



Measurement of the W boson properties in the forward region in pp collisions at $\sqrt{s} = 5.02$ TeV

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Declaration of Authorship

This thesis is submitted to the University of Warwick in support of my application for the degree of Doctor of Philosophy. It has been composed by myself and has not been submitted in any previous application for any degree. Data analysis was performed using data collected by the LHCb experiment at the European Organisation for Nuclear Research (CERN). The research presented in this thesis was carried out under the supervision of Dr. Mika Vesterinen and Dr. Miguel Ramos Pernas.

As a non-native English speaker, I have used proof-checking tools to assist in verifying the spelling, structure, grammar, and logical consistency of the text in this thesis. These tools served as an aid in ensuring the accuracy and clarity of the work; however, the responsibility for any errors or omissions remains entirely my own. All writing and editing was done by me, with the exception of a few corrections and suggestions made by my supervisors.

The research presented here is a collaborative effort; while I completed most of it under the guidance of my supervisors, it is essential to include parts completed by other collaborators. This is clearly indicated in the text, and the contributions of others are acknowledged in the relevant sections. All sources of information have been cited and the work presented is my own, except where explicitly stated.

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THE UNIVERSITY OF WARWICK

Abstract

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Measurement of the W boson properties in the forward region in pp collisions at $\sqrt{s} = 5.02$ TeV

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This thesis primarily presents a detailed study of the fiducial cross-sections for W^+ and W^- boson production at $\sqrt{s} = 5.02$ TeV, using data from the LHCb experiment at the Large Hadron Collider at CERN. The dataset corresponds to an integrated luminosity of $99.86 \pm 2.00 \text{ pb}^{-1}$, recorded during 2017. The research focusses on muonic decays of W bosons. The cross-sections are measured differentially in intervals of the muon transverse momentum between 25 and 55 GeV, with the pseudorapidity ranging from 2.2 to 2.4. Integrated over this 2D fiducial sub-volume, the cross-sections are measured to be:

$$\begin{aligned}\sigma_{W^+ \rightarrow \mu^+ \nu_\mu} &= 301.62 \pm 3.00 \pm 3.74 \pm 6.03 \text{ pb}, \\ \sigma_{W^- \rightarrow \mu^- \bar{\nu}_\mu} &= 237.80 \pm 2.72 \pm 2.67 \pm 4.76 \text{ pb},\end{aligned}$$

where the first uncertainty is statistical, the second systematic, and the third is attributed to the luminosity determination. A fit to the differential cross-sections yields the following determination of the W boson mass:

$$m_W = 80355 \pm 128 \pm 21 \text{ MeV},$$

where the first error is experimental (statistical and systematic errors combined) and the second error is theory systematic errors for measuring m_W . It is important to note that the m_W result is currently blinded, and the analysis is in internal review within the international LHCb collaboration. This analysis serves as a proof-of-principle for determining the W boson mass with intermediate measurements of the differential cross-sections in intervals of the muon transverse momentum. The method of investigation involves novel techniques to extract the W boson mass using differential cross-sections in bins of transverse momentum. By eventually applying this approach to the full Run 2 dataset at $\sqrt{s} = 13$ TeV, we expect a precision of around 15 MeV, which would be a competitive measurement that enriches the findings of ATLAS and CMS.

Furthermore, the thesis covers the author's role in the Real-Time Analysis (RTA) project at the LHCb experiment. This involves developing data processing software and optimising workflows for the Run 3 data collection phase.

List of Abbreviations

ALICE	A Large Ion Collider Experiment
ATLAS	A Toroidal LHC ApparatuS
BBLite	Beeston–Barlow Lite
BSM	Beyond the Standard Model
CDF	Collider Detector at Fermilab
CERN	Conseil Européen pour la Recherche Nucléaire (European Organization for Nuclear Research)
CMS	Compact Muon Solenoid
CP	Charge conjugation and Parity symmetries
DEC	Decision Efficiency
ECAL	Electromagnetic CALorimeter
EFF	Event Filter Farm
EW	ElectroWeak
FSR	Final State Radiation
GEC	Global Event Cut
GEM	Gas Electron Multiplier
GPU	Graphics Processing Unit
HCAL	Hadronic CALorimeter
HL–LHC	High Luminosity LHC
HLT	High Level Trigger
HPD	Hybrid Photon Detectors
IP	Impact Parameter
IT	Inner Tracker
L0	Level 0
LEP	Large Electron–Positron Collider
LHC	Large Hadron Collider
LHCb	Large Hadron Collider beauty experiment
LINAC	LINear ACcelerator
LO	Leading Order
LS	Long Shutdown
MaPMT	Multi-anode Photo Multiplier Tube
MC	Monte Carlo simulation
MWPCs	Multi-Wire Proportional Chambers
ndf	number of degrees of freedom
NLO	Next-to-Leading Order
NNLO	Next-to-Next-to-Leading Order
N3LO	Next-to-Next-to-Next-to-Leading Order
NLL	Next-to-Leading Logarithmic
NNLL	Next-to-Next-to-Leading Logarithmic
OT	Outer Tracker
PDF	Parton Distribution Function

PDG	Particle Data Group
PID	Particle IDentification
PLUME	Probe for LUminosity MEasurement
PMT	PhotoMultiplier Tube
pQCD	perturbative Quantum ChromoDynamics
pp	proton–proton
PS	Pre Shower
PS	Proton Synchrotron
PV	Primary Vertex
QCD	Quantum ChromoDynamics
QED	Quantum ElectroDynamics
RMS	Root Mean Square
RTA	Real–Time Analysis
RICH	Ring-Imaging CHerenkov
SciFi	Scintillating Fibre
SM	Standard Model
SMOG	System for Measuring the Overlap with Gas
SPD	Scintillating Pad Detector
SPS	Super Proton Synchrotron
ST	Silicon Trackers
SVD	Singular Value Decomposition
TIS	Trigger Independent of Signal
TOB	Trigger On Both
TOS	Trigger On Signal
TORCH	Timing Of internally Reflected CHerenkov photons
TT	Tracker Turicensis
UT	Upstream Tracker
VELO	Vertex LOcator

In loving memory of my grandmother, Fawzeyya Abdelwahed.

Chapter 1

Introduction

“There is only one way to eat an elephant, a bite at a time.”

– Desmond Tutu

Particle physics represents the epitome of human curiosity and scientific ambition, with the aim of driving our understanding of the universe to its most fundamental level. It is a field that probes both the smallest building blocks of matter and the largest forces shaping the cosmos [1]. By studying elementary particles and their interactions, particle physicists seek to unravel the mysteries of existence, exploring questions about the origin, structure, and fate of the universe. The pursuit of these questions requires not only cutting-edge technology and sophisticated experiments but also a global collaboration of scientists working together toward a common goal. This collective effort exemplifies the heights of human achievement, showing how far we have come in our quest to comprehend the fundamental nature of reality.

The Large Hadron Collider (LHC) at CERN, the world’s largest and most powerful particle accelerator, is a prime example of human ingenuity. The LHC is capable of generating temperatures more than 100,000 times hotter than the core of the Sun during its particle collisions, recreating conditions that existed just fractions of a second after the Big Bang. These extreme conditions allow scientists to probe the fundamental forces and particles that shaped the early universe, offering insights into the nature of matter and energy at the most basic level [2].

At the heart of modern particle physics lies the Standard Model, a theoretical framework that describes electromagnetic, weak, and strong forces; three of the four known fundamental forces in the universe [1]. Developed in the second half of the twentieth century, the Standard Model has been remarkably successful in predicting and explaining a wide array of experimental results. It classifies all known elementary particles, including quarks, leptons, and force-carrying bosons such as photons, gluons, and W and Z bosons [1]. Despite its success, the Standard Model does not answer every question; it does not incorporate gravity (the fourth known fundamental force in the universe) and leaves many fundamental questions unanswered, such as the nature of dark matter and matter–antimatter asymmetry in the universe, the hierarchy problem, and the origin of neutrino masses [1]. As a result, physicists are constantly testing the Standard Model

by looking for deviations from its predictions that could point to “new physics” beyond the current model. The search for such deviations drives much of the research in particle physics today, as breaking the Standard Model could lead to the discovery of new particles, forces, and interactions; unlocking a deeper understanding of the universe.

There are two primary ways to challenge and potentially “break” the Standard Model: through direct searches for new particles, and through precision measurements of known particles and their interactions. Direct searches aim to discover new particles that the Standard Model does not contain, such as those predicted by theories like supersymmetry, which is an extension of the Standard Model that aims to fill some of the gaps it has [1]. Precision measurements also search for new physics by testing the Standard Model’s predictions with the highest possible accuracy, looking for small deviations that might indicate the presence of new physics beyond the Standard Model. This thesis focusses on the latter approach, leveraging precision measurements to probe the validity and limitations of the Standard Model.

Electroweak physics, a cornerstone of the Standard Model of particle physics, governs the interactions between fundamental particles via electromagnetic and weak forces. The theory, initially unified in the 1960s by Sheldon Glashow, Abdus Salam, and Steven Weinberg, was revolutionary in its successful description of these forces as different manifestations of a single underlying interaction [3, 4, 5]. This unification, supported by extensive experimental evidence, solidified the Standard Model as the prevailing framework for understanding particle interactions. Central to electroweak theory are the W and Z bosons, mediators of the weak force, whose precise properties continue to provide critical insights into the structure of the universe. Electroweak spontaneous symmetry breaking, the process by which the Higgs mechanism imparts mass to the weak bosons and fermions, is another key aspect of the theory and was recently confirmed by the discovery of the Higgs boson at the LHC in 2012 [6, 7].

The Large Hadron Collider at CERN, with its four major experiments—ATLAS, CMS, LHCb, and ALICE—is trying to advance our current understanding of electroweak physics. While ATLAS and CMS are general-purpose detectors designed to explore a broad range of phenomena, LHCb is specialised for studying particles containing b - and c quarks. However, LHCb’s unique forward geometry also makes it well suited for precision electroweak measurements in a kinematic regime less accessible to ATLAS and CMS [2].

The main topic of this thesis is the measurement of the W boson cross-section at 5.02 TeV using data from the LHCb experiment, which then leads to the measurement of the W boson mass. We take advantage of the distinctive capabilities of LHCb to contribute to a better understanding of electroweak interactions.

Measurement of the cross-section and mass of the W boson with high precision is crucial for several reasons. First, the W boson mass is a key parameter of the Standard Model, intricately linked to the masses of the Higgs boson and the top quark through higher-order corrections [8]. Any deviation from the calculated mass value could hint at new physics beyond the Standard Model, potentially uncovering phenomena such as supersymmetry or additional Higgs bosons [9, 10].

The W boson mass has traditionally been measured by fitting a model (including detector simulation) to the raw data at the reconstructed level. This poses a significant challenge for later reinterpretation and analysis using different models. Scientists must either standardise the measurements (i.e. correct for the detector response) to facilitate combinations, or alternatively, adjust their theoretical predictions to align with the experimental results, both of which are no easy tasks [11].

This thesis aims to provide a new perspective. The detailed analysis is a proof-of-principle for a new method to measure the W boson mass using the W differential cross-section in bins of muon transverse momentum.

In addition to determining the W mass, cross-section measurements, especially in the forward region observed by LHCb, offer vital data that help improve parton distribution functions (PDFs), which are crucial for understanding the internal structure of the proton [12]. Precise measurements of these quantities also play a vital role in constraining theoretical models and improving the accuracy of predictions for other Standard Model processes.

The thesis explores new methodologies using the LHCb experiment to measure the cross-section and mass of the W boson, highlighting the challenges and innovations associated with these measurements in a forward detector environment. The implications of the new results for our understanding of the Standard Model and potential new physics will also be discussed, providing a comprehensive overview of the significance of electroweak precision measurements in the current landscape of particle physics. The thesis will also delve into the broader context of electroweak physics, discussing the theoretical underpinnings of the Standard Model, the experimental techniques used to probe its predictions, and the future prospects for exploring new physics beyond the established framework. In addition, the thesis includes supplementary technical contributions by the author for the data processing team within the LHCb experiment. This contribution encompasses software development, data analysis, and participation in collaborative efforts to further the goals of the experiment.

In this chapter, we provide an outline of the thesis in Section 1.1, followed by a discussion of some of the key theoretical concepts and conventions used throughout the document in Section 1.2. We then delve into the motivation behind the analysis of the W boson mass in Section 1.3, highlighting the importance of these measurements for testing the Standard Model and exploring new physics beyond it. Following that, we briefly discuss previous measurements of the W boson cross-section and mass, setting the scene for the current landscape in Section 1.4. Finally, we outline the analysis strategy employed in this thesis, focussing on the innovative methods used to measure the cross-section and mass of the W boson in the LHCb experiment in Section 1.5.

1.1 Thesis outline

The thesis is structured as follows:

- **Chapter 2: Theoretical Overview** provides an overview of the theoretical foundations of particle physics, including the Standard Model, electroweak physics, the W boson, and particle distribution functions.
- **Chapter 3: The LHCb Experiment at the LHC** describes the Large Hadron Collider at CERN and LHCb experiment. It highlights its unique features and the detector components relevant to the measurement of the W boson cross-section and mass. It also mentions the upcoming upgrades to the LHCb experiment and their implications for future measurements.
- **Chapter 4: Contributions to the Real-Time Analysis project at LHCb** details the author's contributions to the LHCb experiment, including software development, data analysis, and participation in collaboration activities. Especially, the author's work on the Real-Time Analysis project, which aims to provide fast access to LHCb data for real-time monitoring and analysis.
- **Chapter 5: Analysis Setup** outlines the data selection criteria, event reconstruction, calibrations and corrections, and the methodology used to measure the W boson cross-section and mass in the LHCb experiment.
- **Chapter 6: W boson differential cross-section measurement** outlines the process used to measure the W boson differential cross-section in bins of the muon transverse momentum, focusing on the isolation variable as a key parameter for background subtraction, unfolding for detector biases and inefficiencies, and signal extraction. It also shows the results and comparisons with theoretical predictions.
- **Chapter 7: W boson mass fit** details the fitting procedure used to extract the W boson mass from the unfolded differential cross-section measurements including uncertainties and systematic effects.
- **Chapter 8: Summary and Outlook** summarises what was discussed in the thesis, mentions the implications of the results for the field of particle physics, and outlines future prospects for precision measurements at the LHCb experiment and beyond.

1.2 Definitions and conventions

Further details on all subsequent definitions can be found in Ref. [8, 13].

- We use natural units throughout, so, the speed of light in vacuum, c , is set to unity, and the reduced Planck constant, $\hbar$, is also set to unity. The energy and momentum of the particles are expressed in units of GeV or MeV.
- The symbol $\sqrt{s}$ denotes the center-of-mass energy of a pp collision.
- Detectors like those at the LHC are optimised to measure particles' momentum components perpendicular to the beamline (transverse plane). The longitudinal momentum (along the beamline) is often less accurately measured due to the limitations in detector design and the fact that particles can escape the detection volume in the beam direction.

The transverse momentum (p_T) is also invariant under boosts along the beamline, which means that it does not change if the entire system is moving parallel to the beam. This invariance simplifies the analysis of particle collisions and interactions.

- Rapidity (y) is a measure that approximates a particle's scattering angle with respect to the beam axis at high energies. Its main benefit lies in its invariance under Lorentz boosts along the beam axis, making it essential for analyzing particle kinematic distributions in collider experiments across different reference frames. It is given by $y = \frac{1}{2} \log \left(\frac{E+p_z}{E-p_z} \right)$, where E represents the particle's energy, and p_z is the beam-direction momentum component. This facilitates an understanding of the angular and momentum properties of particles generated in collisions, especially along the beam trajectory.
- Pseudorapidity (η) is a quantity that approximates rapidity in the limit of high energies. It is defined as $\eta = -\log \left(\tan \left(\frac{\theta}{2} \right) \right)$, where θ is the angle of the particle relative to the beam axis. It is often used in particle physics experiments to describe the angular distribution of particles in the detector. For massive particles, it is not invariant under boosts along the beamline. However, pseudorapidity is a useful approximation for rapidity in the high-energy or massless limit, where the particle's energy E is much greater than the longitudinal component of its momentum; $E \gg m$. It is a convenient way to express one degree of freedom of the direction of an ultra-relativistic final-state particle.
- The symbol ϕ represents the azimuthal angle of a particle in the detector.
- The symbol ΔR represents the angular separation between particles in the detector, defined as the Euclidean distance in the $\eta - \phi$ plane: $\Delta R = \sqrt{(\Delta\eta)^2 + (\Delta\phi)^2}$,
- The symbol $\mathcal{I}$ represents the isolation parameter, which characterises the separation between a particle of interest (e.g., a muon) and other particles in the detector. It is defined as the scalar sum of the transverse momenta of neutral and charged particles within a cone of radius $\Delta R < 0.4$ around the particle of interest. We typically express it in units of GeV.
- The symbol m_W represents the mass of the W boson, a fundamental parameter of the Standard Model.
- The symbol σ represents the cross-section of a particle interaction, where we use the barn unit (b) to measure the cross-section.
- The integrated luminosity is the integral of the instantaneous luminosity over time, and it is measured in units of cm^{-2} . The equation for the integrated luminosity is given by $L = \int \mathcal{L} dt$, where $\mathcal{L}$ is the instantaneous luminosity.
- $\mathcal{L}$ represents the luminosity of a particle collider.
- The primary vertex (PV) is the point of interaction between two colliding particles in a collider experiment. The impact parameter (IP) is the perpendicular distance between the PV and a particle's trajectory, evaluated in either the transverse plane (2D) or full 3D space. In tracking, the point of closest approach (POCA) of a particle's trajectory to a reference

point, such as the beamline or PV, is not uniquely defined. The 2D POCA, determined in the transverse plane, generally differs from the 3D POCA, which considers the full spatial geometry. Once the POCA is determined under either definition, the distance of closest approach (DOCA) can also be evaluated in 2D or 3D. At LHCb, selection algorithms use the 3D POCA and corresponding 3D DOCA for IP calculations. The PV and IP are essential for identifying the origins of particles in the detector.

- A track is the trajectory of a charged particle as it passes through the detector. Track fit quality is a measure of how well a reconstructed particle track matches the recorded detector hits from the actual trajectory of a charged particle. It is used to identify and reject poorly reconstructed tracks.
- Particle identification (PID) is the process of determining the type of particle based on its interactions with the detector material. It is crucial for distinguishing between different particle species in a collider experiment.

1.3 Analysis motivation

The W boson is an elementary particle that plays a crucial role in the Standard Model of particle physics. It is a charged boson that mediates the charged weak force, which is one of the four fundamental forces of nature. The mass of the W boson is a fundamental parameter of the Standard Model, and its precise measurement is crucial to testing the theory and exploring new physics effects that may deviate from it. The mass of the W boson determines the range and strength of the weak force and affects other quantities such as the decay rates of the particles [8]. The accuracy of the W boson mass measurement is essential for the “Global Electroweak Fit” [9, 10]. It is also highly sensitive to BSM physics [9, 10].

The W boson mass and the Higgs boson mass are related in the Standard Model, and the Higgs mass constraint dictates that the Higgs boson mass should be less than about 170 GeV to prevent a breakdown of the electroweak–symmetry–breaking mechanism. The precise measurement of the W boson mass can therefore serve as a valuable test of the Standard Model and its predictions [14].

Moreover, the W boson mass can help to search for new physics beyond the Standard Model. Certain extensions of the Standard Model suggest the existence of additional heavy particles that can influence the mass of the W boson. By measuring the mass of W bosons with increasing precision, experiments at the Large Hadron Collider and elsewhere can explore the presence of such new particles and test the validity of the Standard Model. Finally, the mass of the W boson can also be studied in the context of the early universe, where it played a role in the electroweak phase transition that separated the weak force from the electromagnetic force [14].

1.4 Previous measurements and limitations

The differential cross-section of the W boson was measured at LHCb with data from 7 TeV and 8 TeV in 2014 [11] and 2015 [15], respectively. At LHCb, measurements of the W production

cross-section could not be split into bins of p_T because the shape of the p_T distribution was used to subtract the background [11, 15], i.e. comparing models with detector simulation to raw counts of W candidates in the data. This limits the analysis to a single theoretical model and makes it extremely difficult to reinterpret and analyse later with different models. The cross-section was not measured at LHCb with 5.02 TeV and 13 TeV data. However, the cross-section of the W boson was measured at ATLAS with 5.02 TeV data [16].

Over the years, the mass of the W boson has been measured with increasing precision. Theoretical predictions, based on a global fit of electroweak observables that excludes direct measurements of m_W , yield a value of $m_W = 80354 \pm 7$ MeV [17]. Experimentally, several measurements of the mass of the W boson have been made, including the Large Electron-Positron Collider (LEP) measurement with an uncertainty of 33 MeV [18], the latest ATLAS measurement with an uncertainty of 16 MeV [19], and the most precise measurement to date, conducted by the Collider Detector at Fermilab (CDF), with an uncertainty of 9.4 MeV [20]. Furthermore, a proof-of-principle measurement at LHCb, with limited data from 2016, produced a result with an uncertainty of 32 MeV [21]. However, the CDF result is inconsistent with all of the other aforementioned measurements, resulting in a heated debate within the electroweak physics community. Some researchers have pointed out that certain assumptions or choices in the use of the RESBOS program [22] and the associated PDFs could introduce uncertainties that could affect the mass measurement [23]. One of the challenges is that the specific implementation and treatment of theoretical uncertainties of the CDF collaboration are difficult to reproduce exactly, leading to ongoing discussions in the high-energy physics community about the robustness and interpretation of the CDF result [24]. This issue highlights the importance of transparency and reproducibility in such critical measurements. Also, the importance of being model-independent in the measurement of the W boson mass is crucial to avoid biases and to allow for a more direct comparison with theoretical predictions.

All of these aforementioned experimental measurements, including the inconsistencies between some of them, highlight the need for continued efforts to increase the precision of the W boson mass measurement, which is crucial to testing the Standard Model and searching for new physics.

1.5 Analysis strategy

In this study, we examine W bosons that decay into leptons and neutrinos. Assuming a stationary W boson before decay, the energy of the W boson is anticipated to be almost evenly divided between the resulting particles: the lepton and the neutrino. For example, if the W boson mass is about 80 GeV, we would expect the lepton and neutrino to each have an energy of around 40 GeV, given their negligible rest masses in comparison. Since their masses are so small, this energy translates to a momentum close to 40 GeV for each particle. However, neutrinos cannot be "reconstructed" or detected in our LHCb setup, whereas we can reconstruct the leptons. To probe the W boson mass, we focus on the transverse momentum (p_T) of the charged lepton, which is the component of the lepton's momentum perpendicular to the beam axis. This choice is particularly important because our detector has limited coverage along the beam direction,

making it challenging to measure the full three-dimensional momentum of the particles. However, p_T remains well defined and measurable in the detector coverage. Using the charged lepton transverse momentum spectrum (p_T^l), we can probe the W boson mass due to the unique “Jacobian peak” [25]. The p_T^l distribution resulting from the decay of the W boson exhibits a peak near half the mass of the W boson. This characteristic Jacobian peak allows for the extraction of details regarding the W boson mass through the p_T^l spectrum’s shape.

Consider a simple $2 \rightarrow 2$ scattering process, where a fermion-antifermion pair ($f\bar{f}$) produces a vector boson V , which then decays back into fermions ($f\bar{f}$). We also assume that the particles involved are massless, which means that there are only two degrees of freedom: the scattering angle θ and the energy transfer. To simplify further, we can assume the energy transfer is fixed by considering a boson with zero decay width. Thus, we are left with only one degree of freedom: the scattering angle.

Now, consider the transverse momentum (p_T) of the boson decay products (for example, a muon and a neutrino). Since both are massless, the transverse momentum of both the muon and the neutrino is given by $p_T = \frac{M_V}{2} \sin \theta$, where M_V is the mass of the vector boson and θ is the scattering angle. In this case, the differential cross-section with respect to the scattering angle θ is proportional to $\frac{d\sigma}{d\cos\theta} \propto 1 \pm \cos^2 \theta$.

However, in many cases, we are more interested in how the cross-section depends on the transverse momentum p_T of the decay products, rather than the scattering angle. To transform the differential cross-section from θ to p_T , we need to perform a chain-rule operation $\frac{d\sigma}{dp_T} = \frac{d\sigma}{d\cos\theta} \cdot \frac{d\cos\theta}{dp_T}$.

The final distribution in p_T then becomes:

$$\frac{d\sigma}{dp_T} \propto \frac{d\sigma}{d\cos\theta} \frac{1}{\sqrt{1 - 4\frac{p_T^2}{M_V^2}}}.$$

The final term presents a singularity at $p_T = m_V/2$, resulting in the distinctive maximum known as the Jacobian peak within the transverse momentum distribution of W -boson decay products. The location of this peak enables the measurement of vector boson masses, including the mass of the W boson [26].

The Jacobian peak can be seen in Figure 1.1. The theoretical distribution has no smearing; however, in reality, the observed peak is smeared by detector effects and other inaccuracies. Additionally, the width of the peak is affected by the finite decay width of the vector boson, which we had originally assumed to be zero. As a result, the observed peak has a larger width compared to the ideal distribution [26, 25].

Hence, this analysis aims to improve upon previous analyses by examining the p_T^l spectrum in $W \rightarrow \mu\nu_\mu$ decays and investigating the feasibility of background subtraction to determine the mass of the W boson by measuring the differential cross-section in intervals of transverse momentum. Background subtraction involves constructing a model of the data using simulation templates of the signal and various backgrounds. These templates are combined and adjusted to

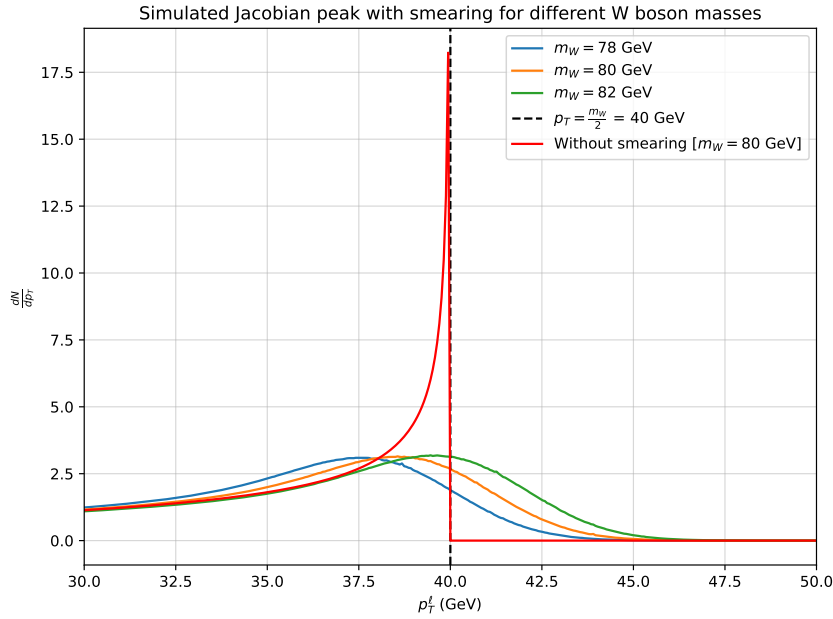


FIGURE 1.1: The graph illustrates the Jacobian peak of a theoretical W boson particle, with a mass of 80 GeV. It also includes several instances of a smeared lepton transverse momentum (p_T^l) distribution.

match the data, allowing us to separate the W boson signal from the background contributions. We additionally concentrate on “unfolding” to mitigate detector biases, aiming to derive the “true”, unbiased p_T^l distribution. Further details on unfolding are provided in Section 6.2. This approach is therefore model-independent and allows for an easier way to compare the results with different theoretical models and other experiments. This analysis is a proof-of-principle for a new method for measuring the W boson mass, proving that a different approach is indeed possible. We used a small dataset to test the method, at $\sqrt{s} = 5.02$ TeV, which is statistically limited. However, this signals the potential for an analysis with a larger dataset and signals to the scientific community that this is a viable method for future measurements.

To reiterate, our primary goal is to create a technique for background subtraction and detector response correction that does not rely on any model of the charged-lepton p_T distribution in W production to obtain a measurement of the differential cross-section of the W boson, which then can be used to measure the W boson mass. Hence, the use of a variable other than p_T is necessary to obtain a separation between the signal and background distributions. That variable is the isolation variable, which was previously defined in Section 1.2. Our signal exhibits higher isolation (indicated by a smaller isolation value) compared to the backgrounds, aiding in discrimination. After performing the background subtraction and unfolding steps, we can extract the differential cross-section in intervals of p_T . We then perform an additional fit using the unfolded differential cross-section to extract the mass of the W boson.

