

Search for New High-Mass Phenomena in Events with Two Muons using the ATLAS Detector at the Large Hadron Collider

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Abstract

Elementary particles and their interactions are extremely well modeled by the Standard Model of particle physics. However, experimental observations such as indirect detection of dark matter and theoretical problems such as the hierarchy of energy scales cannot be explained entirely by this theory. Many extensions of the Standard Model which solve these shortcomings predict the existence of new phenomena at high energies. In particular, there are numerous new resonance models, such as Grand Unified Theories, and contact interaction models leading to dimuon final states.

This dissertation presents a search for new high-mass phenomena in events with two muons using the ATLAS detector at the Large Hadron Collider. The search results are found to be consistent with the Standard Model prediction. Interpretations are carried out in the context of both resonant and non-resonant new physics models. In particular, lower limits on the mass of hypothetical Z' bosons are set between 4.0 TeV for the Z'_{SSM} model and 3.3 TeV for the Z'_ψ model, and lower limits on the contact interaction energy scale Λ are set between 18 TeV and 30 TeV, depending on the chiral structure of the contact interaction.

In addition to data analysis at the energy frontier, the performance of muon reconstruction and identification within the ATLAS experiment is detailed. More precisely, calculations of muon trigger efficiencies for high-momentum muons using events containing a leptonically decaying W boson and jets are presented. A new muon identification working point is also investigated.

Finally, as the ATLAS experiment enters its second long shutdown, the first layer of the endcap regions of the muon spectrometer will be replaced with the New Small Wheels (NSWs), in order to improve both the triggering and tracking capabilities of the ATLAS detector. One of the two main technologies used in the NSW is small-strip Thin Gap Chambers (sTGCs). Results of various test beam campaigns carried out at Fermilab and at CERN, aimed at characterizing sTGCs, are presented. Position resolution measurements of less than $50\mu\text{m}$ are obtained. Measurements using the latest electronics readout chain of the sTGC detectors under realistic conditions are also presented.

Lay Summary

Four known fundamental forces act on the matter in our universe: gravity, electromagnetism, the strong nuclear force, and the weak nuclear force. While gravity is described by Einstein's theory of General Relativity, the three other forces are described by the Standard Model of particle physics. These theories are very successful, and have been experimentally validated to high degrees of accuracy. However, they are incomplete. For example, they offer no explanation for the existence of dark matter, or the unification of forces. To solve these shortcomings, extensions of the Standard Model predicting new forces of nature have been posited. This dissertation analyzes the second dataset collected by the ATLAS experiment at the Large Hadron Collider, with a view to finding new forces of nature. While the search is found to be consistent with the Standard Model prediction, constraints placed on new physics scenarios contribute to the advancement of our knowledge of the universe.

Preface

The ATLAS collaboration comprises over 3000 scientists from 183 institutions worldwide, representing 38 countries. The work presented in this dissertation would not have been possible without the efforts of all members of the collaboration both from a scientific standpoint and a computing resources standpoint. In addition, the successful operation of the LHC is crucial to the realization of the physics experiments carried within the ATLAS collaboration at CERN. In many cases, the material presented in this dissertation relies on the work of many individuals.

Figures with no reference in the caption have been produced by myself, at times in collaboration with other researchers named below. Figures labeled “**ATLAS**”, “**ATLAS Online**”, or “**ATLAS Preliminary**”, are public results released by the ATLAS collaboration. Figures labeled “**ATLAS Work In Progress**” are not yet published but use either real or simulated ATLAS data.

The results presented in this dissertation led to the publications listed below. The first publication is the analysis of the full Run 2 dataset. Although no results from this publication are directly presented in this dissertation, I was involved with the analysis group and contributed to the optimization of the muon channel selection, the investigation of the highest invariant mass events, the development of the analysis framework, and the writing of internal documentation. In addition, I presented the latest results on behalf of the analysis team on several occasions during the internal ATLAS review process. The next two publications are related to the work described in Chapters 9 to 11. Throughout my doctoral studies, I acted as the main muon channel analyzer for the search discussed in Chapters 9 to 11. I played a major role in all aspects of the search, including the following:

- Development of the analysis framework
- Production of signal samples
- Optimization of the event selection
- Evaluation of the dominant systematic uncertainties

- Background estimation
- Signal estimation for numerous signal models
- Statistical analysis
- Investigation of the highest dimuon invariant mass events

In collaboration with D. Hayden and Y. Liu, I coordinated the cross-checks required between the different analyzers to ensure mutual consistency. I evaluated the performance of the ATLAS detector with respect to high- p_T muon identification by investigating the highest dimuon invariant mass events with M. Bugge and the various Muon Combined Performance group conveners G. Sciolla, M. Bellomo, S. Rosati, F. Sforza, and S. Zambito. I provided the inputs and ran the limit-setting analysis code with D. Hayden, W. Fedorko, and G. Artoni.

The last publication in the list includes the work presented in Chapter 8. I participated in the test beam campaigns carried out both at Fermilab and CERN. During the Fermilab test beam campaign, I

- Constructed and troubleshot the experimental setup in collaboration with the authors of the publication
- Led the small-strip Thin Gap Chamber standalone analysis with S. Viel
- Obtained residual distributions
- Addressed many referee comments during the review process

During the CERN test beam campaigns, in addition to preparing the experimental setup, I collected and analyzed data with L. Moleri, S. Sun, L. Guan, I. Ravinovich, N. Hod, D. Shaked-Renous, and S. Bressler. I also spent three months at the Weizmann Institute of Science in Isreal thanks to the Mitacs Globalink Research Award. Part of the work I carried out while at the Weizmann Institute of Science led to the results presented in Section 8.3.3 in collaboration with M. Birman.

ATLAS Collaboration *Search for high-mass dilepton resonances using 139 fb^{-1} of pp collision data collected at $\sqrt{s} = 13\text{ TeV}$ with the ATLAS detector*

Physics Letters B, 796, Pages 68-87, (2019)

ATLAS Collaboration *Search for new high-mass phenomena in the dilepton final state using 36 fb^{-1} of proton-proton collision data at $\sqrt{s} = 13\text{ TeV}$ with the ATLAS detector*
Journal of High Energy Physics, 10, 182, (2017)

ATLAS Collaboration *Search for high-mass new phenomena in the dilepton final state using proton-proton collisions at $\sqrt{s} = 13\text{ TeV}$ with the ATLAS detector*
Physics Letters B, 761, Pages 372–392, (2016)

sTGC Collaboration *Performance of a Full-Size Small-Strip Thin Gap Chamber Prototype for the ATLAS New Small Wheel Muon Upgrade*
Nuclear Instruments and Methods in Physics Research Section A, 817, Pages 85-92, (2016)

In addition to the publications listed above, two publications that include results from this dissertation are currently in preparation. The first describes muon identification using the full Run 2 dataset collected at the LHC, and includes the work described in Chapter 6. As the liaison between the ATLAS Exotics and Muon Combined Performance groups, I have had the pleasure of providing support for over 30 searches with muons in the final state. For example, I ensured the technical details regarding the latest Muon Combined Performance tools were propagated to the analyses adequately, and updated the relevant documentation web pages accordingly. I also contributed to the work carried out in the Muon Combined Performance group. Specifically, I used Alignment Effects On Track to improve the muon momentum uncertainty estimation, and optimized the high- p_T muon selection working point for the ATLAS collaboration, as detailed in Section 6.2. I also employed Alignment Effects On Track to categorize muon tracks in different parts of the alignment system of the Muon Spectrometer within ATLAS in order to study the momentum resolution of tracks in different detector regions. The investigation of the gradient Working Point described in Section 6.3 was performed with N. Hod.

Finally, the last publication in preparation relates to the performance of ATLAS muon triggers in Run 2. The efficiency measurements and Scale Factor calculations presented in Chapter 7 were carried out by myself and A. Held. In particular, we wrote the framework which provides high- p_T muon trigger Scale Factors for the entire ATLAS collaboration. I have served as one of the editors for the publication in preparation for the last year.

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Acronyms

AEOT Alignment Effect On Track. vii, 68, 78

ALICE A Large Ion Collider Experiment. 23

ASIC Application-Specific Integrated Circuit. 99

ATLAS A Toroidal LHC ApparatuS. v–vii, xii, xiv–xvi, xviii, xix, 1, 16, 18, 23, 25–29, 31, 33–35, 37–43, 46, 48, 50–53, 55, 56, 58, 60, 72, 85–88, 91–93, 95–97, 106, 113, 116, 117, 152

BDT Boosted Decision Tree. 49

BEE Barrel Endcap Extra. 40, 58, 65, 72

BSM Beyond the Standard Model. 1, 2, 16, 17, 56, 57, 74, 87, 114, 116, 129, 139, 140, 144, 145, 152

CB Combined. 52–54, 57, 65, 68, 71

CERN European Organization for Nuclear Research. v, vi, xiv, xx, 23, 24, 100, 106, 107, 113

CI Contact Interaction. xiv, xxv, 17, 20, 21, 144, 145, 150–152

CL Confidence Level. xiii, xxiv, xxv, 145–151

CMS Compact Muon Solenoid. 1, 18, 23

CSC Cathode Strip Chamber. xv, 37–39, 45, 53, 61, 65, 72

DAQ Data Acquisition. 99

DY Drell-Yan. xiv, xvii, xxi, 13, 14, 21, 69, 70, 116–120, 122, 135, 138

EM Electromagnetic. xv, 3, 5, 16, 28, 29, 33–37, 48, 50, 51, 88, 114

- EW** Electroweak. xxi, xxiv, 17, 118–120, 134, 138
- EWSB** Electroweak Symmetry Breaking. 8, 9, 15, 22
- FEB** Front-End Board. 99
- GRL** Good Run List. 88, 114
- GUT** Grand Unified Theory. 17, 18
- HL-LHC** High-Luminosity Large Hadron Collider. 95, 153
- HLT** High-Level Trigger. 86, 87
- IBL** Insertable B-Layer. 31, 32
- ID** Inner Detector. xv, xvi, xxiii, 28–35, 46, 47, 50–54, 57, 71, 86, 87, 90, 95, 129–131
- IP** Interaction Point. xv, 28, 29, 33, 47, 52
- ITk** Inner Tracker. 95
- L1** Level-1. xix, 86, 87, 96, 97
- LAr** Liquid Argon. xv, 34–37, 88, 95
- LHC** Large Hadron Collider. v, vii, xiv, 1, 2, 10, 20, 23, 25–28, 87, 88, 95, 139, 152, 153
- LHCb** Large Hadron Collider beauty. 23
- LO** Leading Order. xxi, 117, 119, 120
- LS2** Long Shutdown 2. 95, 152
- MC** Monte Carlo. 67, 87–89, 91, 114, 116, 118, 122, 138, 139
- MDT** Monitored Drift Tube. xv, xvi, 37–39, 44, 45, 53
- ME** Muon Extrapolated. 52, 53, 57, 71, 72
- MM** Micromegas. xix, 95, 97, 98

- MS** Muon Spectrometer. vii, xii, xv, xvii, xxiii, 28, 29, 37–40, 42–44, 47, 52–54, 57, 58, 64, 66, 67, 71, 72, 86, 87, 90, 95, 129–131
- NLO** Next-to-Leading Order. xiv, xxi, 11, 12, 116, 118–120
- NNLO** Next-to-Next-to-Leading Order. xxi, 116–119
- NSW** New Small Wheel. xix, 2, 95–98, 106, 113, 152
- PDF** Parton Distribution Function. xii–xiv, xxi, xxiv, 3, 11, 12, 116–120, 129, 132, 135–138
- PI** Photon-Induced. xxiv, 119, 132, 134, 138
- QCD** Quantum Chromodynamics. xiv, xxi, 12, 116, 118–120, 133
- QED** Quantum Electrodynamics. 5, 119
- RMS** Root Mean Square. xviii, xx, 80, 81, 83, 84, 106
- RoI** Region of Interest. 86, 87
- RPC** Resistive Plate Chamber. xv, xvi, 37–39, 42, 43, 95
- RS** Randall-Sundrum. 22
- SCT** Semiconductor Tracker. xv, 30–32, 46, 47, 57, 58
- SF** Scale Factor. vii, xii, xviii, xix, 87, 88, 91–94, 130, 152
- SM** Standard Model. xiv, xv, xxii, xxiii, 1, 3–10, 12–21, 29, 34, 55, 74, 88, 116, 121–128, 139, 140, 147, 152, 153
- SSM** Sequential Standard Model. 17, 18
- sTGC** small-strip Thin Gap Chamber. vi, xix, xx, 2, 95, 97–113, 152
- TDR** Technical Design Report. xvii, 66, 67
- TGC** Thin Gap Chamber. xv, 37–39, 43
- TRT** Transition Radiation Tracker. xv, 30–33, 46, 57, 58
- WP** Working Point. vii, xii, xvi–xviii, 49–51, 53, 54, 57–59, 61, 62, 64, 65, 67, 70, 71, 73–81, 84, 85, 88, 89, 91, 114, 130, 152

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Chapter 1

Introduction

We currently live in one of the most exciting eras in particle physics. At the forefront of high-energy particle physics are two multipurpose particle physics experiments located at the Large Hadron Collider (LHC) [1]: A Toroidal LHC ApparatuS (ATLAS) [2], and the Compact Muon Solenoid (CMS) [3]. With the discovery of the Higgs boson by the ATLAS and CMS experiments in 2012 [4, 5], the last missing particle of the Standard Model (SM) of particle physics [6–9] has been observed. As the LHC nearly doubles its centre-of-mass energy from 8 TeV to 13 TeV during its second data-taking period, so-called Run 2, the landscape of searches for new physical phenomena is ripe with possibilities that could bring to light a more thorough understanding of our universe. To this effect, searches for new high-mass phenomena are one of the most promising avenues. Progress in experimental particle physics often goes hand in hand with the discovery of new resonances. For example, the discoveries of the J/ψ meson [10], the Υ meson [11], and the Z boson [12], have all led to definite breakthroughs in this field.

This dissertation will focus on three main topics, all centred around one common thread: the fundamental particle known as the muon [13]. Chapter 2 presents the theoretical framework that underpins modern particle physics: the SM of particle physics. In addition to the quantum field theory formulation, both the successes and shortcomings of the SM are described. Chapter 3 details various theories that go Beyond the Standard Model (BSM) in attempts to solve its shortcomings. The ATLAS experiment and the LHC are presented in Chapter 4, with the experimental components such as the detector and accelerator technologies briefly outlined. The reconstructed objects used in various physics analyses are described in Chapter 5. Particular emphasis is given to high-momentum muons in Chapter 6, as these objects are of distinct importance for the main search described in this dissertation. The ATLAS trigger system is described in Chapter 7, and the performance of the muon trigger system is evaluated through trigger efficiency measurements. Chapter 8 gives an overview of the upgrade programme planned for the ATLAS detector in the coming

years, and details the New Small Wheel (NSW) upgrade in particular. One of the main technologies involved in the NSW upgrade, small-strip Thin Gap Chambers (sTGCs), is characterized through various test beam campaigns. A search for new high-mass phenomena in events with two muons in the final state is then presented in three distinct chapters: Chapter 9 describes the event selection and presents various kinematic distributions that compare data and simulation; Chapter 10 summarizes the systematic uncertainties considered; and Chapter 11 details the statistical methods used to search for BSM physics signals. The search uses 36.1 fb^{-1} of proton-proton collision data collected at the LHC. Finally, Chapter 12 concludes the dissertation, and presents an outlook for the future of searches for new high-mass phenomena in events with two muons in the final state.

Chapter 2

The Standard Model of Particle Physics

There are four known fundamental forces acting on the matter in our universe: gravity, electromagnetism, the strong nuclear force, and the weak nuclear force. The first of these, while being the most noticeable in our day-to-day lives, plays a relatively insignificant role at the subatomic level. Indeed, gravity is described by Einstein's theory of General Relativity, which deals mostly with objects on a macroscopic scale. The three other forces are best described by the SM of particle physics. In this chapter, the theoretical framework that makes up the SM is described in Section 2.1¹. Section 2.2 presents the parton model, and in particular Parton Distribution Functions (PDFs). To give an idea of the current state of particle physics, the impressive successes of the SM are presented in Section 2.3, and the limitations of the model are described in Section 2.4.

In the following, natural units where $\hbar = c = 1$ are used. In addition, the Einstein summation convention, where repeated indices imply a sum over all the values of the index, is used.

2.1 Theory

The SM is a very successful predictive theory, which explains the fundamental interactions of elementary particles in the universe, with the exception of gravity. This quantum field theory comprises the elementary particles depicted in Figure 2.1. The quarks and leptons, more generally known as fermions, are spin- $\frac{1}{2}$ particles that interact with each other through the exchange of integer-spin particles, the gauge bosons. For reasons yet unknown, the quarks and leptons exist in three so-called generations, with each generation displaying similar properties. The three main forces described by the SM are the strong nuclear force, the weak nuclear force, and the Electromagnetic (EM) force. These forces are mediated by the gluon (g), the weak

¹This section is loosely based on [14], and follows the conventions therein.

bosons ($W^\pm$ and Z^0), and the photon (γ), respectively. Each particle in the SM has a corresponding antiparticle.

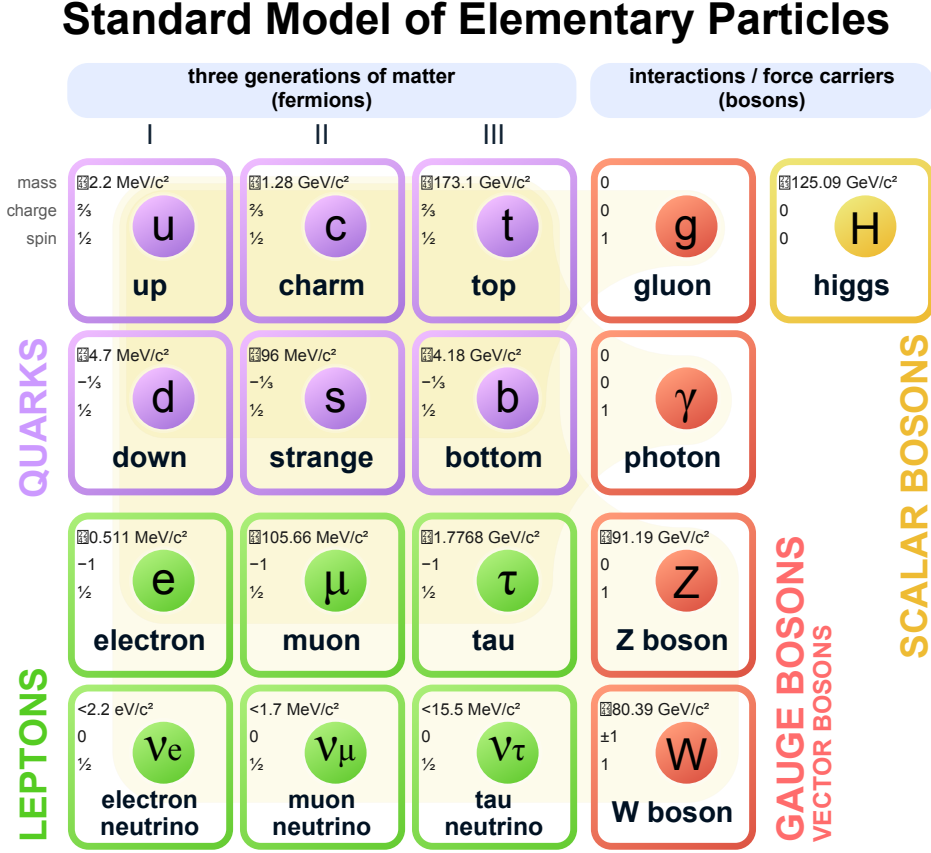


Figure 2.1: Particle content of the SM of particle physics [15].

The SM is a quantum field theory with local gauge invariance. In the group theoretical understanding of the SM, the gauge symmetry of the SM is expressed as the group

$$G_{SM} = SU(3)_C \times SU(2)_L \times U(1)_Y, \quad (2.1)$$

where the subscript C corresponds to the conserved colour charge of the $SU(3)$ group, the subscript L indicates that the $SU(2)$ group acts only on left-handed particles or right-handed antiparticles, and the subscript Y

2.1. Theory

corresponds to the conserved weak hypercharge. The conserved quantity related to the $SU(2)$ group is the weak isospin T_3 . The $SU(3)_C$ symmetry is unbroken and remains at low-energy scales. Quarks can carry one of three colour charges, while gluons carry one colour and one anticolour charge. As described in Section 2.1.3, the remaining gauge symmetry of the SM breaks down as the Higgs field reaches its ground state, generating gauge invariant masses for the weak gauge bosons in the process:

$$SU(2)_L \times U(1)_Y \rightarrow U(1)_Q. \quad (2.2)$$

This $U(1)_Q$ symmetry is the basis of Quantum Electrodynamics (QED), which is used to describe the EM interaction present in nature. The conserved quantity of the resulting $U(1)_Q$ symmetry is the usual electric charge, defined as

$$Q = T_3 + Y. \quad (2.3)$$

The gauge fields related to the $SU(3)_C$, $SU(2)_L$, and $U(1)_Y$ groups are defined as G_μ^α , W_μ^a , and B_μ , respectively. Recalling that the $SU(N)$ ($U(N)$) group has $N^2 - 1$ (N^2) generators, the SM gauge symmetry groups thus have 8 generators for $SU(3)_C$, 3 generators for $SU(2)_L$, and one generator for $U(1)_Y$. Hence, in the notation used throughout this chapter, α runs from 1 to 8, and a runs from 1 to 3.

The Lagrangian describing the fundamental particles in our universe can be written as three separate terms:

$$\mathcal{L}_{\text{SM}} = \mathcal{L}_{\text{gauge}} + \mathcal{L}_{\text{fermions}} + \mathcal{L}_{\text{Higgs}}. \quad (2.4)$$

Each of the three terms in Equation (2.4) will be described in the following sections: $\mathcal{L}_{\text{gauge}}$, $\mathcal{L}_{\text{fermions}}$, and $\mathcal{L}_{\text{Higgs}}$ will be covered in Sections 2.1.1 to 2.1.3, respectively.

2.1.1 Gauge Bosons

Gauge bosons transform under the adjoint representation of their gauge symmetry group. They mediate the SM forces between fermions, allowing the fermions to interact. The Lagrangian characterizing the gauge bosons is

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{4}G_{\mu\nu}^\alpha G^{\alpha\mu\nu} - \frac{1}{4}W_{\mu\nu}^a W^{a\mu\nu} - \frac{1}{4}B_{\mu\nu} B^{\mu\nu}, \quad (2.5)$$

2.1. Theory

where $G_{\mu\nu}^\alpha$, $W_{\mu\nu}^a$, and $B_{\mu\nu}$ are the field strength tensors for the $SU(3)_C$, $SU(2)_L$, and $U(1)_Y$ gauge fields, respectively. The field strength tensors are defined as follows:

$$G_{\mu\nu}^\alpha = \partial_\mu G_\nu^\alpha - \partial_\nu G_\mu^\alpha + g_s f^{\alpha\beta\gamma} G_\mu^\beta G_\nu^\gamma, \quad (2.6)$$

$$W_{\mu\nu}^a = \partial_\mu W_\nu^a - \partial_\nu W_\mu^a + g \epsilon^{abc} W_\mu^b W_\nu^c, \quad (2.7)$$

$$B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu, \quad (2.8)$$

where $f^{\alpha\beta\gamma}$ are the structure constants of the $SU(3)_C$ group and ϵ^{abc} are the structure constants of the $SU(2)_L$ group, i.e., the totally antisymmetric Levi-Civita symbol. The dimensionless coupling constants g_s and g are coupling constants that quantify the strength of the interaction related to the $SU(3)_C$ and $SU(2)_L$ symmetry groups, respectively. It should be noted that $\mathcal{L}_{\text{gauge}}$ contains no explicit mass terms for the gauge boson. As copious experimental evidence has been found implying that gauge bosons have mass [16], this is certainly a problem. The way to solve this issue is known as the Higgs mechanism and is described in Section 2.1.3.

2.1.2 Leptons and Quarks

Most of the observable matter in our universe is composed of fermions. More specifically, leptons and quarks make up the basic building blocks of nature. Fermions exist in three generations, which differ in mass. All three fermion generations interact with gauge bosons in exactly the same manner. For leptons, the electrically charged particles and their associated neutrinos come in the three generations labeled electron (e), muon (μ), and tau (τ). Leptons do not carry any colour charge, and as such transform as a colour singlet under the $SU(3)_C$ symmetry. For the quarks, the generational labels are denoted as follows: up (u) and down (d) for the first generation, charm (c) and strange (s) for the second generation, and top (t) and bottom (b) for the third generation. Quarks carry a colour charge and transform as triplets under the $SU(3)_C$ symmetry. The fermion fields present in the SM are the lepton fields L_L and e_R , and the quark fields Q_L , u_R , and d_R , where the subscripts L and R correspond to left-handed and right-handed chirality, respectively. The left-handed quark fields act as a doublet under the $SU(2)_L$ symmetry; positively-charged quarks have weak isospin $T_3 = \frac{1}{2}$, while negatively-charged quarks have weak isospin $T_3 = -\frac{1}{2}$. They also have weak hypercharge $Y = \frac{1}{6}$. The left-handed lepton fields also act as a doublet

2.1. Theory

under the $SU(2)_L$ symmetry; charged leptons have weak isospin $T_3 = -\frac{1}{2}$, while neutrinos have weak isospin $T_3 = +\frac{1}{2}$. The left-handed leptons have weak hypercharge $Y = -1/2$. All right-handed fermion fields act as singlets under the $SU(2)_L$ symmetry, and hence have weak isospin $T_3 = 0$. Table 2.1 summarizes the various properties of the fermion fields present in the SM for the first generation of fermions.

Fermion Type	Field	T_3	Y	Q	Colour
Leptons	$L_L = \begin{pmatrix} \nu_L \\ e_L \end{pmatrix}$	$+\frac{1}{2}$	$-\frac{1}{2}$	0	Singlet
		$-\frac{1}{2}$		-1	
	e_R	0	-1	-1	Singlet
Quarks	$Q_L = \begin{pmatrix} u_L \\ d_L \end{pmatrix}$	$+\frac{1}{2}$	$+\frac{1}{6}$	$+\frac{2}{3}$	Triplet
		$-\frac{1}{2}$		$-\frac{1}{3}$	
	u_R	0	$+\frac{2}{3}$	$+\frac{2}{3}$	Triplet
	d_R	0	$-\frac{1}{3}$	$-\frac{1}{3}$	Triplet

Table 2.1: Summary of fermion field properties for the first generation of fermions. Fermions in the second and third generation have identical properties.

The Lagrangian describing the fermions of the SM and their interactions is defined as

$$\mathcal{L}_{\text{fermions}} = \sum_{\text{fields}} \sum_{\text{gen}} \bar{\psi} i \gamma^\mu \mathcal{D}_\mu \psi, \quad (2.9)$$

where ψ denotes a generic fermion field and γ^μ represents the gamma matrices. The sum over generations (gen) implies summation over all three fermion generations for each fermion field. The $\mathcal{D}_\mu$ symbol in Equation (2.9) represents the most general form of a covariant derivative that can be written with regard to the symmetry group of the SM:

$$\mathcal{D}_\mu = \partial_\mu - ig_s G_\mu^\alpha t_c^\alpha - ig W_\mu^a t_L^a - ig' B_\mu Y, \quad (2.10)$$

where t_c^α are the generators of the $SU(3)_C$ group and t_L^a are the generators of the $SU(2)_L$ group. The dimensionless coupling constant g' represents the strength of the interaction related to the $U(1)_Y$ symmetry group. The last three terms in Equation (2.10) act only on the fields that transform non-trivially in the appropriate group representation. For example, the term proportional to g_s will not act on lepton fields, as they are colour singlets.

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Similarly to the expression for $\mathcal{L}_{\text{gauge}}$, no explicit mass terms for the leptons and quarks are found in $\mathcal{L}_{\text{fermions}}$, because they would not be gauge invariant. The solution to this discrepancy, i.e., that $\mathcal{L}_{\text{gauge}}$ and $\mathcal{L}_{\text{fermions}}$ predict massless particles but experimental evidence supports massive particles, is obtained via the Higgs mechanism described in Section 2.1.3.

2.1.3 The Higgs Mechanism

As alluded to above, gauge invariance would seem to imply the impossibility of any mass terms for gauge bosons, leptons, and quarks. However, as experimental evidence shows, these particles indeed carry mass. To address this issue, a mechanism¹ is introduced in order to break the underlying $SU(2)_L \times U(1)_Y$ symmetry as in Equation (2.2), and in the process give rise to mass terms in the SM Lagrangian for gauge bosons and fermions [17–22]. As the $SU(2)_L \times U(1)_Y$ symmetry is broken down into a $U(1)_Q$ symmetry, this process is called Electroweak Symmetry Breaking (EWSB). This is accomplished by introducing an $SU(2)_L$ doublet of complex scalar fields with weak hypercharge $Y = \frac{1}{2}$, which acts as an $SU(3)_C$ colour singlet, into the theory:

$$\phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix}. \quad (2.11)$$

The Lagrangian describing the Higgs sector of the SM is then

$$\mathcal{L}_{\text{Higgs}} = (\mathcal{D}_\mu \phi)^\dagger (\mathcal{D}_\mu \phi) - V(\phi) + \mathcal{L}_{\text{Yukawa}}, \quad (2.12)$$

where $\mathcal{L}_{\text{Yukawa}}$ describes the Yukawa interactions between the scalar Higgs field and the fermions, and the Higgs potential is defined as

$$V(\phi) = \mu^2 \phi^\dagger \phi + \lambda (\phi^\dagger \phi)^2. \quad (2.13)$$

If $\lambda = 0$, no symmetry breaking occurs. If λ is negative, the Higgs potential is unbounded from below, and no stable vacuum state exists. As this is considered unphysical, λ is assumed to be positive. If the value of μ^2 is positive, the potential has its minimum at $|\phi| = 0$. This corresponds to the Higgs potential before EWSB. However, if the value of μ^2 is negative, the Higgs potential minimum is not at $|\phi| = 0$. This corresponds to the Higgs potential after EWSB. Figure 2.2 compares the Higgs potential as a function of $|\phi|$ for the two possible cases discussed, i.e., when $\mu^2 > 0$ and $\mu^2 < 0$.

¹For brevity, this mechanism is denoted Higgs mechanism in this dissertation.

2.1. Theory

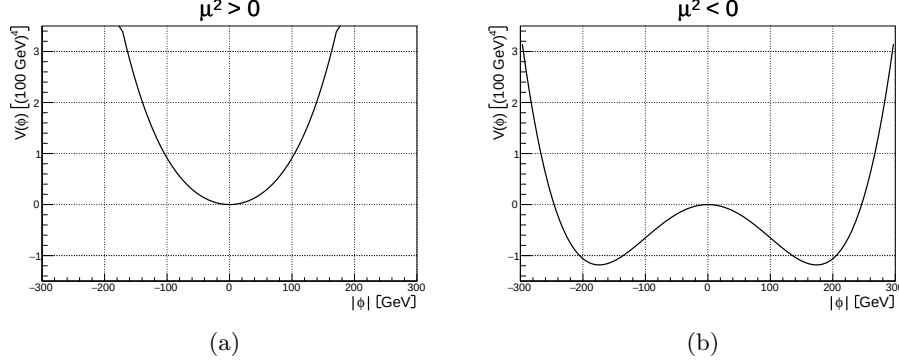


Figure 2.2: Higgs potential $V(\phi)$ as a function of $|\phi|$ for $\mu^2 > 0$ (left) and $\mu^2 < 0$ (right). The values used for the SM parameters of the Higgs potential are $\lambda \simeq 0.129$ and $|\mu^2| \simeq (88.4 \text{ GeV})^2$.

It is convenient to write the Higgs field in a gauge which minimizes the number of scalar degrees of freedom, i.e., the unitary gauge, as

$$\phi = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + h \end{pmatrix}, \quad (2.14)$$

where v represents the vacuum expectation value of the real Higgs field h . From Equation (2.13), it can be seen that the Higgs potential is minimized for $\phi^\dagger \phi = -\mu^2/2\lambda$. The vacuum expectation value is thus defined as

$$v = \sqrt{\frac{-\mu^2}{\lambda}}, \quad (2.15)$$

where μ^2 is negative. Inserting this definition into the first two terms of Equation (2.12) gives rise to mass terms for the massive gauge bosons in the SM Lagrangian. No mass term appears for the photon, and as such it remains massless. In addition, a mass term for a spin-0 scalar particle, the Higgs boson, is obtained. Finally, interaction terms between the massive bosons are also obtained. Once the $SU(2)_L \times U(1)_Y$ symmetry is broken, the usual electric charge becomes the conserved quantity, and is defined in Equation (2.3). After EWSB, the gauge fields of the $SU(2)_L \times U(1)_Y$ group mix to yield the physical states corresponding to the weak bosons ($W_\mu^\pm$ and Z_μ) and the photon (A_μ):

$$W_\mu^\pm = \frac{W_\mu^1 \mp iW_\mu^2}{\sqrt{2}}, \quad (2.16)$$

2.2. Parton Distribution Functions

$$Z_\mu = \cos \theta_W W_\mu^3 - \sin \theta_W B_\mu, \quad (2.17)$$

$$A_\mu = \sin \theta_W W_\mu^3 + \cos \theta_W B_\mu, \quad (2.18)$$

where the weak mixing angle θ_W is defined via

$$\cos \theta_W = \frac{g}{\sqrt{g^2 + g'^2}} \text{ and } \sin \theta_W = \frac{g'}{\sqrt{g^2 + g'^2}}, \quad (2.19)$$

and the coupling constants g and g' have already been defined. The masses of the bosons are given by

$$m_W^2 = \frac{g^2 v^2}{4}, \quad m_Z^2 = \frac{v^2(g^2 + g'^2)}{4}, \quad m_h^2 = 2\lambda v^2. \quad (2.20)$$

The last term in Equation (2.12) is what allows the SM fermions to have mass. The Lagrangian describing Yukawa interactions between the fermions and the Higgs field is defined as

$$\mathcal{L}_{\text{Yukawa}} = \sum_{\text{fields}} \sum_{\text{gen}} (y_i \bar{\psi}_R \phi^\dagger \psi_L + h.c.), \quad (2.21)$$

where ψ denotes a generic fermion field, the sum over fields includes all relevant fermion fields present in the SM, the sum over generations (gen) implies summation over all three fermion generations for each fermion field, and $h.c.$ denotes the Hermitian conjugate of the term before it. The Yukawa couplings y_i represent the strength of the interaction between the Higgs field and the relevant fermion. Inserting the Higgs field given in Equation (2.14) into Equation (2.21) gives rise to fermion mass terms and interaction terms between the fermions and the Higgs field.

2.2 Parton Distribution Functions

Protons are composite objects. As such, each proton-proton interaction at the LHC actually implies interactions between the underlying constituents of the proton, so-called partons. Each proton is composed of three valence quarks: two up quarks and one down quarks. In addition, quark-antiquark pairs are continuously produced and annihilated through interactions with gluons. Figure 2.3 illustrates the complex nature of the proton structure.

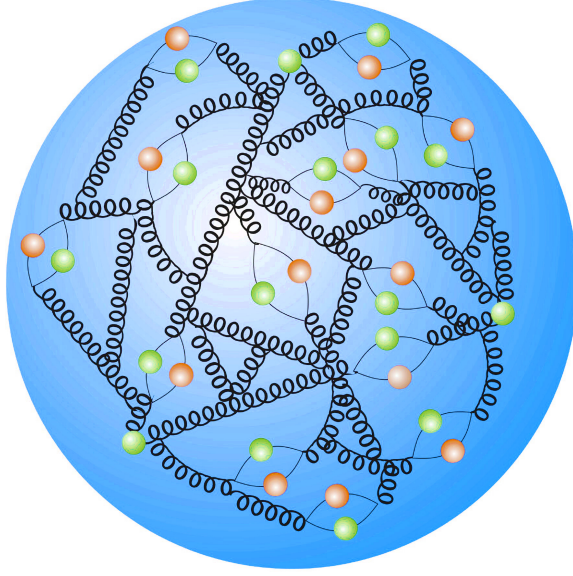


Figure 2.3: Illustration of the constituents of a proton, so-called partons. Quark-antiquark pairs are dynamically produced and annihilated through interactions with gluons [23].

Within the context of proton-proton collisions, the probability of a given parton within a proton to interact with another parton of a different proton is given by PDFs. PDFs are obtained through deep inelastic scattering experiments, and depend on the momentum transfer of the collision, Q , and the fraction of the proton momentum carried by each parton involved in the collision, x . Figure 2.4 depicts the PDFs of proton constituents at NLO as a function of x for $Q^2 = 10 \text{ GeV}^2$ and $Q^2 = 10^4 \text{ GeV}^2$ obtained by the MSTW collaboration [24].

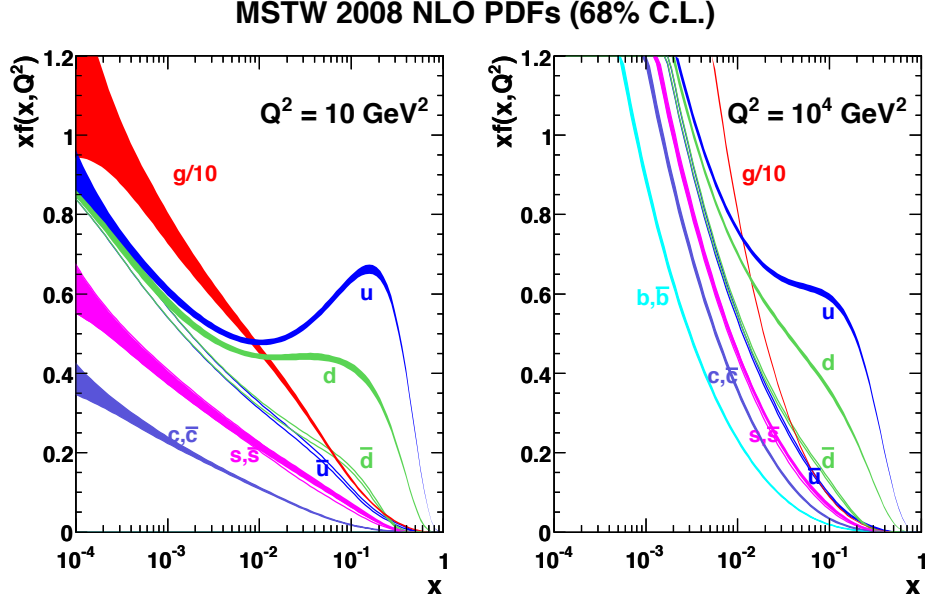


Figure 2.4: PDFs of proton constituents at NLO in QCD as a function of the momentum fraction x of the parton relative to the proton momentum, for momentum transfers $Q^2 = 10 \text{ GeV}^2$ (left) and $Q^2 = 10^4 \text{ GeV}^2$ (right) obtained by the MSTW collaboration [24].

2.3 Successes of the Standard Model

The SM has been experimentally verified to astounding degrees of precision. From the values of the coupling constants for the various forces to the production cross-sections of numerous processes, agreement between data and simulated predictions is impressive. Figure 2.5 gives a summary of various total production cross-section measurements performed by the ATLAS experiment at a centre-of-mass energy corresponding to $\sqrt{s} = 7 \text{ TeV}$, $\sqrt{s} = 8 \text{ TeV}$, and $\sqrt{s} = 13 \text{ TeV}$. The number of orders of magnitude spanned by the various measurements is a testament to the validity of the SM across a wide range of energy regimes.

2.3. Successes of the Standard Model

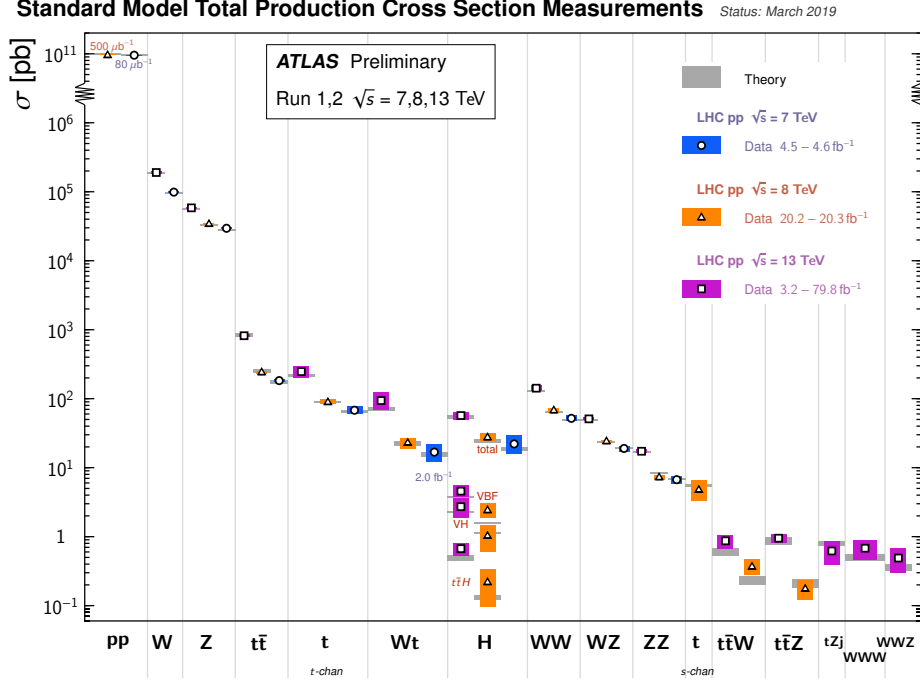


Figure 2.5: Summary of several SM total production cross-section measurements [25].

The third bin from the left in Figure 2.5 provides a parameter of particular relevance for the results presented in this dissertation: the production cross-section for the $pp \rightarrow \gamma/Z$ process. One of the most relevant SM processes to the search presented in this dissertation is the DY process, where two incoming quarks interact to produce an intermediate Z boson or photon, which subsequently decays to two outgoing muons. The DY process within the SM corresponds to the Feynman diagram depicted in Figure 2.6 at leading order.

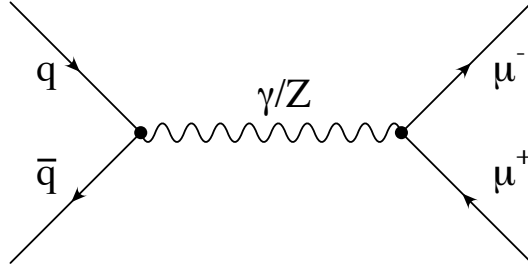


Figure 2.6: Feynman diagram of the SM DY process.

2.4 Limitations of the Standard Model

While the SM is an excellent theory in many particle physics regimes, there are certain experimental measurements that cannot be explained solely by the SM. For example, the existence of dark matter and massive neutrinos is not understood within the context of the SM. In addition, certain theoretical issues arise that forbid the SM to be extrapolated to arbitrarily high energies. In particular, a renormalizable quantum theory of gravity does not exist. Also, the vastly different energy scales present in the SM, while not necessarily a problem, is certainly an aesthetically displeasing issue at the very least. These shortcomings will be discussed in Sections 2.4.1 to 2.4.4

2.4.1 Quantum Gravity

One of the most flagrant shortcomings of the SM is that it simply does not describe the gravitational force. Indeed, gravitational effects at the quantum level are not considered in the majority of particle physics experiments. The gravitational force is extremely weak at the subatomic level. Consequently, it is usually negligible compared with the other forces described by the SM. Even though it is possible to write down a quantum theory of gravity, this theory is non-renormalizable, and thus is valid only up to a certain energy scale. This theory would break down, e.g., when considering extremely massive objects such as black holes, where gravitational effects at the quantum level become important. Hence, a renormalizable theory that reconciles General Relativity and quantum mechanics up to energy scales of the reduced Planck mass, $M_{\text{Pl}} \simeq 2.4 \times 10^{18}$ GeV, is not provided by the SM. A possible explanation for why gravitational effects are so small compared with the forces described by the SM is that gravity propagates in extra

spatial dimensions. Such a theory is briefly discussed in Section 3.3.

2.4.2 The Hierarchy Problem

As the Higgs boson is a scalar particle, its mass undergoes divergent quantum loop corrections. This is contrary to fermion or gauge boson masses, which are protected against divergent quantum loop corrections due to chiral symmetry and gauge invariance, respectively. In order to obtain a finite mass for the Higgs boson when including these corrections, counterterms must be added when computing the mass, in order to renormalize it. This is possible, but if the corrections required are larger than the bare Higgs boson mass without any corrections applied, this would imply that the theory is certainly not “natural”. In fact, to obtain a Higgs boson mass on the order of 125 GeV, which is the experimentally measured value, the counterterms required imply a cancellation of terms up to the 30th decimal place. This large amount of fine-tuning required for the Higgs boson mass seems to imply new physics, which would obviate the need for such fine-tuning of the corrected Higgs mass. That the quantum corrections to the Higgs boson mass are larger by several orders of magnitude than the bare mass is inexplicable, as one would expect the corrections to be of the same magnitude in a natural theory.

In addition to the Higgs mass issue, the fact that the particle content of the SM has masses spanning several orders of magnitude, while not necessarily an issue, is quite puzzling. Finally, even though the SM does not describe gravity, it is not unreasonable to expect that the four fundamental forces acting on the matter in our universe be approximately similar in strength; the fact that the weak nuclear force is $\sim 10^{24}$ times as strong as gravity is not understood with current models. Theoretically, it would be very satisfying if some underlying principle could explain these seemingly extremely large differences in a natural way.

2.4.3 Neutrino Masses

Neutrino oscillations have been experimentally observed [26–28]. These observations imply that neutrinos have non-zero, albeit very small, masses. As the SM does not include right-handed neutrinos, no mass terms are generated for these particles during EWSB. Hence, the SM predicts that neutrinos are explicitly massless. Neutrino oscillation measurements are thus certainly contradictory to SM predictions. In addition, the question of the different scales between e.g., the neutrino masses and their corresponding

lepton masses, which are orders of magnitude larger, remains unanswered. This is clearly an indication of BSM physics as it is currently defined.

2.4.4 Dark Matter and Dark Energy

Many experimental measurements point to the conclusion that the observable matter in our universe consists of only a small portion of the total matter present. The matter not visible through conventional measurements is named dark matter. Its presence is inferred by many different sources. These include gravitational lensing [29], rotational velocities of galaxies [30], and even the cosmic microwave background [31]. A possible explanation for dark matter is the existence of a new particle that interacts via the gravitational force, and potentially the weak nuclear interaction, but has neutral EM and colour charge, so does not interact via the EM and strong nuclear force. Such weakly interacting massive particles would account for the experimental measurements alluded to above, but have eluded detection thus far. Various BSM theories posit dark matter candidates through the existence of new particles, and experimentally detecting such new particles with the ATLAS experiment could help explain a vast array of cosmological problems. Finally, while the visible matter and dark matter comprise 4.9% and 26.8% of the energy density of the universe respectively, a staggering 68.3% of this energy density, named dark energy, is completely unaccounted for within the framework of the SM. The existence of dark energy is inferred by e.g., the measurements of the expanding universe carried out by the Planck collaboration [31]. Therefore, the SM is certainly not a complete description of nature.

