+, \tilde{W}^-$) mix and give rise to two electrically charged mass eigenstates of charge ± 1 called *charginos* ($\tilde{\chi}_i^\pm$ with $i = 1, 2$). In the MSSM framework a new multiplicative quantum number can be introduced, the R-parity, that, if it is conserved, explains the stability of the proton and is defined as:

$$R = (-1)^{3(B-L)+2S} \quad (3.2.1)$$

where B and L are the baryon and lepton numbers respectively and S is the spin. It can be seen that all SM particles have $R = +1$ while supersymmetric particles have $R = -1$. The conservation of R-parity has important phenomenological consequences:

- The Lightest Supersymmetric Particle (LSP) is stable.
- The final decay products of SUSY particles are an odd number.
- In collider experiments sparticles can only be produced in pairs.

If the LSP is stable and electrically neutral (dark) and interacts only weakly with matter it makes a good candidate for non-baryonic dark matter [41]. In the MSSM, good LSP candidates are either the neutralino or the gravitino, the

supersymmetric partner of the graviton (the hypothetical gauge boson associated to gravitation).

No sparticle has been observed at particle colliders so far, hinting that they must have a larger mass than their SM counterparts. If they had the same mass they would have already been produced at particle colliders, this imply that SUSY is a *broken symmetry*. Detecting one of these supersymmetric particles would shed light on some of the SM shortcomings and thus many searches at the current colliders focus their attention on finding SUSY particles. It is possible to estimate, from Eq. (3.1.2), the scale at which new physics is expected. Using $m_H = 125$ GeV [42], we get that $\Lambda \approx 1$ TeV. If the naturalness criterion holds, we thus expect the two main experiments at LHC, ATLAS and CMS, to find signal for new physics at the TeV scale.

3.2.1 Status of the Experimental Search for Supersymmetry

As mentioned earlier, supersymmetry could solve the hierarchy problem and provide a dark matter candidate, for this reason it has been the subject of great experimental attention. Since the start of the LHC the ATLAS and CMS experiments have performed a large number of searches for SUSY in many possible final states. Figure 3.1 illustrates the large number of searches performed by the ATLAS experiment, it can be seen in the first box that the squark and the gluino production reaches the highest masses.

In the channel pursued in this thesis where two squarks are produced together and decay directly to neutralinos, the 2-6 jets analysis excludes squarks up to 1.6 TeV. The $\tilde{q} \rightarrow q + \tilde{\chi}_1^0$ ($q = u, d, c, s$) (compressed) search listed in Figure 3.1 is the one presented in Chapter 7 of this thesis. Other searches focus on third generation squarks, the stop and sbottom where exclusions up to about 900 GeV can be reached and on searches for production of SUSY gauge bosons via the electroweak interaction. In this situation the production cross sections are much lower and the mass limits on these particles are generally below 1 TeV.

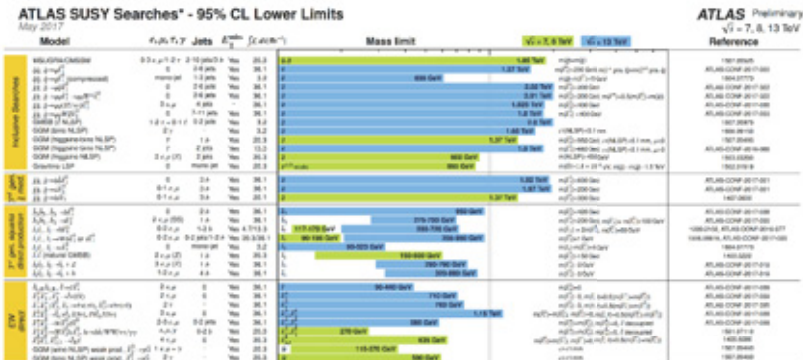


Figure 3.1: Representative selection of the mass reach in searches for supersymmetry with the ATLAS detector [43].

3.3 Large Extra Dimensions

In 1919 Theodor Kaluza proposed that General Relativity (GR) in $4 + 1$ dimensions could describe both, gravity and electromagnetism in a unified geometrical description [44]. In order to account for the four dimensional character of the space-time the *cylindrical conditions* are introduced: in no reference frame the metric and the first four coordinates depend on the fifth dimension. In this context it is possible to express the metric tensor as:

$$g_{MN} = \left(\begin{array}{c|c} g_{\mu\nu} + \phi A_\mu A_\nu & \phi A_\nu \\ \hline \phi A_\nu & g_{55} \equiv \phi \end{array} \right) \quad (3.3.1)$$

where ϕ is a scalar field and A^μ is a gauge vector that was proposed to be identified with the electromagnetic vector potential. The line element is given by:

$$ds_{5D}^2 = g_{\mu\nu} dx^\mu dx^\nu + \phi (A_\nu dx^\nu + dx^5)^2 \quad (3.3.2)$$

and the invariant volume element is:

$$d^5x \sqrt{-\det(g_{\mu\nu})} = d^4x dx^5 \phi^{1/2} \sqrt{-\det(g_{\mu\nu}^{4D})}. \quad (3.3.3)$$

With this notation it is possible to write the GR Hilbert-Einstein action in five dimensions as:

$$S^{5D} = -\frac{1}{16\pi G_5} \int d^5x \sqrt{-\det(g_{MN})} R^{5D} \quad (3.3.4)$$

where G_5 is the gravitational constant in five dimensions and R^{5D} is the Ricci scalar. Since the other coordinates do not depend on the fifth, the integration on the latter can be factorized:

$$S^{5D} = -\frac{1}{16\pi G_5} \int dx^5 \int d^4x \sqrt{-\det(g_{MN})} R^{5D}. \quad (3.3.5)$$

The value of the integral in Eq. (3.3.5) depends on the value of x^5 thus for it to have a finite value, x^5 must be an interval or, how it was proposed by Oskar Klein in 1926 a circle [45]. Some of the consequences of the Kaluza-Klein (KK) theory are briefly outlined in the next section.

3.3.1 Consequences of the Kaluza-Klein Theory

One important fact that naturally arises from the theory developed by Kaluza is the charge conservation; to better illustrate this, consider the momentum in five dimensions:

$$P_A \equiv \frac{\partial \mathcal{L}}{\partial \dot{x}^A} \quad (3.3.6)$$

where

$$\mathcal{L} = \frac{1}{2} m g_{MN} \dot{x}^M \dot{x}^N \quad (3.3.7)$$

is the lagrangian of a free particle. Since $\partial_5 g_{MN} = 0$ then $\partial \mathcal{L} / \partial \dot{x}^5 = 0$ and then P_5 is a constant of motion. Solving for the remaining four components leads to:

$$P_\mu = m g_{\mu\nu} \dot{x}^\nu + P_5 A_\mu \quad (3.3.8)$$

where the first term is the momentum in four dimensions and, identifying P_5 with the electric charge ($P_5 \equiv e$), the expression of the canonically conjugate momentum is recovered. Thus, writing $P_M \equiv (P_0 \equiv E, \vec{P}, P_5 \equiv e)$, the charge conservation is a consequence of the space-time geometry [46].

Another remarkable result is obtained considering the geodesics equation (the motion of a free particle with mass m) in five dimensions:

$$\ddot{x}^M + \Gamma_{AB}^M \dot{x}^A \dot{x}^B = 0 \quad (3.3.9)$$

where Γ_{AB}^M are the connection coefficients. Partitioning the equation with $A = (\mu, 5)$ and solving separately for the usual four dimensions and for the fifth one gives:

$$\ddot{x}^\mu + \Gamma_{\alpha\beta}^\mu \dot{x}^\alpha \dot{x}^\beta = \frac{e}{m} g^{\mu\nu} F_{\nu\rho} \dot{x}^\rho \quad (3.3.10)$$

where $F_{\nu\rho}$ is the electromagnetic field strength tensor. Eq. (3.3.10) is the Lorentz equation for a charged particle of mass m propagating in a curved space-time which shows how gravity in five dimensions manifests itself in four dimensions as motion governed by electromagnetic and gravitational interaction [46].

Finally if we consider a mass-less scalar field $\Phi(x^\mu, x^5)$ that propagates in a space where the fifth dimension is a circle, as proposed by Klein, we can impose the periodic condition:

$$\Phi(x^\mu, x^5 = 0) = \Phi(x^\mu, x^5 = 2\pi R) \quad (3.3.11)$$

where R is the radius of the circle, to the Fourier expansion of the field ($\Phi(x^\mu, x^5) = \sum_{-\infty}^{+\infty} \Phi_n(x^\mu) e^{iP_5 x^5}$) it is possible to see that P_5 is quantized, in particular $P_5 = n/R$. Solving then the equation of motion $\square_5 \Phi(x^\mu, x^5) = 0$ where $\square = g^{MN} \nabla_M \nabla_N$ one finds in four dimensions infinite wave functions, one for each Fourier mode of $\Phi(x^\mu, x^5)$ where only the 0 mode describes a mass-less scalar field while the others, called Kaluza-Klein states, have a charge and a mass which are quantized by construction thanks to the compactification and periodicity of the fifth dimension.

3.3.2 The Arkani-Hamed, Dimopoulos, Dvali Model

An Effective Field Theory (EFT) is a tool used to describe physics phenomena of a complete theory at energies much lower than a cut-off scale ($E \ll \Lambda$). In this low energy limit it is possible to integrate out of the action the interactions above Λ while retaining their information in the couplings of the low energy Lagrangian [47].

The ADD model [3] is an effective theory for gravity where the cut-off scale is given by the electroweak scale m_{EW} instead of the Planck scale (M_{Pl}). In this context the strength of the gravitational interaction ($1/M_{Pl}$) is recovered postulating the existence of n extra compact spatial dimensions of radius R . The Planck scale of this $(4+n)$ dimensional space is taken to be approximately at the electroweak scale m_{EW} and denoted as M_D . In this framework, the gravitational flux lines for two masses, m_1 and m_2 , at a distance $r \ll R$ are

allowed to propagate to the extra-dimensions leading to a potential (given by Gauss' law):

$$V(r) \approx \frac{m_1 m_2}{M_D^{n+2}} \frac{1}{r^{n+1}}, \quad (3.3.12)$$

when $r \gg R$, the gravitational flux lines do not propagate to the extra-dimensions any more and the $1/r$ potential is retrieved:

$$V(r) \approx \frac{m_1 m_2}{M_D^{n+2} R^n} \frac{1}{r}. \quad (3.3.13)$$

The effective Planck scale is thus $M_{Pl} \approx M_D^{n+2} R^n$ where M_D^{n+2} is the scale of the fundamental complete theory in $4 + n$ dimensions. The compactification radius R can be estimated by recalling that in this framework $M_D \approx m_{EW}$ and by setting M_{Pl} to the observed value:

$$R \approx 10^{\frac{30}{n}-17} \text{ cm} \times \left(\frac{1 \text{ TeV}}{m_{EW}} \right)^{1+\frac{2}{n}}. \quad (3.3.14)$$

In the context of the ADD framework, the weakness of gravity is understood in terms of the graviton propagating in the extra-dimensions. The hierarchy problem is also solved by setting the cut-off scale and thus the order of the corrections to the Higgs mass close to the electroweak scale.

3.3.3 Status of the Experimental Search for ADD

Due to the high production cross section associated with processes involving the strong interaction, the monojet signature is a favored experimental signature in the search for large extra dimensions in the ADD scenario. In the previous version of this analysis with the ATLAS detector [7], upper limits on the effective Planck scale M_D have been set. Values of M_D up to 6.6 and 4.3 TeV for $n = 2$ and 6 dimensions respectively have been excluded. Figure 3.2 shows the expected and observed limits on M_D as a function of the number of extra dimensions in the ADD scenario carried out in the 3.2 fb^{-1} collected by the ATLAS detector in 2015 [7].

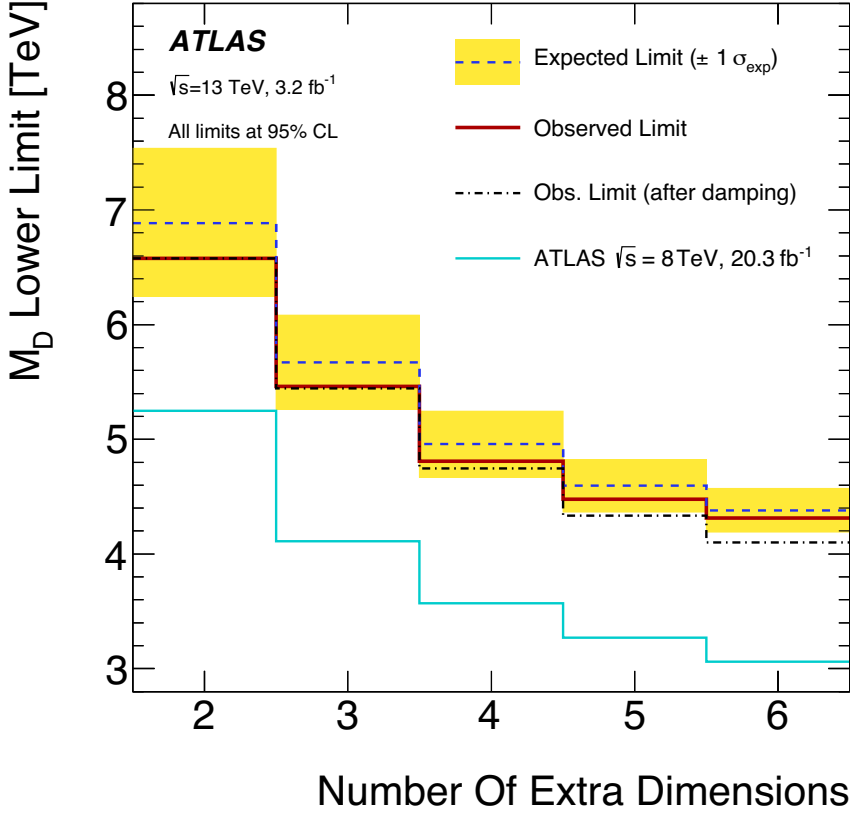


Figure 3.2: Observed and expected lower limits at 95% CL on the fundamental Planck scale in $4 + n$ dimensions, M_D , as a function of the number of extra dimensions. The dashed blue line is the expected limit using 3.2 fb^{-1} , the yellow band is the $\pm 1\sigma$ uncertainty on the estimate. The black dashed line shows the observed limit after the suppression of events with $\hat{s} > M_D^2$. Previous results are also reported for comparison [7].

Chapter 4

Experimental Apparatus

4.1 The Large Hadron Collider

The LHC [48] is a two ring superconducting hadron accelerator and collider located at the CERN.

The performance of a collider is evaluated in terms of its available *center of mass energy*, $\sqrt{s}$ and the *instantaneous luminosity* $\mathcal{L}$. The former defines the accessible energies for particle production. The latter is defined as the interaction rate per unit cross section of the colliding beams (collisions / (cm² s)).

The LHC is designed to operate at $\sqrt{s} = 14$ TeV in the center of mass although it started off at 7 TeV in 2010 and 2011, 8 TeV in 2012 and 13 TeV in 2015 and 2016 after the long shutdown in 2013 and 2014.

There are six experiments at LHC: ATLAS [49], CMS [50], ALICE [51], LHCb [52], LHCf [53] and TOTEM [54]. ATLAS and CMS are designed to work with the maximum luminosity that LHC can provide $\sim 10^{34}$ cm⁻² s⁻¹. This requirement, due to the low efficiency production, excludes the use of anti-proton beams and therefore the LHC is designed to be a proton-proton (pp) and heavy ions collider. The protons are organized in bunches, accelerated by LINAC2 to an energy of 50 MeV and subsequently injected in the Proton Synchrotron Booster (PSB). Here they are further accelerated to an energy of 1.4 GeV and fed to the Proton Synchrotron (PS) where they reach the energy of 25 GeV to be then passed to the Super Proton Synchrotron (SPS) which accelerate them to an energy of 450 GeV. They are finally injected in the LHC in opposite directions where they reach the nominal energy. There are four interaction points where the four main experiments (ATLAS, CMS, ALICE, LHCb) are located, at these locations, every 25 ns, the bunches cross and interact with each other (*bunch crossing*). A schematic view of the injection chain is depicted in Figure 4.1.

The instantaneous luminosity depends on the beam parameters and is given by:

$$\mathcal{L} = \frac{N_b^2 n_b f_{rev} \gamma}{4\pi \epsilon_n \beta^*} F \quad (4.1.1)$$

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 is the relativistic gamma factor, ϵ_n the normalized transverse beam emittance, the beta function is a measure of the transverse beam size and β^* is the value of 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 (IP) [48]. The integrated luminosity is given by:

$$L = \int \mathcal{L} dt \quad (4.1.2)$$

and the integral is carried over data taking periods of the detector. The integrated luminosity can be related to the total number of events of a certain process by:

$$N_{events} = L \sigma_{events} \quad (4.1.3)$$

where N_{events} is the total number of events, L is the integrated luminosity and σ_{events} is the cross section for the process in units of barn ($1 \text{ b} = 10^{-28} \text{ m}^2$). In 2015 ATLAS recorded an integrated luminosity of 3.2 fb^{-1} and in 2016 of 32.9 fb^{-1} .

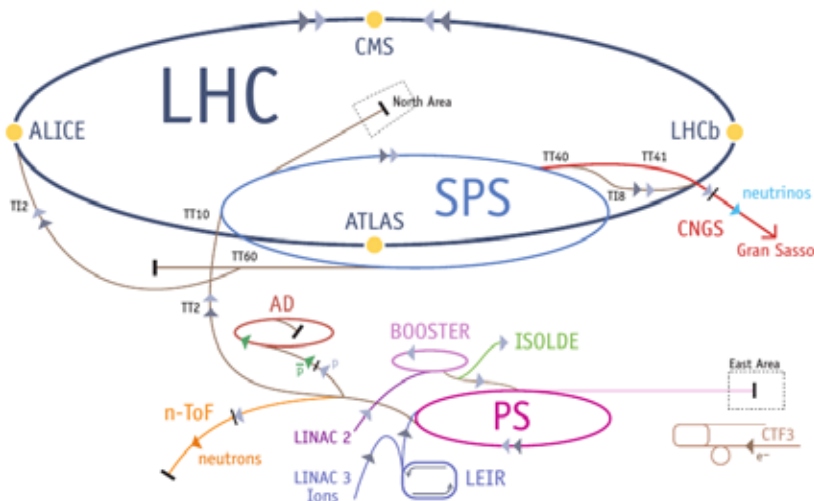


Figure 4.1: The LHC injection chain [55].

4.2 The ATLAS Detector

ATLAS is a multi purpose detector designed to be sensitive to a large physics signatures (supersymmetry and dark matter, briefly introduced in Section 3.2 and Section 3.1.2) and to fully take advantage of the LHC potential. It is capable of identifying photons, electrons, muons, taus, jets and missing energy, Figure 4.2 shows a schematic view of the interaction of the different kind of particles with the ATLAS sub-detectors while Figure 4.3 shows the ATLAS detector with its subsystems. In the following sections a brief overview of

the various subsystems that allow particle identification and reconstruction is presented.

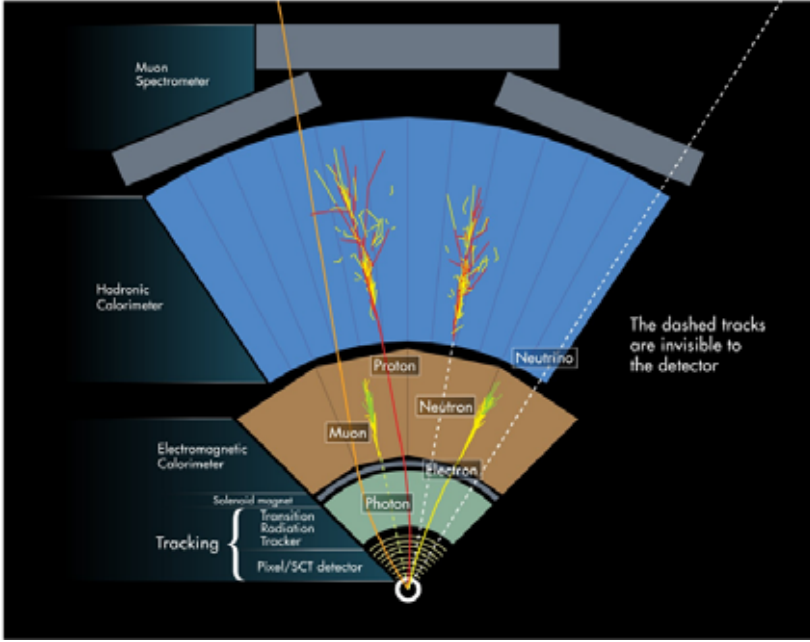


Figure 4.2: Section of the ATLAS detector showing the interaction of different particle types with the sub-detectors [56].

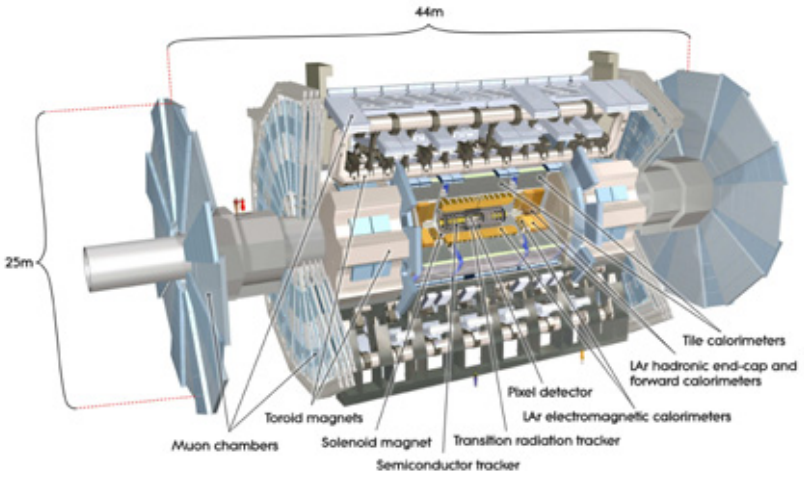


Figure 4.3: Overview of the ATLAS detectors with its main sub-detectors [49].

4.2.1 The Coordinate System

A right handed coordinate system is defined for the ATLAS detector, the origin is at the geometric center of ATLAS with the z -axis oriented along the beam direction and the xy plane orthogonal to it. The positive x -axis points to the center of the LHC ring while the positive y -axis is pointing upwards. The A-side of the detector is defined as that with a positive z -axis while the C-side has the negative z -axis.

The LHC beams are unpolarized and thus invariant under rotations around the beam line axis, a cylindrical coordinate system is particularly convenient to describe the detector geometry where:

$$r = \sqrt{x^2 + y^2}, \quad \phi = \arctan \frac{y}{x}. \quad (4.2.1)$$

A momentum dependent coordinate, the *rapidity*, is commonly used in particle physics for its invariance under Lorentz transformations along the z -axis. The rapidity of a particle is defined as:

$$y = \frac{1}{2} \ln \frac{E + p_z}{E - p_z} \quad (4.2.2)$$

where E is the energy of the particle and p_z its momentum along the z -axis. Rapidity intervals are Lorentz invariant under boost along the z -axis. In the relativistic limit or when the mass of the particle is negligible, the rapidity only depends on the production angle of the particle with respect to the beam axis,

$$\theta = \arctan \frac{\sqrt{p_x^2 + p_y^2}}{p_z}. \quad (4.2.3)$$

This approximation is called pseudorapidity (η) and is defined as:

$$y \xrightarrow{p \gg m} \eta = -\ln \left(\tan \frac{\theta}{2} \right). \quad (4.2.4)$$

A value of $\theta = 90^\circ$, perpendicular to the beam axis, corresponds to $\eta = 0$. The spatial separation between particles in the detector is commonly given in terms of a Lorentz invariant variable defined as:

$$\Delta R = \sqrt{\Delta\phi^2 + \Delta\eta^2}. \quad (4.2.5)$$

Other quantities used to describe the kinematics of the pp interaction are the transverse momentum (p_T) and the transverse energy (E_T) defined as $p_T = p \sin \theta$ and $E_T = E \sin \theta$ respectively.

4.2.2 The Inner Detector

The Inner Detector (ID) [49] is designed to provide good track reconstruction, precise momentum resolution and both primary and secondary vertex measurements (see Section 6.1) above a nominal p_T threshold of 0.5 GeV and within the pseudorapidity $|\eta| < 2.5$. The ID is 6.2 m long and has a radius of about 1.1 m,

it is surrounded by a solenoidal magnetic field of 2 T. Its layout is schematized in Figure 4.4 and, as can be seen, it is composed of three sub-detectors.

At the inner radius the *pixel detector* measures charged particles with silicon sensors with a minimum and maximum size of $50 \times 400 \mu\text{m}^2$ and $50 \times 600 \mu\text{m}^2$ respectively.

In the middle of the ID the SemiConductor Tracker (SCT) is designed to give eight precision measurements per track which contribute to determine the primary and secondary vertex position and momentum measurements. The silicon sensors are $80 \mu\text{m}$ pitch micro strips.

The last layer of the ID is the Transition Radiation Tracker (TRT), it contributes to tracking and identification of charged particles. It consists of drift (straw) tubes, 4 mm in diameter with a $31 \mu\text{m}$ wire in the center of each straw, filled with a gas mixture. These tubes substantially act like proportional counters where the tube is the cathode and the wire is the anode and set to ground. When a charged particle cross one tube, leaves a signal; the set of signals in the tubes, reconstructs to a track which represents the path of the crossing object. The space between the straw tubes is filled with material with a different dielectric constant than the inside, this causes charged particles crossing the boundaries to emit transition radiation thus leading to some straw to have a much stronger signal. The transition radiation depends on the Lorentz γ factor which in turn depends on the energy and the mass of the particles. Lighter particles will have higher transition energy and stronger signal in the straw

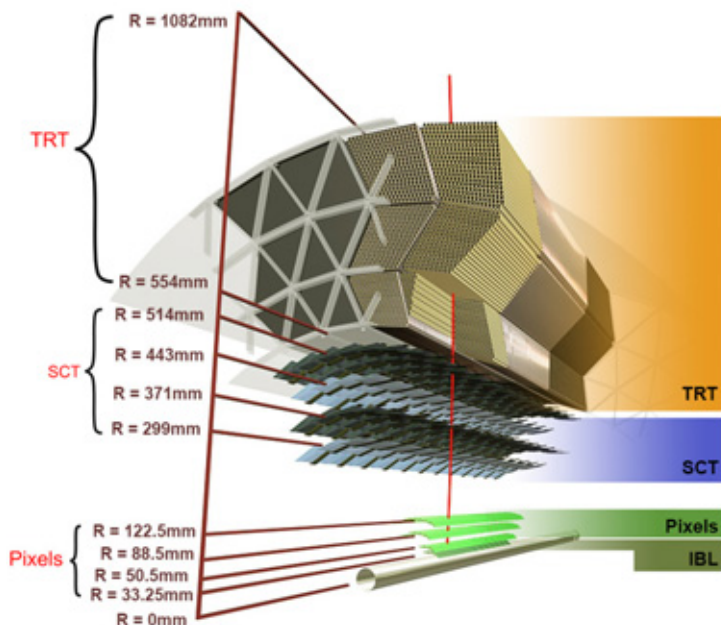


Figure 4.4: Schematic view of a charged track of $10 \text{ GeV } p_T$ that traverses the different ID sub-detectors. After traversing the beryllium pipe, the track passes through the IBL, the three cylindrical silicon-pixel layers, the four layers of silicon-microstrip sensors (SCT) and the approximately 36 straws contained in the TRT within their support structure [57].

tubes. This allows to distinguish between electrons (the lightest charged particle) and single hadrons (pions), for instance, tracks with several strong signal straw, can be identified as belonging to electrons.

An additional layer, the Insertable B-Layer (IBL), was added in the region between the beam pipe and the inner pixel layer (B-layer) for the run 2 LHC phase. It is designed to increase the tracking robustness by replacing damaged parts of the pixel B-layer and increasing the hit redundancy with higher luminosity. In addition, being closer to the beam pipe it increases the impact parameter measurement precision [58].

4.2.3 The Calorimeter

The main purpose of a calorimeter is to measure the energy of electrons, photons and hadrons by mean of materials capable of completely absorbing the energy of the incoming particles transforming it in some measurable quantity. Calorimeters can be classified in two categories, electromagnetic (EM) and *hadronic* depending on the particle they are designed to detect. The EM calorimeters are used to detect photons and electrons while the task of hadronic calorimeters is to identify hadrons. Both types of calorimeters can be further divided into *sampling calorimeters* and *homogeneous calorimeters*. Sampling calorimeters alternate layers of a dense material used to absorb the energy of incident particles (absorber) and an active material to collect the signal. The interaction between the particles and the absorber produces a shower of secondary particles with progressively degraded energy which is deposited in the active material in form of charge or light that can be converted into energy. Homogeneous calorimeters use only one material that serves both as an absorber and an active material [59].

The ATLAS calorimeter is a sampling calorimeter covering up the $|\eta| < 4.9$ region the large η coverage, ensures a good missing transverse momentum measurement (see Section 6.8); an illustration of the system is shown in Figure 4.5.

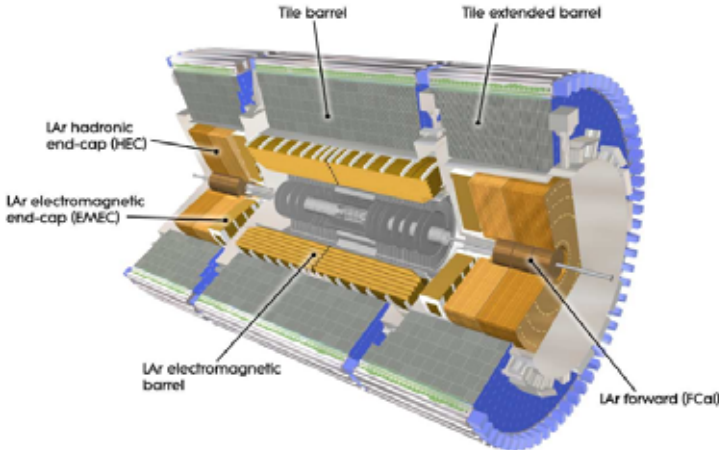


Figure 4.5: Cut-away view of the ATLAS calorimeter system [49].

The EM calorimeter has a barrel and two end-caps, covering the $|\eta| < 1.475$ and $1.375 < |\eta| < 3.2$ region respectively. It uses Liquid Argon (LAr) as active material and lead as absorber in an accordion geometry that provides ϕ symmetry without azimuthal cracks. In the region $|\eta| < 1.8$ a presampler consisting of a LAr active region is used to correct for electrons and photons energy loss upstream of the calorimeter.

There are three hadronic calorimeters: the Tile Calorimeter (TileCal), the Hadronic End-cap Calorimeter (HEC) and the LAr Forward Calorimeter (FCal). The TileCal barrel and extended barrels cover the $|\eta| < 1.0$ and $0.8 < |\eta| < 1.7$ and uses steel as absorber and scintillating tiles connected to photomultiplier tubes through wavelength shifting fibers for readout as an active material. The HEC covers the $1.5 < |\eta| < 3.2$ region and, to avoid drops in material density at the transition, it overlaps slightly with the FCal that covers the $3.1 < |\eta| < 4.9$.

4.2.4 The Muon Spectrometer

The Muon Spectrometer (MS) is designed to identify muons and measure their momentum. It is divided in four sub-detectors, the Monitored Drift Tubes (MDT), the Cathode Strip Chambers (CSC), the Resistive Plate Chambers (RPC), and the Thin Gap Chamber (TGC). Figure 4.6 shows a cut-away view of the muon spectrometer. The sub-detectors are immersed in a magnetic field generated by three different toroidal magnets, a barrel toroid covering the $|\eta| < 1.4$ region and two end-caps magnets at $1.6 < |\eta| < 2.7$, which produces a field almost perpendicular to the muon tracks.

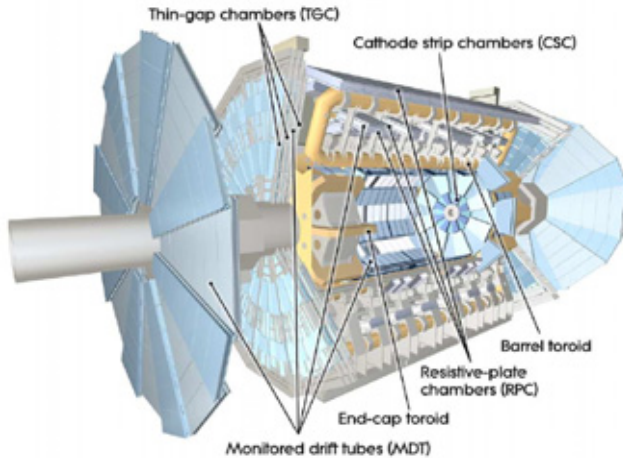


Figure 4.6: Cut-away view of the ATLAS muon spectrometer [49].

The MDT covers the $|\eta| < 2.7$ region and provides a precise measurement of the track coordinates in the principal bending direction of the magnetic field. It uses drift tubes to reconstruct the muon trajectory and the drift time of the ionized charges is used to determine the minimum distance between the wire and the muon. The CSC covers the $2.0 < |\eta| < 2.7$ region and is a multi-wire

proportional chamber with cathodes segmented in strips, one perpendicular to the anode wire, providing the precision coordinate, and the other parallel to it (giving the transverse coordinate).

The RPC and the TGC cover the $|\eta| < 1.05$ and $1.05 < |\eta| < 2.7$ regions respectively. They contribute to the Level 1 trigger providing bunch crossing identification, it allows to select high and low p_T tracks and measure the muon coordinate in the direction orthogonal to that determined by MDT and CSC.