This gives more power to the community, particularly theorists, to fit m_W without being constrained by the uncertainties of the detector. Additionally, measuring the W boson mass in this

way allows for a more straightforward comparison with other experiments and theoretical predictions, and it can be a significant contribution toward an LHC-wide combination of W boson mass measurements.

Would such an approach be viable as a scientifically better way to measure m_W with the full LHCb Run 2 dataset and even larger future datasets? That is what we intend to study.

Chapter 2

Theoretical Overview

“Theories are nets cast to catch what we call ‘the world’: to rationalise, to explain, and to master it. We endeavour to make the mesh ever finer and finer.”

– Karl Popper

In this chapter, we will examine the theoretical foundations of Particle Physics that play an essential role in shaping and informing the context of this project. Our exploration will cover key concepts and principles in Particle Physics, providing a comprehensive understanding of the theoretical framework that underpins the research and development undertaken in this thesis.

In Sections 2.1, we will investigate the theoretical basis of Particle Physics and the Standard Model. This includes discussions of fundamental particles, forces, and the role of Lagrangians. In Section 2.2, we provide a more detailed and formal account of the Standard Model, including the mathematics behind theories such as quantum electrodynamics, quantum chromodynamics, and more complex “higher-order” theoretical frameworks. Section 2.3 will revolve around Parton Distribution Functions (PDFs), detailing their role in high-energy physics, Bjorken x , QCD factorisation, and distinctions between experiments such as ATLAS, CMS, and LHCb at the Large Hadron Collider (LHC) experiment at CERN. We will investigate the Drell–Yan process and angular coefficients, as well as how to generate simulations for them. In addition, we will examine the software used to create these simulations. We will also discuss measuring the cross-section and mass of the W boson, examining production, decay, and their connections to other fundamental constants and particle masses. We will also discuss the implications of new physics and sources of theoretical uncertainty.

2.1 A simple description of the Standard Model of Particle Physics

The Standard Model of Particle Physics, often referred to simply as the Standard Model (SM), is a foundational theoretical framework that describes the fundamental building blocks of the universe and their interactions. It encompasses three of the fundamental forces of nature: electromagnetism, the strong-nuclear force, and the weak-nuclear force, unifying them under a common theoretical framework. The Standard Model categorises particles into two main classes:

fermions, which make up matter, and bosons, which mediate the forces between them. This elegant theory has been remarkably successful in explaining and predicting the behaviour of subatomic particles, though it leaves some fundamental questions unanswered, such as the nature of dark matter and the unification of all forces under a single theory. The main sources for this section and the following are [1, 27].

According to the Standard Model, matter is composed of twelve fermions. These particles possess an inherent characteristic called “spin”, which for these fermions is defined by the value $1/2$ ¹. As far as we know, they are elementary particles, rather than being composed of other particles. These fermions are divided into six leptons and six quarks, which are further divided into three generations.

There are charged and neutral leptons. Charged leptons have whole-number charges, and quarks have fractional charges. In each generation of quarks, there is an “up-type” quark (u, c, t) with an electric charge of $2/3$ and a “down-type” quark (d, s, b) with a charge of $1/3$. As for leptons, each generation includes a down-type charged lepton (-1) (e, μ , τ) and an up-type electrically neutral neutrino (ν_e , ν_μ , ν_τ). Furthermore, each fermion in the Standard Model has a corresponding anti-fermion partner with the same mass and spin, but with an opposite charge. The fundamental fermions and bosons that are included in the Standard Model are summarised in Figures 2.1 and 2.2, respectively.

As mentioned earlier, when it comes to forces, only three of the four known fundamental forces are included in the Standard Model: electromagnetism, the strong-nuclear force, and the weak-nuclear force. These forces are thought to be mediated by spin-1 vector bosons.

- **Electromagnetic force** (Photon – γ): The electromagnetic force is responsible for interactions between charged particles, such as electrons and protons. It is mediated by a massless vector boson called the “photon” (γ). Photons are responsible for electromagnetic phenomena such as electricity, magnetism, and the transmission of light as electromagnetic radiation. In the Standard Model, the electromagnetic force is described by Quantum Electrodynamics (QED).
- **Strong-Nuclear force** (Gluon – g): The strong-nuclear force binds quarks together to form protons, neutrons, and other so-called “hadrons”. It is mediated by massless vector bosons called “gluons” (g). Unlike photons, gluons themselves carry a colour charge (associated with the strong force) and interact with quarks and other gluons. Quantum Chromodynamics (QCD) is the theory that describes the strong-nuclear force in the Standard Model. Note that this force distinguishes between quarks and leptons, as only quarks can interact via the strong force.
- **Weak-Nuclear force** (Weak-current bosons – $W^\pm$, Z): The weak-nuclear force is responsible for processes such as β decay (for example, the decay of a neutron into a proton, electron, and antineutrino). It is the force responsible for the decay of heavy particles and also crucial for enabling particle transformations that are essential in stellar processes, such as

¹As a reminder, we are using natural units with $\hbar$ set to unity.

Leptons					Quarks		
	Particle	Q	mass/GeV		Particle	Q	mass/GeV
First generation	electron (e^-)	-1	0.0005		down (d)	-1/3	0.003
	neutrino (ν_e)	0	$< 10^{-9}$		up (u)	+2/3	0.005
Second generation	muon (μ^-)	-1	0.106		strange (s)	-1/3	0.1
	neutrino (ν_μ)	0	$< 10^{-9}$		charm (c)	+2/3	1.3
Third generation	tau (τ^-)	-1	1.78		bottom (b)	-1/3	4.5
	neutrino (ν_τ)	0	$< 10^{-9}$		top (t)	+2/3	174

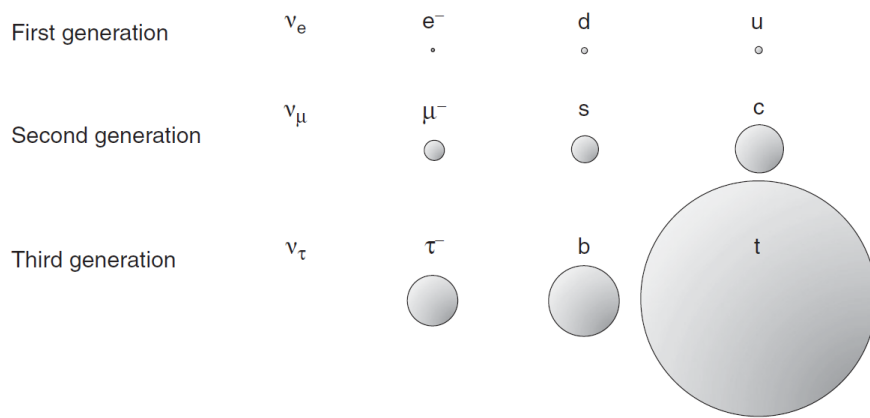


FIGURE 2.1: Fundamental fermions divided into quarks and leptons with the inclusion of their charges and masses. Reproduced from Ref. [1]

Force	Strength	Boson		Spin	Mass/GeV
Strong	1	Gluon	g	1	0
Electromagnetism	10^{-3}	Photon	γ	1	0
Weak	10^{-8}	W boson	$W^\pm$	1	80.4
		Z boson	Z	1	91.2
Gravity	10^{-37}	Graviton?	G	2	0

FIGURE 2.2: The four known forces of nature and the boson carriers of these forces. Note that Gravitons are just hypothetical particles and have not been observed, just proposed as the force carriers of the gravitational force. Reproduced from Ref. [1]

					strong	electromagnetic	weak
Quarks	down-type	d	s	b	✓	✓	✓
	up-type	u	c	t			
Leptons	charged	e^-	μ^-	τ^-		✓	✓
	neutrinos	ν_e	ν_μ	ν_τ			✓

FIGURE 2.3: The forces experienced by fermions. Reproduced from Ref. [1]

the proton–proton chain in stars such as the Sun. This force is mediated by massive vector bosons, the $W^\pm$ (typically labelled W) and the Z bosons. The exchange of W and Z bosons is responsible for inducing weak interactions, which are inherently short–ranged because of the large masses of the bosons. The theory governing the weak–nuclear force is often referred to as “Electroweak Theory” and involves the unification of electromagnetic and weak forces at high energies. The neutral leptons, neutrinos, can only interact via the weak force.

- **Gravitational force** (Graviton – hypothesised): While not part of the Standard Model, the gravitational force is one of the four fundamental forces in the universe. It is significantly weak on everyday scales (though it can become dominant in areas with high mass concentration, like near black holes). A consistent quantum theory of gravity is yet to be defined.

A diagram showing the forces experienced by different fermions can be found in Figure 2.3. The last particle to be included in the Standard Model was the Higgs Boson (H), which was eventually discovered at CERN in 2012 and was often referred to as “the final missing piece of the Standard Model”. It is unique in that it is the only fundamental scalar (spin-0) particle observed in nature. The Higgs boson is critical to the process of providing mass to other fundamental particles in the universe, which will be discussed in detail in Section 2.2.4.

2.2 Formal description of the Standard Model

2.2.1 The Standard Model formalism

The theoretical basis of the Standard Model of Particle Physics relies on the Lagrangian framework, fields, and gauge symmetry principles. The Lagrangian, an essential mathematical construct, captures the dynamics and interactions of particles in the theory. When discussing Lagrangians, we generally refer to the Lagrangian density, which is represented by $\mathcal{L}$. However, in practice, it is more common to use the term “Lagrangian” for simplicity. This is the term that we shall use hereafter. A field is also a concept that represents a physical quantity represented over space and time. The terms in the Lagrangian density include derivatives and products of the fields. Unlike classical particles that have concrete positions, a field has points everywhere in space–time. Fields give rise to particles, where particles can be described as excitations (quanta) in these fields. So, for example, an electromagnetic field can give rise to photons, and a Higgs field gives rise to the Higgs boson. We can have scalar fields, vector fields, or even “spinor”

fields, depending on their properties under transformations such as rotations and translations in space–time. A “quantum field theory” conforms to fundamental symmetry patterns, the so-called “gauge” symmetries. A gauge symmetry is a type of symmetry in a physical theory that allows specific transformations of fields without altering the physical predictions of the system; which is obviously crucial because the laws of physics must be consistent and the same throughout the universe. Hence, the contents of the Lagrangian are tightly constrained by the gauge symmetries. These symmetries encompass gauge invariance, responsible for the behaviour of forces, Lorentz invariance, which accounts for relativistic effects, and Poincaré invariance, encompassing space–time symmetries.

The Standard Model gauge symmetries are encapsulated by the groups $SU(3) \times SU(2) \times U(1)$, where $SU(3)$ describes the strong–nuclear force, $SU(2)$ characterises the weak force, and $U(1)$ represents electromagnetism. The “ $\times$ ” symbol signifies the direct product of the two gauge groups, indicating that the group elements of each group commute with each other. These gauge symmetries are crucial to understanding the fundamental forces and interactions among particles, providing a robust framework for our understanding of the subatomic world. Gauge symmetries allow us to perform field transformations so that the Lagrangian of the system remains invariant. We will talk more about these groups in the coming sections.

2.2.2 Dirac fermions

First, we start with the description of fermions. The Lagrangian density for a free massive fermion is described by a so-called “Dirac field”, which is a mathematical representation of particles with spin- $\frac{1}{2}$. It is given by:

$$\mathcal{L} = \bar{\psi}(i\gamma^\mu \delta_\mu - m)\psi, \quad (2.1)$$

where γ^μ are the Dirac gamma matrices (a set of four complex 4×4 matrices with indices from 0 to 3) which account for spin, ψ is the wavefunction (the field) of the particles, and m is the mass of the fermion. We define $\bar{\psi} = \psi^\dagger \gamma^0$, where $\psi^\dagger$ is the hermitian conjugate of ψ . The term $i\gamma^\mu \delta_\mu$ is the kinetic term, which describes the dynamics of the fermion field. The mass term $m\psi$ describes the mass of the fermion.

According to Dirac theory, particles exhibit chirality, enabling them to be divided into left-handed and right-handed components. Chirality is related to the asymmetry between left-handed and right-handed particles, where the handedness of a particle is determined by its behaviour under Lorentz transformations, with left-handed particles and right-handed antiparticles behaving differently from their right-handed and left-handed counterparts respectively. A Lorentz transformation is a mathematical operation that describes how the coordinates of an event change when viewed from different inertial frames of reference.

In its current form, this Lagrangian describes the dynamics of free Dirac fermions but does not account for interactions with the electromagnetic field. To incorporate gauge interactions, we need to consider how the theory changes when local gauge symmetries are imposed.

2.2.3 QED

Quantum Electrodynamics (QED) is a gauge theory based on the $U(1)$ group, a fundamental concept in quantum field theory. This group is characterised by a single parameter which governs the transformations of the quantum fields in the theory; the phase angle, which represents a rotation in the complex plane of a quantum field. The QED reliance on the $U(1)$ symmetry implies that the laws of physics remain unchanged under transformations that are governed by a single continuous phase angle, demonstrating invariance under certain operations.

However, the initial Lagrangian for the free Dirac field is not invariant under local $U(1)$ transformations:

$$\mathcal{L} \rightarrow \mathcal{L}' = \psi e^{ie\alpha(x)}, \quad (2.2)$$

where e is the electric charge value and $\alpha(x)$ is a phase that freely varies. A local transformation is a transformation that you can apply at any point in your field. For example, you can apply a gauge transformation to the field at one location and another different transformation at another location in space-time. To explain this mathematically, a global transformation would be described by $\psi \rightarrow \psi e^{ie\alpha}$, where α is a constant. On the other hand, a local transformation is described by $\psi \rightarrow \psi e^{ie\alpha(x)}$, where $\alpha(x)$ is a function of space-time.

The local transformation $\psi \rightarrow \psi e^{ie\alpha(x)}$ introduces additional terms that involve derivatives of phase $\alpha(x)$. Hence, we see that the local or gauge transformation of the fermion field does not leave the Lagrangian invariant. Evidently, the Lagrangian in its current form will not do.

To restore the invariance to the Lagrangian, we need to introduce a new field that transforms in a specific way under local gauge transformations in order to cancel out the additional terms introduced by the gauge transformation of the fermion field. We introduce a vector field $A_\mu(x)$, representing the electromagnetic field, that transforms as $A_\mu \rightarrow A_\mu - \delta^\mu \alpha$. This is referred to as a gauge field. This gauge field A_μ mediates the electromagnetic interaction by coupling to the fermion field ψ , and all charged particles. To do this, we replace the ordinary derivative in the kinetic term with a “covariant derivative” that includes the gauge field; $D_\mu = \partial_\mu + ieA_\mu$.

We obtain a new gauge invariant Lagrangian:

$$\mathcal{L} = \bar{\psi}(i\gamma^\mu D_\mu - m)\psi, \quad (2.3)$$

Expanding this Lagrangian reveals a term that is proportional to $e\bar{\psi}\gamma^\mu\psi A_\mu$. The presence of this term indicates the strength of its connection, denoted by the coupling constant e . This term signifies an interaction between the Dirac field ψ and the gauge field A^μ . A^μ is also known as the photon field, and the photon, as the gauge boson, mediates the electromagnetic force between charged particles. This means that the electron interacts with the photon field and the photon field interacts with the electron.

Additionally, the full QED Lagrangian includes the electromagnetic field tensor, $F^{\mu\nu} = \delta_\mu^\nu A_\nu - \delta_\nu^\mu A_\mu$, which captures the strength and direction of the electromagnetic field at each point in space–time. We construct this Lagrangian by combining the kinetic term for the fermion field, the mass term, the interaction term between the fermion field and the gauge field, and the kinetic term for the gauge field,

$$\mathcal{L}_{QED} = i\bar{\psi}\gamma_\mu\delta^\mu\psi - m\bar{\psi}\psi + e\bar{\psi}\gamma_\mu\psi A^\mu - \frac{1}{4}F^{\mu\nu}F_{\mu\nu}, \quad (2.4)$$

where $\frac{1}{4}F^{\mu\nu}F_{\mu\nu}$ is the gauge field kinetic term.

The gauge symmetry of QED is crucial to the theory’s consistency and invariance under local transformations. The electromagnetic force is a direct consequence of this symmetry, with the photon mediating interactions between charged particles. The other fundamental forces all originate in the same manner, yet with a somewhat more intricate gauge transformation.

2.2.4 Electroweak

In processes such as muon decay and nuclear β decay, the weak force is essential, causing the particles to change into other particles due to weak interactions. For example, muon decay involves a muon transforming into an electron, a muon neutrino, and an electron antineutrino. During β^- decay, a neutron within an atomic nucleus is converted into a proton, an electron, and an electron antineutrino. This conversion plays a vital role in processes such as element formation in stars and maintaining the stability of atomic nuclei.

Similarly to how the electromagnetic force is associated with the conservation of the electric charge, the weak force is related to the conservation of weak isospin. The $SU(2)$ symmetry group deals with the weak isospin, which is analogous to the concept of spin in quantum mechanics; i.e., it is an internal degree of freedom of particles that determines their behaviour under the weak force. On the other hand, $U(1)$ symmetry is associated with the weak hyper–charge, which is a quantum number related to the electric charge and the third component of the weak isospin. Both weak isospin and weak hypercharge serve as the gauge charges in the electroweak interaction, analogous to electric charge in electromagnetism.

In the mid–20th century, efforts to comprehend the fundamental forces of nature resulted in many advances as well as many obstacles. Electromagnetism, which governs long–range interactions and is mediated by photons, was well–understood through Quantum Electrodynamics (QED). However, the weak interaction, responsible for processes such as β decay, presented a conundrum. Unlike electromagnetism, the weak interaction was experimentally observed to be short–ranged and exhibited a preference for coupling with left–handed particles and right–handed antiparticles. This difference in behaviour between weak and electromagnetic forces is known as “parity violation”, which is a fundamental aspect of the weak interaction. Parity violation refers to the breaking of parity symmetry, where the laws of physics do not remain invariant under mirror reflections. Parity violation was first observed in the 1950s in experiments involving β decay, where the emitted electrons were found to be predominantly left–handed. This was contrary to the expectation that all fundamental interactions should exhibit parity conservation.

The evidence of chirality in weak interactions suggested that the weak force might be described by a theory involving chiral symmetries. This led to the hypothesis that weak interactions could be unified with other forces through gauge symmetries. As such, this discovery was a significant milestone in the understanding of the weak force and led to the development of the electroweak theory, which was developed by Sheldon Glashow, Abdus Salam, and Steven Weinberg in the late 1960s.

In the framework of electroweak theory, the $SU(2)$ symmetry is more accurately referred to as $SU(2)_L$, with the subscript “ L ” indicating the left-handedness of the particles involved. This distinction is essential to comprehend the behaviour of particles under the weak force and how electroweak theory incorporates this asymmetry. Similarly, the $U(1)_Y$ symmetry describes the gauge group associated with the weak hypercharge, Y .

The electroweak theory proposed that the electromagnetic and weak interactions are different aspects of a single fundamental force, unified under a gauge theory framework. This theory is based on the gauge group $SU(2)_L \times U(1)_Y$ and describes the behaviour of particles under the weak force, which involves fermions forming doublets of isospin. As mentioned earlier, for example, quarks and leptons are organised into these doublets. The weak force accounts for the transformations between these different flavours of particles, such as the transformation of one type of quark into another type. For instance, in a β decay, a neutron transforms into a proton, where a down-quark transforms into an up-quark through the exchange of a W boson. This demonstrates the role of the weak force in particle transformations. The three W bosons, W^1 , W^2 , and W^3 , are the force carriers associated with the $SU(2)_L$ group, and mediate interactions that change the weak isospin of the particles.

The $U(1)_Y$ gauge group yields one gauge boson, denoted B^0 , which mediates the interaction associated with weak hypercharge. As an Abelian group, $U(1)_Y$ does not change the weak hypercharge of particles, and the B^0 gauge boson itself is uncharged. Similarly, the $SU(2)_L$ gauge group contributes three gauge bosons, denoted W^1 , W^2 , and W^3 , which are associated with weak isospin. These four gauge bosons (B^0 and the W^i) form the generators of the combined gauge symmetry.

The unification of electromagnetic and weak interactions in electroweak theory does not arise merely from the coexistence of $U(1)_Y$ and $SU(2)_L$. Rather, it is the process of spontaneous symmetry breaking, via the Higgs mechanism, that mixes the B^0 and W^3 gauge bosons. This will be discussed in more detail in Section 2.2.5. The resulting mixing leads to two new physical bosons: the photon (γ), which corresponds to the unbroken $U(1)$ electromagnetic symmetry, and the Z^0 boson, which is massive and mediates the neutral weak interaction. This symmetry breaking and the mixing of the gauge bosons are what unify –or more accurately, mix – the electromagnetic and weak interactions.

The weak-mixing angle, θ_W , also known as the Weinberg angle, is a fundamental parameter in the electroweak theory that quantifies the mixing between the gauge bosons of the weak and electromagnetic interactions. It emerges from the unification of the $SU(2)_L$ and $U(1)_Y$ gauge groups into the electroweak theory.

It describes how the gauge bosons of the weak force mix to form the physical W and Z bosons. Specifically, the physical $W^\pm$ bosons are combinations of W^1 and W^2 , while the Z boson and the photon (γ) arise from a mixing of W^3 and B^0 . This mixing is described by:

$$W^\pm = \frac{1}{\sqrt{2}} (W^1 \mp iW^2) \quad (2.5)$$

$$Z = \cos \theta_W W^3 - \sin \theta_W B^0 \quad (2.6)$$

$$\gamma = \sin \theta_W W^3 + \cos \theta_W B^0 \quad (2.7)$$

This relationship highlights how the electromagnetic and weak forces are interconnected within the electroweak theory.

2.2.5 Electroweak Symmetry Breaking and the Higgs mechanism

The connection between the masslessness of bosons in the original electroweak theory and their acquisition of mass through the Higgs mechanism can be understood through spontaneous symmetry breaking of the $SU(2)_L \times U(1)_Y$ symmetry. Before symmetry breaking, the $W^\pm$ and Z bosons, which mediate the weak force, were considered massless. However, with spontaneous symmetry breaking, the Higgs field's vacuum expectation value introduces a mechanism by which the $W^\pm$ and Z bosons acquire mass, while the photon remains massless, representing the electromagnetic force.

The electroweak theory combined the electromagnetic and weak–nuclear forces, demonstrating their underlying symmetry. At high energies, electromagnetic and weak forces are indistinguishable, regulated by a unified gauge symmetry represented by the $SU(2)_L \times U(1)_Y$ group. However, at lower energies, the experimental data reveal that the interplay between the two forces is diluted, suggesting the presence of electroweak symmetry breaking.

In the 1960s, the Higgs mechanism was proposed as a solution to the problem of how the weak bosons acquire mass while still adhering to the principle of gauge invariance. This mechanism postulates the existence of a scalar field, the Higgs field, which interacts with the weak bosons and provides them with mass. A scalar field is a field that has a single value at each point in space, as opposed to a vector field which has both magnitude and direction, like an electromagnetic field.

We begin with a gauge theory involving a complex doublet, denoted by $\Phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix}$, where ϕ^+ and ϕ^0 are complex scalar fields. The Lagrangian density is given by:

$$\mathcal{L}_H = (D_\mu \Phi)^\dagger (D^\mu \Phi) - V(\Phi), \quad (2.8)$$

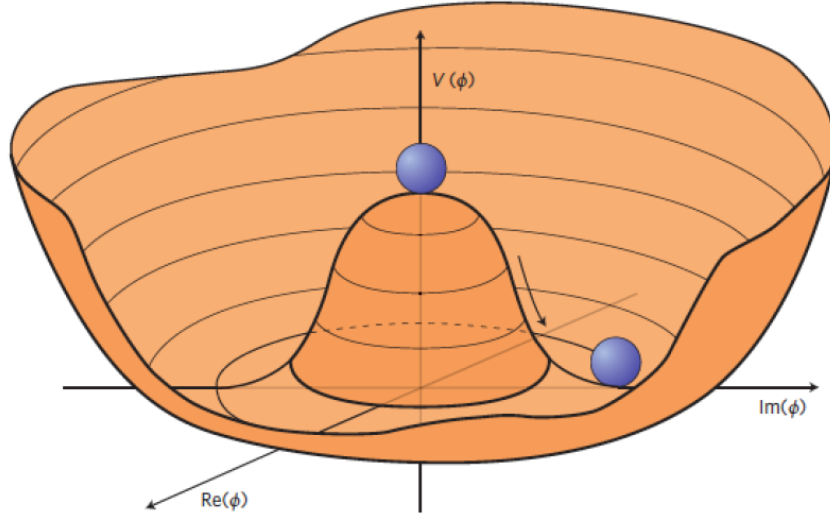


FIGURE 2.4: An illustration of the Higgs potential (5) in the case that $\mu^2 < 0$, in which case the minimum is at $|\Phi|^2 = -\mu^2/2\lambda$. Choosing any of the points at the bottom of the potential breaks spontaneously the rotational $U(1)$ symmetry. Reproduced from Ref. [28]

where Φ signifies the complex Higgs field as a $SU(2)$ doublet, and the covariant derivative, D_μ , is given by:

$$D_\mu = \partial_\mu - ig \frac{\sigma^a}{2} W_\mu^a - ig' \frac{Y}{2} B_\mu, \quad (2.9)$$

where g and g' are the coupling constants associated with the gauge fields $SU(2)$ and $U(1)$, respectively, W_μ^a and B_μ represent the gauge fields $SU(2)$ and $U(1)$, respectively, σ^a are the three Pauli matrices (with $a = 1, 2, 3$), and Y is the weak hypercharge.

The $V(\Phi)$ in equation 2.8 is the Higgs potential, given by:

$$V(\Phi) = \mu^2 \Phi^\dagger \Phi + \lambda (\Phi^\dagger \Phi)^2. \quad (2.10)$$

This potential contains two terms, one involving the mass term, characterised by μ^2 , and the other representing the Higgs self-interaction, characterised by the quartic coupling constant, λ .

The field Φ has a vacuum state, which is its lowest energy state. This state is determined by the minimum potential $V(\Phi)$. If μ^2 is positive, the minimum occurs at $\Phi = 0$. However, if μ^2 is negative, the minimum of the potential occurs for $\Phi^\dagger \Phi = \frac{-\mu^2}{2\lambda} \equiv \frac{v^2}{2}$. This means that the field does not vanish even when there are no particles present, and it acquires a nonzero vacuum expectation value of v . This potential also has the distinct shape of a “Mexican hat”, where the minimum lies along a circle in field space.

The symmetry of the potential allows for two distinct vacuum states. This selection of the vacuum state is called “spontaneous symmetry breaking”. When the Lagrangian of the theory is expressed by expanding the Higgs field around the vacuum state, $\Phi(x) = v + h(x)$, the gauge fields that were previously massless acquire mass terms, thus allowing the gauge bosons to have mass:

$$m_W = \frac{gv}{2} \text{ and } m_Z = \frac{v\sqrt{g^2 + g'^2}}{2}. \quad (2.11)$$

Nevertheless, the photon remains massless. This is central to the theory and can be understood through the potential’s shape, often illustrated as a Mexican hat. One of the scalar fields, out of the Higgs doublet, corresponds to a massless degree of freedom, which can be visualised as rolling around the rim of the Mexican hat. This massless degree of freedom corresponds to the photon. When a non-zero ground state is introduced, it results in the spontaneous breaking of gauge symmetry. Consequently, the original gauge symmetry of the Lagrangian is no longer preserved, leading to the photon remaining massless while other gauge bosons acquire mass. However, the theory as a whole, prior to the introduction of the ground state, is still gauge-invariant.

The Higgs mechanism can also be used to give mass to fermions, by coupling to them with coupling constants g_i , referred to as the Yukawa coupling for fermion i . This coupling is responsible for the different masses of the various fermions.

Quantum fluctuations in the Higgs field are responsible for the emergence of the Higgs boson, which was observed as a short-lived resonance in particle collisions. This is similar to how vibrations in a field, such as those in the electromagnetic field, can appear as particles like photons. The Higgs boson was discovered at CERN’s Large Hadron Collider in 2012, and is often referred to as the “final piece” of the Standard Model of particle physics.