Chapter 3

Beyond the Standard Model

The SM has been experimentally verified to extremely high degrees of precision, and provides an exceptional description of nature. However, it is known to be an effective theory in the sense that it is valid only in a low-energy regime, called the EW scale, and that it does not account for many observed experimental results. For example, it does not offer a satisfying explanation for dark matter. Essentially, the SM cannot be extrapolated to the high energy regime of the Planck scale, i.e., the scale at which quantum effects of gravity become strong, in a straightforward manner. Furthermore, from a theoretical point of view, the SM fails when it comes to explaining the hierarchy problem, the unification of forces, etc. Hence, it is clear that to fully understand and explain nature, a theoretical framework that goes beyond the SM is required. While high-mass resonances and Contact Interactions (CIs) do not offer a direct solution to the problems mentioned above, many BSM theories predict their existence. Therefore, finding such new phenomena would help validate these theories, which do offer solutions to the aforementioned problems. In this chapter, two types of BSM theories will be discussed: Section 3.1 will discuss models postulating new gauge symmetries, and Section 3.2 will discuss CI models. Finally, Section 3.3 briefly presents extra-dimensional models.

3.1 New Gauge Symmetries

Various BSM models predict the existence of new fundamental symmetries of nature. These symmetries would manifest themselves as high-mass resonances, i.e., peak-like excesses in the dimuon invariant mass distribution. Of particular interest for this dissertation are the Sequential Standard Model (SSM), Grand Unified Theories (GUTs), and minimal Z' models.

3.1.1 Sequential Standard Model

The SSM [32] is the simplest extension to the SM. It assumes the existence of heavier copies of the W and Z bosons present in the SM, described in

Chapter 2, which would have the same couplings to SM fermions as their counterparts. These new bosons are named W'_{SSM} and Z'_{SSM} , respectively. Contrary to the SM W and Z bosons, however, the couplings of the W'_{SSM} and the Z'_{SSM} bosons to each other are assumed to be zero. In addition, the couplings between the SM W and Z bosons and the new SSM gauge bosons are also assumed to be zero. While not very well motivated theoretically, the SSM is historically used as a benchmark model in order to compare results, e.g., between different experiments such as ATLAS and CMS [33, 34].

3.1.2 Grand Unified Theories

Other types of models that predict the existence of new gauge bosons are GUTs. GUTs often extend the gauge group of the SM, G_{SM} , to include additional $U(1)$ symmetries. These additional symmetries provide a unifying picture of gauge interactions, in which all fermions belong to the same multiplet of the group at the energy scale where the extended symmetries are unbroken. One interesting example of a GUT is the superstring-motivated E_6 model [35]. The E_6 gauge group can follow the Georgi-Glashow symmetry-breaking pattern [36]

$$E_6 \rightarrow SO(10) \times U(1)_\psi \rightarrow SU(5) \times U(1)_\chi \times U(1)_\psi, \quad (3.1)$$

where $SU(5)$ would break down into G_{SM} , and the additional $U(1)$ symmetries would manifest themselves via the production of Z' gauge bosons. These two neutral gauge bosons can mix with an angle θ_{E_6} , to yield a physical state defined by

$$Z'(\theta_{E_6}) = Z'_\chi \cos \theta_{E_6} + Z'_\psi \sin \theta_{E_6}, \quad (3.2)$$

where $\theta_{E_6} \in [-\pi, \pi]$. Table 3.1 lists the most commonly used values of θ_{E_6} and their corresponding physical states, according to the convention used in [32].

3.1. New Gauge Symmetries

Model Name	Z' State	θ_{E_6} Value
ψ Model	Z'_ψ	$\pi/2$
η Model	Z'_η	$\pi - \arctan \sqrt{5/3}$
Neutral N Model	Z'_N	$\arctan \sqrt{15}$
Inert Model	Z'_I	$\arctan \sqrt{3/5}$
Secluded Sector Model	Z'_S	$\arctan(\sqrt{15}/9)$
χ Model	Z'_χ	0

Table 3.1: Summary of commonly motivated Z' models.

Another possible symmetry-breaking pattern is given by the Pati-Salam model [37]:

$$\begin{aligned}
E_6 &\rightarrow SO(10) \times U(1)_\psi \rightarrow SU(4)_C \times SU(2)_L \times SU(2)_R \\
&\rightarrow SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \quad (3.3) \\
&\rightarrow SU(3)_C \times SU(2)_L \times U(1)_Y,
\end{aligned}$$

where the last line in Equation (3.3) corresponds to G_{SM} . The left-right symmetry is made explicit in this model due to the $SU(2)_L \times SU(2)_R$ structure. The model predicts a Z'_{3R} boson coming from the $SU(2)_R$ group or a Z'_{B-L} boson coming from the breaking of $SU(4)_C$ into $SU(3)_C \times U(1)_{B-L}$, where $3R$ stands for the right-handed third component of weak isospin, and B (L) stands for the baryon (lepton) number, with $(B - L)$ representing the conserved quantum number [38–40]. Both the Z'_{3R} and the Z'_{B-L} arise in the minimal Z' models discussed in Section 3.1.3. The models above all assume that the coupling of the new Z' gauge bosons to SM W and Z gauge bosons is zero.

3.1.3 Minimal Z' Models

An alternative approach to parameterising new gauge bosons is minimal Z' models [41]. In these models, the new resonances are characterized by three parameters: g_{BL} , g_Y , and the mass of the hypothetical Z' boson. Both g_{BL} and g_Y are effective coupling constants representing the coupling of a new Z' boson to the $(B - L)$ current and weak hypercharge Y , respectively. To simplify the notation, it is convenient to define

$$\tilde{g}_{BL} \equiv g_{BL}/g_Z \text{ and } \tilde{g}_Y \equiv g_Y/g_Z, \quad (3.4)$$

3.2. Contact Interactions

where $g_Z = 2m_Z/v$, and m_Z and $v = 246$ GeV are the SM values corresponding to the mass of the Z boson and the vacuum expectation value of the Higgs boson, respectively. From these new quantities, two independent parameters can be defined, in order to specify completely a particular minimal Z' model: γ' and θ_{Min} . These two parameters are defined by the relations

$$\tilde{g}_{BL} = \gamma' \cos \theta_{\text{Min}} \text{ and } \tilde{g}_Y = \gamma' \sin \theta_{\text{Min}}. \quad (3.5)$$

In essence, θ_{Min} corresponds to the mixing between the generators of the $(B - L)$ and weak hypercharge Y gauge groups, while γ' measures the strength of the Z' boson coupling relative to that of the SM Z boson. Table 3.2 summarizes the values of γ' and θ_{Min} for three specific minimal Z' models, which each have a specific Z' boson: Z'_{B-L} , Z'_χ , and Z'_{3R} .

	Z'_{B-L}	Z'_χ	Z'_{3R}
γ'	$\sqrt{\frac{5}{8}} \sin \theta_W$	$\sqrt{\frac{41}{24}} \sin \theta_W$	$\sqrt{\frac{25}{12}} \sin \theta_W$
$\cos \theta_{\text{Min}}$	1	$\sqrt{\frac{25}{41}}$	$\frac{1}{\sqrt{5}}$
$\sin \theta_{\text{Min}}$	0	$-\sqrt{\frac{16}{41}}$	$-\frac{2}{\sqrt{5}}$

Table 3.2: Summary of γ' and θ_{Min} values for three specific minimal Z' models: Z'_{B-L} , Z'_χ , and Z'_{3R} .

3.2 Contact Interactions

New phenomena in the dimuon invariant mass spectrum can also be observed by considering new interactions that occur at much higher energies than the energies in reach at the LHC. Such phenomena could, e.g., be four-fermion interactions, which would indicate that the interacting fermions, in this case quarks and leptons, are composite objects. A particularly interesting new physics scenario would be that of CIs. CIs, depicted in Figure 3.1, would manifest themselves as broad excesses in the dimuon invariant mass distribution. The energy scale at which the CI occurs, denoted Λ , represents the binding energy required to keep the fermion constituents together.

3.2. Contact Interactions

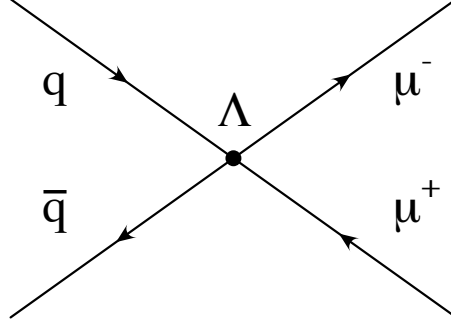


Figure 3.1: Feynmann diagram representing a CI with energy scale Λ .

The Lagrangian describing such CIs is given by [42, 43]

$$\begin{aligned} \mathcal{L}_{\text{CI}} = \frac{4\pi}{\Lambda^2} & [\eta_{\text{LL}}(\bar{q}_L \gamma_\mu q_L)(\bar{\ell}_L \gamma^\mu \ell_L) + \eta_{\text{RR}}(\bar{q}_R \gamma_\mu q_R)(\bar{\ell}_R \gamma^\mu \ell_R) \\ & + \eta_{\text{LR}}(\bar{q}_L \gamma_\mu q_L)(\bar{\ell}_R \gamma^\mu \ell_R) + \eta_{\text{RL}}(\bar{q}_R \gamma_\mu q_R)(\bar{\ell}_L \gamma^\mu \ell_L)], \end{aligned} \quad (3.6)$$

where Λ is the CI scale, and $q_{L,R}$ and $\ell_{L,R}$ are left-handed and right-handed quark and lepton fields, respectively. The symbol γ_μ represents the gamma matrices, and the parameters η_{ij} , where $i, j \in [L, R]$ with L (R) referring to left (right), define the chiral structure of the new CI. The sign of η_{ij} determines whether the interference between the SM DY process and the CI process is constructive ($\eta_{ij} = -1$) or destructive ($\eta_{ij} = +1$). By setting all but one of the η_{ij} parameters in Equation (3.6) to zero, and setting the remaining η_{ij} parameter to ± 1 , different chiral structures of the CI are investigated. The total cross-section for the process $q\bar{q} \rightarrow \mu\bar{\mu}$ can then be written as

$$\sigma_{\text{tot}}(m_{\mu\mu}) = \sigma_{\text{DY}}(m_{\mu\mu}) - \eta_{ij} \frac{F_I}{\Lambda^2} + \frac{F_C}{\Lambda^4}, \quad (3.7)$$

where $\sigma_{\text{DY}}(m_{\mu\mu})$ is the cross-section of the SM DY process shown in Figure 2.6 and the last two terms represent the interference between the SM DY process and the CI process, and the pure CI process, respectively. F_I and F_C are both functions of the differential cross-section with respect to $m_{\mu\mu}$, and do not depend on Λ .

3.3 Extra-Dimensional Models

A different approach to solving the hierarchy problem is to introduce extra spacetime dimensions. New extra dimensions would reconcile the very different energy scales of the Planck scale and the EWSB scale, by allowing gravity to propagate in these extra dimensions. By propagating in these extra dimensions, gravity would thus be much weaker. One relevant example of these types of models is the Randall-Sundrum (RS) model [44]. The RS model posits the metric

$$ds^2 = e^{-2kr_c\phi} \eta_{\mu\nu} dx^\mu dx^\nu + r_c^2 d\phi^2, \quad (3.8)$$

where k is a scale on the order of the Planck scale, x^μ are the usual spacetime coordinates, $\phi \in [0, \pi]$ is the angular coordinate parameterising the extra dimension, and r_c , known as the compactification radius, sets the size of the finite interval of ϕ . In other words, r_c sets the size of the extra dimension. Though outside the scope of this dissertation, it is possible to solve the Einstein field equations for this metric and derive predictions for the mass spectrum and couplings of the graviton Kaluza-Klein modes. In particular, a physical implication of this theory is the appearance of Kaluza-Klein graviton excitations with masses given by [45]

$$m_n = kx_n e^{-kr_c\pi}, \quad (3.9)$$

where x_n are roots of the Bessel function of order 1, i.e., $x_0 = 0, x_1 = 3.8317, x_2 = 7.0156, \dots$. Hence, the mass hierarchy between the physical mass m_n and the Planck scale is generated by an exponential function of the compactification radius. These excitations could give rise to high-mass resonances at the TeV scale with various final states, including two muons.

Chapter 4

The ATLAS Detector at the Large Hadron Collider

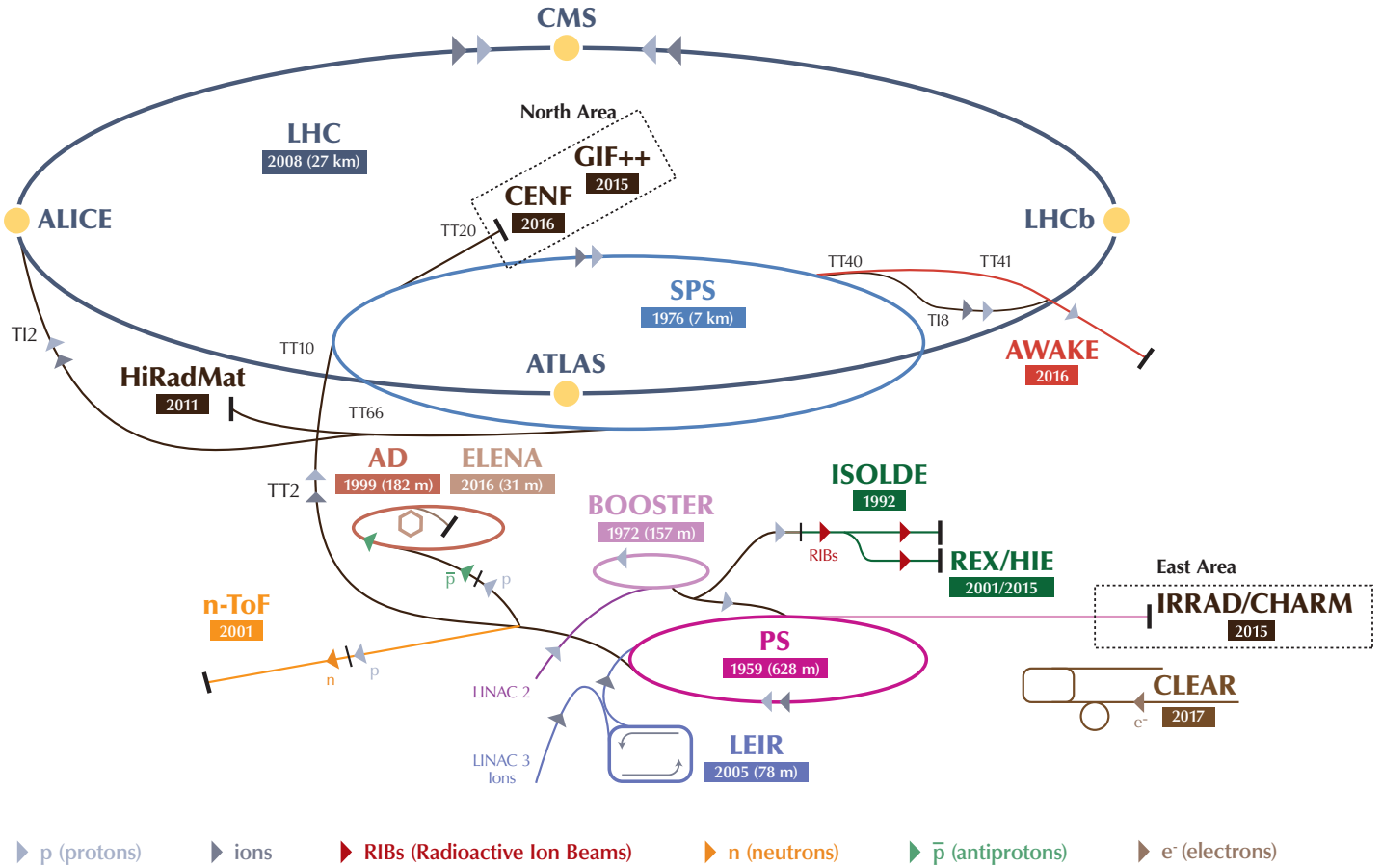
The results presented in this dissertation use data collected by the ATLAS detector at the LHC in the years 2015-2018, during the so-called Run 2 of the LHC. In this chapter, Section 4.1 gives an overview of the LHC accelerator complex, and Section 4.2 describes the ATLAS detector in detail.

4.1 The Large Hadron Collider

The LHC is a proton-proton collider located at the European Organization for Nuclear Research (CERN), on the Franco-Swiss border near Geneva, Switzerland. Two beams of protons are accelerated in opposite directions around a 27 km ring located 100 m underground, before colliding at one of four interaction points. Each of these interaction points is surrounded by a multi-purpose particle detector: ATLAS, CMS, A Large Ion Collider Experiment (ALICE) [46], or Large Hadron Collider beauty (LHCb) [47]. The beams are steered along the ring using an 8.3 T dipole field formed by superconducting magnets made of Nb-Ti. In order to bring the dipole magnets to their superconducting state, liquid helium is used to cool the dipole magnets down to 1.9 K. The LHC's design centre-of-mass energy is $\sqrt{s} = 14$ TeV. During Run 2, the LHC operated with a centre-of-mass energy of $\sqrt{s} = 13$ TeV.

Figure 4.1 shows the CERN accelerator complex, with the LHC being the central component. Initially accelerated by a linear accelerator to 50 MeV, a series of smaller accelerators are subsequently used to inject bunches of protons into the LHC: the proton synchrotron booster accelerates the protons to 1.4 GeV; the proton synchrotron accelerates them to 25 GeV; the super proton synchrotron accelerates them to 450 GeV. Finally, the proton bunches are injected into the LHC, and accelerated to their maximum energy of 6.5 TeV. The spacing between bunches of protons, the so-called bunch spacing, is 25 ns. This equates to a bunch crossing rate of 40 MHz.

The CERN accelerator complex *Complexe des accélérateurs du CERN*



LHC - Large Hadron Collider // SPS - Super Proton Synchrotron // PS - Proton Synchrotron // AD - Antiproton Decelerator // CLEAR - CERN Linear Electron Accelerator for Research // AWAKE - Advanced WAKEfield Experiment // ISOLDE - Isotope Separator OnLine // REX/HIE - Radioactive Experiment/High Intensity and Energy ISOLDE // LEIR - Low Energy Ion Ring // LINAC - LINear ACcelerator // n-ToF - Neutrons Time Of Flight // HiRadMat - High-Radiation to Materials // CHARM - Cern High energy AcceleRator Mixed field facility // IRRAD - proton IRRADIation facility // GIF++ - Gamma Irradiation Facility // CENF - CERN Neutrino platForm

Figure 4.1: The CERN accelerator complex [48].

When proton bunches collide at the centre of the ATLAS detector, the primary interaction of interest is the so-called “hard scatter”, i.e., the interaction with the largest momentum transfer. A hard scatter takes place when a parton from one proton interacts with a parton from another proton. In addition to the hard scatter, other parton-parton interactions take place during each bunch crossing. Pile-up is defined as the mean number of interactions per crossing, μ . Figure 4.2 shows the pile-up distributions for the Run 2 data-taking period.

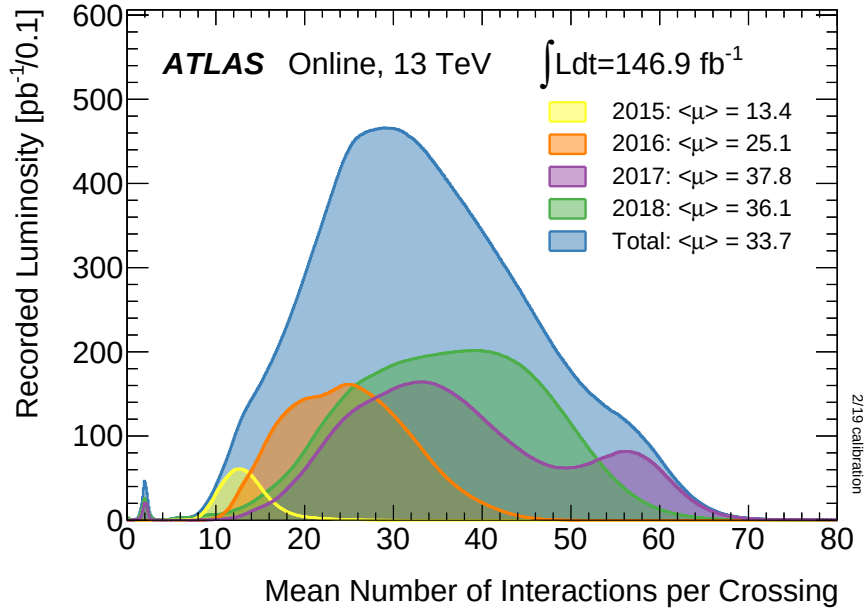


Figure 4.2: Mean number of interactions per LHC bunch crossing [49].

One of the key concepts for the LHC is that of luminosity. The number of expected events N for a particular process characterized by a cross-section σ can be obtained by considering the instantaneous luminosity L of the LHC:

$$N = \sigma \times \int dt L. \quad (4.1)$$

The instantaneous luminosity is measured in $\text{cm}^{-2}\text{s}^{-1}$, and can be expressed as [50, 51]

$$L = \frac{n_b f_r n_1 n_2}{2\pi \Sigma_x \Sigma_y}, \quad (4.2)$$

where n_b is the number of bunch pairs colliding per revolution, f_r is the revolution frequency, n_1 and n_2 correspond to the number of protons per bunch in each proton beam, and Σ_x and Σ_y characterize the horizontal and vertical convolved beam widths. An integrated luminosity of 1 fb^{-1} corresponds to approximately 10^{12} proton-proton collisions. During the first run of data-taking at the LHC, Run 1, the ATLAS experiment recorded an integrated luminosity totaling 25 fb^{-1} with a centre-of-mass energy of $\sqrt{s} = 7 \text{ TeV}$ and $\sqrt{s} = 8 \text{ TeV}$ and a bunch crossing rate of 20 MHz. During the second run of data-taking at the LHC, Run 2, the centre-of-mass energy increased to $\sqrt{s} = 13 \text{ TeV}$. The bunch crossing rate initially remained at 20 MHz, but was increased to the design value of 40 MHz near the beginning of July 2015. A total of 36.1 fb^{-1} of data have been collected in 2015-2016, and the entire Run 2 dataset totals 139 fb^{-1} . Figure 4.3 shows the evolution of the luminosity delivered to the ATLAS detector over time. Figure 4.4 shows the evolution of the total integrated luminosity recorded by the ATLAS detector during Run 2.

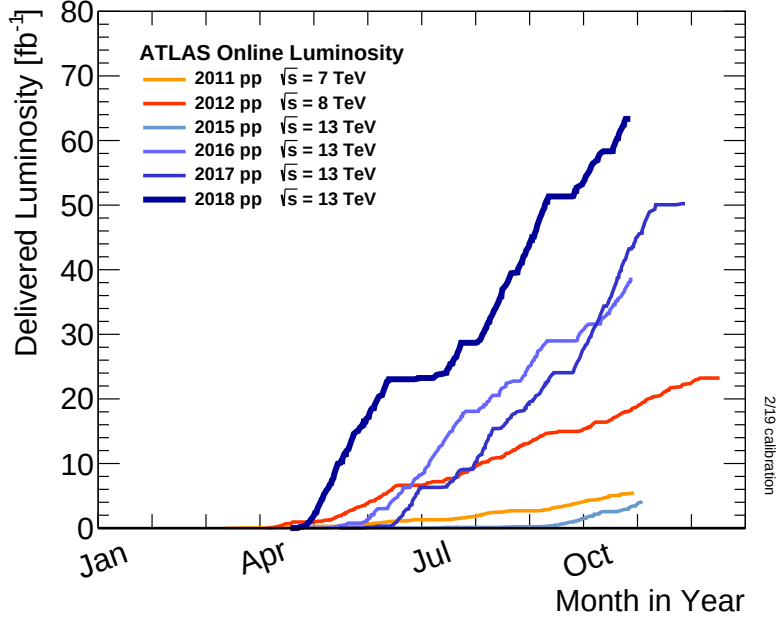


Figure 4.3: Luminosity delivered to the ATLAS detector as a function of time [49].

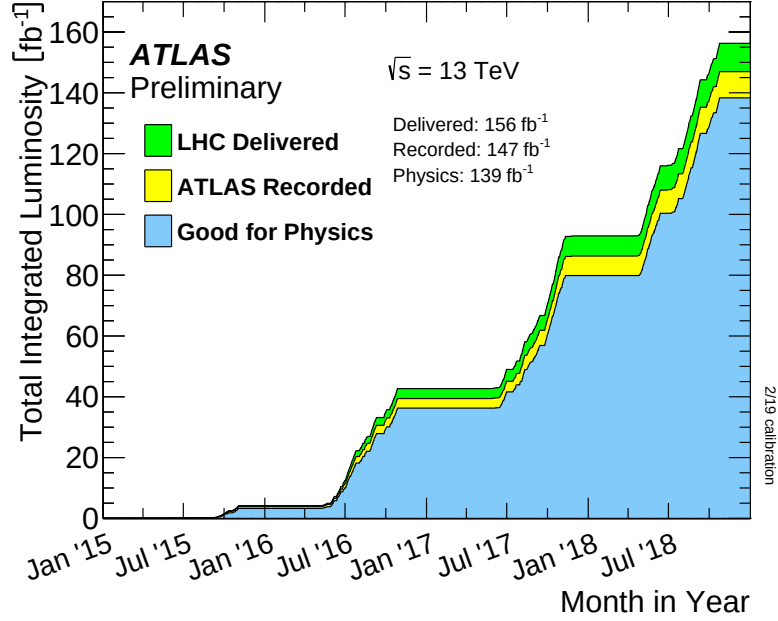


Figure 4.4: Total integrated luminosity recorded by the ATLAS detector as a function of time [49].

In summary, two proton beams are accelerated in opposite directions to nearly the speed of light before colliding at the centre of multi-purpose particle detectors. In particular, the detector used for the search described in this dissertation is the ATLAS detector.

4.2 The ATLAS Detector: Overview

ATLAS is a multi-purpose particle detector used at the LHC. A detailed schematic of the ATLAS detector is shown in Figure 4.5.

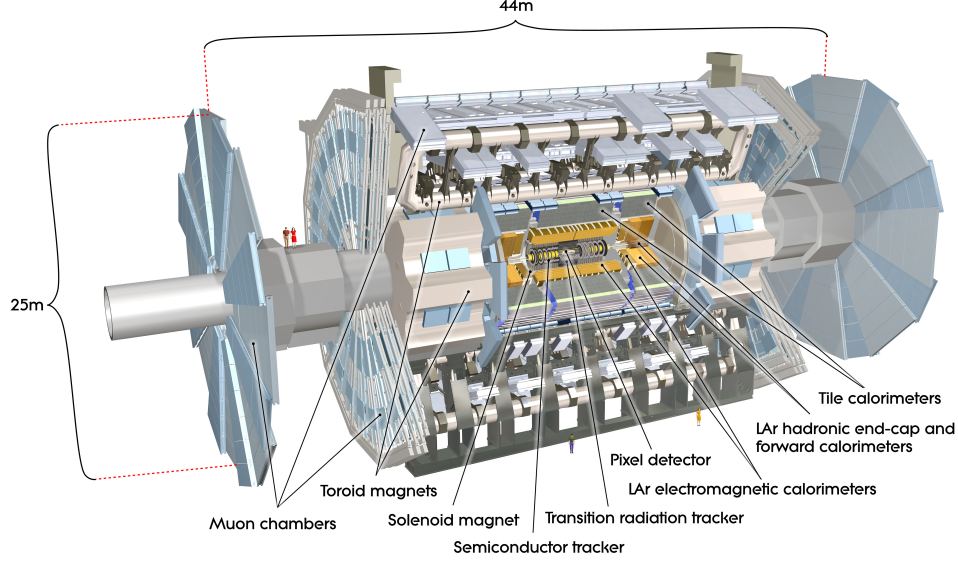


Figure 4.5: Schematic diagram of the ATLAS detector [2].

The coordinate system used in ATLAS is a right-handed coordinate system with its origin at the nominal Interaction Point (IP) in the centre of the detector and the z-axis along the beam pipe. The x-axis points from the IP to the centre of the LHC ring, and the y-axis points upwards. Cylindrical coordinates (r, ϕ) are used in the transverse plane, ϕ being the azimuthal angle around the z-axis. A convenient quantity to note is the pseudorapidity,

$$\eta = -\ln(\tan(\theta/2)), \quad (4.3)$$

where θ is the polar angle. The angular distance in the ATLAS detector is measured in units of $\Delta R = \sqrt{(\Delta\eta)^2 + (\Delta\phi)^2}$.

The ATLAS detector is composed of three main parts: the Inner Detector (ID) described in Section 4.3, the EM and hadronic calorimeters described in Section 4.4, and the Muon Spectrometer (MS) described in Section 4.5. Figure 4.6 illustrates a transverse slice (in the $R-\phi$ plane) of the detector and the various detector technologies used to differentiate between the different types of particles present in the detector during a collision event. Two magnetic fields are present in the ATLAS detector: a 2 T solenoid field in the ID bends charged particles in the ϕ direction, and a toroid field starting after the hadronic calorimeter and covering the entire MS bends charged

4.2. The ATLAS Detector: Overview

particles in the η direction. All three toroid fields, i.e., one in the barrel region and two in the endcap regions of the detector, are formed by 8 coils assembled radially and symmetrically around the beam axis.

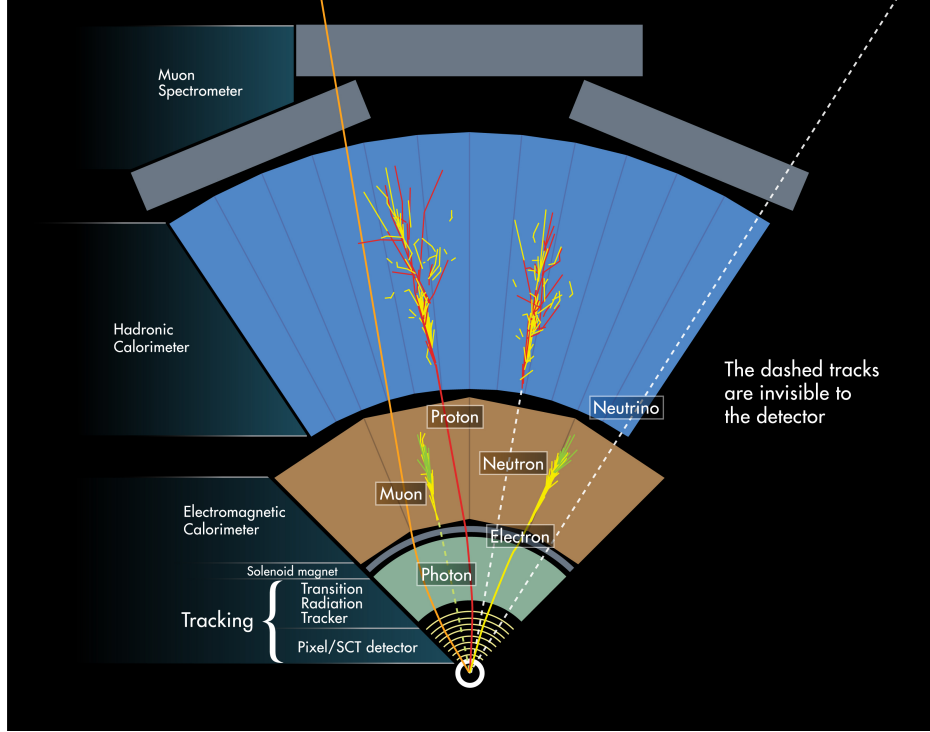


Figure 4.6: Slice in the $R - \phi$ plane of the ATLAS detector. Starting from the IP, a particle will traverse the ID, the EM calorimeter, the hadronic calorimeter, and the MS. Different particles are reconstructed and identified using all relevant sub-detectors. SM neutrinos are invisible to the ATLAS detector [52].

Different particles are reconstructed and identified using all relevant sub-detectors in ATLAS. Electrons are reconstructed from ID tracks and combined with energy deposits in the EM calorimeters. Muons are reconstructed from tracks in the ID, calorimeter, or MS, depending on the quality of the muon. Quarks and gluons produced from proton-proton interactions will hadronize into colour-neutral states producing a shower of particles, collectively called a jet. Jets are reconstructed using energy deposits in the calorimeters. A detailed description of object reconstruction in the ATLAS detector is presented in Chapter 5.

4.3 Inner Detector

The ID is immersed in a 2 T solenoid magnetic field and is composed of three separate detector technologies: silicon pixel trackers described in Section 4.3.1, silicon micro-strip Semiconductor Trackers (SCTs) described in Section 4.3.2, and Transition Radiation Trackers (TRTs) described in Section 4.3.3. Arranged in concentric cylinders in the barrel and in disks perpendicular to the beam axis in the endcaps, the detectors are used for the tracking of charged particles. These allow for electron and pion identification in the range $|\eta| < 2.5$. Figures 4.7 and 4.8 show the transverse and longitudinal views of the ID, respectively.

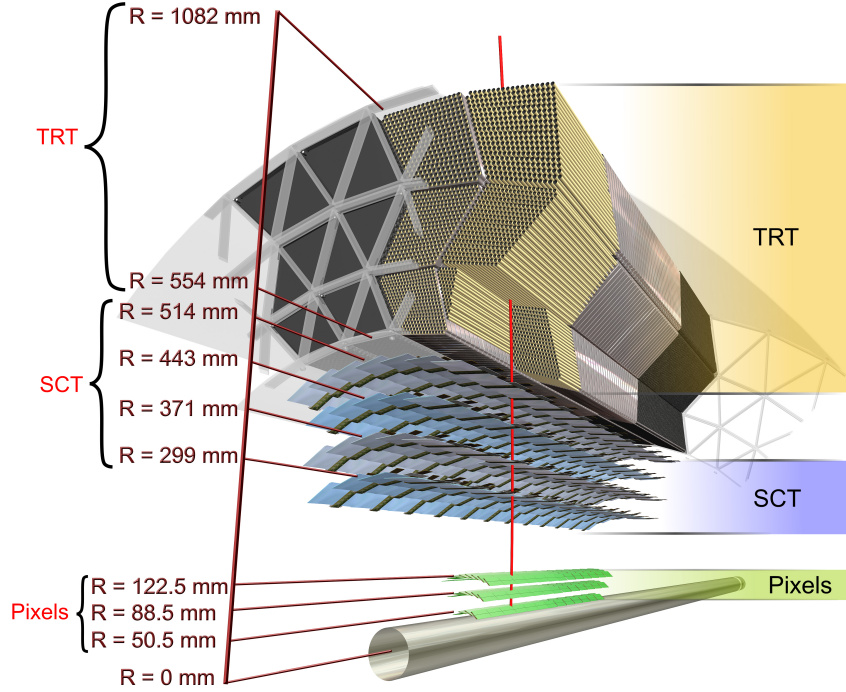


Figure 4.7: The Inner Detector as seen in the transverse $R - \phi$ plane. Particles traverse the beam pipe, three cylindrical silicon pixel layers, four SCT layers, and approximately 36 TRT straw tubes [2].

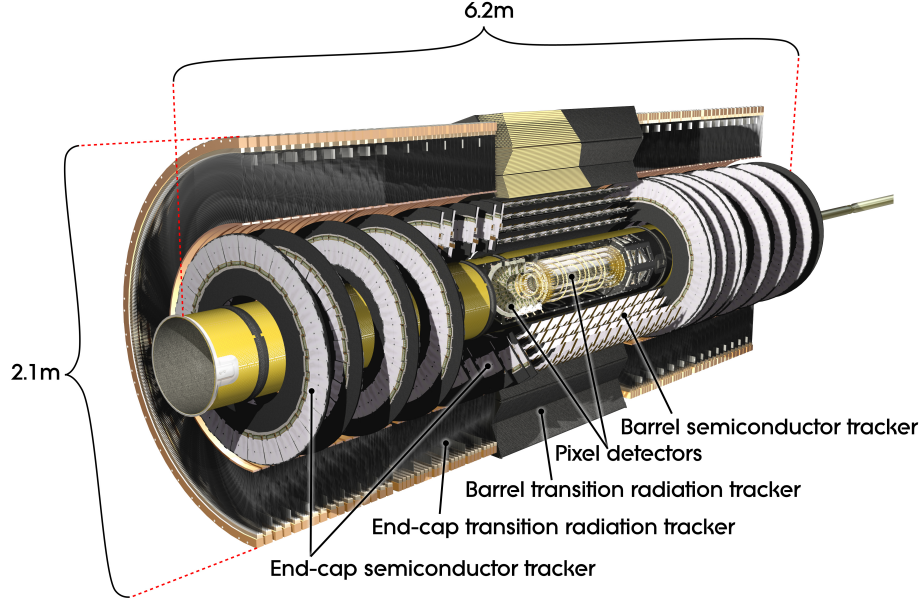


Figure 4.8: The Inner Detector as seen in the longitudinal plane. Immersed in a 2 T solenoid magnetic field, the main components are the pixel and SCT, covering $|\eta| < 2.5$, and the TRT, covering $|\eta| < 2.0$ [2].

4.3.1 Pixel Detector

The pixel detector is the innermost part of the ATLAS detector, and allows for precise tracking using 1774 identical silicon pixel sensor modules. Each module has 46080 readout channels, equating to approximately 80 million pixel readout channels throughout the pixel detector. Approximately 90% of the pixels have dimensions of $50\text{ }\mu\text{m}$ by $400\text{ }\mu\text{m}$, while the remaining 10% have dimensions $50\text{ }\mu\text{m}$ by $600\text{ }\mu\text{m}$. All pixels have a thickness of $250\text{ }\mu\text{m}$. The intrinsic resolution of the pixel detector in the barrel region is $10\text{ }\mu\text{m}$ in $R - \phi$ and $115\text{ }\mu\text{m}$ in z , while in the endcap regions it is $10\text{ }\mu\text{m}$ in $R - \phi$ and $115\text{ }\mu\text{m}$ R . There are three concentric cylindrical layers of pixels in the barrel region, and three disks perpendicular to the beam in each of the endcap regions. A typical track will cross three pixel layers. A new layer of pixel detectors, the Insertable B-Layer (IBL), was commissioned for Run 2. Located only in the barrel region, the IBL consists of an additional cylindrical layer of pixel detectors with a radius of 3.2 cm. The IBL provides a fourth layer of pixel detectors in the ID region and improves the precision

of the ID tracking capabilities. The pixel dimensions for the IBL are $50\text{ }\mu\text{m}$ by $250\text{ }\mu\text{m}$.

4.3.2 Semiconductor Tracker

The SCT consists primarily of two-sided silicon microstrip detector modules. The dimensions of one pixel microstrip is $80\text{ }\mu\text{m}$ by 12.8 cm . Each module consists of two microstrips glued back-to-back at a 40 mrad stereo angle. The intrinsic resolution in the barrel is $17\text{ }\mu\text{m}$ in $R - \phi$ and $580\text{ }\mu\text{m}$ in z , and in the endcap regions the intrinsic resolution is $17\text{ }\mu\text{m}$ in $R - \phi$ and $580\text{ }\mu\text{m}$ in R . The barrel region consists of 4 cylindrical layers of modules, with the length of the strips parallel to the beam axis. The endcap regions have 9 planar discs per side, with the length of the strips extending radially from the beam axis. Cooled to between -5°C and -10°C , the SCT has over 6 million readout channels in total. A typical track will cross four SCT layers, resulting in 8 SCT hits for a typical track.

4.3.3 Transition Radiation Tracker

The last section of the ID to be traversed by particles is the TRT. The basic detector element of the TRT consists of a straw tube with 4 mm diameter. The tube has a 0.03 mm diameter gold-plated tungsten wire in the centre, and it contains a gas mixture of Xenon, CO_2 and O_2 . When a charged particle passes through a straw tube, it ionizes the gas. The ions drift towards the cathode straw wall, while the electrons drift towards the anode wire to be read out. The TRT has a total of 350,000 readout channels; the barrel portion of the TRT consists of 50,000 straws of length 144 cm that are read out at each end, and the endcaps portion consists of 250,000 straws of length 39 cm . A typical track will cross 36 TRT straws. The tubes are interleaved with polypropylene fibres in the barrel, and polypropylene foils in the endcap regions. These polypropylene materials are what generate transition radiation when charged particles traverse them. Transition radiation is a form of electromagnetic radiation emitted when a charged particle passes through inhomogeneous media, such as a boundary between two different media. The transition radiation produced when charged particles traverse the polypropylene is exploited in order to distinguish electrons from pions. Because more transition radiation is produced for particles with a large gamma factor, i.e., more relativistic particles, more transition radiation is produced for electrons compared to pions or even muons. As the TRT measures the number of high-threshold hits produced by this transition

radiation, i.e., hits that are above a certain threshold, electrons traversing a tube will produce more transition radiation on average compared with heavier particles such as pions, and hence the probability of generating high-threshold hits for electrons will be larger. Figure 4.9 depicts how this feature of transition radiation can be used in order to distinguish different particles.

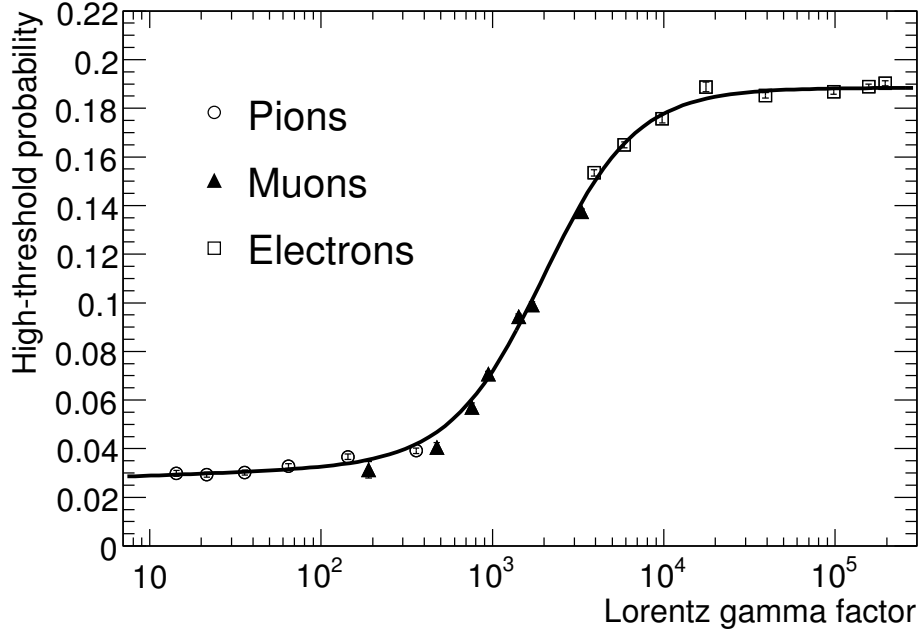


Figure 4.9: Average probability of a high-threshold hit in the barrel TRT as a function of the Lorentz gamma factor for electrons (open squares), muons (full triangles), and pions (open circles) in the energy range 2 – 350 GeV [2].

4.4 Calorimeters

After having traversed the ID, particles coming from the IP reach the calorimeters of ATLAS. Calorimeters are composed of dense layers of absorber material, and active layers of material designed to read out the signals generated by the interaction of the particles with the material. The calorimeters in ATLAS provide coverage in the range $|\eta| < 4.9$, and consist of two main components: the EM calorimeter, used to detect particles interacting via the EM force, such as electrons and photons, and the hadronic calorimeter, used to detect particles interacting via the strong force, such

4.4. Calorimeters

as quarks and gluons. The EM calorimeter is described in Section 4.4.1, and the hadronic calorimeter is described in Section 4.4.2. For the range $|\eta| < 2.5$, i.e., the range also covered by the ID, the EM calorimeter has a fine granularity optimally suited for precision measurements of electrons and photons. Once particles pass through the EM calorimeter, they enter the hadronic calorimeter. Ideally, all SM particles are contained within the calorimeters of ATLAS, with the exception of muons and neutrinos, in order to limit punch-through into the muon system and obtain a better energy measurement from the particles. Figure 4.10 depicts the ATLAS calorimeter system. Table 4.1 summarizes the various components of the ATLAS calorimeter system.

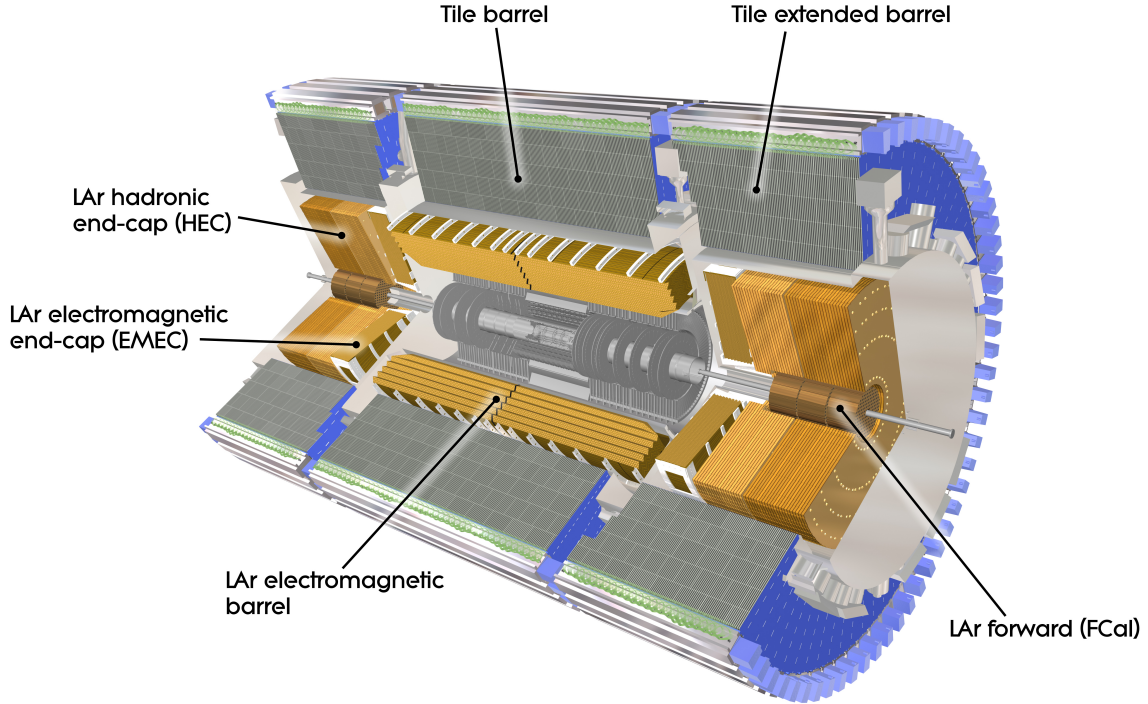


Figure 4.10: The ATLAS calorimeter system. The EM calorimeter, composed of alternating layers of lead and Liquid Argon, is closest to the beam pipe, and the hadronic calorimeter, comprising the tile calorimeter, the hadronic endcap calorimeter, and the forward calorimeter, is farther from the beam pipe [2].

4.4. Calorimeters

Calorimeter	Coverage	Active Material	Passive Material
EM Barrel	$ \eta < 1.475$	LAr	Lead
EM Endcap	$1.375 < \eta < 3.2$	LAr	Lead
Tile Barrel	$ \eta < 1.0$	Scintillating Tiles	Steel
Tile Barrel Extended	$0.8 < \eta < 1.7$	Scintillating Tiles	Steel
Hadronic Endcap	$1.5 < \eta < 3.2$	LAr	Copper
Forward	$3.1 < \eta < 4.9$	LAr	Copper & Tungsten

Table 4.1: Summary of the ATLAS calorimeter system components.