4.2.5 The Forward Detectors

The ATLAS forward region is covered by three smaller detectors: the LUMinosity measurement using Cerenkov Integrating Detector (LUCID), the Absolute Luminosity For ATLAS (ALFA) and the Zero-Degree Calorimeter (ZDC). LUCID [60] is located at ± 17 m from the IP, it is designed to monitor the relative luminosity by detecting the inelastic pp scattering. The ZDC [60] is located at ± 140 m from the IP, it consists of alternating layers of quartz rods and tungsten plates designed to measure neutron at $|\eta| < 8.2$, its purpose is to measure the centrality in heavy-ion collisions. ALFA [60] is located at ± 240 m from the IP and is designed to measure the absolute luminosity via elastic scattering at small angles.

4.2.6 The Trigger System

The bunch crossing rate at LHC is 40 MHz for a bunch spacing of 25 ns (about 7 meters). Each event recorded by ATLAS requires ≈ 1.4 MB of disk space, with approximately 20 to 50 collisions per bunch crossing, the storage space required to record all the events in a second would be ≈ 60 TB. This is not feasible thus only the most interesting events are selected and stored on disk. The *trigger system* decides whether to keep or not a collision event for later studies, it consists of a hardware based Level One (L1) trigger and a software based High Level Trigger (HLT).

The L1 trigger determines Region of Interest (RoIs) in the detector using custom hardware and coarse information from the calorimeter and the muon system. The L1 trigger is capable of reducing the event rate to 100 kHz with a decision time for a L1 accept of $2.5 \mu\text{s}$. The RoIs from the L1 trigger are sent to the HLT where different algorithms are run using the full detector information and reducing the L1 output rate to 1 kHz with a processing time of about 200 ms [61].

4.3 The 2015 and 2016 ATLAS Datasets

In the search for SUSY in the compressed spectrum scenario presented in Chapter 7 of this thesis, the HLT_xe70 trigger is used, it receives an L1 accept that selects events with a missing energy (see Section 6.8) greater than 50 GeV computed at the L1 trigger. No muons are used in the reconstruction of the missing energy at the trigger level. The events that survive L1 are then passed to the

HLT level, where events with a missing energy (calculated from cell energy information) greater than 70 GeV are selected.

Figure 4.7 shows the delivered luminosity as a function of the average number of interactions per bunch crossing $\langle\mu\rangle$ during the 2015 and 2016 data taking periods. Due to the increasing luminosity and pile-up (see Section 6.1), in the search for large extra dimensions in the ADD model scenario presented in Chapter 8 of this thesis the trigger had to be modified. A combination of four different triggers is used, they all pass the level one trigger selection if the event have a missing energy of more than 50 GeV at L1, again no muon information is used in the missing energy reconstruction at this level. For the HLT the missing energy threshold is increased in steps of 10 GeV from 80 GeV up to 110 GeV for the four different triggers. The information on the missing energy takes advantage of the Missing Hadronic Trigger (MHT) [62] algorithm or the Local Cluster Weighting (LCW) [63] calibration scheme depending on the trigger used. In the MHT algorithm the missing energy is calculated as the negative vector sum of the transverse momentum ($-\sum p_T^{\text{jets}} \equiv -H_T$) of all the jets reconstructed with the anti- k_t jet finding algorithm (see Section 6.6.1) from topoclusters (see Section 6.5). The LCW calibration scheme uses the shape of the cluster and the energy density to classify the topocluster energy deposits as electromagnetic or hadronic improving the resolution of the missing energy trigger.

Table 4.1 summarizes the different trigger combination for the 2016 dataset used in the analysis. The increase with time of the maximum instantaneous luminosity $\mathcal{L}^{\text{max}}$ forced the use of a higher missing energy threshold in the trigger. Figure 4.8 shows the trigger efficiency for the four different data taking periods in the 2016 analysis as a function of the missing energy as evaluated offline (i.e. the same variable used in the analyses) estimated using $W(\rightarrow \mu\nu) + \text{jets}$ events. It can be seen that the trigger is fully efficient from approximately 200 GeV. The monojet analyses presented in Chapters 7 and 8 require a missing transverse energy of at least 250 GeV and are therefore not affected by the loss of efficiency below 200 GeV.

Trigger Used in the 2016 Analysis		
Run Range	Trigger	$\mathcal{L}^{\text{max}}(10^{30} \text{ cm}^{-2} \text{ s}^{-1})$
296939-302393	HLT_xe90_mht_L1XE50 or HLT_xe80_tc_lcw_L1XE50	8761
302737-302872	HLT_xe90_mht_L1XE50	9854
302919-304008	HLT_xe100_mht_L1XE50 or HLT_xe110_mht_L1XE50	10261
304128-310216	HLT_xe110_mht_L1XE50	13716

Table 4.1: The table reports the trigger used in the 2016 data taking. The maximum instantaneous luminosity $\mathcal{L}^{\text{max}}(10^{30} \text{ cm}^{-2} \text{ s}^{-1})$ is also reported. The constant increasing of $\mathcal{L}^{\text{max}}$ justifies the choice of higher E_T^{miss} thresholds for the trigger.

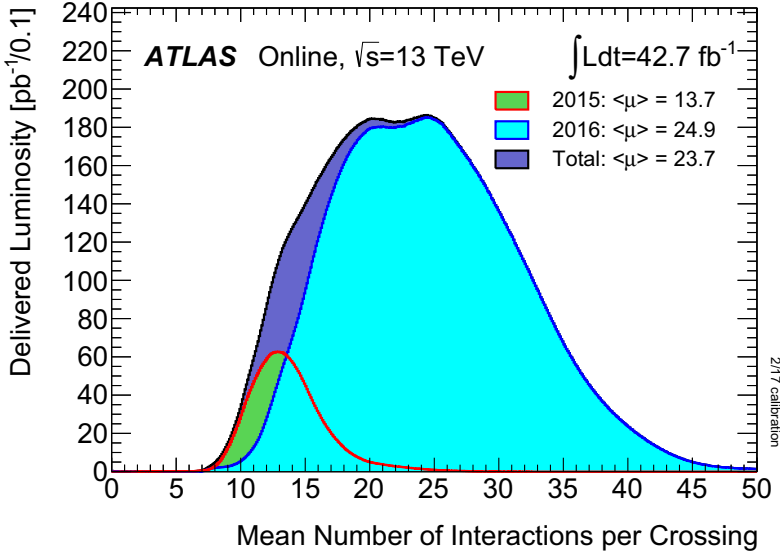


Figure 4.7: Delivered luminosity as a function of the average number of interactions per bunch crossing $\langle\mu\rangle$ during the 2015 and 2016 data taking periods at $\sqrt{s} = 13$ TeV [64].

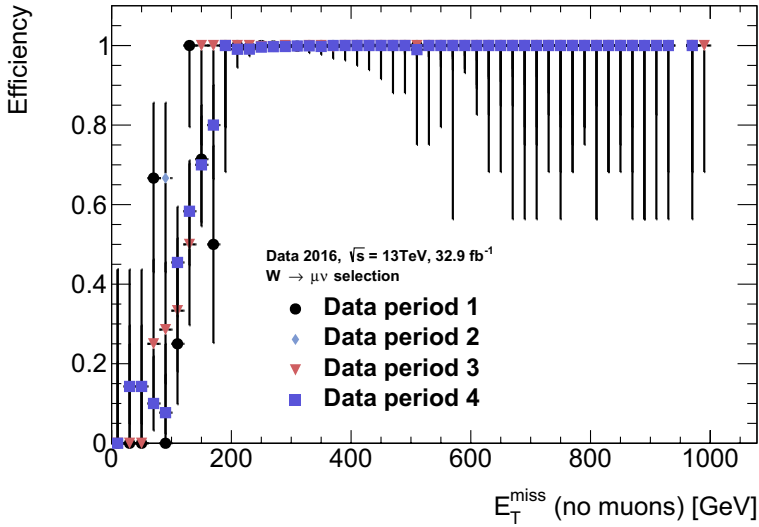


Figure 4.8: Trigger efficiency curve as a function of the missing energy for the 2016 dataset of 32.9 fb⁻¹ estimated in $W(\rightarrow \mu\nu) + \text{jets}$ events with selected requiring exactly one good muon as defined in Section 6.4 and vetoing any other muon or electron, at most four jets in the event, a leading jet p_T of at least 250 GeV and a $\Delta\phi$ between any jet and the E_T^{miss} of less than 0.4. The four different periods correspond to the run intervals defined in Table 4.1. The trigger is fully efficient from approximately 200 GeV.

Chapter 5

Noise Studies with the Tile Calorimeter

5.1 Calorimetry

Particles lose energy interacting with matter. The particle's energy and its type determines the processes causing the energy loss; these can be of two kind, *electromagnetic* and *hadronic*. In this section a brief overview of the physics behind these interactions is given.

5.1.1 Electromagnetic Shower

High energy electrons and photons lose energy mainly by *radiation* and *conversion* respectively. When electrons with energies greater than ~ 10 MeV interact with the electromagnetic field of the absorber nuclei *Bremsstrahlung* can occur. High energy photons produce mostly electron-positron pairs.

Electrons and photons with a sufficient amount of energy interacting with an absorber, produce secondary photons through Bremsstrahlung or secondary electrons and positrons by pair production. These secondary particles will produce more particles through the same mechanisms giving rise to a shower of particles with progressively lower energies. When the energy loss in the shower is dominated by ionization and thermal excitation of the active material atoms, the number of particles in the shower stops growing and there is no further shower development. The above process goes on until the energy of the electrons falls below a critical energy, ϵ , where ionization and excitation becomes the dominant effects [59].

5.1.2 Hadronic Shower

Hadrons lose energy through strong interaction with the calorimeter material. The strong interaction is responsible for the production of energetic secondary hadrons with momenta typically at the GeV scale and nuclear reactions such

as excitation or nucleon spallation in which neutrons and protons are released from the nuclei with a characteristic energy at the MeV scale.

These energetic hadrons are protons, neutrons and pions. On average, 1/3 of the pion produced are neutral pions which decay to photons ($\pi \rightarrow \gamma\gamma$). The photons produced this way will initiate an electromagnetic shower as described in Section 5.1.1 transferring energy from the hadronic part to the electromagnetic shower inside the hadronic shower. The electromagnetic component of the shower does not contribute any more to hadronic processes. The nucleons released by excitation or nuclear spallation, require an energy equal to their binding energy to be released and are not recorded as a contribution to the calorimeter signal thus producing a form of *invisible energy*. Some detectors can compensate for the loss of invisible energy, these are called *compensating calorimeters* [59].

5.1.3 Energy Resolution

The energy resolution of a detector measures its ability of distinguishing between particles of different energies; the better the energy resolution, the better it can separate energy peaks belonging to different decays.

The energy resolution can be written as:

$$\frac{\sigma_E}{E} = \frac{a}{\sqrt{E}} \oplus \frac{b}{E} \oplus c, \quad (5.1.1)$$

where the $\oplus$ symbol indicates a quadratic sum. The first term (a) in the equation is the *stochastic term*, it is mainly due to fluctuations related to the physical evolution of the shower. In homogeneous calorimeters, this term is small because the energy deposited in the active volume by a monochromatic beam of particles is constant for each event. In a sampling calorimeter, the active layers are interleaved with absorber layers thus the energy deposited in the active material fluctuates event by event. These are called *sampling fluctuations* and, in sampling electromagnetic calorimeters, represents the greatest limitation to energy resolution due to the variation in the number of charged particles which cross the active layers. The second term (b) in Eq. (5.1.1) is called the *noise term*, it comes from the electronic noise of the detector readout chain. Sampling or homogeneous calorimeters which collect the signal in the form of light, using for example a photo-multiplier tube with a high gain multiplication of the signal with a low electronic noise, can achieve low levels of noise. Calorimeters that collect the signal in form of charge, must use an pre-amplifier having thus a higher level of noise. In sampling calorimeters, the noise term can be further reduced by increasing the sampling fraction, this way there is a larger signal coming from the active material and a higher noise-to-signal ratio. The last term (c) of the equation is the *constant term*, it does not depend on the energy of the particles but includes all the non uniformities in the detector response such as instrumental effect, imperfections in the calibration of different parts of the detector, radiation damage, detector aging, or the detector geometry [59].

5.2 The ATLAS TileCal

TileCal is the central hadronic calorimeter of the ATLAS experiment covering the $|\eta| < 1.7$ region. It is designed for energy measurement of hadrons, jets, tau particles and also contributes to the measurement of missing transverse energy (see Section 6.8). TileCal is a scintillator steel non compensating sampling calorimeter, the scintillation light produced in the tiles is transmitted by Wavelength Shifting Fibers (WSFs) to PhotoMultiplier Tubes (PMTs). The analog signals from the PMTs are amplified, shaped and digitized by sampling the signal every 25 ns. The TileCal front end electronics read out the signals produced by about 10000 channels measuring energies ranging from 30 MeV to 2 TeV. The readout system is responsible for reconstructing the data in real time. The digitized signals are reconstructed with the Optimal Filtering algorithm (see Section 5.2.1), which computes for each channel the signal amplitude, time and quality factor at the required high rate.

TileCal is designed as one Long Barrel (LB) covering the $|\eta| < 1.0$ range and two Extended Barrel (EB) in the $0.8 < |\eta| < 1.7$ range. The barrels are further divided, according to their geometrical position on the z -axis, in partitions called EBA, LBA, EBC and LBC (see Section 4.2.1). Each partition consists of 64 independent wedges (see Figure 5.1) called *modules* assembled in azimuth (ϕ). The LBA and EBA partitions are shown in Figure 5.2.

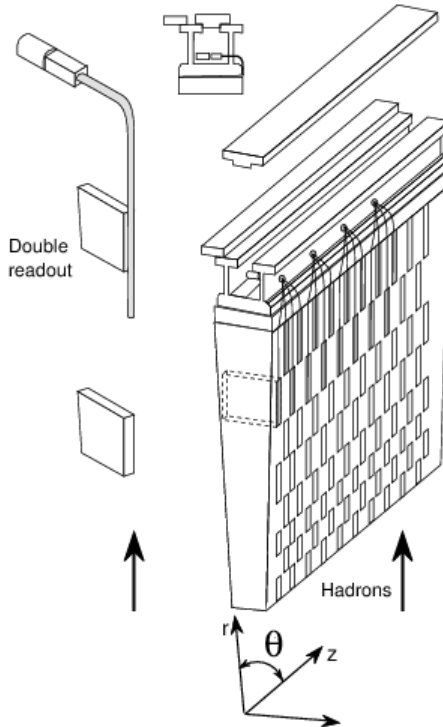


Figure 5.1: Cut away showing an individual TileCal module along with the optical read out and design of a TileCal module [65].

Between the LB and the EB there is a 600 mm gap needed for the ID and the LAr cables, electronics and services. Part of the gap contains the Intermediate Tile Calorimeter (ITC), a detector designed to maximize the active material while leaving enough space for services and cables. The ITC is an extension of the EB and it occupies the $0.8 < |\eta| < 1.6$ region. The combined $0.8 < |\eta| < 1.0$ part is called *plug* and in the $1.0 < |\eta| < 1.6$ region, for space reasons, the ITC is not interleaved with an absorber and is only composed of scintillator material. The scintillators between $1.0 < |\eta| < 1.2$ are called *gap scintillators*, while those between $1.2 < |\eta| < 1.6$ are called *crack scintillators*. The plug and the gap scintillators mainly provide hadronic shower sampling while the crack scintillator, which extends to the region between the barrel and the end-cap cryostats, samples the electromagnetic shower in a region where the normal sampling is impossible due to the dead material of the cryostat walls and the ID cables.

TileCal is also divided in longitudinal layers, the A, BC and D layers as shown in Figure 5.2. The two innermost layers have a $\Delta\eta \times \Delta\phi$ segmentation of 0.1×0.1 while in the outermost, the segmentation is 0.1×0.2 . Each layer is logically divided into *cells* (also shown in Figure 5.2) by grouping together in the same PMT the fibers coming from different scintillator tiles belonging to the same radial depth. The gap/crack scintillators are also called E layer cells.

The energy resolution for jets in ATLAS is:

$$\frac{\sigma_E}{E} = \frac{50\%}{\sqrt{E}} \oplus 3\% \quad (5.2.1)$$

for $|\eta| < 3$. The 3% constant term becomes dominant for high energy hadrons where an increase in energy resolution is expected [66].

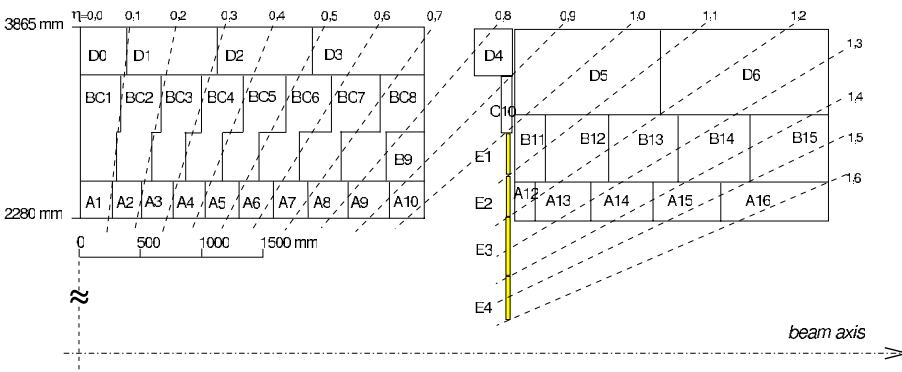


Figure 5.2: Schematic view of the TileCal layer and cell structure in a plane containing the beam axis z [67].

5.2.1 Signal Reconstruction

The TileCal cells are read out by two PMTs with the exception of the E layer cells that are connected to only one photomultiplier tube using WSF. Each PMT is associated to an electronic read-out channel with its own shaper,

preamplifier and Analog to Digital Converter (ADC). The current pulse from the PMTs is shaped and amplified by the 3-in-1 card. There are two possible gains: High Gain (HG) and Low Gain (LG), with an amplification ratio of 64. The 3-in-1 card forms the front-end electronics of the read-out chain and provides three basic functions: shaping of the pulse, charge injection calibration and slow integration of the PMT signals for monitoring and calibration [66]. Up to twelve 3-in-1 cards are serviced by a motherboard that provides power and individual control signals. The amplified signal is sent to two ADCs synchronous with the 40 MHz LHC clock thus sampling the signal every 25 ns. For optimization and efficiency reasons, 7 samples for each pulse are taken and sent to the ReadOut Drivers (RODs) if an L1 trigger accept is received.

Optimal Filtering

The seven samples are used to reconstruct the amplitude of the pulse using the Optimal Filtering (OF) method. The estimate of the amplitude is given by:

$$\hat{A} = \sum_{i=0}^7 a_i S_i \quad (5.2.2)$$

where S_i are the digitized samples expressed in ADC counts and a_i are computed weights that minimize the effect of the electronic noise on the amplitude reconstruction. The procedure minimizes the variance of the amplitude distribution. In order to make the amplitude reconstruction independent from phase and signal baseline due to electronic noise (*pedestal*), the following constraints are used:

$$\sum_{i=0}^7 g_i a_i = 0 \quad (5.2.3)$$

$$\sum_{i=0}^7 g'_i a_i = 0 \quad (5.2.4)$$

$$\sum_{i=0}^7 a_i = 0 \quad (5.2.5)$$

where g_i and g'_i are the pulse shape function from the shaper and its derivative [68].

The TileCal Calibration

The energy deposited in the calorimeter cell is proportional to the reconstructed amplitude. The amplitude is originally measured in ADC counts and needs to be converted in GeV for physics analysis using the formula:

$$E[GeV] = \hat{A}[ADC] \times C_{ADC \rightarrow pC} \times C_{\text{laser}} \times C_{Cs} \times C_{pC \rightarrow GeV} \quad (5.2.6)$$

where $\hat{A}[ADC]$ is the amplitude estimate in ADC counts, $C_{ADC \rightarrow pC}$ is determined using the Charge Injection System (CIS), $C_{pC \rightarrow GeV}$ is measured during

testbeam using electrons with a well defined energy and converts the deposited charge to energy in GeV, the laser system allows to determine the value of the C_{laser} constant while the Cesium sets the C_{Cs} factor.

The CIS calibrates the read out electronics by injecting a known charge and measuring the resulting response of the electronics. The *laser* system main purpose is to monitor the photomultipliers tubes stability and the downstream electronics. Well calibrated light pulses are sent to the PMTs and by reconstructing the signal it is possible to extract the PMTs' gain. The *cesium* system, circulates a Cs source through each scintillating tile using an hydraulic system, the PMTs signal is continuously read out through an integrator. The cesium system allows to equalize the calorimeter cell response to that measured during test beams. Figure 5.3 depicts a schematic representation of the ATLAS TileCal calibration chain.

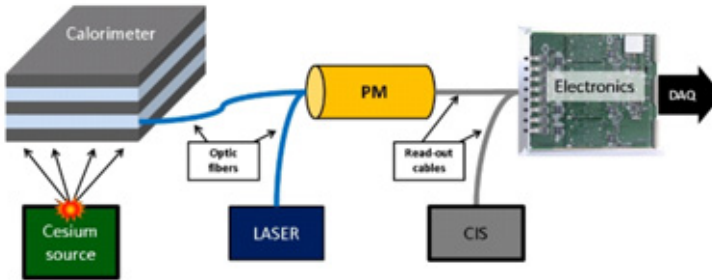


Figure 5.3: The ATLAS TileCal calibration chain [69].

5.2.2 Electronic Noise

TileCal periodically records sets of events (referred to as *runs*) with no signal in the PMTs, called *pedestal runs*. During these runs, each channel is read out using both, HG and LG for about 100000 events. These events are sampled every 25 ns in 7 samples as in normal physics runs and are normally distributed around a mean value called *pedestal*. The Root Mean Square (RMS) of the pedestal is the *noise*. Pedestal runs are used to calculate different parameters called *noise constants* that allow to describe the electronic noise. Two different sets of noise constants are computed: *Digital Noise* (or *Sample Noise*) and *Cell Noise*.

Digital Noise

The digital noise is measured in ADC counts for each PMT in both gains, HG and LG. The noise constants together with all detector conditions are stored in the ATLAS Condition Database (COOL). These constants are the RMS of the seven samples within each event, also called High Frequency Noise (HFN) and the RMS of the first digitized sample in each event or Low Frequency Noise (LFN). The digital noise is used for monitoring the electronics and for Monte Carlo (MC) noise simulation of the calorimeter response.

Cell Noise

Excluding the E layer cells that are connected only to one PMT, the cell noise is the combination of the two readout channels of a cell where the digital noise from the PMTs is added quadratically and converted in MeV using the calibration constants and Eq. (5.2.6). There are four possible gain combinations: High Gain - High Gain (HGHG), Low Gain - Low Gain (LGLG), Low Gain - High Gain (LGHG) and High Gain - Low Gain (HGLG). When the energy deposit is large, the two channels belonging to the cell are readout in low gain thus giving rise to the LGLG combination, for small signals the HGHG combination is used. It can also occur that for energy deposits in the intermediate range, one PMT is readout in LG and the other in HG resulting in a HGLG combination. The cell noise is used to identify the seed cells in the topocluster algorithm (see Section 6.5).

Figure 5.4 shows a comparison between the cell noise and the fitted σ parameter of a normal distribution. The ratio $\text{RMS} / \sigma = 1$ indicates a perfect agreement between the measured and the fitted amplitude distribution for a single Gaussian hypothesis. The blue square in the plot indicate the comparison for an old model of Low Voltage Power Supply (LVPS) while the red square refers to the currently used LVPSs. It can be seen that with the old model of LVPSs the ratio RMS / σ can be large indicating a non Gaussian behavior of the electronic noise. For this reason a double Gaussian distribution is used to

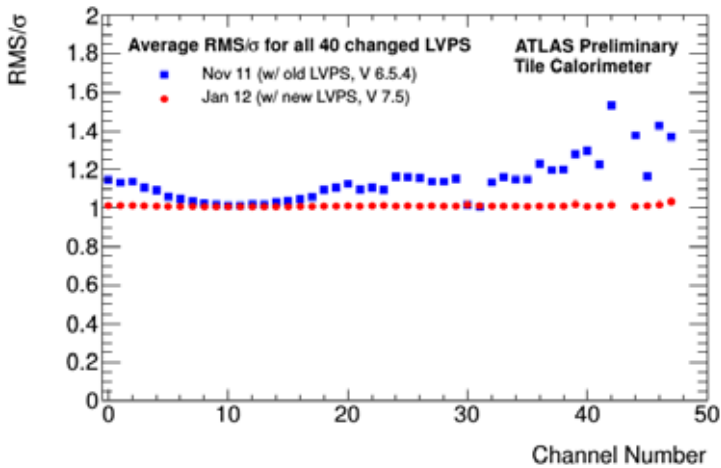


Figure 5.4: Comparison between the TileCal electronic noise, measured as the RMS of the reconstructed amplitude distribution in pedestal runs and the σ of the Gaussian fit of the distribution for the old and new LVPS [70].

fit the energy distribution with the probability density function defined as:

$$f_{2g} = \frac{1}{1+R} \left(\frac{1}{\sqrt{2\pi}\sigma_1} e^{-\frac{x^2}{2\sigma_1^2}} + \frac{R}{\sqrt{2\pi}\sigma_2} e^{-\frac{x^2}{2\sigma_2^2}} \right) \quad (5.2.7)$$

where R is the relative normalization of the two Gaussians and σ_1, σ_2 and R are independent parameters. These three are used to define the region $\sigma_{\text{eff}}(E)$

where the significance for the double Gaussian is the same as the one σ region for a single Gaussian, i.e. $\int_{-\sigma_{\text{eff}}}^{\sigma_{\text{eff}}} f_{2g} = 0.68$ [71]. In terms of σ_{eff} , for an energy deposit E , the significance can be expressed as:

$$\frac{E}{\sigma_{\text{eff}}(E)} = \sqrt{2} \operatorname{Erf}^{-1} \left(\frac{\sigma_1 \operatorname{Erf} \left(\frac{E}{\sqrt{2}\sigma_1} \right) + R\sigma_2 \operatorname{Erf} \left(\frac{E}{\sqrt{2}\sigma_2} \right)}{\sigma_1 + R\sigma_2} \right) \quad (5.2.8)$$

where Erf is the error function. Eq. (5.2.8) is the input to the calorimeter cell clustering algorithm discussed in more details in Section 6.5, moreover this definition allows to use the same unit to describe the noise for both the TileCal and LAr calorimeters. The region $\sigma_{\text{eff}}(E)$ is commonly referred to as *cell noise* and together with the three double Gaussian parameters (σ_1 , σ_2 and R) is stored in the COOL database.

5.3 The 2011 TileCal Run 1 Reprocessing

As new information about the detector becomes available, an update of the calibration constants and thus of the reconstructed energy must be performed; this procedure is called *reprocessing*. In 2011 an update of the laser and cesium calibration analysis procedure were performed together with an update of which cells were considered good or bad. These updates required a recalculation of the cell noise.

5.3.1 Results

The cell noise for the reprocessed data has been calculated in the different gain combinations. In Figures 5.5 to 5.7 the cell noise values have been calculated using all the calibration runs used for the 2011 run 1 reprocessing, each fitted using the optimal filtering method (see Section 5.2.1). The plots show the η -dependence of the ϕ -averaged RMS of the noise in these runs. Each point is an average for a given cell over all modules containing this cell type. This is done for a given gain combination of the readout channels: High Gain-High Gain (Figure 5.5), Low Gain-Low Gain (Figure 5.6) or High Gain-Low Gain (Figure 5.7). The plots separate the different cell types: A, BC, D and E (gap and crack scintillators). Note that gap and crack scintillators have only one readout channel, so are plotted only in HG or LG readout. The transition between the long and extended barrels can be seen in the range $0.7 < |\eta| < 1.0$.

Figures 5.5 to 5.7 exhibit some η -dependence of the noise. There are two main factors behind this, first the low voltage power supplies, which are the main noise source in TileCal, are located approximately at $|\eta| \approx 1$ leading to higher noise in the cells in that region. Secondly the noise in these plots is in MeV, cells with the same noise expressed in ADC counts can have different noise levels when converted in MeV. This is due to the different ADC to GeV calibration factors for cells with different geometries.

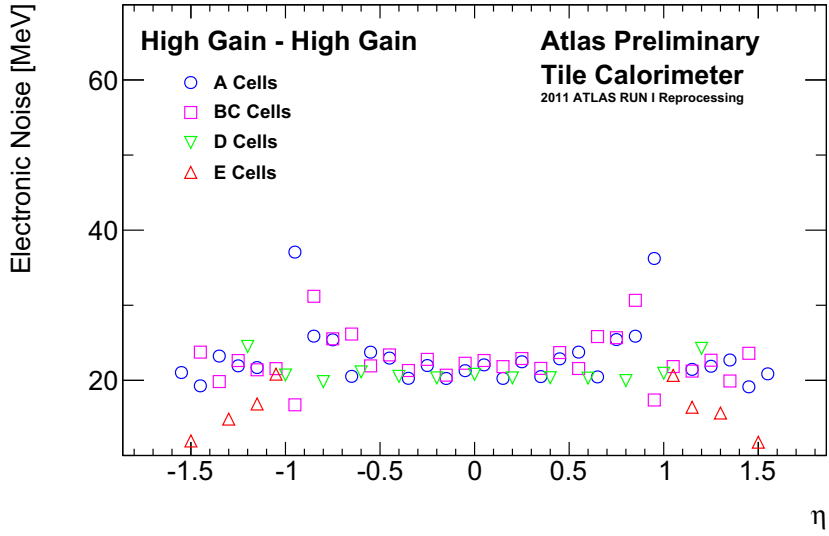


Figure 5.5: ϕ -averaged RMS of electronic cell noise as a function of η of the cell, with both readout channels in High Gain. For each cell the average value over all modules is taken. Values have been extracted using all the calibration runs used for the 2011 RUN I reprocessing. The different cell types are shown separately, A, BC, D, and E (gap/crack). The transition between the long and extended barrels can be seen in the range $0.7 < |\eta| < 1.0$. HGHG combination is relevant when the energy deposition in the cell is $\lesssim 15$ GeV [72].

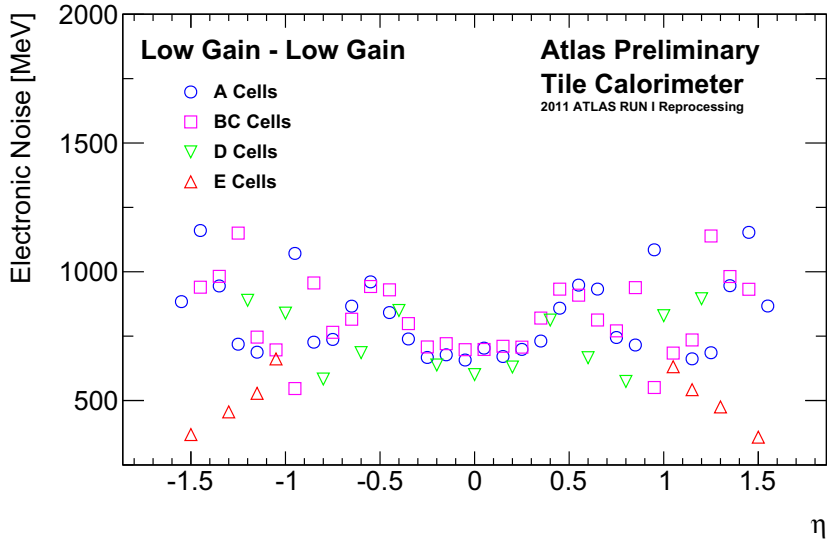


Figure 5.6: ϕ -averaged RMS of electronic cell noise as a function of η of the cell, with both readout channels in Low Gain. For each cell the average value over all modules is taken. Values have been extracted using all the calibration runs used for the 2011 RUN I reprocessing. The different cell types are shown separately, A, BC, D, and E (gap/crack). The transition between the long and extended barrels can be seen in the range $0.7 < |\eta| < 1.0$. LGLG combination is relevant when the energy deposition in the cell is $\gtrsim 15$ GeV [72].

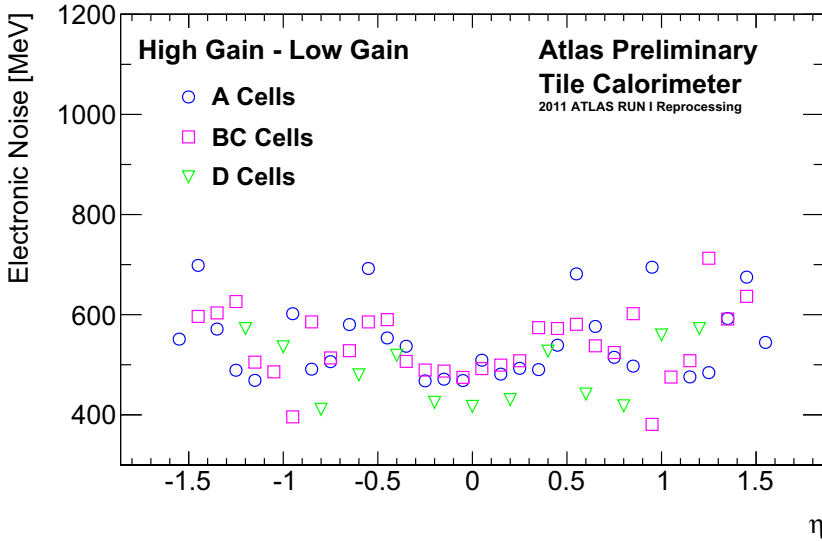


Figure 5.7: ϕ -averaged RMS of electronic cell noise as a function of η of the cell, with one readout channel in High Gain and the other in Low Gain. For each cell the average value over all modules is taken. Values have been extracted using all the calibration runs used for the 2011 RUN I reprocessing. The different cell types are shown separately, A, BC, D. E cells do not appear in this plot since they have only one readout channel. The transition between the long and extended barrels can be seen in the range $0.7 < |\eta| < 1.0$. When the energy deposition in the cell is between $10 \text{ GeV} \lesssim E \lesssim 20 \text{ GeV}$ it can happen that the cell is read out with different voltage on the readout channels [72].