2.2.6 QCD

The $U(1)$ group we discussed earlier is just one example from a larger collection known as the $U(n)$ groups. These $U(n)$ groups are responsible for representing unitary matrices that have n dimensions and maintain the length of vectors when operating on them. Now, our attention turns to the $SU(n)$ group, which stands for “special unitary” and comprises the subset of $SU(n)$ matrices that satisfy the special condition of having a determinant 1. In the case of these $SU(n)$ matrices, we encounter a total of $n^2 - 1$ independent parameters that define the transformations they represent. This is because the matrices are unitary with determinant 1, which imposes constraints on the number of independent parameters. As a result, the $SU(n)$ group has $n^2 - 1$ generators, which generate all possible unitary transformations with determinant 1 of n -component complex vectors.

In Quantum Chromodynamics (QCD), the $SU(3)$ symmetry is used, which has 8 generators. These generators signify the colour charge degrees of freedom that quarks possess. Unlike the $U(1)$ symmetry of Quantum Electrodynamics, the $SU(3)$ symmetry of Quantum Chromodynamics allows for the exchange of three distinct colour charges: red, green, and blue. The colour

charges are fundamental to strong interactions, as they determine how quarks interact with each other via the exchange of gluons, the force carriers of QCD. The combinations of these colour charges are critical for forming colour-neutral (or "white") particles, such as protons and neutrons, which are made up of three quarks, each carrying a different colour charge.

The gluons, the mediators of the strong force, themselves carry colour charge and thus are self-interacting, leading to the unique phenomenon of gluon self-coupling. The QCD colour charge is analogous to the QED electric charge in that it is a conserved quantity in strong interactions. However, unlike the electric charge, which is carried by particles such as electrons and protons, gluons, the force carriers of QCD, are unique in that they carry both a colour charge and an anti-colour charge, which allows them to interact with each other in addition to interacting with quarks.

The strong force is characterised by the coupling parameter α_s , which is a fundamental element of QCD. This parameter depends on the energy scale of the interaction and is a key factor in determining the behaviour of the strong force between quarks and gluons. Its value is not constant but changes with the momentum transfer, making it a "running coupling".

2.2.7 Perturbation theory and higher-order corrections

Consider a scenario where two electrons interact by exchanging a photon, as depicted in Figure 2.5. In this illustration, the paths of the electrons are shown as straight lines, and the photon that is exchanged is indicated by a wavy line. These diagrams, called Feynman diagrams, serve as a graphical depiction of particle interactions within quantum field theory, providing an accessible way to examine processes like Bhabha scattering (two electrons scattering). Feynman diagrams effectively demonstrate how particles use force carriers, such as the photon in this instance, to alter each other's paths [29].

This specific diagram showing the simplest interaction between the two electrons represents a "leading-order" calculation. In leading-order calculations, we consider only the most basic diagrams, which provide a fundamental understanding of the interaction. However, to achieve greater accuracy in theoretical predictions, we must also consider higher-order corrections. We have an infinite number of these higher-order diagrams that you could theoretically add up to get the correct result. These corrections involve additional diagrams that account for more complex interactions, such as multiple photon exchanges or virtual particle loops. An example of such "escalation" is shown in Figure 2.6.

Since we cannot add up an infinite amount of contributions from higher-order diagrams, these higher-order corrections are typically performed via "perturbative" calculations, which involve expanding the cross-section in a series of terms, with each term representing a different order of the interaction. It is similar to how we expand a function in a Taylor series, where each term represents a different order of the expansion.

Perturbation theory is a fundamental technique in quantum field theory for approximating the behaviour of systems with weak interactions. The interacting theory is expressed as a series expansion around a solvable free theory, with the interaction term treated as a small correction.

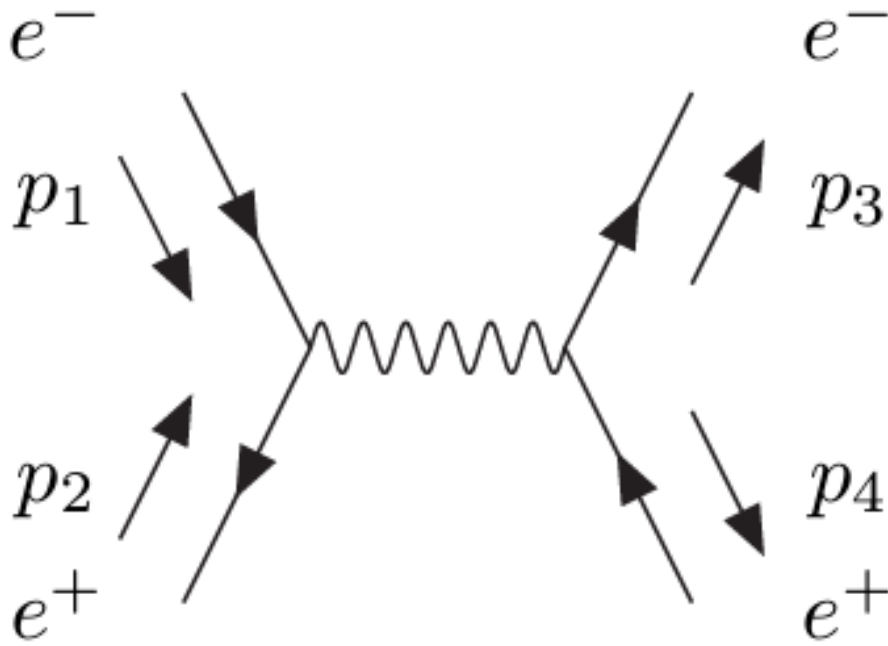


FIGURE 2.5: A Feynman diagram depicting Bhabba scattering – where two electrons exchange a photon. This is a leading-order diagram. Reproduced from [30].

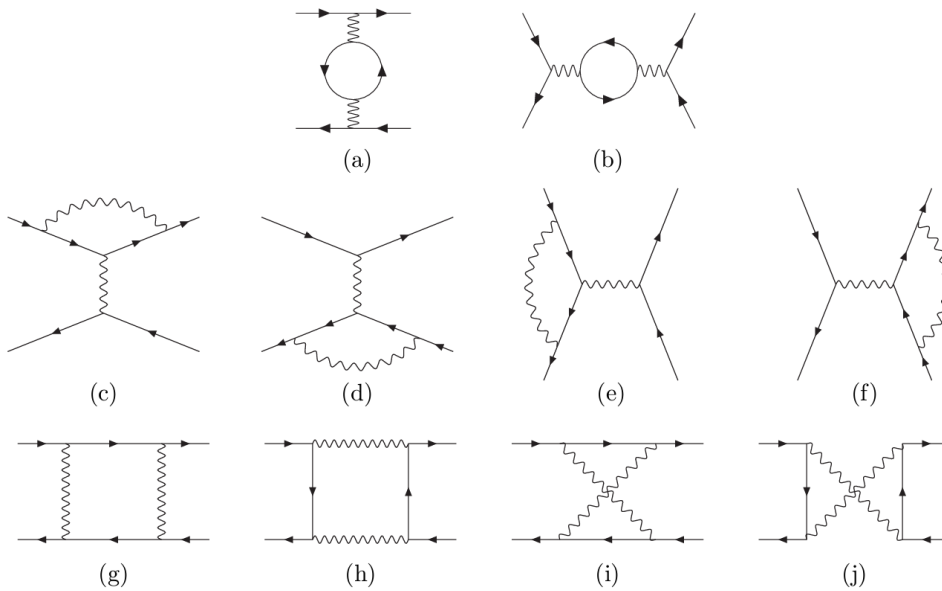


FIGURE 2.6: A few Feynman diagrams depicting Bhabba scattering – where two electrons exchange a photon. These are higher-order one-loop diagrams. Reproduced from [30].

Physical quantities, such as scattering amplitudes or cross-sections, are computed as a power series in terms of a coupling constant. This gives a fairly good approximation for the final result.

For example, in QED, the coupling constant is the fine-structure constant $\alpha \approx 1/137$, which governs the strength of electromagnetic interactions. The cross-section σ for a particular process can be written as an expansion in powers of α :

$$\sigma = \sigma_0 + \alpha\sigma_1 + \alpha^2\sigma_2 + \dots \quad (2.12)$$

where σ_0 represents the tree-level contribution (leading order), and the higher-order terms represent loop corrections. It is evident that each successive term becomes less significant compared to the previous one as they diminish by the factor α , approximately $1/137$. Consequently, we can disregard most of the higher-order terms, while still obtaining a reasonably accurate approximation.

In QCD, the coupling constant is α_s , which depends on the energy scale. At high energies, α_s becomes small, and perturbative techniques can be applied. For example, the perturbative expansion for a quantity R in QCD can be written as:

$$R = R_0 + \alpha_s R_1 + \alpha_s^2 R_2 + \dots \quad (2.13)$$

where R_0 represents the leading-order (LO) term, and the subsequent terms correspond to “next-to-leading-order” (NLO) and “next-to-next-to-leading-order” (NNLO) corrections[29]. These higher-order corrections are corrections to tree-level processes that involve the emission of additional particles, or loops, which are virtual particles that appear in the intermediate states of a process. Additionally, the running of α_s is a consequence of these higher-order corrections; it accounts for the energy dependence of the strong coupling constant and reflects how α_s varies with the energy scale of the interaction.

Thus, perturbative expansions in both QED and QCD depend on the size of the coupling constants α and α_s . At high energies, the series converges well, with each successive term contributing less than the previous [29].

In processes where both QED and QCD interactions are involved, interference terms between the two forces can arise, leading to mixed terms in the perturbative expansion. These terms are proportional to the products of the respective coupling constants, such as $\alpha\alpha_s$. For example, in processes involving quarks, which are both electrically charged (subject to QED) and colour-charged (subject to QCD), one might encounter contributions to the cross-section or amplitude of the form:

$$\sigma = \sigma_0 + \alpha\sigma_1^{\text{QED}} + \alpha_s\sigma_1^{\text{QCD}} + \alpha\alpha_s\sigma_1^{\text{QED-QCD}} + \dots \quad (2.14)$$

The term $\alpha\alpha_s\sigma_1^{\text{QED-QCD}}$ represents an interference effect between the electromagnetic and strong interactions. These mixed terms become important in hadronic processes involving photon emissions or when both electromagnetic and strong forces contribute to the same scattering event. The relative size of these interference terms depends on the energy scale and the magnitude of α and α_s , with α_s generally dominating at low energies (due to the strength of QCD), while α remains small but significant across a broad range of scales [29].

Furthermore, the coupling “constant” α_s is determined by the renormalisation group equation and varies with the reference energy scale μ , as shown by the following expression:

$$\alpha_s(\mu^2) = \frac{12\pi}{(33 - 2n_f) \ln(\mu^2/\Lambda_{\text{QCD}}^2)} \quad (2.15)$$

where n_f is the number of active quark flavours, and Λ_{QCD} is the QCD scale parameter. For this analysis, the reference energy scale is chosen to be the Z pole energy scale, which is approximately 91.2 GeV. At this energy, the electroweak interactions are maximised, and it serves as a crucial point for precision measurements in particle physics.

The running of α_s can also be described by Equation 2.16 at a certain energy, Q , that is different from the reference energy, μ . The dynamics of α_s can be expressed by the equation:

$$\alpha_s(Q^2) = \frac{\alpha_s(\mu^2)}{1 + \frac{b_0\alpha_s(\mu^2)}{2\pi} \ln\left(\frac{Q^2}{\mu^2}\right)}, \quad (2.16)$$

where b_0 is defined as:

$$b_0 = \frac{33 - 2n_f}{12\pi}. \quad (2.17)$$

As the energy scale Q increases, $\alpha_s(Q^2)$ decreases, making perturbation theory more accurate at high energies, while at low energies α_s grows large, signaling the breakdown of perturbative methods.

This equation reveals the intriguing behaviour of α_s , giving rise to the concept of “asymptotic freedom”, implying that at high energies, the value of α_s becomes exceedingly small, so quarks and gluons behave as nearly free particles. This phenomenon has enabled successful perturbative calculations in QCD in high-energy regimes, which ultimately resulted in the awarding of a Nobel Prize in 2004. [31]. However, at lower energies, this force increases in strength, leading to the intriguing phenomenon of confinement, where quarks and gluons are confined within colour-neutral bound states, such as protons and neutrons.

Learning about the behaviour of α_s is essential to understand the dynamics of QCD and the complex character of the strong-nuclear force. This QCD knowledge can provide a thorough understanding of nuclear matter and the behaviour of particles under extreme conditions, such as those found in high-energy collisions and the early universe. Dealing with QCD can be a

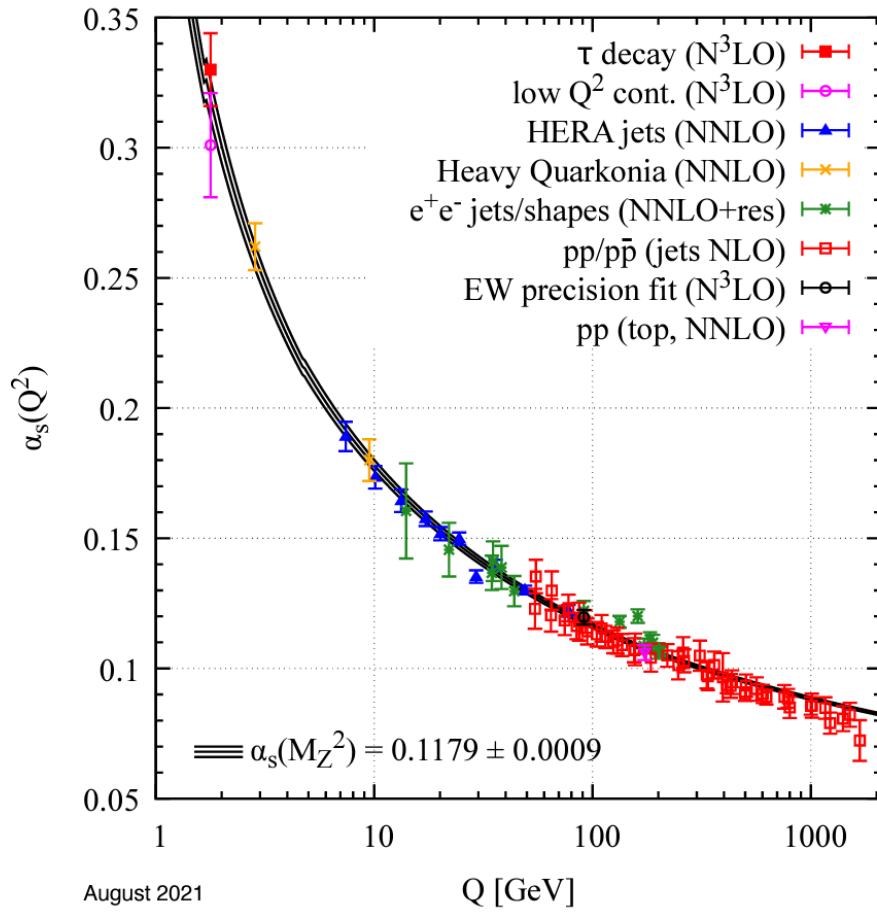


FIGURE 2.7: α_s measurements as a function of the energy scale, Q , at various levels of perturbation theory. Reproduced from Ref. [32]

difficult endeavour. Its non-perturbative effects at low energies make exact calculations difficult, requiring the use of sophisticated computational techniques, such as large-scale simulations. One other method that we will encounter is the resummation of logarithms, which is a technique used to improve the accuracy of perturbative calculations by summing up the leading logarithmic terms that arise in the perturbative expansion of the cross-section. This is particularly useful in the context of QCD, where the running of α_s leads to the appearance of large logarithms that can be resummed to improve the precision of theoretical predictions. For example, the next-to-leading-order (NLO) correction to the differential cross-section for a generic deep inelastic scattering can be expressed as:

$$\frac{d\sigma}{d\ln Q^2} \approx \frac{\alpha_s(Q^2)}{2\pi} \left[C_1 \ln \left(\frac{Q^2}{\mu^2} \right) + C_2 \right], \quad (2.18)$$

where $\frac{d\sigma}{d\ln Q^2}$ is the differential cross-section with respect to the logarithm of the energy scale Q^2 , $\alpha_s(Q^2)$ is the strong coupling constant at the scale Q^2 , and C_1 and C_2 are coefficients that depend on the specific process and order of perturbation. This simplified equation illustrates the appearance of significant logarithmic terms in perturbative QCD (pQCD) calculations, which are often resummed to enhance precision.

The intricate nature of these calculations, along with the numerous variables involved, requires the use of high-performance computing resources. To comprehend QCD in the strong coupling regime and the complexities of the strong force, sophisticated theoretical approaches and computational tools are needed, making the study of QCD a difficult but essential pursuit in modern physics. A challenging computational method is “lattice QCD”, a technique that segments space-time into a lattice to explore QCD in the non-perturbative regime [33]. Non-perturbative regimes are areas where, as shown in Figure 2.7, the coupling strength of the strong interaction becomes so high that traditional perturbative methods fail. Lattice QCD is used, for example, by the BMW collaboration, who perform high-precision calculations of hadron masses, decay constants, and other key observables [34].

Another form of higher-order correction to account for is QED final-state radiation (FSR). This concerns the final states in our process of interest, which involve muons. Because muons are electrically charged, they can emit photons according to quantum electrodynamics. In a simplified model, these processes are predicted without considering such radiation and are referred to as “Born”-level processes. These represent the most basic Feynman diagrams that depict the leading-order (LO) approximation of the process. In reality, however, such photon emissions occur and must be included for precise measurements. This results in a “bare” final-state, where the muon is accompanied by radiated photons. The energy of the bare final-state muon will be lower than in the Born-level approximation because some energy is carried away by the emitted photon(s). The “dressed” state (muon + photon) represents what is typically observed with a detector. This is a result of higher-order correction. Monte Carlo generators often give bare leptons as final-state objects, independently of the level of the calculations. The dressing of the final-state particles is separable from the rest of the event, so that the QED FSR can be computed and modelled independently.

These examples of higher-order corrections are used to improve the precision of theoretical predictions that are essential for comparing with experimental data; especially when it comes to precision measurements at high-energy colliders. In Section 2.3.5, we discuss how we can calculate these higher-order corrections in practice for experimental analyses using specialised software tools.

2.3 Hadron collider phenomenology and the Drell–Yan process

2.3.1 Partons

Partons, a theoretical concept in particle physics, serve as a fundamental concept to clarify the intricate composition of hadrons, such as protons and neutrons. Hadrons are postulated to consist of “valence quarks”, the primary stable building blocks explicitly involved in the quark model representation of these particles. For instance, a proton is characterised by two up quarks and one down quark (uud), while a neutron is characterised by one up quark and two down quarks (udd).

The mass of the proton was found to be considerably greater than the aggregate mass of its three valence quarks. The masses of these valence quarks are known from precise measurements in high-energy collisions or from theoretical calculations within the Standard Model framework. However, this aggregate is insufficient to explain the entire proton mass. The discrepancy emerges due to the strong binding energy within the proton, driven by the surrounding quantum vacuum that constantly produces virtual quark–antiquark pairs, known as “sea quarks”, along with gluons. The sea quarks and gluons interact by exchanging virtual particles, forming a dynamic system that amplifies the proton’s observed mass. The distinction between valence and sea quarks is crucial in understanding phenomena like the spin structure of the proton and is supported by experimental evidence, particularly from deep–inelastic scattering experiments (e.g. $e^-p \rightarrow e^- X$) [35]. This duality underscores the intricate interplay of quark and gluon dynamics governed by the strong force within the framework of quantum chromodynamics. Partons are used to explain high-energy scattering processes, particularly in the context of QCD. This makes partons a very important concept in particle physics experiments that rely on hadrons, e.g. proton–proton colliders like the Large Hadron Collider (LHC) at CERN, because such experiments fundamentally depend on high-energy scattering of partons when the hadrons collide. They also help in the calculation of cross-sections and various other observables related to the strong interactions.

The notion of partons is based on the concept of “asymptotic freedom” in QCD, which was discussed in Subsection 2.2.6. This concept suggests that at high energies, quarks and gluons act like nearly-free particles, exhibiting characteristics similar to those of non-interacting particles. This approximation enables physicists to treat the components of hadrons as quasi-free entities during high-energy interactions.

Isolated quarks cannot be seen in experiments because of confinement, a phenomenon that stops individual quarks and gluons from being seen on their own. Instead, experiments detect hadrons

that are created when these quarks and gluons undergo hadronisation; the process by which these particles combine to form composite states like protons, neutrons, and other hadrons.

2.3.2 Particle distribution functions

Understanding hadron interactions at high energies requires insight into the internal structure of hadrons, which are complex assemblies of quarks and gluons. To compute and predict reactions in high-energy particle physics experiments, we need to know the probability of finding a particular parton—quark or gluon—carrying a specific fraction of the total momentum of the hadron. This probability is quantified by parton distribution functions (PDFs).

In particular, Bjorken- x is a variable that represents the fraction of the hadron’s momentum carried by a parton. To compute observables such as the cross-section for hadron–hadron collisions, we use the following expression:

$$d\sigma = \sum_{a,b} \int dx_1 dx_2 f_{a/H_1}(x_1, \mu_F) f_{b/H_2}(x_2, \mu_F) d\hat{\sigma}_{ab \rightarrow X}(\sqrt{s}, \alpha_s(\mu_R), \mu_R, \mu_F), \quad (2.19)$$

where the sum $\sum_{a,b}$ runs over all possible parton flavours, including quarks and gluons, within the colliding hadrons. The functions $f_{a/H_1}(x_1, \mu_F)$ and $f_{b/H_2}(x_2, \mu_F)$ represent the parton distribution functions for partons a and b in hadrons H_1 and H_2 at the factorisation scale μ_F , respectively, while μ_R is the renormalisation scale. The factorisation scale is a parameter used to differentiate between long-distance and short-distance physics in the context of factorisation theorems. It is used to distinguish between the non-perturbative effects, which are represented by the PDFs, and the perturbative effects that can be calculated using pQCD. The choice of the factorisation scale can influence the convergence and accuracy of perturbative calculations. However, there is no definite way of setting the factorisation scale, and it usually involves making informed decisions based on the particular process being studied. The renormalisation scale represents the scale at which the coupling constant and the fields are renormalised, effectively absorbing the divergences arising from loop corrections in the theory.

These PDFs describe the probability densities of finding the respective partons with momentum fractions x_1 and x_2 in the hadrons. The term $d\hat{\sigma}_{ab \rightarrow X}(\sqrt{s}, \alpha_s(\mu_R), \mu_R, \mu_F)$ refers to the partonic cross-section, which characterises the interaction between partons a and b resulting in the final state X . This cross-section depends on the partonic centre-of-mass energy $\sqrt{s}$, the strong coupling constant α_s at the renormalisation scale μ_R , and the factorisation and renormalisation scales μ_F and μ_R , respectively.

Since PDFs originate from low-energy interactions inside the proton, they cannot be determined using pQCD from first principles, as the value of α_s on this scale is too high. Therefore, PDFs cannot be calculated directly from the fundamental theory of QCD at these low scales, and their determination often involves a combination of experimental measurements and theoretical models. Various techniques, such as global fits to the experimental data and lattice QCD calculations, are employed to extract information about PDFs from the experimental data. PDFs

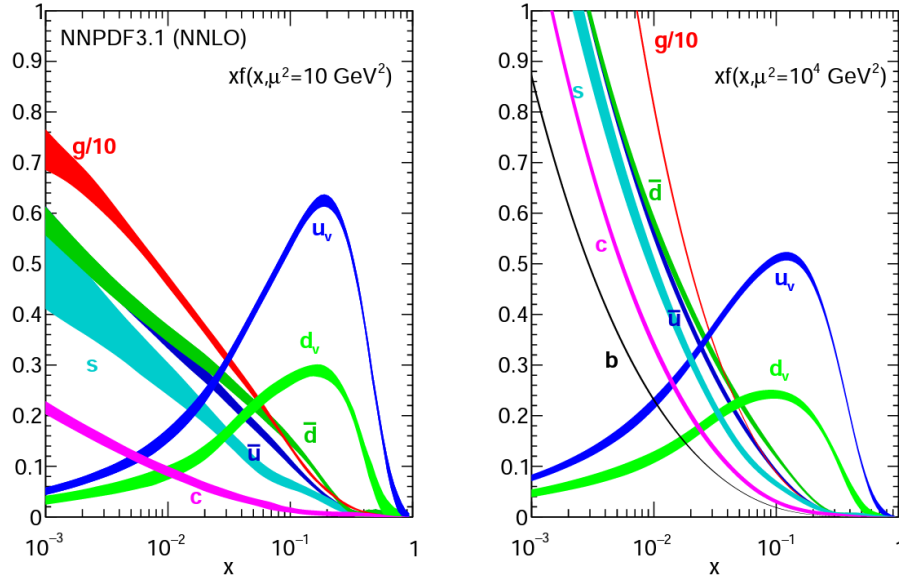


FIGURE 2.8: Parton distribution functions from global fits by the NNPDF Collaboration [36], evaluated at scales of (left) $Q^2 = 10 \text{ GeV}^2$ and (right) $Q^2 = 10^4 \text{ GeV}^2$.

evolve with the factorisation scale μ_F according to the Dokshitzer–Gribov–Lipatov–Altarelli–Parisi (DGLAP) equations. These equations describe how parton densities change as a function of the energy scale, incorporating the effects of parton evolution and interactions. The general form of the DGLAP equation for a parton distribution function $f_a(x, \mu_F)$ is given by:

$$\frac{\partial f_a(x, \mu_F)}{\partial \ln \mu_F^2} = \frac{\alpha_s(\mu_R)}{2\pi} \int_x^1 \frac{dz}{z} P_{ab}(z) f_b\left(\frac{x}{z}, \mu_F\right), \quad (2.20)$$

where:

- $\alpha_s(\mu_R)$ is the strong coupling constant evaluated at the renormalisation scale μ_R .
- $P_{ab}(z)$ is the splitting function that describes the probability of a parton of type b splitting into a parton of type a with a fraction z of the momentum of the parent parton.
- The integral runs from x to 1, reflecting contributions from all possible values of z .

The NNPDF group [36] has extracted the PDFs of the proton at two different energy scales, which are shown in Figure 2.8 as a function of Bjorken- x . The PDFs of the valence quarks (u and d in the proton) are dominant at most x values; however, as the energy scale increases, the PDF of the gluon becomes more prominent, making the LHC essentially a gluon collider. Thus, PDFs are crucial for understanding hadronic interactions and for making predictions about the outcomes of high-energy particle collisions.

The ability to reach varying energy levels for the processes and the momentum transfer is largely influenced by the specific physics experiment. The Large Hadron Collider and its experiments

will be discussed in the following chapter, but for now it is important to note that the LHCb experiment [37] is located in the forward region ($2 < \eta < 5$), while ATLAS [38] and CMS [39] cover the central angles ($\eta < |2.5|$) around the pp collision point.

Collisions with decay products that are propelled forward in LHCb are usually highly asymmetric, with one low x and one high x parton, while ATLAS and CMS have more balanced collisions on average, and thus the different detectors probe different areas in the PDF kinematic space, as demonstrated in Figure 2.9 as a function of x and the squared momentum transfer Q^2 . Comparable PDF-sensitive measurements in LHCb will therefore (to some extent) be anti-correlated to those made in ATLAS or CMS— if the PDFs are greater at low and high x . The level of anti-correlation will depend on factors such as the degree of orthogonality and the acceptance range of η and the significance of the PDFs in the measurement. However, in general, the combination of anti-correlated measurements gives a smaller total uncertainty than the combination of two similar positively correlating measurements. These concepts inspired the proposal for a novel measurement of the W boson mass at LHCb, with a proof-of-principle measurement published in 2021 [21].

2.3.3 Drell–Yan

The Drell–Yan process is a particle physics phenomenon that takes place when a quark from one hadron and an antiquark from another hadron annihilate, creating a virtual photon or Z boson which decays into a pair of oppositely charged leptons, such as electrons or muons. Figure 2.10 depicts the Drell–Yan process. This process is typically described as $pp \rightarrow \gamma^*/Z \rightarrow \ell^+\ell^-$. However, it is also applicable for the W decay, which is what we’re studying in this analysis.