4.4.1 Electromagnetic Calorimeter

The EM calorimeter is used to detect both photons and electrons. It is composed of alternating layers of lead and Liquid Argon (LAr), which are used to initiate particle showers and measure energy deposits, respectively. The LAr measures energy deposits by collecting the charge produced by the particles passing through and ionizing the LAr. In total, the EM calorimeter is approximately 22 radiation lengths thick, where a radiation length is defined as the mean distance over which an electron loses all but $1/e$ of its energy by bremsstrahlung, or $7/9$ the mean free path of a photon before producing an electron-positron pair. An accordion geometry is used throughout the barrel, in order to obtain complete ϕ symmetry without azimuthal cracks. The barrel region of the EM calorimeter consists of two identical half-barrels separated by a 4 mm gap at $z = 0$, and covers $|\eta| < 1.475$. Each endcap region of the EM calorimeter consists of an inner wheel covering $2.5 < |\eta| < 3.2$ and an outer wheel covering $1.375 < |\eta| < 2.5$. Starting after the ID, the EM calorimeter comprises a presampler, used for particles starting to shower before the EM calorimeter, followed by three main layers. The first layer, also known as the strips, is segmented finely along η in order to distinguish single photons from $\pi^0 \rightarrow \gamma\gamma$ decays. The fine segmentation corresponds to 0.0031×0.098 in $\eta \times \phi$. The second layer, where most of the electron energy is deposited, has a granularity of 0.025×0.0245 in $\eta \times \phi$. The third layer of the EM calorimeter is used to capture the tails of the electron shower, and has a segmentation of 0.05×0.0245 in $\eta \times \phi$. Figure 4.11 depicts the granularity of these layers in $\eta \times \phi$.

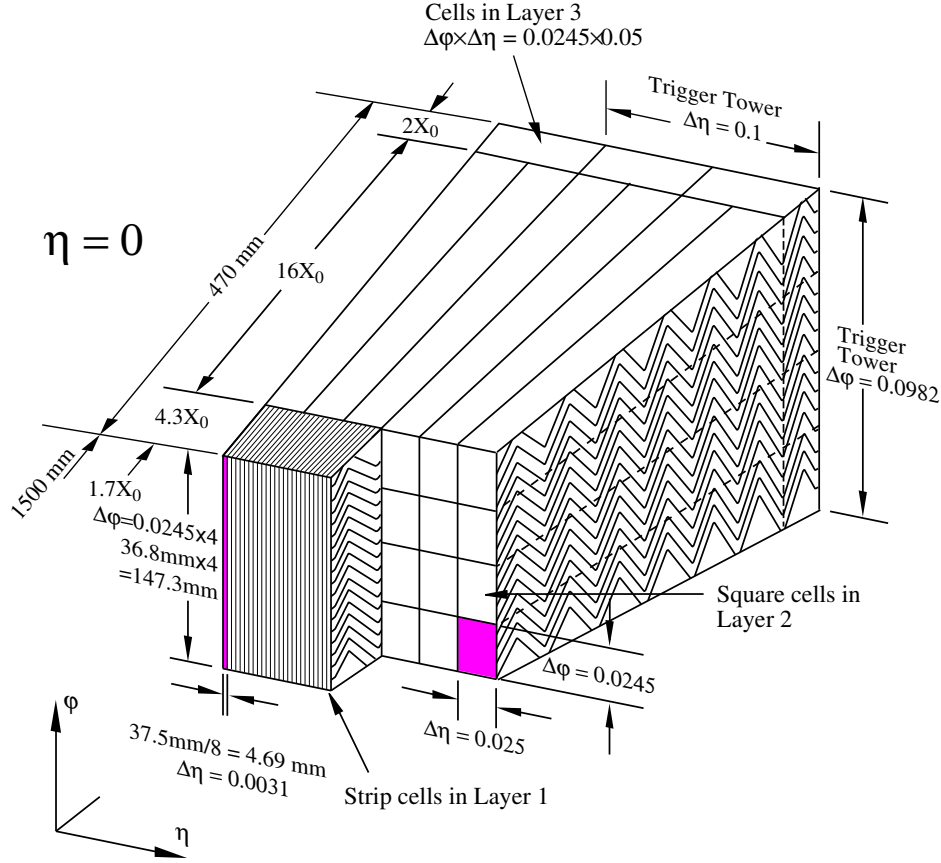


Figure 4.11: Schematic representation of a barrel module in the EM calorimeter. The granularity in η and ϕ of the cells of each of the three layers and of the trigger towers is also shown [2].

4.4.2 Hadronic Calorimeter

The hadronic calorimeter is used to detect strongly interacting particles such as hadrons. Its thickness is 9.7 interaction lengths at $\eta = 0$, where an interaction length is defined as the mean distance travelled by a hadronic particle before undergoing an inelastic nuclear interaction. It can be separated into three main components: a tile calorimeter, a hadronic endcap calorimeter, and a forward calorimeter. The tile calorimeter is a sampling calorimeter using steel as the absorber and scintillating tiles as the active material. Its barrel covers $|\eta| < 1.0$ and its two extended barrels cover $0.8 < |\eta| < 1.7$. The hadronic endcap calorimeter uses LAr as its active material and copper

as its passive material, and covers $1.5 < |\eta| < 3.2$. Finally, the forward calorimeter uses LAr as its active material, and covers $3.1 < |\eta| < 4.9$. Each forward calorimeter consists of three modules in each endcap: the first is made of copper and measures EM interactions, while the last two are made of tungsten and measure mainly hadronic interactions.

4.5 Muon Spectrometer

The MS is used to detect muons. The toroidal magnetic field present in the MS allows for transverse momentum (p_T) measurements by measuring the curvature of the muon tracks magnetically deflected in the MS. The two major regions in the MS are the barrel region covering $|\eta| < 1.1$ and the two endcap regions covering $1.1 < |\eta| < 2.7$. Four different types of detectors are present in the MS: Resistive Plate Chambers (RPCs) and Thin Gap Chambers (TGCs) are used for triggering purposes, and Monitored Drift Tubes (MDTs) and Cathode Strip Chambers (CSCs) are used for precision tracking. Overall, the MS has a muon p_T resolution ranging from 1.7% in the barrel for $p_T \sim 10$ GeV to 4% in the endcaps for $p_T \sim 100$ GeV [53]. More importantly for the analysis presented in this dissertation, the MS provides muon momentum measurements with a relative resolution of up to 10% for muons with $p_T \sim 1$ TeV. In this section, RPCs and TGCs are discussed in Sections 4.5.1 and 4.5.2, and MDTs and CSCs are described in Sections 4.5.3 and 4.5.4. Figure 4.12 depicts the ATLAS muon system.

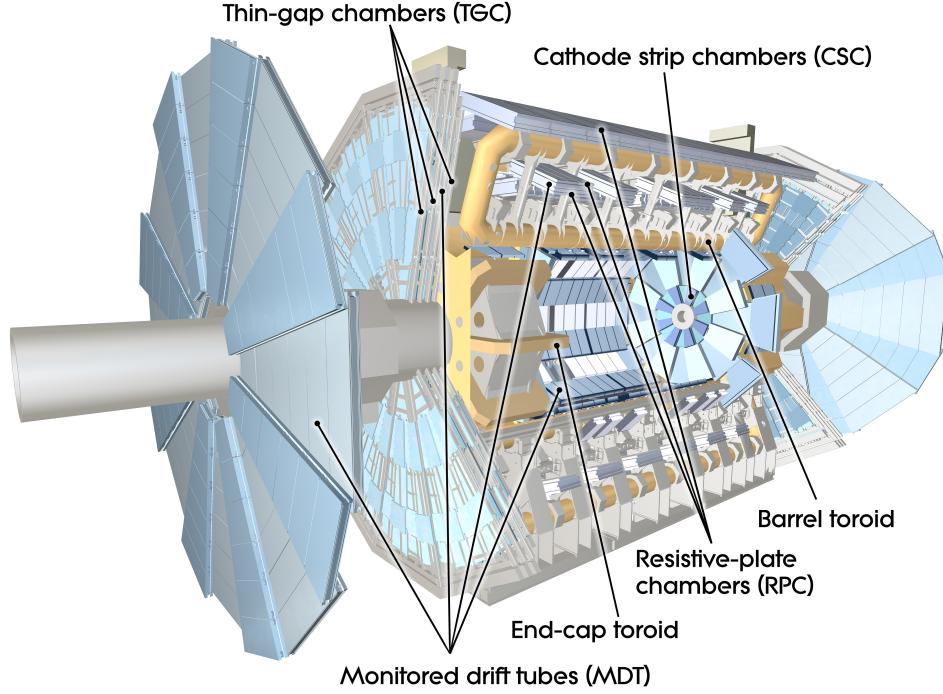


Figure 4.12: The ATLAS muon system. The magnetic field is provided by one barrel toroid and two endcap toroids. Resistive Plate Chambers and Thin Gap Chambers are used for triggering, and Monitored Drift Tubes and Cathode Strip Chambers are used for precision tracking [2].

The general idea of the MS is that each muon passes through a number of so-called stations. In the barrel, three stations are arranged in concentric cylinders around the beam axis. In the endcap regions of the MS, three stations are also present and organized into three planar disks, or wheels. Ranging from 0.15 T to 2.5 T with an average value of 0.5 T in the barrel region, and from 0.2 T to 3.5 T in the endcap regions, the toroidal magnetic field bends electrically charged particles passing through it in order to measure their momentum. The magnetic field integral as a function of $|\eta|$ for two different values of ϕ is shown in Figure 4.13.

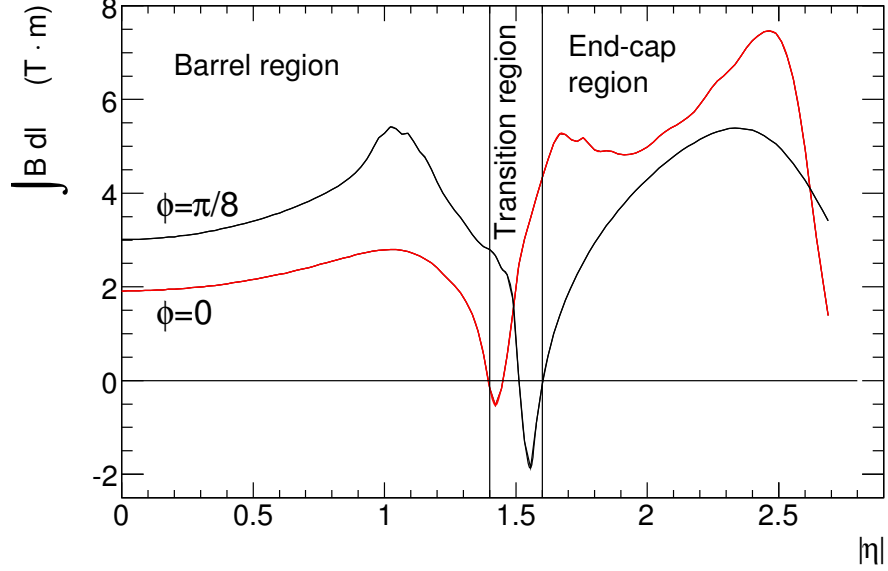


Figure 4.13: Magnetic field integral in the MS as a function of $|\eta|$ for $\phi = 0$ (red) and $\phi = \pi/8$ (black) [2].

Various combinations of detector technologies are used for each station. A summary of the stations and their respective detector compositions is given in Table 4.2. In addition to the three stations described in Table 4.2, additional MDT chambers are installed between the inner and middle endcap stations, in order to improve tracking for muons passing through the barrel-endcap overlap regions. Figure 4.14 depicts a cross-sectional view of one quadrant of the ATLAS muon system and the relevant detector technologies present within it.

Detector Region	Station	Detector Technology
Barrel	Inner	MDT
	Middle	MDT, RPC
	Outer	MDT, RPC
Endcap	Inner	CSC, MDT, TGC
	Middle	MDT, TGC
	Outer	MDT

Table 4.2: Summary of the ATLAS MS stations.

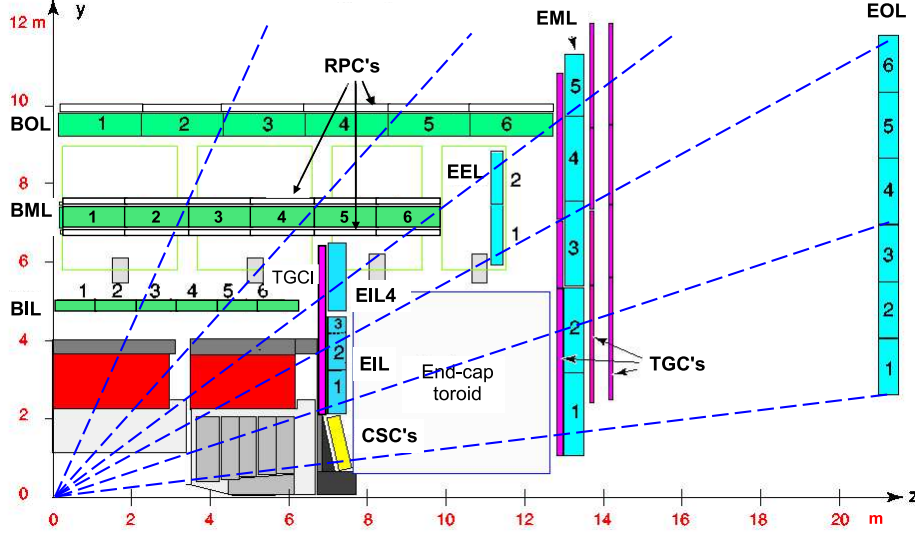


Figure 4.14: Cross-sectional view of one quadrant of the ATLAS muon system [2].

In order to ensure as much coverage as possible, overlapping sectors of large and small chambers are installed throughout the MS. Figure 4.15 gives a cross-sectional view of the same quadrant of the ATLAS muon system shown in Figure 4.14 and highlights the positioning of the small and large chambers. The basic naming convention for the various MS chambers uses three letters to identify specific chambers. The first letter indicates the region of the detector and represents either the barrel (B) region or the endcap (E) regions. The second letter indicates the layer in the MS, and represents either the inner (I), middle (M), or outer (O) layer. Certain chambers have their second letter correspond to the transition region between the barrel and the endcap, i.e., the extra (E) layer. The third letter represents the station type, and can represent either small (S) or large (L) MS chambers. For example, the BIS chamber would refer to the inner chamber within the barrel region with small chamber type. Some exceptions exist, such as the chambers in the barrel extended region mounted on the endcap toroid and elevator shafts, denoted as Barrel Endcap Extra (BEE), and the chambers in the feet region of the detector, denoted as BMG. Figure 4.16 depicts the $R - \phi$ cross-section of the MS, and focuses in particular on the chamber naming of the large and small barrel chambers.

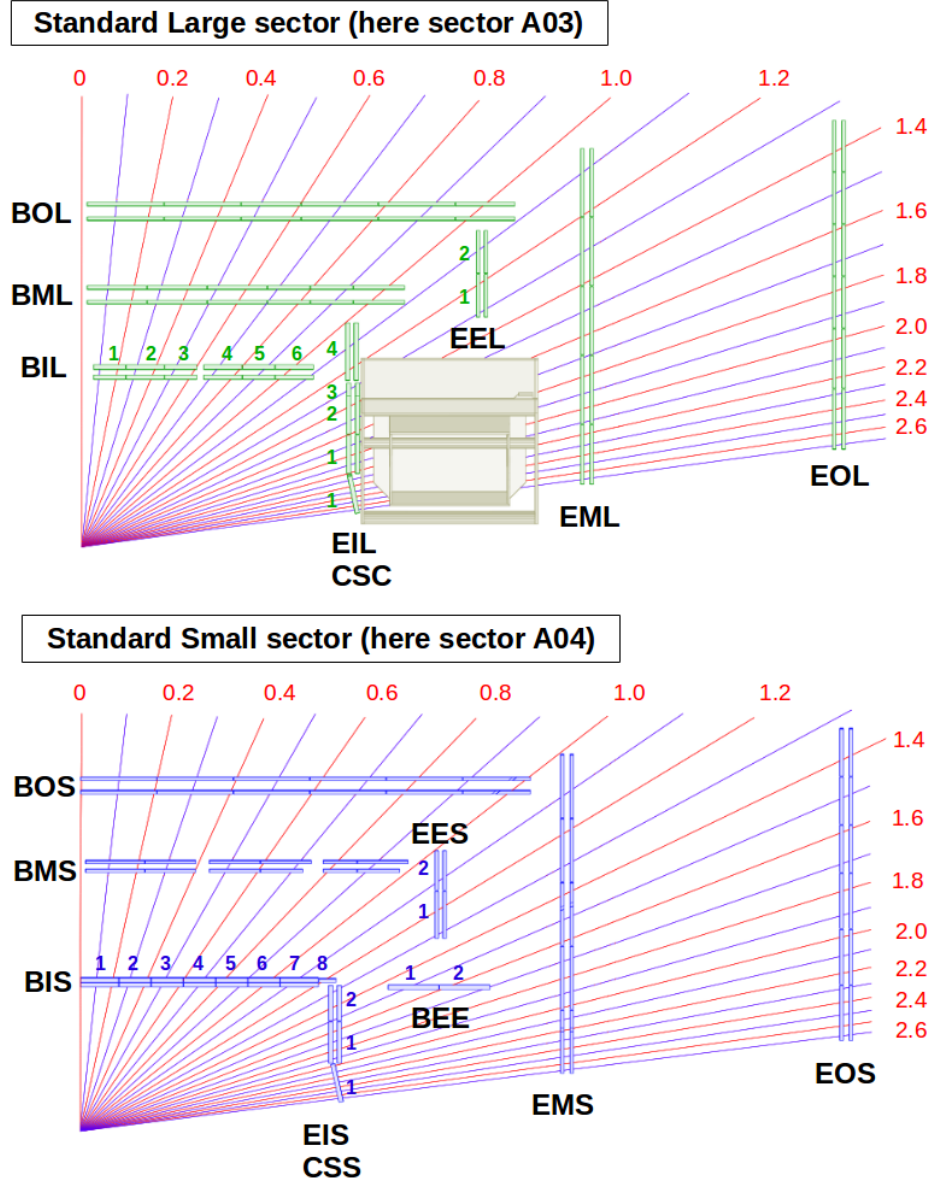
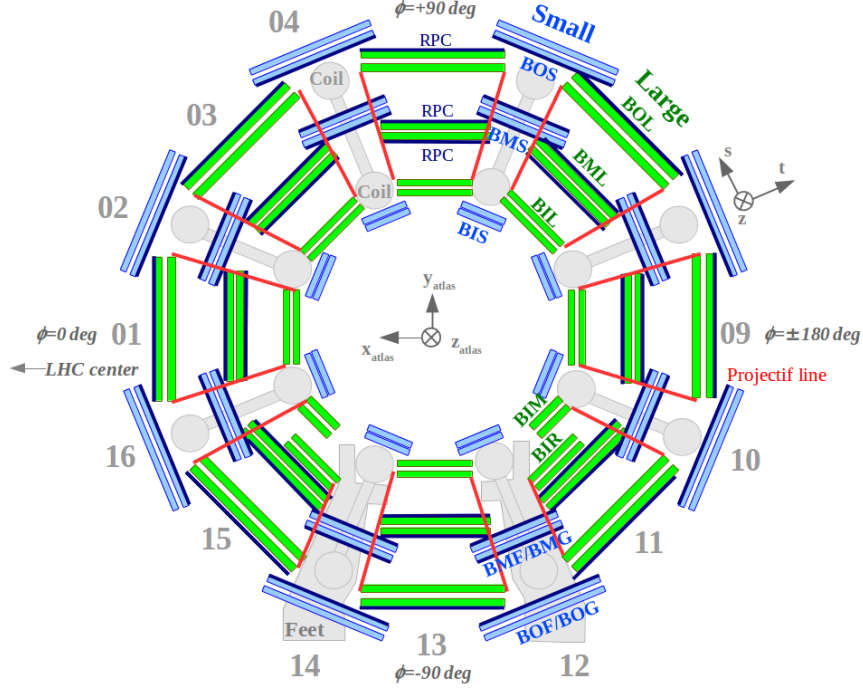


Figure 4.15: Naming convention for the chambers within the ATLAS muon system [54].


 Figure 4.16: $R - \phi$ cross-section of the ATLAS muon system [55].

4.5.1 Resistive Plate Chambers

RPCs are used for triggering on muon tracks contained within the barrel, i.e., $|\eta| < 1.05$. There are three concentric layers of RPCs in the barrel: two layers in the middle station and one layer in the outer station of the MS, as shown in Figure 4.14. Each layer of RPCs has two rectangular gas gaps, allowing for a total of six measurements for a single particle track. Each RPC chamber contains overlapping units, as depicted in Figure 4.17. A single gas gap functions as a gaseous parallel electrode-plate detector, and consists of two parallel resistive plates separated by 2 mm of insulating spacers. The gas gap is filled with a mixture of $\text{C}_2\text{H}_2\text{F}_4$ (94.7%), $\text{Iso-C}_4\text{H}_{10}$ (5%), and SF_6 (0.3%). Metallic strips are mounted on the outer faces of the resistive plates, with one side having the strips segmented in the η direction, and the other having the strips segmented in the ϕ direction. The operating voltage of each gap is 9.8 kV, and charged particles passing through the gap will ionize the gas and immediately create an avalanche of electrons to drift towards the anode plate due to the high electric field. This avalanche

will induce a signal on the strips above and below the ionized region, via capacitive coupling.

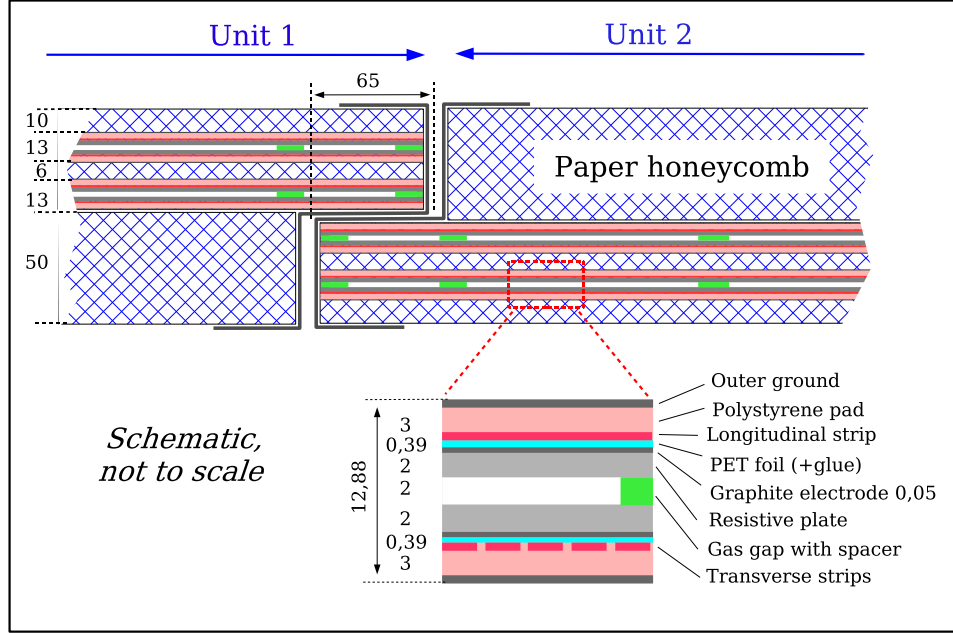


Figure 4.17: Schematic of a Resistive Plate Chamber layer within the ATLAS detector. The dimensions given are in mm [2].

4.5.2 Thin Gap Chambers

With a timing resolution smaller than 25 ns for 99% of tracks, TGCs are used for their triggering capabilities, and provide a muon track measurement in the ϕ direction. They cover the endcap regions $1.05 < |\eta| < 2.7$. They function as multi-wire proportional chambers, with the cathode strips being distributed radially, and the anode wires being distributed perpendicularly to the strips. The gas contained in the TGCs is a mixture of 55% CO₂ and 45% n-pentane. Each gap has a width of 2.8 mm, and the wires are held at a voltage of 2900 V. The width of each wire is 50 μ m, and the distance between consecutive wires within the TGC is 1.8 mm. Due to the absence of a magnetic field between the last two MS stations in the endcaps, the ϕ coordinate in the outer MS station can be extrapolated from the middle station measurement. Hence, no TGCs are present in the outer MS station. Seven layers (one triplet and two doublets) of TGCs are present in the middle

station, while two layers (one doublet) are present in the inner station of the MS.

4.5.3 Monitored Drift Tubes

MDTs provide precision measurements of the η coordinate of a muon track. They are the most pervasive detector technology within the MS, covering $|\eta| < 2.7$. The basic element within the MDT chambers is a pressurized drift tube with a diameter of 29.97 mm. Each tube contains a gas mixture of 93% Ar and 7% CO₂, and a 50 μ m tungsten-rhenium wire at its centre. The wire is kept at a voltage of 3080 V. When a muon traverses the MDT, it ionizes the gas and produces ions, which drift towards the cathode tube, and electrons, which drift towards the anode wire. In order to keep the MDT chambers aligned with each other and with respect to the overall detector, the MDT chambers are monitored by alignment sensors that optically monitor any deviations from straight lines due to, e.g., temperature changes in the environment. Magnetic field sensors are also implemented in the MDTs in order to map the magnetic field. The optimal tracking resolution of an MDT chamber is 30 μ m. A schematic representation of an MDT is shown in Figure 4.18(a). MDT chambers are formed of three to four layers of tubes mounted on each side of a support structure, as depicted in Figure 4.18(b).

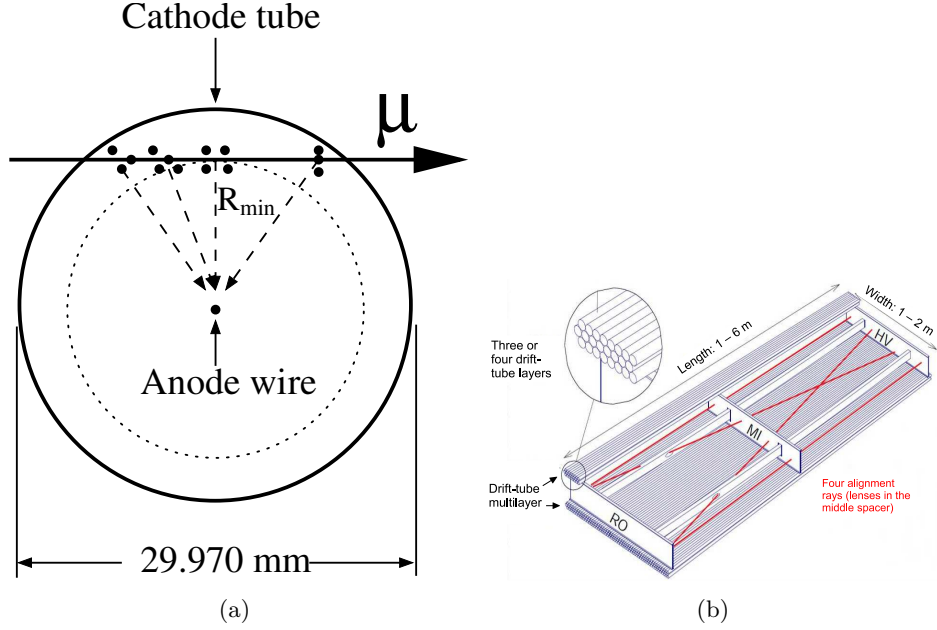


Figure 4.18: Schematic of a single Monitored Drift Tube (left) and a Monitored Drift Tube chamber (right) [2].

4.5.4 Cathode Strip Chambers

The main purpose of the CSCs is to provide precision tracking in the innermost endcap regions where particle rates are beyond the MDT readout limits, i.e., 150 Hz/cm^2 . Able to operate at up to 1 kHz/cm^2 , they cover the range $2.0 < |\eta| < 2.7$. CSCs are multi-wire proportional chambers with anode wires distributed radially and cathode strips in the $\eta - \phi$ plane. One side of the CSC has strips segmented perpendicularly to the wires, allowing for a measurement of the η coordinate, while the other side has its strips segmented parallel to the wires, allowing for a measurement of the ϕ coordinate. Only the cathode strips of the CSCs are read out.

Chapter 5

Object Reconstruction

After having passed the trigger selection described below in Chapter 7, events of interest are fully reconstructed, and relevant physics objects are defined. In this chapter, an overview of the different objects used in physics analyses carried out using the ATLAS experiment will be detailed. In particular, Section 5.1 will describe tracks reconstructed in the ID, and Section 5.2 will cover the definition of vertices within ATLAS. Furthermore, energy deposits in the calorimeters that are reconstructed as jets will be presented in Section 5.3. Next, charged lepton reconstruction will be presented: electrons will be discussed in Section 5.4, and muons will be reviewed in Section 5.5. While not directly relevant to the search for new physics presented in this dissertation, an electron veto is applied in the analysis described in Section 7.3. As muons with large transverse momentum are of particular relevance for the search for new physics presented in this dissertation, an in-depth description of muon reconstruction, with emphasis on high- p_T muons, is presented in Chapter 6. Finally, missing transverse momentum (E_T^{miss}) will be defined in Section 5.6.

5.1 Tracks

Charged particles traversing the ID are bent due to the Lorentz force in the solenoid magnetic field. This is exploited by the reconstruction procedure for tracks, which includes global χ^2 and Kalman filter techniques [56]. The procedure is based on fitting a track model to a set of space point measurements. A space point is defined as a measurement of a particle's position in three dimensions. Tracks are obtained by using track seeds formed from a combination of space points from the pixel layers and the first layer of the SCT. Next, the track seeds are used to find track candidates, which include space points from all SCT layers using a Kalman filter technique, in order to take into account the material traversed by the charged particle. Finally, the track candidates are extended to include information from the TRT detector. Once the extended track candidates are obtained, the full track is refitted with the full information of all three ID components. From the fit

5.2. Vertices

of the track, various measured quantities, M , are obtained:

$$M = (d_0, z_0, \phi, \theta, q/p), \quad (5.1)$$

where d_0 and z_0 are the impact parameters, ϕ and θ are the angles described in Section 4.2, and q/p represents the charge of the particle divided by its momentum. The impact parameters quantify the distance of the particle track to the IP. In particular, the transverse impact parameter d_0 represents the distance between the track and the IP in the plane transverse to the proton beams, and the longitudinal impact parameter z_0 represents the distance between the track and the IP along the z-axis. ID tracks are used in both electron (muon) reconstruction; they are combined with calorimeter (MS) information to obtain the final measurement of the particle parameters.

5.2 Vertices

After all ID tracks are identified, a vertex-finding algorithm based on the reconstructed tracks is used in order to find the vertices of each event [57]. The collection of tracks described in Section 5.1 is refined in order to reject low-quality tracks. First, tracks are required to be consistent with the IP. Consistency with the IP is enforced by requiring the impact parameter $d_0 < 4$ mm, and the impact parameter uncertainties $\sigma_{d_0} < 5$ mm and $\sigma_{z_0} < 10$ mm. Next, each track must have at least four hits in the SCT and at least nine hits overall when considering the combination of SCT and pixel hits. The track must also not contain any pixel holes. A hole is defined as an active detector region where no hit is found, but where a hit would be expected based on the track information. Finally, the tracks used to reconstruct vertices must have $p_T > 400$ MeV. The vertex-finding procedure starts by selecting a seed position for the first vertex consistent with the IP, and the collection of tracks satisfying the above criteria is used to estimate the best vertex position with a fit. Iteratively, all tracks satisfying the criteria above are added to or removed from the vertex, in order to minimize a χ^2 fit. If a track is not associated to the vertex, it is allowed to be associated to another vertex. The procedure is repeated until either all ID tracks are associated to a vertex or no additional vertex is found with the remaining unassociated tracks. When multiple vertices exist in an event, the primary vertex is defined as the vertex within the event with the highest sum of track p_T squared.

5.3 Jets

Quarks and gluons produced from proton-proton interactions will hadronize into colour-neutral states producing a collimated shower of particles, collectively called a jet. The amount of energy deposited in the hadronic calorimeters in comparison with the EM calorimeters allow jets to be distinguished from electrons or photons; jets usually have a much larger fraction of their energy deposited in the hadronic calorimeters. To reconstruct jets, the anti- k_t algorithm [58] is used. This algorithm computes various distances and iterates over available entities. In the case of the ATLAS experiment, entities are either tracks or energy clusters measured in the calorimeters. The distance between two entities labeled i and j , d_{ij} , or the distance between an entity labeled i and the jet beam labeled B , d_{iB} , is computed using

$$d_{ij} = \min(k_{ti}^{-2}, k_{tj}^{-2}) \frac{\Delta_{ij}^2}{R^2} \quad \text{and} \quad d_{iB} = k_{ti}^{-2}, \quad (5.2)$$

where $\Delta_{ij}^2 = (y_i - y_j)^2 + (\phi_i - \phi_j)^2$, and k_{ti} , y_i , and ϕ_i are the transverse momentum, rapidity, and azimuthal angle of particle i , respectively. R is a radius parameter and varies depending on the analysis using the algorithm. A typical value of R is 0.4. The rapidity is defined as

$$y = \frac{1}{2} \ln \left(\frac{E + p_z}{E - p_z} \right), \quad (5.3)$$

where E and p_z are the energy and the momentum along the z-axis of the particle, respectively.

For each entity, the algorithm compares d_{ij} and d_{iB} and selects the smaller of the two distances. If the smaller distance is d_{ij} , then the entity labeled j is combined with the entity labeled i . If the smaller distance is d_{iB} , then the entity labeled i is labeled a jet, and it is removed from the list of entities. The algorithm terminates once the list of entities has been exhausted and all entities have been classified.

5.3.1 b -jets

A particular class of jets are those resulting from the decay of b quarks, which hadronize to b hadrons, so-called b -jets. The relatively long lifetime of b hadrons implies that they will travel within the detector for several millimetres before decaying. The unique topology of such jets is exploited in order to identify b -jets. In particular, a secondary vertex can be identified where the b hadron decayed. Figure 5.1 depicts the typical topology of a b -jet

and highlights the distinction between the primary vertex and the secondary vertex.

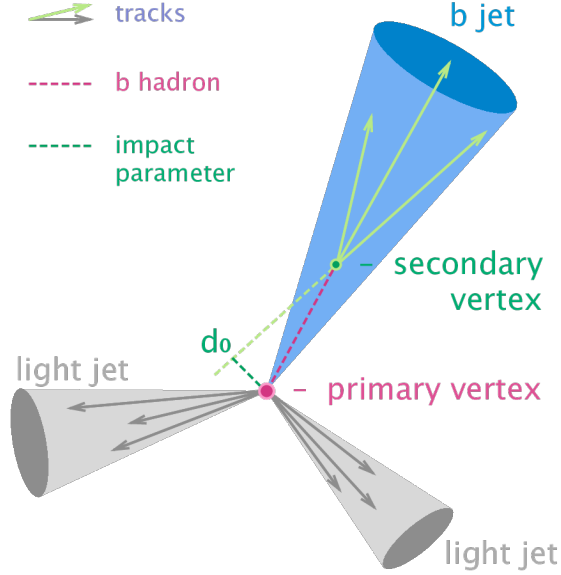


Figure 5.1: Schematic representation of the topology of a typical b -jet compared with jets produced by lighter quarks or gluons [59].

Machine learning algorithms use information relating to the secondary vertex and its associated tracks in order to distinguish b -jets from jets originating from lighter quarks or gluons [60–62]. In particular, the distance between the primary vertex and the secondary vertex, the impact parameters of the tracks associated to the jet, and the information relating to the secondary vertex, are used as inputs to low-level b -jet taggers. The output of these individual low-level taggers are then used in a Boosted Decision Tree (BDT), which enhances the separation power of the individual low-level taggers. The BDT outputs a multivariate discriminant, so-called $mv2c10$, which is used in order to quantify the probability of a jet originating from a b quark, i.e., being a real b -jet. Various Working Points (WPs) are defined based on the $mv2c10$ quantity, with the efficiency of identifying a b -jet as their main feature. For example, the $mv2c10$ 70% WP has a b -jet selection efficiency of 70%. A higher b -jet selection efficiency comes at the cost of a higher fake rate due to jets originating from lighter quarks and gluons. Hence, a balance between the desired purity and selection efficiency of b -jets must be struck. The b -jet veto applied to the analysis described in

Section 7.3 requires the absence of any b -jet satisfying the mv2c10 77% WP.

5.4 Electrons

Information from both the ID and the EM calorimeter is used to reconstruct electrons and obtain measurements of their properties. The EM calorimeter, described in Section 4.4.1, is used in order to measure the transverse energy, E_T , of the electrons. When reconstructing an electron candidate, ID tracks, described in Section 5.1, are matched with EM calorimeter clusters. These clusters are formed from energy deposits in the EM calorimeter. Specifically, the EM calorimeter is divided into so-called towers of dimension $\Delta\eta \times \Delta\phi = 0.025 \times 0.025$, corresponding to the granularity of the second layer of the EM calorimeter. The energy in all three EM calorimeter layers and in the presampler is summed to obtain the total energy of each tower. In total, the ATLAS EM calorimeter comprises 200×256 towers in $\eta \times \phi$. A sliding-window algorithm [63] searches for clusters using a window of size 3×5 towers in $\eta \times \phi$, where the towers included have energy deposit greater than 2.5 GeV. EM calorimeter clusters are then matched to ID tracks to form electron candidates by requiring $|\eta_{\text{cluster}} - \eta_{\text{track}}| < 0.05$ and $-0.10 < q \times [\Delta(\phi_{\text{cluster}}, \phi_{\text{track}})] < 0.05$, where q corresponds to the electron's charge. After being matched to an ID track, the cluster is computed again using an extended window size of 3×7 (5×5) towers in the barrel (endcap) region. Figures 4.7, 4.8 and 4.11 give a summary of the ID and EM calorimeter components, and Figure 5.2 illustrates the sub-detectors involved in the electron reconstruction process within the ATLAS detector.

5.4. Electrons

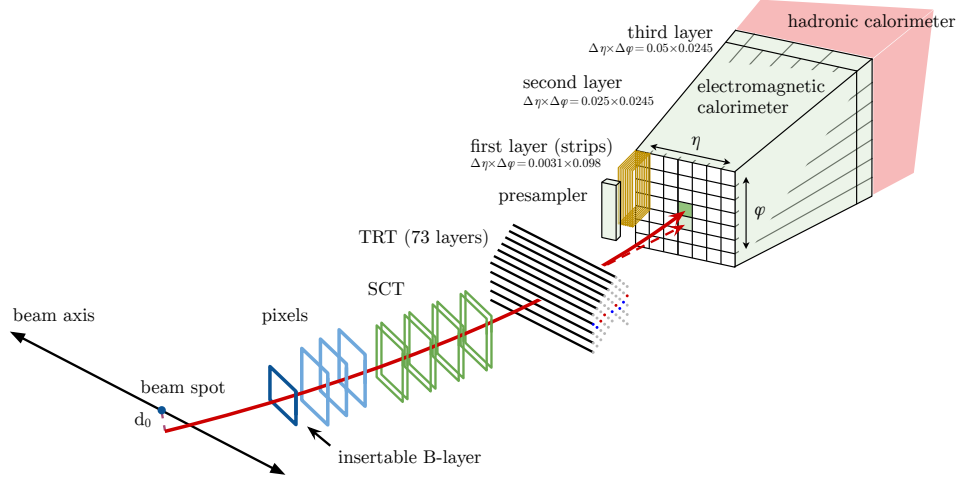


Figure 5.2: Schematic representation of electron reconstruction in ATLAS. The solid red line represents the path of an electron, and the dashed red line corresponds to a photon produced by the interaction of the electron with the material in the ID [64].

The electron veto applied to the analysis described in Section 7.3 uses the electron candidates described above, which satisfy certain quality criteria. In particular, the likelihood-based *Tight* selection WP [64] is required. The selection WP uses probability distributions of the input discriminating variables to construct a likelihood ratio used to separate signal, i.e., real, prompt electrons, from background, i.e., fake electrons, such as those originating from photon conversion inside the detector. The inputs to the likelihood-based discriminant relate to the electron candidate's EM calorimeter cluster and track measurements. The complete list of input variables used in the likelihood-based discriminant can be found in [64]. In addition to the selection WP, the *Gradient* isolation WP must be satisfied. The *Gradient* isolation WP uses information from the EM calorimeter and the ID in order to quantify the level of activity surrounding the electron candidate. Requiring little activity around an electron object helps to reduce the number of fake electrons selected. The efficiency of the electron isolation selection is 90% (99%) at 25 GeV (60 GeV). Finally, electrons in the transition region between the barrel and the endcaps, i.e., $1.37 < |\eta| < 1.52$, are rejected, due to the large amount of inactive material present, which leads to poor identification capabilities and energy resolution.

5.5 Muons

The objects most relevant to the search presented in this dissertation are muons. The muon transverse momentum is defined as

$$p_T = p \sin \theta = \frac{\sin \theta}{|q/p|}, \quad (5.4)$$

where q/p is the ratio of the muon charge to its momentum, and is measured from the track curvature. For muons with track segments coming from all three stations of the MS, the track curvature is obtained from the track sagitta. The sagitta of a curved track of radius r is defined as the distance from the centre of the track arc to the centre of its base. Figure 5.3 depicts the sagitta s for a curved track with radius r . For muons with track segments coming from only two stations of the MS, the track curvature is obtained from the angular difference between the two measured segments.

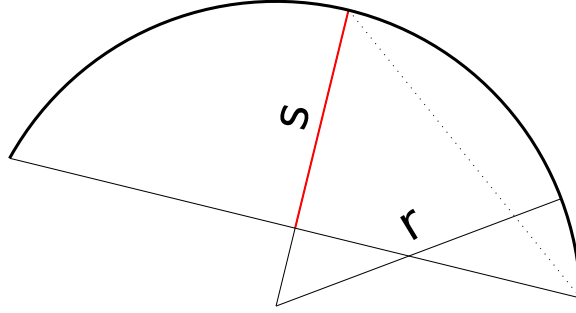


Figure 5.3: Schematic representation of the sagitta s of a curved track with radius r [65].

The main types of muon tracks used within ATLAS are Combined (CB) muon tracks and Muon Extrapolated (ME) tracks [53]. A CB muon track is one where the track reconstruction for the muon is performed independently in the ID and the MS, and a combined track is formed with a global refit that uses the hits from both the ID and MS sub-detectors. An ME muon track is one where the track is reconstructed based only on the MS track, and a loose requirement on the compatibility of it originating from the IP is imposed. ME muons are used mainly to improve the muon acceptance in the $2.5 < |\eta| < 2.7$ region of the detector, where no ID information is available.

5.5.1 *Medium* Quality Muons

The default selection criteria required for muons used in ATLAS analyses are defined as the *Medium* quality WP [53]. A *Medium* quality muon is defined as either a CB track or an ME track, with additional quality requirements. CB muons are required to have at least three hits in at least two MDT layers. CB muons in the central region $|\eta| < 0.1$ are required only to have at least one MDT layer, but no more than one MDT hole layer. Used only in the region $2.5 < |\eta| < 2.7$, ME muons are required to have at least three MDT or CSC layers hit. In order to avoid the selection of hadrons which decay in flight and get misidentified as muons, a requirement on the significance between the ID and MS p_T measurements is placed, requiring $\mathcal{S}_{p_T^{\text{ID}}, p_T^{\text{MS}}} < 7$, where the significance between two measurements A and B is defined as

$$\mathcal{S}_{A:B} = \frac{|A - B|}{\sqrt{\sigma_A^2 + \sigma_B^2}}, \quad (5.5)$$

and σ_A and σ_B represent the uncertainties of the respective measurements. This is effective in rejecting hadrons that would leave a track in the ID with large momentum and subsequently decay into a muon, resulting in a track with low momentum in the MS, e.g., kaons.

5.5.2 Transverse Momentum Corrections

The muon momentum scale and resolution are simulated using $Z \rightarrow \mu\mu$ samples. In order to simulate the scale at the per mille level and the resolution at the percent level, corrections to the simulated momentum are applied. With the aim of describing the collected data as precisely as possible, the corrected transverse momentum in simulation is calculated using

$$p_T^{\text{Cor,Det}} = \frac{p_T^{\text{MC,Det}} + \sum_{n=0}^1 s_n^{\text{Det}}(\eta, \phi) \left(p_T^{\text{MC,Det}}\right)^n}{1 + \sum_{m=0}^2 \Delta r_m^{\text{Det}}(\eta, \phi) \left(p_T^{\text{MC,Det}}\right)^{m-1} g_m}, \quad (5.6)$$

where $p_T^{\text{MC,Det}}$ is the uncorrected transverse momentum in simulation, $\text{Det} \in [\text{ID}, \text{MS}]$, g_m is a normally distributed random variable with zero mean and unit width, $\Delta r_m^{\text{Det}}(\eta, \phi)$ describes the momentum resolution smearing, and $s_n^{\text{Det}}(\eta, \phi)$ describes the scale corrections. These corrections are derived for various detector regions separately, ensuring uniform detector behavior for each set of corrections. The correction parameters are derived from data.

The numerator in Equation (5.6) describes the momentum scale, and the denominator describes the momentum smearing that broadens the p_T resolution. The term $s_0^{\text{Det}}(\eta, \phi)$ models the inaccuracy of simulating the energy loss of muons due to interactions with material. In particular, $s_0^{\text{MS}}(\eta, \phi)$ is used to model the muon energy loss between the interaction point and the MS, mostly due to interactions with the calorimeters. As there is very little material interactions between the interaction point and the ID, $s_0^{\text{ID}}(\eta, \phi) = 0$. $s_1^{\text{Det}}(\eta, \phi)$ is used to account for inaccuracies in the magnetic field integral. The p_T resolution of muons can be parameterised as

$$\frac{\sigma_{p_T}}{p_T} = \frac{r_0}{p_T} \oplus r_1 \oplus r_2 \cdot p_T, \quad (5.7)$$

where $\oplus$ denotes a sum in quadrature. The first term in Equation (5.7) models the fluctuations of the energy loss in the traversed materials. As muon energy loss affects mostly momentum scale, not resolution, $\Delta r_0^{\text{ID}}(\eta, \phi) = \Delta r_0^{\text{MS}}(\eta, \phi) = 0$. The second term accounts for multiple scattering, local magnetic field inhomogeneities, and local radial displacements of the hits. The last term, most important at high p_T , describes intrinsic resolution effects caused by the spatial resolution of the hit measurements and by residual misalignments of the MS.

The final corrected momentum for CB muons is obtained by combining the ID and MS corrected momenta from Equation (5.6) into a weighted average:

$$p_T^{\text{Cor,CB}} = f \cdot p_T^{\text{Cor,ID}} + (1 - f) \cdot p_T^{\text{Cor,MS}}, \quad (5.8)$$

where f is derived using an equation identical to Equation (5.8), but using the uncorrected momenta values and solving for f .