5.4 Time Stability

The calibration constants used in Eq. (5.2.6), as mentioned in Section 5.2.1, are determined with the help of several dedicated calibration systems and runs. Each calibration constant is valid over a period of time called Interval Of Validity (IOV). The cell noise can vary over time for several reasons such as a change in the calibration constants, a variation in the digital noise or the channel status in a particular run. Sudden variations in the noise must be checked and understood.

The different TileCal subsystems (laser, CIS, etc.) all use a common software framework, TileCal Universal Calibration Software (TUCS), to perform validity checks on a number of different studies. To study the stability over time of the updated noise constants, a set of python scripts was developed by the author of this thesis to expand the TUCS functionality. These scripts connect to the ATLAS condition database and allow to visually display the relative change of the cell noise and digital noise constants, the channel status and the ratio between the cell noise and a variable called RMS_{eff} and defined as:

$$\text{RMS}_{\text{eff}} = \sqrt{(1 - R)\sigma_1^2 + R\sigma_2^2} \quad (5.4.1)$$

where σ_1 , σ_2 and R are the free parameters in the double Gaussian model (see Section 5.2.2). The ratio $\sigma / \text{RMS}_{\text{eff}}$, where σ is the cell noise, can be used to

test the goodness of the double Gaussian model: if $\sigma / \text{RMS}_{\text{eff}}$ equals one, the double Gaussian well models the noise, if $\sigma / \text{RMS}_{\text{eff}} > 1$, it means that there is noise that is not well described by it.

Figure 5.8 shows the time evolution plot over the entire reprocessed period for two representative TileCal cells. This represents the typical behavior for most cells over the reprocessing period.

In Figure 5.8a it can be seen that cell number 2 in the BC layer (BC2) of the 41st module in the C side of LB (LBC 41) is stable over several pedestal runs. In Figure 5.8b on the other hand, it is possible to see a variation in the cell noise and of the $\sigma / \text{RMS}_{\text{eff}}$ without a corresponding variation in the calibration, in the digital noise constants or in the channel status. The term *jump* is used in the following to indicate a variation in the cell noise not compatible with a change in the other quantities.

This problem was investigated by re-performing the pedestal noise fit¹ manually and recalculating the noise constants focusing on two specific IOVs, [183110, 183382[(before the jump) and [183382, 183515[(after the jump). Some calorimeter cells without jump were used to validate the noise constants calculated with the fit and checked against the values stored in COOL from automated fits performed by the standard ATLAS software.

Figure 5.9 shows the control cell energy distribution with the double Gaussian fit superimposed for two runs where the jump was present in other cells. The results of the fit and the values stored in the COOL database, both reported in Table 5.1, are in good agreement.

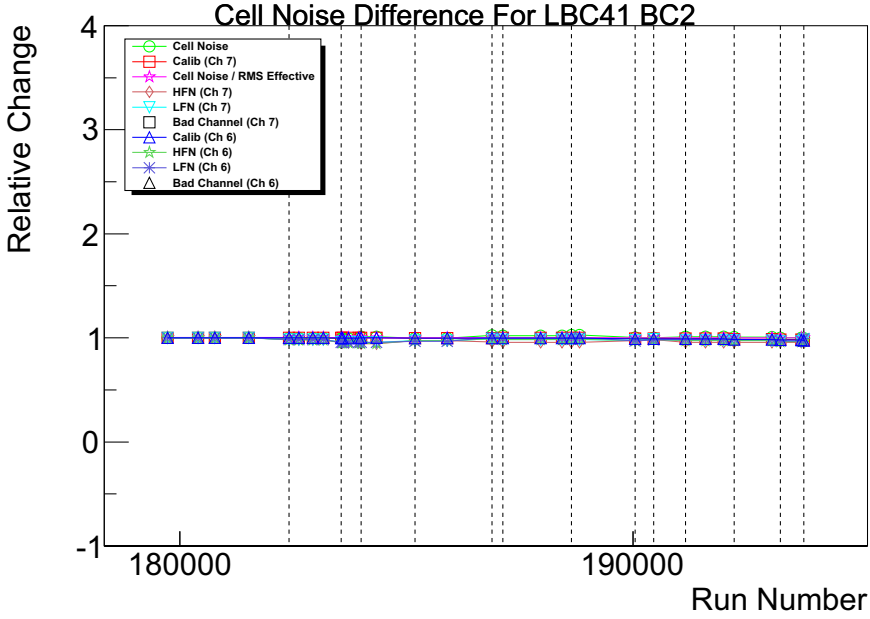
LBC41 BC2 Values Before Jump			LBC41 BC2 Values After Jump		
	Database	Fit		Database	Fit
σ_1 :	19.97	20.08 ± 0.05	σ_1 :	19.94	20.05 ± 0.05
σ_2 :	80.59	77.3 ± 9	σ_2 :	71.41	71.39 ± 10.51
R:	0.00026	0.0003 ± 0.0012	R:	0.00023	0.0002 ± 0.0011

Table 5.1: The table reports the comparison between the double Gaussian parameters stored in the COOL database and those obtained from the fit for two different run numbers corresponding to before and after the jump for a cell where there is no variation in the cell noise.

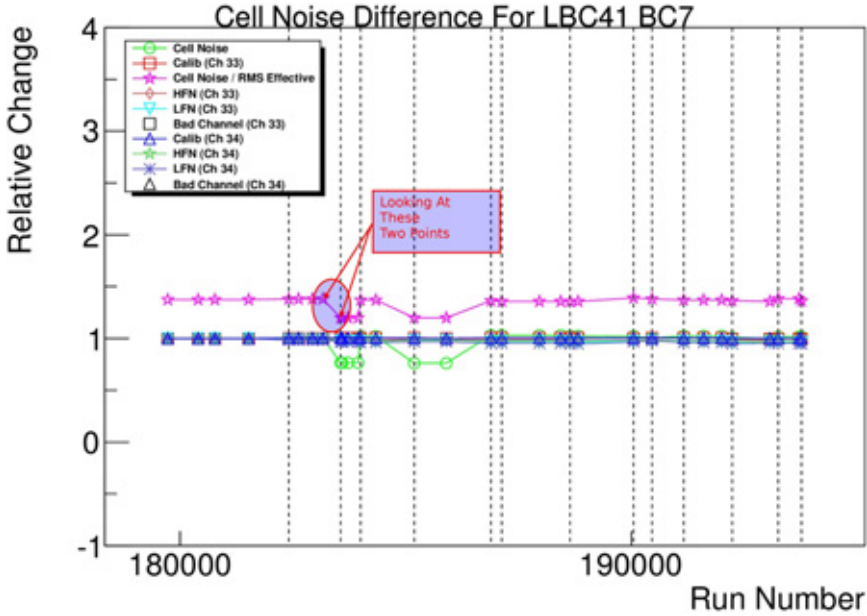
Moving to the investigation of a cell which exhibits the jump in the time evolution the seventh cell of the BC layer (BC7) on the C side of the LB partition of the 41st module (LBC 41) was selected for illustration purposes. Figure 5.10 shows the energy distribution with the double Gaussian fit superimposed. The cell had the jump under investigation (see Figure 5.8) and this is reflected in the fit results reported in Table 5.2 together with the values from the database.

Also in this case, the noise constants from the fit, are in agreement with

¹The standard calibration relies on automated fits performed by the ATLAS reconstruction. It was suspected at first that some of these fits were failing



(a) Control cell with no variation.



(b) Cell with a variation.

Figure 5.8: Time evolution for two different representative cells in the calorimeter over the entire reprocessing period. The plot shows the change relative to the first run considered for several quantities for different IOVs (vertical dashed lines). If a channel is off due to some problems (Bad channel), this is reported in the plot with a black square.

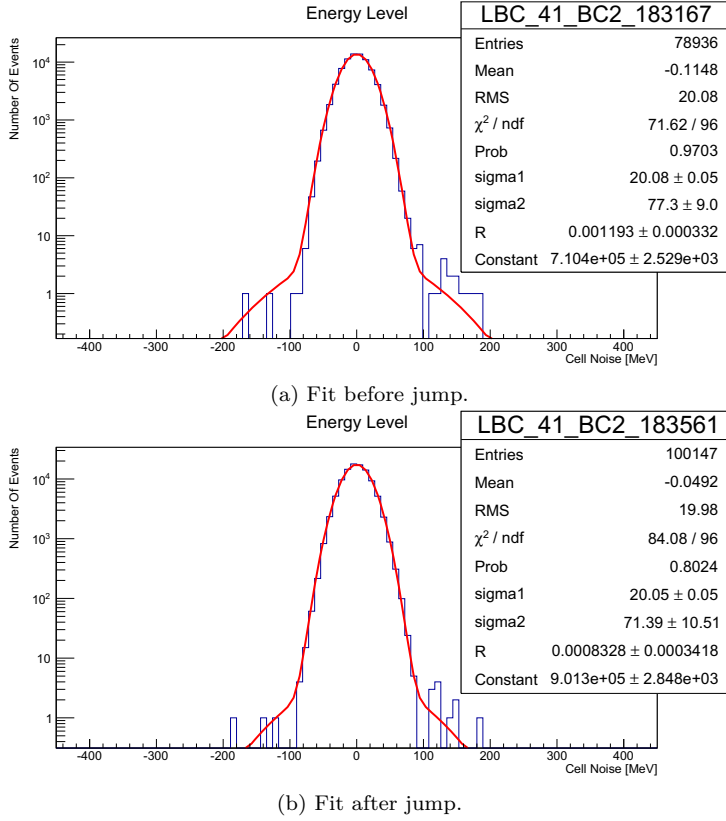


Figure 5.9: Fit of the reconstructed pulse shape on a control cell with no variation (jump) in the cell noise.

LBC41 BC7 Values Before Jump			LBC41 BC7 Values After Jump		
	Database	Fit		Database	Fit
σ_1 :	24.25	23.56 ± 0.1	σ_1 :	24.42	24.4 ± 0.1
σ_2 :	99.16	98.36 ± 0.9	σ_2 :	94.56	94.55 ± 1.34
R:	0.037	0.042 ± 0.036	R:	0.014	0.014 ± 0.015

Table 5.2: The table reports the comparison between the double Gaussian parameters stored in the COOL database and those obtained from the fit for two different run numbers corresponding to before and after the jump for a cell where a variation in the cell noise was spotted.

those stored in the COOL database. From this study it was concluded that bad fits inside the automated ATLAS pedestal and noise reconstruction were not the cause of the jumps. However, the χ^2 of the distribution imply that the double Gaussian model is not a good model for the second type of cells.

After further investigations it was discovered that this behavior is caused by the Tile Noise Filter (TNF). The electronic noise in nearby channels can be

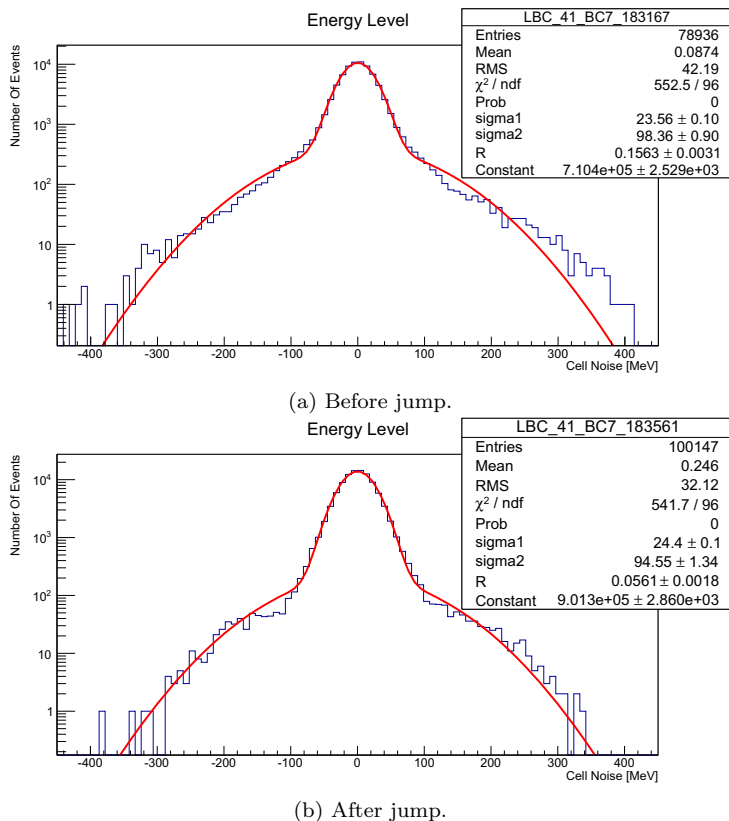


Figure 5.10: Fit of the reconstructed pulse shape on a cell with variation (jump) in the cell noise non compatible with a change in the calibration, digital noise or channel status.

correlated. The version of low voltage power supplies that were used during the 2011 data taking lead to significant coherent noise among several channels, particularly those close to the power supply (see Figure 5.4). Coherent noise means that the signal in several cells can vary coherently and can thus alter the jet and missing transverse energy reconstruction. In order to remedy this a Tile Noise Filter was employed. The TNF calculates for each motherboard the pedestal average:

$$d = \frac{\sum_i^N d_i}{N} \quad (5.4.2)$$

where the index i runs over all the channels belonging to a certain motherboard, d_i is the i -th channel amplitude in ADC counts and N is the number of channels. Variations in this average baseline can be regarded as an estimation of the coherent noise and is subtracted from the channel data of each channel ($d_i - d$) on an event-by-event basis. In order to be able to perform this noise filter calculation only channels without signal from physics are taken into account in the sum of Eq. (5.4.2).

The cell noise was recalculated without noise filter and the corresponding distribution re-fitted. Figure 5.11 show the double Gaussian fit applied to

another cell where the jump was present, in Figures 5.11a and 5.11b the noise filter is on while in Figures 5.11c and 5.11d the TNF was removed. It can be seen that the fit improves when there is no noise filter. This shows that the noise jump arises from bad fits when the TNF is somehow malfunctioning and transforms a double Gaussian distribution into a distribution with more complicated shape.

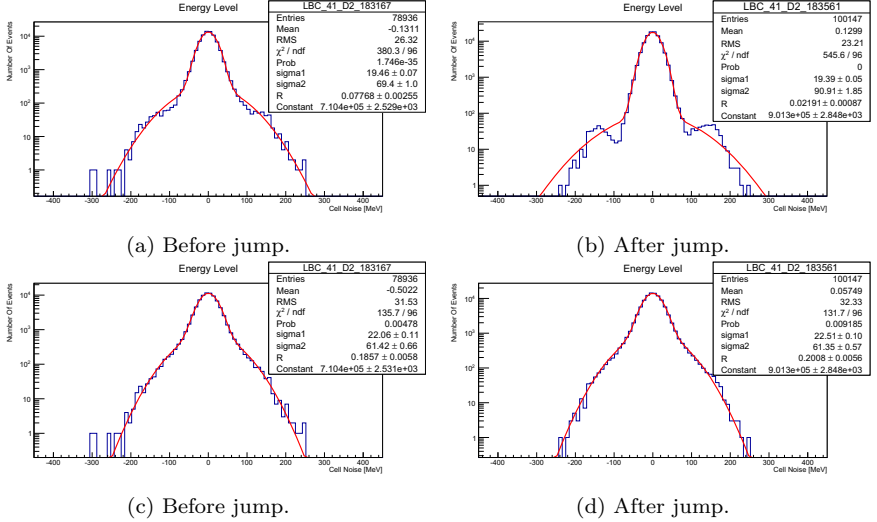


Figure 5.11: Comparison of the reconstructed pulse shape of a cell with the cell noise variation without a corresponding variation in the calibration constants, digital noise or channel status (jump) with the double Gaussian fit superimposed with the noise filter (Figures 5.11a and 5.11b) and without noise filter (Figures 5.11c and 5.11d).

5.5 Conclusions

In 2013 the ATLAS 2011 data was reprocessed with improved algorithms and calibrations. In TileCal the methods to produce Cs and laser calibration constants were improved and required a re-computation of the TileCal noise constants. Changing the calibration constants, according to Eq. (5.2.6), changes the reconstructed energy in the cell thus affecting the jet identification (see Section 6.5). A comparison between the ϕ -averaged RMS of electronic cell noise as a function of η of the cell before and after reprocessing for the high statistic run 192130 with both channels in high gain is shown in Figure 5.12. There is a general increase of the cell noise of about 0.5% in the EBC, LBC and EBA partitions. The LBA partition had module 22 running in *emergency mode*, i.e. operated with ~ 50 V less high voltage on the PMTs, and had the cesium calibration constants updated. This change is likely to have lowered the average noise in this partition.

The change of the cell noise between IOVs was monitored using a special software within the TUCS environment. A problem with the TNF has been

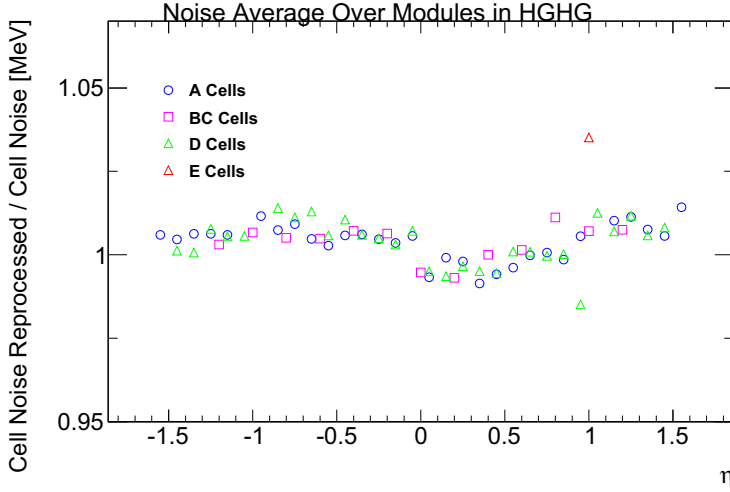


Figure 5.12: Comparison between the ϕ -averaged RMS of electronic cell noise as a function of η of the cell before and after reprocessing for the high statistic 192130 run with both channels in high gain. The different layers are shown separately, A, BC, D and E (gap/crack). The transition between the long and extended barrels can be seen in the range $0.7 < |\eta| < 1.0$.

identified in cells which exhibits a variation in the cell noise without a corresponding variation in other relevant quantities were present. The number of cells affected by this problem is investigated constructing the distribution of the ratio $\sigma / \text{RMS}_{\text{eff}}$ shown in Figure 5.13. The bulk of the distribution have $\sigma / \text{RMS}_{\text{eff}}$ close to one, this corresponds to cells for which the double Gaussian noise is a good model. There is nevertheless a big tail of cells with $\sigma / \text{RMS}_{\text{eff}}$ up to 1.7, with about 5% of the cells with $\sigma / \text{RMS}_{\text{eff}}$ larger than 1.2. After the installation of the new LVPS the contribution from a second wider Gaussian to the electronic noise is much smaller and the effect of the TNF is expected to be smaller as well. However the effect of the TNF with the new power supplies remains to be checked.

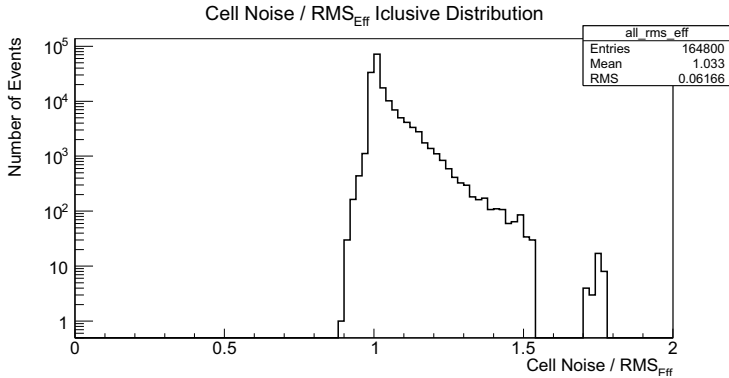


Figure 5.13: Distribution of the ratio $\sigma / \text{RMS}_{\text{eff}}$ where the σ and RMS_{eff} values are taken from all IOVs considered for the 2011 TileCal reprocessing.

Chapter 6

Objects Reconstruction for the Monojet Analysis

Object reconstruction is the process that associates the signal left in the detector by charged particles to physical objects through a series of algorithms. In 2015 this analysis used electrons, muons, jets and missing transverse momentum (E_T^{miss}) with the addition for the 2016 analysis of b -jets. Two types of electrons, muons and jets are also defined: *baseline* and *good*, where the former one is used for removal of overlapping objects and preselection while the latter for selecting the objects used to define the signal and control regions. In the following a brief introduction to the identification criteria of these objects with differences between the 2015 and 2016 analysis is presented.

6.1 Primary Vertex

In pp collisions with high luminosity, multiple interactions occur on a given bunch crossing, the spatial location of the a pp collision is called Primary Vertex (PV), the one that has the highest $\sum p_T^2$ of constituent tracks is known as Hard Scatter (HS) vertex, the rest are called *pile-up vertices*. The reconstruction of the PV generally happens in two stages that are often not distinguishable from each other. The *vertex finding* associates reconstructed tracks to a particular vertex candidate and the *vertex fitting* reconstructs the actual vertex position, refits the tracks and estimates the quality of the fit [73].

In the current analysis, events are required to have at least one PV with two associated tracks.

6.2 Track reconstruction

Charged particles that move through the solenoidal magnetic field that surrounds the ATLAS tracking system, follow helical trajectories. The projection of a helix on the xy plane is a circle and, in order to uniquely parametrize the particle track in three dimensions, five parameters are needed. A common

choice is to use the *perigee* parameters, where the perigee is the point of closest approach to the beam axis. With this choice, the five parameters are:

- The signed curvature C of the helix, defined as $C = q/2R$ where q is the particle charge and R is the radius of the helix. This is related to the transverse momentum $p_T = qB/C$, where B is the magnetic field measured in Tesla, C is measured in m^{-1} and p_T in GeV.
- The distance of closest approach d_0 in the xy plane measured in millimeters.
- The z coordinate of the point of closest approach, denoted by z_0 and measured in millimeters.
- The azimuthal angle ϕ_0 of the tangent to this point.
- The polar angle θ to the z -axis.

The perigee and the track parameters are schematized in Figure 6.1.

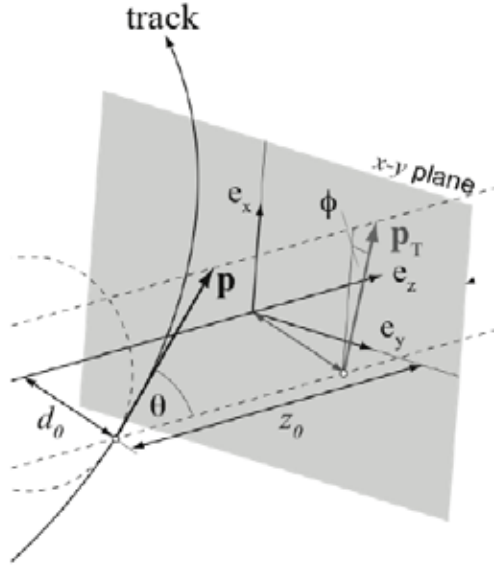


Figure 6.1: Track parameters at the perigee. In the figure $\mathbf{p}$ is the momentum of the incoming particle, e_x , e_y , e_z are the coordinate system defined in Section 4.2.1 and θ and ϕ , also defined in Section 4.2.1 are the polar and azimuthal angle respectively.

6.3 Electrons

Electrons are identified in the central part of the ATLAS detector ($|\eta| < 2.47$) by an energy deposit in the electromagnetic calorimeter and an associated track in the inner detector. Signal electrons are defined as prompt electrons coming from the decay of a W , Z boson or a top quark while background non-prompt

electrons come from hadron decays, photon conversion, semi-leptonic heavy flavor hadron decay and highly electromagnetic jets. To differentiate between signal and background electrons a likelihood discriminant is formed using different information: the shower shape in the EM calorimeter, the track-cluster matching, some of the track quality distributions from signal and background simulation and cuts on the number of hits in the ID. Cuts that depend on $|\eta|$ and E_T on the likelihood estimator allow to distinguish between signal and background electrons.

To further enhance the discrimination from background electrons, signal electrons are required to be separated from other detector activity. This is done by requiring the energy deposited by the electron in the calorimeter or the scalar sum of the transverse momentum of the tracks in the ID contained in a cone of some specified size ΔR to be less than a threshold value that depends on the tightness of the electron identification required. These track or calorimeter based criteria are commonly referred to as *isolation*.

Electron identification efficiencies are measured in pp collision data and compared to efficiencies measured in $Z \rightarrow ee$ simulations. Signal electrons can furthermore be selected with different sets of cuts for the likelihood-based criteria with $\sim 95\%$, $\sim 90\%$ and $\sim 80\%$ efficiency for electrons with $p_T = 40$ GeV. The different criteria are referred to as *loose*, *medium* and *tight* operating points respectively [74] where, for example, the tight criterion leads to a higher purity of signal electrons but has a lower efficiency than the looser criteria.

In this analysis, the *baseline electrons* are selected requiring a transverse energy $E_T > 20$ GeV, $|\eta| < 2.47$, they need to satisfy the *loose* likelihood selection criteria, it is required that no dead EM calorimeter Front-End Board (FEB) or High Voltage (HV) channels in the calorimeter cluster are present and that the baseline electron survives the overlap removal with other particles as described in Section 6.7. The baseline electron criteria is used to veto electrons used in the muon control regions and the signal region definition. In addition to all the baseline criteria, the *good electron* definition requires the electrons to satisfy the *tight* likelihood selection criteria, the electron track $d_0/\sigma_{d_0} < 5$ and $|z_0| < 0.5$ mm and the *LooseTrackOnly* electron isolation criteria which is based only on tracks and is set to be 99% efficient. Table 6.1 summarizes the selection criteria for the electrons in the monojet analysis.

Electron Definition	
Baseline electron	Good electron
$E_T > 20$ GeV	<i>baseline</i>
$ \eta < 2.47$	<i>tight</i> working point
<i>loose</i> working point	$d_0/\sigma_{d_0} < 5$
No dead FEB in the EM calo cluster	$ z_0 < 0.5$ mm
No dead HV in the EM calo cluster	<i>LooseTrackOnly</i>
passes the overlap removal	

Table 6.1: Electron definition for the monojet analysis.

6.4 Muons

Muons are identified using different criteria from the information provided by the ID and the MS leading to four different types of muons. The Standalone (SA) muons use only the MS information to reconstruct the muon’s trajectory; the Combined (CB), where the track is independently reconstructed in the ID and the MS and then combined; the Segment Tagged (ST) are identified as muons only if the track in the ID is, after being extrapolated to the MS, associated to at least one local track segment in the MDT or CSC chambers and finally the Calorimeter Tagged (CT) where tracks in the ID are associated to an energy deposit in the calorimeter compatible with a minimum ionizing particle. CB candidates perform best in terms of muon purity and momentum resolution.

Muons are identified using quality requirements specific to each of the four type of muons aiming at rejecting those coming from pion and kaon decays and guarantee a robust momentum measurement. The *loose* identification criteria maximize the reconstruction efficiency and provide good muon tracks; the *medium* criteria minimize the systematic uncertainties associated to muon reconstruction and calibration; the *tight* muons maximize the purity of the sample and the *high p_T* selection maximize the momentum resolution for tracks with transverse momenta above 100 GeV [75].

This analysis uses the CB muons that pass the medium identification criteria, moreover the *baseline muons* are required to have $p_T > 10$ GeV and $|\eta| < 2.5$, the E_T^{miss} definition and in the lepton veto used to define the signal and control regions. The *good muons* are required to pass the baseline selection criteria, moreover $d_0/\sigma_{d0} < 3$, $|z_0 \sin \theta| < 0.5$ mm. The good muons are used in the one muon and di-muon control regions. The definition of the baseline and good muons are summarized in Table 6.2.

Muon Definition	
Baseline muon	Good muon
CB muon	<i>baseline</i>
<i>Medium</i> id. criteria	$d_0/\sigma_{d0} < 3$
$p_T > 10$ GeV	$ z_0 \sin \theta < 0.5$ mm
$ \eta < 2.5$	
passes the overlap removal	

Table 6.2: Muon definitions for the monojet analysis.

6.5 Topocluster

The cell noise described in Section 5.2.2 is used in the *topological clustering* algorithm [76], for identification of the real energy deposits in the calorimeters. The algorithm assumes that the noise in all the calorimeter cells is normally distributed with significance (the ratio between the deposited energy and the

parameter σ used to describe the cell noise) expressed in units of Gaussian sigmas. In the algorithm, cluster of cells called *topoclusters*, are formed by comparing the energy deposit in a cell for a significant incompatibility with a noise only hypothesis. The algorithm starts by finding the *seed cells* with $E > 4\sigma$ where σ is the measured RMS of the energy distribution for every cell in the pedestal run. The second step is to add to the seeds neighbor cells that satisfy the $E > 2\sigma$ condition. Finally an additional level of cells with $E > 0$ is added to the perimeter of the cluster and the algorithm is ended. At this point the splitting algorithm is applied to separate the topoclusters based on the local energy maxima.

6.6 Jets

The quarks and gluons that carry color charge and are created in the scattering process, hadronize producing collimated bunches of colorless hadrons (jets) which keep track of the energy and direction of the originating parton. Jets in ATLAS are reconstructed as massless particles using the *anti- k_t* algorithm, calibrated and corrected for pile-up contamination. These steps are briefly outlined in the next sections.

6.6.1 The *anti- k_t* Algorithm

The anti- k_t algorithm [77] is a sequential recombination algorithm. It defines two distances d_{ij} and d_{iB} . The distance d_{ij} between the physical objects i and j is defined as:

$$d_{ij} = \min(k_{ti}^{-2}, k_{tj}^{-2}) \frac{\Delta_{ij}^2}{R^2} \quad (6.6.1)$$

where $\Delta_{ij}^2 = (\eta_i - \eta_j)^2 + (\phi_i^2 - \phi_j^2)$ and η_i is the rapidity, ϕ_i is the azimuthal angle, k_{ti} is the transverse momentum of the object i and R is the radius parameter that controls the size of the jet. The distance d_{iB} between the object i and the beam (B) defined as:

$$d_{iB} = k_{ti}^{-2}. \quad (6.6.2)$$

This distance is meant to distinguish between hard and soft terms. The algorithm identifies the smallest of the two distances and, if it is d_{ij} , it recombines the i and j objects while if it is d_{iB} , it calls i a jet and removes it from the list of inputs to the algorithm. The distances are recalculated and the procedure reiterated until there are no more objects.

6.6.2 The Jet Vertex Tagger

The excess transverse energy coming from pile-up jets is generally subtracted on average from the signal energy, however, due to localized fluctuations in the pile-up, some of it remains in the p_T of the reconstructed jet. The Jet Vertex Fraction (JVF) is a variable that uses information from the track associated to each jet to identify the origin vertex of each jet and rejects it if not coming from a hard-scatter vertex [78]. The JVF can be regarded as a measure of the

fraction of the jet momentum associated with a primary vertex and is defined as:

$$\text{JVF} = \frac{\sum_k p_T^{\text{trk}_k}(\text{PV}_0)}{\sum_l p_T^{\text{trk}_l}(\text{PV}_0) + \sum_{n \geq l} \sum_k p_T^{\text{trk}_k}(\text{PV}_n)} \quad (6.6.3)$$

where PV_0 is the hard scatter vertex and PV_n with $n \geq 1$ is any other pile-up PV in the same bunch crossing thus $\sum_k p_T^{\text{trk}_k}(\text{PV}_0)$ is the scalar p_T sum of the tracks associated with the jet and originating from the hard scatter vertex while $\sum_{n \geq l} \sum_k p_T^{\text{trk}_k}(\text{PV}_n)$ is the scalar p_T sum of the tracks associated with the pile-up vertexes.

Since the JVF denominator increases with the number of reconstructed PV, this introduces a pile-up dependence on the number of PV when minimal JVF selections are imposed in rejecting pile-up jets. To address this problem, two new variables to separate between HS and Pile Up (PU) jets are introduced: corrJVF and R_{pT} . The former is defined as:

$$\text{corrJVF} = \frac{p_T^{\text{HS}}}{p_T^{\text{HS}} + p_T^{\text{PU,corr}}} \quad (6.6.4)$$

where $p_T^{\text{HS}} = \sum_k p_T^{\text{trk}_k}(\text{PV}_0)$ is the scalar sum of the p_T of the tracks associated with the jet that comes from the hard scatter vertex and $p_T^{\text{PU,corr}} = \sum_{n \geq l} \sum_k p_T^{\text{trk}_k}(\text{PV}_n) / (kn_{\text{trk}}^{\text{PU}})$ is the scalar sum of the associated tracks originating from a pile-up vertex. Since the average p_T^{PU} increases linearly with the number of pile-up tracks, $n_{\text{trk}}^{\text{PU}}$, p_T^{PU} is divided by $(kn_{\text{trk}}^{\text{PU}})$ where $k = 0.01$ is the slope of the $\langle p_T^{\text{PU}} \rangle$ dependence with $n_{\text{trk}}^{\text{PU}}$ [79]. The corrJVF corrects the N_{PV} dependence in the JVF denominator.