This process allows us to study the internal structure of protons. In terms of Quantum Chromodynamics (QCD), Drell–Yan is a “ $2 \rightarrow 0$ ” process, meaning that the initial state has two coloured partons, but the final-state particles are colourless. This makes it the simplest process involving QCD at hadron colliders, and it is consequently very well understood theoretically. The production of charged leptons results in a clean final-state, and the cross-section is also very high in comparison to other Standard Model processes, meaning that there are large data samples available for analysis. This combination of factors makes Drell–Yan the ideal process for rigorously testing the SM.

The Drell–Yan process is advantageous for making predictions because its differential cross-section can be factorised into angular and unpolarised components, which can be calculated separately. This is only applicable if electroweak corrections, such as Quantum Electrodynamics Final-State Radiation (QED FSR), are not taken into account. Photons produced in the final-state because of electromagnetic interactions can disrupt the simple factorisation scheme. We will use this in our favour to simulate the W boson differential cross-section.

The production and decay of vector bosons have five degrees of freedom. Initially, there are six (three components of momentum for each parton), but one is eliminated due to the cylindrical symmetry of the interaction around the beam axis. The remaining five are usually expressed in terms of the vector boson mass M , rapidity y and transverse momentum p_T^V , and two angles that

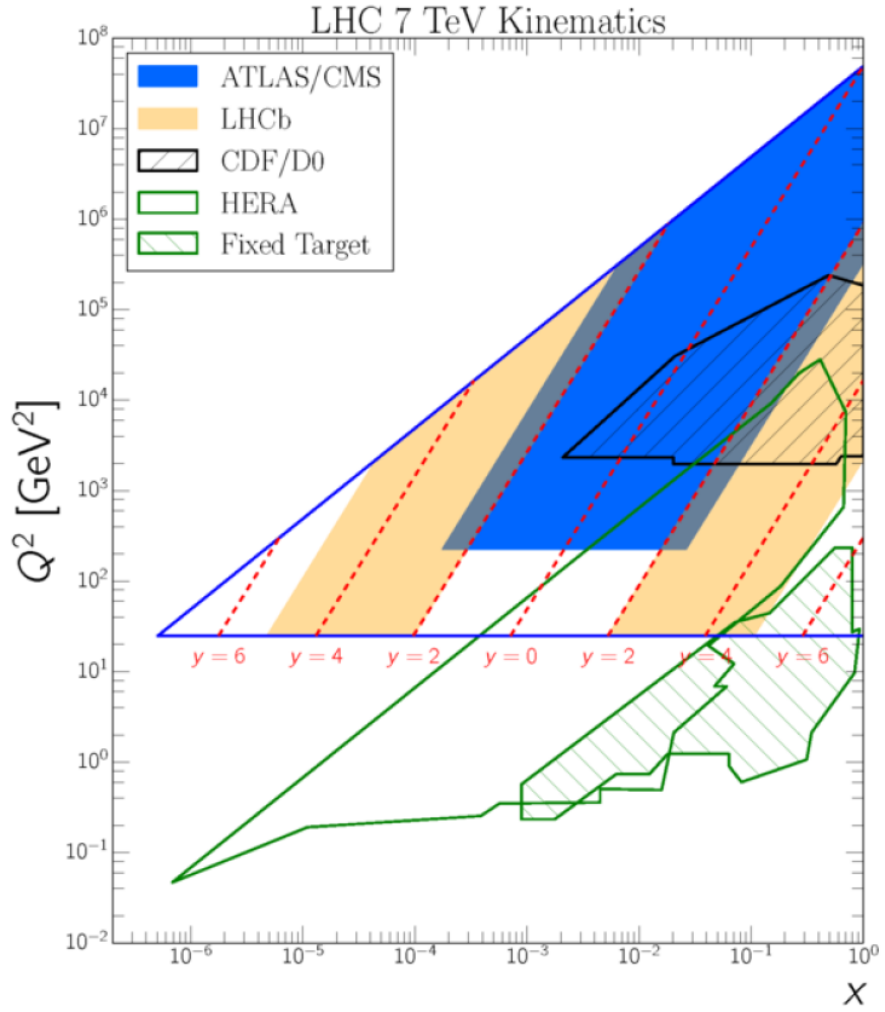


FIGURE 2.9: Coverage of high-energy colliders versus Bjorken x and momentum transfer Q^2 . Reproduced from Ref. [40]

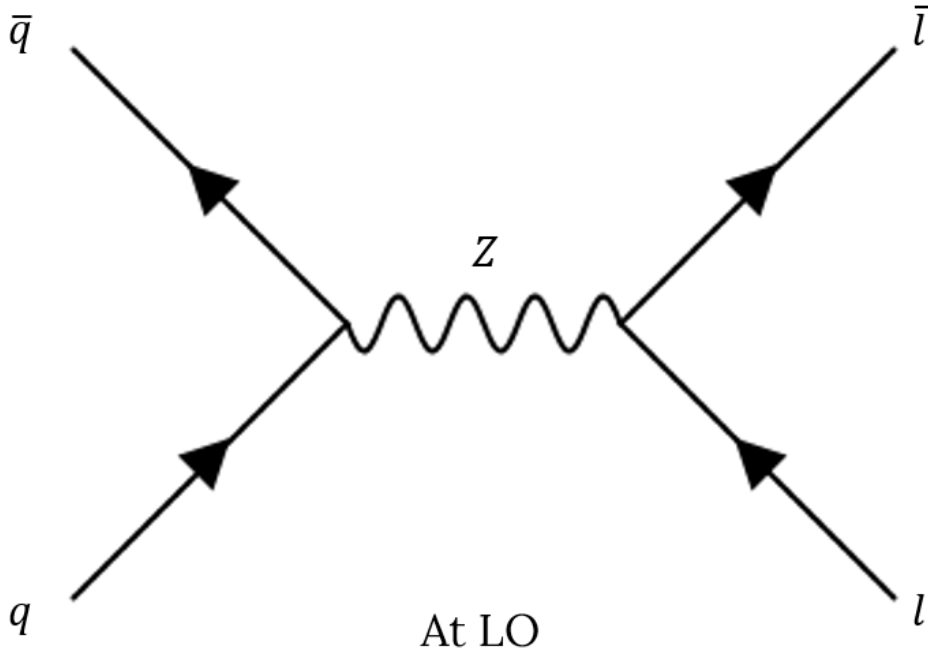


FIGURE 2.10: The Drell–Yan process, where a quark and an antiquark annihilate to produce a virtual photon or Z boson, which then decays into a pair of oppositely charged leptons. The process is shown at leading order in QCD (without any additional gluon radiation or other higher–order effects).

determine the direction of the final–state leptons. The angular distribution of the lepton system is impacted by the polarisation of the vector boson, which can be separated from the unpolarised part of the differential cross–section σ^{unpol} , resulting in the full 5D differential cross–section being

$$\begin{aligned} \frac{d\sigma}{dp_T^V dy dM d\cos\theta d\phi} &= \frac{3}{16\pi} \frac{d\sigma^{\text{unpol}}}{dp_T^V dy dM} \\ &\quad \left\{ (1 + \cos^2\theta) + A_0 \frac{1}{2} (1 - 3\cos^2\theta) + A_1 \sin 2\theta \cos\phi \right. \\ &\quad \left. + A_2 \frac{1}{2} \sin^2\theta \cos 2\phi + A_3 \sin\theta \cos\phi + A_4 \cos\theta \right. \\ &\quad \left. + A_5 \sin^2\theta \sin 2\phi + A_6 \sin 2\theta \sin\phi + A_7 \sin\theta \sin\phi \right\}, \end{aligned} \quad (2.21)$$

where the lepton decay angles θ and ϕ are defined in the Collins–Soper rest frame of the lepton pair [41] as shown in Figure 2.11. The angular distribution is expressed in terms of spherical harmonics, with each term having an angular coefficient A_i . These coefficients are ratios of helicity cross–sections and also depend on p_T^V , y and M . The angular cross–section comes from the representation in spherical harmonics with $l \leq 2$, where l is the quantum number of the orbital angular momentum. The restriction $l \leq 2$ arises from the spin–1 nature of vector bosons, such as the W and Z bosons, which decay into a pair of fermions. The total angular momentum

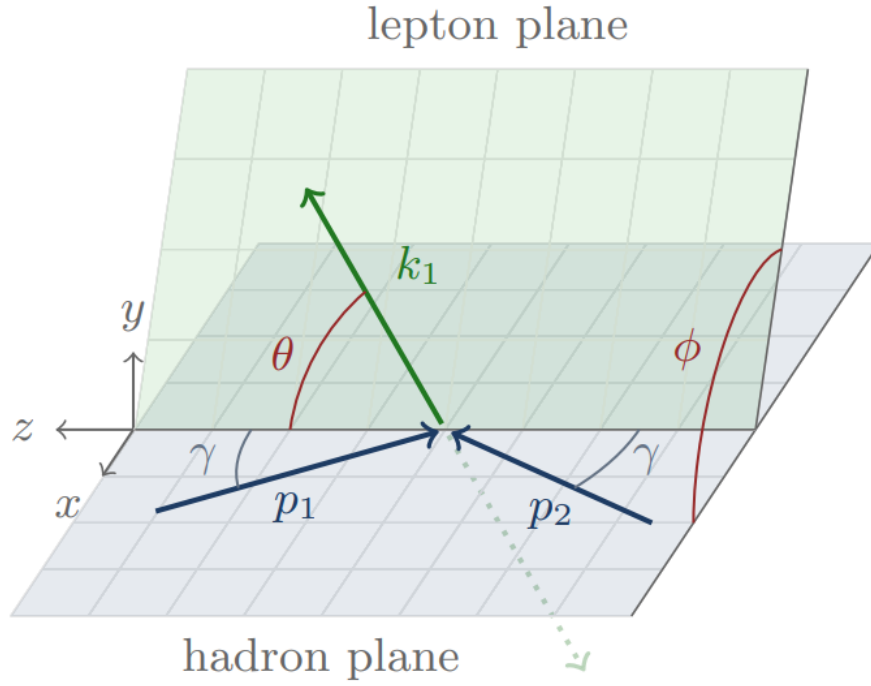


FIGURE 2.11: The definition of the Collins–Soper angles in the dilepton rest frame. In this frame, p_1 and p_2 are the momentum vectors of the colliding hadrons. k_1 is the momentum vector of one of the leptons. Reproduced from Ref. [41].

J is a combination of the spin and orbital angular momentum l . Since these vector bosons have spin 1, their helicity states allow for $J = 0, 1, 2$. Therefore, the angular distribution of the decay products can be described by spherical harmonics with $l = 0$ (s-wave), $l = 1$ (p-wave), and $l = 2$ (d-wave). Higher l values correspond to more complex angular distributions that are suppressed at leading-order, making $l \leq 2$ sufficient to describe the angular cross-section. In higher-order corrections, terms with $l > 2$ can appear, arising from more complex interaction dynamics such as loop corrections or higher-order interference effects. These contributions, while possible, are typically suppressed in comparison to the lower-order terms and therefore contribute less significantly to the overall angular distribution [29].

Polarisation is defined as the orientation of the vector boson spin with respect to its momentum. The polarisation of the vector boson can hence be determined by measuring the angular distribution of the final-state leptons.

2.3.4 W boson mass at higher-orders

In the case of the mass of the W boson, the tree-level expression involves the Higgs vacuum expectation value and the weak isospin coupling, offering an initial approximation of the mass of the W boson, as seen earlier in Equation 2.11.

However, to achieve greater accuracy, it is crucial to consider the effects of quantum corrections

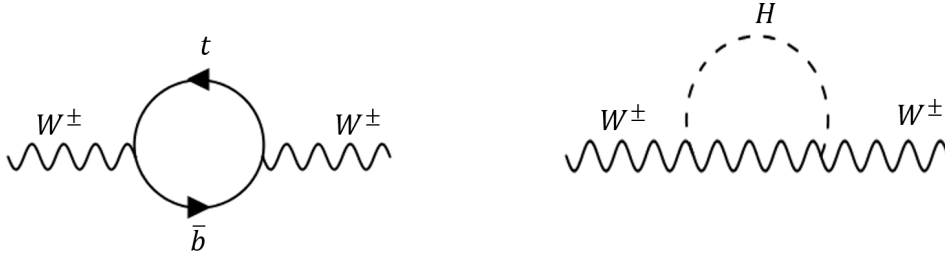


FIGURE 2.12: Corrections to the W boson propagator from the top quark (left) and the Higgs boson (right).

that arise from loop diagrams involving virtual particles. These corrections, known as “loop-level corrections” or “higher-order corrections”, account for the intricate quantum-mechanical interactions that influence the physical properties of particles. At loop-level, virtual particles temporarily emerge and contribute to the mass and coupling constants, resulting in corrections to the tree-level predictions. The W mass can be linked to the Z boson mass (m_Z), the fine-structure constant (α) and the Fermi constant (G_F) by the following equation:

$$m_W^2 \left(1 - \frac{m_W^2}{m_Z^2}\right) = \frac{\pi\alpha}{\sqrt{2}G_F} (1 + \Delta), \quad (2.22)$$

where the term Δ encompasses all the extra contributions to m_W that come from corrections of a higher-order to the W boson propagator [42].

At this higher-order, the mass of the W boson is influenced by the masses of the top quark and the Higgs boson through radiative corrections encapsulated in Δ , which affect the overall behaviour and properties of the W boson. Additionally, new particles and interactions, which are often implied by theories beyond the Standard Model could have a non-negligible effect on m_W through these corrections. Therefore, this iterative process of considering increasing higher-order corrections plays a critical role in advancing our understanding of particle physics and in testing the consistency of theoretical models with experimental data.

The connection between the masses of the W boson, the Higgs boson, and the top quark is illustrated in Figure 2.13. Enhancing the accuracy of the W boson mass measurement in conjunction with the already precise measurements of the masses of the Higgs boson and the top quark can be used to validate the Standard Model.

As briefly mentioned in Chapter 1, the Standard Model’s prediction of the mass of the W boson is constrained by the global fit of the electroweak parameters. At the time of writing this thesis, a global fit of electroweak observables [10, 9], excluding measurements of m_W , yields a prediction of

$$m_W^{\text{pred}} = 80354 \pm 7 \text{ MeV}. \quad (2.23)$$

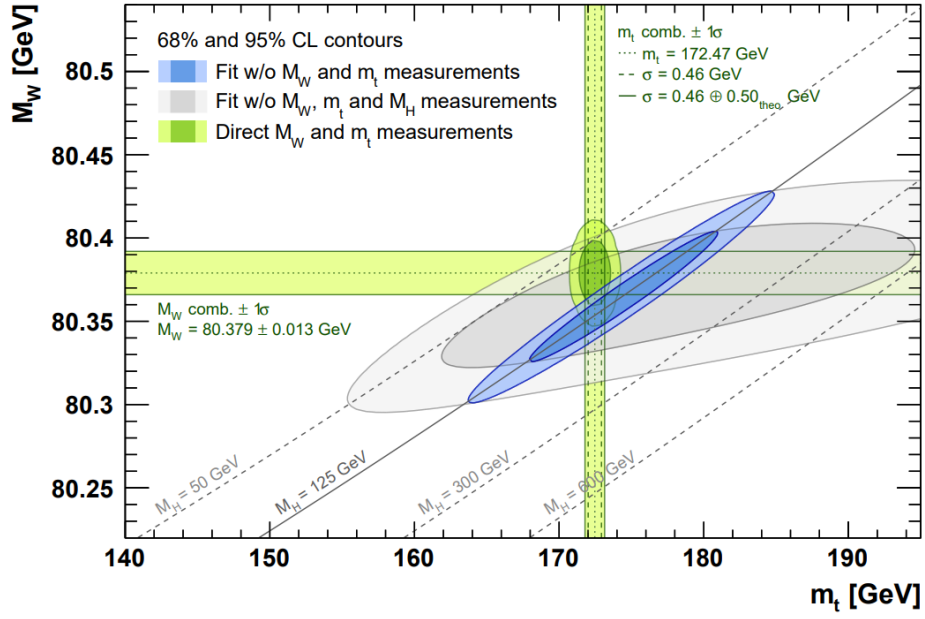


FIGURE 2.13: The relation between the masses of the top quark (m_t), the Higgs boson (m_H) and the W boson (m_W) in the Standard Model is depicted by the grey contour. The direct measurements of these parameters limit the range to the green and blue bands, which is narrower than the phase-space accessible without the measurements. Taken from Ref. [10].

Additionally, the Particle Data Group (PDG), the most prominent organisation for compiling published results for particle physics, currently quotes an average of

$$m_W^{\text{PDG}} = 80377 \pm 12 \text{ MeV}, \quad (2.24)$$

combining experimental results from the LHC, Tevatron and LEP collaborations [32].

The comparison between the theoretical prediction of m_W from global electroweak fits and the experimental measurements is crucial for validating the predictions of the Standard Model and for searching for potential deviations that might indicate new physics beyond the Standard Model.

2.3.5 Higher-order corrections calculations for the analysis

We mainly use the DYTURBO software [43] to produce unpolarised and angular components for our analysis. It can provide such predictions with an accuracy up to NNLO. To enhance our leading-order simulation initially generated by Pythia, we employ DYTURBO predictions, implementing a reweighting strategy. For our analysis, we only use up to “NLO” for the angular component, as it is good enough for our purposes. Additionally, it is prohibitively expensive (computationally) to go beyond this level of accuracy. We do, however, calculate the unpolarised

components using NNLO perturbative accuracy and provide a cross-check using N³LO (next-to-next-to-next-to-leading order) predictions, which are also prohibitively expensive and not yet commonly used by most experiments.

We can also use parton shower algorithms to improve the simulation precision by incorporating higher-order QED corrections such as QED FSR with software programs such as PYTHIA [44] and PHOTOS [45]. This combination provides a more comprehensive description of both hard scattering and subsequent radiation effects. Combining parton showers with fixed-order calculations is a process known as matching, which helps to create a more accurate and realistic simulation of high-energy particle interactions.

The reason why we use the DYTURBO software is that it is specifically designed for the Drell–Yan process and is able to provide predictions at a high level of accuracy. It is important to have precise predictions to compare with the experimental data. The software is also able to provide predictions for the angular components of the Drell–Yan process, which is essential for our analysis, as we are interested in studying the angular distribution of the final-state leptons.

DYTURBO relies on a parameter referred to as the "g-factor," which controls the non-perturbative component of the cross section. This is similar to the intrinsic k_T parameter found in other simulators, such as POWHEG [46], which represents the small transverse momentum partons carry within the incoming hadrons and acts as a non-perturbative factor. This g-factor will be used alongside the value of α_s to fit m_W . Further details on this process can be found in Section 7.

Chapter 3

The LHCb experiment at the LHC

“It doesn’t matter how beautiful your theory is, it doesn’t matter how smart you are. If it doesn’t agree with experiment, it’s wrong.”

– Richard P. Feynman

The pursuit of unraveling the fundamental components of the universe has driven the evolution of increasingly sophisticated particle accelerators. Among these scientific marvels, the Large Hadron Collider (LHC) stands as an exceptional achievement, pushing the boundaries of human knowledge in particle physics. This chapter focusses on the specifics of the Large Hadron Collider beauty experiment (LHCb) at the LHC, providing a comprehensive overview of its design, operation, and the significance of the high-energy particle collisions it generates in advancing our understanding of the subatomic world. Sections 3.1 and 3.2 introduce LHC and LHCb, respectively. We then review the tracking system and particle identification in LHCb in Sections 3.3 and 3.4. The trigger systems in LHCb are discussed in Section 3.5, followed by an introduction to data processing and reconstruction in Section 3.6. We include a brief outline of the simulation process in LHCb in Section 3.7, and we end with Section 3.8 discussing recent advancements and future developments for the LHCb experiment.

3.1 The LHC

The Large Hadron Collider (LHC) stands out as a remarkable achievement in particle physics. Located beneath the Franco–Swiss border near Geneva, Switzerland, the LHC is an expansive particle accelerator and collider complex managed by the European Organisation for Nuclear Research (CERN). With a circumference of 26.7 km, it operates at a depth varying between 45–175 m below the Earth’s surface. Comprising a complex network of tunnels and powerful magnets, the LHC accelerates protons and heavy ions to nearly the speed of light, facilitating high-energy collisions. These collisions occur at four main detectors positioned around the LHC’s ring: ATLAS (A Toroidal LHC Apparatus) [38] and CMS (Compact Muon Solenoid) [39], serving as “general purpose” detectors; ALICE (A Large Ion Collider Experiment) [47], focused

on exploring quark–gluon plasma through ion collisions; and LHCb [37] (Large Hadron Collider beauty), the focus of our detailed discussion in the following sections.

The LHC operates at temperatures colder than those of outer space, approaching absolute zero, and the vacuum conditions within the accelerator are more stable than those of outer space [2]. This controlled environment enables scientists to recreate conditions just moments after the Big Bang, allowing them to explore the fundamental nature of matter, identify new particles, and scrutinise their interactions, ultimately advancing our understanding of the universe at its most basic level.

The approval for the creation of the LHC was obtained in 1994 by a resolution of the CERN council. With substantial financial backing from a diverse array of contributors spanning both member and non-member states, a resounding commitment was made to erect this groundbreaking apparatus, engineered to achieve a staggering centre-of-mass energy of 14 TeV [2].

Construction began after the shutdown of the Large Electron–Positron Collider (LEP) in November 2000. Remarkably, the LHC cleverly repurposed the preexisting experimental infrastructure of LEP, harnessing the cost-effective strategy of employing proton injection into its operational framework. This strategic choice not only expedited the development of the LHC but also optimised resource allocation, culminating in a testament to scientific and fiscal efficiency [2].

The LHC can be used to accelerate and collide both protons and ions. We will focus on proton–proton (pp) collisions within this thesis. The experiment process begins with hydrogen atoms being ionised to produce protons [2]. These protons then undergo an acceleration using a series of particle accelerators.

Beginning at the linear accelerator (LINAC 2), protons gain energy up to 450 MeV. They then go through circular synchrotron accelerators, starting with the Proton Synchrotron Booster. Once they reach 1.4 GeV, the Proton Synchrotron (PS) takes over, raising their energy to 25 GeV. In the PS, the proton beam is collimated into groups called “bunches”, each containing around 1.5×10^{11} protons. These bunches are spaced at least 25 ns apart. The LHC is engineered to support a total of 2808 such bunches per beam, underscoring its substantial experimental capabilities. The protons are then passed to the Super Proton Synchrotron (SPS), which acts as the final stage before entering the LHC at an energy of 450 GeV. This process takes around 16 minutes [2]. The LHC accelerator complex and experimental setup are demonstrated in Figure 3.1.

The LHC employs a set of eight radio-frequency cavities to allow beams to reach their final energy levels. The beams’ control and acceleration are based on 1232 dipole magnets, each generating an 8.4 T magnetic field. These magnets are chilled to 1.9 K for the required strong field, achieved through superfluid helium cooling. The LHC beams navigate ultra-high vacuum conditions better than 10^{-8} mbar.

To focus the beams onto the detectors, quadrupole magnets narrow them down to roughly $14 \mu\text{m}$ [2].

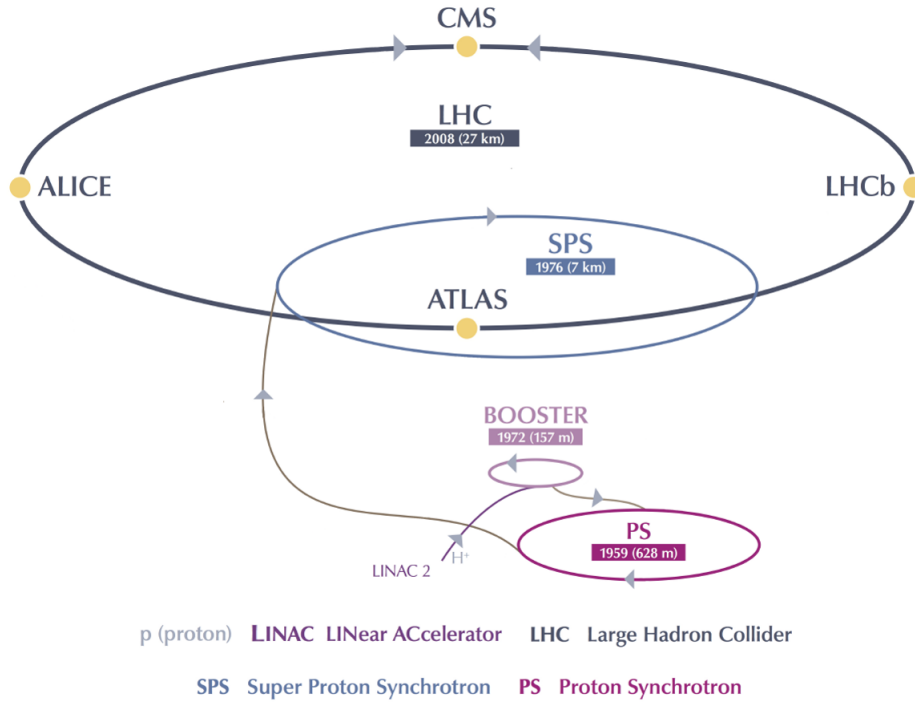


FIGURE 3.1: Schematic diagram of CERN's accelerator complex. The accelerator chain used for pp collisions is: LINAC 2 $\rightarrow$ BOOSTER $\rightarrow$ PS $\rightarrow$ SPS $\rightarrow$ LHC. Reproduced from Ref. [48].

The initial operation of the LHC was planned for 2008. However, in September 2008, a mechanical failure disrupted these plans when, merely nine days after the collider's first beam circulation, a cable connecting two of the beam-focussing quadrupole magnets experienced a sudden and excessive rise in temperature, leading to its melting. Due to this incident, the liquid helium intended for magnet cooling inadvertently boiled off, leading to the rupture of the machine. Consequently, approximately 1 tonne of gas was released into the LHC tunnels. In response, nearly 50 magnets needed replacement [49, 50].

As a result of this catastrophe, the first proper run of the LHC had to wait until November 2009. Run 1 was initiated, involving collisions at 7 TeV, progressing to 8 TeV in 2012. This phase marked the onset of physics measurements, notably culminating in the renowned discovery of the Higgs boson through the ATLAS and CMS experiments, an announcement made in July 2012 [6, 7].

The first long shutdown occurred between 2013 and 2015, focussing on enhancements and preparations to achieve a collision energy of 13 TeV. Subsequently, Run 2 began in April 2015 and continued through December 2018, successfully achieving the nominal instantaneous luminosity (defined as the number of events collected per unit of time with a given interaction cross-section) of $10^{34} \text{ cm}^{-2} \text{ s}^{-1}$ with a bunch crossing rate of up to 40 MHz. This luminosity is not the one used by LHCb — that is discussed later. The dataset used in this thesis is, in fact, derived from this second operational phase, particularly during the year 2017. A table with the LHCb performance numbers for the integrated luminosity over the years can be found in Table 3.1.

Year	Centre-of-mass energy (TeV)	Integrated luminosity (fb^{-1})
2011	7	1.0
2012	8	2.0
2015	13	0.30
2016	13	1.6
2017	13	1.7
2017	5.02	0.1
2018	13	2.1

TABLE 3.1: Integrated luminosities for each year with their corresponding center of mass energies. Reproduced from Ref. [51].

The second long shutdown, initially planned from late 2018 to 2021, was significantly prolonged due to the disruptive impact of the COVID-19 pandemic, extending well into 2022. This second downtime was mainly allocated to gear up for the high-luminosity upgrade of the LHC (HL-LHC) in preparation for Run 4.

Run 3 began in July 2022 at centre-of-mass energy $\sqrt{s} = 13.6$ TeV and is planned to continue for the next four years. All 4 main LHC experiments expect more collisions to be recorded during the third run than the previous two combined. The LHCb experiment, having undergone a complete overhaul, seeks to increase its data-taking rate by a factor of ten, while ALICE aims to increase the number of recorded collisions by a staggering fifty times [52].