5.5.3 Transverse Momentum Uncertainty

The uncertainty on the transverse momentum is of prime importance, as it is used in the high- p_T selection WP described in Section 6.1. From the track measurements in Equation (5.1), the covariance matrix for the track measurement is

$$\text{cov}(M) = \begin{pmatrix} \sigma_{d_0}^2 & \sigma_{z_0, d_0} & \sigma_{\phi, d_0} & \sigma_{\theta, d_0} & \sigma_{q/p, d_0} \\ \sigma_{d_0, z_0} & \sigma_{z_0}^2 & \sigma_{\phi, z_0} & \sigma_{\theta, z_0} & \sigma_{q/p, z_0} \\ \sigma_{d_0, \phi} & \sigma_{z_0, \phi} & \sigma_{\phi}^2 & \sigma_{\theta, \phi} & \sigma_{q/p, \phi} \\ \sigma_{d_0, \theta} & \sigma_{z_0, \theta} & \sigma_{\phi, \theta} & \sigma_{\theta}^2 & \sigma_{q/p, \theta} \\ \sigma_{d_0, q/p} & \sigma_{z_0, q/p} & \sigma_{\phi, q/p} & \sigma_{\theta, q/p} & \sigma_{q/p}^2 \end{pmatrix}. \quad (5.9)$$

5.6. Missing Transverse Momentum

We can expand p_T in a Taylor series in order to get p_T expressed as a linear combination of θ and q/p :

$$p_T \approx p_T^0 + \frac{\partial p_T}{\partial \theta} \theta + \frac{\partial p_T}{\partial q/p} q/p. \quad (5.10)$$

In general, the error for a linear combination

$$f = \sum_i^n a_i x_i, \quad (5.11)$$

where a_i are the linear combination coefficients and x_i are the measured values, is

$$\sigma_f^2 = \sum_i^n a_i^2 \sigma_i^2 + \sum_i^n \sum_{j(j \neq i)}^n a_i a_j \sigma_{ij}. \quad (5.12)$$

Now that p_T is a linear combination of θ and q/p , we can use the propagation of uncertainty method for linear combinations of correlated variables to obtain

$$\sigma_{p_T}^2 = \left(\frac{\partial p_T}{\partial \theta} \sigma_\theta \right)^2 + \left(\frac{\partial p_T}{\partial q/p} \sigma_{q/p} \right)^2 + 2 \left(\frac{\partial p_T}{\partial \theta} \right) \left(\frac{\partial p_T}{\partial q/p} \right) \sigma_{\theta, q/p}. \quad (5.13)$$

This result will be important when considering quantities that depend on the uncertainty of muon momentum measurements, such as the uncertainty of the dimuon invariant mass presented in Equation (9.3).

5.6 Missing Transverse Momentum

The ATLAS detector cannot directly detect SM neutrinos. In order to account for these particles in analyses, their presence is inferred using so-called missing transverse momentum, defined as

$$E_T^{\text{miss}} = - \left[\sum_i^{\text{all objects}} p_T^i + E_T^{cl} \right], \quad (5.14)$$

where the sum includes the transverse momenta p_T^i of all leptons, jets, and photons, and the E_T^{cl} term accounts for energy clusters in the calorimeters over threshold which are not associated to any of the objects already included in the sum. It should be noted that other unknown particles have

5.6. *Missing Transverse Momentum*

the potential to leave the ATLAS detector without being detected. This would leave a final state signature of large E_T^{miss} , and certain searches for such BSM physics scenarios are undertaken within the ATLAS collaboration.

Chapter 6

High-Momentum Muons

Objects with large transverse momentum are used in many searches for BSM physics. In addition to the muons described in Section 5.5, of particular importance for the search presented in this dissertation are muons with large transverse momentum. In particular, high- p_T muons are key when searching for new high-mass resonances, because of the very nature of the search; hints of new physics will manifest themselves as an excess of events in the high-mass tail of the dimuon invariant mass distribution. Because higher invariant mass events are characteristic of new and interesting physics, the main candidates for the final state objects are inherently high- p_T muons. This chapter describes the treatment of high- p_T muons and their relevance to the search described in this dissertation. The high- p_T WP is described in Section 6.1, and the optimization of a component of the high- p_T WP, the BadMuon veto, is presented in Section 6.2. Finally, the gradient WP, which investigates whether efficiency losses due to the high- p_T WP selection can be recovered without compromising the p_T resolution, is discussed in Section 6.3.

6.1 High- p_T Working Point

The high- p_T WP consists of a set of criteria that a muon must satisfy. In addition to being a CB muon of quality *Medium* as described in Section 5.5, muons selected by the high- p_T WP must satisfy the following criteria [53]:

- At least three hits in at least three precision layers of the MS
- Not excluded by geometrical $\eta - \phi$ chamber vetoes, described below
- ID q/p measurement agrees with ME q/p measurement within 7 standard deviations, i.e., $\mathcal{S}_{q/p^{\text{ID}}:q/p^{\text{ME}}} < 7$
- ID track has at least one pixel hit, at least five SCT hits, at least 10% of the TRT hits originally assigned to the track before the combined

6.1. High- p_T Working Point

fit (in the region of full TRT acceptance $0.1 < |\eta| < 1.9$), and fewer than three pixel or SCT holes

- Satisfies the BadMuon veto, described below

While requiring three MS stations reduces the reconstruction efficiency by approximately 20%, this requirement improves the p_T resolution by approximately 30% for muons with $p_T > 1.5$ TeV. The geometrical vetoes are imposed in order to avoid chambers in the MS with poor alignment, which would lead to p_T mismeasurements. The vetoed chambers include the region where the barrel and endcaps overlap, i.e., $1.01 < |\eta| < 1.1$, the BIS7/8 chambers, and the BMG chambers located in the region accommodating the supporting feet of the ATLAS detector. A muon passing through the BEE chambers is rejected if its track has fewer than four layers hit. Finally, if the muon passes through the overlap region between small and large MS chambers, i.e., if hits are found in both small and large chambers, the muon is rejected. The BadMuon veto is a cut on the relative q/p error measurement, and has been optimized in order to reject individual muon candidates with poorly measured p_T . The BadMuon veto depends on both the η and p_T of the muon considered. The optimization of the BadMuon veto is discussed in Section 6.2.

This muon selection has a p_T -dependent efficiency. In particular, at larger values of p_T the efficiency degrades in part because of catastrophic muon energy loss, i.e., the emission of energetic photons in the dense material traversed by muons. This effect in turn leads to an η -dependent selection efficiency loss of approximately 3% per TeV on average for muons with $p_T > 200$ GeV, and is folded into the high- p_T WP as a systematic uncertainty. Table 6.1 summarizes the efficiency losses due to catastrophic muon energy loss for muons with $p_T > 200$ GeV passing the high- p_T WP in various η ranges.

6.1. High- p_T Working Point

η Range	Drop [%/TeV]
$0.1 < \eta < 1.0$	2.74 ± 0.09
$1.0 < \eta < 1.2$	1.74 ± 0.28
$1.2 < \eta < 1.3$	4.09 ± 0.66
$1.3 < \eta < 1.4$	3.18 ± 0.75
$1.4 < \eta < 1.7$	12.03 ± 0.52
$1.7 < \eta < 2.0$	6.20 ± 0.70
$2.0 < \eta < 2.5$	12.25 ± 1.22

Table 6.1: Drop in selection efficiency per TeV due to catastrophic muon energy loss for muons with $p_T > 200$ GeV passing the high- p_T WP in different η ranges.

Figures 6.1 and 6.2 depict the acceptance times efficiency and the dilepton mass resolution as a function of Z' pole mass, respectively. The dielectron channel is beyond the scope of this dissertation, but it can be seen that the efficiency in this channel is much higher compared with the dimuon channel described in this dissertation.

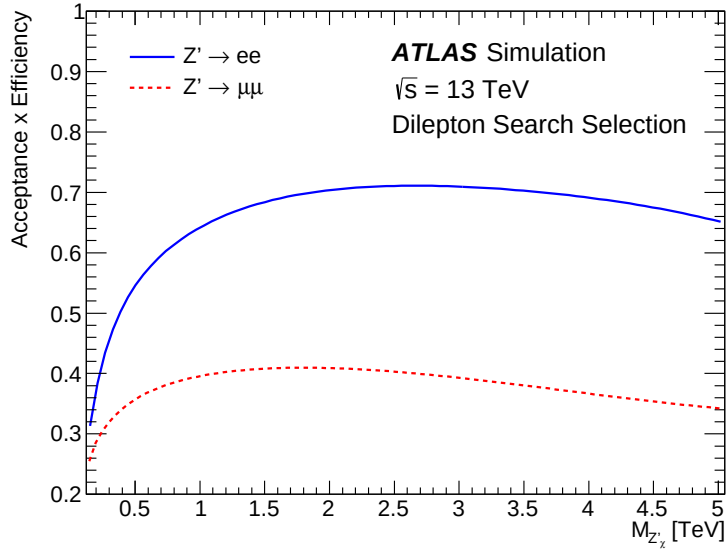


Figure 6.1: Acceptance times efficiency for the dielectron (solid blue) and dimuon (dashed red) selections as a function of the Z' pole mass. Both plots are shown after the full dilepton selection described in Section 9.1.

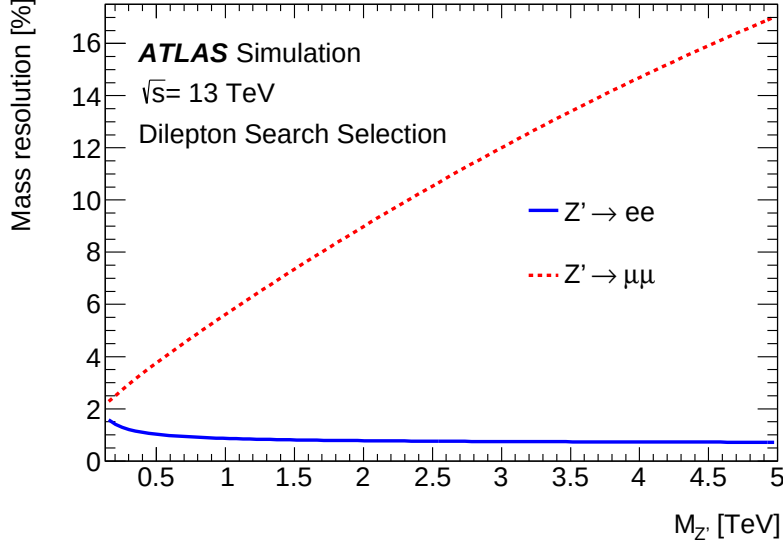


Figure 6.2: Z' mass resolution for the dielectron (solid blue) and dimuon (dashed red) selections as a function of the Z' pole mass. Both plots are shown after the full dilepton selection described in Section 9.1.

6.2 BadMuon Veto Optimization

The goal of the BadMuon veto is to reject pathological muons residing in the tails of the p_T resolution distributions. The efficiency of this veto diminishes with increasing p_T because of a cut on the relative error of the q/p measurement. This cut is optimized in order to continue rejecting pathological muons, but improve the efficiency of the selection as much as possible. Dedicated samples produced with realistic detector misalignments are used for these optimizations. The expected resolution of the momentum measurement is parameterised as a function of p_T in 5 different η regions following the parameterisation described in Equation (5.7). The η regions are chosen to coincide with specific regions of the ATLAS detector having varying p_T resolution, and are defined as follows: $|\eta| \leq 1.05$, $1.05 < |\eta| \leq 1.3$, $1.3 < |\eta| \leq 1.7$, $1.7 < |\eta| \leq 2.0$, and $|\eta| > 2.0$. Assuming σ is the expected resolution of a muon candidate as expressed in Equation (5.7) given its η and p_T measurements, a cut on the relative error of the q/p measurement

of the muon is then applied:

$$\frac{\sigma_{q/p}}{q/p} < C(p_T) \cdot \sigma, \quad (6.1)$$

where $C(p_T)$ is a p_T -dependent coefficient equal to 1.8 if $p_T \leq 1$ TeV which linearly decreases, or tightens, with p_T if $p_T > 1$ TeV. Two possible optimizations are investigated: the p_T value at which the tightening starts, and the initial value of $C(p_T)$. The tightening point is increased from 1.0 TeV to 3.5 TeV in increments of 500 GeV. Increasing this value shows minimal selection efficiency improvements, so a value of 1.0 TeV is kept. The initial value of $C(p_T)$ represents the multiple of the expected resolution allowed when starting the tightening of the p_T -dependent cut on the relative error of the q/p measurement. This value is increased from 1.8 to 2.0, 2.5, and 3.0, while maintaining the tightening point discussed above at 2.5 TeV. In all cases, the cut on the relative error of the q/p measurement is linearly tightened as a function of p_T until the cut reaches a value of 1σ at $p_T = 5$ TeV. Increasing this value improves the cut efficiency while still effectively removing the tails of the resolution distributions, but diminishing returns are observed when increasing the initial value of $C(p_T)$ from 2.5 to 3.0. In the transition region, $1.3 < |\eta| \leq 1.7$, and the CSC region, $|\eta| > 2.0$, the cut is conservatively relaxed to allow for an initial value of $C(p_T)$ equal to 2.0. In all other detector regions, $C(p_T)$ is relaxed to start at 2.5. In summary, the new definition of the BadMuon veto is one where the p_T value at which the tightening of the cut starts is 1.0 TeV, and the initial value of $C(p_T)$ is 2.5 everywhere, except in the transition region and in the CSC region, where it is 2.0. The optimized BadMuon veto removes the tails of the p_T resolution distributions, defined as $(p_T/p_T^{\text{tru}}) - 1$, where p_T is the reconstructed transverse momentum and p_T^{tru} is the truth-level p_T . Various η and p_T configurations are investigated. In particular, the resolution distributions of the q/p measurements for the high- p_T WP, the high- p_T WP and the new implementation of the BadMuon veto, and the high- p_T WP and the old implementation of the BadMuon veto, in the region $1.05 < |\eta| \leq 1.3$ with $700 \text{ GeV} < p_T < 1.3 \text{ TeV}$, are shown in Figure 6.3. Evidently, the old implementation of the BadMuon veto and the optimized version of the BadMuon veto are both removing unwanted events in the tail of the q/p residual distribution, as intended. In addition, the optimized BadMuon veto implementation is recovering events in the core of the residual distribution, allowing for an improved selection efficiency.

6.2. BadMuon Veto Optimization

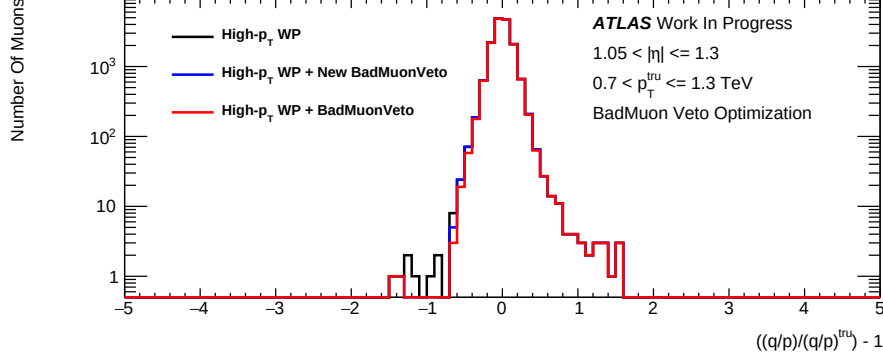


Figure 6.3: q/p residual distribution for muons in the region $1.05 < |\eta| \leq 1.3$ with $700 \text{ GeV} < p_T \leq 1.3 \text{ TeV}$ satisfying the high- p_T WP (black), the high- p_T WP and the new implementation of the BadMuon veto (blue), and the high- p_T WP and the old implementation of the BadMuon veto (red).

Figures 6.4 and 6.5 give a summary of the efficiency gain for the optimized BadMuon veto, and the increase in the fraction of events in the tails of the p_T resolution distributions. An efficiency gain of up to 5% with respect to the previous definition of the BadMuon veto, i.e., starting the tightening at 1.0 TeV with an initial value of $C(p_T)$ of 1.8, is seen in certain η regions while only increasing the tail fraction at the per mille level. The tail fraction is defined as the fraction of events in the p_T resolution distributions above one.

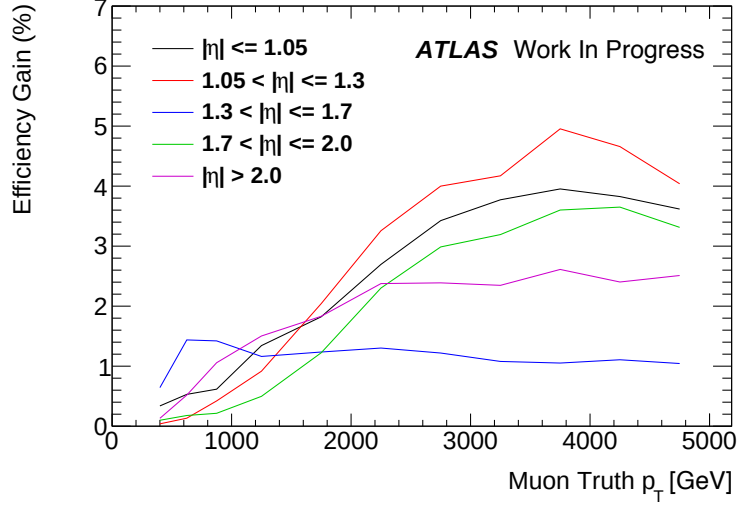


Figure 6.4: Summary of the efficiency gain of the optimized BadMuon veto with respect to its previous definition in the five $|\eta|$ regions studied.

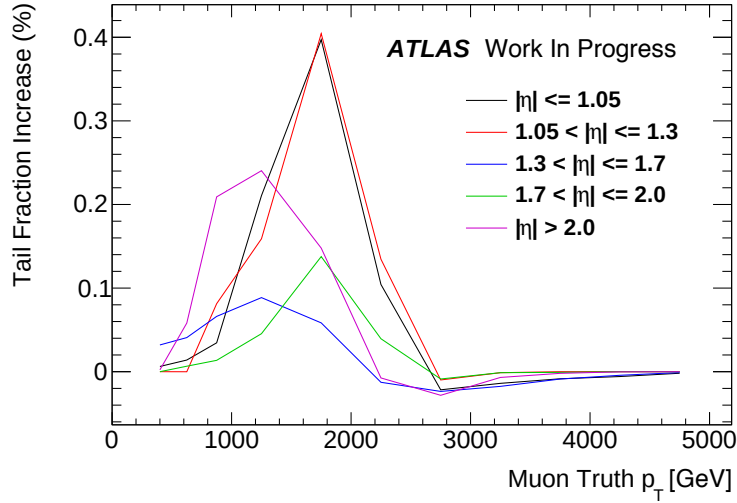


Figure 6.5: Summary of the increase in the fraction of events in the tails of the p_T resolution distributions of the optimized BadMuon veto with respect to its previous definition in the five $|\eta|$ regions studied. The tail fraction is defined as the fraction of events in the p_T resolution distributions above one.

6.3 Gradient Working Point

Until 2015-2016, the analysis group searching for new high-mass phenomena in events with two muons used the high- p_T WP described in Section 6.1. This WP results in a relatively good single-muon p_T resolution and dimuon mass resolution, but has an inherently low efficiency, particularly when compared with the dielectron efficiency, as seen in Figures 6.1 and 6.2.

The lower efficiency in the dimuon channel seen in Figure 6.1 is a direct result of the low efficiency of the single-muon identification WP. As can be seen in Figure 6.2, the dimuon mass resolution, which is inherently worse compared with the dielectron resolution, is poor in absolute terms. The idea behind the high- p_T WP is to maximize selection efficiency while maintaining adequate single-muon p_T resolution with a series of geometrical vetoes, a veto of 2-station muons, i.e. requiring at least three hits in at least three precision layers of the MS, and a cut on the relative error of the q/p measurement, i.e., the BadMuon veto. While the high- p_T WP p_T resolution is good, the WP has a very low efficiency. This section investigates whether the significant losses of efficiency due to the high- p_T WP vetoes can be largely recovered without compromising the p_T resolution. The simulated samples used are dedicated samples produced with realistic detector misalignments: one sample with two muons per event having $300 < p_T < 1000$ GeV, one single-muon sample with $1.0 < p_T < 3.0$ TeV, and one single-muon sample with $3.0 < p_T < 5.0$ TeV. The procedure is as follows:

1. Define a single-muon truth-level acceptance selection, the so-called truth acceptance.
2. Require a baseline selection consisting of simple kinematic and quality cuts on single-muon reconstruction-level quantities.
3. Require full truth-to-reconstruction matching between reconstructed baseline muons and truth muons in acceptance.
4. Plot the two-dimensional distribution of the reconstructed quantities $\sigma_{q/p}/(q/p)$ as a function of p_T for the fully matched sample of muons.
5. Derive the 95% quantile of the above distribution.
6. Fit the 95% quantile. This fit is used as the gradient cut.
7. Remove all muons above the 95% quantile fit, i.e., the gradient cut, in the baseline sample where no truth-level cuts are applied.

8. Apply the integrity cuts described below.
9. Inspect the resulting acceptance times efficiency ($\mathcal{A} \times \epsilon$) curve and resolution of the sample passing the baseline and gradient selection.
10. Compare with the high- p_T WP $\mathcal{A} \times \epsilon$ and resolution.

The procedure is carried out separately for each of the geometrical or conceptual vetoes defined in the high- p_T WP.

6.3.1 Analysis

Truth-level acceptance selection

The truth-level acceptance is a set of kinematic cuts on truth-level muons to ensure they are within the geometrical and kinematic acceptance of the detector. The selection is presented below.

- $p_T^{\text{tru}} > 30 \text{ GeV}$
- $|\eta_{\text{tru}}| < 2.5$

Baseline selection

The baseline selection is a set of standard kinematic, quality, and identification requirements.

- $p_T > 30 \text{ GeV}$
- $|\eta| < 2.5$
- Muon is type CB
- Muon is of quality *Medium*
- Muon is not in the CSC region vetoed by the high- p_T WP
- Muon is not in the BEE region vetoed by the high- p_T WP

Truth-level to reconstruction-level matching

The new gradient algorithm is based on a tight truth-level to reconstruction-level matching scheme from which the quantiles are derived. First, in addition to truth-level acceptance selection, spatial and charge matching are required; $\Delta R(\mu_{\text{tru}}, \mu_{\text{rec}}) < 0.01$, $q_{\text{tru}} \times q_{\text{rec}} > 0$. Next, the reconstructed muon p_T must match the p_T of the truth muon assuming an ideal detector, i.e., considering the designed intrinsic resolution. This resolution is digitized from the MS Technical Design Report (TDR) [66], and is shown in Figure 6.6. The digitized points in Figure 6.6 are fitted and extrapolated above 1 TeV. The fits are done separately for the barrel and the endcap regions using higher-order polynomials. Since this is a digitized version of the TDR figure, the errors of the points are unavailable and are fixed to an arbitrarily small value. The p_T matching scheme thus follows the ideal resolution. The requirement is obtained from the fit, i.e.,

$$\frac{|p_T^{\text{tru}} - p_T|}{p_T^{\text{tru}}} < \sigma_{\text{TDR}}(p_T^{\text{tru}}), \quad (6.2)$$

where $\sigma_{\text{TDR}}(p_T^{\text{tru}})$ is the fit of the TDR resolution at a given value of p_T^{tru} .

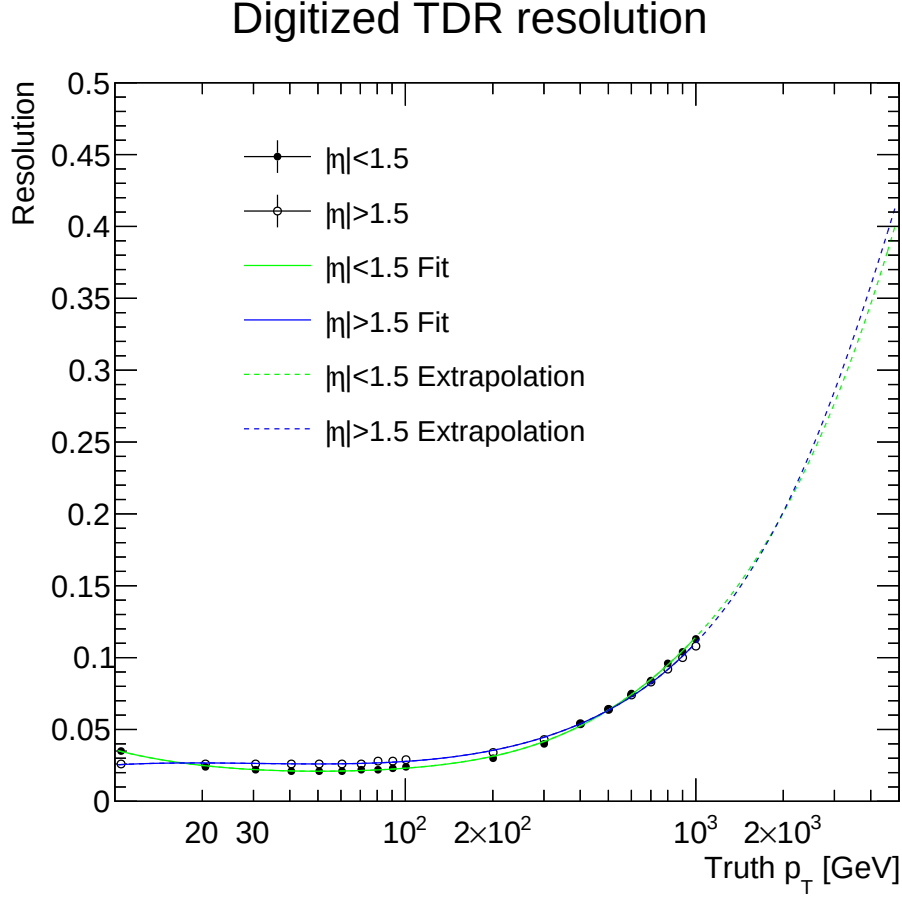


Figure 6.6: Digitized version of the MS TDR resolution, taken as the ideal scenario to define the matching. The fits are extrapolated above 1 TeV.

Derivation of the gradient cut

The gradient WP exploits the correlation between the truth-level to reconstruction-level response, $\frac{|q/p - q/p^{\text{tru}}|}{q/p^{\text{tru}}}$, and the relative error on the q/p measurement, $\frac{\sigma_{q/p}}{q/p}$. The response quantity is available in Monte Carlo (MC) simulation but not in data as truth information is only available in simulated samples. Hence, the gradient cut is derived for the measurable quantity $\frac{\sigma_{q/p}}{q/p}$. To derive the correlation between $\frac{\sigma_{q/p}}{q/p}$ and p_T , the two-dimensional distribution of the baseline-selected fully truth-matched muons is used. This distribu-

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tion represents a sample of muons having well-determined momentum by construction. An unbiased relative error, $\frac{\sigma_{q/p}}{q/p}$, is obtained by taking the non-calibrated and non-smeared q/p or p_T value in the denominator, as the errors are not calibrated or smeared.

During Run 1, the final uncertainties related to muon tracks were inflated to account for alignment uncertainties. This amounted to a deweighting of the chambers in the CB track fit, and was only carried out in specific critical situations such as the barrel-endcap overlap regions or the small-large overlap regions. In Run 2, a more realistic error on the q/p measurement is obtained by fitting the alignment discontinuities. In particular, the CB track fit is performed using a Gaussian constraint on the chamber hits, where alignment uncertainties are used as Gaussian widths. This is accomplished using so-called Alignment Effects On Track (AEOTs). Each AEOT specifies a position and angle uncertainty on the chamber hits to which it is associated.

The 95% quantile¹ of the $\frac{\sigma_{q/p}}{q/p}$ dimension of the two-dimensional distribution is calculated in each bin of p_T . The result is a set of points in the two-dimensional space, below which 95% of the fully truth-matched muons are found. The remaining 5% of the muons above these points naturally have the largest $\frac{\sigma_{q/p}}{q/p}$ values. The set of points is then fit, in order to obtain a smooth description in the full p_T range. This fit is used as the gradient cut. Figure 6.7 shows the two-dimensional distribution of $\frac{\sigma_{q/p}}{q/p}$ as a function of p_T for baseline-selected fully truth-matched muons in the barrel-endcap overlap region with the 95% quantile overlaid. Dedicated samples of single muons generated with a flat η and ϕ distribution with p_T between 300 GeV and 5 TeV produced with realistic detector misalignments are used.

¹For reference, the 50% quantile is simply the median.

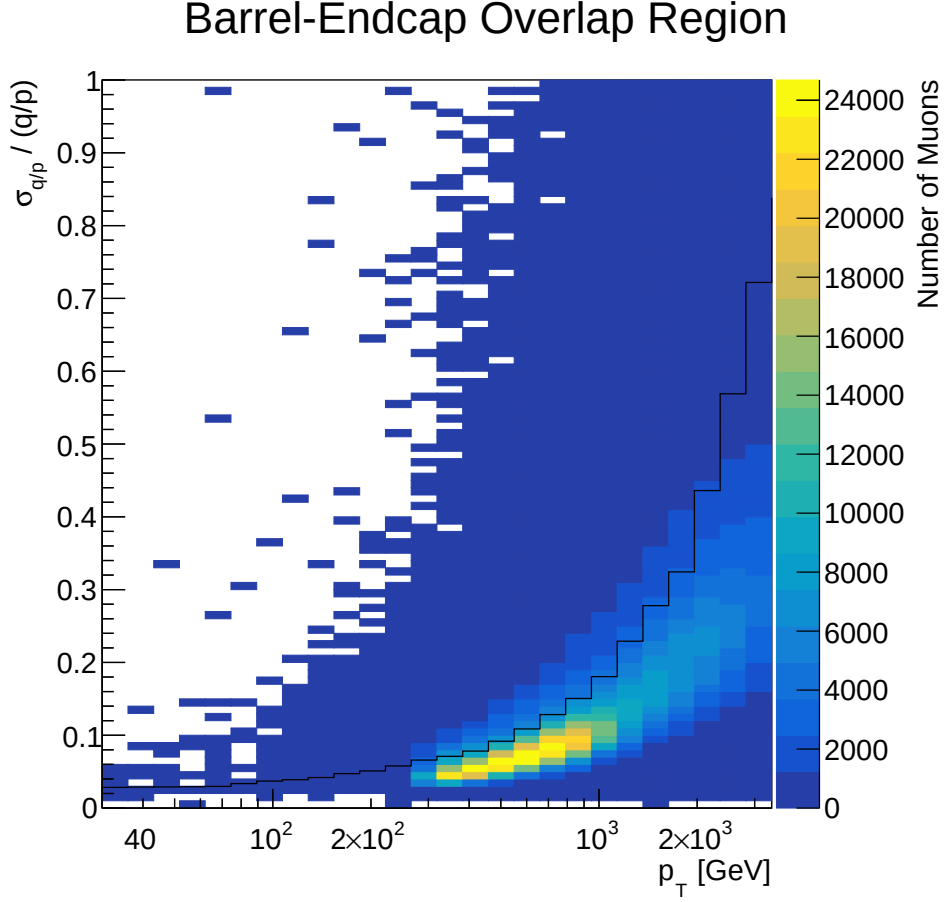


Figure 6.7: Two-dimensional distribution of $\frac{\sigma_{q/p}}{q/p}$ as a function of p_T for baseline-selected fully truth-matched muons in the barrel-endcap overlap region with the 95% quantile line overlaid in black.

Figure 6.8 shows the summary of the 95% quantile fit for muons in the barrel-endcap overlap region. The quantiles and fits of three samples are overlaid: a single-muon sample with realistic detector misalignments, a perfectly aligned single-muon sample, and a DY mass-sliced sample with Z' signal samples, assuming perfect alignment. The perfectly aligned samples are not used in the gradient cut fits and are shown only for comparison. However, in the range below ~ 300 GeV in muon p_T , where there are very few or no entries in the misaligned samples, the quantiles of the DY and Z' samples are used. In this range of p_T , the impact of misalignments is assumed to be negligible, and hence the procedure is assumed to be valid.

The misaligned sample's quantiles are fit with higher-order polynomials in the different regions considered.

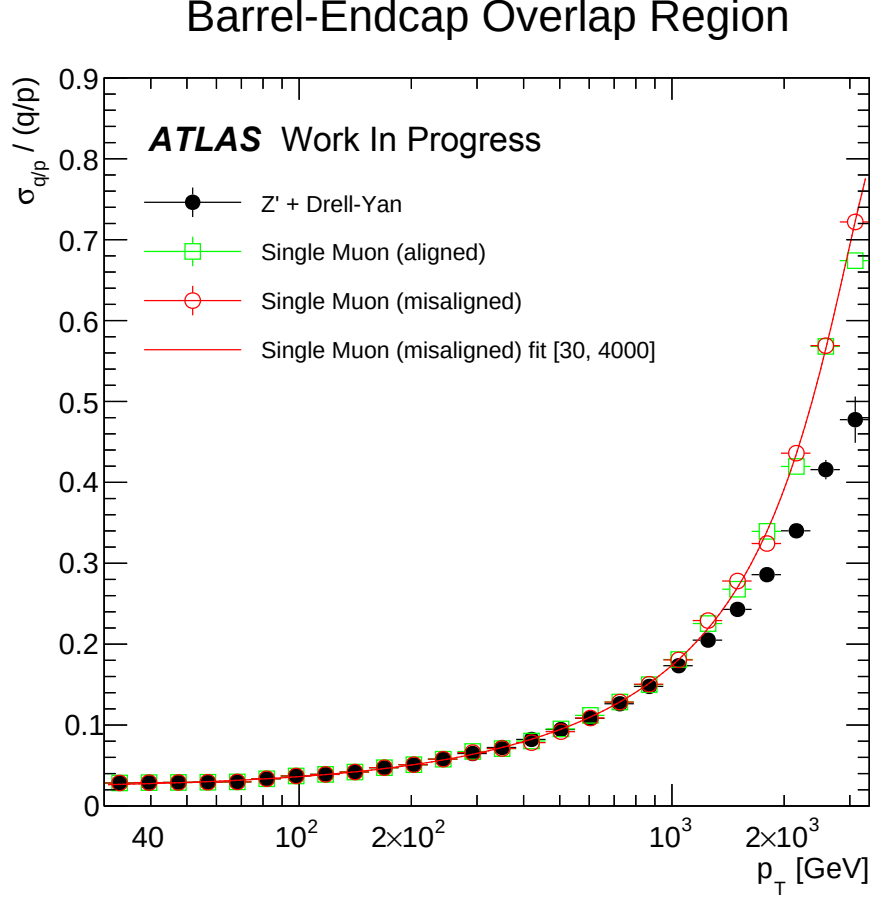


Figure 6.8: Summary of the 95% quantile fit in the barrel-endcap overlap region. The quantiles of three samples are overlaid: the misaligned single-muon sample (red), the perfectly aligned single-muon sample (green), and the perfectly aligned DY mass-sliced samples and Z' signal samples (black).

In addition to the gradient cut, it is essential to independently ensure that the p_T measurement of the selected muons is accurate in the previously vetoed regions of the high- p_T WP. This is accomplished by the integrity cuts discussed below.

Integrity cuts

In addition to the gradient cut described above, integrity cuts are defined in order to ensure that the various p_T measurements for a given muon are consistent with each other when considering the previously-enforced vetoes of the high- p_T WP. This set of cuts is inspired by the existing cut on $\mathcal{S}_{q/p^{\text{ID}}:q/p^{\text{ME}}}$ present in the high- p_T WP. The cuts are tuned using simulated samples with realistic detector misalignments. In addition, 36.1 fb^{-1} of data taken in 2015-2016 is used to understand potential problems. The selection used in this limited data-only study is the dimuon selection described in Chapter 9 with the muon identification set to the *Medium* WP.

The cuts make use of a set of significance variables using the four p_T measurements: ID, ME, MS, and CB. The significance between two measurements is defined in Equation (5.5). The selection criteria applied are such that a muon is rejected if any of the criteria below are met.

- $\mathcal{S}_{p_T^{\text{CB}}:p_T^{\text{ID}}} > 10$ to ensure consistent p_T measurements between the ID measurement and the CB measurement
- $\mathcal{S}_{p_T^{\text{CB}}:p_T^{\text{ME}}} > 10$ to ensure consistent p_T measurements between the ME measurement and the CB measurement
- $\mathcal{S}_{p_T^{\text{MS}}:p_T^{\text{ME}}} > 10$ for $p_T < 1 \text{ TeV}$ or $\mathcal{S}_{p_T^{\text{MS}}:p_T^{\text{ME}}} > 20$ for $p_T > 1 \text{ TeV}$
- $\Delta_{\text{ID:ME:CB}} > 20$, where $\Delta_{\text{ID:ME:CB}} = \frac{|p_T^{\text{ID}} - p_T^{\text{ME}}|}{\sigma_{p_T^{\text{CB}}}}$
- $\mathcal{S}_{p_T^{\text{WA}}:p_T^{\text{CB}}} > 10$, where $p_T^{\text{WA}} = \frac{\sum_i p_T^i / \sigma_{p_T^i}^2}{\sum_i 1 / \sigma_{p_T^i}^2}$, $\sigma_{p_T^{\text{WA}}}^2 = \frac{1}{\sum_i 1 / \sigma_{p_T^i}^2}$, $i \in [\text{ID}, \text{ME}]$

An example distribution of one of these variables, $\mathcal{S}_{p_T^{\text{WA}}:p_T^{\text{CB}}}$, is shown in Figure 6.9. It can be seen that the integrity cuts already reduce the number of events with large $\mathcal{S}_{p_T^{\text{WA}}:p_T^{\text{CB}}}$ even before the cut on this variable. Hence, it can be concluded that the set of integrity cuts applied are consistently removing events in the tails of the variables under consideration. The simulated samples used contain realistic detector misalignments, and the selection is the baseline selection described above. In addition to the baseline selection, the distribution for the baseline selection with all integrity cuts except that of the specific variable plotted, i.e., the so-called $N - 1$ distribution, is shown. The integrity cuts are optimized considering events within the vetoed regions of the high- p_T WP leading to the largest efficiency

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losses: 2-station muons, muons in the small-large and barrel-endcap overlap regions, and muons in the BIS7/8 region. The CSC and BEE vetoes are kept as part of the baseline selection, as they lead to very small efficiency losses. While the MS p_T measurement is never used independently in muon reconstruction for most analyses in ATLAS, the $\mathcal{S}_{p_T^{\text{MS}}:p_T^{\text{ME}}}$ variable is important for determining the integrity of 2-station muons. The idea is to require the MS and ME measurements to be consistent with each other, in order to ensure that the MS alone, although having only two stations participating in the measurement, is capturing the muon's true p_T . The true p_T should be close to the one obtained by the ME algorithm because it accounts for energy losses in the calorimeter.

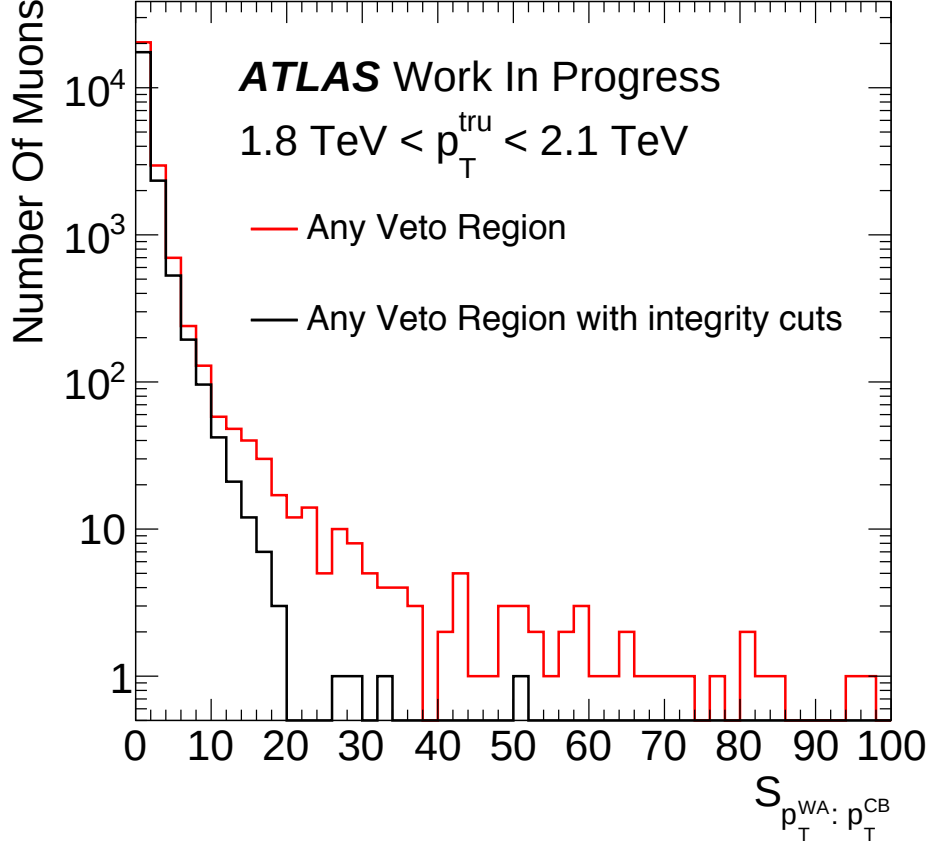


Figure 6.9: $\mathcal{S}_{p_T^{WA}: p_T^{CB}}$ integrity variable in any of the regions vetoed by the high- p_T WP for the baseline selection (red), and the baseline and all integrity cuts except the cut on $\mathcal{S}_{p_T^{WA}: p_T^{CB}}$ (black). The cut is at $\mathcal{S}_{p_T^{WA}: p_T^{CB}} > 10$.

The impact of the integrity cuts is tested on 36.1 fb^{-1} of data taken during 2015-2016 using the sample of highest-mass dimuon events, i.e., events having a dimuon invariant mass larger than 1 TeV, passing the event selection described in Chapter 9, but where both muons pass the *Medium* quality criteria, and at least one of the muons fails the high- p_T WP. Figures 6.10 and 6.11 show the impact of the integrity cuts on the number of muons passing the event selection after the application of the integrity cuts. The events appearing because of the new selection, which would otherwise be rejected due to the high- p_T WP, show no indication of muons with highly

inconsistent p_T measurements. These events are of particular importance, as they allow for a gain in selection efficiency in the phase space most relevant to the search presented in Chapter 9. Indeed, even a few events at high dimuon invariant mass where the contribution from background processes is small enough can lead to an excess of data compared to the SM prediction, which could indicate the presence of BSM physics. The fact that the p_T measurements of the muons recovered from the high- p_T WP fall within a few hundred GeV indicates that no pathological muons are preferentially accepted as a result of the gradient WP procedure.

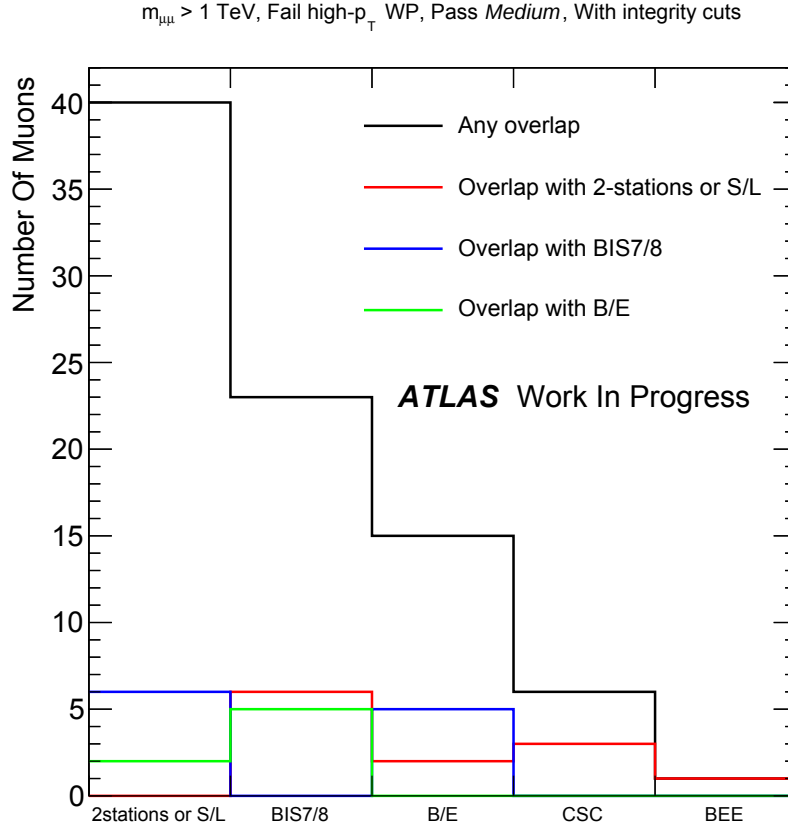


Figure 6.10: Categories of muons after the application of the integrity cuts on data events taken during 2015-2016 with dimuon invariant mass greater than 1 TeV. All muons here are recovered from the high- p_T WP selection.

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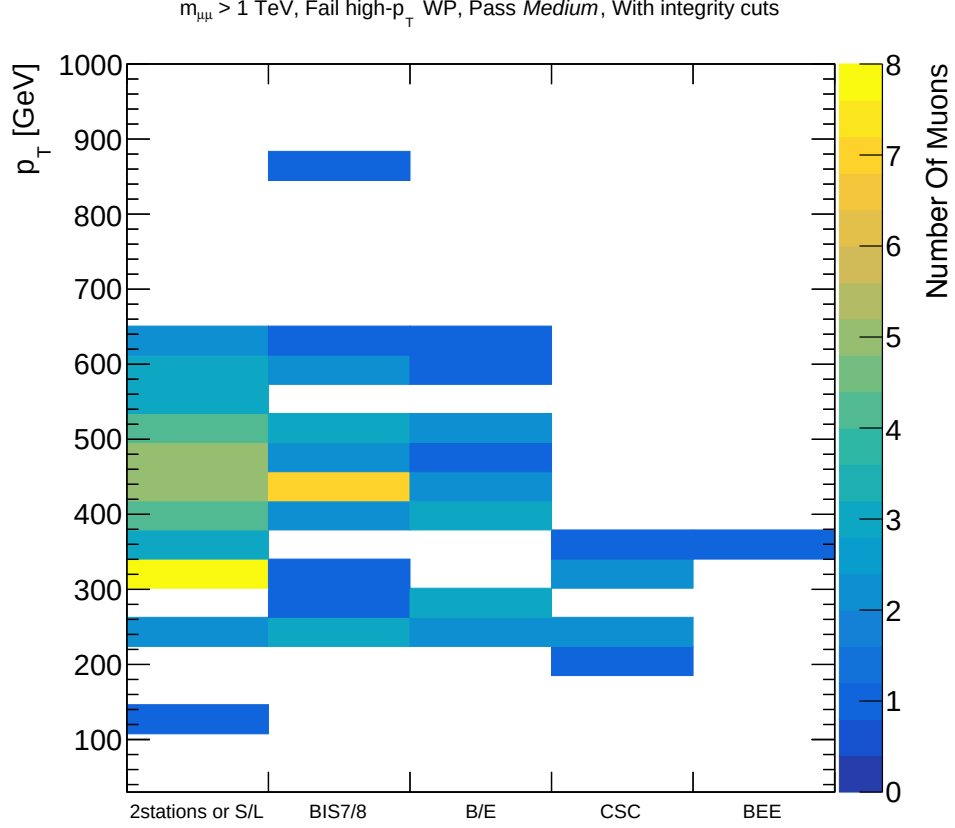


Figure 6.11: Muon p_T as a function of the category of muons after the application of the integrity cuts on data events taken during 2015-2016 with dimuon invariant mass greater than 1 TeV. All muons here are recovered from the high- p_T WP selection.