Finally the quantity R_{pT} is defined for each jet as:

$$R_{pT} = \frac{\sum_k p_T^{\text{trk}_k}(\text{PV}_0)}{p_T^{\text{jet}}} \quad (6.6.5)$$

where p_T^{jet} is the fully calibrated jet p_T . This variable is defined using tracks associated with the vertex, it is at first order independent on the N_{PV} [80].

The Jet Vertex Tagger (JVT) is constructed from the corrJVF and R_{pT} by forming a two dimensional likelihood based on the k-nearest neighbor algorithm. Using the JVT algorithm, the HS jet efficiency is stable within 1% up to 35 interactions per bunch crossing [79]. Figure 6.2 shows the hard scatter jet efficiency dependence as a function of the average number of interactions per bunch crossing μ for a target signal efficiency of 95%. It can be seen that the JVT distribution, within statistical error, is flat and performs better than JVF at high pile-up as expected.

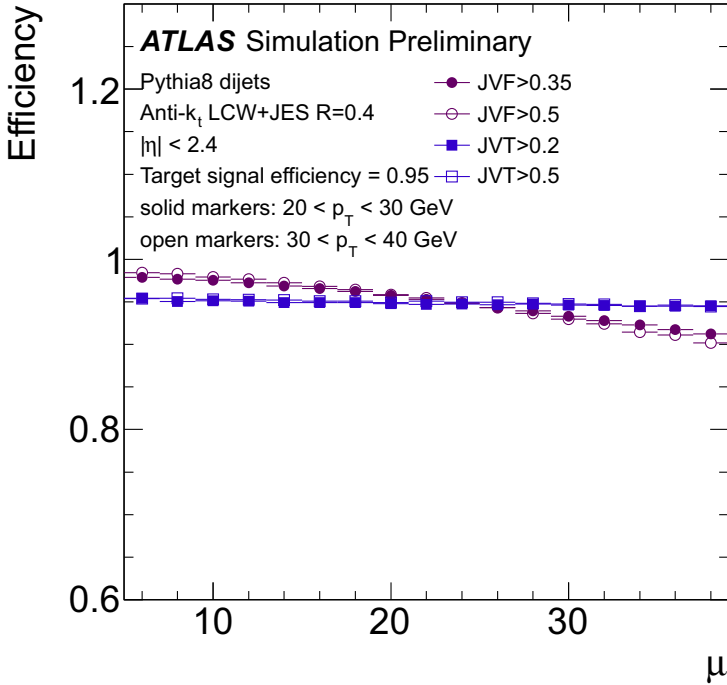


Figure 6.2: Dependence of the hard scatter jet efficiency as a function of the average number of interactions per bunch crossing μ . The solid markers refers to jets with $20 < p_T < 30$ GeV while the open markers to jets with $30 < p_T < 40$ GeV. Fixed cuts on JVT (blue) and JVF (purple) are imposed such that the inclusive efficiency is 95% [79].

6.6.3 Jet Calibration

The *electromagnetic scale* is the baseline signal scale of the ATLAS calorimeters, it is established during test beam measurements with electrons and it accounts correctly for the energy deposited by particle interacting electromagnetically. The topoclusters introduced in Section 6.5 are evaluated at the electromagnetic scale and used as an input to the anti- k_t algorithm outlined in Section 6.6.1. Due to detector effects such as the non compensating nature of the ATLAS hadronic calorimeter (see Section 5.1.2), energy loss due to inactive regions in the detector or particle showers not fully contained in the calorimeter, the energy of the jets measured at the electromagnetic scale is lower than the true energy of the original jet. These effects are calibrated using MC simulations and a correction, referred to as the Jet Energy Scale (JES), is applied in order to recover the correct energy scale of the jets [81].

The basic jet calibration scheme that applies the JES correction to the EM scale is usually referred to as EM + JES and uses p_T and η dependent corrections derived from simulations. The correction in this case is derived by

matching jets measured in the calorimeter with truth jets reconstructed in simulations, the ratio between the true jet energy and the electromagnetic scale matching one is taken as the correction factor [82]. The LCW briefly introduced in Section 4.2.6 is another calibration method that clusters topologically connected cells, classifying them as electromagnetic or hadronic and deriving the energy corrections from single pion MC simulation and dedicated studies to account for the detector effects. The jets calibrated with this method are referred to as LCW + JES. The Global Cell Weighting (GCW) calibration uses the fact that electromagnetic and hadronic showers leave a different energy deposition in the calorimeter cells, with the electromagnetic shower more compact than the hadronic one. The energy correction are then derived for each calorimeter cell within the jet and minimizing the energy resolution. The Global Sequential (GS) method starts from EM + JES calibrated jets and corrects for the fluctuation in particle content of the hadronic shower using the topology of the energy deposits. The corrections are applied in a way that leaves unchanged the mean jet energy [83].

6.6.4 Jet Selection

In order to distinguish jets coming from pp collisions from those from a non-collision origin, two jet selection criteria, *loose* and *tight*, are available. The loose selection criteria, provides an efficiency for selecting jets coming from pp collisions above 99.5% for $p_T > 20$ GeV, the tight criteria can reject even further background jets [84].

In this analysis the jets are reconstructed using the anti- k_t algorithm with the radius parameter $R = 0.4$. The *baseline jets* are required to have $|\eta| < 2.8$ and to make sure they come from a HS vertex, they need to satisfy any of the following:

A - The $p_T > 50$ GeV.

B - They have $20 < p_T < 50$ GeV and $|\eta| > 2.4$.

C - They have $20 < p_T < 50$ GeV, $|\eta| < 2.4$ and JVT > 0.64 .

Furthermore, events in which the jets, after the overlap removal (see Section 6.7) is applied, fail the loose selection criteria are disregarded. Finally the most energetic jet in the event (the *leading jet*) is required to pass the tight selection criteria. The *good jets* require an increased p_T threshold of 30 GeV and at most 4 HS jet in the event.

In the 2016 analysis the A, B and C conditions above are replaced with:

A - $p_T > 30$ GeV and $|\eta| < 2.8$.

B - $p_T > 60$ GeV or $|\eta| > 2.4$ or JVT > 0.59 .

The selection criteria for both analyses are summarized in Table 6.3.

Jet Definition	
Baseline jet	Good jet
$R = 0.4$	<i>baseline</i>
$ \eta < 2.8$	at most 4 jets
<i>Loose</i> selection criteria	$p_T > 30$ GeV (2015)
<i>Tight</i> on the leading jet	
passes the overlap removal	
any of (2015):	
• $p_T > 50$ GeV	
• $20 < p_T < 50$ GeV, $ \eta > 2.4$	
• $20 < p_T < 50$ GeV, $ \eta < 2.4$, JVT > 0.64	
any of (2016):	
• $p_T > 30$ GeV and $ \eta < 2.8$	
• $p_T > 60$ GeV or $ \eta > 2.4$ or JVT > 0.59	

Table 6.3: Jet definition in the monojet analysis. The year in parentheses refers to the selection applied in different versions of the analysis.

6.6.5 Identification of b and c Jets

Jets coming from the hadronization of a b quark are commonly referred to as b -jets and their identification as b -tagging. A multivariate discriminant, the MV2, is used in ATLAS to discriminate b -jets from light- (u , d , s or gluon) and c -jets (jets coming from the hadronization of the c quark). Information from secondary vertexes, b -hadron decay chains and a log likelihood ratio based on the probability of the tracks to originate from given flavor quark are combined in the multivariate discriminant [85]. In the 2016 analysis the bottom quarks are identified using the MV2c10 algorithm where the b -jets are treated as signal and the background is a mixture of 7% c quarks and 93% light quarks. The presence of c quarks as background in the training sample is low in order to enhance the charm rejection. If a jet has a $p_T > 30$ GeV and $\eta < 2.5$ it is within the b -tagging acceptance and its MV2c10 discriminant can be calculated, the jet is identified as a b -jet if this value is above 0.9349 corresponding to a 60% efficiency working point.

6.7 Overlap Removal

During object reconstruction, it may happen that different algorithms identify the same track and cluster as different types of particles, this results in a duplicate object. In physics analyses a decision must be made on which interpretation to give to the reconstructed object, this process is called Overlap Removal (OR) [86].

In this analysis, an overlap removal is applied to electrons, muons and jets that pass the baseline criteria and the following objects are removed:

- A - Remove jet in case any pair of jet and electron satisfies $\Delta R(j, e) < 0.2$.

- B - Remove electron in case any pair of jet and electron satisfies $0.2 < \Delta R(j, e) < 0.4$.
- C - Remove muon in case any pair of muon and jet with at least 3 tracks satisfies $\Delta R(j, \mu) < 0.4$.
- D - Remove jet if any pair of muon and jet with less than 3 tracks satisfies $\Delta R(j, \mu) < 0.4$.

In the 2016 analysis the A condition is replaced by:

- E - In case any pair of jet and electron or muon satisfy $\Delta R(j, e \text{ or } \mu) < 0.2$, the b -tagging is adjusted to the 85% efficiency working point and:
 - The jet is not a b -jet: remove the jet and keep the electron or muon.
 - The jet is a b -jet: since the jet is likely coming from a semi-leptonic b decay, keep the jet and remove the electron or the muon.

6.8 Missing Transverse Energy

Due to the conservation of momentum and the fact that the proton bunches are parallel to the z -axis, the sum of the momenta of the collision products in the transverse plane should sum to zero. Any energy imbalance is known as missing transverse momentum (E_T^{miss}), it may indicate weakly interacting stable particles (neutrinos within the SM, new particles in beyond SM models) or non reconstructed physical objects that escape the detector acceptance like for example a muon that goes into a hole of the muon system (e.g. supports for ATLAS) cannot be detected and give rise to fake E_T^{miss} . Physical objects that are fully reconstructed and calibrated such as electrons, photons, hadronically decaying tau-leptons, jets or muons are called *hard objects* and are used to compute the missing transverse momentum in an event [87]. The x and y components of the E_T^{miss} can be written as:

$$E_{x(y)}^{\text{miss}} = E_{x(y)}^{\text{miss, e}} + E_{x(y)}^{\text{miss, } \gamma} + E_{x(y)}^{\text{miss, } \tau} + E_{x(y)}^{\text{miss, jets}} + E_{x(y)}^{\text{miss, } \mu} + E_{x(y)}^{\text{miss, soft}} \quad (6.8.1)$$

where the terms for jets, charged leptons and photons are the negative sum of the momenta of the respective calibrated object while the *soft term* is reconstructed from the transverse momentum deposited in the detector that is not already associated to hard objects. It may be reconstructed by means of calorimeter-based methods, the so called Calorimeter Soft Term (CST), or using track-based methods known as Track Soft Term (TST).

The CST is reconstructed using energy deposits in the calorimeters which are not associated to hard objects, it arise from soft radiation accompanying the hard scatter event and from underlying event activity. A downside of the CST is its vulnerability to pile up.

The TST is built from tracks not associated to any hard object, tracks can be associated to hard scatter vertices and thus to a particular pp collision, making this method robust against pile-up. This method is, however, insensitive to soft terms coming from neutral particles that do not leave a track in the

ID, thus the TST E_T^{miss} is combined with calorimeter-based measurements for hard objects.

Due to its stability against pile-up, this analysis uses the TST E_T^{miss} term. Moreover in the monojet analysis the muons are treated as invisible particles in the E_T^{miss} reconstruction so that the missing energy is calculated as:

$$E_{x(y)}^{\text{miss}} = E_{x(y)}^{\text{miss, e}} + E_{x(y)}^{\text{miss, } \gamma} + E_{x(y)}^{\text{miss, } \tau} + E_{x(y)}^{\text{miss, jets}} + E_{x(y)}^{\text{miss, soft}} \quad (6.8.2)$$

where $E_{x(y)}^{\text{miss, } \mu} = 0$ is used in Eq. (6.8.1).

Chapter 7

Search for Compressed Supersymmetry Models with 2015 ATLAS Data

Many experimental SUSY searches are not sensitive to models where the mass difference between some of the supersymmetric particles is too small. The search carried out using the data collected by the ATLAS detector in 2015 and presented in this chapter focus on this situation, in particular for squark pair production, where the neutralinos are only slightly lighter than the squark.

7.1 Missing Transverse Energy and SUSY Particle Masses

There are two possible cases that results in an energy imbalance in the detector, the first one occurs in beyond Standard Model physics, that involves the presence of particles that interact weakly or not at all with normal matter. These particles are not detected thus leave an energy imbalance in the detector. In the second case, the decay products in the final state involve neutrinos that are not detectable by ATLAS. To better understand the first category of events, one considers Figure 7.1a that shows the decay topology of squark pair production with a neutralino and two jets in the final state. Using the two body decay energy and momentum relations [42]:

$$E_q = \frac{M_{\tilde{q}}^2 - m_{\tilde{\chi}_1^0}^2 + m_q^2}{2M_{\tilde{q}}}, \quad (7.1.1)$$

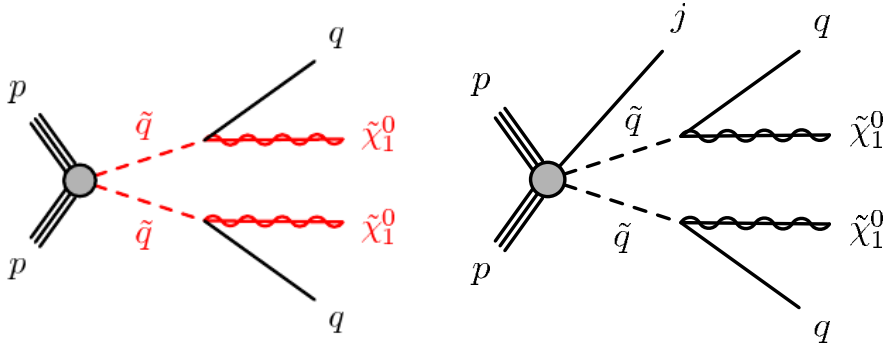
$$|\vec{p}_q| = |\vec{p}_{\tilde{\chi}_1^0}| = \frac{\left[\left(M_{\tilde{q}}^2 - (m_q + m_{\tilde{\chi}_1^0})^2 \right) \left(M_{\tilde{q}}^2 - (m_q - m_{\tilde{\chi}_1^0})^2 \right) \right]^{1/2}}{2M_{\tilde{q}}} \quad (7.1.2)$$

where $M_{\tilde{q}}$ is the squark center of mass energy, $m_{\tilde{\chi}_1^0}$ is the neutralino mass and m_q is the quark mass. Neglecting the quark mass ($m_q = 0$) we get that:

$$E_q = \frac{M_{\tilde{q}}^2 - m_{\tilde{\chi}_1^0}^2}{2M_{\tilde{q}}}, \quad (7.1.3)$$

$$|\vec{p}_q| = |\vec{p}_{\tilde{\chi}_1^0}| = \frac{M_{\tilde{q}}^2 - m_{\tilde{\chi}_1^0}^2}{2M_{\tilde{q}}}. \quad (7.1.4)$$

The quark will hadronize and the corresponding hadron will shower in the calorimeter resulting in a jet that can be detected by the ATLAS detector. If for example the mass of the squark is $M_{\tilde{q}} = 450$ GeV and the mass of the neutralino is $m_{\tilde{\chi}_1^0} = 445$ GeV then it can be seen from Eq. (7.1.4) that the quark and neutralino momenta are given by $|\vec{p}_q| = |\vec{p}_{\tilde{\chi}_1^0}| \simeq 5$ GeV. Since the neutralino escape detection, this results in low E_T^{miss} thus when the mass of the neutralino approaches the mass of the quark this results in a low energy jet and E_T^{miss} . Since the missing energy resolution of the ATLAS detector is approximately 15 GeV in data as can be seen from Figure 7.2, such events can be difficult to trigger on and cannot be extracted from the multijet background using missing transverse energy. This means that there is no sensitivity to SUSY models with compressed mass spectra (when the mass difference between the particles is small). This problem applies to several SUSY production channels.



(a) Event without initial state radiation [88]. (b) Event with initial state radiation [89].

Figure 7.1: Event topology of squark pair production resulting in a neutralinos with two jets final state with (Figure 7.1b) and without (Figure 7.1a) initial state radiation.

Figure 7.3 illustrates this effect for the search for squark pair production in the case of the squark decaying directly to a quark and a neutralino through the mechanism illustrated in Figure 7.1a. This search uses a classical multijet + E_T^{miss} analysis, it can be seen that there is no sensitivity close to the diagonal (dashed line) in the region $400 < M_{\tilde{q}} < 600$ GeV.

If an initial state radiation jet is present in the event, as depicted in Figure 7.1b, the squark-squark system gets boosted in the opposite direction thus increasing the momentum of the decay products and the missing energy leading

to a signature of a high p_T jet on one side and additional jets and E_T^{miss} on the other side of the event.

Events with an energetic jet p_T and large E_T^{miss} in the final state constitute a clean signature for new physics searches at hadron colliders. Other BSM signals that can be studied with this experimental signature include the production of WIMPs, the ADD model for large extra dimensions and SUSY.

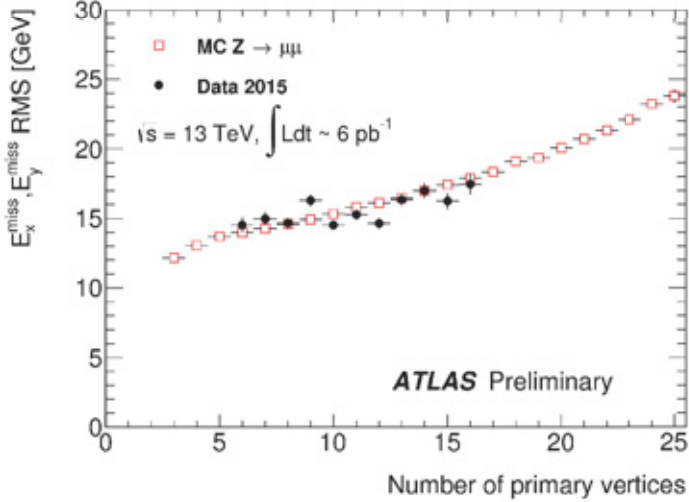


Figure 7.2: Distribution of the x and y components of the TST E_T^{miss} resolution as a function of the number of primary vertices in $Z(\rightarrow \mu^+\mu^-) + \text{jets}$ events [90].

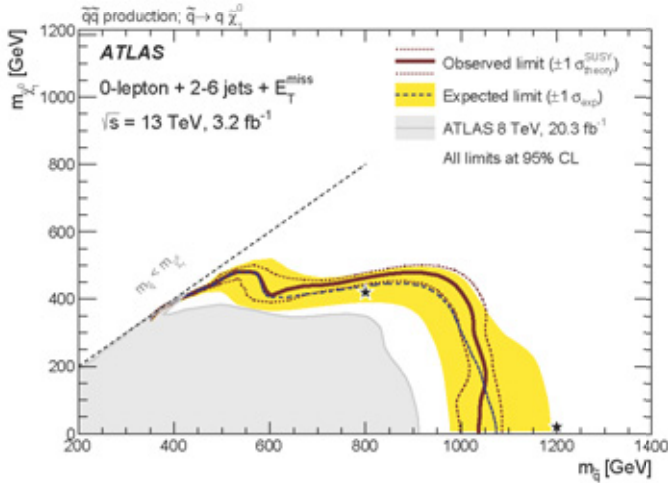


Figure 7.3: Exclusion limits for direct production of squark pairs where the squark decays into a quark and a neutralino. The x -axis represents the mass of the squark and the y -axis represents the mass of the lightest neutralino. The black stars represent a benchmark model as explained in more details in Ref. [88].

7.2 Event Selection

The search for squark pair production with compressed mass spectrum is carried out in pp collisions using the data collected by the ATLAS experiment during the 2015 Run II corresponding to a total integrated luminosity of 3.2 fb^{-1} . The signal region is defined by the following selection criteria:

- A - The HLT_xe70 trigger was used in the whole 2015 dataset.
- B - In order to ensure that the event originated from a pp collision, a primary vertex with at least two associated tracks with $p_T > 0.4 \text{ GeV}$ is required.
- C - Events in which any jet fails the *loose* jet cleaning criteria are rejected. This suppress noise from non-collision background.
- D - The most energetic jet in the event (the *leading jet*) must have $p_T > 250 \text{ GeV}$ and $|\eta| < 2.4$. Moreover, in order to reject beam-induced and cosmic particles background, the event is rejected if the leading jet fails the *tight* cleaning criteria.
- E - In order to suppresses $Z(\rightarrow \ell\bar{\ell}) + \text{jets}$ and $W(\rightarrow \ell\nu)$ background, events with an identified electron of $p_T > 20 \text{ GeV}$ or muon of $p_T > 10 \text{ GeV}$ are rejected, see Sections 6.3 and 6.4 for the lepton definitions.
- F - In order not to overlap with other ATLAS SUSY searches, events with more than four jets are rejected, see Section 7.3 for more details on this selection.
- G - The E_T^{miss} trigger can select multi-jet events in case of a mis-reconstructed jet. In these cases the missing transverse momentum points in the direction of one of the jets, this background can be suppressed by imposing a minimum azimuthal angle separation between the missing transverse momentum and any jet of $\Delta\phi(\text{jets}, E_T^{\text{miss}}) > 0.4$.
- H - A $E_T^{\text{miss}} > 250 \text{ GeV}$ requirement is imposed in order to be able to test different BSM signals with different sensitivities to the missing energy.
- I - Events with an identified electron of $p_T > 20 \text{ GeV}$ or muon of $p_T > 10 \text{ GeV}$ are vetoed in order to suppress the $Z(\rightarrow \ell\bar{\ell}) + \text{jets}$ and $W(\rightarrow \ell\nu)$ background. See Sections 6.3 and 6.4 for the lepton definitions.

Inclusive (IM1-IM7) and exclusive (EM1-EM6) Signal Regions (SRs) are defined in the monojet analysis with increasing E_T^{miss} thresholds from 250 GeV to 700 GeV, see Table 7.1 for the exact definition of the signal regions. These different E_T^{miss} bins are defined in order to address different BSM signals tested with the monojet signature. In this chapter special emphasis is placed on the compressed squark-neutralino model.

Event Selection Criteria							
HLT_xe70 trigger Primary Vertex $E_T^{\text{miss}} > 250$ GeV Leading jet with $p_T > 250$ GeV and $ \eta < 2.4$ At most 4 jets $\Delta\phi(\text{jets}, E_T^{\text{miss}}) > 0.4$ Jet quality requirements No identified muon with $p_T > 10$ GeV or electron of $p_T > 20$ GeV							
Inclusive SRs	IM1	IM2	IM3	IM4	IM5	IM6	IM7
E_T^{miss} [GeV]	> 250	> 300	> 350	> 400	> 500	> 600	> 700
Exclusive SRs	EM1	EM2	EM3	EM4	EM5	EM6	
E_T^{miss} [GeV]	[250-300]	[300-350]	[350-400]	[400-500]	[500-600]	[600-700]	

Table 7.1: Definition of the signal regions.

7.3 Jet Veto

In order to reduce the overlap with other ATLAS data analyses, especially those containing multijet final states, an upper cut on the number of HS jets is necessary. This procedure is denominated *jet veto*. In order to define hard scatter jets, cuts on the jet p_T and on the JVT (see Section 6.6.2) have been studied. The aim of these cuts is to render the analysis independent from pile-up as much as possible but stay signal efficient, too hard cuts would lead to disregarding possible signal events while, on the other hand, too loose cuts let pile-up jets in the signal region and allow significant overlap with other ATLAS searches.

Due to the jets coming from the squark decays, the SUSY compressed squark-neutralino models have more jets compared to other signals considered in the monojet analysis and are therefore the most sensitive to jet veto efficiency and to the definition of hard scatter jet. To study the pile-up contamination for different p_T thresholds a *figure of merit* is defined and its dependence from the average number of proton-proton collisions per bunch crossing studied. The figure of merit is referred to as *jet veto efficiency* and defined as:

$$\frac{\text{N (events) with baseline cuts + at most N (jet) HS jets}}{\text{N (events) with baseline cuts}} \quad (7.3.1)$$

where the baseline cuts are the event selection given in Section 7.2 except the cut F (the study presented here uses a symmetric E_T^{miss} and leading jet p_T cut at 250 GeV) and the hard scatter jets are defined as:

- Jets with $p_T > 50$ GeV is always considered coming from the hard scatter.
- Due to the lacking of a tracking system in the $|\eta| > 2.4$ region and thus the impossibility of using a JVT criteria, jets belonging to the $|\eta| > 2.4$ (forward jets) are always considered as coming from the hard scatter.

- In order to be considered coming from the hard scatter proton-proton collision, the jets with $p_T^{\text{thresh}} < p_T < 50$ GeV must additionally satisfy a $\text{JVT} > \text{JVT}^{\text{thresh}}$ selection criteria where different p_T^{thresh} and $\text{JVT}^{\text{thresh}}$ are studied, in particular:

A - $p_T^{\text{thresh}} \in \{30 \text{ GeV}, 40 \text{ GeV}, 50 \text{ GeV}, 70 \text{ GeV}\}$,

B - $\text{JVT}^{\text{thresh}} \in \{0.14, 0.64, 0.92\}$.

As can be seen from Figure 7.1b, at least three jets in the event are required for the SUSY compressed spectra models, on the other hand other searches [91] already explore signatures with four jets or more. For these reasons and to increase the acceptance of the SUSY compressed models, events with at most three or four jets are studied with different p_T thresholds. Figures 7.4a and 7.4c shows the jet veto efficiency for the different values of p_T^{thresh} and number of jets in the case where no JVT cut is applied. With a maximum number of three 30 GeV jets the signal efficiency loss is as large as 20%. The JVT cut reduces the loss of efficiency for the signal, by preventing pile-up jets to trigger a veto of the signal events. For three jets at 30 GeV the loss of signal is reduced from about 20% in Figure 7.4a to approximately 15% in Figure 7.4b and becomes less than 5% for at most four jets. From this study and similar

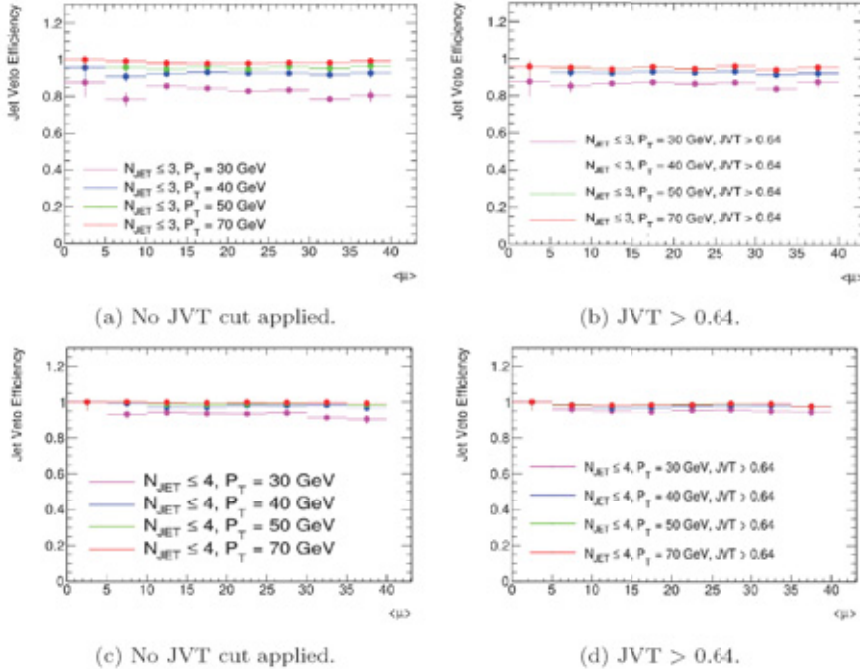


Figure 7.4: Jet veto efficiency for different jet p_T thresholds and number of jets as a function of the average number of interactions per bunch crossing (μ) for a compressed spectra model point $m_{\tilde{q}} = 450$ GeV $m_{\tilde{\chi}_1^0} = 435$ GeV. In Figures 7.4a and 7.4c no JVT cut is applied and there is some drop of the efficiency at high pile-up. In Figures 7.4b and 7.4d a JVT > 0.64 cut is applied, the dependence from pile-up is reduced.

ones on other signals explored by this analysis, it was decided to select events with at most four jets above 30 GeV. Even though harder cuts on the p_T of the jets provide a better efficiency, such a choice would overlap with the analysis given in Ref. [91]. Figure 7.5 shows the same behavior described above on the $Z \rightarrow \nu\bar{\nu}$ background. Also in this case an improvement of the jet veto efficiency as function of the average number of interaction per crossing is seen especially for $p_T = 30$ GeV. The selected definition of the hard-scatter jets, with $N(\text{jets}) \leq 4$ is pile-up independent.

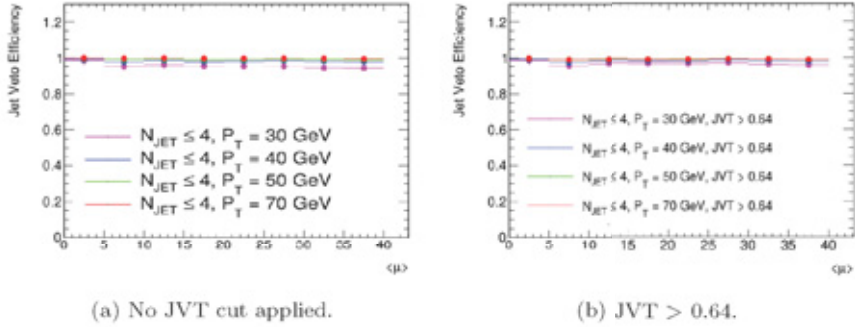


Figure 7.5: Jet veto efficiency for different jet p_T thresholds and $N(\text{jets}) \leq 4$ as a function of the average number of interactions per bunch crossing $\langle\mu\rangle$ for the $Z \rightarrow \nu\bar{\nu}$ background. In Figure 7.4c no JVT cut is applied and there is some drop of the efficiency at high pile-up. In Figure 7.4d a JVT > 0.64 cut is applied, the dependence from pile-up is reduced.

7.4 Sources of Background

The Standard Model process $Z \rightarrow \nu\bar{\nu} + \text{jets}$ where the neutrinos escape detection generating large E_T^{miss} is experimentally similar to signal events in the search for BSM physics with a high energy jet and large missing energy signature in the final state. This process constitutes an irreducible background and is the largest background in the analysis. The key to estimate the $Z \rightarrow \nu\bar{\nu} + \text{jets}$ background is to be able to predict the momentum spectrum of the Z boson and, since this is assumed to be equal to the E_T^{miss} , the corresponding amount of missing transverse energy. The control regions defined in later sections provide ways to derive the Z boson momentum using data control regions.

The second biggest source of background with a monojet topology is the $W(\rightarrow \ell\nu) + \text{jets}$ and in particular $W(\rightarrow \tau\nu) + \text{jets}$, in which the τ decays hadronically, gives the highest contribution to this class of events. The $W(\rightarrow \ell\nu) + \text{jets}$ events with $\ell = e$ or μ can fake a monojet signature if the lepton escapes the detector acceptance or has a quality that is lower than the quality of the lepton veto criteria.

Other background processes are diboson (WW , WZ and ZZ production), $t\bar{t}$ and single top processes. Multi-jet production from QCD processes where one or more jets are mis-reconstructed leading to high E_T^{miss} represent a small source of background. Finally Non Collision Background (NCB) coming from cosmic particles, detector noise and beam-induced background can give rise

to fake jets and consequently to E_T^{miss} . As discussed in more details in the upcoming sections the control regions will be used to effectively scale the Monte Carlo predictions for the $V + \text{jets}$ processes using normalization factors (see Section 7.5.4).

7.5 Estimation of the $Z + \text{jets}$ and $W + \text{jets}$ backgrounds

A Control Region (CR) is a region of the phase space where the signal contribution is negligible but the event selections are similar to those of the signal region so that the normalization of the backgrounds in the signal region can be deduced from observed number in the control regions. The main reason to define control regions is to compute and check the agreement in shape and normalization between MC simulations and data in reconstructed kinematic quantities. The $V + \text{jets}$, where V is either a W or a Z vector boson, constitutes the main background of the monojet analysis. The Monte Carlo simulations for the $V + \text{jets}$ processes are generated using SHERPA [92] with 0, 1 or 2 jets at Next to Leading Order (NLO) and 2 or 4 jets at Leading Order (LO). A pure MC estimation of these processes, suffers from theoretical uncertainties like the limited knowledge of the Parton Distribution Functions (PDFs), limited order accuracy of the Monte Carlo generators and experimental uncertainties related to the jet energy scale and luminosity determination. In order to estimate the contribution of these backgrounds in the SR, a *data driven* technique is used. The method aims at reducing the systematic uncertainties by relying on information from data in the CRs rather than on MC simulations. It can be divided in three major steps:

- Define CRs to select $V + \text{jets}$ event in data.
- Calculate a transfer factor from MC predicted events in the CR to background estimates in the SR.
- Apply the transfer factor to the observed events in the CR to obtain the estimate number of events from the process in the SR.

The CRs used to constrain the $V + \text{jets}$ backgrounds have an event selection that differs from the SR only in the lepton veto and the missing transverse momentum calculation. Owing to the different lepton selections the CRs and SRs are orthogonal and a minimum contribution from a monojet-like signal is expected.