3.2 The LHCb detector overview

The LHCb detector [37] is a specialised instrument dedicated to the exploration of heavy-flavour physics. This is achieved by investigating the generation and decay processes of particles that encompass beauty (b) or charm (c) quarks. In pp collisions, the primary mechanism for the production of b or c quarks is gluon fusion [37]. There is a stochastic preference for the two interacting gluons to possess an uneven momentum distribution. Also, given $m_b \ll \sqrt{s}$, there is a lot of energy left over to produce these b hadrons. Therefore, the pairs $b\bar{b}$ and $c\bar{c}$ are boosted along the beam line. Figure 3.2 presents the relative production cross-section of the produced $b\bar{b}$ pairs as a function of the polar angle of each quark, where the coverage of the LHCb experiment is highlighted in red. The LHCb detector was designed to adequately cover the detection of numerous pair-produced B-hadrons. Functioning as a single-arm forward spectrometer, it covers the pseudorapidity range of $2 < \eta < 5$. The pseudorapidity was introduced, among other useful variables in Section 1.2. Similar to other detectors, the LHCb geometry also employs a right-handed Cartesian coordinate system along with spherical polar coordinates (r, ϕ, θ). The origin is established at the particle interaction point, where the z-axis extends along the beam axis toward the detector, the y-axis is vertical, and the x-axis is horizontal, pointing toward the outside of the LHC ring. The layout and geometry of the detector are illustrated in Figure 3.3.

The LHCb detector operates with a moderate luminosity around $10^{32} \text{ cm}^{-2} \text{ s}^{-1}$, deliberately chosen in contrast to the nominal luminosity of the LHC of around $10^{34} \text{ cm}^{-2} \text{ s}^{-1}$, which is typically used in the ATLAS and CMS experiments. This choice offers numerous benefits; reduced

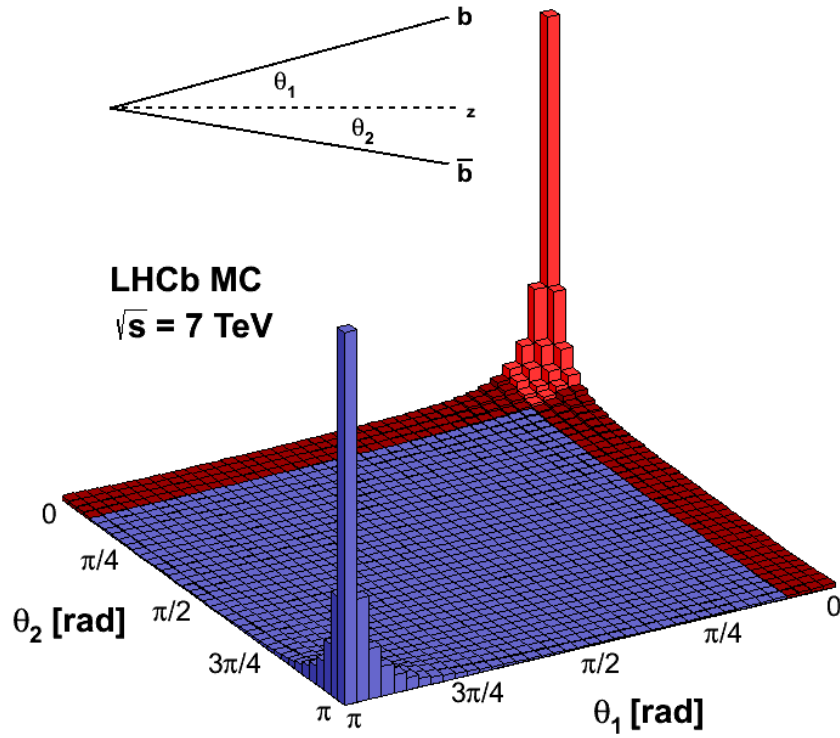


FIGURE 3.2: Angular distribution of $b\bar{b}$ pairs at the LHC under Run 1 conditions, reproduced from Ref. [53]. The red segment represents the acceptance region of the LHCb detector, with the lighter red area indicating instances where both b -hadrons originate within LHCb's acceptance limits. The transition to $\sqrt{s} = 13 \text{ TeV}$ for Run 2 yields no substantial qualitative variation in the diagram.

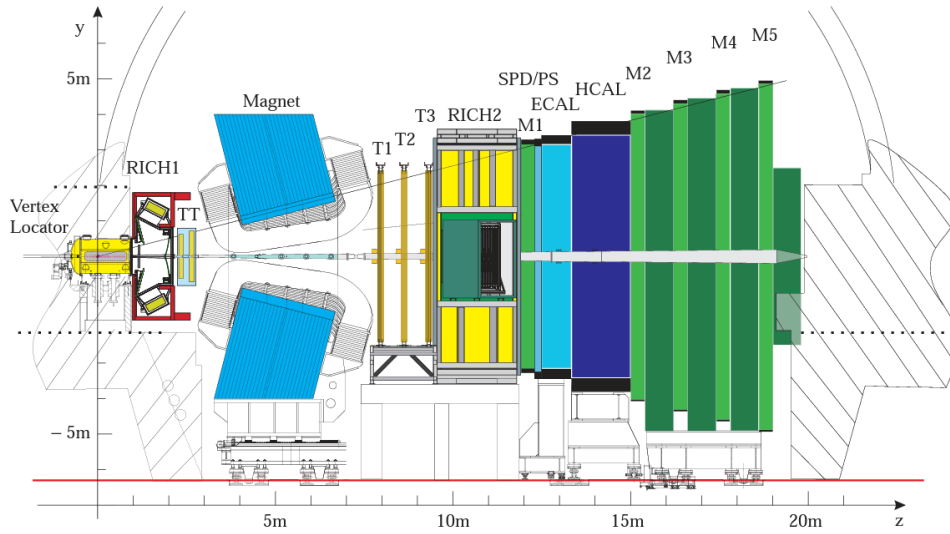


FIGURE 3.3: Schematic depiction of the LHCb detector, adapted from Ref. [54]. The diagram shows the different subsystems that make up the detector, as well as the utilisation of Cartesian coordinates, with the x -axis extending into the plane of the page.

luminosity maintains low detector occupancy and lessens radiation damage. Practical reasons also played a role, primarily concerning the trigger. Running at increased luminosity would not result in a greater number of recorded B signals (specifically those with hadronic decays), as the trigger system would have to impose stricter criteria to limit the rate of data taking to a manageable level. The applied hardware trigger reaches its limit and saturates at the Run 2 luminosity. More details on the trigger are discussed later in 3.5. For beauty (b) and charm (c) quarks, which produce heavy hadrons with relatively longer lifetimes, lower luminosities prove helpful. These heavy hadrons travel short distances before decaying into other particles. By operating at lower luminosities, LHCb effectively examines these decays, minimising multiple interactions within a bunch crossing, and simplifying decay analysis. In Run 1 and Run 2, an average LHCb event predominantly involves only one pp interaction [55], where an event is defined as a single bunch crossing. This is in stark contrast to the ATLAS and CMS experiments, which typically encounter around 20–40 interactions per event.

The LHCb detector accumulated a combined integrated luminosity of 3 fb^{-1} during Run 1 and 5.7 fb^{-1} during Run 2, resulting in a cumulative total of 8.7 fb^{-1} . The integrated luminosity for pp collisions at LHCb is depicted in Figure 3.4.

The LHCb detector consists of a precise tracking setup and a set of subdetectors, which are mainly used to identify different types of particles. The tracking system, which a particle encounters after being produced in a collision, includes components such as the Vertex Locator (VELO), a silicon-strip detector called the “Tracker Turicensis” (TT), a magnet, and three stations (T1–3) of silicon-strip detectors and straw drift tubes.

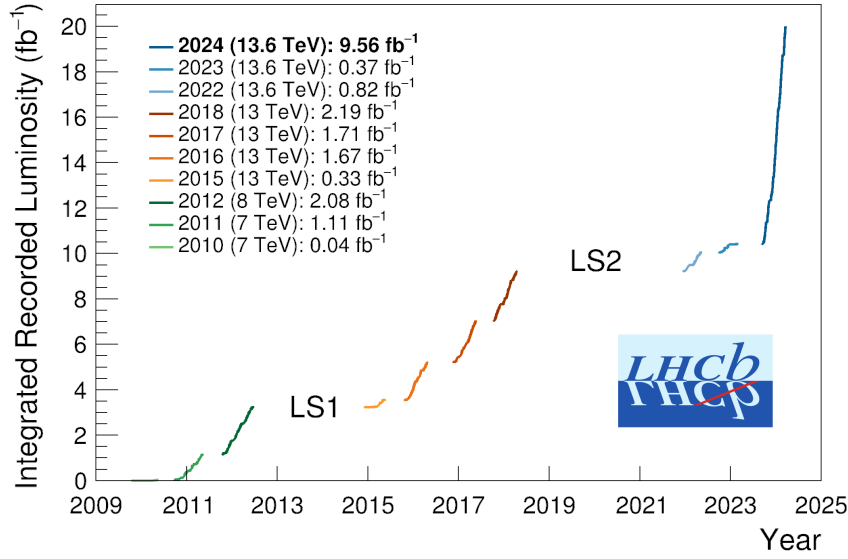


FIGURE 3.4: LHCb's integrated luminosity for pp collisions for 2010 to 2023 runs. Reproduced from Ref. [56]

To distinguish between various types of hadrons, the particle identification (PID) system comprises two Ring-Imaging Cherenkov (RICH) detectors; RICH1 and RICH2, positioned before and after the magnet, respectively. This system helps classify different types of hadrons. Further discrimination among photons, electrons, and hadrons is achieved by a four-part calorimeter system: the Scintillating Pad Detector (SPD), PreShower (PS), Electromagnetic Calorimeter (ECAL), and Hadronic Calorimeter (HCAL). Finally, there is a muon identification system which is made up of alternating layers of iron absorber and multiwire proportional chambers (M1–M5).

The relevant collision events undergo a filtering process that uses a two-stage trigger system in Run 2, consisting of hardware and software components. Subsequently, a complete event reconstruction is performed, enabling complicated selections prior to storage on servers. Reconstruction is a crucial step in the analysis process, as it involves the identification of particles and the determination of their momenta and energies. The reconstruction process is discussed in more detail in Section 3.6.

Although the primary aim of the LHCb experiment is to seek indirect evidence of new physics through investigations of violation of CP and rare decays, it is also capable of conducting precise studies in electroweak physics, which is what we explore in the analysis discussed in this thesis. This aspect will be further emphasised in the following sections.

3.3 Tracking

The VELO, Tracker Turicensis (TT), dipole magnet, and Tracking stations (T1–T3) collectively form the foundation of the LHCb tracking system, responsible for mapping the 3D trajectory of

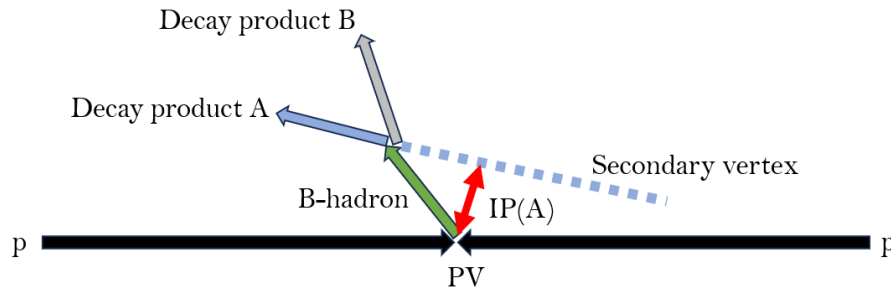


FIGURE 3.5: Illustration depicting the decay of a B-hadron into its products A and B. The impact parameter (IP) of A represents its closest distance to the primary vertex (PV).

charged particles and quantifying their momentum.

3.3.1 The Vertex Locator

The Vertex Locator (VELO) [57] plays a crucial role in the LHCb experiment by providing precise tracking of charged particles produced in high-energy pp collisions. Located closest to the interaction point, the VELO consists of several layers of silicon microstrip detectors arranged concentrically around the beam pipe [57]. These detectors allow for the accurate determination of particle trajectories, enabling the reconstruction of primary vertices (PV) with a spatial resolution of around 10 μm and secondary vertices with a resolution of approximately 50 μm . This is particularly important in LHCb since many analyses search for secondary vertices where a b-hadron decays, whose lifetimes of around 1 ps allow them to travel a few mm before decaying. Of equal importance is the calculation of impact parameters (IP), the distance of closest approach between tracks and their PVs. The VELO's close proximity to the collision point facilitates the study of short-lived particles, particularly those containing beauty and charm quarks.

A diagram that clarifies the significance of PV and IP is presented in Figure 3.5. The PV resolution along the z axis for four years of data collection, as a function of the number of tracks associated with the primary vertex, is shown in Figure 3.6.

Another interesting feature of the VELO subdetector is that it is built in two parts that move together when the particle beam is collimated. This is necessary because when the protons enter the LHC, they are not close enough together. If the VELO were closed immediately, it would be damaged by radiation. During the initial beam injection, the VELO is about 3 cm downstream of the particle beam. Once the beam becomes stable, the two parts of the VELO come together, maintaining a distance of approximately 5 mm from the beam while surrounding it. Additionally, the left and right VELO modules are intentionally overlapped to enhance the precision of detector alignment and to avoid collision between the components in case of material expansion or the loosening of certain parts over time. You can see a diagram of the VELO layout in Figure 3.7.

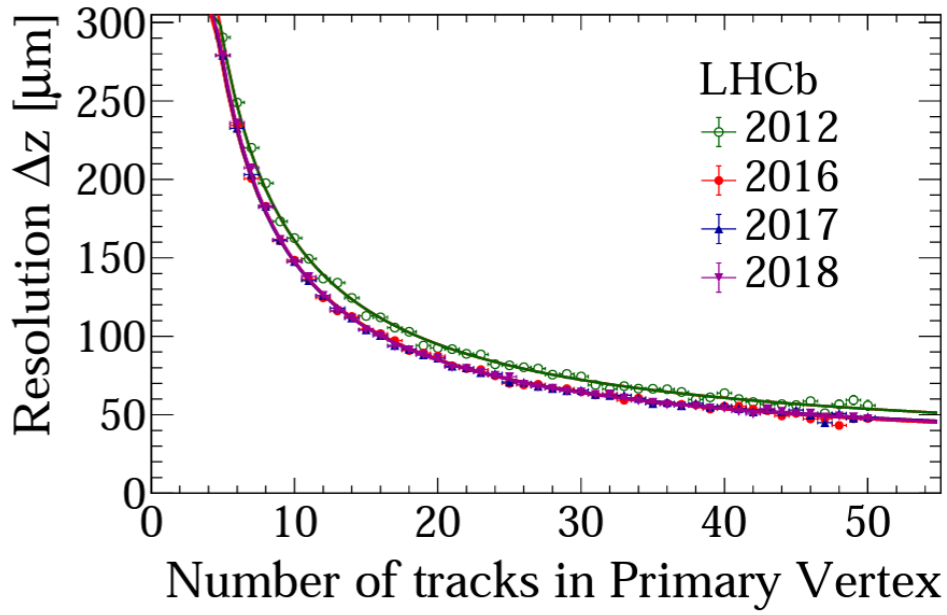


FIGURE 3.6: The primary vertex resolution along z (right) over four years of data taking, as a function of the number of tracks in the PV [59]. Note that the lines from the years 2016, 2017, and 2018 are masking each other. Reproduced from Ref. [58]

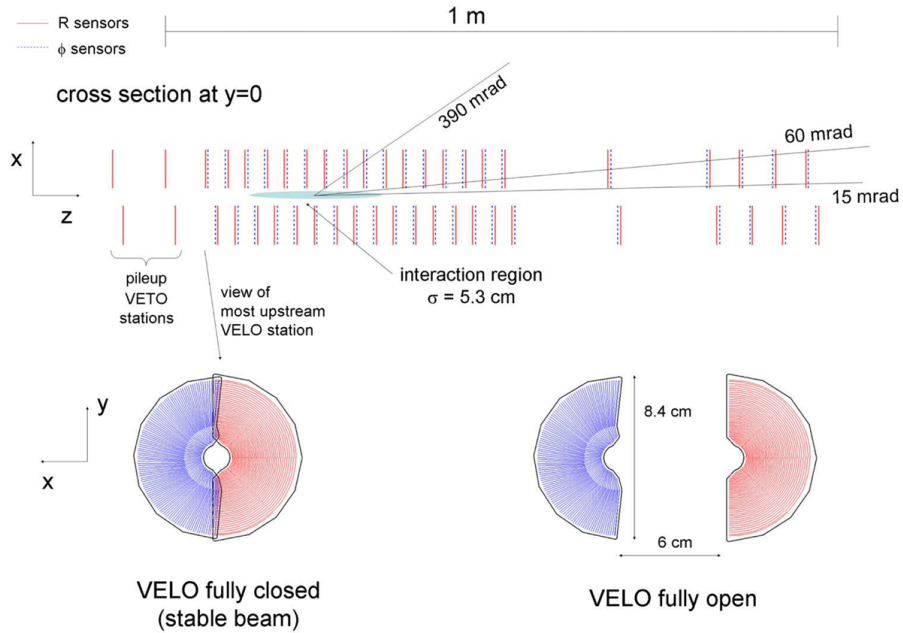


FIGURE 3.7: Schematic diagram of the Vertex Locator detector (VELO) used in Run 1 and Run 2. Reproduced from Ref. [37]

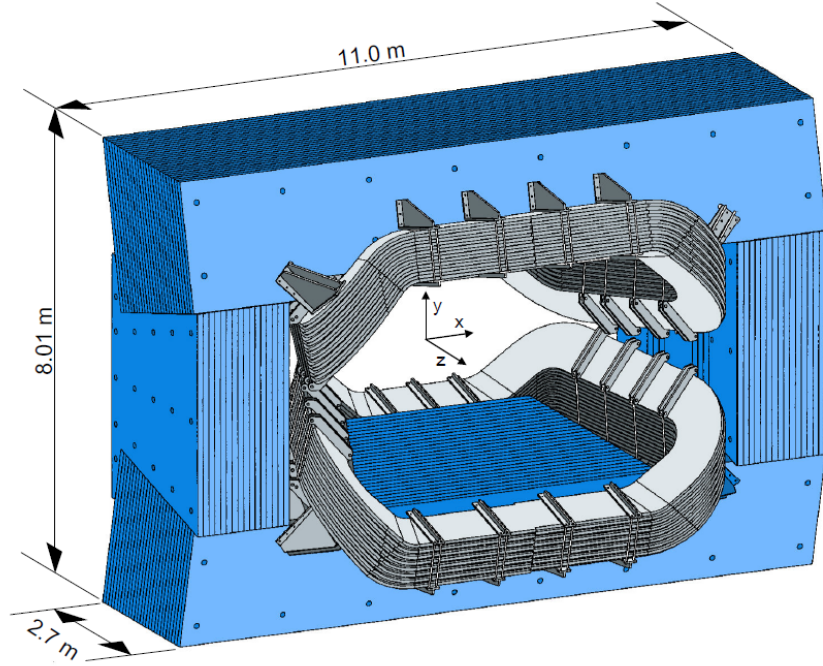


FIGURE 3.8: Engineering diagram of the LHCb dipole magnet, reproduced from Ref. [37].

3.3.2 Magnet

To gauge the charge and momentum of charged particles, a “warm” (non-superconducting) dipole magnet [59] is strategically located after (downstream) of both the VELO and the initial RICH detector, RICH1, as previously illustrated in Figure 3.3. The direction of curvature of a particle in a magnetic field provides information about its charge, whereas the radius of curvature corresponds to its momentum. This dipole magnet comprises two sets of aluminium coils encased within an iron yoke and has a peak magnetic field of 1.1 T. The magnet setup is illustrated in Figure 3.8.

The magnet covers the entire acceptance of LHCb, with the field aligned in the $\pm y$ -direction, so the particles are deflected in the x - z plane. The direction of the magnetic field can be changed at will. This helps in analysing asymmetries, since taking measurements with different field directions and averaging results can cancel out many potential errors, and it also allows for performing cross-checks for detector biases. For instance, we can verify the accuracy of the magnetic field calibration when it is reversed and assess any charge asymmetry that may be induced by this reversal.

The effectiveness of the VELO and RICH detectors is influenced by external magnetic fields. Straight tracks in the VELO enhance vertex reconstruction and event triggering. In the same way, magnetic fields alter the images captured by the RICH photon detectors. To mitigate this impact, both detectors are shielded from magnetic interference. Additionally, the magnet’s design ensures a lower magnetic field where these detectors are positioned, while it is stronger around the tracking stations. The magnetic field strength is shown in Figure 3.9

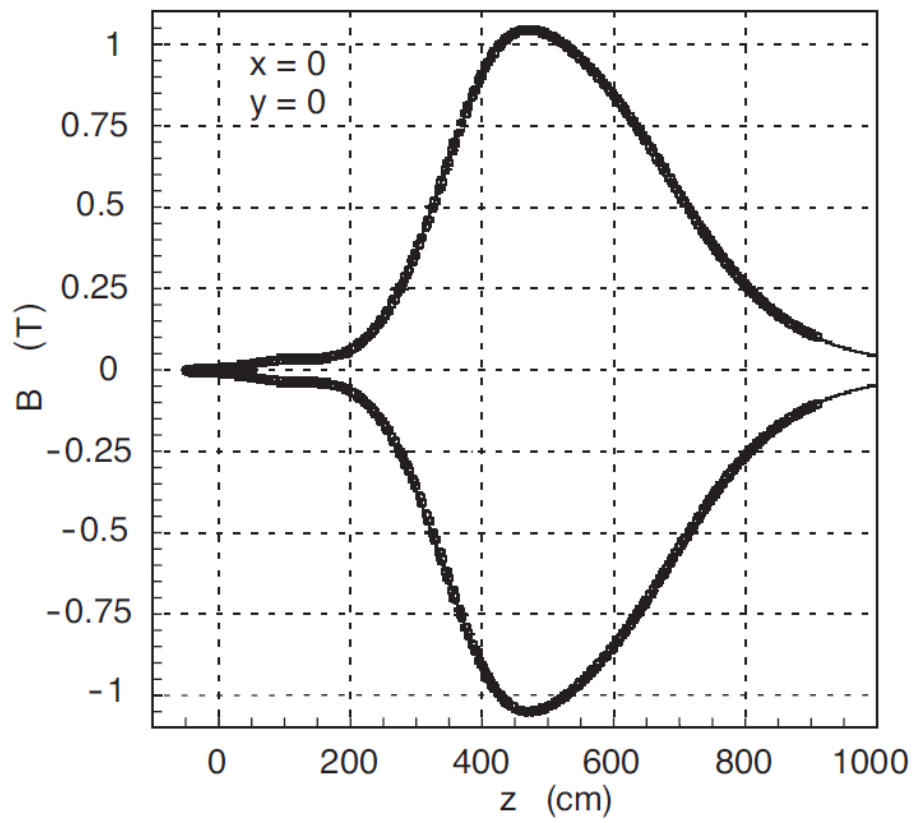


FIGURE 3.9: Vertical component of the magnet's field strength as a function of z . Note that fields are shown for both polarities. Reproduced from Ref. [37].

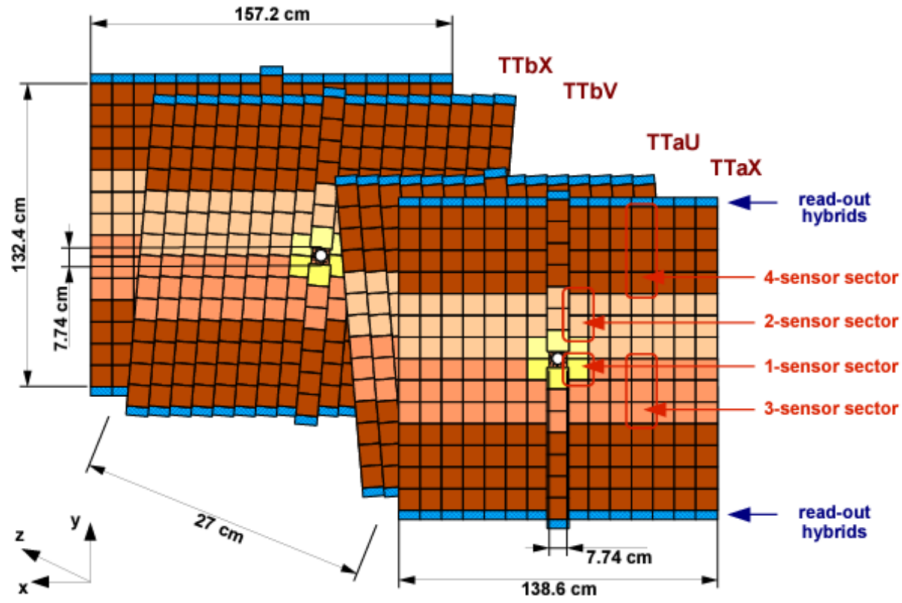


FIGURE 3.10: The TT layout features the LHC beam pipe running through a central aperture within the detection layers. These four layers are named TTaX, TTaU, TTbV, and TTbX. Reproduced from Ref. [60].

3.3.3 Tracker Turicensis

The Tracker Turicensis (TT) was a 150 cm wide and 130 cm high planar tracking station that was positioned right before (upstream of) the LHCb dipole magnet. The TT consisted of four layers of silicon microstrip sensors, offering a precise spatial resolution of 50 mm for single-hit measurements across the detector's acceptance range. The layout of the TT subdetector is shown in Figure 3.10.

To handle high-track densities, the modules nearest to the beamline are divided into three sectors: L, M, and K. The remaining modules are grouped into two sectors, L and M. Furthermore, the two inner layers are positioned with a slight offset of -5° and 5° relative to the vertical axis along the x - y plane. This arrangement aims to enhance the resolution of the y coordinates. The TT was replaced by a new silicon strip detector for Run 3 [60].

3.3.4 Tracking stations

Following the dipole magnet, the trio of tracking stations (T1–T3) assumes the task of determining the track trajectories [60]. This role is critical to achieve both elevated momentum precision and optimal track reconstruction efficiency. Moreover, they can effectively trace long-lasting particles that decay after passing through the VELO. To deal with the high multiplicity of tracks in the region close to the beam line, each station is divided into two parts: the Inner Tracker (IT) uses the same crossed-silicon micro strips found in the TT, while the Outer Tracker (OT) is equipped with drift tube detectors. Since both IT and TT have silicon designs, they are collectively called the “silicon trackers” (ST).

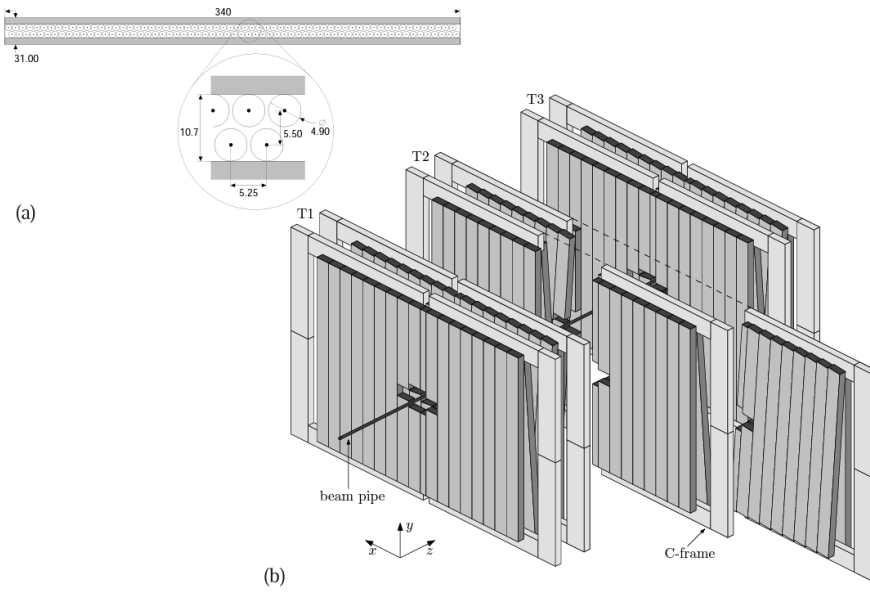


FIGURE 3.11: Configuration of straw-tube modules in Outer Tracker (OT) layers and stations. Reproduced from Ref. [61]

Each station is composed of four layers with the two internal layers rotated by -5° and 5° relative to the vertical axis along the x - y plane. Similarly to the TT case, this helps to enhance the sensitivity of the detector in the y -coordinate.

The OT covers the large area of LHCb acceptance that is not occupied by the IT, with 4 layers of vertically orientated drift tubes, arranged in the same geometry as the ST. This setup is illustrated in Figure 3.11. Each layer has 64 drift tubes of 4.9 mm in diameter, filled with a gas mixture composed of 70% argon, 28.5% carbon dioxide and 1.5% oxygen. This ensures a drift time < 50 ns. The track reconstruction procedure is discussed in more detail in Section 3.6.