6.3.2 Results

The figures of merit used to evaluate the performance of the gradient WP are its impact on the dependence of $\frac{\sigma_{q/p}}{q/p}$ as a function of p_T , its muon selection efficiency, and its p_T resolution in comparison with both the baseline selection and the high- p_T WP selection. In order to evaluate how the different selections act on the same sample, it is informative to look at the two-dimensional distribution of $\frac{\sigma_{q/p}}{q/p}$ as a function of p_T for the given selec-

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tions. Figures 6.12 and 6.13 show the impact of the selections in the $\frac{\sigma_{q/p}}{q/p}$ as a function of p_T plane for the baseline selection, the gradient WP, and the high- p_T WP.

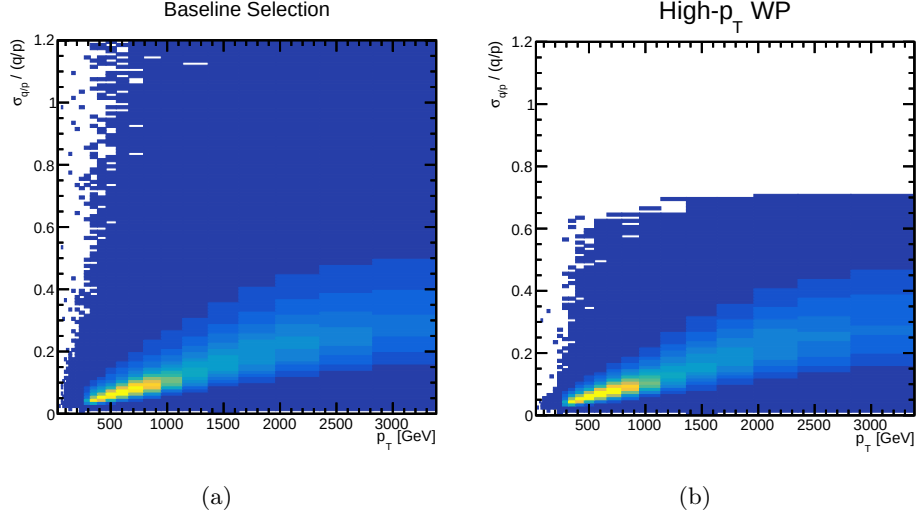


Figure 6.12: Impact of the baseline selection (left) and high- p_T WP (right) in the $\frac{\sigma_{q/p}}{q/p}$ as a function of p_T plane. While the baseline selection does not impose any requirements on the $\frac{\sigma_{q/p}}{q/p}$ as a function of p_T , the high- p_T WP clearly sculpts the distribution through the use of the BadMuon veto described in Section 6.1.

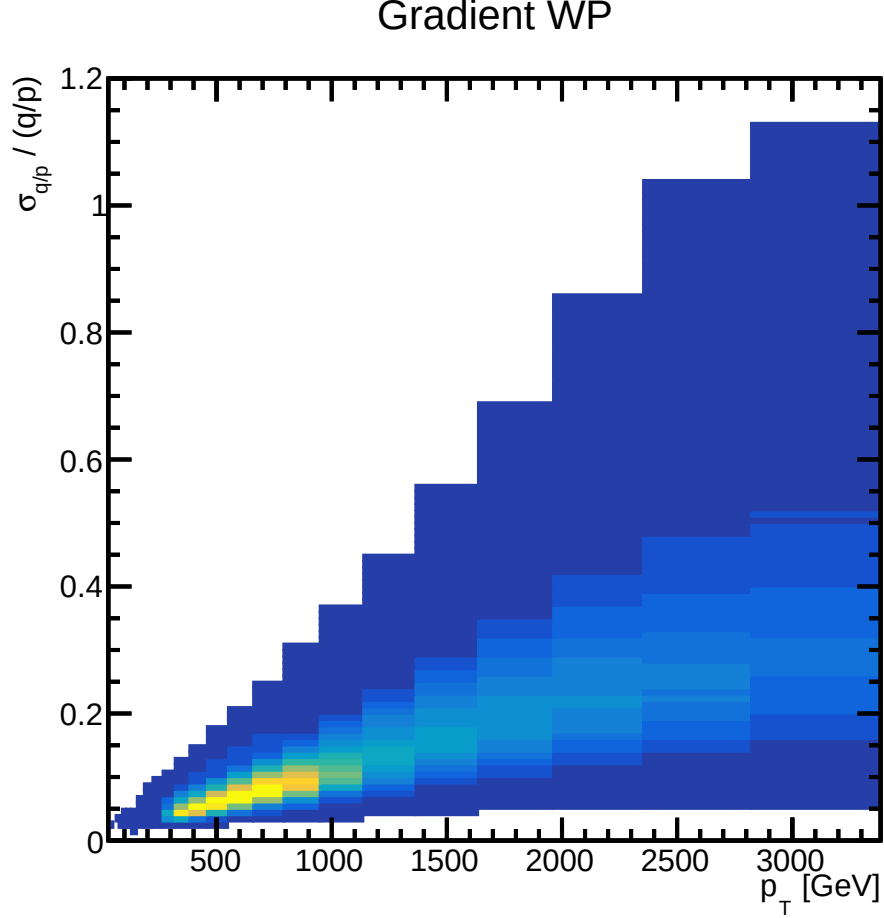


Figure 6.13: Impact of the gradient WP in the $\frac{\sigma_{q/p}}{q/p}$ as a function of p_T plane. As the muon p_T increases, larger values of $\frac{\sigma_{q/p}}{q/p}$ are allowed in order to accommodate the worsening p_T resolution of the detector.

The most important figures of merit are the selection efficiencies and p_T resolutions. Figures 6.14 and 6.15 show a comparison of the selection efficiency as a function of truth ϕ , η , and p_T . Figure 6.16 shows the absolute gain in efficiency of the gradient WP, including the integrity cuts, with respect to the high- p_T WP. A notable feature in Figure 6.14(b) is the lower selection efficiency for the gradient WP compared to the high- p_T WP in the region $|\eta| \sim 1.5$, corresponding to the region where the toroid magnetic field

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transitions from the barrel toroid to the endcap toroids. This is attributable to the fact that the relative error assigned to the q/p measurement in this region is larger, and as such, the 95% quantile cut on the relative q/p error will select fewer muons compared to the high- p_T WP due to the nature of the gradient cut. This phenomenon indicates that the new implementation of the q/p uncertainty using AEOTs discussed in Section 6.3.1 is performing as desired.

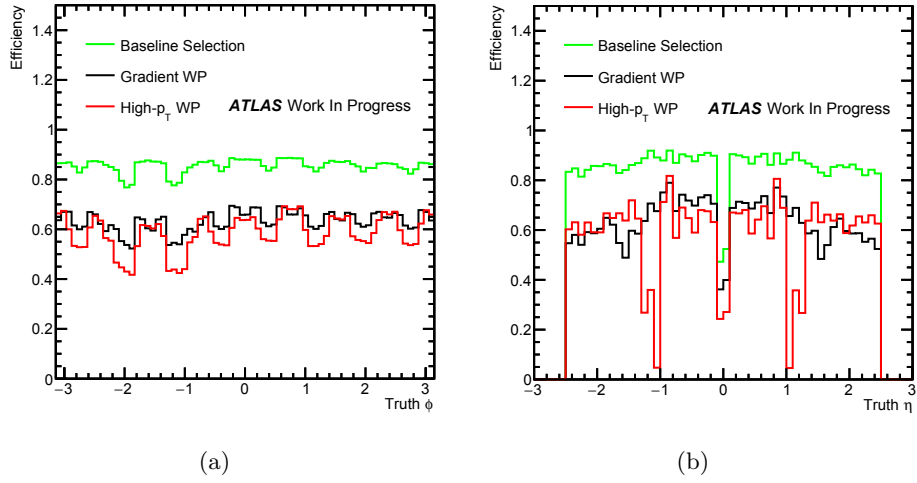


Figure 6.14: Efficiency as a function of truth ϕ (left) and truth η (right) for the baseline selection (green), gradient WP (black), and high- p_T WP (red).

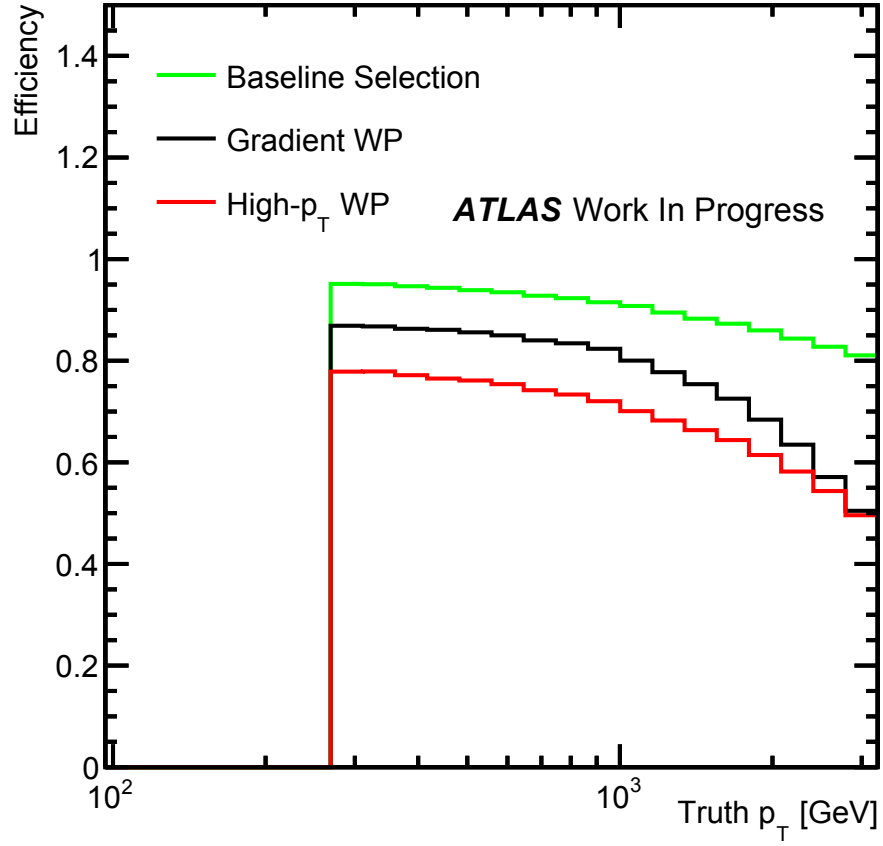


Figure 6.15: Efficiency as a function of truth p_T for the baseline selection (green), gradient WP (black), and high- p_T WP (red).

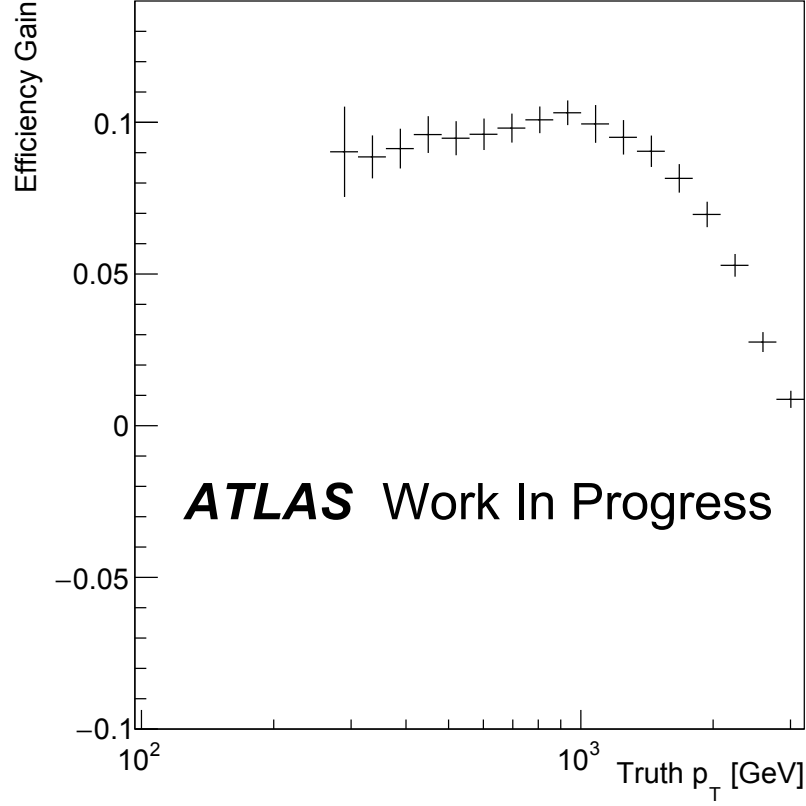


Figure 6.16: Absolute gain in efficiency of the gradient WP, which includes the integrity cuts, with respect to the high- p_T WP.

While Figure 6.17 illustrates an example of the q/p response distribution, Figures 6.18 and 6.19 show a comparison of the Root Mean Square (RMS) of the response distributions as a function of truth p_T for the baseline selection, the gradient WP, and the high- p_T WP. The chosen figure of merit is the RMS of the response divided by the mean of the ratio $\frac{q/p}{q/p^{\text{tru}}}$, as this quantity gives the relative resolution of the distribution in each bin of truth p_T , and avoids any momentum scale bias. The motivation to divide by the mean of the ratio $\frac{q/p}{q/p^{\text{tru}}}$ instead of taking the RMS of the response distribution as the resolution is as follows. If we scale q/p as $q/p \rightarrow f \times q/p$, which is unphysical and does not represent the correct momentum scale, but is

a possible scaling nonetheless, the mean of the ratio $\frac{q/p}{q/p^{\text{tru}}}$ will be f times the mean of the original ratio before the transformation, and the RMS will transform as q/p does. Thus, if we use the RMS as the resolution, it would be artificially higher or lower after this operation, depending on the exact value of f . When looking at the response RMS over the ratio mean, this effect is not present. The shift in the RMS is present whenever the mean of the ratio $\frac{q/p}{q/p^{\text{tru}}}$ is not equal to 1. This occurs, e.g., in each of the three curves in Figure 6.18(b).

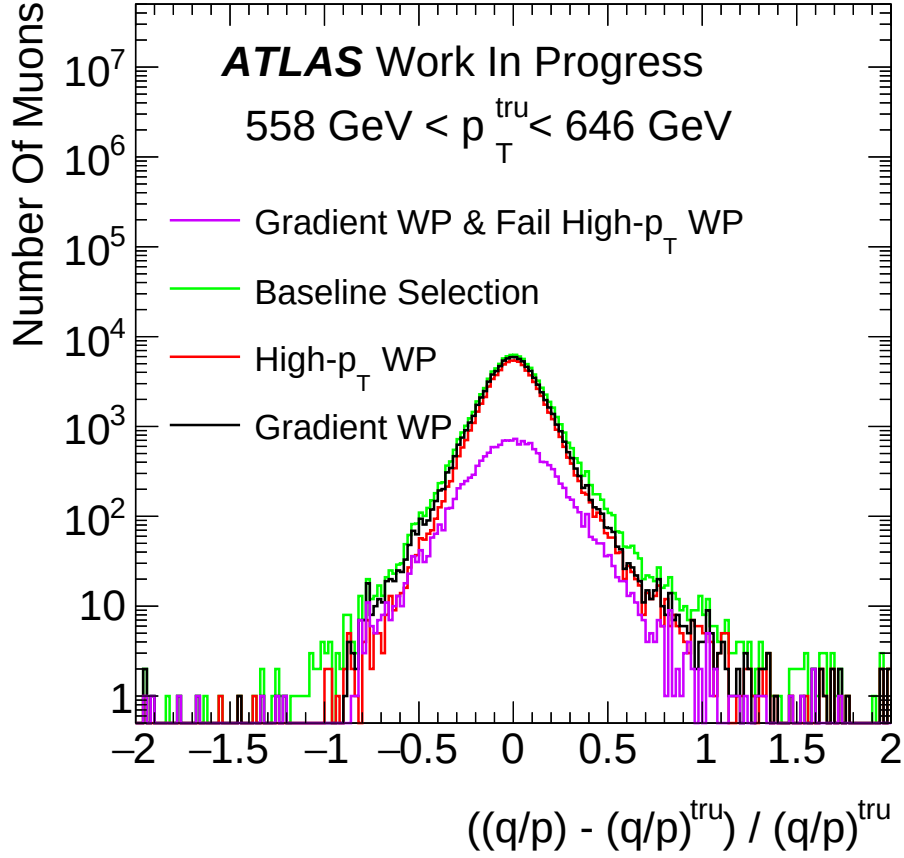


Figure 6.17: Example q/p response distribution for the baseline selection (green), high- p_T WP (red), and gradient WP (black). The purple line represents events passing the gradient WP, but failing the high- p_T WP.

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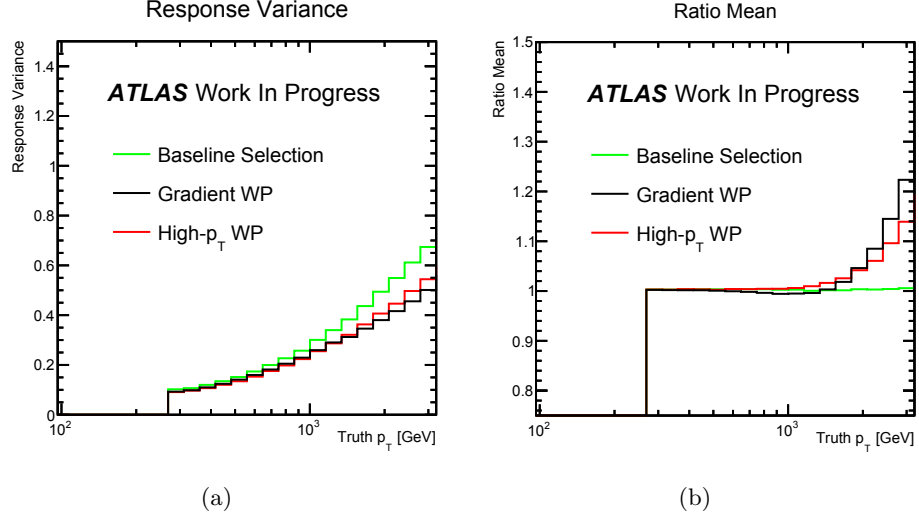


Figure 6.18: Variance of the response $\frac{|q/p - q/p^{\text{tru}}|}{q/p^{\text{tru}}}$ (left) and mean of the ratio $\frac{q/p}{q/p^{\text{tru}}}$ (right) used to compute the figure of merit shown in Figure 6.19.

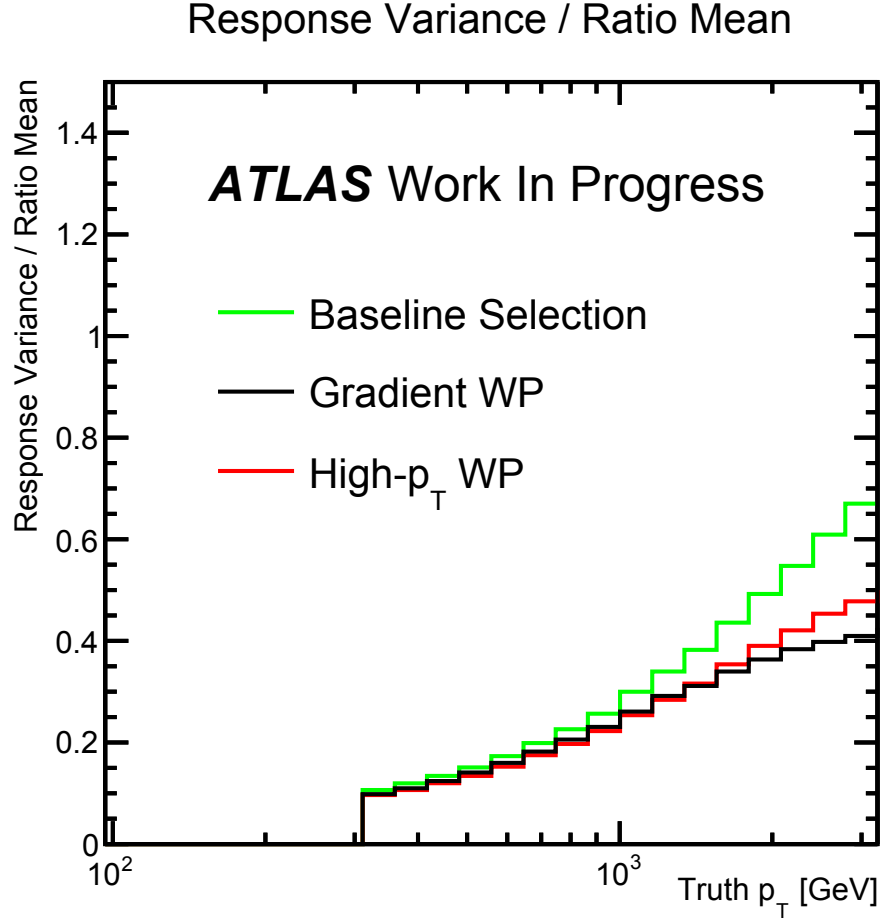


Figure 6.19: Resolution derived from the RMS of the truth-level to reconstruction-level p_T response distributions as a function of truth p_T . The figure of merit is the response variance over the ratio mean. A more precise estimation of the resolution, obtained with Gaussian fits to the response itself, in bins of p_T , is given in Figure 6.21.

Figure 6.20 shows a similar comparison to that of Figure 6.18, where the resolution is derived from fits to the response distributions in bins of truth p_T . The response distributions are fit with a Gaussian between the distribution mean minus $1.5 \times \text{RMS}$ and the distribution mean plus $1.5 \times \text{RMS}$. The width of the Gaussian fit, σ , as a function of truth muon p_T , is shown

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in Figure 6.20(a). The shifted mean of the Gaussian fit, $\mu + 1$, is shown in Figure 6.20(b). Given that $\mu + 1$, which is equivalent to the ratio $\frac{q/p}{q/p^{\text{tru}}}$, is not equal to 1 throughout the entire range for all three selections, the ratio of the width to $\mu + 1$ is taken in order to correct for different scales, as explained in the text above for the numerical RMS procedure. The corrected resolution is seen in Figure 6.21. The dip in resolution as the muon p_T gets very large is obtained by construction; the 95% quantile gradient cut selects more muons at high p_T values compared to the high- p_T WP. In turn, the residual distributions for higher values of p_T will have more events in the bulk of the distribution for the gradient WP, resulting in a better resolution. Indeed, by looking at Figures 6.12 and 6.13, it is clear that the gradient WP selects more muons at high p_T values compared to the high- p_T WP.

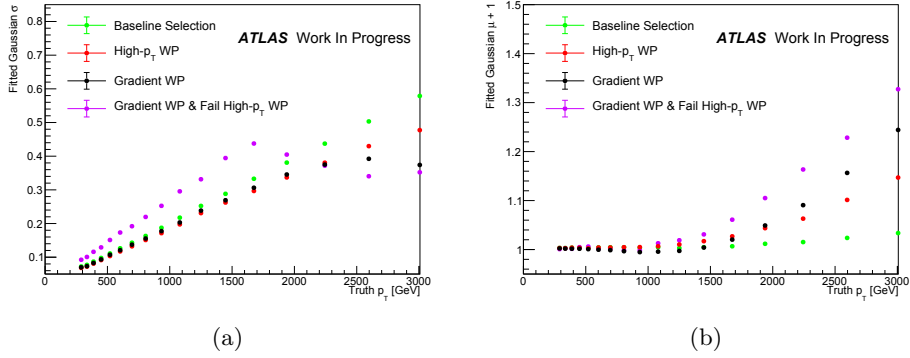


Figure 6.20: Width σ (left) and shifted mean $\mu + 1$ (right) obtained from the Gaussian fitting of the q/p response distributions in various p_T bins.

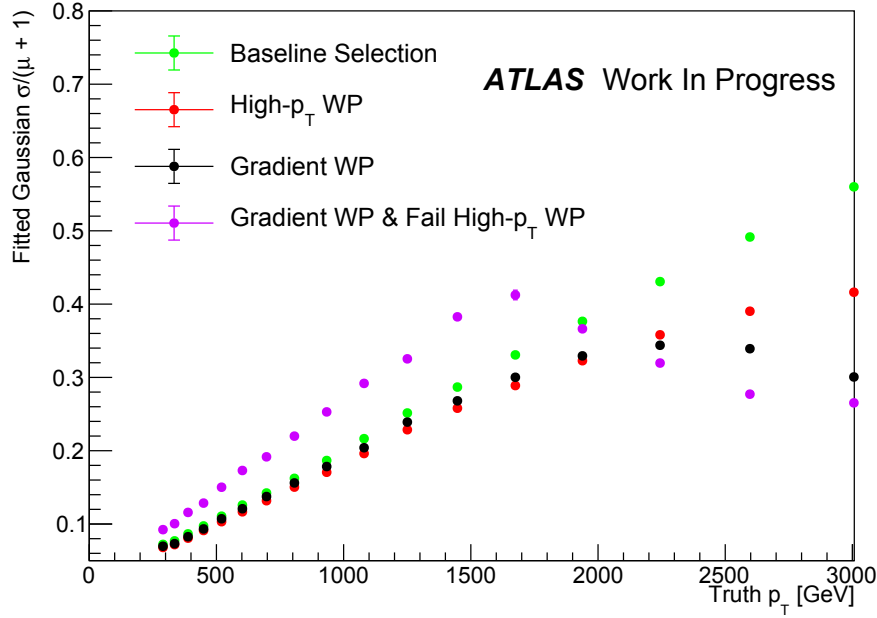


Figure 6.21: Corrected q/p resolution $\frac{\sigma}{\mu+1}$ obtained from the Gaussian fitting of the q/p response distributions in various p_T bins.

6.3.3 Summary

A new algorithm is presented as an alternative to the current high- p_T muon identification WP, which provides good resolution but is inefficient. The new algorithm comprises two independent components: a gradient cut and a set of integrity cuts. The gradient cut is applied in the previously vetoed regions of the high- p_T WP and in the regions outside of these vetoes, while the integrity cuts are applied only in the previously vetoed regions. Overall, the new algorithm delivers a better efficiency and similar resolution throughout the entire p_T range. The algorithm can be used by many other analyses in ATLAS other than the analysis presented in this dissertation. For example, the gradient WP is ideal for analyses which work at lower values of muon p_T and want to use a more robust selection WP compared to the *Medium* WP, but cannot use the high- p_T WP because of the associated efficiency loss. The new algorithm is thus a good compromise in terms of robustness and efficiency for these analyses.

Chapter 7

ATLAS Trigger System

Proton bunches collide within the ATLAS detector at a rate of 40 MHz. Recording all collision events is impossible, both from a data collection bandwidth standpoint and from a data storage capacity standpoint. Hence, the ATLAS experiment uses a two-level trigger system in order to select events of interest and reject events with little potential to be of interest from a physics point of view. This chapter introduces the two levels of the ATLAS trigger system: the Level-1 (L1) trigger is described in Section 7.1, and the High-Level Trigger (HLT) is described in Section 7.2. Particular consideration is given to the muon trigger system in Section 7.3, because of the relevance of muon triggers and their efficiencies to the analysis containing two high-momentum muons in the final state described in this dissertation.

7.1 Level-1 Trigger

The fast hardware-based L1 trigger enables the reduction of the event rate from 40 MHz to ~ 100 kHz. The L1 trigger uses detectors that are fast enough to accurately identify the bunch crossing of an event, thus allowing to filter out uninteresting events each time a proton-proton collision occurs. The information used by the L1 trigger is obtained from the trigger chambers of the MS and from the calorimeters with reduced-granularity; no ID information is used for the L1 trigger. The MS trigger chambers identify high- p_T muons, while the calorimeters allow the detection of electrons, photons, jets, τ leptons, and E_T^{miss} . The decision time allocated for the L1 trigger is $2.5 \mu\text{s}$. If an event passes the L1 trigger, one or several coarse Regions of Interest (RoIs) are defined with preliminary measurements in η and ϕ using a subset of detectors. Events passing the L1 trigger are sent to get processed by the software-based HLT.

7.2 High-Level Trigger

The HLT employs dedicated offline algorithms to reconstruct particles and estimate their parameters. Starting from the RoIs obtained by the L1 trigger, the full granularity of the calorimeters and the MS is exploited to measure the properties of the particles traversing the detector within the RoIs. The average processing time for objects found in the ROIs is 40 ms. Next, track matching with the ID is carried out. Finally, events are fully reconstructed at the HLT in approximately 4 s, and a final decision is made whether to write the event to permanent storage for offline analysis. The event rate after the HLT is reduced from the ~ 100 kHz accept rate coming from the L1 trigger to ~ 1 kHz. In cases where the accept rate of a particular trigger is still too large to handle after having passed the HLT, e.g., J/ψ meson production events, the trigger must be prescaled, i.e., only one in a number of events passing the trigger is kept. However, whenever possible, triggers that are not prescaled are used, in order to make the most of the valuable data produced by the LHC.

7.3 Muon Trigger Efficiency Measurements

Measuring the performance of high- p_T muon triggers used in the ATLAS experiment is vital for understanding the detector. Their performance needs to be evaluated in data to determine whether MC simulations are reliable enough for use in physics analyses. This is achieved by performing measurements of the trigger efficiencies in events containing a leptonically decaying W boson and jets (W +jets), and events with muons from top quark pair production ($t\bar{t}$). From these measurements, data-to-simulation correction factors, so-called Scale Factors (SFs), which will be applied to physics analyses to correct for any mismodelling of the triggers in simulation, can be derived. This is of direct relevance in searches for BSM physics, in particular in searches for high-mass phenomena, because the final states considered have two high-momentum muons in each event.

The method used to calculate the efficiencies of muon triggers is the so-called “tag and probe” method. This approach uses a “tag” trigger to select desired events. Once events pass the tag trigger, the various “probe” trigger efficiencies are evaluated in both data and MC by calculating the fraction of events passing both the tag and the probe triggers in the sample of events passing the tag trigger. For low- p_T (medium- p_T) muons, i.e., muons with $p_T \lesssim 10$ GeV ($10 \lesssim p_T \lesssim 100$ GeV), this method is used with events where

7.3. Muon Trigger Efficiency Measurements

a SM J/ψ meson (Z boson) decays into a pair of muons. In this case, one of the muons causes the tag trigger to record the event, and the other muon is used to measure the probe trigger performance without bias. For high- p_T muons, i.e., muons with $p_T \gtrsim 100$ GeV, two distinct types of events are utilized for the tag-and-probe method: $t\bar{t}$ events and W +jets events. The tag trigger used in this study is a trigger based on selecting events with high E_T^{miss} , defined in Equation (5.14). The scale factors are then calculated according to

$$\text{SF} = \frac{\epsilon^{\text{data}}}{\epsilon^{\text{MC}}}, \quad (7.1)$$

where ϵ^X is defined as the ratio of the number of probe muons matched to at least one trigger object to the total number of probe muons:

$$\epsilon^X = \frac{N_{\text{probe}}^X}{N_{\text{tag}}^X + N_{\text{probe}}^X}, \quad (7.2)$$

where $X \in [\text{data}, \text{MC}]$.

In the following, Section 7.3.1 presents the samples and event selection applied in order to perform efficiency measurements of high- p_T muon triggers. Section 7.3.2 describes the systematic uncertainties considered when calculating the trigger SFs. The efficiency measurements and trigger SF calculations are presented in Section 7.3.3.

7.3.1 Event Selection

Both the data and simulated samples undergo a series of cuts in order to select a high-purity sample of W +jets or $t\bar{t}$ events. Cuts are applied both at the event level and at the object level. The object-level requirements for the selected muons are that each pass the *Medium* quality WP with no isolation requirements. At the event level, data events are required to be in the so-called Good Run List (GRL) as defined by the ATLAS data quality group. An event is contained in the GRL if it occurred during a period where the proton beams in the LHC were stable and all relevant detector systems were functioning nominally. In addition, events must not be flagged as incomplete, and must pass LAr and Tile error requirements. This provides protection against noise bursts and data corruption in the EM and hadronic calorimeters, respectively. Each event must contain a primary vertex, defined as the vertex in the event with the highest sum of track p_T squared, and must pass either the tag or the probe trigger. The tag trigger

7.3. Muon Trigger Efficiency Measurements

used, HLT_xe100_L1XE60, requires at least 100 GeV of E_T^{miss} . The probe triggers, HLT_mu50 or HLT_mu26_ivarmedium, require at least one muon with $p_T > 50$ GeV or one isolated muon with $p_T > 26$ GeV, respectively. Events must contain exactly one muon with $p_T > 26$ GeV, and contain no electrons. The E_T^{miss} present in the event must also be greater than 200 GeV, in order to avoid the turn-on curve of the E_T^{miss} tag trigger. The requirements above can be applied in the context of high- p_T trigger efficiency studies for both $t\bar{t}$ and W +jets event topologies. Specific to the W +jets selection, events are required to have between 1 and 4 jets with $p_T > 25$ GeV inclusively, and to contain no b -jets, which are defined using the mv2c10 77% WP defined in Section 5.3.1. For the $t\bar{t}$ selection, events are required to have at least 4 jets with $p_T > 25$ GeV, and at least one b -jet. The b -jet requirement ensures the orthogonality of the W +jets and $t\bar{t}$ measurements. The samples used to study high- p_T muon triggers, and their associated MC generators, are listed in Table 7.1.

Process	Generator
W +jets	SHERPA 2.2.1 [67]
Z +jets	SHERPA 2.2.1 [67]
Diboson	SHERPA 2.1 [68]
Single Top	POWHEG [69] + PYTHIA 6 [70]
$t\bar{t}$	POWHEG [69] + PYTHIA 8 [71]
$t\bar{t}$ + V	MadGraph5/aMC@NLO [72] + PYTHIA 8 [71]

Table 7.1: Summary of the samples used to evaluate the performance of high- p_T muon triggers.

The collected Run 2 data are compared to the simulated samples in various kinematic regions, in order to validate the agreement between data and simulation. Figure 7.1 shows an example of a relevant kinematic distribution for events passing the tag trigger requirement and the offline E_T^{miss} cut of 200 GeV: the η distribution of the muons for the $t\bar{t}$ selection. The disagreement between data and simulation is within systematic uncertainties.

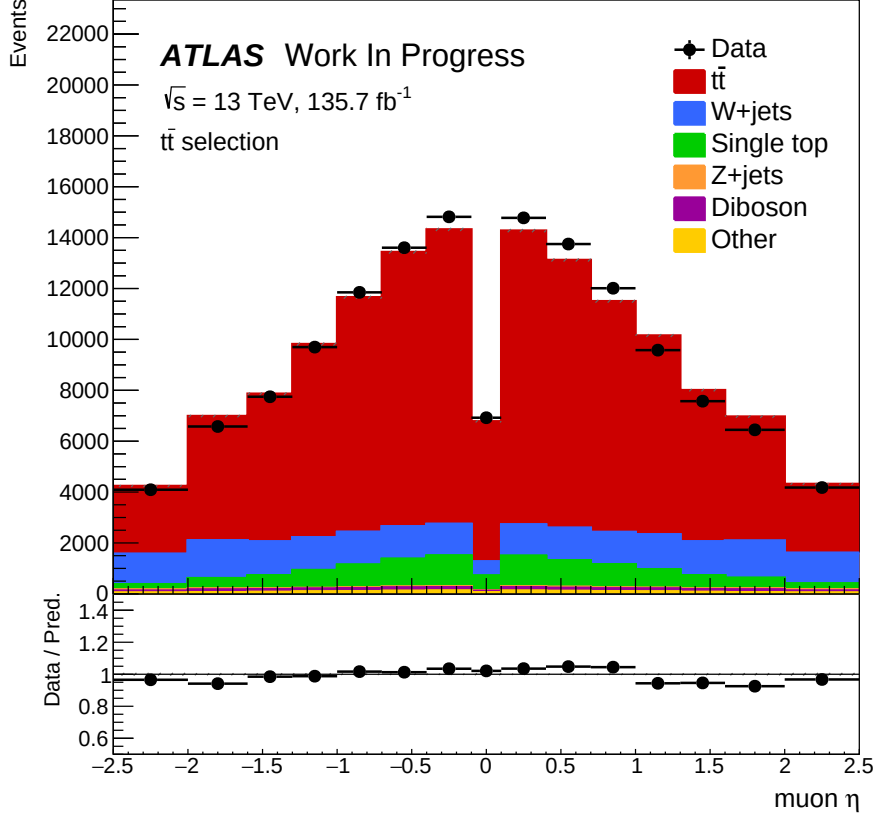


Figure 7.1: Muon η distribution for the $t\bar{t}$ selection. The disagreement between data and simulation is within systematic uncertainties.

7.3.2 Systematic Uncertainties

Various sources of systematic uncertainty were considered for the efficiency measurements presented in Section 7.3.3. All contributions are added in quadrature to obtain the total systematic uncertainties. In addition to the sources of systematic uncertainty described below, systematic variations on the smearing of the muon ID tracks, the smearing of the muon MS tracks, and the scale of the muon momentum were also studied and found to be negligible.

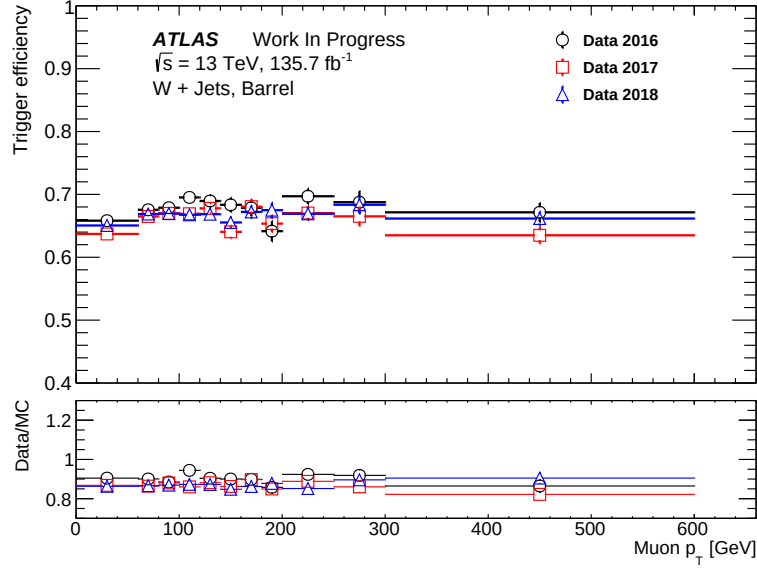
The identification of b -jets can affect the efficiency measurements be-

cause of the b -jet requirements described in Section 7.3.1. This effect is estimated by measuring the muon trigger efficiency while changing the object-level selection for b -jets from the `mv2c10 77% WP` to the `mv2c10 70% WP`. Depending on how the E_T^{miss} is reconstructed, different efficiencies can be obtained for the high- p_T muon triggers. The systematic uncertainty associated with the E_T^{miss} reconstruction is estimated by measuring the muon trigger efficiency while changing the E_T^{miss} event-level selection from 200 GeV to 150 GeV. Because of the large number of jets present in the final state considered for the efficiency measurements, the p_T cut applied to selected jets will affect the overall muon trigger efficiencies. By varying the p_T cut applied to selected jets from 25 GeV to 30 GeV, the systematic uncertainty related to the jet p_T cut is evaluated. Muon selection criteria are also considered when calculating systematic uncertainties. In particular, the muon isolation and selection WPs are varied. For the isolation, the systematic uncertainty is estimated by measuring the efficiency when applying the *FCTight* or the *FCTightTrackOnly* isolation [53]. The larger difference between either of these settings and the nominal efficiency, calculated with no isolation criteria, is used as an uncertainty. For the selection WP, the systematic uncertainty is estimated by changing the muon quality WP from the *Medium* WP to the high- p_T WP described in Section 6.1. The largest source of systematic uncertainty is the choice of the selection WP. All systematic uncertainties are below 3%.

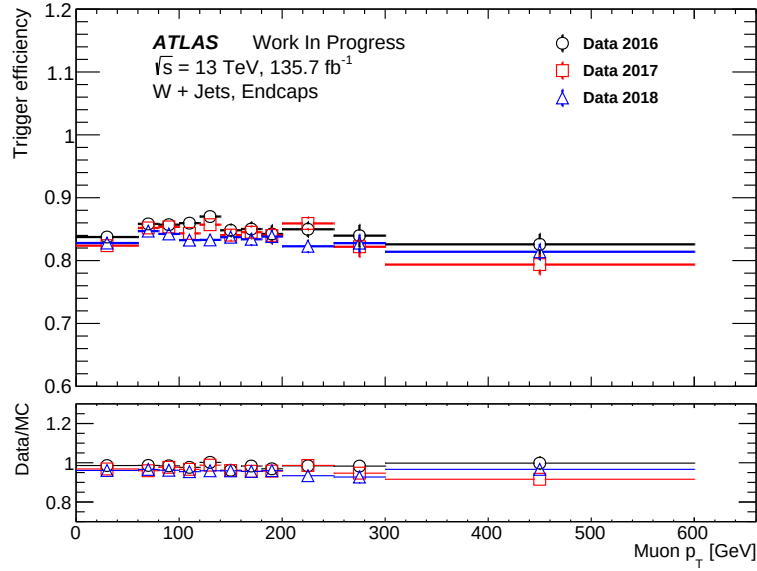
7.3.3 Efficiency Measurements and Trigger Scale Factors

The efficiency described in Section 7.3 is measured for the relevant high- p_T muon triggers used within ATLAS: `HLT_mu50` or `HLT_mu26_ivarmedium`. Sample distributions of trigger efficiencies as a function of muon p_T are shown in Figure 7.2 for the W +jets measurements and in Figure 7.3 for the $t\bar{t}$ measurements. In particular, the efficiencies plotted are those of a muon passing either the `HLT_mu50` trigger or the `HLT_mu26_ivarmedium` trigger. The final SFs are obtained by taking the ratio between the data and MC simulation efficiencies.

7.3. Muon Trigger Efficiency Measurements



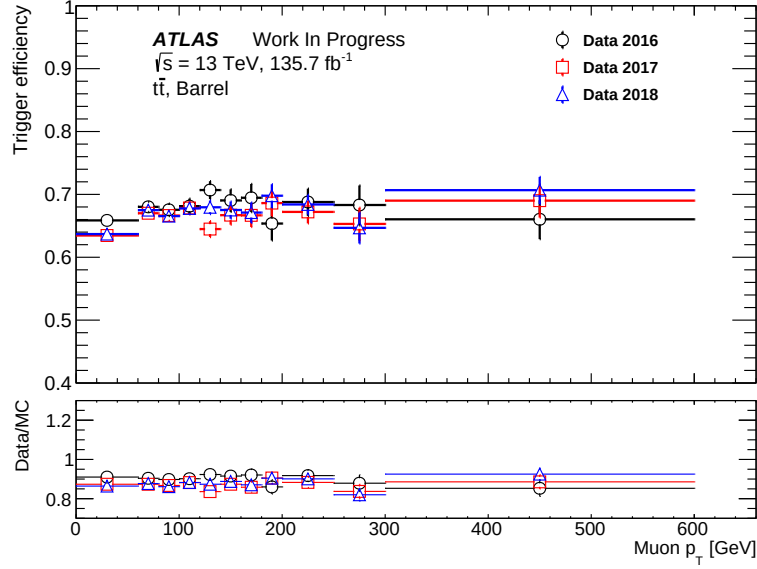
(a)



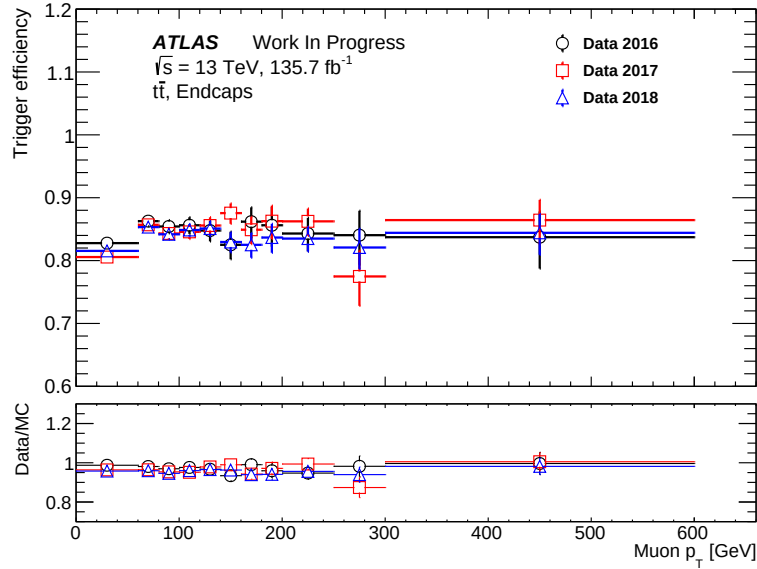
(b)

Figure 7.2: Trigger efficiencies for muons to pass either the HLT_mu50 trigger or the HLT_mu26_ivarmedium trigger in the barrel (top) and endcap (bottom) regions of the ATLAS detector as a function of muon p_T for the $W + \text{jets}$ selection. The ratio plot in the lower panels give the SFs.

7.3. Muon Trigger Efficiency Measurements



(a)



(b)

Figure 7.3: Trigger efficiencies for muons to pass either the HLT_mu50 trigger or the HLT_mu26_ivarmedmedium trigger in the barrel (top) and endcap (bottom) regions of the ATLAS detector as a function of muon p_T for the $t\bar{t}$ selection. The ratio plot in the lower panels give the SFs.

7.3. Muon Trigger Efficiency Measurements

Table 7.2 summarizes the SF calculations and their associated uncertainties. The total uncertainty is the sum in quadrature of the statistical and systematic uncertainties, with the systematic uncertainty being equal to the sum in quadrature of the various effects discussed in Section 7.3.2.

Year	Selection	Region	SF (%)	Unc. (%)	Stat. (%)	Syst. (%)
2016	$t\bar{t}$	Barrel	90.5	1.2	0.9	0.8
	$t\bar{t}$	Endcaps	96.7	2.6	0.9	2.5
	W +jets	Barrel	90.7	1.1	0.8	0.8
	W +jets	Endcaps	98.2	0.9	0.5	0.7
2017	$t\bar{t}$	Barrel	86.9	1.0	0.8	0.6
	$t\bar{t}$	Endcaps	96.7	1.1	0.8	0.7
	W +jets	Barrel	86.8	0.7	0.6	0.4
	W +jets	Endcaps	96.6	1.1	0.5	1.0
2018	$t\bar{t}$	Barrel	88.3	1.7	0.7	1.6
	$t\bar{t}$	Endcaps	95.8	1.7	0.7	1.5
	W +jets	Barrel	86.9	1.2	0.5	1.1
	W +jets	Endcaps	95.5	0.8	0.4	0.7

Table 7.2: Summary of the high- p_T muon trigger SF calculations for muons passing either the HLT_mu50 trigger or the HLT_mu26_ivarmedium trigger. The last two columns give the statistical and systematic uncertainties that comprise the total uncertainty.