7.5.1 The Muon Control Region

The muon control region (CR_{wmn}) is used to estimate the $W(\rightarrow \mu\nu) + \text{jets}$ and the $Z(\rightarrow \nu\bar{\nu}) + \text{jets}$ background contribution in the SR. The $W(\rightarrow \mu\nu) + \text{jets}$ events can enter the SR if the muon is outside the detector acceptance or it fails the muon selection criteria presented in Section 6.4. $W(\rightarrow \mu\nu) + \text{jets}$ events are selected and used in order to estimate the $Z(\rightarrow \nu\bar{\nu}) + \text{jets}$ contribution in

the signal region. The muon is treated like a neutrino in the E_T^{miss} calculation, in this way the E_T^{miss} is a measurement of the W momentum which is later translated into the Z boson momentum to estimate the $Z(\rightarrow \nu\bar{\nu}) + \text{jets}$ background. Moreover, since the muon energy deposit in the calorimeter is tiny, the calorimeter activity in $W(\rightarrow \mu\nu) + \text{jets}$ events is similar to $Z(\rightarrow \nu\bar{\nu}) + \text{jets}$ events. In addition to cuts from A to H defined in Section 7.2 the CR_{wmn} region selects events with:

- Exactly one good muon.
- Other baseline muons and baseline electrons are vetoed.
- The transverse mass, defined as:

$$m_T = \sqrt{2p_T^\mu p_T^\nu (1 - \cos(\phi_\mu - \phi_\nu))} \quad (7.5.1)$$

and determined by the muon p_T (p_T^μ) and neutrino p_T (p_T^ν), is required to be $30 < m_T < 100$ GeV, consistent with W boson production. The neutrino p_T is calculated assuming that $p_T^\nu = E_T^{\text{miss}}$ where the missing energy is calculated according to Eq. (6.8.1).

The transverse mass cut suppress the $W(\rightarrow \tau\nu) + \text{jets}$ processes in this region. The measured E_T^{miss} and leading jet p_T distributions after the background only fit procedure described in Section 7.5.4 are shown in Figure 7.6. The agreement between data and MC is good.

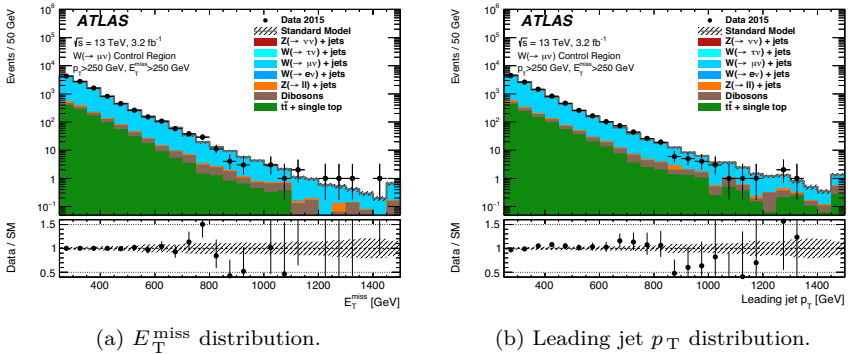


Figure 7.6: Observed and predicted E_T^{miss} and leading jet p_T distributions after the background only fit in the single muon CR_{wmn} for the $E_T^{\text{miss}} > 250$ GeV selection. The error bands include the statistical and systematic errors.

7.5.2 The Electron Control Region

The $W(\rightarrow e\nu) + \text{jets}$ process enters the signal region, thus contributing to the background, in case the electron is not identified in the detector. In the $W(\rightarrow \tau\nu) + \text{jets}$ process the τ particle can decay hadronically in 65% of the cases resulting in additional jets that can help this background mimic the signal. To address these backgrounds an electron control region (CR_{ele}) is built. It

is designed to constrain both $W(\rightarrow e\nu) + \text{jets}$ and $W(\rightarrow \tau\nu) + \text{jets}$ processes thanks to the decays of τ leptons into electrons. In order to efficiently reject the multi-jet background, the electron is retained in the E_T^{miss} calculation, the missing transverse momentum in this case measures the momentum of the escaping neutrino. In addition to cuts from A to H defined in Section 7.2 the CR_{ele} region selects events with:

- Exactly one good electron.
- No additional baseline electrons or baseline muons in the final state.

In order to enhance the $W(\rightarrow \tau\nu) + \text{jets}$ contribution in this region, no m_T cut is applied. Figure 7.7 shows the observed and predicted E_T^{miss} and the leading jet p_T distribution in this control region. The overall agreement between data and MC is good and improved after the global likelihood fit procedure described in Section 7.5.4.

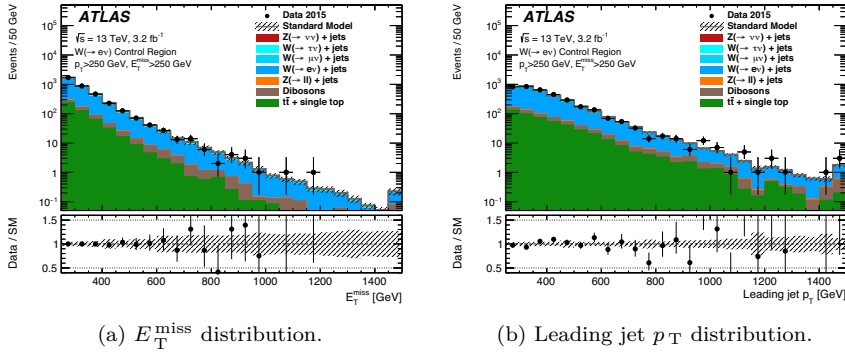


Figure 7.7: Observed and predicted E_T^{miss} and leading jet p_T distributions after the background only fit in the electron CR_{ele} for the $E_T^{\text{miss}} > 250$ GeV selection. The error bands include the statistical and systematic errors.

7.5.3 The Di-Muon Control Region

This region (CR_{zmm}) is designed to select the $Z(\rightarrow \mu^+\mu^-) + \text{jets}$ events in order to estimate their contribution to the background in the signal region. In addition to cuts from A to H defined in Section 7.2, events in the CR_{zmm} region are required to have:

- Exactly two good muons.
- Electrons are vetoed.
- An invariant mass in the Z boson mass window range $66 < m_{\mu\mu} < 116$ GeV.

Figure 7.8 shows a good agreement between data and MC for the measured E_T^{miss} and leading jet p_T distributions also in this region.

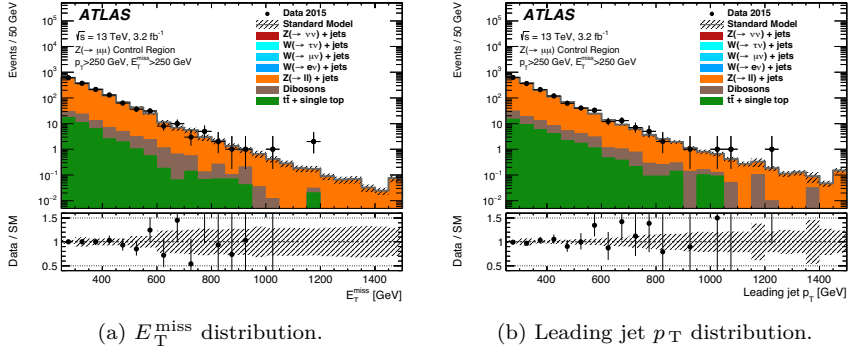


Figure 7.8: Observed and predicted E_T^{miss} and leading jet p_T distributions after the background only fit in the di-muon CR_{Zmm} for the $E_T^{\text{miss}} > 250$ GeV selection. The error bands include the statistical and systematic error.

7.5.4 The Global Simultaneous Likelihood Fit

A pure MC based prediction of the background (BG) contamination in the signal region yield large theoretical and experimental systematic uncertainties in the shape and normalization of the predicted distributions. For this reason a data driven approach, where the different backgrounds are normalized using the data in CRs that are orthogonal to the SR, is used. To extrapolate from CR to SR prediction a *transfer factor* computed in MC is used. This factor is a ratio of MC predictions in regions with similar selections on missing transverse energy and jets thus most of the systematic uncertainties either cancel out or are significantly reduced. The extrapolation from one kinematic region to another can be itself a source of uncertainties. To avoid this, the selection in terms of kinematic in the CRs and SRs is chosen to be the same. The expected contributions of a background process BG_i , which is extracted from a control region CR_j to a signal region SR_k , is given by:

$$N_{\text{BG}_i}^{\text{SR}_k} = \frac{(N_{\text{data}}^{\text{CR}_j} - N_{\text{non BG}_i, \text{MC}}^{\text{CR}_j})}{N_{\text{BG}_i, \text{MC}}^{\text{CR}_j}} \times N_{\text{BG}_i, \text{MC}}^{\text{SR}_k} \quad (7.5.2)$$

where $N_{\text{BG}_i}^{\text{SR}_k}$ is the control region driven predicted number events for background i events in the signal region SR_k , $N_{\text{data}}^{\text{CR}_j}$ is the observed number of events in the control region j , $N_{\text{non BG}_i, \text{MC}}^{\text{CR}_j}$ is the estimated number of the background contamination coming from other processes in the given control region, $N_{\text{BG}_i, \text{MC}}^{\text{CR}_j}$ is the MC prediction of the background i in the control region j and $N_{\text{BG}_i, \text{MC}}^{\text{SR}_k}$ is the number of the background i events predicted in MC simulation in the signal region. The ratio:

$$C_{\text{BG}_i}^{\text{CR}_j \rightarrow \text{SR}_k} = \frac{N_{\text{BG}_i, \text{MC}}^{\text{SR}_k}}{N_{\text{BG}_i, \text{MC}}^{\text{CR}_j}} \quad (7.5.3)$$

is the transfer factor used to extrapolate from CR_j to SR_k . The factors used in the normalization of the background expectation in the SRs are bin dependent

and given by:

$$\mu_j = \frac{(N_{\text{data}}^{\text{CR}_j} - N_{\text{non BG}_i, \text{MC}}^{\text{CR}_j})}{N_{\text{BG}_i, \text{MC}}^{\text{CR}_j}}. \quad (7.5.4)$$

The different normalization factors are not independent since different processes can enter several CRs and thus the background subtraction term, $N_{\text{non BG}_i, \text{MC}}^{\text{CR}_j}$, gets contributions from the other normalization factors. To properly treat the correlations, the normalization factors are obtained from a simultaneous fit of all the CRs referred to as the *global fit*. A summary of the different processes and the corresponding CR used to extract the normalization factor is given in Table 7.2.

The global fit can be performed in two different ways:

1. The background only hypothesis, fits only the control regions in order to predict the background in the signal region. This fit is used to set model independent limits.
2. The signal plus background hypothesis, fits both the signal and control regions with a sum of background and specific signal. The normalization of the specific BSM signal is a free parameter. This fit is used to exclude specific models.

Control Region	Background Process	Normalization Factor
CR_{wmn}	$W(\rightarrow \mu\nu), Z(\rightarrow \nu\bar{\nu})$	μ_{wmn}
CR_{ele}	$W(\rightarrow e\nu), W(\rightarrow \tau\nu), Z/\gamma^*(\rightarrow \tau^+\tau^-)$	μ_{ele}
CR_{zmm}	$Z/\gamma^*(\rightarrow \mu^+\mu^-)$	μ_{zmm}

Table 7.2: Summary table of the different background processes and the corresponding control regions used to evaluate the bin dependent normalization factors.

7.6 Other Backgrounds

7.6.1 The Multi-Jet Background

The multi-jet background events are mainly due to mis-reconstruction or loss in some dead part of the calorimeter of jets or to the presence of neutrinos in some heavy-flavor hadronic decay. The acceptance in the SR for multi-jet events is low, nevertheless the large cross section of this process could potentially lead to a high contamination in the SR. The use of MC simulation in order to estimate the contribution to BG of this process is very difficult due to the very large MC samples that would be required and the detailed modeling of any calorimeter defects. For these reasons a data-driven technique, the *jet smearing method*, is used. It addresses event topologies with large $E_{\text{T}}^{\text{miss}}$ originating from jet mis-reconstruction. This method creates large sample of well measured low $E_{\text{T}}^{\text{miss}}$ jets called *seeds* which are *smeared* using a function that quantify the p_{T} fluctuation of a measured reconstructed jet (the *response function*) to create

E_T^{miss} in the event. This process is reiterated many times to create a *pseudo-data* sample that is used in the SR analysis selection to estimate the distribution of variables defining the control and signal regions. The procedure is described in more details in Ref. [93]. The multi-jet control region is defined by inverting the $\Delta\phi_{\min}(E_T^{\text{miss}}, \text{jet})$ and applying the inclusive and exclusive SR E_T^{miss} cuts. For the EM1 and IM1 the multi-jet background constitutes the 0.5% of the total BG and it is negligible in other signal regions.

7.6.2 The Non-Collision Background

Non collision background is a term used to refer to BG processes coming from cosmic particles, beam induced muons resulting from proton-gas inelastic interaction or beam halo protons intercepting the LHC collimators and detector noise. The characteristic signature of NCB is that of a jet recoiling against invisible energy thus resembling the monojet final state signature. The jet quality selection criteria mentioned in Section 7.2 manage to reduce the rate of jet coming from cosmic muons to a negligible amount compared to the rate of data in the SR thus the main source of NCB is Beam Induced Background (BIB). In order to estimate the BIB contribution the two-sided no-time method [94] was used. This method tries to match a calorimeter energy cluster with a muon segment in both A and C side of the detector. This topology corresponds to a particle moving parallel to the beam line but several meters away from the beam axis and thus presumably arising from beam background. An estimate of the number of NCB events in the SR is obtained by correcting the number of events tagged as BIB inside the signal region for the efficiency of the method. The efficiency of the tagger is estimated in a sample of events in the SR failing the jet tight cleaning criteria that is dominated by BIB jets.

The NCB contribution in the IM1 results in 112 ± 23 events and only 19 ± 9 events in the EM3, this constitutes about 0.5% of the total background in these regions. For $E_T^{\text{miss}} > 500$ GeV there is no NCB contribution in the signal region.

7.7 Background Systematic Uncertainties

Systematic uncertainties can originate from the theory, caused for example by our limited knowledge of the parton distribution function, or from experimental sources such as the absolute jet energy scale and resolution or production cross section for various processes determination. The systematic uncertainties are treated as Gaussian nuisance parameters in the global likelihood fit described in Section 7.5.4.

7.7.1 Theoretical Uncertainties

The normalization factor for the $Z(\rightarrow \nu\bar{\nu}) + \text{jets}$ is constrained primarily from the single muon control region. Following the run 1 analysis strategy [9], a conservative 3% uncertainty is used. It was obtained by taking the maximum variation of the number of $Z(\rightarrow \nu\bar{\nu}) + \text{jets}$ events in bins of the leading jet

p_T and E_T^{miss} predicted by SHERPA [92] and ALPGEN [95], with two different parton showering algorithms. This 3% uncertainty was therefore a simple way to estimate generator uncertainty and parton shower uncertainty.

In addition an electroweak correction that takes into account the theoretical differences between the W and Z bosons p_T as described in Ref. [96] is added in quadrature to the flat 3% introduced earlier leading to a $\pm 3.5\%$ and $\pm 6\%$ uncertainty in the IM1 and IM7 region respectively.

Top quark production processes get their uncertainties from the $t\bar{t}$ and single-top production cross section, levels of initial and final state radiation, the model used to generate the parton shower. This introduces a $\pm 2.7\%$ and $\pm 3.3\%$ variation in the top background estimation in the IM1 and IM7 signal region respectively.

Uncertainties coming from *diboson* processes are estimated using different MC generators and account for an uncertainty between $\pm 0.05\%$ and ± 0.4 on the number of events in the signal region.

7.7.2 Experimental Uncertainties

Experimental systematic uncertainties can be divided in the following broad categories:

- **Jet/ E_T^{miss} :** This category includes uncertainties on the E_T^{miss} and jet energy scale and resolution, jet quality, pile-up estimation and p_T measurement.
- **Leptons:** This category includes uncertainties related to the identification, simulation, reconstruction efficiencies and energy and momentum resolution of electrons and muons.
- **Luminosity:** Is the uncertainty on the integrated luminosity.

A $\pm 5\%$ uncertainty is assigned to the *integrated luminosity*. Since the efficiency plateau of the HLT_xe70 is reached below 250 GeV E_T^{miss} values, no uncertainty is assigned to the trigger. Uncertainties in the jet and E_T^{miss} *energy scale* and *resolution* contribute to the total background with a variation of the number of predicted events of $\pm 0.5\%$ and $\pm 1.6\%$ in the IM1 and IM7 respectively. The jet quality requirement, pile-up estimation and correction to the measured E_T^{miss} and jet p_T results in an uncertainty in the estimation of the total background of $\pm 0.2\%$ in the IM1 and $\pm 0.9\%$ in the IM7 control region. Lepton *identification* simulation and *reconstruction efficiency* lead to a $\pm 0.1\%$ uncertainty in the IM1 and IM7 regions while energy and momentum *resolution* results in a $\pm 1.4\%$ variation in the IM1 selection and a $\pm 2.6\%$ uncertainty in the IM7 signal region. A flat $\pm 100\%$ uncertainty is assigned to *multi-jet* and *NCB* that translates to a $\pm 0.2\%$ variation in the total background in the IM1 signal region.

7.8 Signal Systematic Uncertainties

The theoretical uncertainties associated to the SUSY compressed signal models arise from the uncertainties on the parton distribution function, the renormalization and factorization scales, the tuning of the parton shower in the Monte Carlo generator and the initial and final state radiation modeling and uncertainties associated to the parton shower matching. The estimation of the theoretical uncertainties is presented in the following sections.

7.8.1 Parton Distribution Function Uncertainties

For precision measurement of the physics processes at the LHC, an accurate description of the PDF is important. The parton distribution functions are non-perturbative and extracted from global fits to hard-scatter data. Two kind of PDFs uncertainties are considered:

- Intra-PDF uncertainty: This is the uncertainty within a given specific PDF set.
- Inter-PDF uncertainty: This is the uncertainty when a certain PDF is replaced with another.

Three different collaborations provide equally precise parton distribution functions making it arbitrary to use one over the others to generate the signal MC samples. For this reason the recommendation of the PDF4LHC [97] working group for a correct estimation of the PDF uncertainty is to combine the uncertainties from all three groups and thus the CT10 [98], the NNPDF [99] and the MMHT [100] PDF sets are used.

The MC signal samples for the SUSY compressed models are generated using the NNPDF2.3 set. In principle in order to evaluate the PDF uncertainties for CT10 and MMHT it would be necessary to re-generate the signal samples using these PDF sets. It is common practice to instead re-weight each event in the original sample to the distributions of the incoming parton momenta x_1 , x_2 it would have had if generated with an alternative PDF. This is done using the LHAPDF [101] tool.

The CT10 and the MMHT sets provide error sets to estimate the intra-PDF uncertainty. The idea is that each PDF has a number of uncorrelated error parameters that can be varied independently by $\pm 1\sigma$ and a new, varied, PDF can be calculated. This procedure is then repeated for each error parameter in the PDF set resulting in a set of PDFs. This procedure is also done using the LHAPDF tool. The Hessian [102] method is then used to estimate the intra-PDF uncertainty on the sets of PDFs.

For CT10 the symmetric Hessian method with 52 error sets is used and the uncertainty is given by:

$$\delta X = \sqrt{\sum_{k=1}^{n_{\text{err}}} (X^{(k)} - X^{(0)})^2} \quad (7.8.1)$$

where $X^{(0)}$ is the number of events using the best PDF fit, $X^{(k)}$ is the number of events using the PDF calculated varying the k -th error parameter and n_{err} is the provided number of error sets.

For MMHT the asymmetric Hessian method with 50 error sets is used and the uncertainty is given by:

$$\delta x^{\text{up}} = \sqrt{\sum_{k=1}^{n_{\text{err}}} \max(0, X_{2k} - X_0, X_{2k-1} - X_0)^2} \quad (7.8.2)$$

$$\delta x^{\text{dw}} = \sqrt{\sum_{k=1}^{n_{\text{err}}} \max(0, X_0 - X_{2k}, X_0 - X_{2k-1})^2} \quad (7.8.3)$$

where X_0 is the number of events obtained with the best PDF fit, X_{2k} (X_{2k-1}) is the number of events calculated with the $2k$ -th ($2k-1$ -th) parameter variation and n_{err} is the number of error sets.

The NNPDF group provides an ensemble of 100 different and equally good PDFs. The one used to generate the signal samples (the *nominal* set) is then obtained by calculating the mean of the 100 error PDFs while the associated intra-PDF systematic uncertainty is given by their standard deviation.

The PDF systematic uncertainty is decomposed in:

- An uncertainty $\Delta\sigma$ on the overall normalization or cross section, given by the variation of the total number of signal events.
- An uncertainty ΔA on the signal acceptance, which is defined by the variation of the ratio of number of signal events determined at the Monte Carlo truth level in a particular bin of the signal region divided by the total number of signal events.

This decomposition is done since the ATLAS standard is to depict the uncertainty on the overall normalization as an uncertainty on the expected limits, while the uncertainty on acceptance is to be incorporated in the uncertainty on the observed limits. Figure 8.9 presents an example of the results in terms of weighted signal yields for the acceptance as a function of the PDF family and error PDF set. It shows the inter- and intra-PDF uncertainties of the three families. The final PDF uncertainty on the number of signal events is the envelope that contains the error bands of the three families.

The PDF uncertainties affect both the overall SUSY compressed signal normalization or cross section and the signal acceptance (the probability for a signal events to have kinematic characteristics that fits the signal region definition or of a particular bin of the signal region). These uncertainties are collected in Table 7.3. The two main features are that the systematic uncertainty grows with higher $E_{\text{T}}^{\text{miss}}$ bin and with higher squark masses. The increase of the uncertainty on the cross section with the mass of the squark is due to the higher momentum fraction of the partons required for its production. At these energy regimes the PDFs are less well known. A similar behavior is seen for the acceptance at higher $E_{\text{T}}^{\text{miss}}$ values which probes less precisely known regions of the PDFs. At low squark masses and in the low $E_{\text{T}}^{\text{miss}}$ bin the systematic uncertainty increases due to low Monte Carlo statistics.

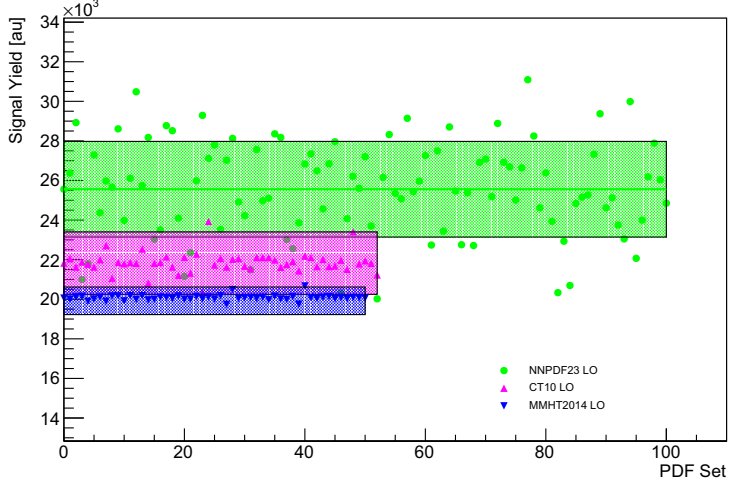


Figure 7.9: Variation in SUSY compressed weighted event yields for a model with $m_{\tilde{q}} = 450$ GeV and $m_{\tilde{\chi}_1^0} = 445$ GeV in the first E_T^{miss} bin of the signal region ($250 < E_T^{\text{miss}} < 300$ GeV) for three PDF sets NNPDF, CT10 and MMHT. The x -axis represents the different systematic uncertainties associated to each PDF set.

PDF Uncertainty							
$m_{\tilde{q}}, m_{\tilde{\chi}_1^0}$ [GeV]	400, 395	450, 445	500, 495	550, 545	600, 595	650, 645	700, 695
$\Delta\sigma$	19	20	21	22	24	25	26
ΔA ($250 < E_T^{\text{miss}} < 300$)	25	26	7	6	7	5	7
ΔA ($300 < E_T^{\text{miss}} < 350$)	5	7	5	4	4	4	4
ΔA ($350 < E_T^{\text{miss}} < 400$)	4	4	5	9	3	6	2
ΔA ($400 < E_T^{\text{miss}} < 500$)	6	2	6	5	2	7	9
ΔA ($500 < E_T^{\text{miss}} < 600$)	7	9	9	10	10	8	10
ΔA ($600 < E_T^{\text{miss}} < 700$)	8	10	6	14	13	10	9
ΔA ($700 > E_T^{\text{miss}}$)	10	9	16	15	17	16	14

Table 7.3: PDF systematic uncertainties in % on the SUSY compressed models. The uncertainty is the envelop that contains the signal yields from the three PDF families, and their error bands. The first row indicates the systematic uncertainty on the overall normalisation. The following rows show uncertainty on the acceptance in the signal region E_T^{miss} bins.

7.8.2 Scale Uncertainties

The cross section can be separated in a short distance partonic component coming from the hard scatter process and long distance terms describing the original hadron [103].

The *factorization scale* is the scale that separates the long distance effects from the short distance ones [104]. To absorb all very-short time physics and ultraviolet divergences renormalization is applied and the coupling constant becomes dependent on the scale at which it is computed. The *renormalization scale* is the scale at which the running coupling constant is evaluated. In QCD the choice of renormalization scale is ambiguous and introduces a non-physical parameter. The uncertainty due to the choice of these two scales affects the cross section and is derived in Ref. [105] by varying them simultaneously by a

factor 2 and $1/2$. The uncertainty is found to be 13% when taking an envelope of the observed variations. For the monojet analysis this is a conservative approach due to the higher value of squark masses investigated by the authors of Ref. [105]. This uncertainty is added in quadrature to the cross section uncertainty on the PDFs collected in Table 7.3.

7.8.3 Matching, Tune, Initial and Final State Radiation Uncertainties

Minimum Bias (MB) events are inelastic pp collisions characterized by low momentum exchange, typically below approximately 2 GeV. They make up the vast majority of the LHC collisions. Proper Monte Carlo modeling of MB processes is therefore important for pile-up prediction of event yields in signal and control regions. Perturbative QCD is commonly used to describe parton interactions when possible and non-perturbative QCD for low p_T processes. The free parameters of these models (*tune parameters*) are tuned in the MC generators in order to replicate experimental MB data. The variation of these parameters may lead to different number of generated events in the MC samples or different kinematic characteristics of the simulated events and thus to an uncertainty on the signal and background predictions. Following the ATLAS recommendation given in Ref. [106] five tune parameters have been varied in order to account for uncertainties from:

- Underlying event effects.
- Jet structure effects.
- Aspects of the MC generation that provide extra jet production.

The presence of large numbers of high p_T jets in the final state could hide signal for new physics that involves the decay of heavy colored particles. In order to control this kind of multijet events the theory calculations need to describe in a precise way the matrix element used to describe the hard process and the subsequent parton shower into jets of hadrons in the calorimeter. A particular N-jet event can arise from soft-radiation evolution of a N-parton final state (matrix element jets) or from an (N-1)-parton configuration where the extra jet appears from hard emission during the shower evolution (parton shower jet). A *matching scheme* is applied on an event-by-event basis in order to

avoid double counting and decide which jets should be taken from the matrix element and which ones should be taken from parton shower [107]. The SUSY compressed signal samples are produced with MADGRAPH [108], thus an additional matching scale systematic uncertainty is considered.

Practically the matching, tune, initial and final state radiation uncertainties are estimated by first generating 16 alternative MC samples for each squark mass with different settings and evaluating the relative variation in number of events in the signal regions as observed at the Monte Carlo truth level. For a given squark and neutralino mass, the different relative variations are added in quadrature. Finally the envelope of the systematics obtained over all squark

and neutralino mass is derived in order to obtain a single systematic uncertainty for all SUSY compressed mass points. The results are presented in Table 7.4.

Initial and Final State Radiation							
E_T^{miss} [GeV]	250-300	300-350	350-400	400-500	500-600	600-700	> 700
ΔA	11	18	11	8	7	9	8

Table 7.4: Theoretical uncertainty in % on the SUSY compressed spectra signal region acceptance as function of the E_T^{miss} bin in the signal region, from tune, matching, initial and final state radiation systematic uncertainties. The final value is a common envelope valid for all the SUSY compressed models.

7.9 Results and Interpretations

The number of events observed in data in the signal regions was compared with the background only fit. The results of the comparison between the observed number of events in data and the background calculations including the background only fit is shown in Figure 7.10. No significant excess of data over background is observed. The results of the global fit procedure described in Section 7.5.4 are used to set model independent exclusion limits (see Section 7.9.3) and interpreted in terms of squark pair production with $\tilde{q} \rightarrow q + \tilde{\chi}_1^0$ ($q = u, d, c, s$) (see Section 7.9.4).

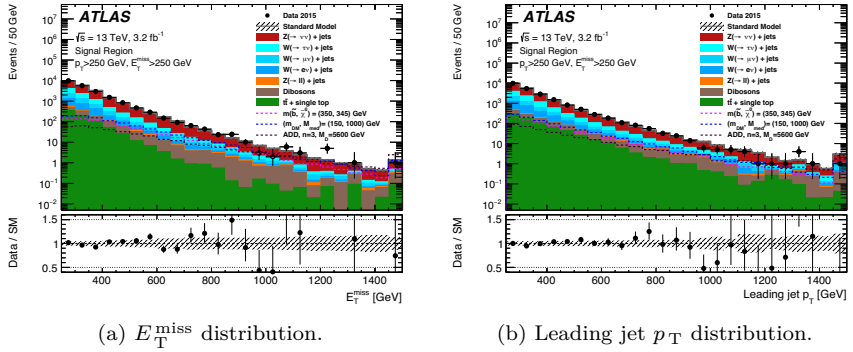


Figure 7.10: Distribution of the E_T^{miss} and the leading jet p_T for IM1 signal region compared with the background estimates from the background only fit in the control regions. The distributions of different signal models are superimposed for comparison. The contribution from the multi-jet and NCB background is negligible and not reported in the plot. In the ratio window the error bars include experimental and systematic uncertainties.

7.9.1 The Statistical Procedure

The p -value is the quantity used in the LHC frequentist approach [109] to quantify the level of agreement between the observed data and some hypothesis H . Assuming that H is true, the p -value is defined as the probability to observe

data of equal or greater incompatibility with H . In order to *discover* a new physics signal, the *null hypothesis*, H_0 , is defined to include only known SM processes (background) and tested against an *alternate hypothesis*, H_1 , that describes both the looked after signal and the background. If the signal cannot be discovered the roles of the two hypotheses are inverted (the *signal plus background* hypothesis becomes the null hypothesis and the *background only* hypothesis becomes the alternate one) and *exclusion limits* are set instead. The conventionally accepted p-value to claim a discovery is set at 2.87×10^{-7} while to exclude a signal hypothesis a threshold p-value of 0.05 is used.

In an experiment that measures a variable x for each event in the signal sample and constructs an histogram $\mathbf{n} = (n_1, \dots, n_N)$ out of these measurements, the expectation value of n_i can be written as:

$$E[n_i] = \mu s_i + b_i \quad (7.9.1)$$

where s_i and b_i are the mean number of signal and background entries in the i -th bin and μ is the *signal strength* parameter of the signal process with $\mu = 0$ indicating the background only hypothesis and $\mu = 1$ the nominal signal one. The parameters of the physical model, unknown properties of the response of the detector as well as other parameters that are not of interest for the physics search are called *nuisance parameters* and indicated with $\boldsymbol{\theta}$. These can be evaluated constructing histograms of some kinematic variable of interest in a control region where mostly background events are expected. Indicating as $\mathbf{m} = (m_1, \dots, m_M)$ one of such histograms, the expectation value of the i -th bin, m_i , is given by:

$$E[m_i] = u_i(\boldsymbol{\theta}) \quad (7.9.2)$$

where u_i are calculable quantities that depend on the nuisance parameters $\boldsymbol{\theta}$. With these definitions, the *likelihood* function can be written as the product of Poisson probabilities for the different bins:

$$L(\mu, \boldsymbol{\theta}) = \prod_{j=1}^N \frac{(\mu s_j + b_j)^{n_j}}{n_j!} e^{-(\mu s_j + b_j)} \prod_{k=1}^M \frac{u_k^{m_k}}{m_k!} e^{-u_k} \quad (7.9.3)$$

The p-value is determined at LHC using the *test statistics* [109]:

$$t_\mu = -2 \ln \lambda(\mu) \quad (7.9.4)$$

where $\lambda(\mu)$ is the *profile likelihood ratio* defined as:

$$\lambda(\mu) = \frac{L(\mu, \hat{\boldsymbol{\theta}})}{L(\hat{\mu}, \hat{\boldsymbol{\theta}})}. \quad (7.9.5)$$

Here $\hat{\mu}$ and $\hat{\boldsymbol{\theta}}$ are the maximum likelihood estimators, $L(\hat{\mu}, \hat{\boldsymbol{\theta}})$ is the maximized likelihood function and $\hat{\boldsymbol{\theta}}$ is the value of the parameters $\boldsymbol{\theta}$ that maximizes the likelihood for a fixed value of the signal strength μ . It can be seen that $0 \leq \lambda(\mu) \leq 1$ with a value close to one corresponding to a good agreement between the data and μ , this means that high values of the test statistic in Eq. (7.9.4) correspond to bigger incompatibility between data and the hypothesized μ .