3.4 Particle identification

Particle identification (PID) is fundamental for the LHCb physics program. Measurements of the energy of charged hadrons, electrons and photons are provided by a calorimeter system consisting of scintillating pad (SPD) and PreShower (PS) detectors, an electromagnetic calorimeter (ECAL) and a hadronic calorimeter (HCAL). The muon system consists of five muon stations (M1–M5) and is fundamental for the identification and trigger of particle decays into final states containing muons. Muon detection is essential for our research on the decay of Z and W bosons into muons presented in this thesis.

3.4.1 The Ring Imaging Cherenkov detectors

When a particle moves through a material with a refractive index greater than 1, known as a radiator, at a velocity faster than the phase velocity of light within the material, it generates a cone of photons recognised as Cherenkov radiation. True to their name, the Ring Imaging Cherenkov

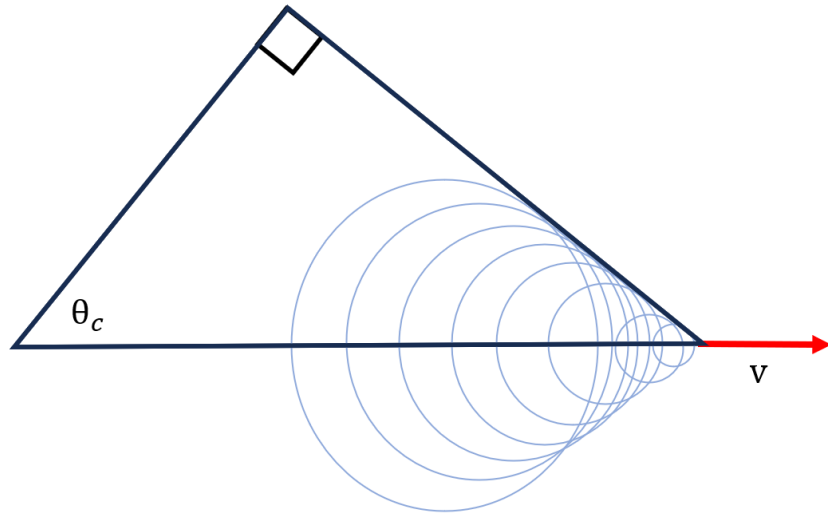


FIGURE 3.12: Geometry of Cherenkov radiation is shown. A particle moving with speed v is represented by the red arrow, and the Cherenkov cone's development is illustrated by the blue rings.

detectors (RICH1 and RICH2) detect and measure the angle of emission of Cherenkov radiation released by charged particles as they pass through [37]. The velocity of the particle, v , can then be determined by

$$v = \frac{c}{n \cos(\theta_c)}, \quad (3.1)$$

where c is the speed of light in vacuum, n is the refractive index of the radiator, and the angle θ_c is defined in Figure 3.12.

Combining these data with momentum details from the tracking systems, we can identify the charged hadrons (pions - $\pi^\pm$, kaons - $K^\pm$, protons - p). Hybrid photon detectors (HPD) are used to detect Cherenkov photons in the wavelength range 200–600 nm. The optics of an HPD are shown in Figure 3.13.

Located before the magnet, RICH1 encompasses the region of low-momentum charged particles, approximately 1–60 GeV/ c , employing aerogel and C_4F_{10} radiators. On the other hand, the subsequent detector, RICH2, is located after the dipole magnet and is responsible for the high-momentum range, starting at around 15 GeV/ c and extending beyond 100 GeV/ c , using a CF_4 radiator. Figure 3.14 shows the Cherenkov angle versus the particle momentum for the RICH radiators.

Both RICH detectors use a mix of spherical and flat mirrors to guide Cherenkov light, ensuring that it spans the acceptance of the spectrometer. In RICH1, the optical design is orientated vertically, while in RICH2, it is configured horizontally. This helps to improve the efficiency of

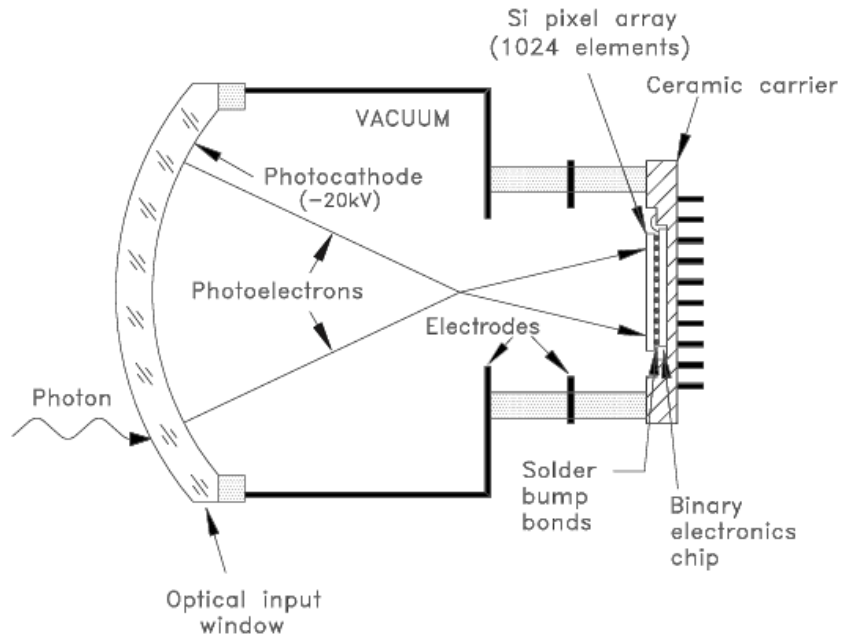


FIGURE 3.13: The HPD's focusing optics play a crucial role in detecting Cherenkov radiation. When a photo-electron is generated at the photo-cathode, it gets accelerated towards the silicon sensor chip, as shown in the diagram. Reproduced from [62].

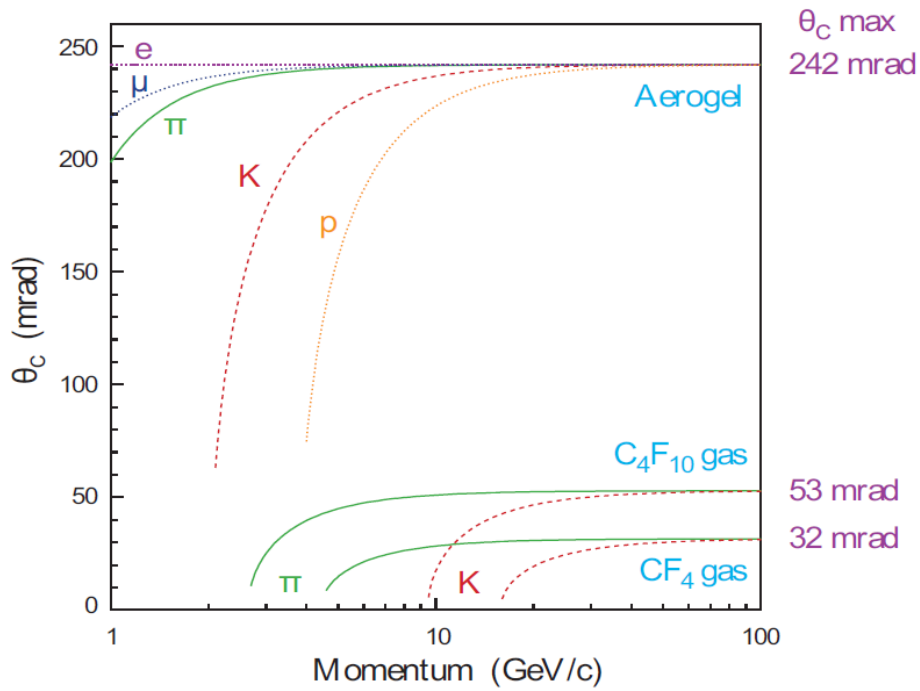


FIGURE 3.14: Cherenkov angles, for different particle species and RICH radiators, as a function of momentum. Reproduced from Ref. [37].

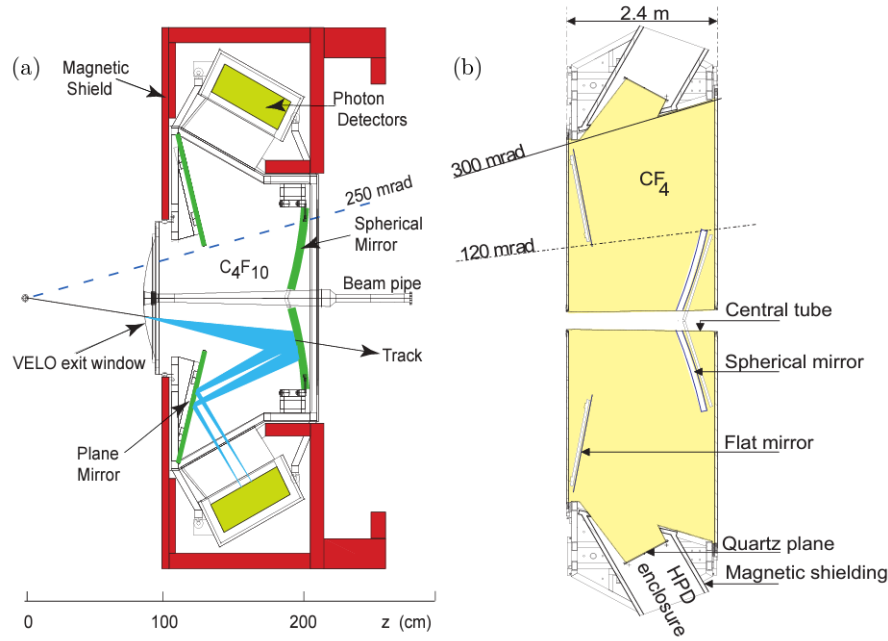


FIGURE 3.15: Optical layout of RICH1 (a) and RICH2 (b). Reproduced from Ref. [62].

photon detection by making the photons approximately normally incident on the HPD planes. Figure 3.15 illustrates both RICH detectors.

3.4.2 Calorimetry

In Run 2, the layout of the LHCb calorimeter system encompassed several components: starting with the Scintillating Pad Detector (SPD), followed by the PreShower (PS) detector, and finally, the electromagnetic and hadron calorimeters (ECAL and HCAL, respectively). A schematic showing the layout of the calorimetry systems is shown in Figure 3.16.

The Scintillating Pad Detector (SPD) and PreShower (PS) systems in the LHCb experiment represent critical elements of its particle detection setup. The SPD, located at the front of the detector, plays a pivotal role in rapidly identifying charged particles emerging from high-energy collisions. Using scintillating material, the SPD efficiently converts the energy of the particles into detectable light signals. The SPD transmits binary data per cell, depending on whether the deposited energy exceeds a particular threshold. This differentiation serves to distinguish between charged particles (that produce a signal) and neutral particles (which produce no signal) upon impact [63]. As such, the SPD is used to discriminate between electrons and photons that occur via bremsstrahlung radiation and pair production. Bremsstrahlung radiation is the electromagnetic radiation emitted by a charged particle when it decelerates. Pair production is the creation of an electron and a positron from a photon, which typically occurs when the photon's energy is greater than 1.02 MeV (twice the electron's rest mass energy).

Following the SPD, the PS system further enhances the particle identification process. It consists of two alternating layers of lead and scintillating material, creating a dual-purpose mechanism.

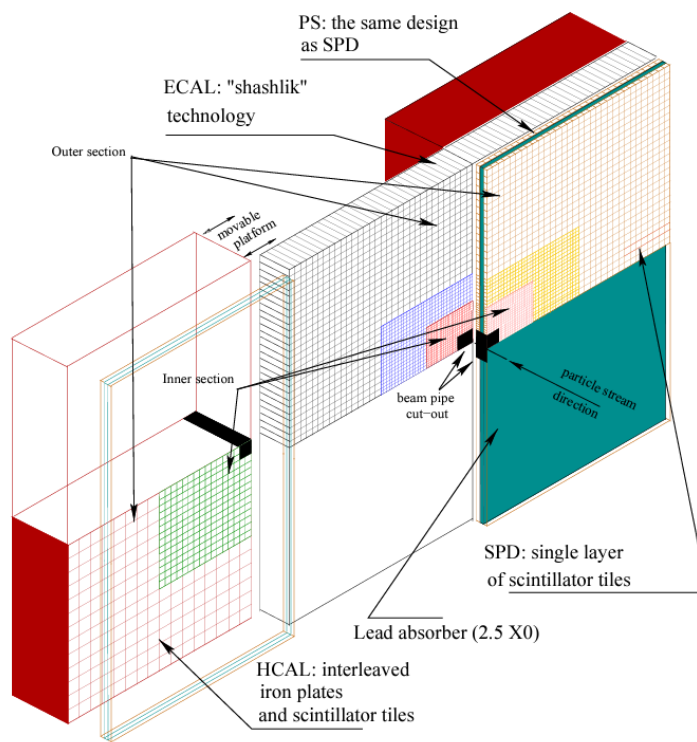


FIGURE 3.16: Layout of the calorimeter system. Reproduced from Ref. [63].

As neutral particles traverse this setup, they generate electromagnetic showers within the lead, which are subsequently detected by the scintillating material. The charged particles created in the showers are then measured in the PS, which enables the separation of electrons and charged pions.

The ECAL system within the LHCb experiment is designed with a “Shashlik” structure, employing alternating layers of 2 mm thick lead absorbers and 4 mm thick plastic scintillating material. As the particles traverse this arrangement, the lead layers initiate electromagnetic showers, which generate light signals in the scintillating material. These signals are then converted into electrical pulses for measurement. This dynamic interplay enables an accurate determination of the energy of electromagnetic particles such as electrons and photons, which is essential for studying the properties of particle collisions. The ECAL is divided into 6016 cells that span the inner, middle, and outer regions. Every cell is monitored by a single Photomultiplier Tube (PMT), ensuring precision. This configuration is pivotal for studying particles such as electrons and photons in collision events.

Similarly to the ECAL, the HCAL also functions as a sampling calorimeter. It uses alternating layers of 3 mm scintillator tiles and 1 cm iron absorbers. These layers connect to photomultiplier tubes through wavelength-shifting absorbers. However, the plates are positioned perpendicular to the ECAL plates along the x -axis. Since hadronic showers are larger, the HCAL does not require the same level of detail as the ECAL for its granularity. Figure 3.17 displays an illustration of both an ECAL and an HCAL cell.

3.4.3 Muon system

Highly efficient reconstruction and identification of energetic muon tracks plays a crucial role in numerous analyses of LHCb physics programs, including studies related to the W boson. The muon system consists of 5 rectangular tracking stations, named M1–M5, and covers the full acceptance region of LHCb [37]. Muons are one of the only particles that can pass through the entire calorimeter system without being intercepted (as well as neutrinos), so the muon detection system is the last subdetector downstream, apart from the first muon station, M1, which is located after RICH2 and before the calorimeter system and provides additional information about the transverse momentum. The other muon stations, M2–M5, are used to identify and trace muons for online and offline analysis and are placed after HCAL.¹ M2–M5 are also interspersed with iron absorbers 80 cm thick, so only muons with a momentum greater than 6 GeV can pass through the entire muon system; anything less will be blocked. This is important because the muon system is used to identify muons from the decay of beauty and charm hadrons, which have high momenta, as well as muons from the decay of the other particles such as the W boson. A diagram of the muon stations is shown in Figure 3.18.

Each station contains 276 chambers, divided into four distinct quadrants (R1–R4), with R1 having the finest resolution and more precise division, since it is closest to the beam pipe, which is

¹There is more discussion of online and offline analysis in Section 3.6.

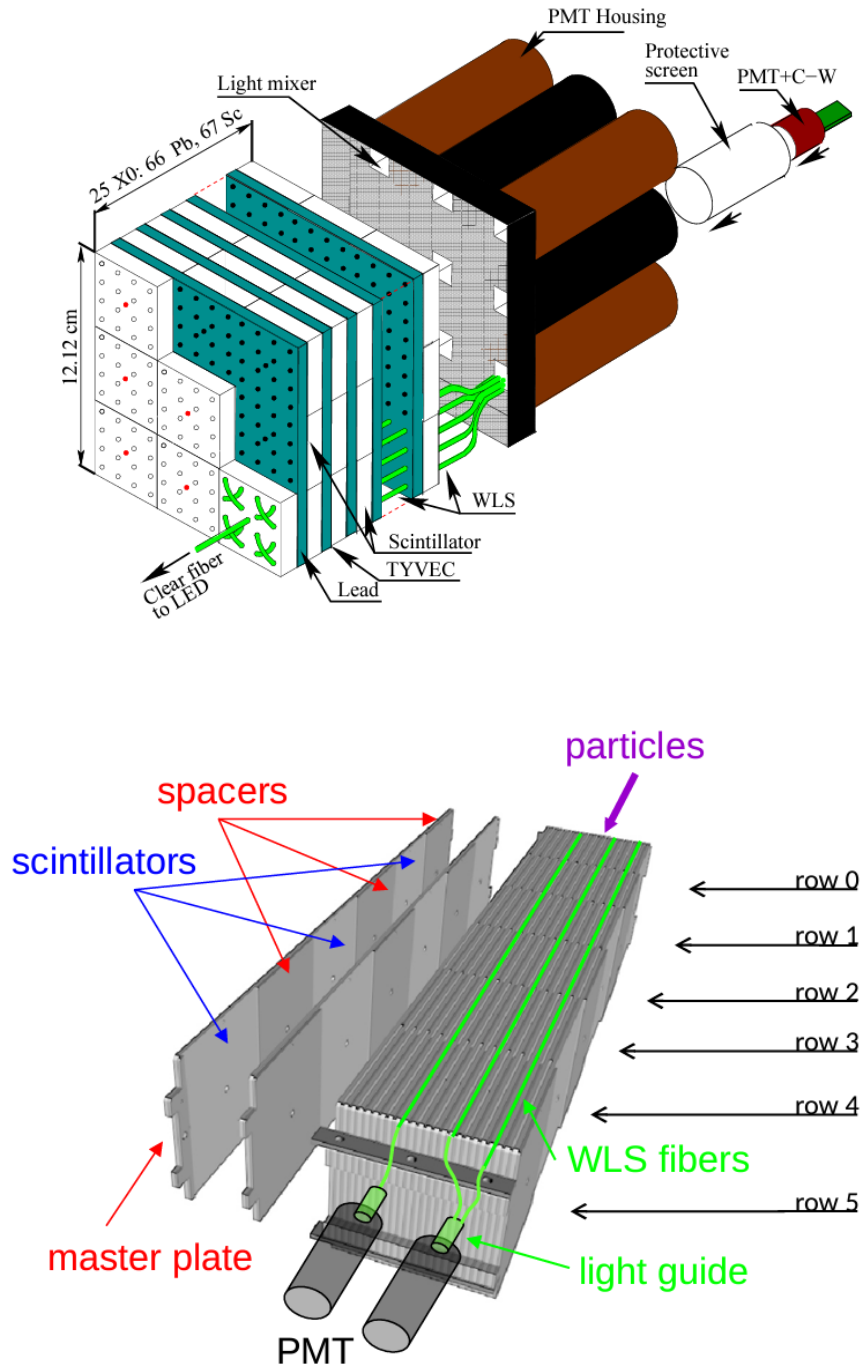


FIGURE 3.17: Diagram of an ECAL cell (on the top) and a HCAL cell (on the bottom). The PMTs located at the rear of each cell are collectively read out as a single detector channel. Reproduced from Ref. [63].

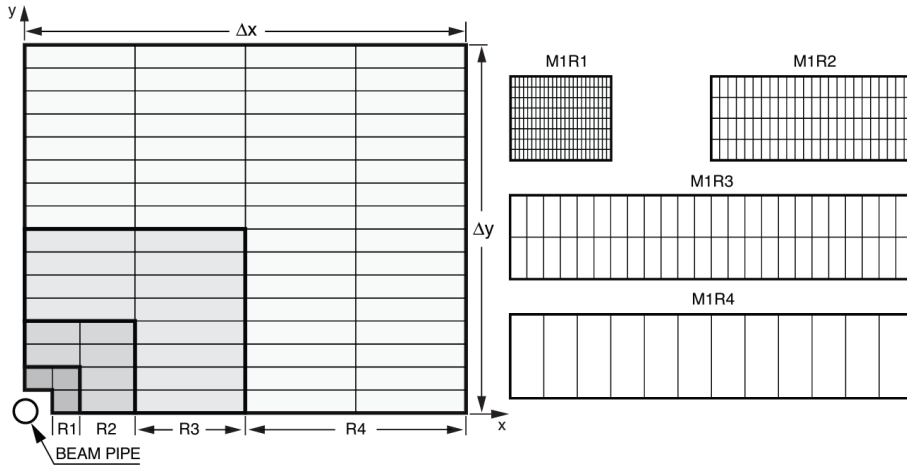


FIGURE 3.19: Left: A quadrant of a muon station is depicted from the front. Each of the rectangles in the image represents a single chamber with a total of 276 chambers per station. Right: this illustrates the separation of compartments in each quadrant, depending on their proximity to the beam pipe. We use the M1 station for this demonstration. Reproduced from Ref. [37].

exposed to a larger number of tracks. A diagram of these quadrants of the muon station is shown in Figure 3.19.

Muon stations use multi-wire proportional chambers (MWPCs) to track particle movement. The only exception is the central area of the M1 station, which instead uses gas-electron multipliers (GEMs) near the beamline due to the greater multiplicity of tracks and radiation levels in this area [64]. In the same way as previously mentioned with the OT, incident muons create ionisation in the drift tubes, and the ionisation electrons move and are amplified when an electric field is applied until they are collected at the anode. A multi-wire proportional chamber has numerous anode wires, which provide positional sensitivity, and the multiplication of the electric field is increased as it gets closer to the wire. GEMs are electrically conductive foils with holes that are precisely cut into them. These holes guide electric field lines through them, thus intensifying the electric field and allowing for multiplication.

3.5 Trigger

The LHCb detector is designed to study the properties of heavy particles that are produced in rare interactions. The LHCb experiment operates at a luminosity lower than that of the other experiments at the LHC. It has been running at a luminosity around $10^{32} \text{ cm}^{-2} \text{ s}^{-1}$, which is about 1/25 of the design luminosity of LHC.

This means that most of the bunch crossings contain only a few pp interactions, which is beneficial for triggering and reconstruction since occupancy is lower than at general-purpose detectors. However, there are still too many interactions to save every event. The LHCb trigger is therefore responsible for reducing the data taken from the maximum interaction rate of 40 MHz

to a manageable rate (around 3–5 kHz) that can be recorded [65]. This is advantageous for permanent data storage, as the average size of an event is usually around 100 kB. It is not practical to store every event, as this would require several TB/s. Given the average event size, the suitable rate can be determined for each stage of the trigger. For HLT2, that would be ~0.5 GB/s of output.

The LHCb trigger uses all the subsystems we introduced earlier in Sections 3.3 and 3.4. Until Run 2, this system consisted of two stages: the Level 0 trigger (L0); which is implemented in hardware, and the High Level trigger; (HLT) which uses software. These two components identified and stored events that contain interesting physics signatures, as well as selected events for calibration. A general overview of the LHCb trigger system in Run 2 is shown in Figure 3.20. After Run 2, L0 was eliminated and graphics processing units (GPU) were used for the first stage of the HLT. Hence, the LHCb trigger system now consists of only the HLT (based entirely on software).

3.5.1 L0 hardware trigger

Since b-hadrons are massive, this results in their decay products usually having a high transverse momentum compared to the abundant background of “soft” QCD processes. The L0 trigger takes advantage of this decay pattern by using a fast readout from the calorimeter and muon systems to pick out events with high transverse momentum. This allows it to reduce the event rate from the original 40 MHz of bunch crossing to 1 MHz. This is accomplished by dividing the ECAL and HCAL into 2×2 cell clusters with transverse energy:

$$E_T = \sum_{i=1}^4 E_i \sin(\theta_i), \quad (3.2)$$

where E_i is the energy deposit in the i -th cell and θ_i is the angle of the cell relative to the beam axis and the average collision point. Events with at least one cluster with measured transverse energy above a certain threshold are selected.

The L0 trigger has distinct subsystems that are executed in parallel. These subtriggers are referred to as **L0Muon** for muons, and **L0Electron** (**L0Photon**) for electrons (photons).

Events with a high multiplicity that would take too long to be processed by the high-level trigger are also eliminated by rejecting those with more than 1000 SPD hits. Multiplicity is simply the number of hits in a detector. In this case, the SPD is used to measure the number of charged particles in the event, which is important to reject events with too many particles to be processed in a reasonable amount of time. This requirement is known as a global event cut (GEC). For most LHCb analyses, the global event cut (GEC) has a minimal effect. However, for Electroweak (EW) analyses, particularly those measuring cross-sections, the GEC has a considerable and intricate influence on the trigger efficiencies of particularly high p_T EW events, which is not satisfactory. A discussion of the removal of GECs by default from the HLT1 system is discussed in Section 4.3.3

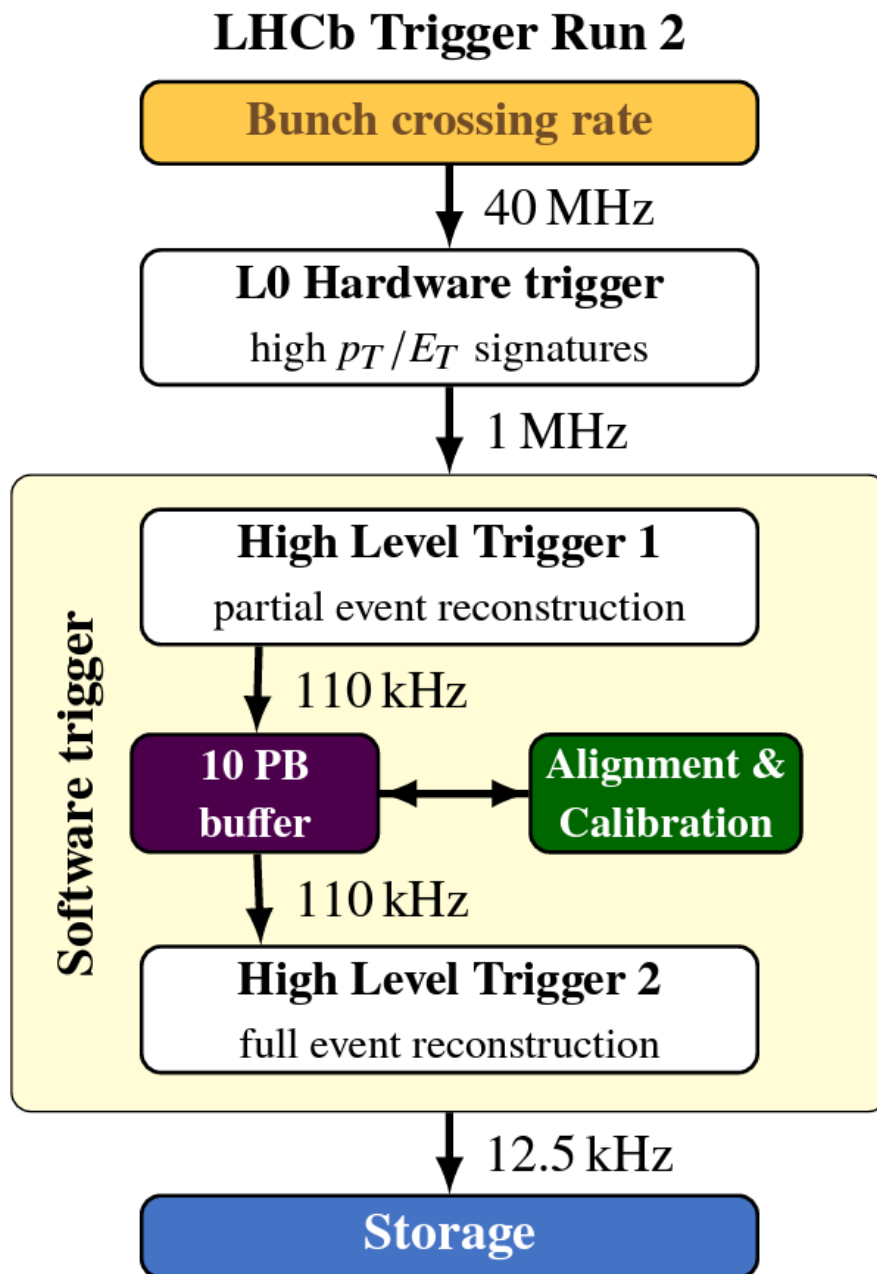


FIGURE 3.20: An overview of the LHCb trigger system in Run 2. Reproduced from Ref. [66].