Chapter 8

ATLAS Detector Upgrade

The LHC entered the Long Shutdown 2 (LS2) at the end of 2018, which will last approximately two years. During Run 2, the design instantaneous luminosity of $1 \times 10^{34} \text{ cm}^{-2}\text{s}^{-1}$ was attained, and surpassed, reaching a peak instantaneous luminosity of $2 \times 10^{34} \text{ cm}^{-2}\text{s}^{-1}$. Following the third run of data-taking of the LHC in 2021 – 2023, so-called Run 3, the instantaneous luminosity is expected to increase by up to a factor of three compared with the Run 2 conditions. The data-taking period after Run 3, the so-called High-Luminosity Large Hadron Collider (HL-LHC), promises to deliver an astounding 3000 fb^{-1} of proton-proton collision data, and is scheduled to start in 2026. To cope with these harsher luminosity conditions, a series of upgrades are foreseen to be implemented within the ATLAS detector. In particular, the upgrade programme consists of two main phases: Phase-I upgrades will be installed during LS2, and Phase-II upgrades will be installed during the third long shutdown. The largest Phase-II upgrade is the Inner Tracker (ITk) [73, 74], which will replace the existing ID with a new, entirely silicon-based, detector. The ITk will provide higher granularity and increased radiation hardness with respect to the current ID. Among others, anticipated Phase-II upgrades include the replacement of the electronics related to the trigger system for the LAr calorimeters [75], and the installation of additional RPC chambers in the MS [76]. These upgrades will allow the ATLAS detector to remain compatible with the higher trigger rates expected during the HL-LHC and gain efficiency in the barrel, respectively.

The largest Phase-I upgrade to the ATLAS detector is the replacement of the inner-most station of the MS, the Small Wheel. Described in Section 8.1, the NSW will be integrated into the ATLAS detector during LS2. One of the two technologies used in the NSW is sTGCs. This technology, and results of test beam campaigns undertaken with the goal of characterizing sTGCs, will be presented in Sections 8.2 and 8.3, respectively. The other detector technology employed in the NSW are Micromegas (MM). A detailed description of MM is beyond the scope of this dissertation.

8.1 The New Small Wheel

The L1 muon trigger rate in the forward region of the ATLAS detector is currently very high, and comprises approximately 90% fake muon triggers, as seen in Figure 8.1. With the upcoming increase in instantaneous luminosity, the situation will only get worse.

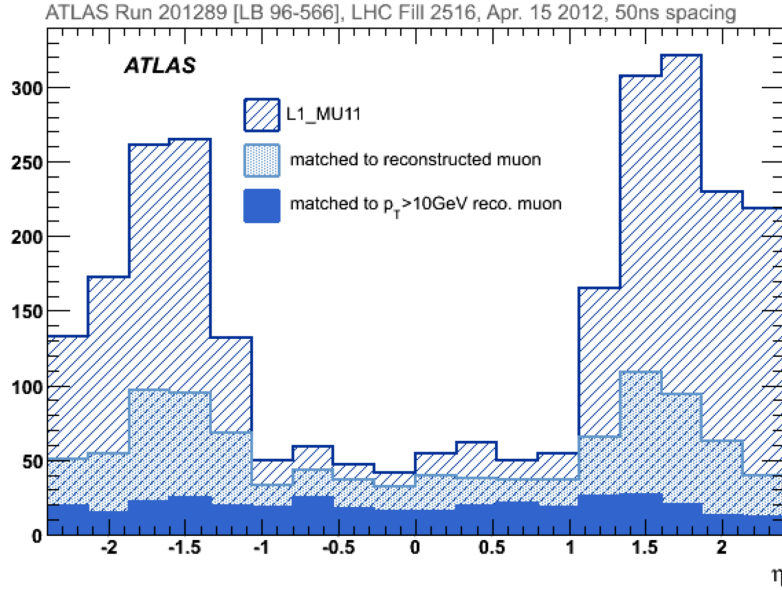


Figure 8.1: L1 muon trigger signal (hashed blue) as a function of η compared with the subset of matched muon candidates within a cone of $\Delta R = 0.2$ with $p_T > 3$ GeV (dotted blue) and offline-reconstructed muons with $p_T > 10$ GeV (solid blue) [77].

The NSW is an upgrade to the ATLAS detector, aimed at reducing fake muon triggers while maintaining an acceptable trigger bandwidth and improving the tracking performance in the forward region of the detector. In addition, the NSW will allow the p_T thresholds of the muon triggers to remain low. This is accomplished by providing an additional means of rejecting fake triggers coming, e.g., from activation of material in the ATLAS detector toroids. Figure 8.2 depicts a cross-sectional view of one quadrant of the ATLAS detector with the NSW installed. The current trigger system within ATLAS accepts all tracks in the diagram, i.e., tracks labeled A, B, and

C. The desired tracks that should fire the trigger are only the tracks labeled A. With the installation of the NSW, the required coincidence between the big wheel trigger and small wheel trigger will help to reduce the fake triggers coming from tracks labeled B and C. The tracks labeled B will be rejected because no corresponding track in the NSW is found. The tracks labeled C will be rejected because the corresponding NSW track does not point back to the interaction point.

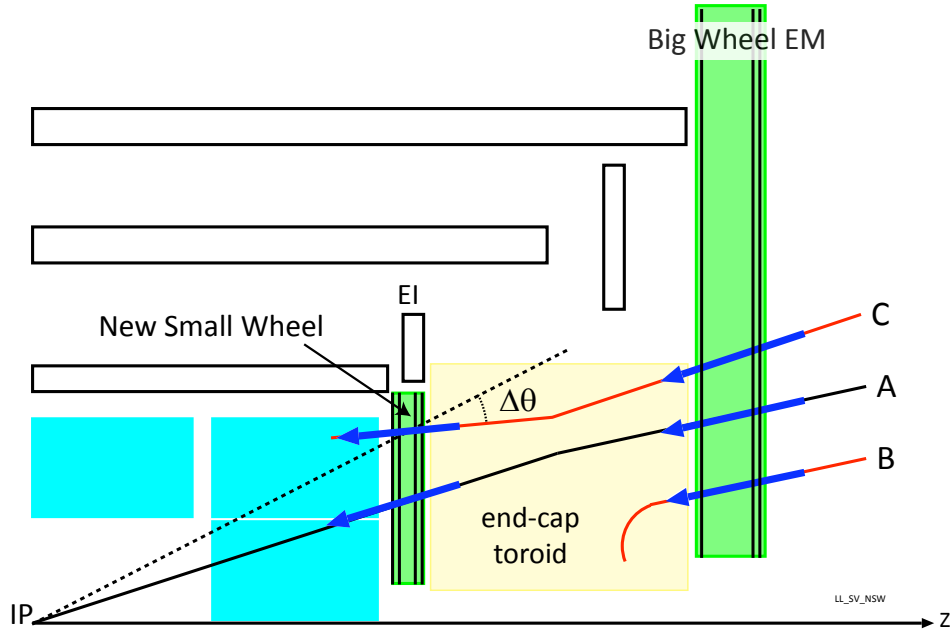


Figure 8.2: Cross-sectional view of one quadrant of the ATLAS detector with the NSW installed. Tracks labeled A will cause the trigger to fire, while tracks labeled B and C will not [77].

One NSW will be installed on either side of the interaction point. Each NSW consists of 16 detection layers: 8 sTGC layers and 8 MM layers. Starting from the interaction point, a particle will cross four layers of sTGCs (a quadruplet), eight layers of MM, and another four layers of sTGCs. Each sTGC-MM-sTGC grouping is referred to as a sector. There are two types of sectors: small and large sectors. Figure 8.3 depicts the arrangement of small and large sectors of the NSW detector components. The NSW will reconstruct muon tracks with high precision and provide information for the L1 trigger using sTGCs as the primary trigger system and MM as the primary

8.2. Small-strip Thin Gap Chambers

precision tracking detectors. The performance requirements for the NSW are 1 mrad angular resolution and 100 μm position resolution per detector layer.

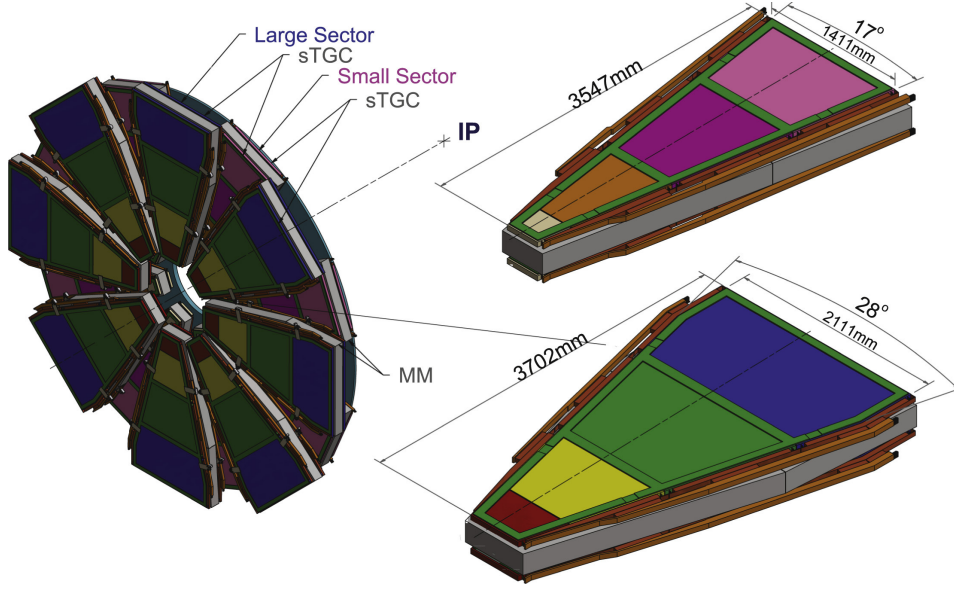


Figure 8.3: Schematic diagram of the small and large sectors comprising the NSW. Each sector consists of eight MM planes sandwiched between two sTGC quadruplets [78].

8.2 Small-strip Thin Gap Chambers

sTGCs are similar to multi-wire proportional chambers. The anode plane, realized with 50 μm gold-plated tungsten wires with a 1.8 mm pitch, lies between two cathode planes, at a distance of 1.4 mm, which have a resistive graphite layer to reduce sensitivity to field fluctuations and prevent sparking. One cathode plane uses thin readout strips with a 3.2 mm pitch. The other cathode plane uses readout pads to determine which strips to read out for the trigger. The wires are held at a voltage of approximately 3 kV. The chamber is filled with a gas mixture of 55% O_2 and 45% n-pentane. Charged particles passing through the sTGC ionize the gas and create an avalanche of electrons on the high-voltage wire, which induces a signal related to the charge deposited on the pads and strips. The internal structure of sTGCs is shown in Figure 8.4.

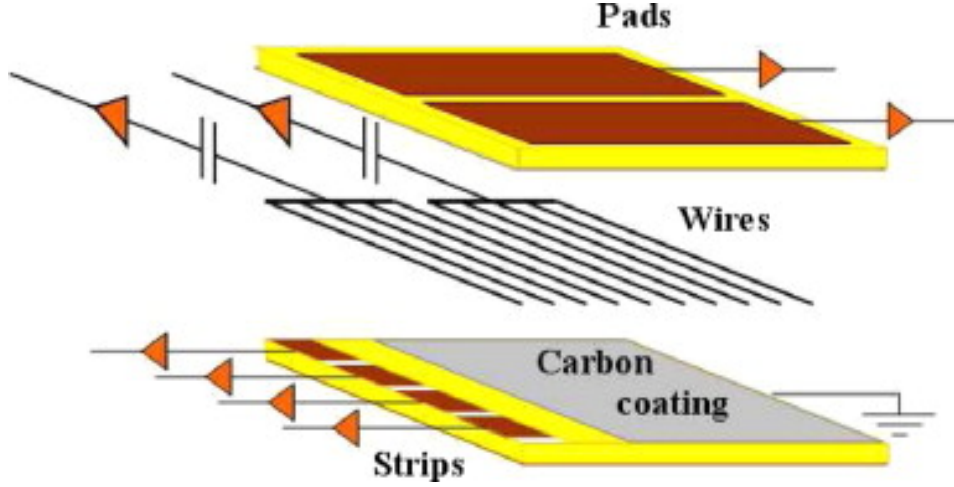


Figure 8.4: Internal structure of a small-strip Thin Gap Chamber. The high-voltage anode wires are sandwiched between a strips cathode plane and a pads cathode plane, which have a resistive graphite layer to reduce sensitivity to field fluctuations and prevent sparking [77].

Four individual chamber layers, or gaps, are assembled into quadruplets, or modules. The naming convention for sTGC modules follows the convention “QXY”, where $X \in [S, L]$ refers to small (S) or large (L) sectors, and $Y \in [1, 2, 3]$ refers to the distance from the beam pipe, with 1 being closest to the beam pipe and 3 being farthest away from the beam pipe. Each module has a trapezoidal shape, as shown in Figure 8.3, with surface areas between 1 m^2 and 2 m^2 .

8.2.1 Readout Electronics

The readout chain of sTGCs comprises multiple electronic components. An Application-Specific Integrated Circuit (ASIC), the so-called VMM chip, is used as the first step in the readout chain. Multiple versions of the VMM ASIC are used throughout the test beam campaigns discussed in this chapter. The VMM chips are placed on Front-End Boards (FEBs), and are used to amplify and digitize the hits coming from the pads, strips, and wires of the sTGC chamber. The data from the FEBs are then sent to the Data Acquisition (DAQ) system through Twinax cables. The DAQ system handles triggering, configures the FEBs, reads all data, and forwards hit data to computerized storage for offline analysis. Finally, event building is done offline after all hits are recorded.

8.3 Test Beam Results

In order to characterize the sTGC modules and understand their performance and limitations, a series of test beam campaigns are carried out. The goals of the campaigns vary, but in general the tracking and timing resolution are important quantities to measure. The readout electronics of the chambers must also be understood, as several versions of the readout chips are produced. Section 8.3.1 details the test beam campaign carried out at Fermilab, and Section 8.3.2 presents results obtained during subsequent test beam campaigns carried out at CERN.

8.3.1 Fermilab Test Beam Campaign

During the summer of 2014, a test beam campaign was undertaken at Fermilab, and aimed at characterizing the spatial resolution of a prototype QS2 sTGC detector built in Israel [78]. The prototype consists of four sTGC strip and pad layers. The position resolution measurements are done with the strip layers. A 32 GeV pion beam is provided by Fermilab and used to irradiate the sTGC detector prototype. The detector is placed between the two arms of an external silicon pixel telescope with intrinsic resolution of $5.3\mu\text{m}$. Each arm of the pixel telescope comprises three pixel detector planes, for a total of six pixel detector planes overall. Figure 8.5 gives a schematic overview of the experimental setup used during the Fermilab test beam campaign and the coordinate system used for the ensuing analysis. In particular, the positive z-axis is in the same direction as the incoming beam, the y-axis is perpendicular to the strips in the chamber, and the x-y plane corresponds to the surface of the sTGC strip layer under study.

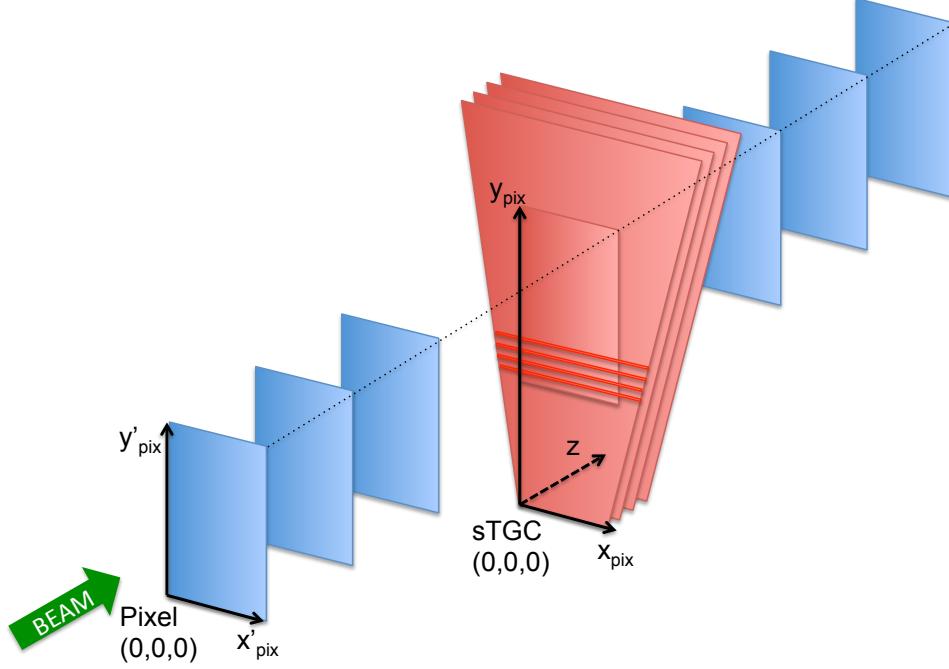


Figure 8.5: Schematic overview of the experimental setup used during the Fermilab test beam campaign. The blue planes correspond to the pixel detector telescope, and the red planes correspond to the four sTGC layers of the quadruplet. The dimensions are not to scale [78].

Strip clustering is done by fitting a Gaussian to the charge distribution coming from the set of strips having a signal above threshold and using the mean of the fitted Gaussian as the position of the cluster, $y_{sTGC,0}$. The strip-cluster multiplicities considered in this analysis are 3, 4, and 5. In addition to alignment corrections, a correction accounting for differential non-linearity effects in the detector is applied, based on the number of strips in each cluster. The need for this correction stems from the discrete nature of the strips; the spatial coordinate measured by the sTGC results from a discrete sampling of measured charge on the strips. The correction is derived by fitting the residual distribution with respect to the pixel telescope track as a function of relative y position, measured in number of strip pitches. The corrected position y_{sTGC} is then given by

$$y_{sTGC} = y_{sTGC,0} - a_i \sin \left(2\pi y_{sTGC,0}^{\text{rel}} \right), \quad (8.1)$$

where $y_{sTGC,0}$ is the uncorrected position of the cluster, $y_{sTGC,0}^{\text{rel}}$ is the

8.3. Test Beam Results

relative position of the cluster with respect to the closest inter-strip gap centre, and the values of a_i depend on the number of strips in the cluster, with $i \in [3, 4, 5]$ corresponding to the strip-cluster multiplicity. The values of a_i are obtained by fitting the residual distribution as a function of $y_{\text{sTGC},0}^{\text{rel}}$ with a sine wave and extracting the fitted amplitudes for the various strip-cluster multiplicities. Figure 8.6 shows the residual distribution as a function of $y_{\text{sTGC},0}^{\text{rel}}$ before and after applying the differential non-linearity correction. Table 8.1 summarizes the amplitudes used for the differential non-linearity corrections applied to the strip-cluster sizes considered in the position resolution analysis. As all amplitudes are consistent, a universal correction is applied using the amplitude for the strip-cluster multiplicity of 3 for all strip-cluster multiplicities.

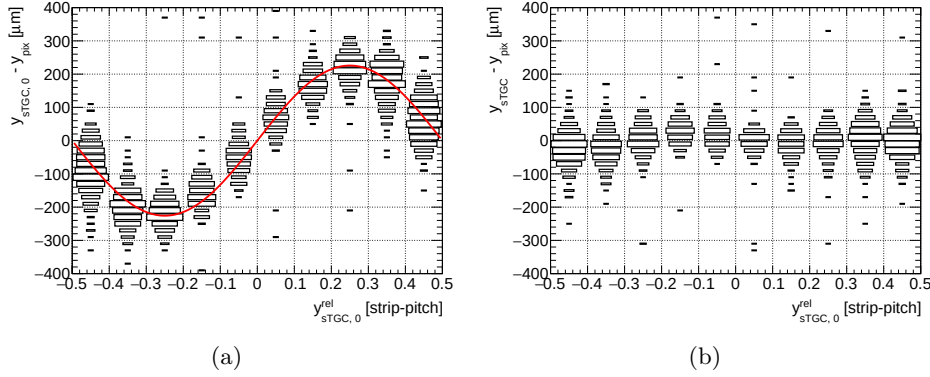


Figure 8.6: Residual distribution as a function of $y_{\text{sTGC},0}^{\text{rel}}$ before (left) and after (right) applying the differential non-linearity correction [78].

Strip-Cluster Multiplicity i	Amplitude Parameter a_i
3	205 ± 9
4	206 ± 4
5	211 ± 5

Table 8.1: Amplitudes of the differential non-linearity correction applied to the three strip-cluster sizes considered in the position resolution analysis.

In order to avoid multiple scattering effects within the sTGC chambers, only straight tracks are selected for the analysis. This is accomplished by requiring that the track quality parameter of the selected pixel telescope

8.3. Test Beam Results

tracks be less than 10. The track quality parameter is based on the χ^2 of the pixel telescope track fit, and is depicted in Figure 8.7

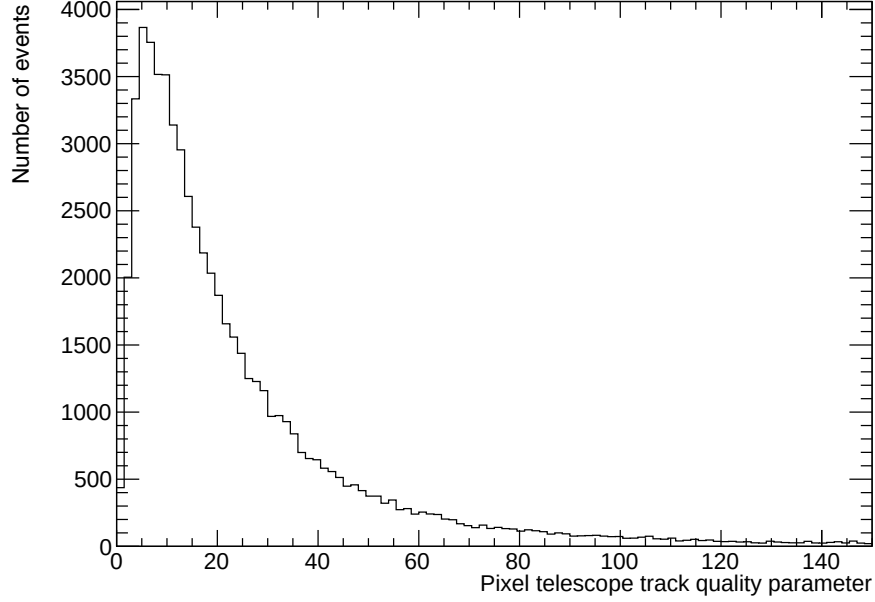


Figure 8.7: Track quality parameter of pixel telescope tracks. Tracks used in the analysis are required to have a track quality parameter less than 10.

The main goal of the test beam campaign at Fermilab is to measure the position resolution of the first full-size sTGC prototype detector. The intrinsic position resolution for each sTGC layer is estimated in two ways.

First, the extrapolated beam particle trajectory reconstructed by the telescope, y_{pix} , is compared with the measurements in each of the four sTGC quadruplet planes. A Gaussian model is adopted to fit the residual distribution between the sTGC position measurement and the pixel telescope position measurement, and the width of the Gaussian fit, σ , is quoted as the resolution of the sTGC layer. Figure 8.8(a) shows the residual distribution obtained with the pixel telescope analysis for a representative run of data-taking.

The second position resolution measurement uses the difference between two independent measurements of the beam particle position determined in two different sTGC layers. This standalone analysis assumes that two

8.3. Test Beam Results

sTGC strip layers are identical. The differential non-linearity correction derived from the pixel telescope analysis is also applied to the standalone analysis. The width of the residual distribution of pairwise sTGC strip layer position measurements, σ_{pair} , is then used to compute the position resolution of a single layer, $\sigma = \sigma_{\text{pair}}/\sqrt{2}$. Figure 8.8(b) shows the residual distribution obtained with the standalone analysis for a representative run of data-taking.

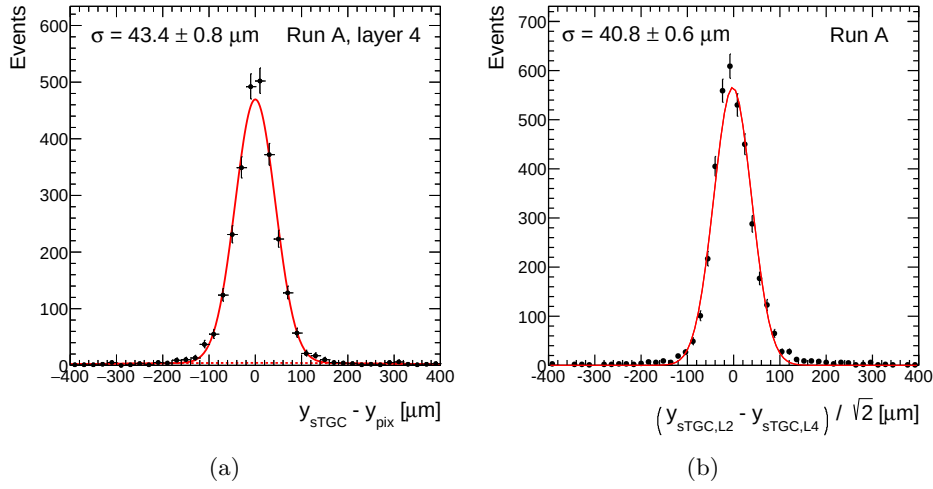


Figure 8.8: Residual distributions for the pixel analysis (left) and the standalone analysis (right). The intrinsic sTGC detector position resolution is quoted as the width of the fitted Gaussian, σ .

The data collected throughout the test beam campaign are analyzed by both methods. Data-taking runs are labeled by a letter corresponding to a particular location on the sTGC quadruplet where the beam is impinging for that data-taking run. Figure 8.9 depicts the positions of the pion beam for the various runs of data-taking during the test beam campaign. Runs G to N do not have synchronized data between the sTGC chamber and the pixel telescope, so only standalone analysis residual measurements are available for these runs.

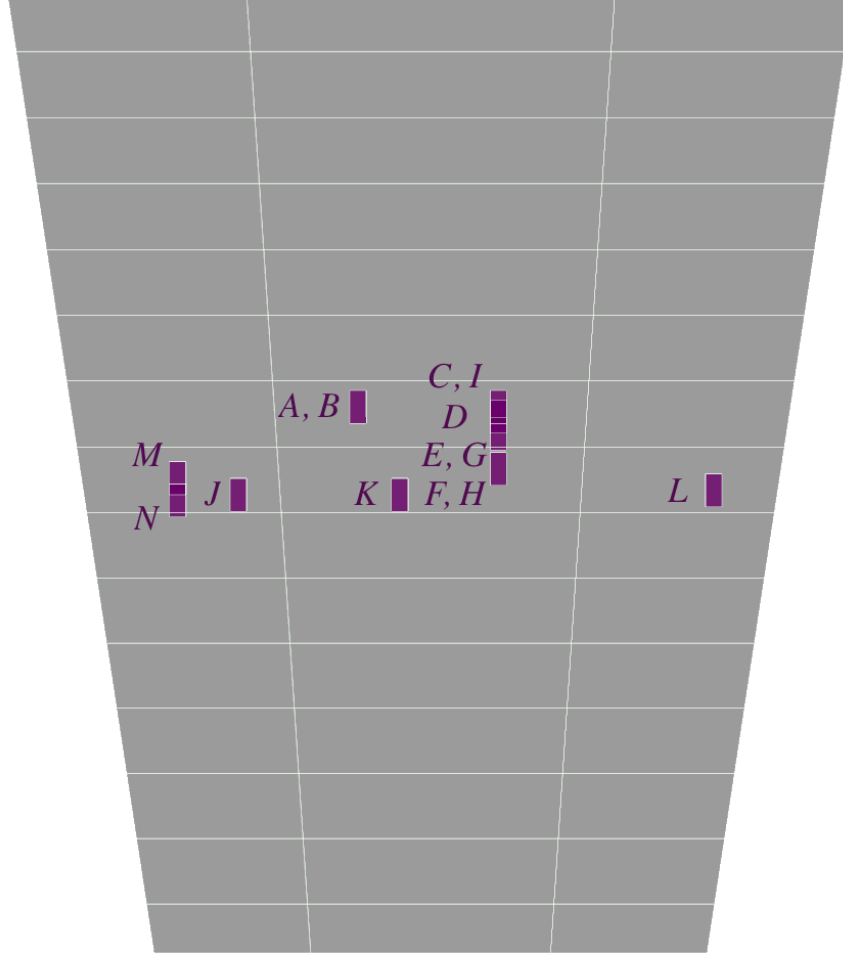


Figure 8.9: Summary of the pion beam position for the Fermilab data-taking runs analyzed. Runs G to N do not have synchronized data between the sTGC chamber and the pixel telescope, so only standalone analysis residual measurements are available for these runs [78].

Figure 8.10 shows the intrinsic sTGC resolution σ for various periods of data-taking during the Fermilab test beam campaign for both the pixel telescope analysis and the standalone analysis. The black markers indicate points where data are taken with no expected degradation due to sTGC detector support structure or calibration. The white markers indicate the

8.3. Test Beam Results

resolution measurements for the remainder of the runs, and are not included in the calculation of the average position resolution. In both the pixel telescope and sTGC standalone analysis, the average position resolution is about $45\text{ }\mu\text{m}$, with an RMS spread of $8\text{ }\mu\text{m}$, well below the $100\text{ }\mu\text{m}$ required by ATLAS for the NSW upgrade project.

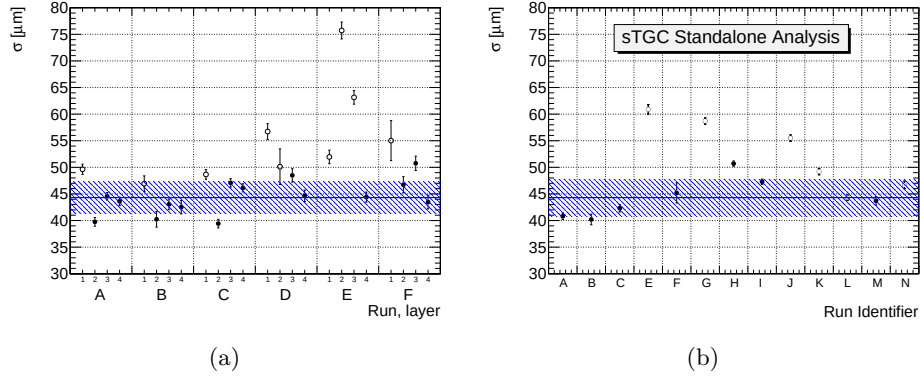


Figure 8.10: Intrinsic sTGC position resolution measured using the pixel telescope analysis (left) and standalone analysis (right), for various periods of data-taking during the Fermilab test beam campaign. Results for runs with no expected degradation due to sTGC detector support structure or calibration are shown as black filled circles. The horizontal line represents the average resolution for these runs, and the hashed band represents the RMS spread. Results for the remaining runs are shown as open circles.

8.3.2 CERN Test Beam Campaigns

Following the initial test beam campaign described in Section 8.3.1, a series of subsequent test beam campaigns was carried out at CERN. In particular, a campaign with a QS2 module was carried out in October 2017, and a campaign with a QL1 module was carried out in September 2018. Both campaigns were held at the H8 beamline at CERN. The H8 beamline provides 150 GeV muons to the sTGC module under test. While conceptually similar to the experimental setup of the test beam campaign described in Section 8.3.1, the experimental setup used at CERN does not include a pixel telescope, as shown in Figure 8.11. Triggering is performed by requiring the coincidence between two scintillators placed before and after the sTGC chamber under test.

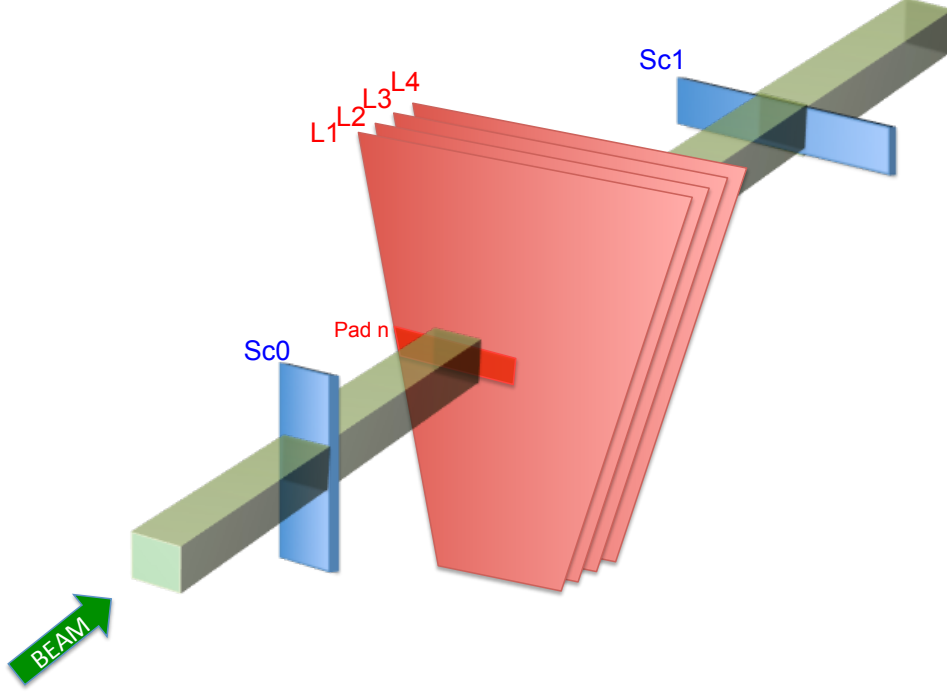


Figure 8.11: Schematic overview of the experimental setup used during the CERN test beam campaigns. The blue planes correspond to scintillators used for triggering, and the red planes correspond to the four sTGC layers of the quadruplet. The dimensions are not to scale [78].

The coordinate system used for the analysis of the test beam data is such that the positive z -axis is in the same direction as the incoming beam, the y -axis is perpendicular to the strips in the chamber, and the x - y plane corresponds to the surface of the sTGC strip layer under study.

The main goal of these test beam campaigns, in addition to measuring the spatial resolution of the prototype detectors, is to test the operation of full-sized sTGC quadruplets with complete 4-layer VMM readout from pads and strips. In particular, the efficiency of the pads is investigated, and the performance of sTGC chambers is characterized under high-background environments. Two distinct clustering algorithms are employed during the offline analysis of the test beam data. The first clustering algorithm is similar to the one described in Section 8.3.1, where the strip charge distribution is fit with a Gaussian function, and the mean of this Gaussian is taken as the y_{sTGC} position. The second clustering algorithm uses the centroid of the

strip charge distribution, defined as

$$y_{\text{sTGC}} = \frac{\sum_{i=0}^{N_{\text{hits}}} C_i y_i}{\sum_{i=0}^{N_{\text{hits}}} C_i}, \quad (8.2)$$

where N_{hits} is the total number of hits in a given sTGC layer, and C_i is the charge of the i^{th} hit with position y_i . No major difference is observed between the two clustering algorithms. As no pixel telescope is available to calculate the correction parameters, the differential non-linearity correction is not applied in these test beam campaigns. Figure 8.12 shows the beam profile recorded by the QL1 module pads. In particular, the beam is concentrated on one pad within the sTGC layer.

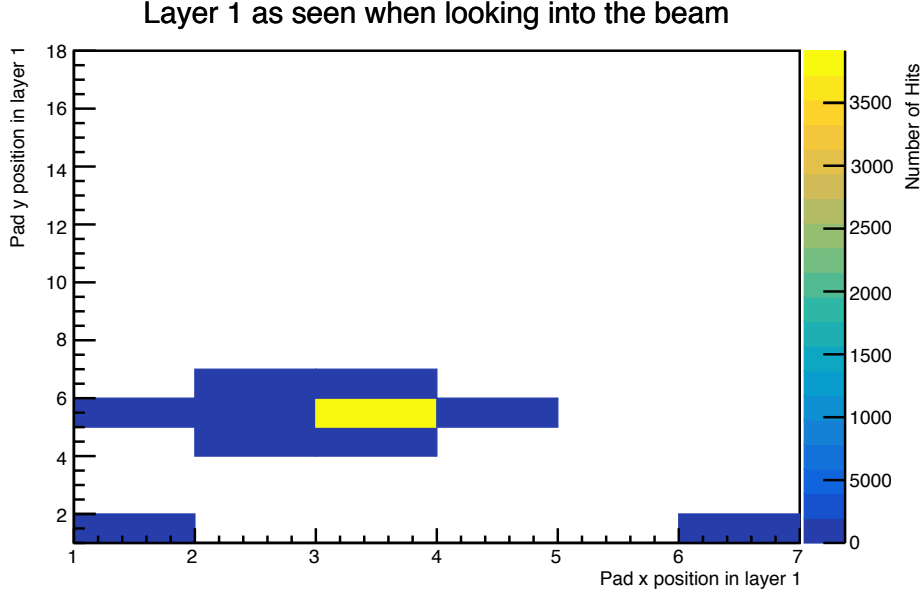


Figure 8.12: Beam profile from the QL1 module pads operating at 2.8 kV. The beam is concentrated on one pad within the sTGC layer. The other pads with content 0 are instrumented with readout electronics, but do not register any hits. The pads with no entries (white) are not instrumented with readout electronics.

The same beam can also be profiled using the sTGC strips, as shown in Figure 8.13. The dips in the profile at strip y positions of 220 and 280 are caused by the spacer buttons within the sTGC chamber.

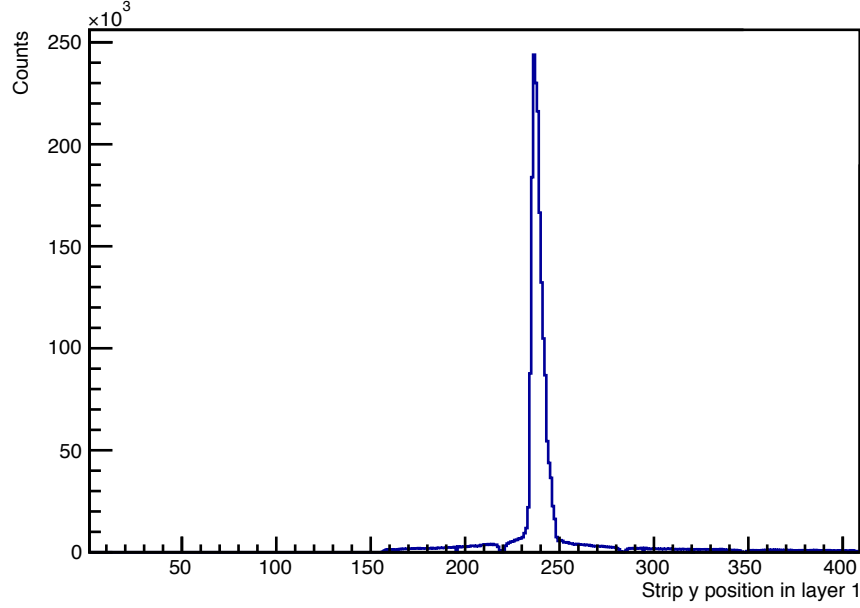


Figure 8.13: Beam profile from the QL1 module strips operating at 2.9 kV. The dips in the profile at strip y positions of 220 and 280 are caused by the spacer buttons within the sTGC chamber.

The signals induced by charged particles traversing the sTGC chamber are related to the charge deposited within the chamber by the particles. The charge spectrum should follow a Landau distribution [79]. Figure 8.14 shows the charge spectrum recorded by the QL1 module pads.

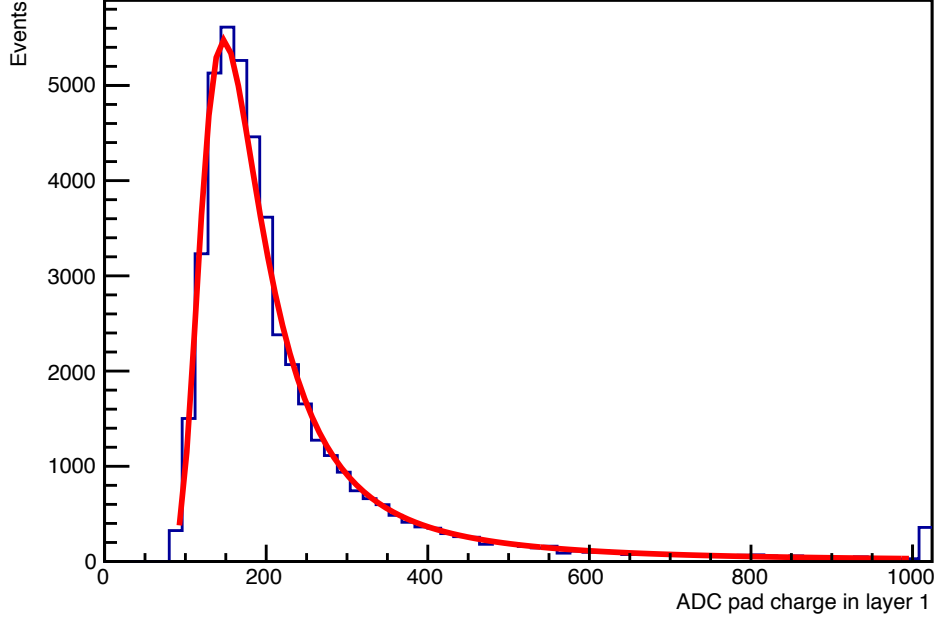


Figure 8.14: Charge spectrum from the QL1 module pads operating at 2.8 kV. The red line corresponds to an overlaid Landau distribution.

In order to reconstruct tracks, a feature extraction technique used in image analysis, the Hough transform [80, 81], is employed. In particular, tracks are parameterised as straight lines with two parameters:

$$y_{\text{sTGC}} = mz + b, \quad (8.3)$$

where the two parameters are the slope m and the intercept b of the line. For each layer in the quadruplet, the y_{sTGC} and z positions are measured. From these measurements, all possible values of the slope m and intercept b which parameterise a line compatible with the y_{sTGC} and z measurements for a given sTGC layer are iterated over, and the so-called accumulator is filled with the corresponding (b, m) pair. Each point in the accumulator space corresponds to a line going through the measured y_{sTGC} and z positions for a given sTGC layer. The point in the accumulator space where the lines of all four sTGC layers intersect corresponds to the values of the (b, m) pair of the straight line passing through all four measured points. Figure 8.15 illustrates the conceptual procedure of the Hough transform using the (θ, r) parameterisation of a straight line instead of the (b, m) parameterisation

used in this analysis:

$$r = z \cos \theta + y_{\text{sTGC}} \sin \theta. \quad (8.4)$$

Each coloured line represents a particular combination of θ and r that is compatible with the measured values of y_{sTGC} and z . Possible straight lines (coloured lines) are iterated over for each data point (black points), in order to fill an accumulator. Clearly, the line best describing all three black points, i.e., independent measurements of y_{sTGC} and z , is the blue line.

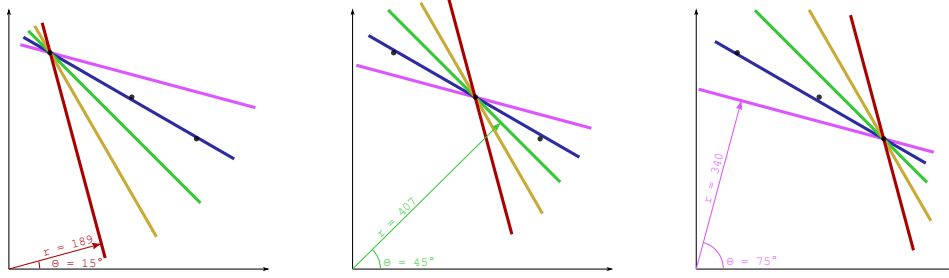


Figure 8.15: Illustration of the Hough transform technique used to reconstruct tracks during the sTGC test beam campaigns. Each diagram (left, centre, right) represents the various straight line possibilities (red, yellow, green, blue, magenta) for a particular measurement of y_{sTGC} and z (black points) [82].

8.3.3 Pulser Board Tests

As mass production of sTGC modules progresses, it is imperative to test the quality of each detector module. In particular, it is crucial to test the electrical connectivity and functionality of all readout elements of a quadruplet: pads, strips, and wires. This is accomplished by using the so-called pulser board. The experimental setup used at the Weizmann Institute of Science in Israel in order to carry out pulser board tests is shown in Figure 8.16.

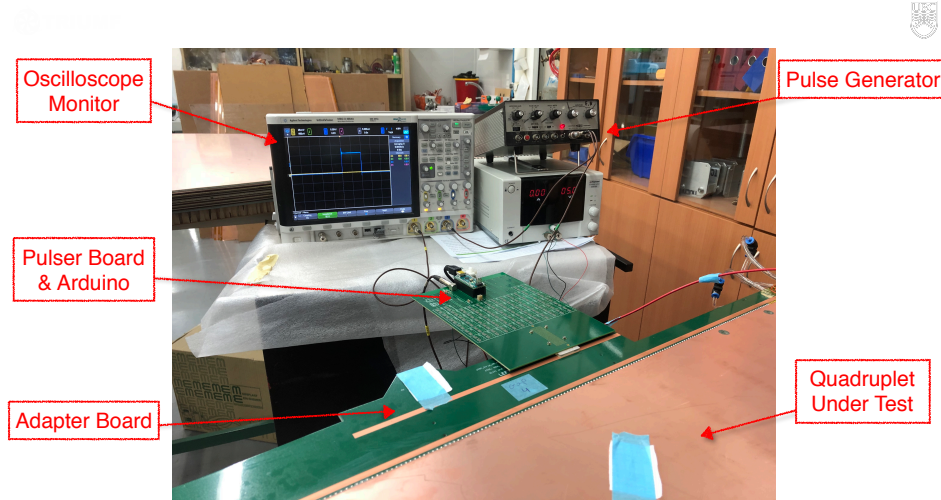


Figure 8.16: Experimental setup used for pulser board tests at the Weizmann Institute of Science.

A high-voltage signal is injected into the chamber. The signal is a square wave with a peak-to-peak amplitude of 20 V. The signal propagates through the chamber, and the response signal from the external adapter boards is measured. The signals coming from the various quadruplet elements are processed through the pulser board and digitized with an oscilloscope. As the analog signal coming from the chamber is inherently noisy, the digitized signal produced by the oscilloscope remains noisy, as shown in Figure 8.17(a). The pulser board software includes an averaging function that smooths the noisy waveform from the oscilloscope and produces a cleaner signal for analysis use, as shown in Figure 8.17(b). The pulser board tests provide the ability to detect bad channels in the sTGC module by distinguishing good pulses from pulses such as the one shown in Figure 8.17(c) and reporting this to the user. In particular, the amplitude, mean, and variance of each channel's pulse is computed, and good channels will tend to have clear regions of high variance and large amplitude.

8.3. Test Beam Results

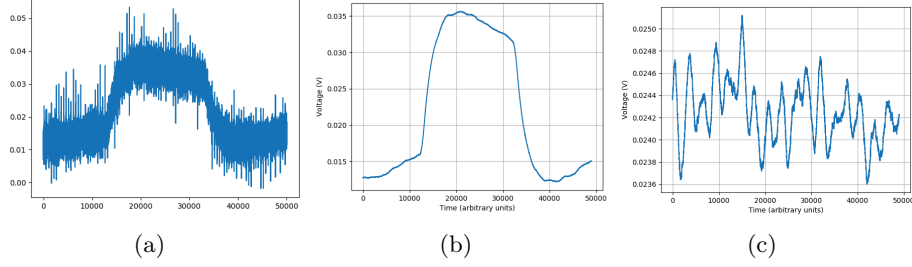


Figure 8.17: Examples of pulses obtained during pulser board tests: noisy pulse coming from the oscilloscope before the averaging procedure (left), clean pulse after averaging procedure (centre), and signal from dead channel with no clear peak (right).