For the discovery of new physics signals, the test statistics in Eq. (7.9.4) is modified to:

$$q_0 = \begin{cases} -2 \ln \lambda(0) & \hat{\mu} \geq 0 \\ 0 & \hat{\mu} < 0 \end{cases} \quad (7.9.6)$$

where in the model it is assumed that $\mu \geq 0$ and the likelihood ratio in Eq. (7.9.5) is evaluated for $\mu = 0$ (the background only hypothesis). The p-value is given in this case by:

$$p_0 = \int_{q_0, \text{obs}}^{\infty} f(q_0|0) dq_0 \quad (7.9.7)$$

where $f(q_0|0)$ is the PDF of q_0 under the background only hypothesis. Excluding the $\mu = 0$ hypothesis corresponds to the discovery of a new signal. In the large sample limit it is possible to use the result from Wald [110]:

$$-2 \ln \lambda(\mu) = \frac{(\mu - \hat{\mu})^2}{\sigma^2} + \mathcal{O}(1/\sqrt{N}), \quad (7.9.8)$$

where N is the sample size and $\hat{\mu}$ follows a Gaussian distribution with mean μ' and standard deviation σ , to approximate the test statistics in Eq. (7.9.6) to:

$$q_0 = \begin{cases} \hat{\mu}^2/\sigma & \hat{\mu} \geq 0 \\ 0 & \hat{\mu} < 0 \end{cases} \quad (7.9.9)$$

where $\hat{\mu}$ is distributed with a Gaussian shape with mean μ' and standard deviation σ . From this result it is possible to show [109] that the p-value for the background only hypothesis ($\mu = 0$) is given by:

$$p_0 = 1 - F(q_0|0) \quad (7.9.10)$$

where $F(q_0|0) = \Phi(\sqrt{q_0} - \mu'/\sigma)$ is the cumulative distribution function of $f(q_0|0)$ and Φ is the cumulative distribution function of a standard Gaussian with zero mean and unit standard deviation.

In order to exclude a signal strength μ upward fluctuations of the data are not regarded as less compatible with μ than the obtained data and for this reason the test statistic is set to zero for $\hat{\mu} > \mu$ and the base test statistic given in Eq. (7.9.4) is modified to:

$$q_\mu = \begin{cases} -2 \ln \lambda(\mu) & \hat{\mu} \leq \mu \\ 0 & \hat{\mu} > \mu \end{cases} \quad (7.9.11)$$

where $\lambda(\mu)$ is the likelihood ratio of Eq. (7.9.5). The p-value can be calculated with:

$$p_\mu = \int_{q_\mu, \text{obs}}^{\infty} f(q_\mu|\mu) dq_\mu \quad (7.9.12)$$

where $f(q_\mu|\mu)$ is the PDF of q_μ under the hypothesis μ . Using the result from Wald [110], (7.9.11) can be approximated to a chi-square distribution with one degree of freedom [109]:

$$q_\mu = \begin{cases} \frac{(\mu - \hat{\mu})^2}{\sigma^2} & \hat{\mu} < \mu \\ 0 & \hat{\mu} > \mu \end{cases} \quad (7.9.13)$$

where $\hat{\mu}$ follows a normal distribution centered around μ' and standard deviation σ . In this approximation the p-value for the μ hypothesis is [109]:

$$p_\mu = 1 - F(q_\mu|\mu) = 1 - \Phi(\sqrt{q_\mu}) \quad (7.9.14)$$

where $F(q_\mu|\mu)$ is the cumulative distribution function of $f(q_\mu|\mu)$ and Φ is the cumulative distribution function of a Gaussian centered around zero and with unit variance. If the calculated p-value p_μ is below some threshold α :

$$p_\mu < \alpha \quad (7.9.15)$$

then the hypothesis μ is said to be excluded with a $1 - \alpha$ Confidence Level (CL). As mentioned earlier, as a convention, for exclusion purposes, the value of α is taken to be 0.05 and the excluded μ is thus quoted with a 95% CL. The upper limit on μ is the largest μ that satisfies $p_\mu \leq \alpha$. It can be calculated analytically by setting $p_\mu = \alpha$ and solving for μ obtaining [109]:

$$\mu_{\text{up}} = \hat{\mu} + \sigma\Phi^{-1}(1 - \alpha). \quad (7.9.16)$$

Since σ depends on μ [109], it is common to calculate the upper limit numerically as the value of μ for which $p_\mu = \alpha$.

The CL_S Method

In the search for new physics signal it may happen that the experiment does not have the sensitivity to exclude a particular model. This can happen for example when testing the hypothesis of some SUSY particles with a mass greater than what can be produced at LHC. In this case the frequentist approach described in Section 7.9.1 for setting upper exclusion limits when $p_\mu < \alpha$ will reject the model with a probability of α or greater [111]. In general this kind of situation happens whenever the signal plus background distribution is very similar to the background only one.

The CL_S method [112] can be used in such situations. This more conservative method prevents the rejection of models where the sensitivity is low reverting to the frequentist procedure introduced in Section 7.9.1 for models with higher sensitivity. The method defines the quantity:

$$\text{CL}_S = \frac{p_{s+b}}{1 - p_b} \quad (7.9.17)$$

where p_{s+b} and p_b correspond to the p-value for the signal plus background ($\mu = 0$) and background only ($\mu = 0$) hypothesis respectively. The CL_S value is then used in place of the p-value and the inequality $\text{CL}_S \leq \alpha$ is evaluated instead. Figure 7.11a allows to understand qualitatively how the method works, in the figure the test statistics Q is similar both for the signal plus background and for the background only hypotheses. If Q_{obs} is such that $p_{s+b} < \alpha$, in the frequentist approach described in Section 7.9.1 the signal plus background hypothesis would be rejected while using the CL_S method it is possible to see that, since p_b is big and thus $1 - p_b$ is small, the ratio in Eq. (7.9.17) will be greater than p_{s+b} and the model is not rejected. In the case of high sensitivity,

the signal plus background and the background only hypothesis distributions are well separated as depicted in Figure 7.11b. In this case if Q_{obs} is such that $p_{s+b} < \alpha$ then p_b is small and, being the denominator in Eq. (7.9.17) close to one, the CL_S method becomes similar to the usual frequentist approach of Section 7.9.1 where the p-value p_{s+b} is tested [112].

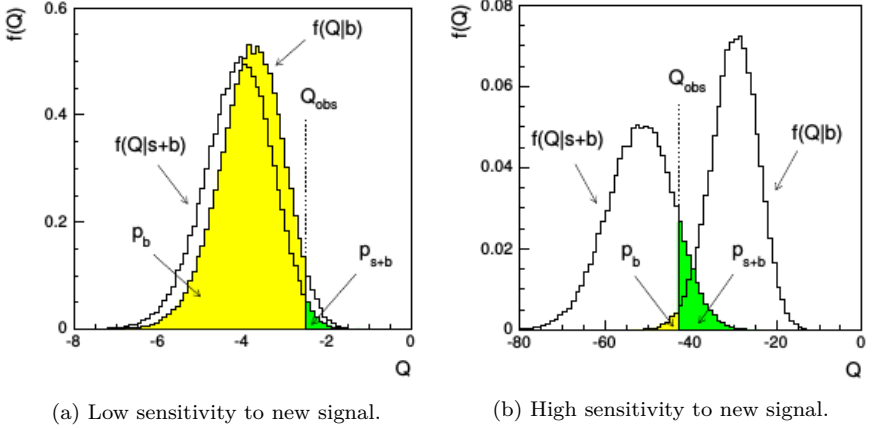


Figure 7.11: Graphical illustration of the CL_S method. In Figure 7.11a the low sensitivity would lead to the rejection of the p_{s+b} hypothesis if the CL_S method is not used. In the case of two well separated distributions like in Figure 7.11b the use of the CL_S method is equivalent to the usual frequentist approach [111].

7.9.2 Expected Limits

Since in the compressed SUSY squark-neutralino model the E_T^{miss} and jets recoil against an Initial State Radiation (ISR) jet, the average jet momenta are lower than the missing energy, thus requiring an asymmetric cut on the leading jet momentum and E_T^{miss} to capture this signal. Prior to deciding the signal region selections presented in Section 7.2 a signal region (AM1) with $E_T^{\text{miss}} > 700$ GeV and leading jet $p_T > 300$ GeV was defined and compared with the IM1 signal region. Since the final integrated luminosity cannot be known in advance, the studies on AM1 and IM1 were carried out on a foreseen integrated luminosity of $L = 3.3 \text{ fb}^{-1}$. The results for these two signal regions is presented in Figure 7.12. On these figure a model is excludable if the signal strength μ , depicted by the color scale in the picture, is lower than one. The sensitivity improves for the signal region AM1 to reflect the asymmetric kinematics.

In the AM1 signal region, models with mass gaps $\Delta m = 5$ GeV between the squark mass and the neutralino mass, can be excluded up to squark masses of 650 GeV. At a larger mass gap of $\Delta m = 25$ GeV, squark masses up to 580 GeV are excludable.

The performance of the shape fit method was also tested on the SUSY compressed squark-neutralino model and the result is also shown in Figure 7.12.

The shape fit method, thanks to its flexibility, is capable of adapting to different signal models and thus improves the expected limits of the SUSY squark-neutralino model especially for the largest mass gap of $\Delta m = 25$ GeV between the squark and the neutralino masses.

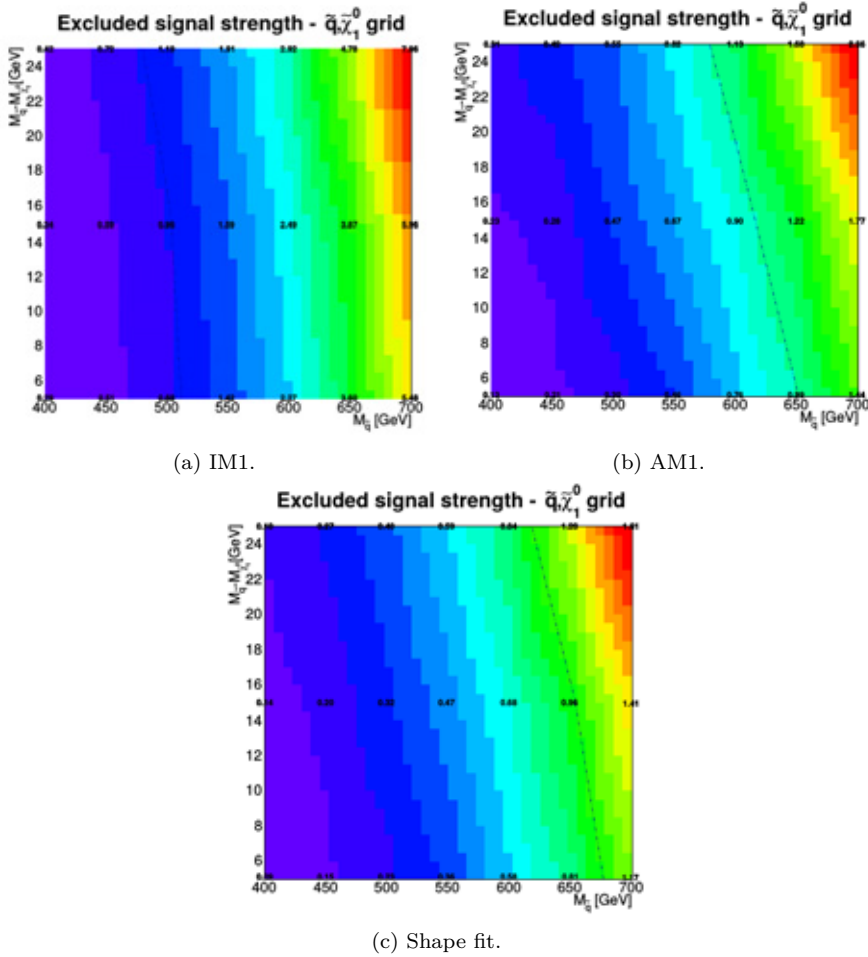


Figure 7.12: Sensitivity to the compressed SUSY models with an integrated luminosity of $L = 3.3 \text{ fb}^{-1}$ in the plane defined by the squark mass $m_{\tilde{q}}$ and the mass difference between the squark mass and the lightest neutralino mass $\Delta m = m(\tilde{q}) - m(\tilde{\chi}_1^0)$. The color scale reflects the lowest excludable signal strength for a particular mass point. The values indicated on the map indicate the actual excludable μ values at a some specific SUSY models. The dashed line shows the limit between excludable and not excludable models. The theoretical uncertainties in the signal are not included.

7.9.3 Model Independent Limits

The *visible cross section* represents the number of BSM events in the signal region per unit of luminosity. It is related to the signal production cross section by:

$$\sigma_{\text{vis}} = \sigma \times A \times \epsilon \quad (7.9.18)$$

where σ is the production cross section for the signal process in pp collisions, A is the selection acceptance and ϵ is the selection efficiency. The CL_S modified frequentist approach introduced in Section 7.9.1 is used to set 95% CL exclusion limit on the *visible cross section* for models with a monojet final state experimental signature. The results are reported in Table 7.5, values of σ_{vis} above $\langle\sigma\rangle_{\text{obs}}^{95}$ are excluded at 95% CL. These results can be understood as follows: BSM models with more than 19 events per inverse femtobarn of integrated luminosity in the IM7 signal region are excluded by the ATLAS data.

Result Table			
Signal Region	$\langle\sigma\rangle_{\text{obs}}^{95}$ [fb]	S_{obs}^{95}	S_{exp}^{95}
IM1	533	1773	1864^{+829}_{-548}
IM2	308	988	1178^{+541}_{-348}
IM3	196	630	694^{+308}_{-204}
IM4	153	491	401^{+168}_{-113}
IM5	61	196	164^{+63}_{-45}
IM6	23	75	84^{+32}_{-23}
IM7	19	61	48^{+18}_{-13}

Table 7.5: Results on the expected and observed upper limits on the number of events, S_{exp}^{95} and S_{obs}^{95} respectively and on the visible cross section at 95% CL.

7.9.4 SUSY Compressed Spectra Interpretation

The results of model independent limits are interpreted in terms of limits computed for squark pair production with $\tilde{q} \rightarrow q + \tilde{\chi}_1^0$ ($q = u, d, c, s$) in a SUSY compressed scenario. The sensitivity to these models is estimated by performing a global signal plus background fit that includes the CR_{ele} , CR_{wmn} , CR_{zmm} control regions, the signal regions and all the systematic uncertainties. The expected limits are derived with the same procedure but by replacing the data with the background prediction.

Figure 7.13 shows the result limits on the SUSY compressed squark-neutralino model for the run 2 2015 data. This is the result of the shape fit with a luminosity of 3.2 fb^{-1} with all theoretical signal systematic uncertainties included in the observed and expected exclusion contour and using the CR_{ele} , CR_{wmn} , CR_{zmm} control regions. The theoretical uncertainty on the signal cross section is used to derive the uncertainty on the observed limit.

Models with a mass gap between the squark and the neutralino of $\Delta m = 5$ GeV are excluded up to squark masses of 608 GeV. For the larger mass gap of $\Delta m = 25$ GeV, squark masses up to 532 GeV are also excluded. The results of this study are combined with the more general SUSY searches adding sensitivity to the region close to the diagonal (dashed line) in Figure 7.3.

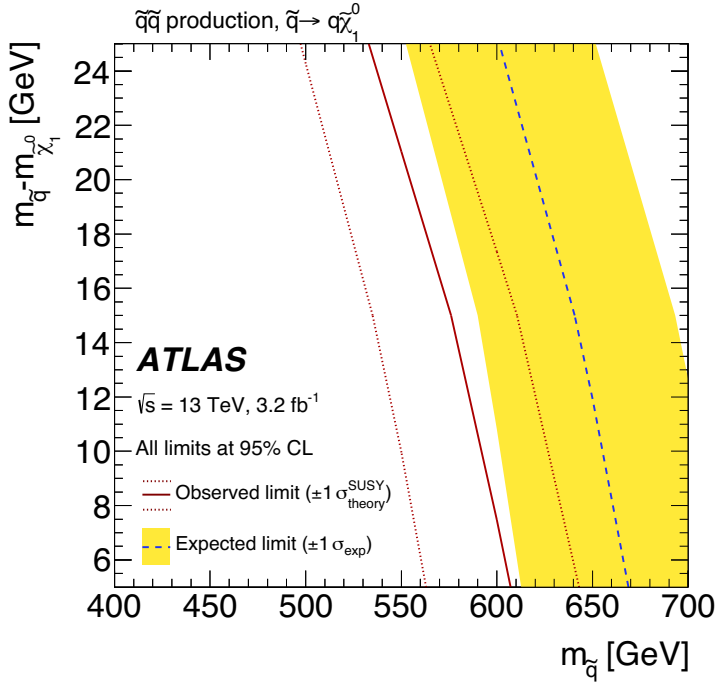


Figure 7.13: Expected and observed limits on the SUSY compressed models using the run 2 2015 data, in the plane defined by the squark mass on the x -axis and the mass difference between the squark and lightest neutralino mass on the y -axis. All experimental and theoretical systematic uncertainties are included.

Chapter 8

Search for Large Extra Dimensions with 2015 and 2016 ATLAS Data

The large gap between the electroweak scale and the Planck scale is solved in the ADD scenario for LED by a lowered Planck mass set to be at the electroweak scale and by introducing n extra spatial dimensions of radius R in order to recover the strength of the gravitational interaction ($1/M_{Pl}$). The search carried out using the dataset collected by the ATLAS detector in 2015 and 2016 and presented in this chapter focuses on testing the existence of two to six extra spatial dimensions.

8.1 Phenomenology of the ADD Model

Gravitons could be produced in proton-proton collisions at the LHC. The interaction Lagrangian is given by [113]:

$$\mathcal{L} = -\frac{1}{\bar{M}_P} G_{\mu\nu}^{(n)} T^{\mu\nu} \quad (8.1.1)$$

where $\bar{M}_P = M_P/\sqrt{8\pi}$ is the reduced Planck mass, $G_{\mu\nu}^{(n)}$ is the graviton field in n dimensions and $T^{\mu\nu}$ is the stress-energy tensor, and thus the graviton production cross section is suppressed by a factor $1/\bar{M}_P^2$. Nevertheless the mass splitting of the Kaluza-Klein modes is given by [113]:

$$\Delta m \approx \frac{1}{R} = M_D \left(\frac{M_D}{\bar{M}_P} \right)^{2/n} \approx \left(\frac{M_D}{\text{TeV}} \right)^{\frac{n+2}{2}} 10^{\frac{12n-31}{n}} \text{ eV} \quad (8.1.2)$$

where n represents the number of extra dimensions which for $M_D = 1 \text{ TeV}$ and $n = 4, 6, 8$ gives $\Delta m = 20 \text{ keV}, 7 \text{ MeV}$ and 0.1 GeV . For $n > 6$ only few KK modes can be produced and the total production cross section becomes negligible [113] while only for a large number of extra dimensions the mass

splitting energy becomes comparable with the experimental energy resolution and the number of KK modes produced is capable of compensating the $1/\bar{M}_P^2$ suppression factor in the production cross section. The graviton lifetime of $\tau_G \approx \bar{M}_P^2/m^3 \approx (\text{TeV}/m)^3 \approx 10^3$ seconds [113] makes it stable over ATLAS detection time. The $1/\bar{M}_P^2$ suppression factor can be interpreted as the low probability that a graviton propagating in the extra dimensions crosses the usual three spatial dimensions [113]. The Kaluza-Klein graviton only interacts weakly through the gravitational interaction with the Standard Model particles and thus behaves like a non interacting, massive and stable particle which, escaping detection, leads to an energy imbalance in the detector which can be investigated using the monojet signature. Figure 8.1 shows the three main graviton production modes at LHC. These production modes are implemented at LO in the Monte Carlo generator PYTHIA8 [114] with the ATLAS2014 (A14) tune [115].

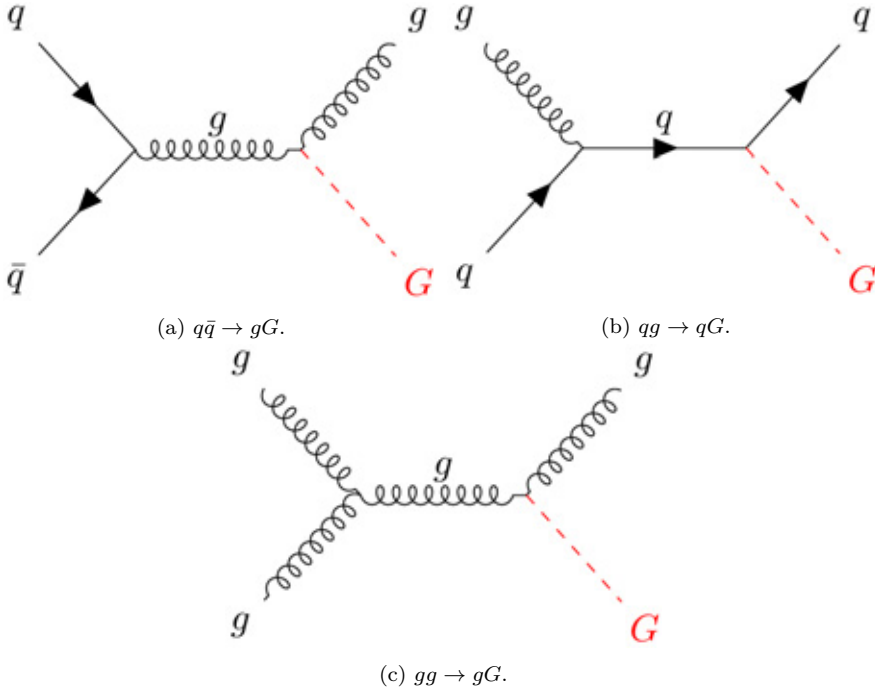


Figure 8.1: The main leading order diagrams for graviton plus quark and graviton plus gluon production in the ADD model at LHC.

As mentioned earlier, collider searches do not have sensitivity to ADD scenarios with more than six extra dimensions, for this reason in this analysis only models with six or less extra dimensions are considered. Moreover, since from Eq. (3.3.14), using $n = 1$ leads to $R \approx 10^{11}$ m which is empirically excluded [113], only models with two extra dimensions or more are considered. Finally in the case with two and three extra dimensions, values of M_D lower than 110 TeV and 5 TeV respectively are excluded by cosmological arguments [116].

The signal samples used to calculate the exclusion limits on M_D are generated with PYTHIA8. The following relation between the production cross section and M_D [113] was verified with PYTHIA samples and is exploited in order to generate only a single Monte Carlo sample per number of dimensions:

$$\sigma_{\text{nom}} = \frac{C_n}{(M_D^{\text{nom}})^{n+2}} \quad (8.1.3)$$

where σ_{nom} and M_D^{nom} are the nominal values of the cross section and the Planck scale used in PYTHIA8 to generate the MC signal sample while C_n is a constant term for a given center of mass energy and number of extra dimensions that is determined from Eq. (8.1.3) as:

$$C_n = \sigma_{\text{nom}} (M_D^{\text{nom}})^{n+2}. \quad (8.1.4)$$

Thus a single Monte Carlo sample is generated for each n and the ADD graviton cross section can be deduced for any M_D from that relation. The samples generated along with their values of σ_{nom} and M_D^{nom} used in this analysis are reported in Table 8.1. Once the value of C_n is determined, Eq. (8.1.3) can be regarded as a function $\sigma(M_D)$ and the cross section is set by the Planck scale. In particular, for some excluded value of M_D^{excl} :

$$\frac{\sigma_{\text{nom}}}{\sigma_{\text{excl}}} = \left(\frac{M_D^{\text{excl}}}{M_D^{\text{nom}}} \right)^{n+2}. \quad (8.1.5)$$

From the definition of excluded signal strength $\mu_{\text{excl}} = \sigma_{\text{excl}}/\sigma_{\text{nom}}$ where σ_{excl} is the production cross section of some BSM model, we have that:

$$\sigma_{\text{excl}} = \mu_{\text{excl}} \sigma_{\text{nom}} \quad (8.1.6)$$

from which:

$$\mu_{\text{excl}} = \left(\frac{M_D^{\text{nom}}}{M_D^{\text{excl}}} \right)^{n+2} \quad (8.1.7)$$

and finally:

$$M_D^{\text{excl}} = \frac{M_D^{\text{nom}}}{\mu_{\text{excl}}^{1/(n+2)}}. \quad (8.1.8)$$

With these relations upper limits on the signal strength are translated into upper limits on the fundamental Planck scale M_D .

In the previous version of the analysis it was shown that low missing energy signal regions have no experimental sensitivity to ADD models. Therefore the ADD samples used in this analysis are generated with with a 150 GeV phase space cut at the truth level on the transverse momentum of the parton and the graviton ($p_T > 150$ GeV). The cross section in Table 8.1 include the efficiency of the truth level cut on the leading jet.

Effective Planck Scale and Nominal Cross Section		
Model	$M_{\text{D}}^{\text{nom}}$ [GeV]	σ [pb]
ADD $n = 2$	5300	1.55
ADD $n = 3$	4100	1.82
ADD $n = 4$	3600	2.29
ADD $n = 5$	3200	4.18
ADD $n = 6$	3000	7.22

Table 8.1: Nominal values for the cross section and the Planck scale used to generate the ADD signal samples at 13 TeV. The samples are generated using PYTHIA8 with NNPDF23LO [117] PDF set and the A14 tune. The number of extra dimensions n and the effective Planck scale M_{D} are free parameters of the model. Theoretical systematic on the cross section are discussed in Section 8.6.

8.2 Validity of the Effective Field Theory

As mentioned in Section 3.3.2 in the ADD model, the fundamental scale of the gravitational interaction, M_{D} , is brought close to the electroweak scale. At these energies, due to the large center of mass energy available at LHC, the predictions of the EFT become unreliable. For this reason two different sets of limits have been calculated, one where the full cross section is used and the other where it is weighted down by a factor $M_{\text{D}}^4/\hat{s}^2$ for events in which $\hat{s} > M_{\text{D}}^2$, where $\hat{s}$ is the center of mass energy of the initial partons involved in Figure 8.1 [118]. The general idea of weighting the events is therefore introduced. It consists in re-evaluating the limit on M_{D} in a conservative way, by prescribing that all events with $\hat{s}$ exceeding the validity limit should only weakly contribute to the limit on M_{D} .

The selection of the missing transverse energy changes the $\hat{s}$ distribution as discussed below. Therefore the validity of the EFT needs to be evaluated in each $E_{\text{T}}^{\text{miss}}$ bin used in the analysis. As an illustration of the problem Figure 8.2 shows the $\hat{s}$ distribution for ADD events surviving a signal region selection with $250 < E_{\text{T}}^{\text{miss}} < 300$ GeV and $700 < E_{\text{T}}^{\text{miss}} < 800$ GeV for the $n = 3$ and 6 extra dimensions models. The straight line indicates the excluded M_{D} value before any weighting of the events. This figure shows that a fraction of the selected events in the signal regions have a value of $\hat{s}$ that exceeds the limit of validity of the effective field theory. Figure 8.2 shows that the bulk of the $\hat{s}$ distribution moves towards higher energy values both in the case when the number of extra dimensions and the $E_{\text{T}}^{\text{miss}}$ are increased. The first dependence is understood considering the growth of the graviton mass with the number of extra dimensions showed in Figure 8.3: the production of gravitons with higher mass requires higher $\hat{s}$ of the initial partons. A similar argument holds when the $E_{\text{T}}^{\text{miss}}$ selection is increased, higher values of missing energy imply the production of higher momentum gravitons. As shown in Figure 8.2, only a small fraction of events is beyond the EFT validity limit. The visible event yields are recomputed in each $E_{\text{T}}^{\text{miss}}$ bin (as defined in Section 8.3) and for each n and, as hinted by Figure 8.2, the effect of the re-weighting is negligible.

Figure 8.4 shows the visible cross section (see Section 7.9.3) as a function

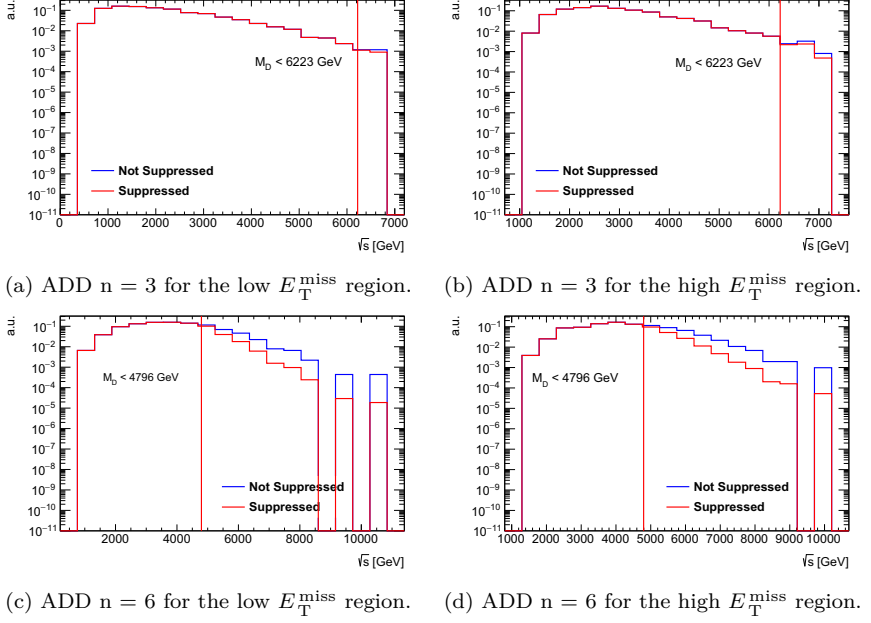


Figure 8.2: Generated $\hat{s}$ distribution with (blue) and without (red) weighting of the events for which $\hat{s} > M_D^2$ for the signal regions where $250 < E_T^{\text{miss}} < 300$ GeV and $700 < E_T^{\text{miss}} < 800$ GeV for the ADD $n = 3$ and 6 models for the 13 TeV PYTHIA8 Monte Carlo samples. A vertical line indicating the value of the excluded value of M_D is also reported in the figure.

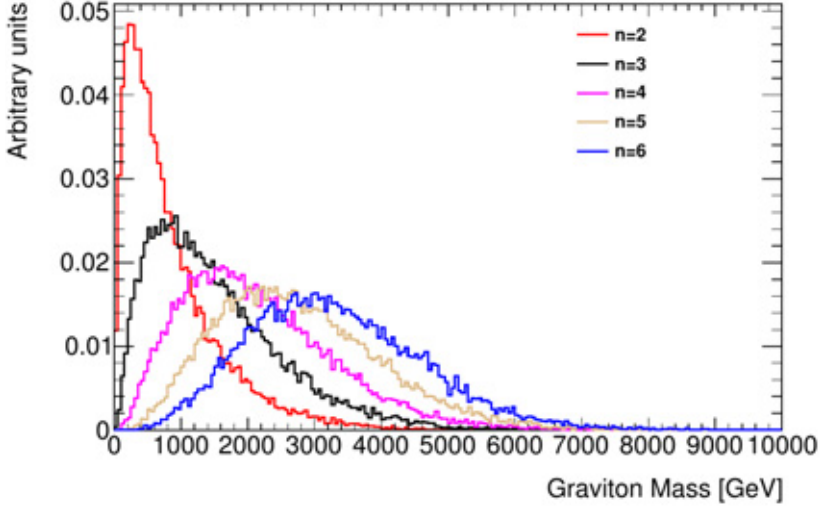


Figure 8.3: Simulated heavy graviton masses for 2 to 6 extra dimensions in ADD models computed with PYTHIA8 using the NNPDF23LO PDF set and the A14 tunes [119].

of the fundamental Planck scale M_D in the $250 < E_T^{\text{miss}} < 300$ GeV and $700 < E_T^{\text{miss}} < 800$ GeV regions for the ADD $n = 3, 6$ models. The solid line is the visible cross section when considering the entire $\hat{s}$ range, while in the dashed

line a $M_D^4/\hat{s}^2$ weighting factor is applied for events where $\hat{s} > M_D^2$. It can be seen that for $M_D = 0$ all events would get weighted down since $\hat{s} > M_D^2 = 0$. The black square is the nominal M_D value of the generated sample listed in Table 8.1. The ratio of the dashed red curve over the blue is used to scale down the exclusion limits in each bin of E_T^{miss} in order to take into account the effect of the validity of the EFT. This effect is however smaller than the uncertainties on the exclusions and is thus neglected. The $\hat{s}$ distribution is different for each E_T^{miss} bin which therefore receives a different suppression. The lowest E_T^{miss} bins are suppressed the least.

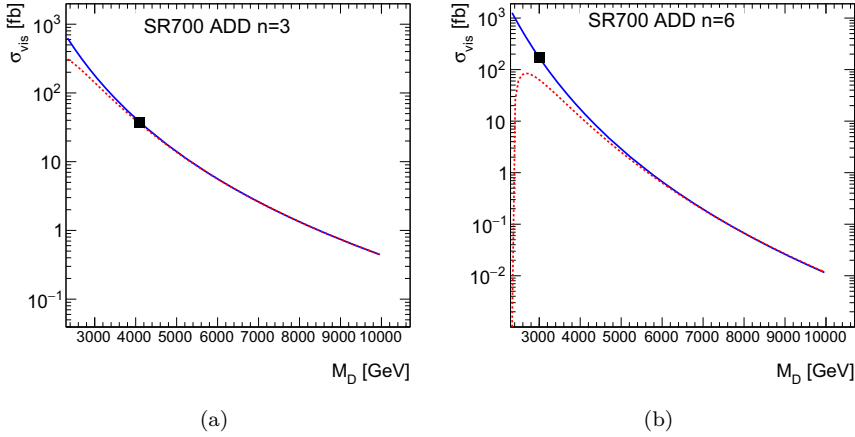


Figure 8.4: Visible cross section as a function of M_D for the signal region where $250 < E_T^{\text{miss}} < 300$ GeV and $700 < E_T^{\text{miss}} < 800$ GeV for the ADD $n = 3$ and 6 models. The solid line is the visible cross section when considering all events, while the dashed curve is built only from events where a $M_D^4/\hat{s}^2$ weighting factor is applied for events where $\hat{s} > M_D^2$. The black square is the nominal M_D value of the generated sample as listed in Table 8.1.