The analysis in this thesis uses the **LOMuonEW** trigger line, which is designed for high-momentum muons in EW processes. The line accepts only muon tracks with a transverse momentum of more than 6 GeV [66]. A unique feature of the **LOMuonEW** line is that it does not apply a GEC, which helps avoid systematic uncertainties in efficiency calculations.

3.5.2 High level software trigger

The second trigger level, HLT, is a software-based application executed in the Event Filter Farm, which is a large ensemble of CPUs. It uses information from the entire detector to reduce the 1 MHz L0 output to a bandwidth of 12.5 kHz. This is achieved in two steps: HLT1 where the events with tracks that have large impact parameters or matching hits in the muon arm are reconstructed, and HLT2 where the events are fully reconstructed.

HLT1 reduces the event rate from 1 MHz to 40–80 kHz by partially reconstructing events using information from the VELO, tracking stations and muon stations. HLT1 reconstructs the trajectories of charged tracks with a transverse momentum of more than 500 MeV that traversed the entire tracking system (also known as “long tracks”). It also performs a precise reconstruction of the primary vertex. Due to timing constraints, most particle identification algorithms cannot be executed in HLT1, except for a simple muon identification, which is possible due to the clear signature left by muons in the detector. The HLT1 muon PID algorithm requires hits in M2 and M3 for tracks with a transverse momentum of less than 6 GeV, at least one additional hit in M4 or M5 for tracks with a transverse momentum between 6 and 10 GeV, and hits in M2–M5 for tracks with a transverse momentum greater than 10 GeV. Once this stage is complete, the events are temporarily stored on a buffer disk where the alignment and calibration of the entire detector is performed in real-time before being sent to HLT2. This real-time processing of the data is a major success of the Run 2 trigger system and led to the later establishment of the “Real-Time Analysis” project (RTA), which will be covered in detail in Chapter 4.

A trigger line is a predefined condition or set of conditions that, when satisfied, triggers a response or action within a system, typically related to data acquisition or event selection in collider experiments. HLT1 had approximately twenty trigger lines during Run 2. The most relevant line for this thesis was the **Hlt1SingleMuonHighPt** line, which selected events with an identified muon that had a transverse momentum greater than 4.8 GeV [67].

The alignment and calibration of the detector involves a number of steps (e.g. aligning the VELO, tracking stations, MUON, and RICH mirrors; calibrating the global time in the OT and the refractive index of the RICH detector gas, etc.). Data for each subtask are obtained from the HLT1 output. Periodically, the alignment and calibration of the system is performed and the constants associated with alignment and calibration are only adjusted if they are significantly different from the current values [68].

The second stage of the software trigger (HLT2) is a sophisticated set of algorithms that are used to perform a complete event reconstruction. Once HLT1 reduces the trigger rate to about 50 kHz, HLT2 can select events based on the complete reconstruction of the events. This reconstruction is similar to the complete offline reconstruction (discussed in Section 3.6). The main

difference is that a simplified geometry is used for track alignment and RICH information is provided only when required for a limited number of events.

The HLT2 trigger lines are more specific than those in HLT1 and are used to select events for particular physics studies. In Run 2, there were approximately 500 HLT2 lines, including inclusive and exclusive selections. The analysis in this thesis concentrates on the **Hlt2SingleMuonHighPT line**, which requires a transverse momentum of at least 10 GeV. The HLT2 stage filters and reduces the rate of events that are stored on the data disk to 12.5 kHz.

3.6 Data processing

3.6.1 Track reconstruction

The process of track reconstruction and particle identification begins with energy depositions from the VELO and T1–3 stations, which have relatively straight tracks due to the weaker magnetic field at these stations. A Kalman filter is then employed to take into account multiple scattering and energy loss as additional data points are added to the track in an iterative manner [69].

The tracks are then classified according to the detectors with energy deposits associated with that track:

- **VELO tracks** specifically cover the VELO region, typically exhibit substantial angular deviations from the horizontal plane, and are typically used for vertex reconstruction.
- **Upstream tracks** extend through both the VELO and TT sections, usually carry low momenta and possess less precise momentum measurements, and are redirected beyond the acceptance of the spectrometer by the bending magnet.
- **Long tracks** traverse all tracking stations and tend to offer the most accurate momentum measurements.
- **Downstream tracks** pass through the TT and T stations and typically arise from particles that decay outside the VELO.
- **T tracks** only contain information by passing through the T1–3 stations.

A diagram showing all these different tracks is illustrated in Figure 3.21.

3.6.2 Particle identification reconstruction

Neutrals

Neutral particles are characterised by the lack of tracks linked to energy deposits within the calorimetry system. Detecting them involves extrapolating all reconstructed tracks to the calorimetry system, and those that do not match any energy clusters are selected as neutral particles. In the ECAL, these clusters are essentially defined as 3x3 cell patterns centering around the cell

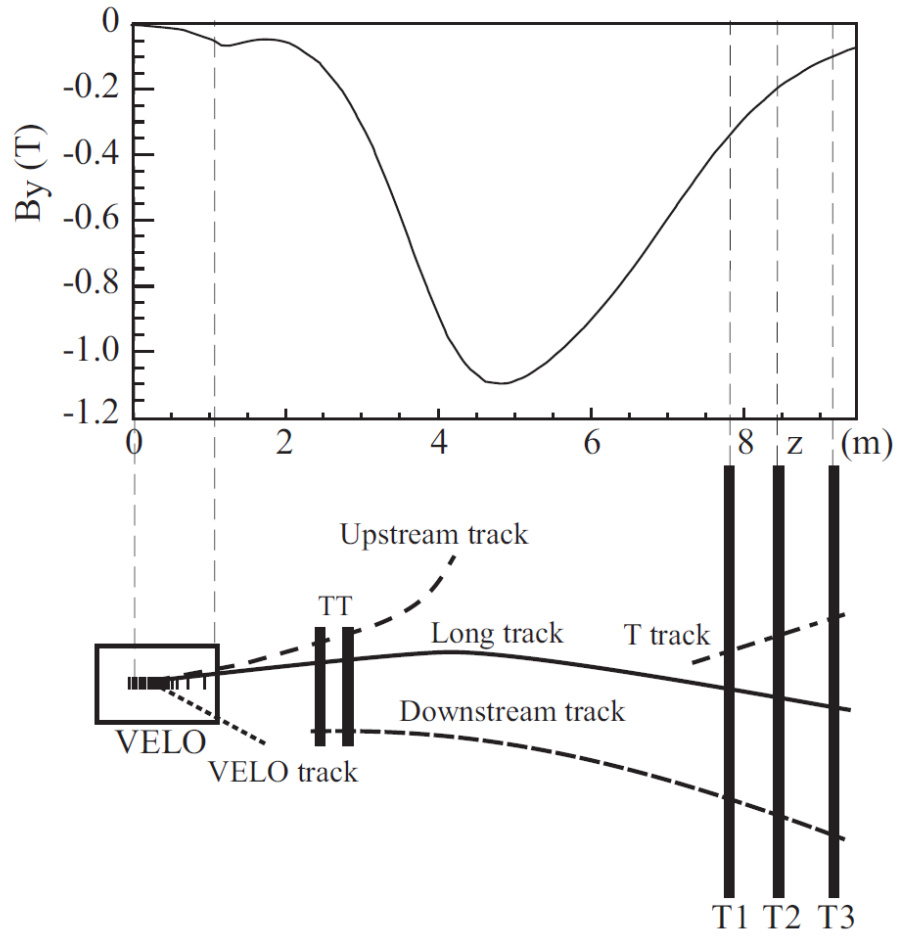


FIGURE 3.21: The illustration at the top displays the B_y component of the B-field as a function of z . Below, the various types of tracks are shown, including long, upstream, downstream VELO and T tracks. This figure is sourced from Ref. [70].

with the highest energy deposit. We evaluate the match using a χ^2_{2D} metric:

$$\chi^2_{2D} = (\mathbf{r}_{\text{track}} - \mathbf{r}_{\text{cluster}})^T (\mathbf{C}_{\text{track}} + \mathbf{S}_{\text{cluster}})^{-1} (\mathbf{r}_{\text{track}} - \mathbf{r}_{\text{cluster}}), \quad (3.3)$$

where the local coordinates of the tracks and clusters, $\mathbf{r}_{\text{track}}$ and $\mathbf{r}_{\text{cluster}}$, respectively, are accompanied by the covariance matrix of $\mathbf{r}_{\text{track}}$, $\mathbf{C}_{\text{track}}$, and the spread matrix of the energy of the cluster, $\mathbf{S}_{\text{cluster}}$.

Neutrals are identified on the basis of clusters with a high value of this metric, typically $\chi^2_{2D} > 4$.

To calculate the energy of non-converted photons, we use the total energy deposited in the ECAL and PS. We also determine the direction of the photon on the basis of where it likely originated. Photon candidates can also arise from the conversion of electron-positron pairs before they encounter the bending magnet. In such cases, we pair the electron and positron to form the photon candidate and adjust the energy measurements to account for any associated bremsstrahlung photons.

Reconstructing neutral pions involves finding two photons that are sufficiently separated. However, when their transverse momentum exceeds 2 GeV, these photon clusters may overlap. To address this, we modify the energy measurements in overlapping cells by fitting the energy distribution of the two photons using simulation data. Additionally, to prevent the incorrect identification of high-energy neutral pions and photons, we utilise a neural network that has been trained on specific particle decay data to distinguish between their unique energy deposit patterns.

Hadrons

Charged hadrons are distinguished primarily using a comprehensive pattern recognition algorithm (not by matching to energy clusters). This algorithm correlates the patterns generated by Cherenkov radiation in the photo-detectors of RICH1 and RICH2 with the anticipated characteristics of a reconstructed track based on a particular mass assumption. The mass hypothesis that yields the highest probability is determined by iteratively considering all potential candidates (pion, kaon, proton, electron, or muon) for each reconstructed track and identifying the best match. For each track, the logarithmic likelihood difference is computed concerning the pion mass assumption, which is the most frequently produced hadron:

$$\Delta \log \mathcal{L}_{X-\pi} = \log \mathcal{L}_X - \log \mathcal{L}_\pi, \quad (3.4)$$

where $\log \mathcal{L}_X$ is the likelihood of a given mass hypothesis X and $\log \mathcal{L}_\pi$ is the likelihood given the π mass hypothesis. This process ensures an accurate identification of hadrons by comparing the computed likelihoods for each candidate hypothesis. The choice is then made by opting for the particle hypothesis that has the greatest likelihood or alternatively by implementing a threshold on the values to differentiate between the candidates. This method improves the probability of accurately identifying the type of particle, using the pion hypothesis as a reference point.

Muons

Since this thesis focusses on W boson decays to muons, the ability of the LHCb experiment to accurately reconstruct muons is of the utmost importance.

Muons are identified by connecting the extrapolated tracks with signals detected in the muon system. The size of the window and the number of stations used in the particle identification process depend on the energy of the particle, since a minimum momentum of 3 GeV is required to reach the M2 and M3 stations, and 6 GeV to reach the five stations. The χ^2_{2D} metric that measures the proximity between the extrapolated track and any associated electromagnetic calorimeter cluster (energy deposits) is minimised, similar to the ECAL cluster–track matching.

A Boolean variable can be used to choose a group of tracks with a high concentration of muons and low contamination; **ISMUON** is set to true if the track is close to hits in the muon chambers at LHCb (a search window is created based on the track’s kinematics) and false otherwise. If the particle has a momentum of more than 10 GeV, then the hits must be found in M2, M3, M4 and M5 to meet the **ISMUON** criterion.

3.6.3 Offline data processing

The offline data processing procedure begins with “stripping”, which can be regarded as an offline trigger system. Initially, the recorded events are filtered to remove those energies, momenta, or particle types that do not align with the study’s goals. The retained samples are then sorted into streams on the basis of their potential relevance to physics studies. During the stripping process, the most recent alignments and calibrations are used to completely reconstruct the recorded events. This provides precise tracking, vertex, and PID data, which are then used to explore different independent decay channels. The reconstruction of composite particles and the application of lenient selection criteria are specified in dedicated stripping lines. This initial selection step serves to reduce the dataset to a manageable size while still allowing the selection criteria to be adapted for a variety of analyses.

There are currently 3 offline pre–selection (stripping) streams [71]:

- **Full stream:** saves raw detector read–out and trigger summary reports while removing trigger reconstruction.
- **Turbo stream:** preserves trigger summary reports and reconstructed candidates for each trigger line, while omitting the raw detector readout.
- **TurCAL stream:** retains both the aforementioned data components.

The Full stream is the usual way of storing data for a high–energy physics experiment, where it is not possible to accurately reconstruct the event in real time, so the online reconstruction of the trigger is replaced by a more advanced offline reconstruction. This was the case for LHCb in Run 1, but by Run 2, the offline and HLT2 reconstructions were of the same quality, allowing the Turbo persistency model to be used. By using the Turbo stream instead of the Full stream, it is possible to reduce the data by a factor of ten. The TurCAL stream was used to calculate

efficiencies and corrections for both online and offline reconstructions [72]. This represented a significant enhancement in efficiency in terms of data collection at LHCb.

After undergoing the full stream and offline reconstruction, thousands of additional stripping lines were run [73]. These lines usually made exclusive selections, and the most important for this thesis were the **WMu** and **Z02MuMu** stripping lines: the former selects a single muon with $p_T > 15$ GeV; and the latter selects an oppositely charged dimuon pair with an invariant mass of $m_{\mu\mu} > 40$ GeV and each muon has $p_T > 15$ GeV.

3.7 Simulation

Monte Carlo simulation and studies play an essential role in the study of signal modes, efficiencies, detector response, and background studies in LHCb analyses. They are used ubiquitously with all LHCb analyses, but are especially important for “template fitting” analyses such as the one covered in this thesis. In template fitting, simulations are used to generate templates—essentially reference distributions—that represent expected signal and background distributions. These templates are then compared with experimental data to extract physical quantities or parameters of interest. The accuracy of template fitting relies heavily on the quality and representativeness of these simulated templates.

The simulation process starts with GAUSS [74], a software application, which then directs an event generation step with PYTHIA [44]. Subsequently, any unstable particles are simulated to decay with EVTGEN [75]. The QED final-state radiation is added using PHOTOS [45].

We aim to ensure that the simulation samples are treated in a similar fashion to the data, so that they can reproduce the same inaccuracies of the data stream (which is a fundamental necessity for providing unbiased efficiency estimations).

To this end, simulation of interactions with the detector material is carried out using GEANT4 [76, 77]. Subsequently, BOOLE [78] is used to simulate the response of the front-end electronics and digitisation. This allows the simulated data and the real data to be treated in the same manner for further processing steps.

3.8 Upgrades

The LHCb detector described in this chapter ended its data-taking operations, at the end of 2018 to begin the period known as Long Shutdown 2 (LS2). In LS2, LHCb went through a significant enhancement in anticipation of a new phase of data collection with an instantaneous luminosity five times higher. In order to handle the increased luminosity, greater granularity and radiation resilience were necessary, leading to the replacement of more than 90% of LHCb’s operational electronics [79]. Moreover, the design of the LHCb system for Run 1–2 did not allow for a significant improvement in statistical accuracy, especially for fully hadronic final-state decays. A major limitation came from the maximum output rate of the first trigger level (L0) [80], where stringent transverse momentum thresholds created a bottleneck for numerous

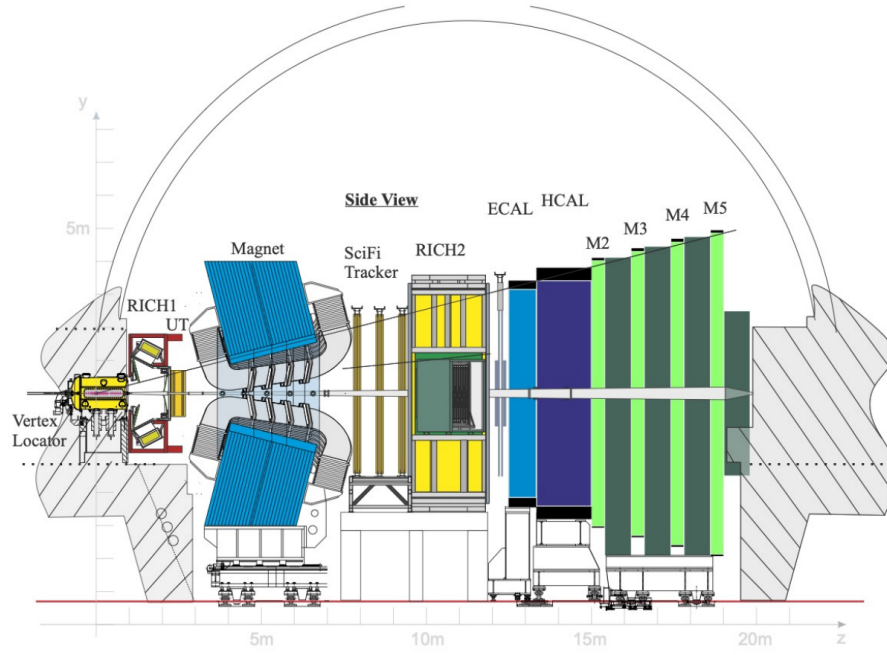


FIGURE 3.22: Diagram of the new LHCb experiment layout for Upgrade I. Reproduced from Ref. [79]. The diagram shows the different subsystems that make up the detector, as well as the use of Cartesian coordinates, with the x -axis extending into the plane of the page.

studies. This necessitated a move towards a software-only trigger system, which the author has contributed to, and the contributions are highlighted in Chapter 4.

The LHCb experiment expects Run 3 to last from 2022 to 2025, with a possible extension to 2026. This section reviews the enhancements made to the LHCb detector to support Run 3, highlighting the most recent improvements and differences from previous runs. Since these upgrades are not directly relevant to the thesis content (since we use data from Run 2), they will be briefly summarised hereafter. The upgrade required for Run 3 is known as LHCb Upgrade I [79].

3.8.1 Upgrade I and Run 3

A diagram showing the upgraded LHCb experiment with the developments of Upgrade I is shown in Figure 3.22.

Given the significant modifications made to the experiment and setup, I will discuss some of the main updates, but definitely not all of them. Most of the descriptions provided here come from Ref. [79].

The upgrade of the LHCb tracking system begins with the upgrade of the VELO, which is now a hybrid silicon pixel detector composed of 52 modules with a radiation tolerance of up to $10^{16} \text{ n}_{\text{eq}}/\text{cm}^2$.

Further downstream, the silicon–strip upstream tracker (UT) is in front of the large–aperture dipole magnet. It consists of 4 planes made of silicon strips, with finer segmentation and better acceptance compared to the older TT. It is mainly used to track charged particles and is a crucial element of the initial processing algorithm within the software trigger. It offers fast p_T determination for track extrapolation, reducing ghost tracks.

Subsequently, three scintillating fibre tracker stations (SciFi) are located downstream to cope with high occupancy, replacing the older straw–tube OT and silicon–strip IT. It is responsible for the tracking of charged particles and momentum measurement. It consists of 128 modules arranged in 12 planes that are cooled to -40°C using a mono–phase liquid cooling system. It has a hit efficiency of better than 99%.

Upgrades were made to the particle identification systems, including improving RICH1 with new mirrors that have a longer focal length to reduce occupancy by half. Furthermore, both RICH1 and RICH2 were equipped with new photon detector chains to support higher readout rates. The previous hybrid photon detectors (HPDs) were substituted with multi–anode photomultiplier tubes (MaPMTs), leading to improved granularity.

The Scintillating Pad Detector (SPD) and the PreShower (PS), which were part of the previous calorimeter system, have been eliminated due to their diminished significance in the new LHCb all–software trigger.

The layout of both the electromagnetic calorimeter (ECAL) and the hadronic calorimeter (HCAL) remains unchanged for the upgrade. However, the front–end electronics boards have been fully redesigned to comply with the 40 MHz readout frequency; in fact, the enhancements made to all subsystems were designed to conform to the 40 MHz readout structure and address the challenges associated with increased luminosity.

In addition, Upgrade I introduced two new subsystems: PLUME and SMOG2. The Probe for Luminosity MEasurement (PLUME) is a novel luminometer designed for both online and offline luminosity measurements. It operates by detecting Cherenkov light generated when charged particles traverse a quartz radiator, which is then captured by a series of photomultiplier tubes. The quantity of photons detected corresponds to the number of charged particles passing through the quartz radiator, which is directly related to the quantity of proton–proton interactions. Positioned 477 m from the interaction point, the PLUME detector is located within the LHCb cavern and integrated into the beam pipe.

The System for measuring the overlap with gas (SMOG) [81], initially designed and applied to accurately calibrate the luminosity of colliding beams, was employed to introduce a noble light gas into the VELO vacuum chamber. This system was used to assess the alignment of the proton and antiproton beams at the collision point and to determine the luminosity of the LHCb experiment. It also offered a method for assessing the interactions between the beam and gas, which were used to determine the beam characteristics and oversee the beam’s consistency. Moreover, SMOG presented a distinctive opportunity to run the LHCb experiment in fixed target mode. Gaseous targets of varying nuclear sizes (helium, neon, and argon) were used along with proton and lead beams at centre–of–mass energies of up to 115 GeV (nucleon–nucleon equivalent), with

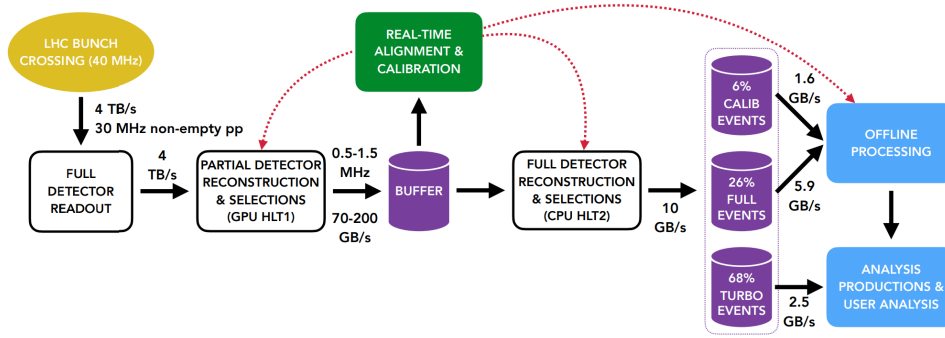


FIGURE 3.23: An overview of the LHCb trigger system in Run 3. Reproduced from Ref. [79].

minimal impact on LHC operations. Based on promising initial findings and future potential, an enhancement of SMOG (referred to as SMOG2) was suggested and implemented. The gas density was increased by up to two orders of magnitude, resulting in significantly higher luminosity. Gas targets may consist of helium, neon, argon, and potentially hydrogen, deuterium, nitrogen, krypton, and xenon.

A recently constructed modular data centre has been established above the experimental site to house all the necessary computing resources for the enhanced readout system and event-filter farm. The data centre is equipped with a high-speed network connection to the CERN data center, which is essential for efficient data transfer to the offline storage system.

Trigger and Real-Time Analysis

As mentioned above, all sub-detectors are read out at 40 MHz. Therefore, the plan for Run 3 was to eliminate the limiting hardware trigger (L0) and move to a fully software trigger with “Real-Time Analysis” capabilities, where “real-time” means that the trigger decision is made within a few milliseconds of the collision. This is a significant change from the previous run and required a lot of work to ensure that the software was ready for this change.

The objective of the trigger system is to reduce the data volume, which reaches 4 TB/s at the nominal instantaneous luminosity in pp collisions, to around 10 GB/s, which can be recorded to permanent offline storage. The software trigger is responsible for reducing the initial input of 30 million collisions per second to an output of approximately 1–2 MHz. This is achieved by swiftly reconstructing tracks and choosing proton–proton collision events using one- and two-track entities solely on GPUs. Figure 3.23 shows a flow diagram of the trigger system in Run 3.

Upgrade II

Moving to the era of High-Luminosity LHC, the LHCb experiment will move to Upgrade II to try to match the high expectations given the increased luminosity and hence the increased data

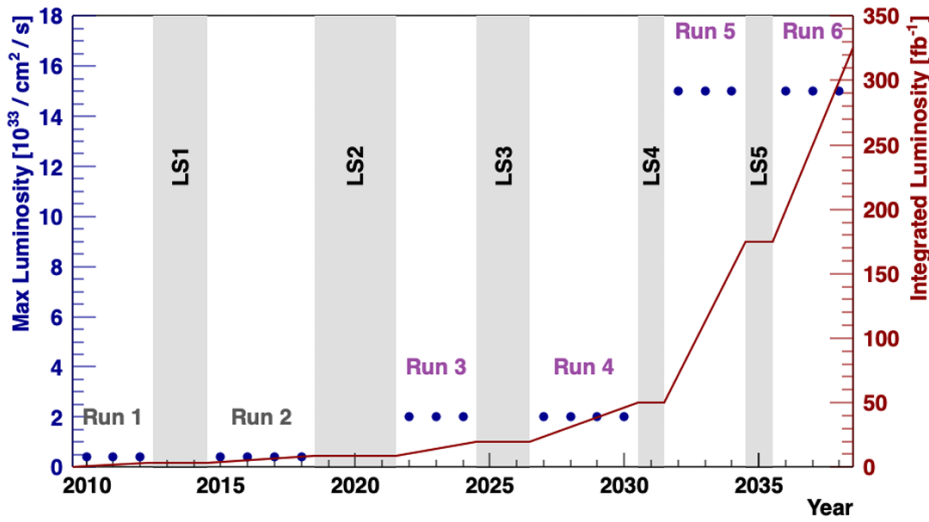


FIGURE 3.24: Integrated luminosity profile for LHCb spanning different runs. Reproduced from Ref. [82].

and radiation [82]. Figure 3.24 demonstrates the integrated luminosity profile for LHCb over the years.

In the course of Long Shutdown 4 (LS4), Upgrade II will be carried out to prepare for Runs 5 and 6. Important preparatory tasks are also planned during Long Shutdown 3 (LS3), which will also yield advantages for the physics performance in Run 4. Upgrade II will enable the experiment to collect data equivalent to at least 300 fb^{-1} and will allow many observables to be measured with a level of precision that cannot be achieved in competing experiments. The plan is to have a general-purpose and high-luminosity flavour experiment.

The main challenges of this new experiment will be the handling of increased pile-up and radiation. Additionally, a significant element missing from the previous detector version was accurate timing. This aspect will be addressed, for example, in the new version of the VELO that incorporates 4D tracking and precise timing. Another interesting addition is the TORCH detector, which will offer high-precision time-of-flight measurements to help improve the detection of low-momentum particles that are difficult to identify with conventional Cherenkov detectors (e.g., RICH).

By evaluating the time difference between particle arrivals and a reference clock, TORCH improves particle identification capabilities and complements existing detectors, providing deeper insights into particle physics phenomena, particularly the identification of kaons and pions. Figure 3.25 shows the improved LHCb experiment with the projected developments of Upgrade II. With extensive preliminary research ongoing, the future of LHCb's particle physics research promises an exciting challenge and a promising outlook. Upgrade II is expected to be up and running in 2029, and the experiment will continue data collection until 2039. The LHCb experiment will continue to offer valuable insights into particle physics, with upgrades that allow the exploration of new frontiers in the field.