Currently used at the Weizmann Institute of Science, the pulser board testing system allows the construction site to assess the quality of constructed sTGC chambers and the functionality of their readout electronics. If channels on the chamber are flagged as faulty by pulser board tests, bad connections can, e.g., be re-soldered in order to ensure a fully functional chamber upon arrival at CERN and integration into the NSW of the ATLAS detector.

Chapter 9

Search for New High-Mass Phenomena

The search presented in this dissertation selects events based on a number of criteria aiming to detect any potential signal of BSM physics. The dimuon final state is advantageous because of its clean signature and simplicity. In this chapter, Section 9.1 describes the event selection applied to both MC simulation and data. Section 9.2 presents the discriminating variable used for this search. Comparisons between MC simulation and data are presented for various kinematic distributions relating to the dimuon system in Section 9.3.

9.1 Event Selection

In order to select events with two muons in the final state, a series of cuts is applied to each event. A successful event satisfies all of these cuts. Three types of cuts are applied in this analysis: event-level cuts, muon-level cuts, and cuts relating to the dimuon system.

In terms of event-level cuts, each event is required to be in the GRL, as defined in Section 7.3.1. Next, each event passes at least one of the muon triggers `mu26_ivarmedium` (`mu26_imedium` for 2015 data) or `mu50`, which require at least one isolated muon with p_T greater than 26 GeV or one muon with p_T greater than 50 GeV, respectively. Event cleaning is then applied, where events are rejected if they are flagged as incomplete. An incomplete event is one where noise bursts or data corruption occurs in the EM or hadronic calorimeter. At least one reconstructed proton-proton interaction vertex is required for each event, with the primary vertex defined as the one with the highest sum of track p_T squared.

Each event is required to contain at least two muons with p_T greater than 30 GeV satisfying the high- p_T WP described in Section 6.1. Muons are also required to fulfill track-based isolation requirements, in order to reduce the background from light and heavy hadron decays inside jets. In

9.1. Event Selection

particular, the quantity $p_T^{\text{varcone30}}/p_T^\mu$ is used to ensure a 99% efficiency that is constant across η and p_T . The quantity $p_T^{\text{varcone30}}$ is the scalar sum of the transverse momenta of the tracks with $p_T > 1$ GeV within a cone of size $\Delta R = \min(10 \text{ GeV}/p_T^\mu, 0.3)$ around the muon of transverse momentum p_T^μ . The dimuon pair is selected from the muons in the event using the pair having the highest scalar sum p_T . The pair must have opposite sign charge, as well as have an invariant mass larger than 80 GeV, because the search is mainly interested in the higher invariant mass spectrum. Figure 9.1 shows the number of dimuon events passing the event selection, normalized to the integrated luminosity, for each data-taking run during the 2015-2016 data-taking period.

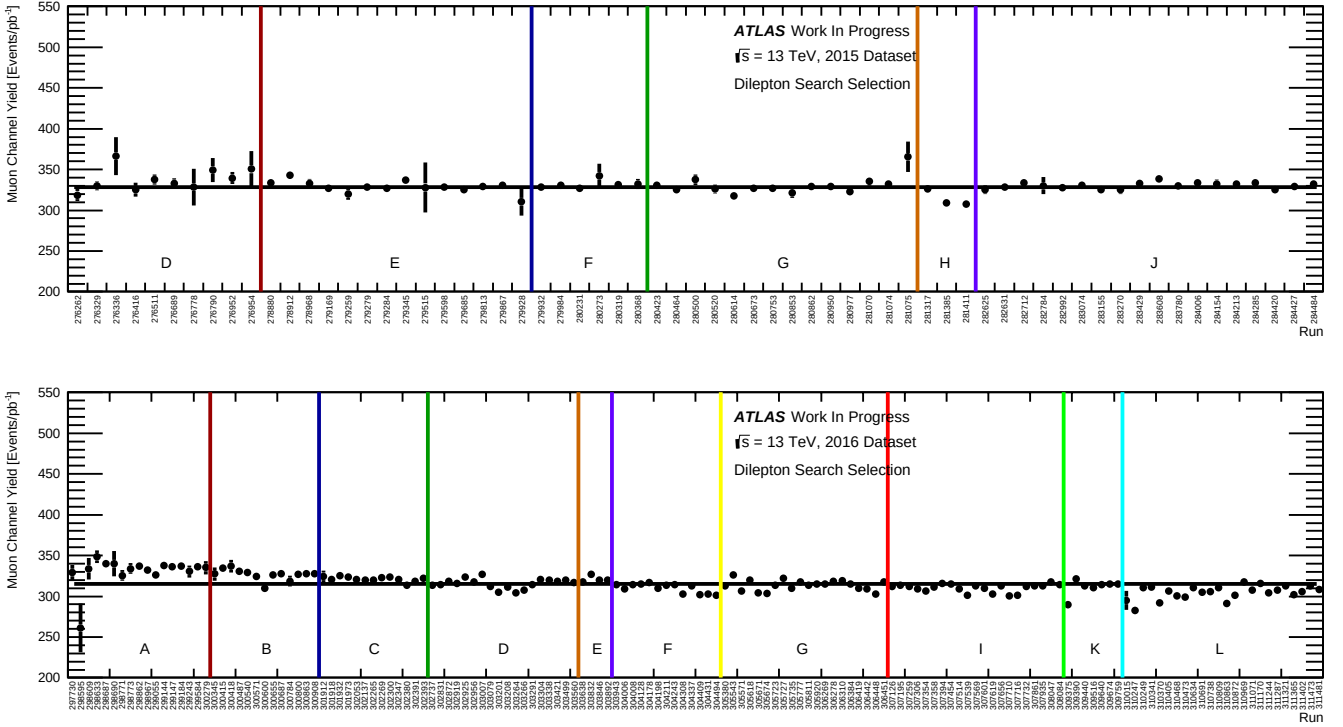


Figure 9.1: Number of events passing the event selection, normalized to the integrated luminosity, as a function of run number during the 2015 (top) and 2016 (bottom) data-taking periods.

9.1.1 Background Processes

The background processes for this analysis are modelled using MC simulations. The three main background contributions to this search are, in order of importance, DY production, top quark production, and diboson (WW , WZ , ZZ) production. The multi-jet and W +jets SM background processes are found to be negligible and are not considered in this search. A summary of the various properties of the background processes considered is presented in Table 9.1. For each background sample, this table presents the MC generator used, the order in QCD at which the sample is generated and corrected, the showering program employed, and the PDFs adopted. In particular, the DY and top quark backgrounds are generated and showered using POWHEG [69] and PYTHIA [70, 71], respectively, while the diboson background is generated and showered using SHERPA [68]. All background processes are generated at NLO with the CT10 [83] PDFs. The DY and top quark backgrounds are corrected to Next-to-Next-to-Leading Order (NNLO) with the CT14NNLO [84] PDFs, while the diboson background is corrected to NLO. In order to gain statistics in regions of phase space of interest, the DY and diboson processes are generated in slices of dimuon invariant mass. Once the background processes are generated, the ATLAS detector response is simulated with GEANT4 [85]. Momentum calibration and resolution smearing are applied to the muon objects in the MC samples, in order to match the performance observed in data [53]. The background simulations then pass through the same event selection discussed in Section 9.1. Finally, data collected by ATLAS is compared with the sum of the simulated backgrounds, in order to look for possible data excesses, which would indicate the presence of new BSM physics.

Process	Generator	Order	Corrected	Parton Shower	PDFs
$q\bar{q} \rightarrow Z/\gamma^* \rightarrow \mu^+\mu^-$	POWHEG v2 [69]	NLO	NNLO	PYTHIA 8.186 [71]	CT10 [83]
$t\bar{t} \rightarrow \mu X, Wt \rightarrow X$	POWHEG v2 [69]	NLO	NNLO	PYTHIA 6.428 [70]	CT10 [83]
$WW, WZ, ZZ \rightarrow \mu X/\mu\nu/\mu\mu$	SHERPA 2.1.1 [68]	NLO	NLO	SHERPA 2.1.1 [68]	CT10 [83]

Table 9.1: Summary of background processes considered in this search. The columns from left to right give the process of interest, generator, matrix element order, parton shower program, and PDFs adopted.

9.1.2 Signal Processes

Various BSM signal processes are generated, in order to compare discriminating variable distributions with collected data, in the hope of finding com-

9.1. Event Selection

patibility between new physics and data. Signal samples for this analysis are obtained by reweighting a LO DY sample. This is validated by comparing the reweighted signal templates to dedicated signal samples generated at specific pole masses. For CI signal samples, the DY background estimate is used below 300 GeV, as a negligible amount of new physics is expected in this range. The simulated signal samples are fed into the same GEANT4 program as the background processes, in order to simulate the ATLAS detector response. A summary of the various properties of the signal processes considered is presented in Table 9.2.

Process	Generator	Order	Corrected	Parton Shower	PDFs
$q\bar{q} \rightarrow Z/\gamma^* \rightarrow \mu^+\mu^-$	PYTHIA 8.186 [71]	LO	NNLO	PYTHIA 8.186 [71]	NNPDF23LO [86]
$q\bar{q} \rightarrow Z' \rightarrow \mu^+\mu^-$	PYTHIA 8.186 [71]	LO	NNLO	PYTHIA 8.186 [71]	NNPDF23LO [86]
CI: $q\bar{q} \rightarrow \mu^+\mu^-$	PYTHIA 8.186 [71]	LO	NNLO	PYTHIA 8.186 [71]	NNPDF23LO [86]

Table 9.2: Summary of signal processes considered in this search. The columns from left to right give the process of interest, generator, matrix element order, parton shower program, and PDFs adopted.

Figure 9.2 shows the dimuon invariant mass distribution for Z'_χ signals of varying pole mass. A notable feature of these distributions is the increasing tails at lower invariant mass for higher Z'_χ pole masses. This is due to the PDFs, hence the moniker parton luminosity tails. Indeed, for particles with a nonzero natural width generated at masses close to the collision energy, the PDFs favor lower-mass collisions. This leads to an increased fraction of off-shell production for signal events as the kinematic production limit is approached. The size of the parton luminosity tails increases with the resonance width. The parton luminosity tails produce a noticeable impact on cross-section limit values presented in Chapter 11 when at least 5 – 10% of events are in the parton luminosity tail.

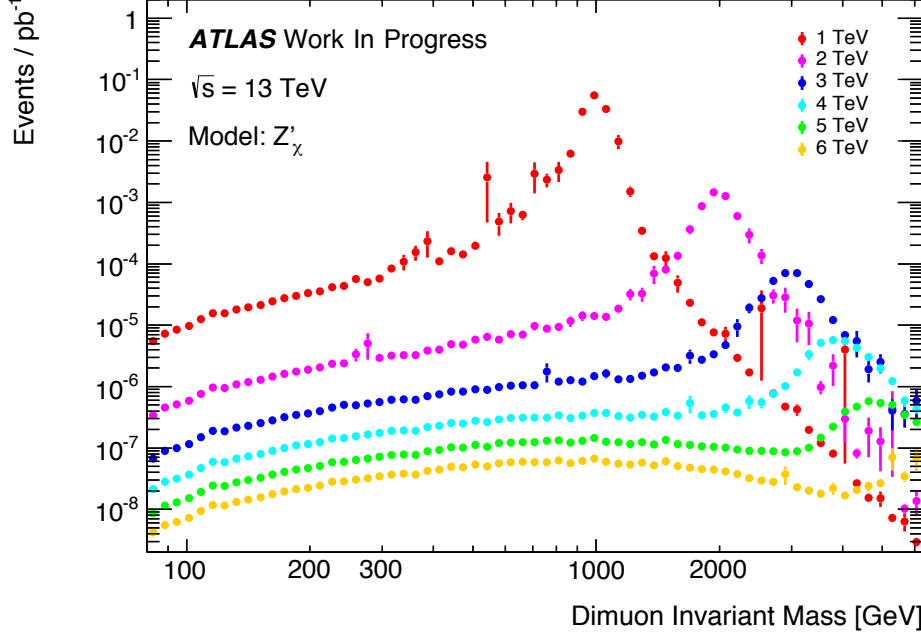


Figure 9.2: Invariant mass distribution of the dimuon system for various Z'_χ pole masses. As the pole masses get larger, the parton luminosity tails at lower invariant mass become more important.

9.1.3 Higher-Order Corrections

In addition to the precise matrix element calculations provided by the MC generators, higher-order corrections to the main DY background process and the various signal processes are desirable, to the extent that they are computationally feasible. In this analysis, mass-dependent K -factors are used in order to obtain a better estimate of the cross-sections of the signal and background processes. In particular, higher-order QCD and EW corrections are applied using K -factors defined as

$$K(m_{\mu\mu}) = \frac{\sigma_{\text{best}}(m_{\mu\mu})}{\sigma_{\text{MC}}(m_{\mu\mu})}, \quad (9.1)$$

where “best” refers to an external NNLO QCD cross-section calculation using the CT14NNLO PDF set, and including missing higher-order EW corrections for the POWHEG samples. The NLO EW corrections are omitted for the PYTHIA samples, as they are not available. The nominal higher-order

9.1. Event Selection

PDFs used to compute the final K -factors are the CT14NNLO PDFs. There are four main sets of K -factors applied to the DY background and various signal processes in this analysis.

First, K -factors are calculated in order to correct the DY background and signal cross-sections from LO to NNLO in QCD with respect to the NNPDF23LO PDFs generated with PYTHIA 8. Second, to correct the DY cross-section from NLO to NNLO in QCD with respect to the CT10 PDFs generated with POWHEG + PYTHIA 8, K -factors are calculated using VRAP0.9 [87]. Third, LO to NLO EW corrections are calculated using MCSANC [88]. Finally, K -factors related to LO Photon-Induced (PI) corrections, computed using MRST2004QED [89], take into account that the generated DY samples do not include these QED effects. Figures 9.3 and 9.4 illustrate the invariant mass dependence of the various higher-order corrections considered in this analysis.

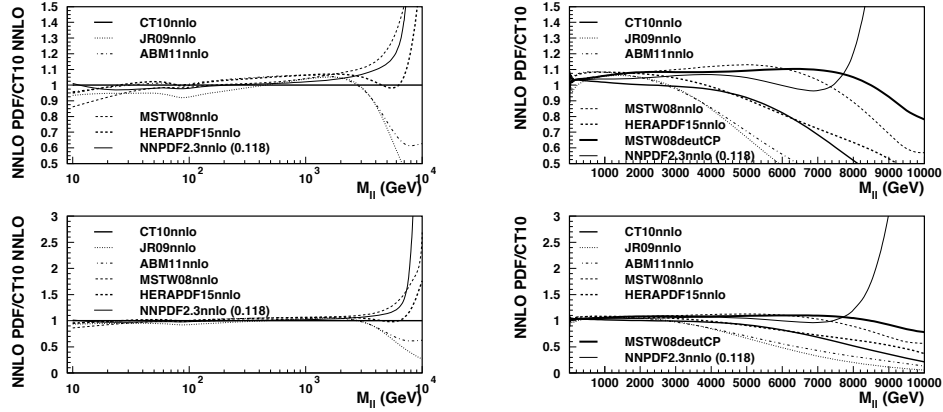


Figure 9.3: Invariant mass dependence of the total DY production cross-section predictions for modern NNLO PDFs with respect to the CT10NNLO PDFs (left) and QCD NNLO-to-NLO K -factors with respect to the CT10NLO PDFs (right). The top and bottom panels show the same quantities, but with different y-axis ranges. The calculations are carried out assuming a centre-of-mass energy of $\sqrt{s} = 14$ TeV and are based on VRAP0.9 and the PDFs as indicated in the legend [90].

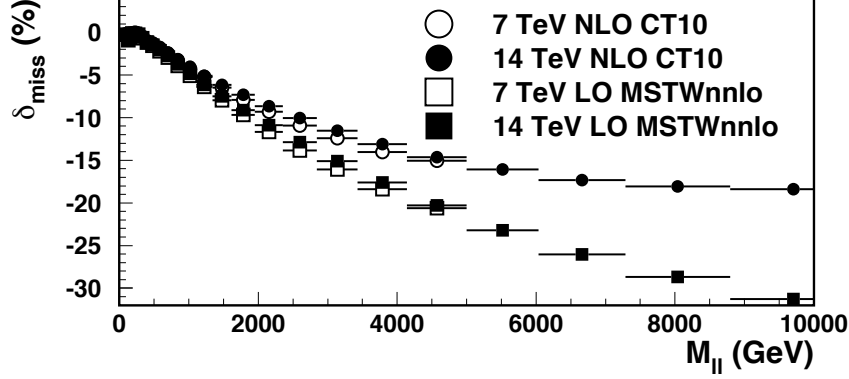


Figure 9.4: Higher-order EW corrections for DY production. The results of the factored (additive) approach based on LO QCD (NLO QCD) for $\sqrt{s} = 7$ TeV are shown with empty squares (empty circles) and for $\sqrt{s} = 14$ TeV with full squares (full circles). All calculations are based on FEWZ 3.1.b2 and the PDFs as indicated in the legend [90].

9.2 Dimuon Invariant Mass

The discriminating variable used in this analysis is the invariant mass of the final state dimuon pair. The invariant mass is calculated by assigning a four-vector to each of the two muons in the final state and calculating

$$m_{\mu\mu}^2 = (p^{\mu_1} + p^{\mu_2})^2, \quad (9.2)$$

where p^{μ_1} and p^{μ_2} are the four-vectors of the final state muons. From Section 5.5.3, the error on $m_{\mu\mu}$ can be approximated as

$$\sigma_{m_{\mu\mu}} = m_{\mu\mu} \sqrt{\left(\frac{\sigma_{p_T^{\mu_1}}}{p_T^{\mu_1}}\right)^2 + \left(\frac{\sigma_{p_T^{\mu_2}}}{p_T^{\mu_2}}\right)^2}. \quad (9.3)$$

Figure 9.5 shows the dimuon invariant mass distribution for the analysis discussed in this dissertation. A summary of the event yields in specific $m_{\mu\mu}$ ranges is given in Table 9.3.

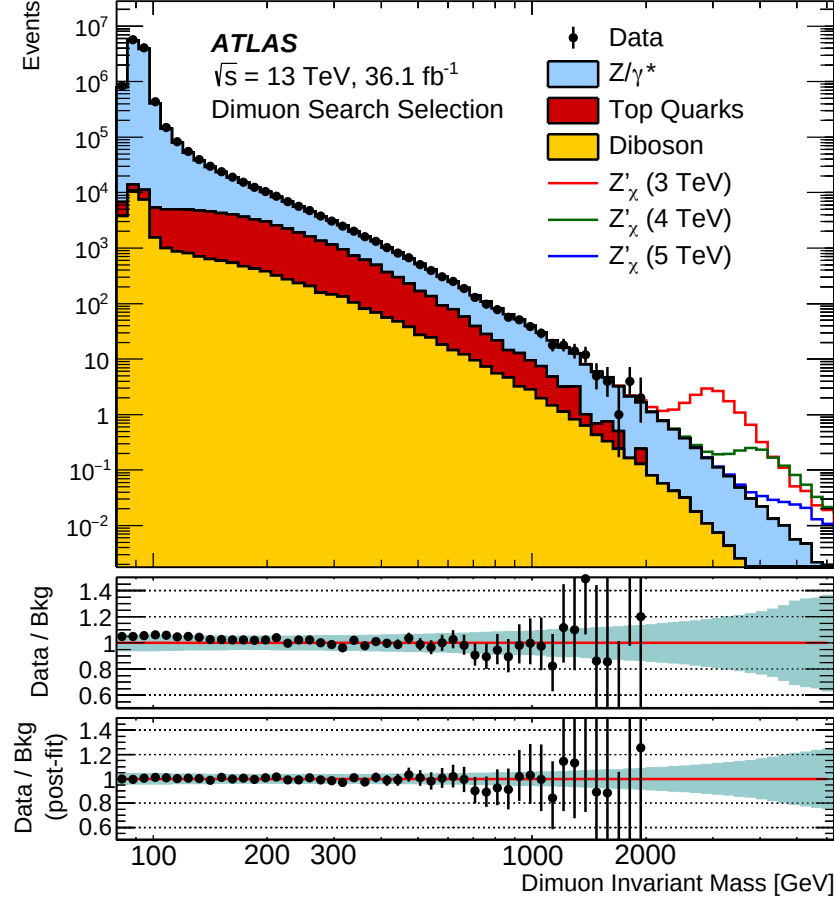


Figure 9.5: Invariant mass distribution of the dimuon system. The data are shown as black markers, while the SM background processes are shown as filled histograms. For illustration purposes, three different Z'_χ signals are added on top of the background distribution. The shaded band in the lower panels illustrates the total systematic uncertainty.

9.3. Kinematic Distributions

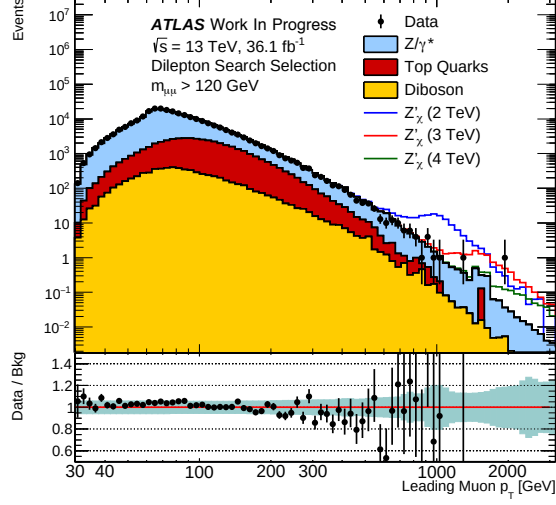
$m_{\mu\mu}$ [GeV]	80–120	120–250	250–400	400–500	500–700
DY	$10\,700\,000 \pm 600\,000$	$177\,900 \pm 10\,000$	$12\,200 \pm 700$	1770 ± 120	1060 ± 80
Top quarks	$24\,700 \pm 1700$	$34\,200 \pm 2400$	6100 ± 500	830 ± 70	401 ± 33
Dibosons	$26\,000 \pm 2800$	5400 ± 600	910 ± 100	155 ± 17	93 ± 11
Total SM	$10\,800\,000 \pm 600\,000$	$218\,000 \pm 10\,000$	$19\,200 \pm 900$	2760 ± 140	1550 ± 90
Data	11 321 561	224 703	19 239	2766	1532
Z'_χ (4 TeV)	0.00873 ± 0.00032	0.0334 ± 0.0015	0.0441 ± 0.0021	0.0246 ± 0.0014	0.052 ± 0.004
Z'_χ (5 TeV)	0.00347 ± 0.00014	0.0137 ± 0.0006	0.0151 ± 0.0007	0.0105 ± 0.0006	0.0176 ± 0.0012
$m_{\mu\mu}$ [GeV]	700–900	900–1200	1200–1800	1800–3000	3000–6000
DY	263 ± 23	110 ± 11	37 ± 4	5.4 ± 0.8	0.30 ± 0.07
Top quarks	68 ± 6	24.5 ± 3.0	5.3 ± 0.9	0.11 ± 0.08	< 0.001
Dibosons	24.3 ± 2.9	9.8 ± 1.2	3.2 ± 0.4	0.45 ± 0.07	0.0184 ± 0.0035
Total SM	355 ± 24	144 ± 11	45 ± 4	6.0 ± 0.8	0.32 ± 0.07
Data	322	141	48	4	0
Z'_χ (4 TeV)	0.0362 ± 0.0026	0.048 ± 0.004	0.067 ± 0.006	0.186 ± 0.022	1.24 ± 0.19
Z'_χ (5 TeV)	0.0153 ± 0.0011	0.0185 ± 0.0015	0.0233 ± 0.0021	0.0258 ± 0.0029	0.118 ± 0.020

Table 9.3: Expected and observed event yields in different $m_{\mu\mu}$ intervals. The quoted errors correspond to the combined statistical, theoretical, and experimental systematic uncertainties. Expected event yields are reported for the Z'_χ model, for two values of the pole mass.

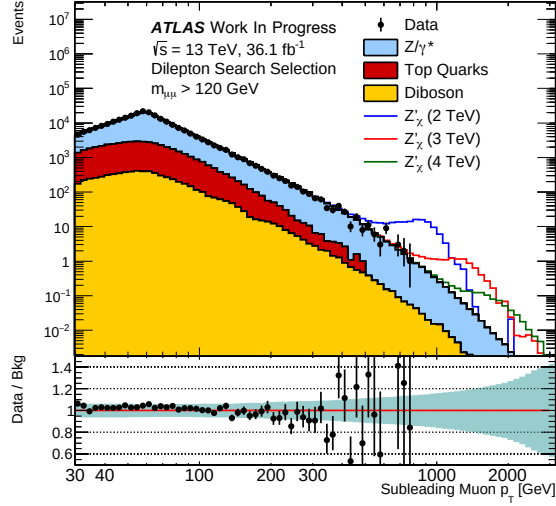
9.3 Kinematic Distributions

Once the event selection is defined, kinematic variables can be examined, in order to compare MC simulations of expected background processes with collected data. Below are the main kinematic distributions used in the analysis. In particular, Figures 9.6 to 9.11 show the kinematic distributions used to validate that MC simulations are in agreement with data.

9.3. Kinematic Distributions



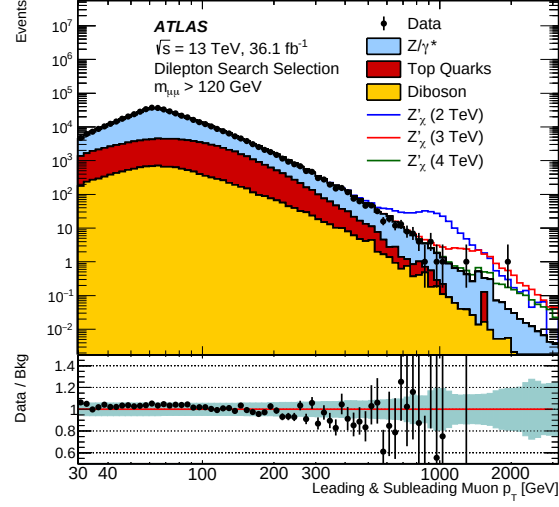
(a)



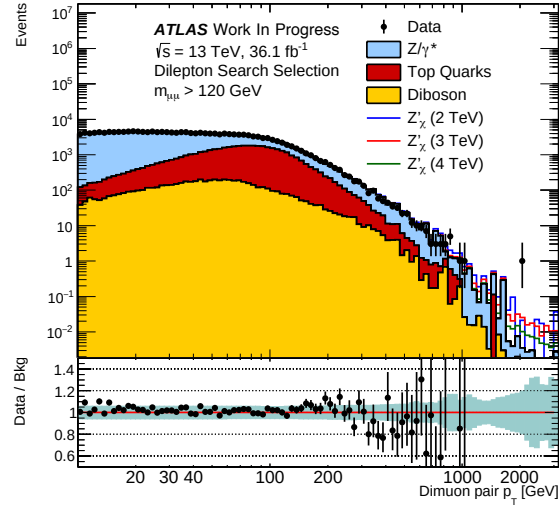
(b)

Figure 9.6: Reconstructed muon p_T distribution for the leading (top) and subleading (bottom) muon after event selection. The distributions are shown for dimuon invariant masses greater than 120 GeV. The data are shown as black markers, while the SM background processes are shown as filled histograms. For illustration purposes, three different Z'_χ signals are added on top of the background distribution. The shaded band in the lower panels illustrates the total systematic uncertainty.

9.3. Kinematic Distributions



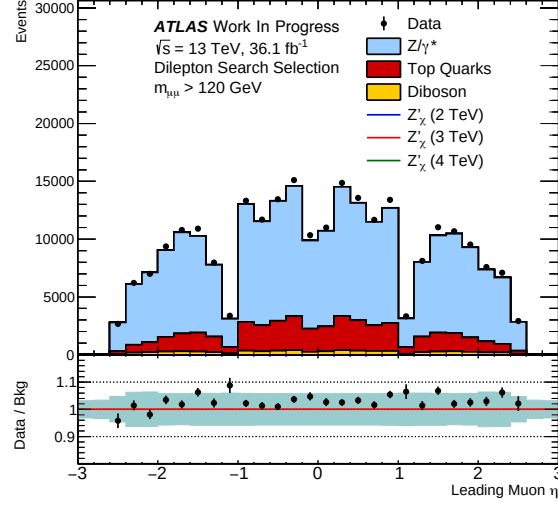
(a)



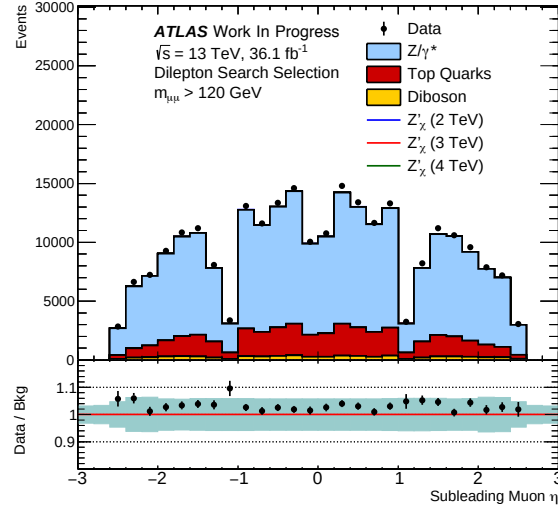
(b)

Figure 9.7: Reconstructed muon p_T distribution for the sum of the leading and subleading muons (top) and dimuon system (bottom) after event selection. The distributions are shown for dimuon invariant masses greater than 120 GeV. The data are shown as black markers, while the SM background processes are shown as filled histograms. For illustration purposes, three different Z'_χ signals are added on top of the background distribution. The shaded band in the lower panels illustrates the total systematic uncertainty.

9.3. Kinematic Distributions



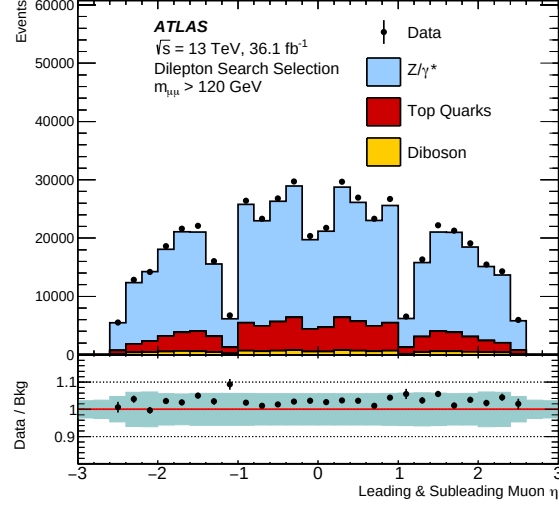
(a)



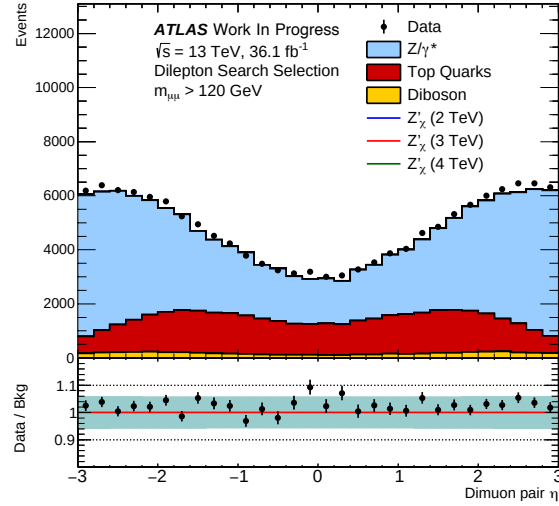
(b)

Figure 9.8: Reconstructed muon η distribution for the leading (top) and sub-leading (bottom) muon after event selection. The distributions are shown for dimuon invariant masses greater than 120 GeV. The data are shown as black markers, while the SM background processes are shown as filled histograms. For illustration purposes, three different Z'_χ signals are added on top of the background distribution. The shaded band in the lower panels illustrates the total systematic uncertainty.

9.3. Kinematic Distributions



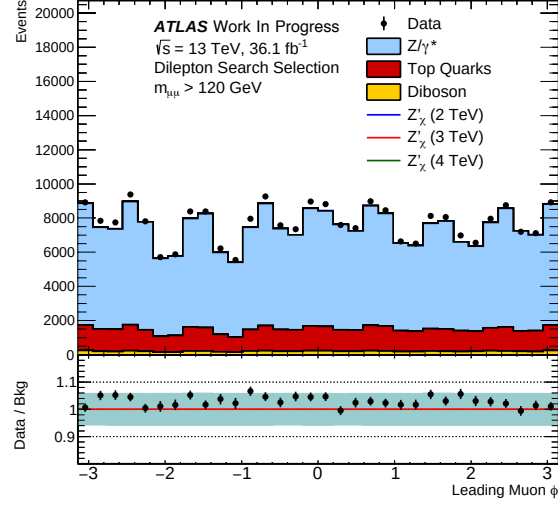
(a)



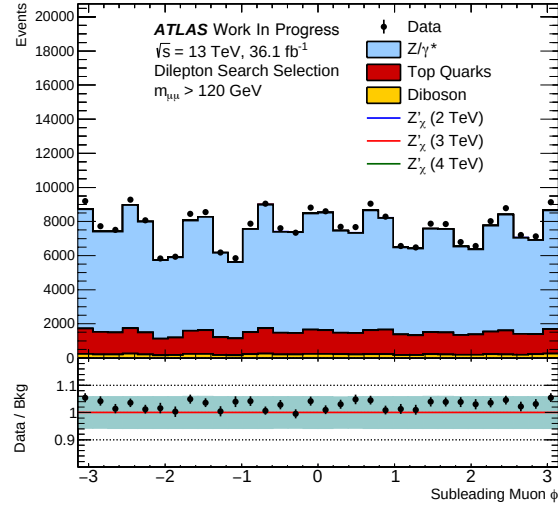
(b)

Figure 9.9: Reconstructed muon η distribution for the sum of the leading and subleading muons (top) and dimuon system (bottom) after event selection. The distributions are shown for dimuon invariant masses greater than 120 GeV. The data are shown as black markers, while the SM background processes are shown as filled histograms. For illustration purposes, three different Z'_χ signals are added on top of the background distribution. The shaded band in the lower panels illustrates the total systematic uncertainty.

9.3. Kinematic Distributions



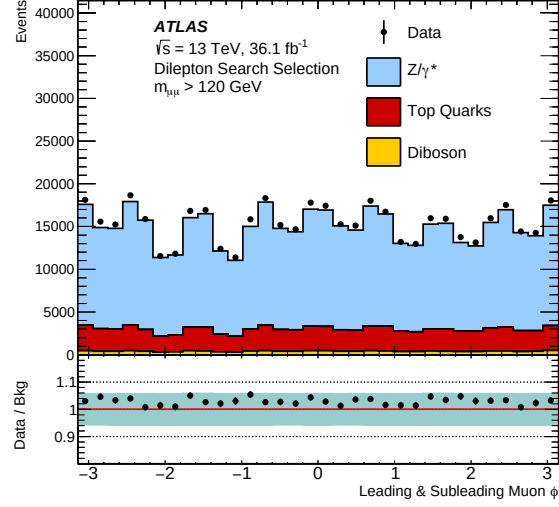
(a)



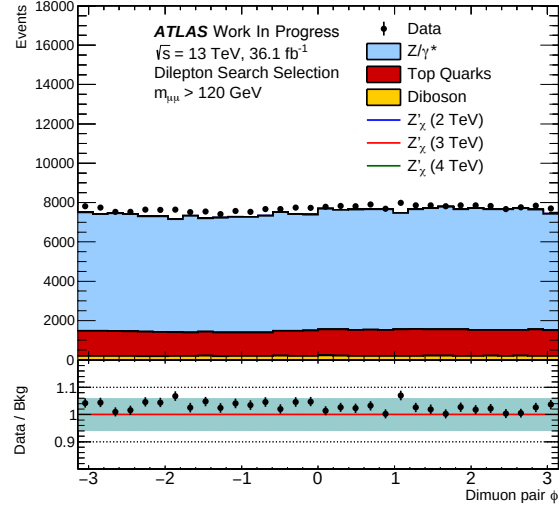
(b)

Figure 9.10: Reconstructed muon ϕ distribution for the leading (top) and subleading (bottom) muon after event selection. The distributions are shown for dimuon invariant masses greater than 120 GeV. The data are shown as black markers, while the SM background processes are shown as filled histograms. For illustration purposes, three different Z'_χ signals are added on top of the background distribution. The shaded band in the lower panels illustrates the total systematic uncertainty.

9.3. Kinematic Distributions



(a)



(b)

Figure 9.11: Reconstructed muon ϕ distribution for the sum of the leading and subleading muons (top) and dimuon system (bottom) after event selection. The distributions are shown for dimuon invariant masses greater than 120 GeV. The data are shown as black markers, while the SM background processes are shown as filled histograms. For illustration purposes, three different Z'_χ signals are added on top of the background distribution. The shaded band in the lower panels illustrates the total systematic uncertainty.

Chapter 10

Systematic Uncertainties

There are three distinct types of systematic uncertainties considered in this analysis: those encapsulating experimental effects, those due to theoretical effects, and those pertaining to the PDFs of the incoming partons. The experimental systematic uncertainties are applied to both signal and background processes, whereas the theoretical and PDF-related systematic uncertainties are applied to the background processes only. The systematic variations are estimated as a function of $m_{\mu\mu}$, and for each $m_{\mu\mu}$ bin the uncertainty applied is the larger between the up and down variations. The overall uncertainty related to each effect is the symmetric up/down variation envelope, expressed as a function of dimuon invariant mass. In all cases, the non-negligible systematic uncertainties are treated as nuisance parameters in the statistical analysis described in Chapter 11. This chapter describes all of the systematic uncertainties considered in the search for new BSM physics presented in this dissertation. Section 10.1 outlines the experimental systematic uncertainties, Section 10.2 presents the theoretical systematic uncertainties, and Section 10.3 discusses the systematic uncertainties related to the PDFs of the incoming partons. Section 10.4 gives a quantitative summary of the systematic uncertainties considered in the analysis presented in this dissertation.

10.1 Experimental Uncertainties

The experimental systematic uncertainties considered in this analysis are applied to both signal and background simulation samples. They include the beam energy uncertainty, the luminosity measurement uncertainty, the pile-up reweighting uncertainty, and the trigger efficiency uncertainty. Furthermore, the uncertainties on the muon momentum scale, the muon momentum resolution (both in the ID and the MS), and the muon isolation efficiency are also evaluated. The uncertainty on the efficiency to reconstruct and select muons is also taken into account.

The dominant experimental systematic uncertainty in this search is due to the reconstruction and selection efficiency of muons. In particular, the

uncertainty related to the high- p_T WP selection is larger with increasing p_T due to the cut on the relative q/p error. The uncertainty related to the muon isolation and trigger is estimated by varying the data-to-simulation SFs applied to simulation by their respective uncertainties. The trigger SFs and their uncertainties are presented in Chapter 7. The systematic uncertainties related to the scale and resolution of the ID and MS p_T measurements are obtained by varying the parameters of the p_T smearing described in Section 5.5.2 by their appropriate uncertainties, and taking the relative difference between the two estimates as the uncertainty. An uncertainty on the beam energy of 0.65% is estimated and included. A flat systematic uncertainty of 3.2% is assessed due to the uncertainty on the luminosity measurement. The uncertainty related to pile-up is estimated by varying the mean number of interactions per bunch crossing in simulation and taking the relative difference between the nominal and varied simulation as the uncertainty.

Figures 10.1 and 10.2 depict the magnitude of each experimental systematic uncertainty source, expressed as a fraction of the signal or background considered. The systematic uncertainty on the trigger efficiency, the muon momentum scale, and the pile-up reweighting are neglected in the statistical analysis, because of their negligible contribution. All other sources of experimental systematic uncertainty are used as nuisance parameters in the statistical analysis described in Chapter 11.

10.1. Experimental Uncertainties

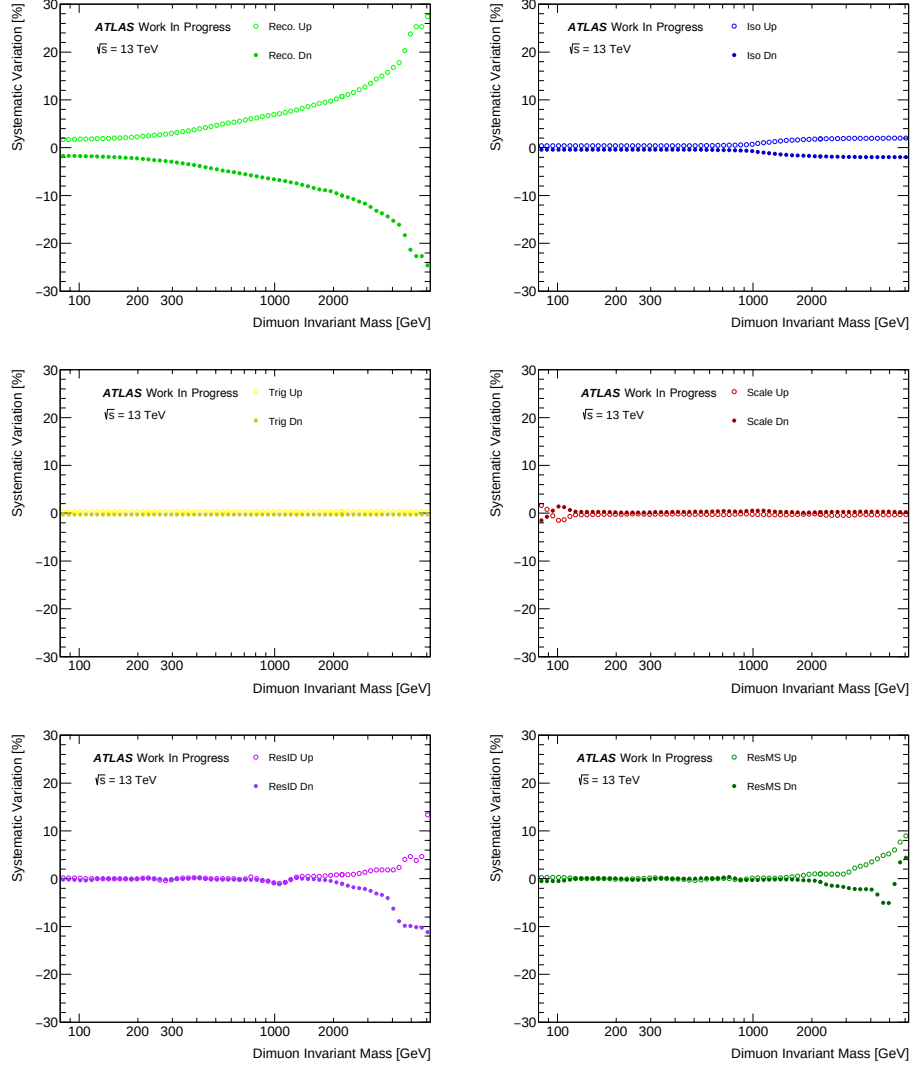


Figure 10.1: Systematic uncertainties due to experimental sources related to muon performance. Shown are, from left to right and top to bottom, the uncertainty related to muon reconstruction and selection, muon isolation, muon trigger, muon momentum scale, muon momentum resolution in the ID, and muon momentum resolution in the MS.

10.2 Theoretical Uncertainties

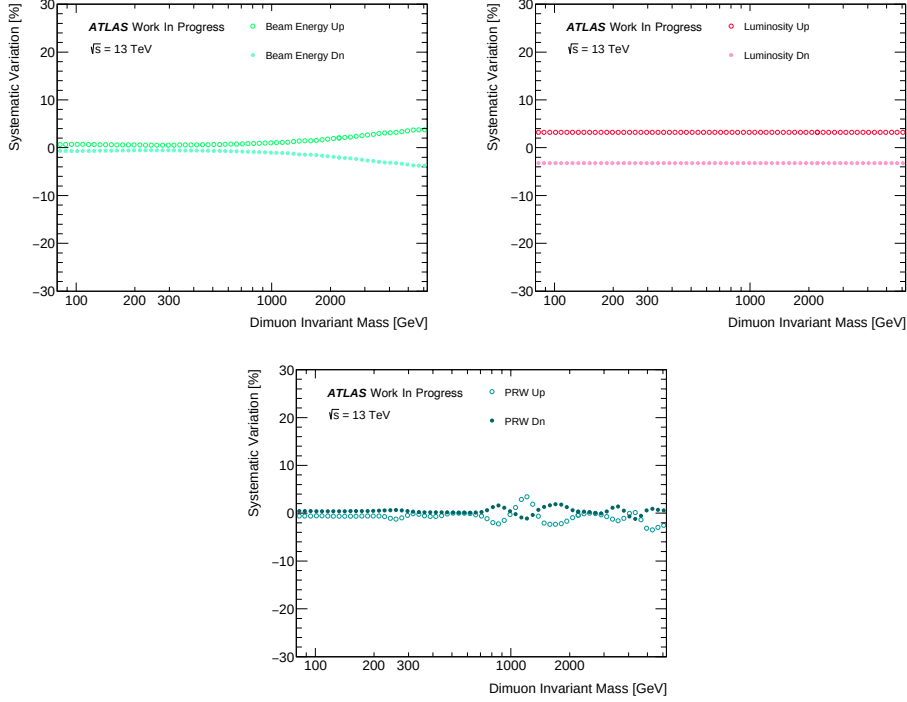


Figure 10.2: Systematic uncertainties due to experimental sources related to accelerator performance. Shown are the uncertainty related to beam energy (top left), luminosity (top right), and pile-up (bottom).

10.2 Theoretical Uncertainties

The theoretical systematic uncertainties considered in this analysis are applied to background simulation samples. They include the variations in the α_s strong coupling constant, the differences between the additive and factored treatment of the electroweak corrections (the additive treatment is used as the nominal treatment in this search), and PI corrections. In addition, the theoretical systematic uncertainties on the normalization of the $t\bar{t}$ and diboson background contributions are evaluated.

The systematic uncertainty related to the strong coupling constant α_s is evaluated using the VRAP program, by varying the nominal input value of $\alpha_s = 0.118$ by ± 0.003 . The uncertainty related to PI corrections is estimated by taking into account the uncertainties on the quark masses and on the photon PDF. This search nominally uses the additive treat-

ment when combining higher-order electroweak and QCD corrections, i.e., $(1 + \delta_{\text{EW}} + \delta_{\text{QCD}})$. The systematic uncertainty resulting from electroweak corrections is estimated by comparing the additive treatment to the multiplicative treatment of this combination, i.e., $((1 + \delta_{\text{EW}})(1 + \delta_{\text{QCD}}))$, and taking the difference between the two treatments as the uncertainty. Additional theoretical systematic uncertainties are assigned to the estimation of the $t\bar{t}$ and diboson backgrounds. They are derived following the PDF4LHC prescription [91]. In particular, the variation of the factorization and renormalization scales are taken into account.