8.3 Event Selection

The search for LEDs in the ADD model is carried out using the pp collision data at $\sqrt{s} = 13$ TeV collected by the ATLAS experiment during 2015 and 2016 corresponding to an integrated luminosity of 36.1 fb^{-1} . The events in the signal region are selected as described in Section 7.2 where in addition to cut A a combination of four different triggers as described in Table 4.1 is used to select the 2016 data. Inclusive (IM1-IM10) and exclusive (EM1-EM9) signal regions are defined in the monojet analysis with increasing E_T^{miss} thresholds from 250 GeV to 1000 GeV, Table 8.2 shows the exact definition of the signal regions. These different E_T^{miss} bins are defined in order to address different BSM signals tested with the monojet signature. For each extra dimension ADD model there is a specific E_T^{miss} bin that has the most sensitivity. Since most bins provide some additional sensitivity they are all combined in the global fit described in Section 7.5.4.

Event Selection Criteria										
HLT_xe70 trigger for 2015 data HLT_xe80_tc_lcw_L1XE50, HLT_xe90_mht_L1XE50, HLT_xe100_mht_L1XE50 and HLT_xe110_mht_L1XE50 triggers for 2016 data ¹ Primary Vertex $E_T^{\text{miss}} > 250$ GeV Leading jet with $p_T > 250$ GeV and $ \eta < 2.4$ At most 4 jets $\Delta\phi(\text{jets}, E_T^{\text{miss}}) > 0.4$ Jet quality requirements No identified muon with $p_T > 10$ GeV or electron of $p_T > 20$ GeV										
Inclusive SRs	IM1	IM2	IM3	IM4	IM5	IM6	IM7	IM8	IM9	IM10
E_T^{miss} [GeV]	> 250	> 300	> 350	> 400	> 500	> 600	> 700	> 800	> 900	> 1000
Exclusive SRs	EM1	EM2	EM3	EM4	EM5	EM6	EM7	EM8	EM9	
E_T^{miss} [GeV]	[250-300]	[300-350]	[350-400]	[400-500]	[500-600]	[600-700]	[700-800]	[800-900]	[900-1000]	

Table 8.2: Definition of the signal regions for the search for ADD gravitons in the 2015 and 2016 data.

8.4 Background Estimation

The sources of background that can mimic the signal are described in Section 7.4. The estimation of the $V + \text{jets}$ and $t\bar{t}$ processes in the signal region is evaluated in a data driven way where observations in regions of the phase space where an enhancement of these processes is expected are used to improve the MC predictions. In the search for SUSY with 3.2 fb^{-1} and described in Chapter 7 the uncertainty on top production was one of the biggest, being approximately 30% for $t\bar{t}$ and single top processes in the signal region for the IM1 selection. In this version of the analysis a dedicated top control region was introduced in order to improve the background estimation. The new control region dedicated to the top background is described in Section 8.4.2. Sections 8.4.1, 8.4.3 and 8.4.4 describe the changes, where applicable, to the remaining control regions that were also used in the 2015 analysis.

The multi-jet and the non-collision backgrounds are estimated with the same technique as the 2015 data analysis and described in Sections 7.6.1 and 7.6.2 respectively. The multi-jet background contribution is largest in the EM1 and IM1 signal regions where it only contributes to about 0.4% and 0.3% of the total background respectively and is negligible in the other signal regions. The NCB contribution to the total background is negligible in all regions.

8.4.1 The Muon Control Region With b -Jets

Top pairs and single top production are generated with the POWHEGBOX [120] generator with CT10 PDFs and the parton shower, fragmentation and underlying event simulated with PYTHIA8 with A14 set of tune parameters. Due to the large Monte Carlo modeling uncertainties on the top background at high E_T^{miss} observed in the analysis of the 2015 data (about 40% for $t\bar{t}$ and single top processes in the signal region for IM6) this process is estimated with a data-driven method in the combined 2015 and 2016 data. To this end the Top control region (CR_{Top}) is introduced to estimate the contribution of $t\bar{t}$ and

single top processes in the signal region. In addition to cuts from A to H as defined in Section 8.3 the CR_{Top} region selects events with:

- Exactly one good muon.
- No baseline electrons.
- The transverse mass defined in (7.5.1) of the $E_{\text{T}}^{\text{miss}} - \mu$ system satisfies $30 < m_{\text{T}} < 100$ GeV compatible with a W boson production.
- At least one of the selected jets in the event is identified as a b -jet (see Section 6.6.5).

Figure 8.5 reports the $E_{\text{T}}^{\text{miss}}$ and the leading jet p_{T} distributions after the background only fit described in Sections 7.5.4 and 8.4.6. It shows, within uncertainties, a good agreement between data and MC.

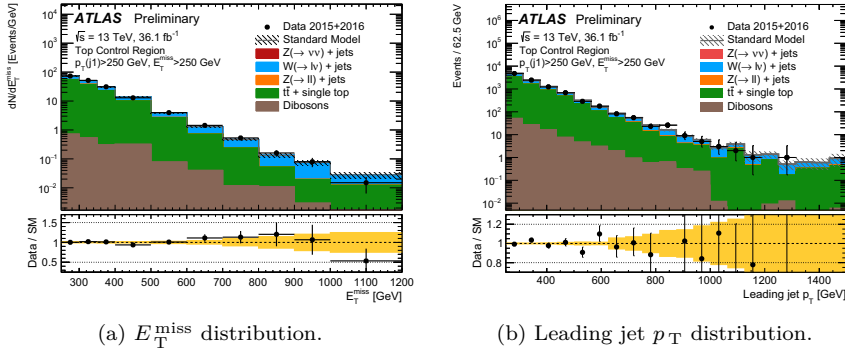


Figure 8.5: Observed and predicted $E_{\text{T}}^{\text{miss}}$ and leading jet p_{T} distributions after the background only fit in the CR_{Top} control region for the $E_{\text{T}}^{\text{miss}} > 250$ GeV inclusive selection. The error bands in the ratio plot on the bottom of the figures include statistical and systematic uncertainties. The negligible contribution of NCB and diboson backgrounds is not reported in the figure.

8.4.2 The Muon Control Region With b -Veto

The CR_{wmn} control region is enriched in $W(\rightarrow \mu\nu) + \text{jets}$ processes and contributes to the estimation of the $V + \text{jets}$ background in the signal region. It is defined in the same way as the CR_{Top} control region defined in Section 8.4.1 but b -jets, identified according to the criteria defined in Section 6.6.5, are vetoed. The $E_{\text{T}}^{\text{miss}}$ and the leading jet p_{T} distributions after the background only fit described in Sections 7.5.4 and 8.4.6 are shown in Figure 8.6. There is a good agreement between data and MC within uncertainties after the background only fit.

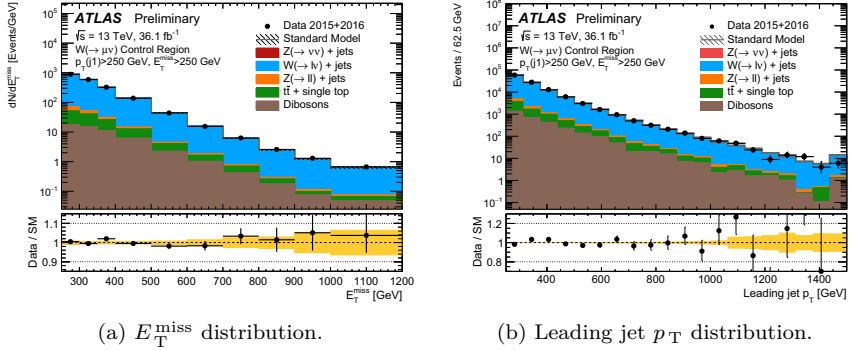


Figure 8.6: Observed and predicted E_T^{miss} and leading jet p_T distributions after the background only fit in the CR_{wmn} control region for the $E_T^{\text{miss}} > 250$ GeV inclusive selection corresponding to IM1. The error bands in the ratio plot on the bottom of the figures include statistical and systematic uncertainties. The negligible contribution of NCB and diboson backgrounds is not reported in the figure.

8.4.3 The Electron Control Region

The CR_{ele} is enriched in $W(\rightarrow e\nu) + \text{jets}$ processes and is used to constrain the $V + \text{jets}$ background contamination in the signal region. Besides cuts from B to H as defined in Section 8.3, this region has the following specific selections:

- The online trigger is required to select events with exactly one electron in the final state:
 - For the 2015 data the event selection requires a low electron p_T threshold of 24 GeV and medium likelihood quality combined with a higher p_T threshold of 60 and 120 GeV with medium and loose quality.
 - For the 2016 data a tight electrons with a p_T threshold of 26 GeV and loose isolation are combined with medium and loose quality electrons with 60 and 140 GeV p_T threshold respectively in order to select the events for the CR_{ele} region.
- Events with baseline muons are vetoed.
- In order to exclude electrons in the crack region (see Section 5.2), exactly one baseline electron with $p_T > 30$ GeV and $|\eta| > 1.52$ or $|\eta| < 1.37$ is selected.
- A tight isolation criteria as defined in Section 6.3 for the selected electrons is required. Multijet QCD events can enter the CR_{ele} region if they have a jet that fakes an electron. A tight isolation requires low calorimeter and track activity around the electron thus removing most jets and single pions faking electrons.

- The transverse mass of the $E_T^{\text{miss}} - e$ system is required to be $30 < m_T < 100$ GeV, compatible with a $W \rightarrow e\nu$ process.
- The events must satisfy: $E_T^{\text{miss}}/\sqrt{H_T} > 5 \text{ GeV}^{1/2}$ (E_T^{miss} significance) where H_T is the scalar sum of the p_T of the electron and the jets in the final state. Multijet QCD events can enter the CR_{ele} region if they have a fake electron and fake E_T^{miss} from mis-measured electrons or jet p_T . In case of fake E_T^{miss} the significance takes low values, that can be suppressed with this cut.

The single electron trigger, the veto of electrons in the crack region, the tighter isolation and the E_T^{miss} significance criteria all aim at reducing the multi-jet contribution in this region. In order to apply the theory corrections on $V + \text{jets}$ discussed in Section 8.4.5 the CRs have to be enriched in $W + \text{jets}$ events thus requiring a uniform cut definition between this control region and the CR_{wmn} . For this reason in this version of the analysis the cut on the transverse mass is introduced. Figure 8.7 shows the E_T^{miss} and leading jet p_T distribution of the CR_{ele} control region. There is a good agreement within uncertainties between data and MC after the background only fit.

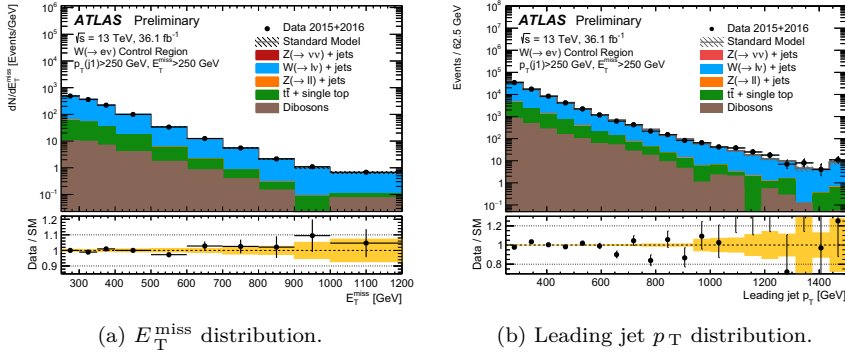


Figure 8.7: Observed and predicted E_T^{miss} and leading jet p_T distributions after the background only fit in the CR_{ele} control region for the $E_T^{\text{miss}} > 250$ GeV inclusive selection corresponding to IM1. The error bands in the ratio plot on the bottom of the figures include statistical and systematic uncertainties. The negligible contribution of NCB and diboson backgrounds is not reported in the figure.

8.4.4 The Di-Muon Control Region

The CR_{zmm} region is defined to select $Z(\rightarrow \mu^+\mu^-) + \text{jets}$ processes and contributes to constrain the $V + \text{jets}$ background contribution in the signal region. It is defined in the same way as in Section 7.5.3. Figure 8.8 reports the E_T^{miss} and leading jet p_T distributions after the background only fit, it can be seen that there is a good agreement between data and MC prediction within uncertainties.

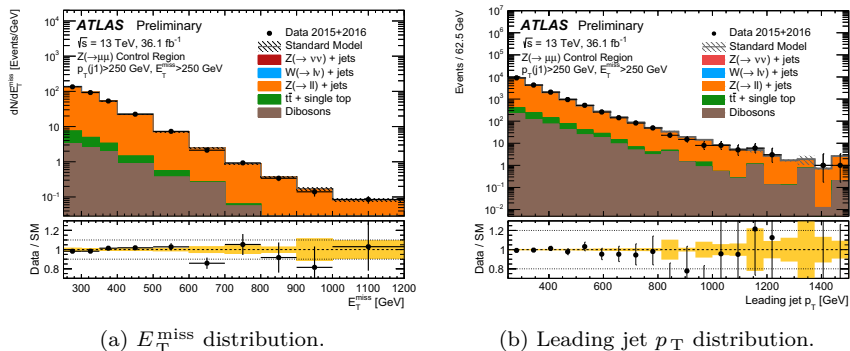


Figure 8.8: Observed and predicted E_T^{miss} and leading jet p_T distributions after the background only fit in the CR_{Zmm} control region for the $E_T^{\text{miss}} > 250$ GeV inclusive selection corresponding to IM1. The error bands in the ratio plot on the bottom of the figures include statistical and systematic uncertainties. The negligible contribution of NCB and diboson backgrounds is not reported in the figure.

8.4.5 Corrections to The $V + \text{jets}$ Processes

The most direct way to estimate the main $Z(\rightarrow \nu\bar{\nu}) + \text{jets}$ background would be from $Z(\rightarrow \ell^+\ell^-) + \text{jets}$ processes where $\ell = e, \mu$. However the smaller branching ratio of the Z boson decaying to leptons compared to the neutrinos makes the $Z(\rightarrow \ell^+\ell^-) + \text{jets}$ statistically limited. For this reason the decay of the Z boson to leptons as a proxy to estimate the $Z(\rightarrow \nu\bar{\nu}) + \text{jets}$ process with a statistical precision of less than 1% can be used up to a p_T of the boson of 600 GeV, beyond this value the $W(\rightarrow \ell\nu) + \text{jets}$ processes are statistically favored and should be used instead [121].

The theoretical extrapolation of the $Z(\rightarrow \nu\bar{\nu}) + \text{jets}$ background from the $Z(\rightarrow \ell^+\ell^-) + \text{jets}$ and $W(\rightarrow \ell\nu) + \text{jets}$ processes at high p_T can be influenced by several factor including the choice of the parton distribution function and the renormalization and factorization scales therefore it is important to include in the calculations higher order correction in QCD and Electroweak (EW) processes. The $V + \text{jets}$ samples used in this analysis are generated at the NLO and LO in QCD for up to two and four partons respectively and at the LO in EW using SHERPA-2.2.1 [122]. The samples are then normalized to an inclusive Next to Next to Leading Order (NNLO) cross section [123].

One possible way to include higher order corrections to the MC samples relies on the re-weighting of the differential cross section of these samples to NLO in QCD and NLO plus Next to Leading Logarithm (NLL) Sudakov approximation² at NNLO in EW. The framework proposed in Ref. [121] for applying

²At high energies EW corrections are dominated by single and double logarithms of the scale of the process over the mass of the vector boson (s/m_V^2 where $V = W, Z$) called Sudakov logarithms [124]. In the limit when all the kinematic invariants are much larger than the electroweak scale (the *Sudakov limit*) the EW one-loop corrections can be predicted in a process invariant way [125]. This Sudakov approximation is computationally easier than the full NLO EW correction and since it includes leading and sub-leading logarithmic terms it is

higher order corrections to the MC samples introduces a single re-weighting factor for all the V + jets Monte Carlo simulations with a set of relative systematic uncertainties that can be implemented as nuisance parameters in the profile likelihood method described in Section 7.9.1. It is based on the formula:

$$\frac{d}{dx} \frac{d}{d\vec{y}} \sigma^{(V)}(\vec{\epsilon}_{\text{MC}}, \vec{\epsilon}_{\text{TH}}) = \frac{d}{dx} \frac{d}{d\vec{y}} \sigma_{\text{MC}}^{(V)}(\vec{\epsilon}_{\text{MC}}) \left[\frac{\frac{d}{dx} \sigma_{\text{TH}}^{(V)}(\vec{\epsilon}_{\text{TH}})}{\frac{d}{dx} \sigma_{\text{MC}}^{(V)}(\vec{\epsilon}_{\text{MC}})} \right] \quad (8.4.1)$$

where the one dimensional parameter x is in our case the transverse momentum of the vector boson $x = p_{\text{T}}^{(V)}$, $\vec{y}$ represents other kinematic variables included in the MC simulation, $\vec{\epsilon}_{\text{MC}}$ and $\vec{\epsilon}_{\text{TH}}$ are the nuisance parameters associated with the Monte Carlo predictions and the theoretical higher order corrections introduced with this calculation. Finally $\sigma_{\text{MC}}^{(V)}$ and $\sigma_{\text{TH}}^{(V)}$ are the production cross sections for the V + jets samples computed from MC and from theory respectively and directly taken from Ref. [121].

The left hand side of Eq. (8.4.1) can be interpreted as the number of V + jets events in a given $E_{\text{T}}^{\text{miss}} = p_{\text{T}}(\text{V})$ and jet p_{T} bin after theory correction. In the right hand side the terms with the MC subscript refer to quantities determined using SHERPA-2.2.1 V + jets simulated events while the TH subscript refers to quantities taken from Ref. [121]. The first factor in the right hand side of Eq. (8.4.1) can be understood as the number of V + jets events in a given $E_{\text{T}}^{\text{miss}} = p_{\text{T}}(\text{V})$ and jet p_{T} bin, directly taken from SHERPA while the last factor, which is a ratio, can be interpreted as a transfer factor to convert the MC SHERPA to the theory estimate of Ref. [121].

The $d\sigma_{\text{TH}}(\vec{\epsilon}_{\text{TH}})/dx$ is first calculated at the LO and subsequently corrected independently at the NLO in QCD and NLO plus the Sudakov approximation at NNLO in EW:

$$\frac{d}{dx} \sigma_{\text{TH}}(\vec{\epsilon}_{\text{TH}}) = k_{\text{NLO}}(x, \vec{\epsilon}_{\text{TH}}) \cdot (1 + k_{\text{EW}}(x, \vec{\epsilon}_{\text{TH}})) \cdot \frac{d}{dx} \sigma_{\text{LO}} \quad (8.4.2)$$

where $k_{\text{EW}} = k_{\text{NLO}}^{\text{EW}} + k_{\text{NNLO}}^{\text{Sudakov}}$. The systematic uncertainties associated with this theoretical procedure are described in Section 8.5.2.

8.4.6 Fit Strategy

The profile likelihood method described in Section 7.9.1 is used to combine the background estimated with the different methods outlined in Section 8.4 and the relative systematic uncertainties described in Sections 8.5 and 8.6 in a global fit. The likelihood function is the product of Poisson probabilities, one for each bin used in the analysis defined in Section 8.3 and Gaussian probabilities that model the systematic uncertainties. The Poisson probabilities include as free parameters the normalization factors mentioned in Section 7.5.4 used to scale the relative contributions of the signal and background predictions. In the SUSY search described in Chapter 7 these factors were bin dependent while in this version of the analysis a unique normalization factor ($k^{\text{V+jets}}$) is used to

at the NLL accuracy level [125].

constrain all the $V + \text{jets}$ samples. This is justified by the use of the theoretical framework introduced in Section 8.4.5. Analogously a single normalization factor (k^t) is used for the $t\bar{t}$ and single top processes. The nuisance parameters used to model the systematic uncertainties and described in Section 8.5 are the free parameters that follow a Gaussian distribution of mean zero and standard deviation one.

The nuisance parameters and the normalization factors are adjusted to the optimal value by the fit which allows for a reduction of the otherwise relatively large theoretical and experimental systematic uncertainties (of the order of 20 - 40% before fit as explained in Section 8.5). Two fitting strategies are implemented in the analysis:

- A - A binned global likelihood fit simultaneous to all the exclusive E_T^{miss} region EM1–EM10. The nuisance parameters introduced in Section 8.5 are treated as bin-wise correlated systematics (the same nuisance parameter will be used in all the CRs and SRs) and treated accordingly in the fit. This fit takes advantage of the shape of the information from the E_T^{miss} distribution in the determination of the normalization of the $V + \text{jets}$ and top processes.
- B - A single bin likelihood fit to the different inclusive E_T^{miss} regions IM1–IM10 (a different fit for each inclusive region). The normalization factors $k^{V+\text{jets}}$, k^t and the nuisance parameters in this case refer to the specific E_T^{miss} region. This fit is used to derive the E_T^{miss} and leading jet p_T distributions in Figures 8.10a and 8.10b.

Table 8.3 summarizes the normalization factors and the control regions most constraining for their estimation. The $k^{V+\text{jets}}$ and k^t together with the signal strength μ are the free parameters of the global fit.

Background Process	Normalization Factor	Most Constraining Region
$V + \text{jets}$	$k^{V+\text{jets}}$	$\text{CR}_{\text{wmn}}, \text{CR}_{\text{zmm}}, \text{CR}_{\text{ele}}$
$t\bar{t}$, single top	k^t	CR_{Top}

Table 8.3: Summary table of the different background processes and the corresponding control regions most constraining for the normalization factors.

In Section 8.1 it was discussed that the ADD Monte Carlo samples generated with PYTHIA are simulated with a cut of 150 GeV on the p_T of the parton and the graviton. This cut is fully efficient only for signal regions with $E_T^{\text{miss}} \gtrsim 350$ GeV. For this reason for the model dependent interpretation of the ADD signal, only the $E_T^{\text{miss}} > 400$ GeV signal regions are considered.

8.5 Systematic Uncertainties

Systematic uncertainties on the knowledge of the shape and normalization of the E_T^{miss} can originate from theoretical and experimental sources. The systematic uncertainties are treated as Gaussian nuisance parameters in the fit

procedure described in Section 8.4.6. A complete list of all the parameters used to model the systematic uncertainties and included in the fit is given in the following sections.

8.5.1 Experimental Systematic Uncertainties

Experimental sources of systematic uncertainties derive from our limited knowledge of the energy and momentum scale of the reconstructed physical objects, on their reconstruction and identification. Different nuisance parameters have been used to model the impact of the experimental systematic uncertainties on the expected signal and background predictions. Correlations of systematic uncertainties among backgrounds or regions are implemented by using the same nuisance parameters while independent systematic uncertainties are implemented by using separate ones. In the following a full list of the nuisance parameters used in the analysis is given.

Luminosity

The integrated luminosity measured in data is used in the normalization of the MC samples, in order to account for uncertainties in the luminosity measurement, a flat 3.2% uncertainty for the 2015 and 2016 combined dataset has been applied to all MC event yields.

Jets Nuisance Parameters

`syst_JET_GroupedNP_i` ($i = 1, 2, 3$): These parameters account for uncertainties in the knowledge of the JES (see Section 6.6.3) and arise from Monte Carlo modeling and statistical uncertainties in in-situ JES determination.

`syst_JET_JER_SINGLE_NP`: This nuisance parameter models the uncertainty on the jet energy resolution in data.

`syst_JET_EtaIntercalibration_NonClosure`: The ATLAS jet energy scale calibration introduced in Section 6.6.3 is p_T and η dependent. In order to have a uniform jet response in different pseudorapidity regions for jets with same p_T an intercalibration is performed [82]. This parameter accounts for the non-closure of the jet intercalibration procedure.

b -tagging Nuisance Parameters

A common way to model disagreements between data and MC simulations in reconstruction and identification efficiencies is to determine correction factors using some auxiliary dataset for MC to data differences. These *scale factors* suffer themselves from systematic uncertainties arising from the specific method and dataset used to derive them.

`syst_FT_EFF_B_systematics`: This parameter models scale factor uncertainties in b -tagging (recognizing as a b -jet a jet originating from a b quark) efficiency.

`syst_FT_EFF_C_systematics`: This parameter models the uncertainty on the probability to b -tag a c -jet.

`syst_FT_EFF_Light_systematics`: This parameter models the uncertainty in mistagging light flavor (u , d , s and gluon) jets.

`syst_FT_EFF_extrapolation`: This parameter accounts for uncertainties due to the extrapolation procedure of the data driven b -tagging calibration efficiency from the low p_T to the high p_T regime [126].

`syst_FT_EFF_extrapolation_from_charm`: This parameter accounts for the extrapolation uncertainties related to the use of the charm jet b -tagging scale factors extrapolated to tau jets.

Electron Nuisance Parameters

`syst_EG_SCALE_ALL`: This parameter models the uncertainty on the electron energy scale.

`syst_EG_RESOLUTION_ALL`: This parameter takes into account the uncertainties on the electron resolution.

`syst_EL_EFF_Reco_TOTAL`: This parameter accounts for the uncertainty on the ratio of electron reconstruction efficiency in data to Monte Carlo.

`syst_EL_EFF_ID_TOTAL`: This parameter models uncertainties on the ratio of data to Monte Carlo efficiencies for the electron identification selections. This refers to the “loose” and “tight” selections which define the baseline and good electron introduced in Section 6.3.

`syst_EL_EFF_Iso_TOTAL`: This parameter models the uncertainties on the isolation efficiency, namely the LooseTrackOnly listed in Table 6.1.

`syst_EL_EFF_TriggerEff_TOTAL`: This parameter account for the uncertainty on the electron trigger efficiency for the trigger used to define the electron control region in Section 8.4.3.

`syst_EL_EFF_Trigger_TOTAL`: The electron reconstruction and identification efficiencies are derived using data selected with electron triggers. This parameter accounts for the effect of uncertainties from the trigger on these efficiencies.

Muon Nuisance Parameters

`syst_MUON_ID`: This nuisance parameter takes into account the impact of the uncertainty from the measurement of the inner detector track momentum on the final muon momentum.

`syst_MUON_MS`: This parameter accounts for the impact of the uncertainty from the measurement of the muon system track momentum on the final muon momentum.

syst_MUON_EFF_STAT (SYS): This parameter accounts for the statistical (systematic) uncertainty on the data to Monte Carlo muon reconstruction scale factor.

syst_MUON_EFF_STAT (SYS) _LOWPT: This parameter accounts for the statistical (systematic) uncertainties on the data to Monte Carlo muon reconstruction scale factor for low p_T muons.

syst_MUON_SCALE: This parameter accounts for the uncertainty on the muon momentum scale.

syst_MUON_BADMUON_STAT (SYS): This parameter models bad muon statistical (systematic) uncertainties for high p_T muons (above 100 GeV) due to the specific muon working point used for high p_T muons [127].

syst_MUON_SAGITTA_RESBIAS (RHO): The muon momentum measurement relies on the precise knowledge of the sagitta of the muon track between muon stations. The calibration constants used to calibrate the muon Monte Carlo are derived from $Z \rightarrow \mu\mu$ and J/ψ data in bins of η and ϕ . These parameters model uncertainties on the calibration procedure and residual charge dependent bias in muon momentum scale.

syst_MUON_TTVA_STAT (SYS): This parameter models the statistical (systematic) uncertainties on the efficiency of the muon track-to-vertex matching efficiency.

Missing Transverse Energy Nuisance Parameters

The soft missing energy ($E_{x(y)}^{\text{miss, soft}}$) is a component of the E_T^{miss} calculation described in Section 6.8 that relies on tracks associated to the primary vertex and not associated with any hard object such as electrons, muons, jets, taus or photons. The detector response to the soft missing energy reconstruction is parametrized by an energy scale, a resolution in the perpendicular direction to the hard scatter and a resolution along the main axis of the hard scatter. The Monte Carlo to data mismodeling with respect to these three parameters that characterize the detector response to soft missing energy is modeled with three systematic uncertainties: syst_MET_SoftTrk_Scale, syst_MET_SoftTrk_ResoPara and syst_MET_SoftTrk_ResoPerp.

Pileup Nuisance Parameters

The generation of Monte Carlo samples is done before or during the data taking period when the pile-up conditions are not yet known. For this reason at the analysis level it is necessary to re-weight the MC pile-up conditions to what is found during the data taking, this procedure is commonly referred to as *pile-up re-weighting*.

syst_PRW_DATASF: Once the Monte Carlo samples have been re-weighted to the same distribution of number of pp interaction per bunch crossing (see Section 7.8.3) as in data, it is observed that the event activity in data

and simulation is still slightly different. Auxiliary datasets are used to derive a scaling of the Monte Carlo activity to match that of data. This nuisance parameter represents the uncertainty on the scaling to match the event activity in Monte Carlo to that of data at equal pile-up.

8.5.2 Background Theoretical Systematic Uncertainties

Sources of theoretical systematic uncertainties include variations due to the choice of the renormalization, factorization and parton shower matching scales and the limited knowledge of the PDFs. In the following sections an overview of the treatment of these uncertainties for $V + \text{jets}$, top and diboson backgrounds is provided.

$V + \text{jets}$ High Order Corrections

In this analysis only a single overall scale factor $k^{V+\text{jets}}$ is considered (see Section 8.4.6), it scales up or down the entire $W + \text{jets}$ and $Z + \text{jets}$ predictions by the same amount in all the E_T^{miss} bins defined in the analysis. This means effectively that the momentum spectrum of the vector boson or its experimental proxy, the missing energy, is taken from Monte Carlo simulation. In the 2015 analysis one independent scale factor for each E_T^{miss} bin was used in order to account for the limited knowledge of the shape of the p_T spectrum of the gauge boson. In this version of the analysis this is accounted for using instead theoretical nuisance parameters taken from Ref. [121] associated with the $V + \text{jets}$ normalization strategy given in Section 8.4.5. These nuisance parameters allow the shape of the boson p_T to change with respect to corrected Monte Carlo simulation within certain allowed theory boundaries and in the global fit implemented in the analysis are labeled as:

`vjets_d1K_NLO`: Accounts for uncertainties in the normalization of the vector boson p_T distribution due to QCD scale variations.

`vjets_d2K_NLO`: This parameter accounts for shape uncertainties of the vector boson p_T distribution due to QCD scale variations.

`vjets_d3K_NLO`: This parameter accounts for the correlation of QCD uncertainties in the different $V + \text{jets}$ processes.

`vjets_d1kappa_EW`: This parameter estimate the uncertainty of unknown high p_T EW effects beyond NNLO.

`vjets_d2kappa_EW_ej`: This parameter accounts for unknown terms in the NLO approximation at high p_T and for NNLO effects that might become relevant when hard contributions dominate. It is used for the $Z + \text{jets}$ processes.

`vjets_d2kappa_EW_evj`: This parameter accounts for the same effects as the previous one but it is used to constrain the $W + \text{jets}$ processes.

`vjets_d3kappa_EW_eej`: This term accounts for the limitations of the Sudakov approximation for the $Z + \text{jets}$ processes.

`vjets_d3kappa_EW_evj`: This term accounts for the limitations of the Sudakov approximation for the $W + \text{jets}$ processes.

`vjets_dK_NLO_mix`: The QCD and EW are treated as separate corrections, this term accounts for unknown interference terms not taken into account.

Since both the `vjets_d2kappa_EW_eej` and `vjets_d2kappa_EW_evj` and both the `vjets_d3kappa_EW_eej` and `vjets_d3kappa_EW_evj` parameters all describe subleading unknown effects they are treated with independent nuisance parameters for each process in the fit. The other parameters are treated as correlated between the $V + \text{jets}$ processes and vector boson p_T but independent from one another. The nuisance parameters are applied by re-weighting each event based on the theory input of Section 8.4.5, in particular up and down variation histograms of the p_T distributions of the vector boson are obtained applying an event weight and treated in a consistent way to other normalization and shape uncertainties.

Due to the limited Monte Carlo statistics, statistical fluctuations are possible in the QCD correction evaluations. For this reason a linear fit is applied to the ratio of the differential cross sections in the SHERPA MC and the $V + \text{jets}$ theory prediction. The “`vjets_QCDSmoothing`” nuisance parameter is then introduced in order to take into account uncertainties coming from this fit.

V + jets Parton Distribution Function Uncertainties

The $V + \text{jets}$ samples are generated using SHERPA [122], the intra-PDF uncertainty is computed taking the standard deviation of the 100 PDFs provided by the NNPDF [99] group.

Two alternative sets of PDFs provided by the CT14 [128] and the MMHT [100] groups are also considered where the maximal difference between the central value of the alternative PDF set with respect to the NNPDF one is taken as the inter-PDF uncertainty.

PDFComb: This parameter accounts for the final parton distribution function uncertainty on the $V + \text{jets}$ samples evaluated as the quadratic sum of the inter and intra-PDF uncertainties.

Table 8.4 reports the resulting relative PDF systematic uncertainties on the $V + \text{jets}$ backgrounds as a function of the exclusive signal region bins EM1-EM10. The uncertainties of the various processes are quoted in the region where that process is dominant. As can be seen the PDF uncertainties are of the order of 1-2% in the lowest missing energy bins but rise with higher E_T^{miss} since these regimes probe more extreme kinematic regions with less well known PDFs.

	PDF Uncertainty									
	EM1	EM2	EM3	EM4	EM5	EM6	EM7	EM8	EM9	EM10
$Z \rightarrow \nu\nu$ (SR)	1.8	1.8	2.2	2.6	2.7	3.2	3.4	3.3	4.2	6.7
$Z \rightarrow \tau\tau$ (SR)	2.3	2.0	2.9	2.1	3.6	2.5	3.5	9.5	11.6	8.3
$Z \rightarrow \mu\mu$ (CR _{zmm})	1.7	1.9	2.2	2.3	2.8	3.3	4.0	4.2	5.6	5.2
$W \rightarrow \mu\nu$ (CR _{wmn})	1.5	1.9	1.8	2.2	2.5	4.1	3.4	4.4	4.5	6.7
$W \rightarrow e\nu$ (CR _{ele})	1.5	1.6	2.0	2.3	2.6	3.2	3.8	3.6	4.8	7.5
$W \rightarrow \tau\nu$ (SR)	1.3	1.1	1.3	1.3	2.5	2.1	3.3	4.9	7.1	12.0

Table 8.4: Systematic uncertainties in % on the $V + \text{jets}$ background as a function of the exclusive signal regions EM1-EM10. The uncertainties on a given process are quoted in the region where this process is dominant.

$V + \text{jets}$ Matching Uncertainties

Uncertainties on the predictions of the $V + \text{jets}$ samples due to the choice of the matching scale (see Section 7.8.3) for the parton showers to the matrix element matching in SHERPA [122] are derived by varying the parton matching scale between 15 and 30 GeV, while 20 GeV is the default for the $V + \text{jets}$ samples. This is practically implemented via event weights derived specifically for the ATLAS $V + \text{jets}$ Monte Carlo samples. It allows to obtain up and down variations of the E_T^{miss} distribution using an event-by-event re-weighting based on the number of jets with $p_T > 20$ GeV and $|\eta| < 2.8$ and the p_T of the Z boson. The nuisance parameter associated with the matching uncertainty is labeled “ckkw”.

Diboson Background

The SHERPA [92] and POWHEG [120] interfaced with PYTHIA8 [114] generators are used to produce the diboson (WW , WZ and ZZ processes) MC samples. After taking into account possible inclusive cross section differences of the two generators, the full difference in yields in each E_T^{miss} bin of the two generators is taken as the systematic uncertainty and added in quadrature with a 6% theory uncertainty on the NLO cross section.

`diboson_Sys`: This parameter accounts for the diboson systematic variations.

Table 8.5 reports the diboson uncertainties evaluated from the comparison between the two generators for the different exclusive regions EM1-EM10 defined in the analysis for the signal region and the different control regions. The diboson systematic uncertainty is observed to grow at higher missing transverse energy, where in particular the two generators can differ by up to 31% in the signal region. The generators are however in very good agreement at lower missing energies. In the 2015 analysis the agreement between the two generators ranged between 5% and 25% for E_T^{miss} bins below 800 GeV increasing up to 50% above 800 GeV. The diboson systematic uncertainties in this version of the analysis are generally lower due to the higher order of the Monte Carlo generators used.

Diboson Uncertainty										
	EM1	EM2	EM3	EM4	EM5	EM6	EM7	EM8	EM9	EM10
SR	3.6	1.6	0.4	3.3	7.3	11.3	15.3	19.3	23	31
CR _{ele}	2.5	1.0	0.4	2.5	5.3	8.2	11.0	13.9	16.7	22.4
CR _{wmn}	0.4	0.2	0.0	0.3	0.8	1.2	1.6	2.0	2.4	3.3
CR _{Top}	1.2	2.1	3.1	4.5	6.5	8.4	10.3	12.2	14.2	18.0
CR _{zmm}	0.9	0.6	0.2	0.2	0.8	1.4	2.1	2.7	3.3	4.5

Table 8.5: Relative systematic uncertainties in % on the diboson background as a function of the exclusive bins EM1-EM10 defined in the analysis for the signal region and the different control regions as evaluated from the comparison between the generators. These values are then added in quadrature with a 6% uncertainty on the NLO diboson cross section.

Top Quark Background

Uncertainties due to different event generators, parton shower models and two models of initial and final state radiation are estimated using dedicated $t\bar{t}$ MC samples and used also for single top processes. This choice of the Monte Carlo samples is motivated by the fact that the $t\bar{t}$ process contributes for approximately 80% of the top background ($t\bar{t}$ and single top) in the signal region. Except for the variation due to the amount of initial and final state radiation where the semi-difference between the higher and lower variations is considered, the change in the MC generator and parton shower model variations are symmetrized and applied considering them uncorrelated. Since similar $E_{\text{T}}^{\text{miss}}$ shapes are observed in signal and control regions and in order to avoid statistical fluctuations, the final uncertainties are evaluated in a single region obtained by summing the yields of the signal and control regions. The total systematic uncertainties on the top background are obtained summing in quadrature each variation in each $E_{\text{T}}^{\text{miss}}$ bin.

The results are reported in Table 8.6 where it can be seen that the uncertainty grows with the missing energy. This growth is mildly observed in both the parton shower and the ISR and Final State Radiation (FSR) model variations while it is more evident in the uncertainty associated with the change of the event generator. The top systematic uncertainty in this version of the analysis are compatible with those calculated in 2015.

The introduction of the CR_{Top} control region (see Section 8.4.1) is motivated by these large systematic uncertainties which are constrained by the “top_Sys” nuisance parameter in the global fit.

Top Uncertainty									
EM1	EM2	EM3	EM4	EM5	EM6	EM7	EM8	EM9	EM10
24	26	28	31	36	42	47	53	59	72

Table 8.6: Total relative systematic uncertainty in % on the top background in each exclusive bin EM1-EM10 used in the analysis evaluated in a single region obtained by summing the yields of the signal and control regions.

8.6 Signal Theoretical Systematic Uncertainties

The theoretical uncertainties associated to the ADD signal derive from the uncertainties on the parton distribution function, the renormalization and factorization scales, the tuning of the parton shower in PYTHIA [114] and the initial and final state radiation modeling. Since the Monte Carlo samples are generated at the leading order, there is no parton matching scale uncertainty. The exact definition of these uncertainties was presented in details in Section 7.8. Here the focus is on the information relevant to the ADD models.

8.6.1 Parton Distribution Function Uncertainties

The ADD signal sample is generated using the PYTHIA8 [114] generator with the NNPDF2.3LO PDF set. To estimate the uncertainty deriving from the limited knowledge of the PDFs the same recommendations from the PDF4LHC group used in the estimation of the PDF uncertainties for the SUSY compressed models and described in Section 7.8.1 are used.

The PDF uncertainties affects both the overall ADD signal normalization or cross section and the signal acceptance. Figure 8.9 presents the observed variations in the number of signal ADD events for intra- and inter-PDF uncertainties. The resulting relative yield variations gives the relative PDF uncertainty on the ADD cross section. The final PDF uncertainty is the envelope that contains the error bands of the three families.

Table 8.7 summarizes the uncertainties on cross section and acceptance. The growth of the uncertainty with both the number of extra dimensions and the missing energy is understood with a similar argument as the one given in Section 8.2. Larger numbers of extra dimensions require heavier gravitons (see Figure 8.3) and thus a higher center of mass energy of the originating partons. In these kinematic regions the PDFs are less well known. A similar argument can be used in the case of higher E_T^{miss} bins where higher momentum gravitons are produced.

	PDF Uncertainty				
	ADD n = 2	ADD n = 3	ADD n = 4	ADD n = 5	ADD n = 6
$\Delta\sigma$	11	18	27	35	43
$\Delta A(400 < E_T^{\text{miss}} < 500)$	8	11	12	13	13
$\Delta A(500 < E_T^{\text{miss}} < 600)$	13	13	15	11	14
$\Delta A(600 < E_T^{\text{miss}} < 700)$	16	16	18	20	8
$\Delta A(700 < E_T^{\text{miss}} < 800)$	18	21	19	15	20
$\Delta A(800 < E_T^{\text{miss}} < 900)$	21	19	20	12	9
$\Delta A(900 < E_T^{\text{miss}} < 1000)$	20	22	21	15	19
$\Delta A(E_T^{\text{miss}} > 1000)$	32	26	28	30	26

Table 8.7: Systematic uncertainties in % on PDFs. The uncertainty is the envelop the contains the signal yields from the three PDF families, plus their error bands.

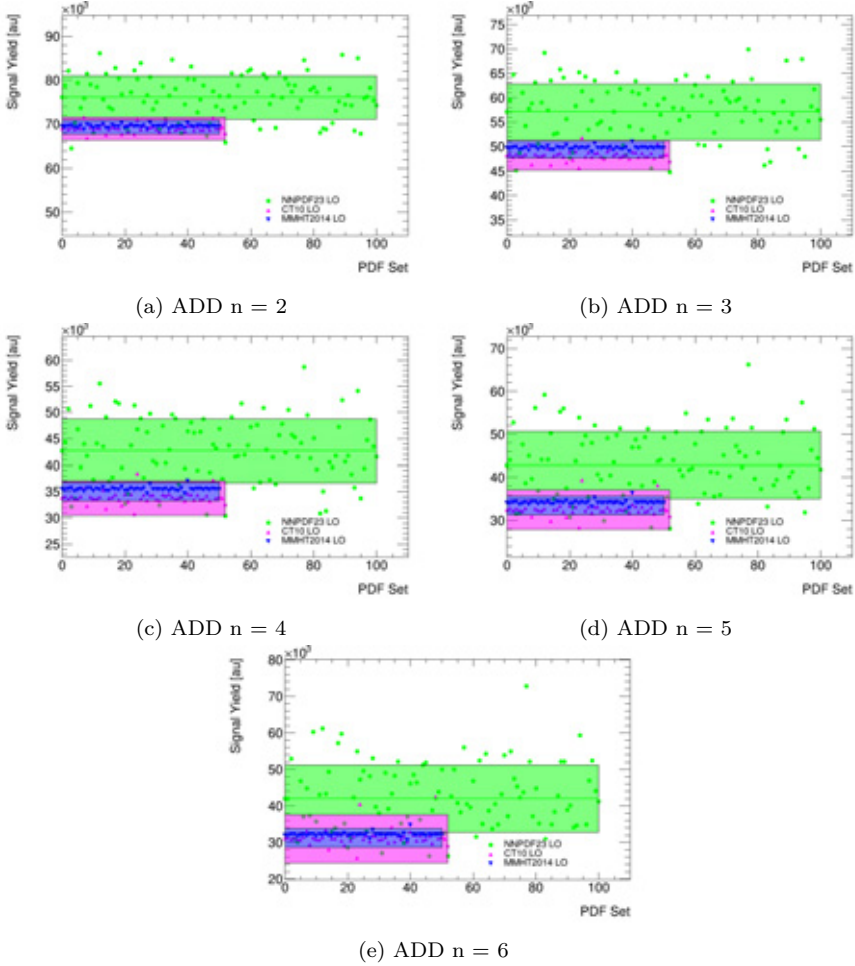


Figure 8.9: Observed variations in un-normalized number of ADD events when the PDF sets are varied by the uncertainties for $n = 2$ to $n = 6$. The event yield before any selection from three different PDF families are compared. The final uncertainty on the ADD cross section is the envelop that contains all three families with their error bands. The same procedure is also applied to the acceptances.

8.6.2 Scale Uncertainties

Similarly to the SUSY compressed scenario, the renormalization and factorization scale uncertainties are obtained by re-generating the signal MC samples varying the two scales by a factor of 2 and 0.5 respectively. The final uncertainty is the average of the up and down variations. This source of uncertainty mostly affects the cross section and the results are collected in Table 8.8.

Renorm./Fact. Scale Uncertainty					
	ADD n = 2	ADD n = 3	ADD n = 4	ADD n = 5	ADD n = 6
$\Delta\sigma$	23	27	30	33	36

Table 8.8: Systematic uncertainty in % on the ADD signal normalization arising from the factorization and renormalization scales. The scales are varied up and down simultaneously, the final uncertainty is the average of these variations.

8.6.3 Tune, Initial and Final State Radiation Uncertainties

The evaluation of the tune, ISR and FSR uncertainties follows what is done for the SUSY compressed models and described in Section 7.8.3. Since the generator contains only the matrix elements for the diagram given in Figure 8.1, there is no parton shower matching uncertainty in this case. Five tune parameters have been varied in order to account for uncertainties from:

- Underlying event effects.
- Jet structure effects.
- Aspects of the MC generation that provide extra jet production.

For each ADD model, ten systematic samples are produced and analyzed at the truth level, the effect of the acceptance is evaluated in different bins of truth missing energy corresponding to those used in the signal regions. The final uncertainty in each bin affects the signal acceptance and is a common envelope valid for the different extra dimensions models (from ADD n = 2 to ADD n = 6). The results are summarized in Table 8.9.

Initial and Final State Radiation							
E_T^{miss} [GeV]	400-500	500-600	600-700	700-800	800-900	900-1000	1000-3000
$\Delta A(\%)$	7	7	10	13	18	13	9

Table 8.9: Tune, initial and final state radiation uncertainties in % on the acceptance of the different signal region E_T^{miss} bins in the analysis. The final value is a common envelope valid for all the ADD models between n = 2 and n = 6 dimensions.

8.7 Results and Interpretations

The fit strategies described in Section 8.4.6 are applied to the control regions in a background only fit in order to compute the background yields and the impact of the systematic uncertainties to the background expectations in the signal region. The results of this fit are presented in Section 8.7.1. For the model independent interpretation of the ATLAS data also the signal regions are used in a background only fit and the results presented in Section 8.7.2. For

the interpretation of the data in terms of specific models the exclusive signal and control regions are used in a signal plus background fit and the results of this fit are presented in Section 8.7.3.

8.7.1 Signal Region Yields After Background Fit

The number of events observed in data and in the Standard Model predictions for a selected set of exclusive signal regions as defined in Section 8.3 after the background only fit to the control regions are reported in Table 8.10. The total uncertainty of the SM predictions range between 2% for EM2 and 4.2% for EM9. A multiplicative factor of 1.26 ± 0.13 and 1.31 ± 0.17 is calculated for the

Analysis Results					
Exclusive Signal Region Observed events (36.1 fb ⁻¹)	EM2 67475	EM4 27843	EM6 2975	EM8 512	EM9 223
SM prediction	67100 ± 1400	27640 ± 610	2825 ± 78	463 ± 19	213 ± 9
$W(\rightarrow e\nu) + \text{jets}$	5510 ± 140	1789 ± 59	147 ± 9	18 ± 1	8 ± 1
$W(\rightarrow \mu\nu) + \text{jets}$	6120 ± 200	2021 ± 82	173 ± 9	21 ± 5	11 ± 1
$W(\rightarrow \tau\nu) + \text{jets}$	13680 ± 310	4900 ± 110	397 ± 11	55 ± 5	29 ± 2
$Z(\rightarrow ee) + \text{jets}$	0.03 ± 0	0.02 ± 0.02	0 ± 0	0 ± 0	0 ± 0
$Z(\rightarrow \mu^+\mu^-) + \text{jets}$	167 ± 8	36 ± 2	2 ± 0.2	0.4 ± 0.1	0.5 ± 0.1
$Z(\rightarrow \tau\tau) + \text{jets}$	185 ± 6	68 ± 4	5.1 ± 0.3	0.3 ± 0.1	0.31 ± 0.04
$Z(\rightarrow \nu\nu) + \text{jets}$	37600 ± 970	17070 ± 460	1933 ± 57	337 ± 12	153 ± 7
$t\bar{t}$, single top	2230 ± 200	848 ± 86	43 ± 6	4 ± 1	1.3 ± 0.4
Diboson	1327 ± 90	874 ± 64	124 ± 16	26 ± 5	10 ± 2
Multijet background	170 ± 160	13 ± 13	1 ± 1	1 ± 1	0.1 ± 0.1
Non-collision background	71 ± 71	18 ± 18	0 ± 0	0 ± 0	0 ± 0

Table 8.10: Number of events in the observed data and Standard Model predictions in a representative set of signal regions as defined in Section 8.3 after the background only fit to the exclusive control regions EM1–EM10. For the SM predictions the quoted uncertainty includes both statistical and systematic uncertainties. The individual uncertainties reported for each background process can be correlated and do not necessarily add in quadrature to the total background uncertainty.

$V + \text{jets}$ and top background respectively. These factors are applied in Figure 8.10 which shows the comparison between the measured distributions and the Monte Carlo predictions for the E_T^{miss} and leading jet p_T in the inclusive signal region with $E_T^{\text{miss}} > 250$ GeV after the fit.

The nuisance parameters defined in Section 8.5 are used to estimate the impact of the systematic uncertainties on the background predictions. They are treated as bin-wise correlated and treated accordingly in the fit. Figure 8.13 shows the value of the nuisance parameters after the background only fit to the data of the control regions. In most cases there is a mild constrain of the parameters which is more evident in the case of the top uncertainty.

After the fit the agreement between data and theory expectations is good. Since no significant excess of data compared to the predictions is observed, the statistical approach described in Section 7.9.1 is used to set 95% CL model independent exclusion limits (see Section 8.7.2) and interpreted in terms of the ADD scenario for large extra dimensions in Section 8.7.3. The observed limits are in general slightly worse than the expected sensitivity due to the small

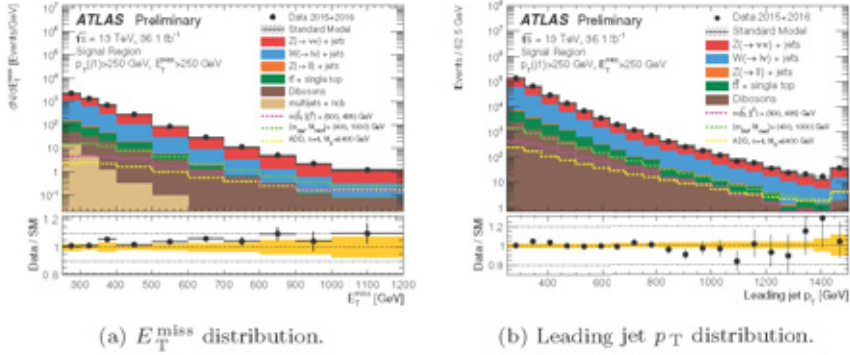


Figure 8.10: Distribution of the E_T^{miss} and the leading jet p_T for IM1 signal region compared with the background estimates from the background only fit to the control regions. The distributions of different signal models are superimposed for comparison. The contribution from the multi-jet and NCB background is negligible and not reported in the plot. In the ratio window the error bars include experimental and systematic uncertainties.

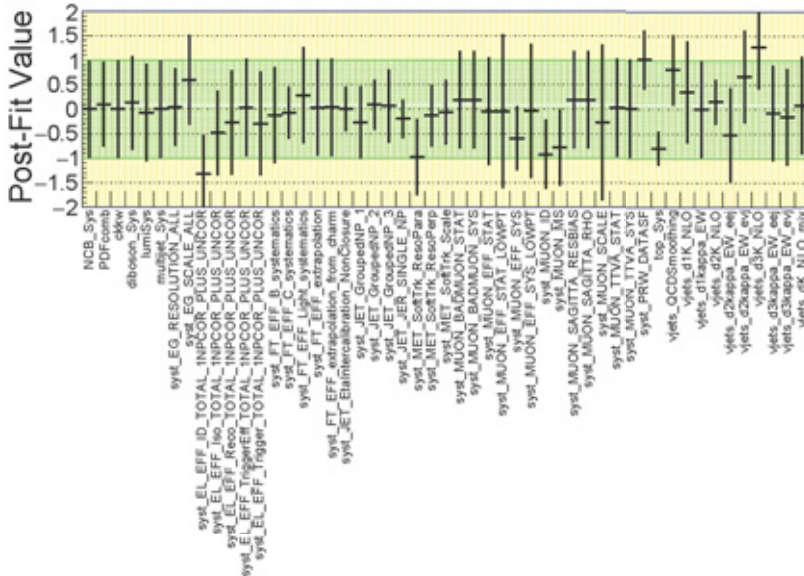


Figure 8.11: Values of the nuisance parameters used in the analysis after the background only fit to the exclusive control regions EM1–EM10.

excess of data compared to the Standard Model expectation as can be seen from Table 8.11.

Exclusive Signal Region Yields		
Region	Predicted	Observed
EM1	111100 ± 2300	111203
EM2	67100 ± 1400	67475
EM3	33820 ± 940	35285
EM4	27640 ± 610	27843
EM5	8360 ± 190	8583
EM6	2825 ± 78	2975
EM7	1094 ± 33	1142
EM8	463 ± 19	512
EM9	213 ± 9	223
EM10	226 ± 16	245

Table 8.11: Total yields for Standard Model predictions and data in the different exclusive signal region selections after the background only fit to the control regions EM1–EM10. The statistical and systematic uncertainties are also quoted on the SM predictions.

8.7.2 Model Independent Limits

Exclusion limits at 95% CL are set on the visible cross section defined in 7.9.3 using the CL_S modified frequentist approach for models with a monojet signature in the final state. Table 8.12 reports the results of the likelihood fit on the inclusive signal and control regions IM1–IM10 defined in Section 8.3, values of σ_{vis} above $\langle\sigma\rangle_{\text{obs}}^{95}$ are excluded at 95% CL. These results can be understood as follow: BSM models that predicts more than 531 and 1.6 events per inverse femtobarn of integrated luminosity in the IM1 and IM10 region respectively are excluded by the data collected with the ATLAS detector.

Result Table			
Signal Region	$\langle\sigma\rangle_{\text{obs}}^{95}$ [fb]	S_{obs}^{95}	S_{exp}^{95}
IM1	531	19135	11700^{+4400}_{-3300}
IM2	330	11903	7000^{+2600}_{-2600}
IM3	188	6771	4000^{+1400}_{-1100}
IM4	93	3344	2100^{+770}_{-590}
IM5	43	1546	770^{+280}_{-220}
IM6	19	696	360^{+130}_{-100}
IM7	7.7	276	204^{+74}_{-57}
IM8	4.9	178	126^{+47}_{-35}
IM9	2.2	79	76^{+29}_{-21}
IM10	1.6	59	56^{+21}_{-16}

Table 8.12: Expected and observed upper limits on the number of events, S_{exp}^{95} and S_{obs}^{95} respectively and on the visible cross section at 95% CL.

8.7.3 ADD Model Interpretation

The analysis results are interpreted in terms of limits on the effective Planck scale in $4 + n$ dimensions M_D in the ADD model [3] for LED. To this end a binned likelihood fit that includes both the signal and background predictions to all the exclusive $E_T^{\text{miss}} > 400$ GeV (see Section 8.4.6) regions EM4-EM10 as defined in Section 8.3 is performed.

Figure 8.12 reports the $V + \text{jets}$ and the top processes normalization factors obtained from the background only fit and applied to the MC predictions. It can be seen that a multiplicative factor of 1.35 ± 0.13 and 0.87 ± 0.23 is calculated for the $V + \text{jets}$ and top background respectively.

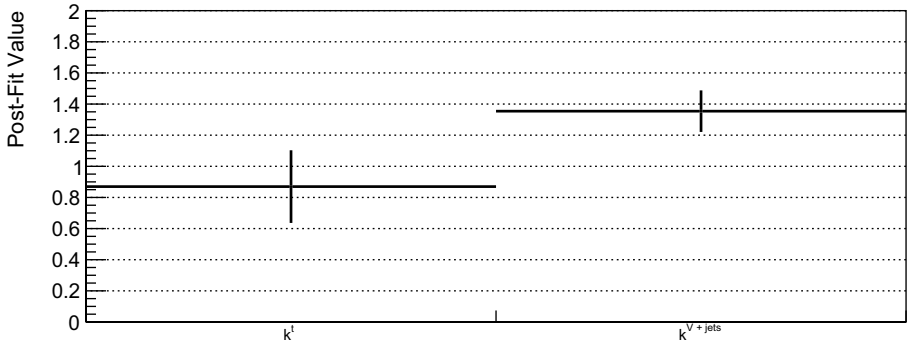


Figure 8.12: Values of the normalization factors for the $V + \text{jets}$ and single top backgrounds after the background only fit to the data that includes both the control and signal regions in all the exclusive E_T^{miss} bins EM1-EM10 defined in the analysis.

The nuisance parameters defined in Section 8.5 and used to estimate the impact of the systematic uncertainties on the background predictions are treated as bin-wise correlated in the fit. Figure 8.13 shows the value of the nuisance parameters after the background only fit to the data of the control and signal regions EM4-EM10. In most cases there is a mild constrain of the parameters. In this case the impact of the fit on the top systematic is less important. This is understood in terms of the better precision of top processes measurements in low E_T^{miss} regions (as can be seen from Table 8.6) that increase the sensitivity to this systematic uncertainty and therefore the constrain of the parameter by the fit.

Table 8.13 reports the number of observed and predicted events in some of the exclusive signal regions defined in the analysis and used in the background only fit to the signal and control regions.

Figure 8.14 show the expected and observed limits at 95% CL on M_D as a function of the large extra dimensions n . The dashed blue line shows the expected limits using 36.1 fb^{-1} with the $\pm 1\sigma$ and $\pm 2\sigma$ error bands (green and yellow band respectively). The previous results obtained with 3.2 fb^{-1} are also reported for comparison, as can be seen a significant improvement to the analysis is expected. The observed limit (solid black line) is significantly worse than the expectations but within uncertainties. This is due to the excess of data compared to SM predictions as reported in Table 8.10.

reported in parentheses. The effect of the the truncation is only noticeable in the ADD $n = 6$ model meaning that the analysis is probing values of M_D where the effective theory can be trusted.

Values of M_D up to 7.7 TeV for $n = 2$ dimensions and 4.8 TeV for $n = 6$ dimensions are excluded at 95% CL improving previous results. The expected sensitivity of this search to M_D has increased significantly thanks to a dataset ten times larger than the previous search and improved techniques for background calculations in particular to constrain the top and the $V + \text{jets}$ backgrounds. The statistical and systematic uncertainties contribute almost equally to the total uncertainty on the background estimation in the EM4 signal region. The systematic uncertainty is dominated by the data statistics in the control regions constraining the $V + \text{jets}$ background and to a lesser extent by its theory systematic. In higher E_T^{miss} regions the statistical uncertainty becomes the largest source of uncertainty on the background while in the lower bins the systematic uncertainties associated with the $V + \text{jets}$ normalization and the control region statistics are dominant. Due to the relation between the cross section and M_D of the form $M_D = \sigma^{1/(n+2)}$ further increase in sensitivity will be difficult without reduced systematics and much larger datasets (of the order of at least 10 times larger).

95% CL Limits on M_D [TeV]		
ADD Model	Expected	Observed (damped)
$n = 2$	$9.2^{+0.8}_{-1.0}$	$7.7^{+0.4}_{-0.5}$ (7.7)
$n = 3$	$7.1^{+0.5}_{-0.6}$	$6.2^{+0.4}_{-0.5}$ (6.2)
$n = 4$	$6.1^{+0.3}_{-0.4}$	$5.5^{+0.3}_{-0.5}$ (5.5)
$n = 5$	$5.5^{+0.3}_{-0.3}$	$5.1^{+0.3}_{-0.5}$ (5.1)
$n = 6$	$5.2^{+0.2}_{-0.3}$	$4.8^{+0.3}_{-0.5}$ (4.8)

Table 8.14: Expected and observed 95% CL lower limits on the fundamental Planck scale M_D in $4 + n$ dimensions as a function of the number of extra dimensions n . The impact of the $\pm 1\sigma$ uncertainty from the theory on the observed limits and the expected $\pm 1\sigma$ range of limits in absence of a signal is reported. The 95% CL observed limits after damping the signal cross section for $\hat{s} > M_D^2$ are also reported in parentheses.

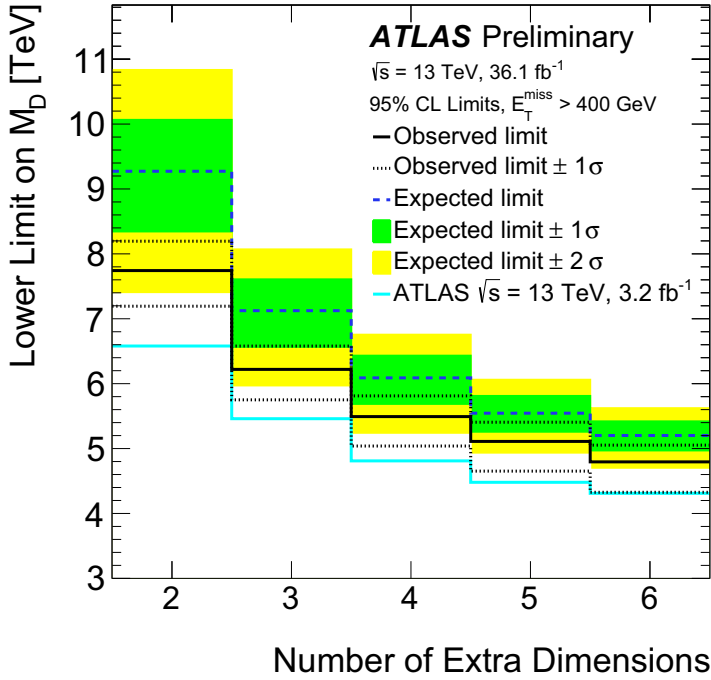


Figure 8.14: Expected and observed limits at 95% CL on the fundamental Planck Scale in $4 + n$ dimensions, M_D , as a function of the number of extra dimensions n . The dashed blue line shows the expected limit using 36.1 fb^{-1} , the green and yellow bands are the $\pm 1\sigma$ and $\pm 2\sigma$ uncertainties on the estimate. The solid black line is the observed limit while the cyan line represents the observed limits in the 2015 analysis using 3.2 fb^{-1} .

Chapter 9

Conclusions

The ATLAS detector is a multipurpose detector used for precision measurements of a wide range of Standard Model processes and the search for new decay channels of the Higgs boson and beyond the Standard Model phenomena such as supersymmetry, extra dimensions or dark matter particles.

The Tile Calorimeter is the hadronic calorimeter covering the most central region of the ATLAS detector up to $|\eta| < 1.7$. It plays a vital role in measuring the energy of jets and missing energy used in this thesis to search for SUSY and extra dimensions. To this end, an understanding of the electronic noise inside the calorimeter is crucial as it affects the signal left by the particles crossing the calorimeter and must be known precisely. The exact noise level per calorimeter cell is used as an input to the algorithm that identifies significant calorimeter energy deposits and then identifies and measures jets, electrons and taus. Part of the work presented in this thesis was to update, monitor and study the noise calibration constants for the Tile Calorimeter to allow for precise reprocessing of the 2011 data and used for precision measurements. During these studies unexpected variations over time in the cell noise were observed. Further investigation led to discover that the tile noise filter, an algorithm used to mitigate the effect of coherent noise, was not behaving as expected in some situations significantly affecting approximately 5% of the cells in the TileCal.

This thesis presented two searches for BSM physics in ATLAS. The first one used the 2015 data corresponding to an integrated luminosity of 3.2 fb^{-1} where for the first time the monojet analysis was used to set limits on compressed light squark-neutralino models. Selections adapted to the specific characteristics of this signal were studied. It was shown that the generic approach provided by the global fit gives better sensitivity to this new signal than a single signal region with asymmetric jet and $E_{\text{T}}^{\text{miss}}$ cuts. Several jet veto criteria were studied since the BSM compressed SUSY scenario probed in this thesis is the most sensitive to this type of cuts. No significant excess in the data compared to the SM predictions has been found thus the results have been translated into model independent upper limits on the visible cross section with a 95% CL and into model dependent limits in terms of squark pair production with $\tilde{q} \rightarrow q + \tilde{\chi}_1^0$ ($q = u, d, c, s$). Squark masses up to 608 GeV for $m_{\tilde{q}} - m_{\tilde{\chi}_1^0} = 5 \text{ GeV}$

are excluded. The study of the 2015 dataset extends limits shown in Figure 7.3 that are provided by the more classical approach to SUSY searches.

The second analysis presented in this thesis uses the combined 2015 + 2016 dataset corresponding to an integrated luminosity of 36.1 fb^{-1} which was used to explore the presence of large extra dimensions in the ADD model scenario. A good agreement between data and Standard Model prediction is observed and no significant excess in data is seen thus the results are translated in 95% CL lower limits on the effective Planck scale M_D . Values of M_D of 7.7 and 4.8 TeV for two and six extra dimension hypotheses are excluded, significantly improving the results obtained by previous work. Future versions of this analysis will benefit from an increased dataset. Including 2017, the data collected by ATLAS during run 2 amounts to approximately 70 fb^{-1} , almost doubling the luminosity used in this thesis. However any significant improvement beyond the limits presented in this work will require very substantial amount of additional data. Assuming that the sensitivity to ADD models is limited by statistical uncertainties on the background, which is an optimistic assumption, we estimate from Eqs. (8.1.6) and (8.1.8) that a doubling of the integrated luminosity would only lead to a 10% or less improvement on the limits on M_D (due to the high exponent in the power dependence between M_D and the production cross section).

The sensitivity to squarks in the compressed squark-neutralino model is dominated by intermediate E_T^{miss} bins, where the systematic uncertainty on the background is more dominant than at high missing energies. Therefore significant improvements in squark-neutralino sensitivity will require better background calculations and reduced systematic uncertainties in addition to a simple increase in integrated luminosity.

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