Figure 10.3 depicts the contribution of each theoretical systematic uncertainty to the total systematic uncertainty used in the statistical analysis described in Chapter 11. All theoretical systematic uncertainties are used as nuisance parameters in the statistical analysis, except the diboson normalization systematic uncertainty, which is neglected.

10.2. Theoretical Uncertainties

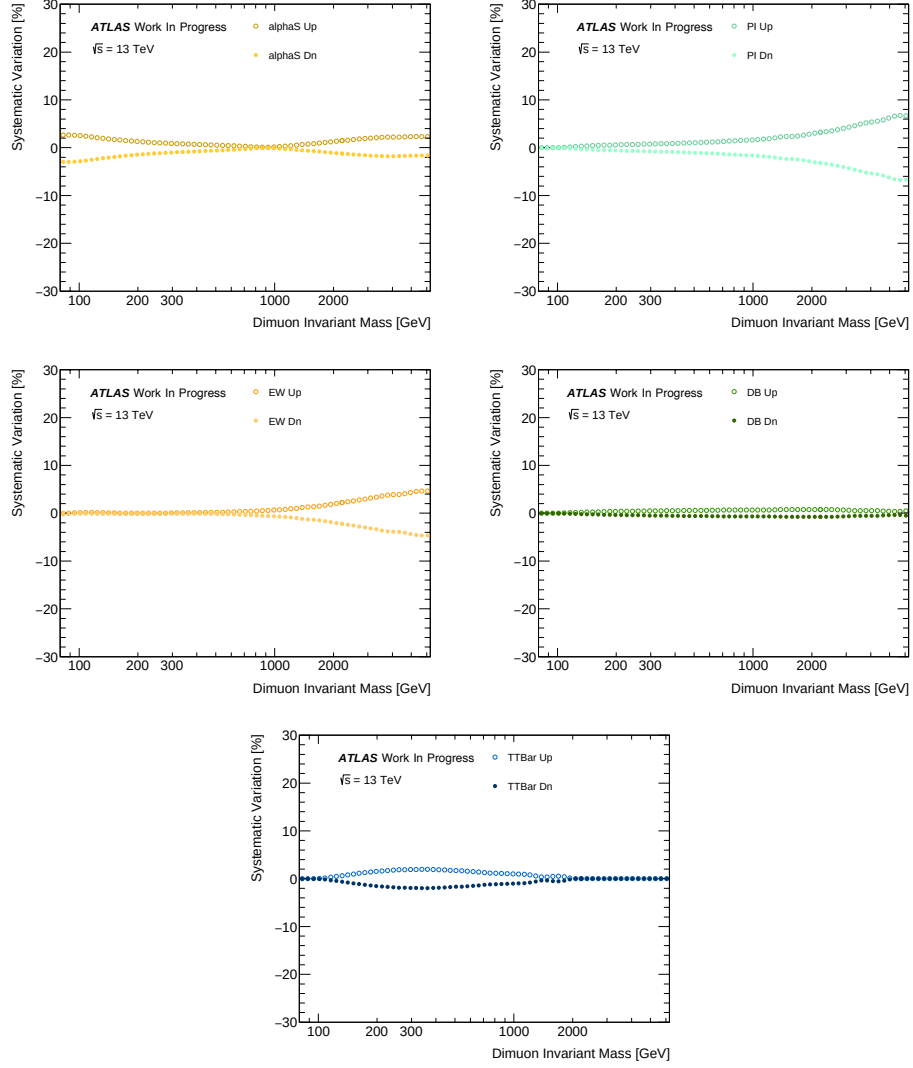


Figure 10.3: Systematic uncertainties due to theoretical sources. Shown are, from left to right and top to bottom, the uncertainty related to the strong coupling constant α_s , PI corrections, EW corrections, diboson background estimation, and $t\bar{t}$ background estimation.

10.3 Parton Distribution Function Uncertainties

Because varying the PDFs of the incoming partons changes the DY cross-section as a function of $m_{\mu\mu}$, theoretical systematic uncertainties pertaining to the PDFs of the incoming partons are considered. In particular, the choice of the PDFs used, the variations of the PDF eigenvector sets, and the variations in the PDF scales are evaluated in this analysis.

The uncertainty on the PDF scales is evaluated by using the VRAP program and varying the nominal CT14NNLO PDF renormalization and factorization scales up or down simultaneously by a factor of 2. The resulting maximum variations are taken as the PDF scale uncertainties. Each PDF has a set of independent parameters associated with it, the so-called eigenvectors of the PDF, which can be varied in orthogonal directions to quantify the systematic uncertainties associated with the PDF variations. For each $m_{\mu\mu}$ histogram bin, the asymmetric uncertainty on the cross-section is calculated using

$$\begin{aligned}\Delta\sigma^+ &= \sqrt{\sum_{i=1}^n (\max(\sigma_i^+ - \sigma_0, \sigma_i^- - \sigma_0, 0))^2}, \\ \Delta\sigma^- &= \sqrt{\sum_{i=1}^n (\max(\sigma_0 - \sigma_i^+, \sigma_0 - \sigma_i^-, 0))^2},\end{aligned}\tag{10.1}$$

where n is the number of PDF eigenvectors, σ_i^+ (σ_i^-) is the cross-section for the higher (lower) value of the i^{th} PDF eigenvector, and σ_0 is the cross-section for the central value PDF. A total of seven eigenvector variations are considered [84]. These seven variations result from the bundling of the 28 eigenvectors of the CT14 PDFs into seven bundles with similar mass dependency. The eigenvector bundles are provided by the CT14 authors, and are particularly useful when considering any future combination effort with other searches due to the orthogonality of the eigenvector bundles between the W and Z bosons. The systematic uncertainty due to the variation of the PDF eigenvectors is then taken as the larger of the positive and negative variation on the cross-section calculated in Equation (10.1). The nominal PDF choice used for the incoming partons in this analysis is the CT14NNLO PDF set [84]. The uncertainty related to this choice is evaluated by comparing the central values of the CT14NNLO PDFs to similar PDFs. In particular, two alternatives for the nominal PDF choice are investigated: MMHT14 [92] and NNPDF3.0 [93]. When the difference between the nominal PDF set pre-

10.3. Parton Distribution Function Uncertainties

dictions and an alternative choice is larger than the uncertainty related to the PDF eigenvector variation, the maximum absolute deviation from the envelope of these comparisons is used as the PDF choice uncertainty.

Figures 10.4 and 10.5 depict the contribution of each systematic uncertainty pertaining to the PDFs of incoming partons to the total systematic uncertainty used in the statistical analysis described in Chapter 11. All systematic uncertainties pertaining to the PDFs of the incoming partons are used as nuisance parameters in the statistical analysis.

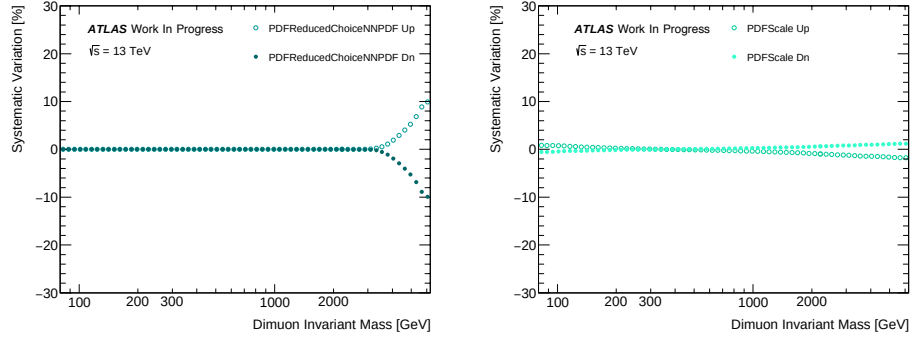


Figure 10.4: Systematic uncertainties pertaining to the choice (left) and scale (right) of the PDFs of the incoming partons.

10.3. Parton Distribution Function Uncertainties

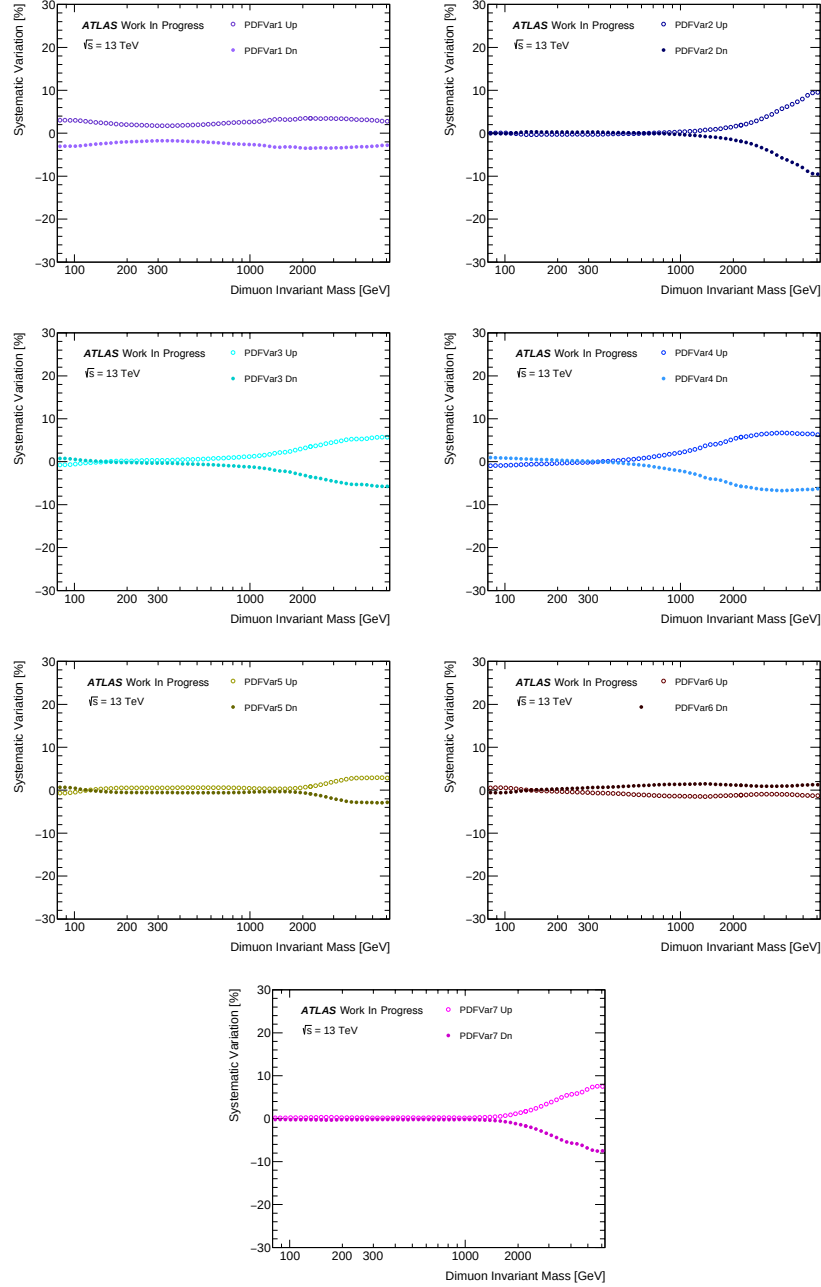


Figure 10.5: Systematic uncertainties pertaining to the eigenvector variations of the PDFs of the incoming partons. Shown are, from left to right and top to bottom, the uncertainty related to the seven eigenvector bundles provided by the CT14 authors.

10.4 Summary

Table 10.1 summarizes the various systematic uncertainties considered and their various contributions. The experimental systematic uncertainties are dominated by the muon reconstruction efficiency, which results in a relative systematic uncertainty of $\sim 10\%$ ($\sim 17\%$) at $m_{\mu\mu} = 2$ TeV ($m_{\mu\mu} = 4$ TeV). The dominant theoretical systematic uncertainty is the DY background PDF eigenvector set variation in this analysis, resulting in a relative systematic uncertainty of $\sim 8\%$ ($\sim 13\%$) at $m_{\mu\mu} = 2$ TeV ($m_{\mu\mu} = 4$ TeV).

Source	Signal [%]	Background [%]
Luminosity	3.2 (3.2)	3.2 (3.2)
MC statistical	<1.0 (<1.0)	<1.0 (<1.0)
Beam energy	1.9 (3.1)	1.9 (3.1)
Pile-up effects	<1.0 (<1.0)	<1.0 (<1.0)
DY PDF choice	N/A	<1.0 (1.9)
DY PDF variation	N/A	7.7 (13)
DY PDF scales	N/A	<1.0 (1.5)
DY α_S	N/A	1.4 (2.2)
DY EW corrections	N/A	2.1 (3.9)
DY PI corrections	N/A	3.0 (5.4)
Top quarks theoretical	N/A	<1.0 (<1.0)
Dibosons theoretical	N/A	<1.0 (<1.0)
Reconstruction efficiency	10 (17)	10 (17)
Isolation efficiency	1.8 (2.0)	1.8 (2.0)
Trigger efficiency	<1.0 (<1.0)	<1.0 (<1.0)
Muon momentum scale	<1.0 (<1.0)	<1.0 (<1.0)
Muon momentum resolution	2.7 (2.7)	<1.0 (6.7)
Total	11 (18)	14 (24)

Table 10.1: Summary of the pre-marginalized relative systematic uncertainties in the expected number of events at dimuon masses of 2 TeV (4 TeV). The values quoted for the background represent the relative change in the total expected number of events in the corresponding $m_{\mu\mu}$ histogram bin containing the reconstructed $m_{\mu\mu}$ mass of 2 TeV (4 TeV). For the signal uncertainties, the values were computed using a Z'_χ signal model with a pole mass of 2 TeV (4 TeV), by comparing yields in the core of the mass peak, within the full width at half maximum, between the distribution varied due to a given uncertainty and the nominal distribution.

Chapter 11

Statistical Analysis

The comparisons between data and MC simulations presented in Chapter 9 allow for the statistical analysis of the available dataset. In particular, the compatibility between the data and the SM expectation can be quantified. Any excess of data above the SM expectation can be interpreted as a manifestation of various BSM theories. In the case of an excess, the LHC has the potential to measure the properties of the new phenomena. On the other hand, if no significant excess is found, constraints on theoretical parameters of BSM theories are placed. In this chapter, Section 11.1 describes the statistical methods used to quantify the agreement between the data and the MC simulations, and Section 11.2 presents the various limits set on the BSM models presented in Chapter 3.

11.1 Search Results

In order to quantify any potential excess of data compared with MC simulations, a search for a resonant signal is performed using the dimuon invariant mass distribution. Starting from Poisson statistics, the likelihood of observing n events given an expected number of events λ is given by

$$\mathcal{L}(n|\lambda) = \frac{\lambda^n e^{-\lambda}}{n!}. \quad (11.1)$$

As this search uses the $m_{\mu\mu}$ distribution presented in Section 9.2, the generalization of this likelihood becomes the product of the Poisson likelihoods across the bins of the $m_{\mu\mu}$ histogram:

$$\mathcal{L}(\mathbf{n}|\boldsymbol{\lambda}) = \prod_{j=1}^{N_{\text{bin}}} \frac{\lambda_j^{n_j} e^{-\lambda_j}}{n_j!}, \quad (11.2)$$

where λ_j represents the expected number of events in bin j and n_j corresponds to the observed number of events in bin j . The expected number of events for this search depends on the number of background events from

11.1. Search Results

SM processes and the number of signal events from the given BSM model under consideration:

$$\lambda_j = \mu s_j + b_j, \quad (11.3)$$

where s_j and b_j are the expected number of signal and background events in bin j , respectively. The parameter μ determines the strength of the signal process, with $\mu = 0$ corresponding to the background-only hypothesis and $\mu = 1$ corresponding to the nominal signal hypothesis.

The systematic uncertainties described in Chapter 10 are accounted for in the search by introducing nuisance parameters θ_s , which parameterise all sources of systematic uncertainty, into the likelihood and by modifying the expected number of events to allow for systematic error fluctuations:

$$\lambda_j \rightarrow \tilde{\lambda}_j = \lambda_j \left(1 + \sum_{s=1}^{N_{\text{sys}}} \theta_s \varepsilon_{sj} \right), \text{ where } \varepsilon_{sj} = \left(\frac{\lambda_{sj}^{\text{var}} - \lambda_j}{\lambda_j} \right). \quad (11.4)$$

In the above equation, $\lambda_{sj}^{\text{var}}$ represents the number of background events in bin j for the systematically varied $m_{\mu\mu}$ distribution corresponding to the systematic uncertainty s . Hence, ε_{sj} quantifies the relative shift in the expected number of background events in bin j from the nominal number of expected background events for the systematic uncertainty s . Adding a Gaussian prior probability density function constraint to the nuisance parameters θ_i , the likelihood of observing the $m_{\mu\mu}$ distribution obtained in data given a signal strength μ and nuisance parameters $\boldsymbol{\theta} = (\theta_1, \theta_2, \dots, \theta_{N_{\text{sys}}})$ is

$$\mathcal{L}(\mathbf{n}|\mu, \boldsymbol{\theta}) = \prod_{j=1}^{N_{\text{bin}}} \frac{\tilde{\lambda}_j^{n_j} e^{-\tilde{\lambda}_j}}{n_j!} \prod_{s=1}^{N_{\text{sys}}} \frac{e^{-\theta_s^2/2}}{\sqrt{2\pi}}. \quad (11.5)$$

From this expression, the maximum-likelihood estimators $\hat{\mu}$ and $\hat{\boldsymbol{\theta}}$ for a given observed $m_{\mu\mu}$ distribution and signal hypothesis can be computed by maximizing the likelihood function. Likewise, the nuisance parameter values that maximize the likelihood function under the background-only hypothesis, $\hat{\boldsymbol{\theta}}_0$, are computed for the observed data.

From these quantities, the test statistic q_0 can be constructed:

$$q_0 = \begin{cases} +2 \ln \frac{\mathcal{L}(\text{data}|0, \hat{\boldsymbol{\theta}}_0)}{\mathcal{L}(\text{data}|\hat{\mu}, \hat{\boldsymbol{\theta}})} & \text{for } \hat{\mu} < 0, \\ -2 \ln \frac{\mathcal{L}(\text{data}|0, \hat{\boldsymbol{\theta}}_0)}{\mathcal{L}(\text{data}|\hat{\mu}, \hat{\boldsymbol{\theta}})} & \text{for } \hat{\mu} \geq 0, \end{cases} \quad (11.6)$$

11.1. Search Results

where a larger value of q_0 corresponds to a larger incompatibility between the data and the background-only hypothesis. Using the log-likelihood ratio test described in [94], the probability of the background fluctuating in such a way as to create a signal-like excess equal to or larger than what is observed in data, the so-called p -value p_0 , can be obtained:

$$p_0 = \int_{q_0^{\text{obs}}}^{\infty} f(q_0|0, \hat{\theta}_0) dq_0, \quad (11.7)$$

where $f(q_0|0, \hat{\theta}_0)$ denotes the probability density function of q_0 under the assumption of the background-only hypothesis. The value of p_0 for $\hat{\mu} < 0$ is capped at 0.5 by setting $q_0 = 0$ for Z' pole masses above 1.5 TeV, in order to avoid negative probabilities arising from the lack of statistics in this region.

The significance, Z , defined such that a Gaussian distributed variable found Z standard deviations (σ) above its mean has an upper-tail probability equal to p_0 , is used to quantify any potential excess in data:

$$Z = \Phi^{-1}(1 - p_0), \quad (11.8)$$

where Φ^{-1} is the inverse of the cumulative distribution function of the standard Gaussian distribution. Figure 11.1 depicts the cumulative distribution function of the standard Gaussian distribution and its inverse.

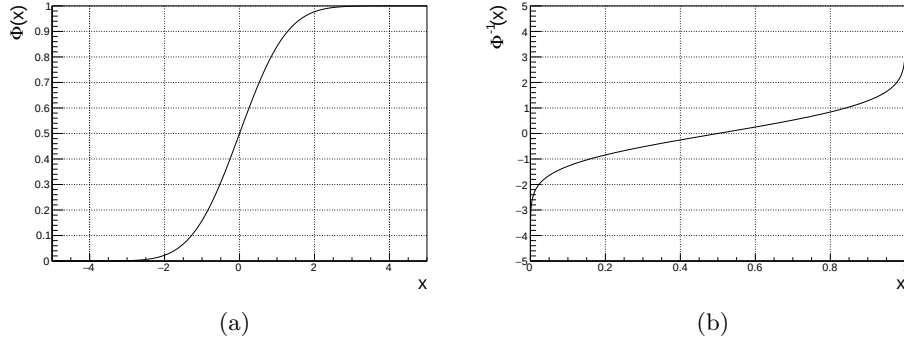


Figure 11.1: Cumulative distribution function (left) and inverse of the cumulative distribution function (right) of the standard Gaussian distribution.

The generally accepted significance value required in order to claim evidence (discovery) of new physics is 3σ (5σ).

11.1.1 Asymptotic Approximation

Figure 11.2 shows the q_0 distribution for background-only pseudo-experiments, or toys, generated using the nuisance parameter values $\hat{\theta}_0$ and event counts following Poisson probabilities. In the limit where the number of observed events in data is large, the distribution for q_0 for background-only datasets follows a χ^2 distribution with one degree of freedom [95].

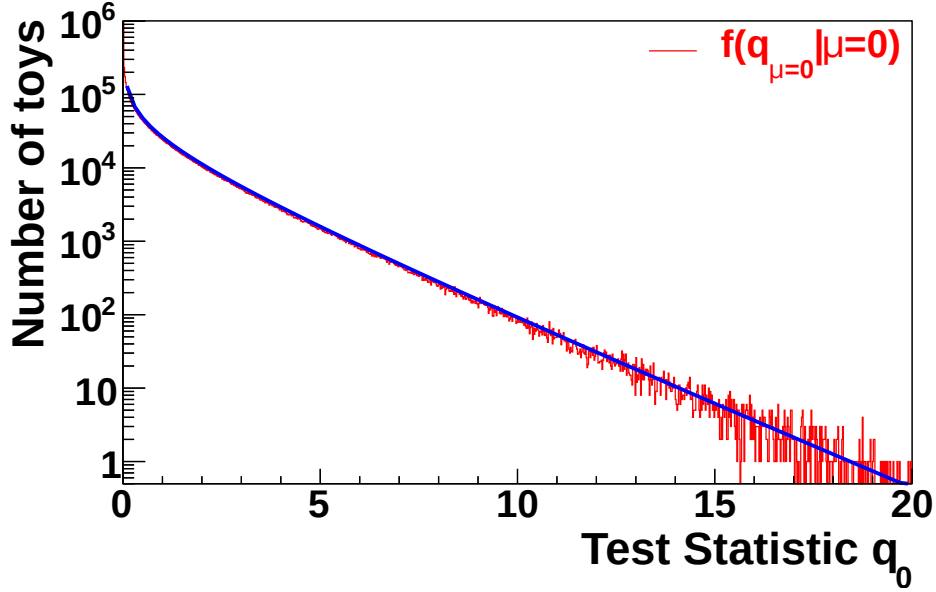


Figure 11.2: Distribution of the q_0 test statistic for background-only toys. The solid blue line shows the χ^2 distribution with one degree of freedom scaled to the integral of the q_0 distribution.

Rather than generating the $f(q_0|0, \hat{\theta}_0)$ distribution for each q_0 calculated, the asymptotic approximation described in [94] is utilized in order to save computational resources. In particular, it has been shown that

$$p_0 = 1 - \Phi(\sqrt{q_0}) \text{ and } Z = \sqrt{q_0}, \quad (11.9)$$

where Φ is the cumulative distribution function of the standard Gaussian distribution.

11.1.2 Look-Elsewhere Effect

When considering p -values, the so-called “look-elsewhere effect” [96] must be taken into account. This effect states that an apparently statistically significant observation may have actually arisen because of the size of the parameter space being searched. Because we are calculating the local p -value over a large dimuon invariant mass range, the look-elsewhere effect cannot be neglected. To account for this, a global p -value, i.e., the probability of measuring a particular local p -value somewhere in the background invariant mass spectrum that is at least as significant as the one observed in data, is calculated. The fluctuations in the background invariant mass distribution are modeled by running an ensemble of pseudo-experiments generated under the background-only hypothesis. The largest local significance for each toy is then used as the means by which to measure the global p -value:

$$p_{global} = \frac{1}{N_{toys}} \int_{\hat{z}_0}^{\infty} n(z_0) dz_0, \quad (11.10)$$

where $n(z_0)$ is the distribution of highest local significances observed in the pseudo-experiments, and $\hat{z}_0$ is the local significance observed in data. The global p -value represents the probability of obtaining a local significance larger than or equal to the local significance observed in data, under the background-only hypothesis. In other words, the global p -value can be interpreted as a p -value of local p -values.

Figure 11.3 shows the p -values obtained from the collected dataset using the methodology described above. The binning used for this p -value scan is an optimized binning corresponding to the dimuon mass resolution, which varies from 60 GeV at $m_{\mu\mu} = 1$ TeV to 200 GeV at $m_{\mu\mu} = 2$ TeV and 420 GeV at $m_{\mu\mu} = 3$ TeV.

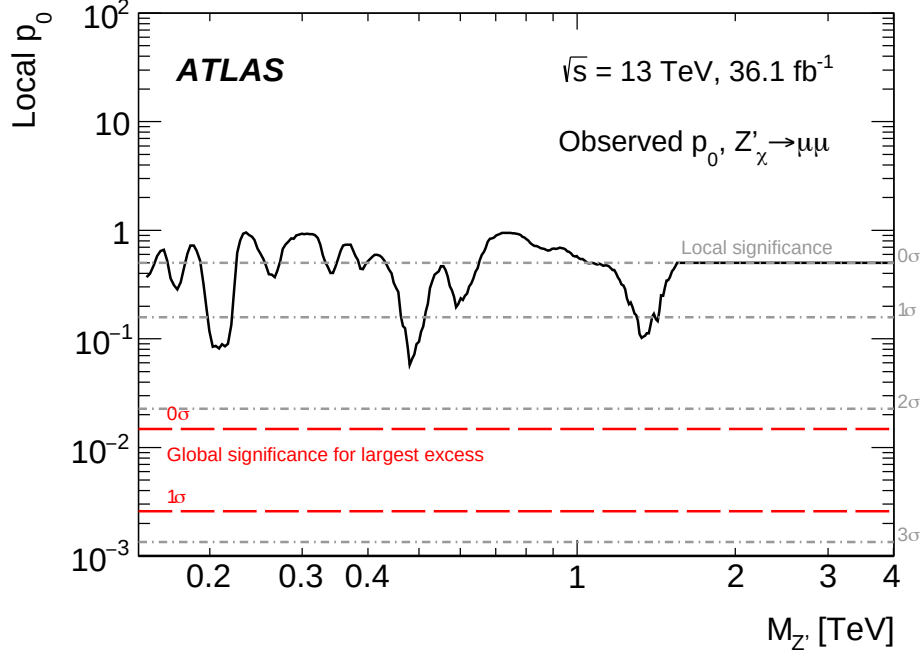


Figure 11.3: Local p -value as a function of pole mass assuming Z'_χ signal shapes. Local (global) significance levels are shown as dashed grey (red) lines.

11.2 Exclusion Limits

If no significant excess is found, exclusion limits on parameters of interest for various BSM scenarios described in Chapter 3 are set. The likelihood function used for this procedure is the same as Equation (11.5), with the expected number of events given by

$$\lambda_j = s_j(\Theta, \boldsymbol{\theta}) + b_j(\boldsymbol{\theta}), \quad (11.11)$$

where $\boldsymbol{\theta}$ still represents the set of nuisance parameters, and $s_j(\Theta, \boldsymbol{\theta})$ is the number of signal events expected from the BSM scenario under consideration, which depends on the model parameter Θ . For the resonant signal limits, Θ corresponds to a choice of the Z' pole mass, whereas in the CI signal limits, Θ corresponds to a choice of the energy scale Λ and interference parameter η_{ij} .

Bayes' theorem is used to relate the marginalized likelihood function to the posterior probability for the parameter Θ given $\mathbf{n}$ observed events:

$$\mathcal{P}(\Theta|\mathbf{n}) = \frac{1}{\mathcal{N}} \mathcal{L}_{\mathcal{M}}(\mathbf{n}|\Theta) P(\Theta), \quad (11.12)$$

where $\mathcal{N}$ is a normalization constant and $P(\Theta)$ is the prior probability density function of the model parameter Θ . For resonant signals, a uniform and positive prior $P(\Theta)$ is assumed. For CI signals, two priors are considered: $1/\Lambda^2$ and $1/\Lambda^4$. $\mathcal{L}_{\mathcal{M}}(\mathbf{n}|\Theta)$ is the marginalized likelihood, i.e., the likelihood after all nuisance parameters have been integrated out:

$$\mathcal{L}_{\mathcal{M}}(\mathbf{n}|\Theta) = \int \mathcal{L}(\mathbf{n}|\Theta, \boldsymbol{\theta}) \prod_{s=1}^{N_{\text{sys}}} d\theta_s. \quad (11.13)$$

The 95% Confidence Level (CL) limit on Θ is then obtained by finding Θ_{lim} such that

$$\int_0^{\Theta_{\text{lim}}} \mathcal{P}(\Theta|\mathbf{n}) d\Theta = 0.95. \quad (11.14)$$

These calculations are performed using the Bayesian Analysis Toolkit [97], which uses a Markov Chain Monte Carlo technique to perform the integration over the nuisance parameters. The observed limit on Θ is calculated using the data distribution $\mathbf{n}$. The expected limit on Θ is obtained by taking the median Θ_{lim} from a large sample of pseudo-experiments generated under the background-only hypothesis for the particular signal of interest under consideration.

11.2.1 Z' Cross-Section and Mass Limits

The results presented in this section use the same $m_{\mu\mu}$ binning as the one presented in Figure 9.5, with the bins in the Z boson peak region 80 – 120 GeV summed together into a single bin in order to help constrain mass-independent components of systematic uncertainties. Upper limits are placed on the product of the cross-section and the branching ratio (σB) of the BSM scenario Z' boson as a function of the Z' pole mass. Lower mass limits can be set on the mass of the theoretical Z' boson using the upper limits on σB , by plotting the observed σB as a function of Z' pole mass and finding where the observed limit intersects the theoretical line. Figure 11.4 gives the 95% CL upper limit on the product of cross-section and branching ratio as a function of Z' pole mass. Table 11.1 summarizes the observed and

expected lower limits on the mass of various Z' gauge boson models at the 95% CL.

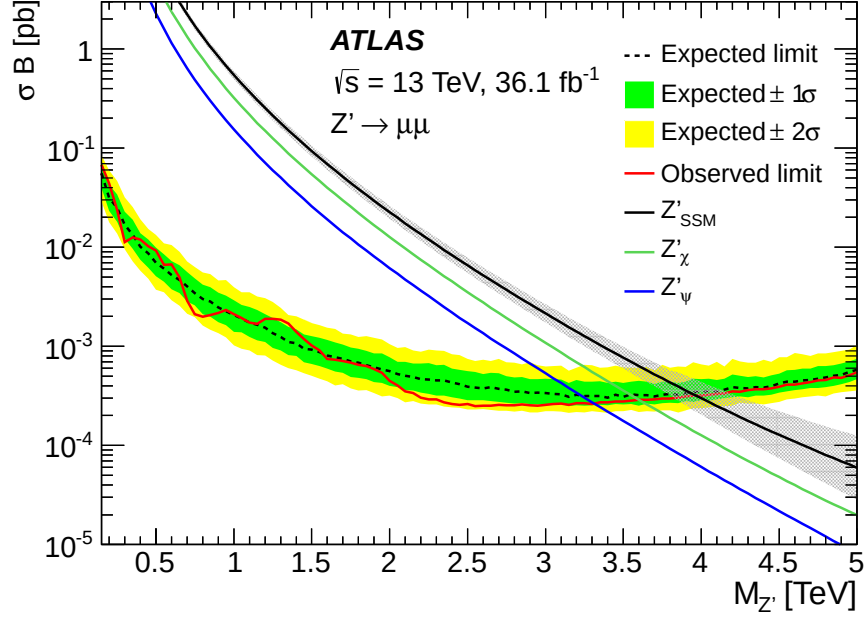


Figure 11.4: Expected (dashed black line) and observed (solid red line) upper 95% CL limits on the product of the Z' production cross-section and branching ratio to two muons as a function of Z' pole mass. Theoretical uncertainties related to the Z'_{SSM} signal are shown as a grey band around the black line for illustration purposes, but are not included in the limit calculation.

11.2. Exclusion Limits

Model	Width [%]	θ_{E_6} [rad]	Lower limits on $M_{Z'}$	
			Observed [TeV]	Expected [TeV]
Z'_{SSM}	3.0	-	4.0	3.9
Z'_χ	1.2	0.50π	3.6	3.6
Z'_S	1.2	0.63π	3.6	3.5
Z'_I	1.1	0.71π	3.5	3.4
Z'_η	0.6	0.21π	3.4	3.3
Z'_N	0.6	-0.08π	3.4	3.3
Z'_ψ	0.5	0π	3.3	3.2

Table 11.1: Observed and expected 95% CL lower mass limits for various Z' gauge boson models. The widths are quoted as a percentage of the resonance mass.

11.2.2 Minimal Z' Model Limits

As described in Section 3.1.3, minimal Z' models can be characterized by the γ' and θ_{Min} parameters, representing the strength of the Z' boson coupling relative to that of the SM Z boson and the mixing between the generators of the $(B - L)$ and weak hypercharge Y gauge groups, respectively. The 95% CL limits are placed on γ' as a function of the Z'_{Min} mass, and on γ' as a function of θ_{Min} . The results are shown in Figures 11.5 and 11.6, and summarized in Table 11.2. Limit curves in the $\gamma' - M_{Z'_{\text{Min}}}$ parameter space are shown for three representative values of the mixing angle, θ_{Min} , between the generators of the $(B - L)$ and the weak hypercharge Y gauge groups. These are: $\tan \theta_{\text{Min}} = 0$, $\tan \theta_{\text{Min}} = -2$, and $\tan \theta_{\text{Min}} = -0.8$, which correspond respectively to the Z'_{B-L} , Z'_{3R} , and Z'_χ models at specific values of γ' . The grey band envelops all observed limit curves, which depend on the choice of $\theta_{\text{Min}} \in [0, \pi]$. The corresponding expected limit curves are within the area delimited by the two dotted lines. Limits in the $\gamma' - \theta_{\text{Min}}$ parameter space are set for several representative values of the mass of the Z' boson in the minimal Z' model. In both cases, the region above each line is excluded.

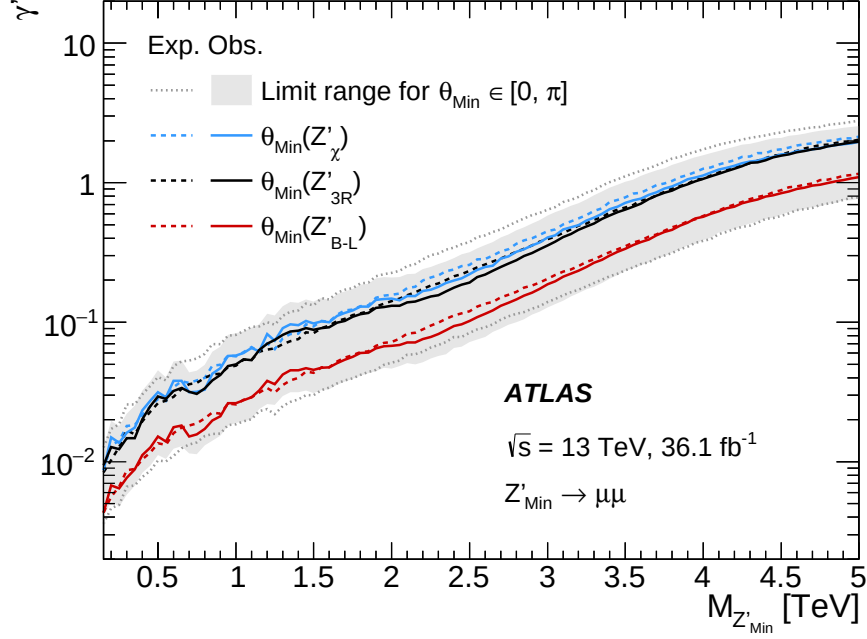


Figure 11.5: Expected (dotted and dashed lines) and observed (filled area and lines) 95% CL limits on the relative coupling strength γ' as a function of the mass of the Z' boson in the minimal Z' model. Limit curves are shown for three representative values of the mixing angle, θ_{Min} , between the generators of the $(B - L)$ and the weak hypercharge Y gauge groups. These are: $\tan \theta_{\text{Min}} = 0$, $\tan \theta_{\text{Min}} = -2$, and $\tan \theta_{\text{Min}} = -0.8$, which correspond respectively to the Z'_{B-L} , Z'_{3R} , and Z'_χ models at specific values of γ' . The region above each line is excluded. The grey band envelops all observed limit curves, which depend on the choice of $\theta_{\text{Min}} \in [0, \pi]$. The corresponding expected limit curves are within the area delimited by the two dotted lines.

11.2. Exclusion Limits

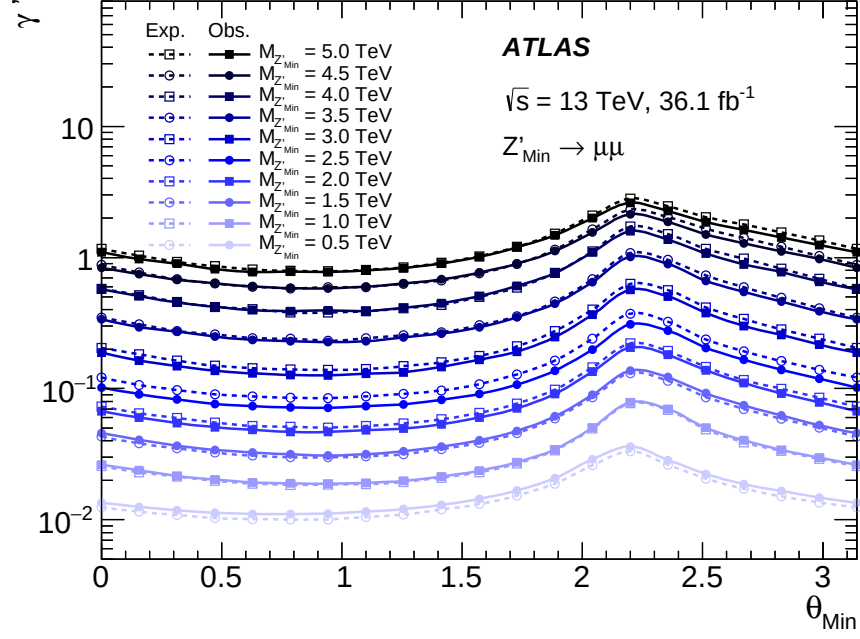


Figure 11.6: Expected (empty markers and dashed lines) and observed (filled markers and lines) limits at 95% CL on γ' as a function of θ_{Min} . The limits are set for several representative values of the mass of the Z' boson in the minimal Z' model. The region above each line is excluded.

Model	γ'	$\tan \theta_{\text{Min}}$	Lower limits on $M_{Z'_{\text{Min}}}$	
			Observed [TeV]	Expected [TeV]
Z'_χ	$\sqrt{\frac{41}{24}} \sin \theta_{\text{Min}}$	$-\frac{4}{5}$	3.4	3.3
Z'_{3R}	$\sqrt{\frac{5}{8}} \sin \theta_{\text{Min}}$	-2	3.6	3.6
Z'_{B-L}	$\sqrt{\frac{25}{12}} \sin \theta_{\text{Min}}$	0	3.6	3.6

Table 11.2: Observed and expected 95% CL lower mass limits for various minimal Z' models.

11.2.3 Limits on Contact Interactions

The binning used to set limits on CI models corresponds to eight bins above $m_{\mu\mu} = 400$ GeV of varying width from 100 GeV to 1500 GeV. The procedure described in Section 11.2 is employed, with the parameter of interest Θ representing the energy scale, Λ , and the interference parameter, η_{ij} , of the CI. All permutations of the chiral structure of the CI are investigated. Two priors are considered: $1/\Lambda^2$ and $1/\Lambda^4$. Figures 11.7 and 11.8 present the lower limits on the CI scale Λ , and the results are summarized in Table 11.3.

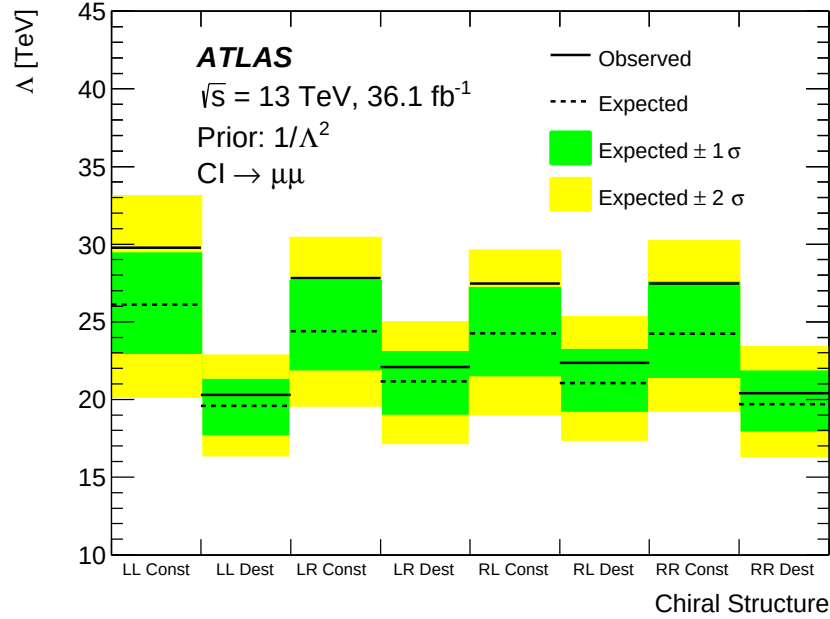


Figure 11.7: Expected (dashed black line) and observed (solid black line) lower limits on the energy scale Λ at 95% CL, for the CI model with constructive (Const) and destructive (Dest) interference for all considered chiral structures using a $1/\Lambda^2$ prior.

11.2. Exclusion Limits

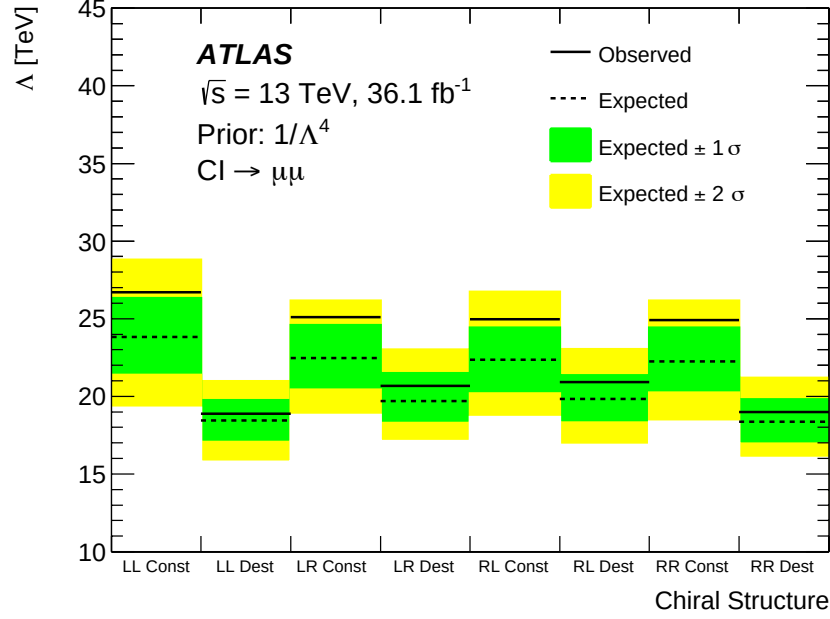


Figure 11.8: Expected (dashed black line) and observed (solid black line) lower limits on the energy scale Λ at 95% CL, for the CI model with constructive (Const) and destructive (Dest) interference for all considered chiral structures using a $1/\Lambda^4$ prior.

Channel	Prior	Lower limits on Λ [TeV]							
		Left–Left		Left–Right		Right–Left		Right–Right	
		Const	Dest	Const	Dest	Const	Dest	Const	Dest
Observed	$1/\Lambda^2$	30	20	28	22	28	22	28	20
Expected		26	20	24	21	24	21	24	20
Observed	$1/\Lambda^4$	27	19	25	21	25	21	25	19
Expected		24	18	23	20	22	20	22	18

Table 11.3: Observed and expected 95% CL lower limits on Λ for the LL, LR, RL, and RR chiral coupling scenarios, for both the constructive (Const) and destructive (Dest) interference cases using a uniform positive prior in $1/\Lambda^2$ or $1/\Lambda^4$. The limits are rounded to the nearest 100 GeV.

Chapter 12

Conclusion and Outlook

The work presented in this dissertation is centred around one common thread: muons.

In particular, the performance of the ATLAS detector with respect to high- p_T muons is evaluated. A new selection WP allowing higher selection efficiency with a modest increase in p_T resolution is investigated. The performance of the muon trigger system of the ATLAS experiment is measured using W +jets and $t\bar{t}$ samples, and trigger SFs are provided for the entire collaboration.

With the expected upgrade of the NSW during LS2 in mind, the sTGC detector technology is characterized through a series of test beam campaigns. An average spatial resolution of $45\,\mu\text{m}$ is obtained, well below the $100\,\mu\text{m}$ required by ATLAS for the NSW upgrade project. The operation of full-sized sTGC quadruplets with complete 4-layer VMM readout electronics from both pads and strips is achieved.

Finally, the main result of this dissertation describes a search for new high-mass phenomena in the dimuon final state. It is motivated by various theories of BSM physics. The search results are found to be consistent with the SM prediction. Thus, lower limits on the mass of hypothetical new particles are presented. The lower limits range from 4.0 TeV for the Z'_{SSM} model, to 3.3 TeV for the Z'_ψ model. In addition, lower limits on the CI scale Λ are set, ranging between 18 TeV and 30 TeV, depending on the chiral structure of the CI.

Future work includes analyzing the entire dataset collected during Run 2 of the LHC. This has been accomplished for the resonant search and has been published in [98]. The next logical step for this search is thus to investigate non-resonant signals using the full Run 2 dataset in order to constrain BSM scenarios predicting, e.g., CIs. This can be accomplished by searching for a broad excess in the $m_{\mu\mu}$ spectrum, or by performing the analysis of angular information of the dimuon system and examining the forward-backward asymmetry of dimuon systems. In addition, exclusive channels could provide hints of new physics that would be buried within the inclusive search signal. These could include searching in final states

containing two muons and E_T^{miss} , or two muons and b -jets.

Beyond the third run of data-taking at the LHC, up to 3000 fb^{-1} of data will be delivered by the HL-LHC. This vast dataset has the potential to uncover extremely rare processes, such as the Higgs boson decaying into muon pairs. Such an analysis would be the first measurement of the Higgs coupling to second generation fermions, allowing to probe whether or not the Higgs boson's properties deviate from the SM prediction. This measurement has the potential to uncover signs of new physics, which could manifest itself via contributions to the Higgs production and decay rates.

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