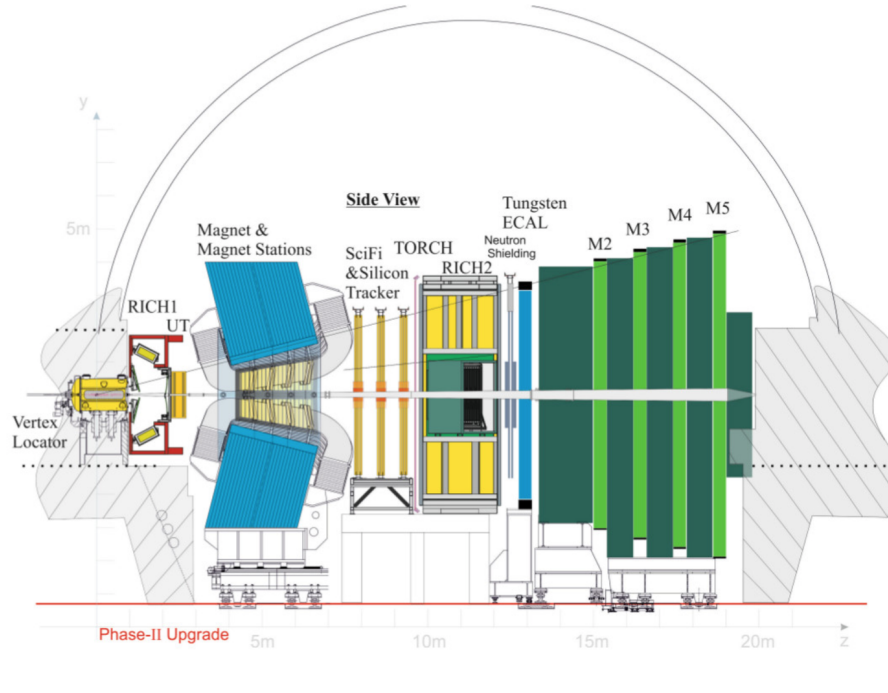


FIGURE 3.25: Diagram of the new LHCb experiment layout for Upgrade II. Reproduced from Ref. [82]. The diagram shows the different subsystems that make up the detector, as well as the use of Cartesian coordinates, with the x -axis extending into the plane of the page.

Implications and challenges

How will these new upgrades enhance the measurements? Essentially, the primary advantage is the ability to use a significantly larger amount of data, which allows us to achieve more precise and reliable measurements. This is particularly valuable for precision measurements, where even small improvements in data quality, such as improved event reconstruction, reduced background noise, and increased particle identification efficiency, can lead to more accurate results.

However, these advancements come with their own set of challenges. For instance, the increased data volume can lead to more pile-up events, where multiple interactions overlap, making it more difficult to disentangle individual events. Additionally, elevated radiation levels associated with higher data rates can affect detector performance and data integrity.

Moreover, the sheer volume of data requires substantial increases in storage capacity, which can be both costly and complex to manage. This demand for storage requires investment not only in physical hardware, but also in robust data management and processing systems to handle the increased load.

To address these challenges, our data acquisition systems and algorithms need to be optimised. This includes enhancing data-recording efficiency, improving algorithms to better handle pile-up, and developing strategies to mitigate the impact of radiation on detector performance. Through these optimisations, we can ensure that the benefits of the upgraded systems are fully realised, leading to more accurate measurements and a better understanding of the underlying physics.

Chapter 4

Contributions to the Real–Time Analysis project at LHCb

“If I have seen further it is by standing on the shoulders of giants.”

– Isaac Newton

This chapter starts with an overview of the Real–Time Analysis (RTA) project at the LHCb experiment in Chapter 4.1. Following that, Section 4.2 reviews the primary software, structured as hierarchical projects. Lastly, Section 4.3 examines the various contributions made by the author to the RTA project.

4.1 Introduction

As a reminder of what was mentioned in Section 3.5, a rapid event selection mechanism, known as a trigger, is essential for examining the data derived from proton–proton (pp) collisions at the LHC. This mechanism is commonly divided into two stages. The initial stage, known as the “online” phase, encompasses the data recording by the detector, its rapid transmission through electronic devices, and its rapid analysis by computer algorithms. Subsequently, a portion of the events is preserved on storage devices such as disks and magnetic tapes. The subsequent stage, termed “offline analysis”, involves establishing condition parameters, such as alignment and calibration, which are contingent on the data collection period. The complete dataset is then re–analysed using these updated parameters before the final data scrutiny is conducted. This process is time–consuming and may result in disparities between entities in the online and offline analyses.

To address these issues, the LHCb Collaboration has introduced a significant advancement to the data collection and analysis process, known as “real–time analysis” (RTA) [71]. This marked a new way of handling the final stages of data processing; where they would be carried out in “real–time”, with the trigger system automatically handling calibration and alignment. Thus, we only ever do one alignment and calibration, so we do not need to reprocess the events. Subsequently, the data becomes promptly available for offline scientific investigations. The term “real–time”

refers to the duration from the moment a collision occurs to that when the corresponding event must be discarded or transferred to offline storage for long-term archival.

During LHCb Run 2 (2015–2018), the triggering mechanism involved a combination of rapid electronic systems (referred to as the “hardware trigger”) and computational algorithms (referred to as the “software trigger”).

In preparation for Run 3 (2022–2025) and subsequent runs, the LHCb detector will function at a luminosity five times higher than that of Run 1 and Run 2 [71]. Consequently, a comprehensive overhaul of the trigger system was necessary. In Run 2, LHCb had to cope with an increase in both luminosity and event rates, making traditional offline analysis methods inefficient. After testing the RTA approach during Run 2 as a way to overcome these issues, the collaboration created the Real-time Analysis project during the long shutdown between Run 2 and Run 3. The hardware trigger was removed, and the complete detector is now read at the full LHC bunch-crossing rate of 40 MHz. This adjustment enables the entire selection process to be carried out using software, in accordance with the real-time analysis (RTA) methodology [71].

The foundation of the current (Run 3) LHCb trigger system is its data acquisition setup, consisting of approximately 160 computer servers equipped with 480 specialised electronic cards. This arrangement is responsible for collecting and merging data from nearly 1 million individual electronic channels across the detector. This data flow amounts to a rate of 5 TB/s, which is subsequently directed to the initial phase of the trigger, known as HLT1 [83].

In HLT1, data from numerous events are transmitted to Graphics Processing Units (GPUs) housed on the same servers, where they consolidate the detector data, thus decreasing the volume of data to be sent over the network. Each GPU can handle approximately 150,000 events per second, while retaining complete detector information for 3% of the most notable events for later analysis. The remaining data amounts to 120 GB/s. The chosen events are then stored in a 40 PB disk buffer, enabling real-time alignment of subdetector positions with micrometre accuracy, as well as calibration of the calorimeter using π mesons. Furthermore, the detector’s capability to distinguish various types of particles is fine-tuned with per mille precision. This real-time alignment and calibration guarantee that LHCb achieves maximum sensitivity for the numerous high-precision measurements that are fundamental to the collaboration’s physics program. The presence of the disk buffer also enables the execution of the second trigger stage (HLT2) even when the LHC is not actively colliding protons, such as during downtime or between different running periods, thus optimising the use of LHCb computational resources [71, 83].

After ensuring the detector’s optimal performance, the alignment and calibration constants are revised and applied in ensuing events by transferring the constants to HLT, which functions on about 3700 computer servers. These servers perform the most precise reconstruction of the entire detector and apply about 1500 different criteria (lines) to select the most intriguing events to further decrease the data volume by a factor of 10. The exceptional quality of LHCb’s real-time reconstruction enables these criteria not only to identify interesting events but also to transform the raw detector data into high-level physics entities instantly. This entirely software-driven process provides LHCb with significant flexibility to choose both the most intriguing events

and the most critical aspects of each event, thus maximising the utilisation of its computational resources throughout its entire physics program. Ultimately, approximately 10 GB of the data are continuously recorded every second and are accessible to physics analysts for offline analysis [83].

4.2 Core software projects

The LHCb core software [84] is organised into hierarchical projects hosted on GitLab, with numerous interconnected projects that span the LHCb software domain. A project comprises various scripts designed to perform particular tasks related to a common theme. Every project oversees a specific segment of LHCb software. Key examples include MOORE [85] for real-time data processing, ALLEN [86] for GPU-based high-performance computing, DAVINCI [87] for data analysis, DETECTOR for managing detector tasks, and MONET [88] for monitoring operations. This structured approach ensures that each project is responsible for a unique and essential component of the LHCb software framework.

The link to all LHCb projects can be found on the LHCb GitLab page [89].

GitLab [90] is used to oversee and manage those software projects; dividing them into different repositories. One of GitLab's key functionalities is the ability to navigate through the source code of individual projects and assess suggested modifications (referred to as merge requests) through a user-friendly web interface.

4.2.1 Moore

The MOORE project [73] has been used since Run 2 to configure and test the LHCb High-Level Trigger application known as MOORE, which is responsible for reducing the input rate of 40 million collisions per second to an output rate of approximately 100 kHz. This process occurs in two stages:

1. HLT1 conducts a rapid track reconstruction and reaches a verdict based on single- and two-track objects.
2. HLT2 performs a detailed reconstruction and reaches a verdict based on the complete detector read-out information.

4.2.2 Allen

The ALLEN project [91] refers to the LHCb HLT1 software that runs on graphics processing units (GPUs). Its main task is to reduce an initial input of 30 million collisions per second to an output of approximately 1–2 MHz. This is achieved by swiftly reconstructing tracks and choosing proton–proton collision events using one- and two-track entities solely on GPUs.

The project is called Allen, based on Frances E. Allen, in accordance with the LHCb tradition of naming software initiatives after distinguished scientists [91].

4.2.3 DaVinci

DAVINCI [73] serves as the offline analysis tool for LHCb, allowing users to generate tuples that contain essential information on reconstructed particles from a specific decay. It is the gateway to accessing the LHCb data. The primary function of DAVINCI is to create tuples. It can be used in Analysis Productions, for submitting custom productions using GANGA, or for executing small tasks on “cvmfs” systems such as “lxplus” at CERN [92, 93]. cvmfs is a global, distributed file system that allows the distribution of software to thousands of nodes around the world. It is used by the LHCb experiment to distribute software to the grid and to the CERN computing facilities. lxplus is a service provided by CERN that allows users to access the CERN computing facilities remotely. It is a Linux server that provides access to the CERN software environment and the CERN computing grid. Moreover, by providing access to detailed information on the reconstructed particles, DAVINCI facilitates more in-depth studies on the Turbo output. The different ways in which we handle offline data are reviewed in Section 3.6. This resource outlines the various components of DAVINCI, which essentially consists of a collection of Python packages that configure algorithms, tools, data flow, and control flow to execute a GAUDI-based application with complete access to Turbo output [73, 92].

4.2.4 Detector

The DETECTOR project encompasses the description of the LHCb geometry presented in DD4hep Compact XML format [94, 95]. DD4hep, a toolkit tailored for high-energy physics, addresses various challenges encountered throughout all phases of an experiment. Drawing insights from the LHCb experiment and advancements in the Linear Collider community, this toolkit leverages established software to create a unified and comprehensive detector description. Key elements include the ROOT geometry package for geometry construction and visualisation, and the GEANT4 simulation toolkit, which can be connected through DD4hep to perform detector simulations in intricate designs [96].

The motivation behind this toolkit is to provide a thorough and centralised description of the detector that covers all stages of the experiment. It includes in-depth details about the detector’s geometry (its physical structure and arrangement), the materials used in its construction, visualisation properties that aid in comprehension and representation, specifications for detector readout (detailing how data are collected), as well as alignment, calibration, and environmental considerations to ensure accurate data interpretation. This toolkit is designed to support all phases of the process, from the initial design and optimisation stages to construction and operation, thereby facilitating seamless transitions between these stages. In addition, it functions as a central information repository, guaranteeing a consistent depiction of the detector for simulation, reconstruction, and analysis purposes throughout the experiment [94, 95, 96].

4.2.5 Monet

MONET, a web-based application, was developed to supervise the quality of the data obtained from the LHCb experiment [88]. The capability to supervise current operations and oversee data

gathering in real-time is essential. Promptly identifying any decline in Data Quality or potential issues is critical. It serves as a significant tool for people working in the LHCb control room.

4.3 Contributions to the Real-Time Analysis group and LHCb

4.3.1 RTA Software Shifter

As noted in the previous section, the integration of numerous GitLab projects is essential for the smooth operation of the LHCb software. Therefore, it is crucial to ensure that the software undergoes thorough testing and validation prior to its deployment for production purposes. This involves executing the software on a designated set of test data and cross-referencing the results with the anticipated results. The software maintainer is tasked with certifying the testing and validation of the software before its implementation, overseeing the software's release procedures, and ensuring its continuous maintenance and support. Assisting the software maintainer is the software shifter, who is responsible for reviewing various merge requests, confirming the rationale behind the modifications, and ensuring proper documentation. The Software Shifter also sends notifications to developers who are falling behind to encourage them to complete their tasks.

The software is tested and validated through the LHCb nightly builds [97], which are executed daily. The Software Shifter can execute tests on merge requests within the LHCb nightly builds. Given the interconnected nature of the projects, it is crucial to verify that the modifications do not disrupt the dependencies. Occasionally, it is necessary to perform tests on several merge requests across various projects to confirm their compatibility.

The LHCb nightly builds are run on multiple operating system platforms and compiler versions. The Software Shifter oversees the validation of tests across diverse platforms and architectures in the nightly builds. They are tasked with verifying that the tests are carried out using different compilers and CUDA versions. This validation process is crucial to guarantee the compatibility of the software with a range of platforms and compilers [93].

In case of test failures, the Software Shifter is tasked with pinpointing the issue and notifying the Merge Request assignee, possibly providing recommendations for resolution. On the other hand, if the tests are successful, the Software Shifter can alert the Software Maintainer to assess the merge request for final approval and merging if all is well. Effective collaboration and communication between the Software Maintainer and the Software Shifter are essential to guarantee thorough testing and validation of the software prior to its deployment in production.

Another crucial responsibility of the Software Shifter is to ensure that the software's physics performance remains intact after modifications. This involves conducting tests on the LHCb nightly builds and contrasting the results with the anticipated results. For example, one can verify that the updated results do not introduce new bugs or lead to a decrease in performance or throughput.

Another important task I performed was the monitoring of various unique distributions and histograms that are produced by the LHCb software. These distributions are crucial for the

LHCb physics analyses and are used to monitor the performance of the software. The Software Shifter is responsible for ensuring that these distributions are produced correctly and that they are consistent with the anticipated results. In case of discrepancies, the Software Shifter must identify the issue and notify the relevant parties for resolution.

At the end of my two-week shift, I successfully assisted in merging over 40 merge requests, witnessing the implementation of the changes I had reviewed in the production environment. Summarising the tasks completed during my shift is crucial, as well as making sure that the next Software Shifter is informed about the updates I have reviewed and integrated. Equally important is to inform them about any issues detected and reported, along with the feedback received. Normally, the shift report is documented on an issue within the MOORE project, serving as a means to convey the completed tasks and the pending work that still needs attention. I was able to contribute to the resolution of various issues and bugs that plagued the LHCb software, and we were able to successfully merge the changes into the production environment, leaving our software in a much better state than when we started.

4.3.2 RTA HLT Piquet

The RTA HLT Piquet has a vital role that includes three primary functions. Firstly, they are responsible for verifying the proper functioning of HLT1 and HLT2. Secondly, they must handle any emerging issues autonomously or by consulting specialists. Thirdly, they are tasked with maintaining efficient communication with the shift leader, experts, and other relevant parties.

This position plays a critical role in the efficient operation of the LHCb experiment by overseeing the proper operation of the HLT and resolving any emerging issues. The Piquet serves as the primary contact for any concerns related to HLT and is tasked with maintaining clear communication with the shift leader, experts, and other relevant individuals.

One of my first assignments included updating the LHCb software stack to incorporate recent enhancements and bug fixes. I noticed that one of the newly implemented software stacks was not used consistently during operations with the RICH detector being activated and included, leading to failures in some ALLEN data collection nodes (the RICH was not part of the data collection process previously). This meant that we were not collecting data from the RICH detector, which was a significant issue that needed to be addressed quickly, so I reached out to experts and specialists to address the issue. We discovered that the problem was with respect to the architecture that was being used involving the use of two GPUs simultaneously for data collection. Collaborating with experts, I was able to address the problem by temporarily switching back to a setup that used only one GPU, which we had setup before. This fixed the problem, allowing us to include the RICH detector in the data collection process once again. Notably, this adjustment also did not affect the speed of data collection, which is a crucial factor in data collection efficiency. The fix served as a temporary measure until the experts had more involved discussions and implemented a more permanent solution that was in line with the new software stack as well as future data collection plans. I was able to document the issue and the steps taken to address it in the shift logbook, ensuring that the information was available for future reference.

and follow-up. I also followed up with the expert discussion about the future plans regarding the more permanent solution and the timeline for its implementation.

4.3.3 Removing Global Event Cuts from HLT1 by default

To protect the reconstruction strategy (Section 3.6) against saturated events, it is necessary to begin by removing events with too high an occupancy in the Vertex Locator (VELO), Outer Tracker (OT) and Inner Tracker (IT). For a recap of the different detectors in LHCb, see Section 3.2. In the case of OT and IT, events are rejected if they contain more than 10,000 and 3,000 hits, respectively, corresponding to an occupancy of about 20%. In the case of VELO, events are rejected if they contain more than 3,000 hits [98].

As mentioned in Section 3.5, in Run 2, most of the L0 trigger lines were equipped with a “global event cut” (GEC) based on hits in the SPD subdetector, which was aimed at discarding events that required significant reconstruction time. Typically, the GEC was set to 1,000 hits in the SPD. In Run 3, the GEC will be determined by the count of clusters in the updated tracking stations before and after the magnet, as the SPD is no longer in use.

Although GECs have minimal impact on most LHCb studies, they are critical for electroweak (EW) analyses, especially in cross-section measurements, due to their significant and intricate influence on the trigger efficiency of high p_T EW events [99]. If not handled correctly, the GEC can lead to a significant reduction in the trigger efficiency of high p_T EW events, on the order of 10–20% [99].

Multiple factors influence the effectiveness of detecting a specific decay channel: acceptance, triggering, reconstruction, and selection efficiencies. The particles within the event under consideration must initially fall within the detector’s acceptance range, then be triggered, reconstructed, and ultimately meet the offline selection criteria.

We have three trigger categories [100]:

1. Triggered On Signal (**TOS**): Events in which the signal’s presence alone is enough to prompt a positive trigger decision.
2. Triggered Independent of Signal (**TIS**): The remaining part of the event is adequate to prompt a positive trigger decision, with the remaining part being determined by a set procedure that involves eliminating the signal and all associated detector hits.
3. Triggered On Both (**TOB**): These events do not fall under the categories of **TIS** or **TOS**. The signal’s presence by itself or the rest of the event alone cannot lead to a positive trigger decision; instead, both elements are required.

We contrast the efficiencies computed using GECs with **TIS–TOS** methods, which distinguishes between events where the trigger signature is fully matched to an offline selection (**TOS**) and events where the trigger signature is only partially matched or not matched at all (**TIS**). The results are shown in Table 4.1.

Trigger efficiency is defined as

$$\epsilon(DEC) = \frac{N_{\text{triggered}}}{N_{\text{pre-selected}}}, \quad (4.1)$$

where $N_{\text{pre-selected}}$ represents the count of simulated signal candidates that satisfy a pre-selection threshold, also known as the “denominator requirement”. This is occasionally denoted as the “decision” efficiency (DEC). The standard denominator is known as **CanRecoChildren**, which mandates that all decay’s end-state children are charged and can be reconstructed as long tracks (with hits at the VELO and downstream T stations) within the pseudorapidity range $2 < \eta < 5$, a criterion that is pertinent for most decays.

Decay mode	GEC	TIS-OR-TOS	TOS
$B^0 \rightarrow K^{*0} \mu^+ \mu^-$	88.9 ± 2.0	90.6 ± 2.0	88.8 ± 2.1
$B^0 \rightarrow K^{*0} e^+ e^-$	84.2 ± 2.7	69.1 ± 3.8	61.7 ± 4.0
$B_s^0 \rightarrow \phi \phi$	83.2 ± 2.6	75.8 ± 3.2	68.5 ± 3.5
$D_s^+ \rightarrow K^+ K^- \pi^+$	82.5 ± 3.6	58.5 ± 5.1	42.6 ± 5.1
$Z \rightarrow \mu^+ \mu^-$	77.8 ± 1.2	99.5 ± 0.2	99.5 ± 0.2

TABLE 4.1: The effectiveness of the complete HLT1 selection process is evaluated. Efficiencies for **TIS-OR-TOS** and **TOS** are calculated using events that meet the GEC criteria. All efficiencies and their corresponding uncertainties are presented as percentages. Signal events are selected based on the following conditions: b and c hadrons are required to possess a transverse momentum $p_T > 2 \text{ GeV}$ and a lifetime $\tau > 0.2 \text{ ps}$. Children of b and c hadrons must have a transverse momentum $p_T > 200 \text{ MeV}$. Children of Z bosons must have a transverse momentum $p_T > 20 \text{ GeV}$. It is evident that the $Z \rightarrow \mu^+ \mu^-$ decay is significantly affected by GECs, and removing them provides considerable advantages. Nonetheless, it is also apparent why GECs might sometimes play a crucial role when it comes to the decays of B and D mesons. Taken from Ref. [91].

The impact of using the GEC on the efficiency of the $Z \rightarrow \mu^+ \mu^-$ process is evident when compared to solely employing **TIS-OR-TOS**. In the past, there were intentions to eliminate GEC for certain lines, such as **Hlt1SingleHighPtMuonDecision**

and **Z0_Hlt1SingleHighPtMuonDecisionTrueSim**. I managed these modifications within the Allen project, where I was tasked with modifying the GECs setup to ignore these lines and only be applied to the remaining lines. This was a valid approach, as the GECs were not necessary for these lines and removing them would have a significant impact on improving the efficiency of the trigger lines. Additionally, there were some changes to relax the GECs and their effects, namely by removing the use of Upstream Tracker (UT) information from the GECs, which I was not involved with. This, of course, helped improve the efficiency of the trigger lines, as the GECs were less restrictive and allowed more events to pass through the trigger selection process.

However, after considering the presence of two GPUs per event building server, it was concluded that there was a sufficient buffer to eliminate GECs as the default option in HLT1 for all trigger lines. Consequently, all HLT1 processes would operate without these cuts being activated at the start of data processing. As a result, the initial strategy was extended and the modifications that I had previously made on Allen were discarded. Nonetheless, this new approach would

undoubtedly be a favourable development, especially for electroweak studies, as the experiment now avoids GECs by default.

I was responsible for implementing these changes across the LHCb software projects, which span multiple repositories on GitLab. The changes had to be compatible with the existing software configuration and in line with other ongoing developments in the LHCb software.

Updating the reference files to align with the new configuration settings was necessary, along with conducting tests on the LHCb nightly builds to verify the compatibility of the changes with various platforms and compilers. As a result of my extensive changes, there were only minor adjustments to the throughput and efficiency of the LHCb online setup. The throughput saw a slight increase, while the efficiency increased specifically for the EW processes. This outcome was anticipated because GECs are used to lower throughput in favour of efficiency.

Efficiency verification was carried out using `HltEfficiencyChecker`, a tool designed to assess the efficiency of the trigger lines and to verify that modifications do not cause a decrease in performance [101]. These parameters are evaluated in simulated events, allowing the calculation to be performed prior to the data collection process. Efficiency tests were performed using the `SingleHighPtMuon` trigger line for two decay processes: $W \rightarrow \mu\nu_\mu$ and $Z \rightarrow \mu^+\mu^-$.

As expected, the efficiency increased for both processes. The results can be examined in Figures 4.1 and 4.2 for the W and Z decays, respectively. The efficiencies observed with GECs seem to be high in the results presented here, as some of the GECs had been relaxed earlier by excluding the use of UT information.

Performance and throughput variations fell within the anticipated range. These assessments are automatically conducted by the GitLab and nightly testing system and are reviewed by the relevant experts.

As previously stated, the configuration of the LHCb software is intricate and consists of various projects. Following initial testing, issues arose in relation to the MOORE project, which resulted from its failure to incorporate recent modifications into the ALLEN project. The MOORE project was updated to incorporate recent modifications in the ALLEN project, specifically, renaming certain sequences, and the tests were reexecuted to verify the compatibility of the changes with the various projects.

The results were also presented to a panel of experts and discussed at the weekly RTA general meeting, and therefore, the changes were eventually approved. The final tests I conducted were successful and the modifications were successfully merged and incorporated into the ALLEN and MOORE projects.

The entire process spanned around 11 months, starting from the initial conversations, involving extensive consultations with experts, presentations to other LHCb specialists, incorporating requested features and adjustments to ensure that my work was aligned with current and future plans, and finally culminating in the implementation and approval of the changes. This highlights the intricate nature of the LHCb software and underscores the importance of comprehensive testing and validation to guarantee the reliability and efficiency of the software. The

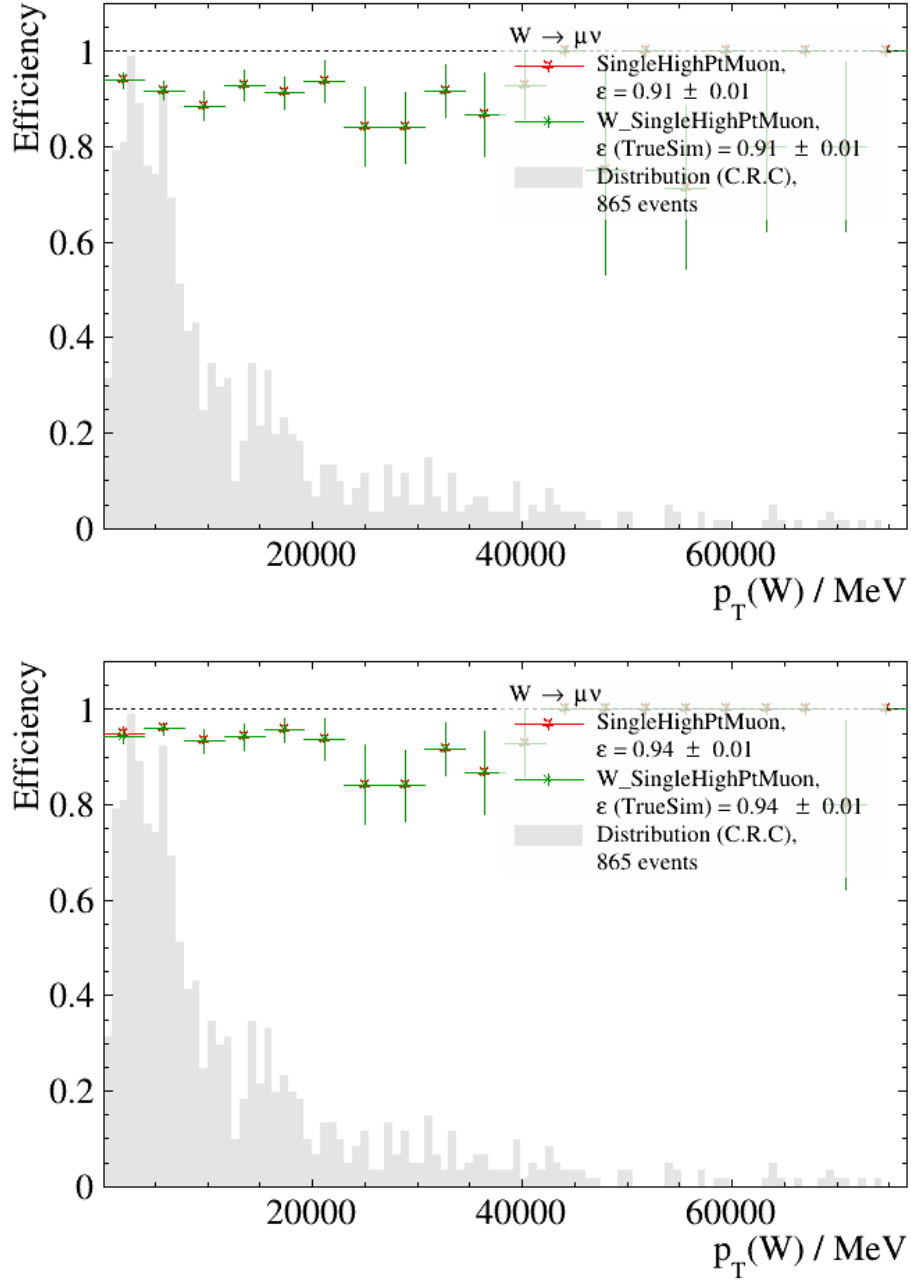


FIGURE 4.1: Trigger efficiencies for simulated $W \rightarrow \mu\nu_\mu$ events (top) before and (bottom) after removing GECs. Produced using **HltEfficiencyChecker**. Efficiency increases from 91% to 94%. C.R.C is the CanRecoChildren denominator distribution.

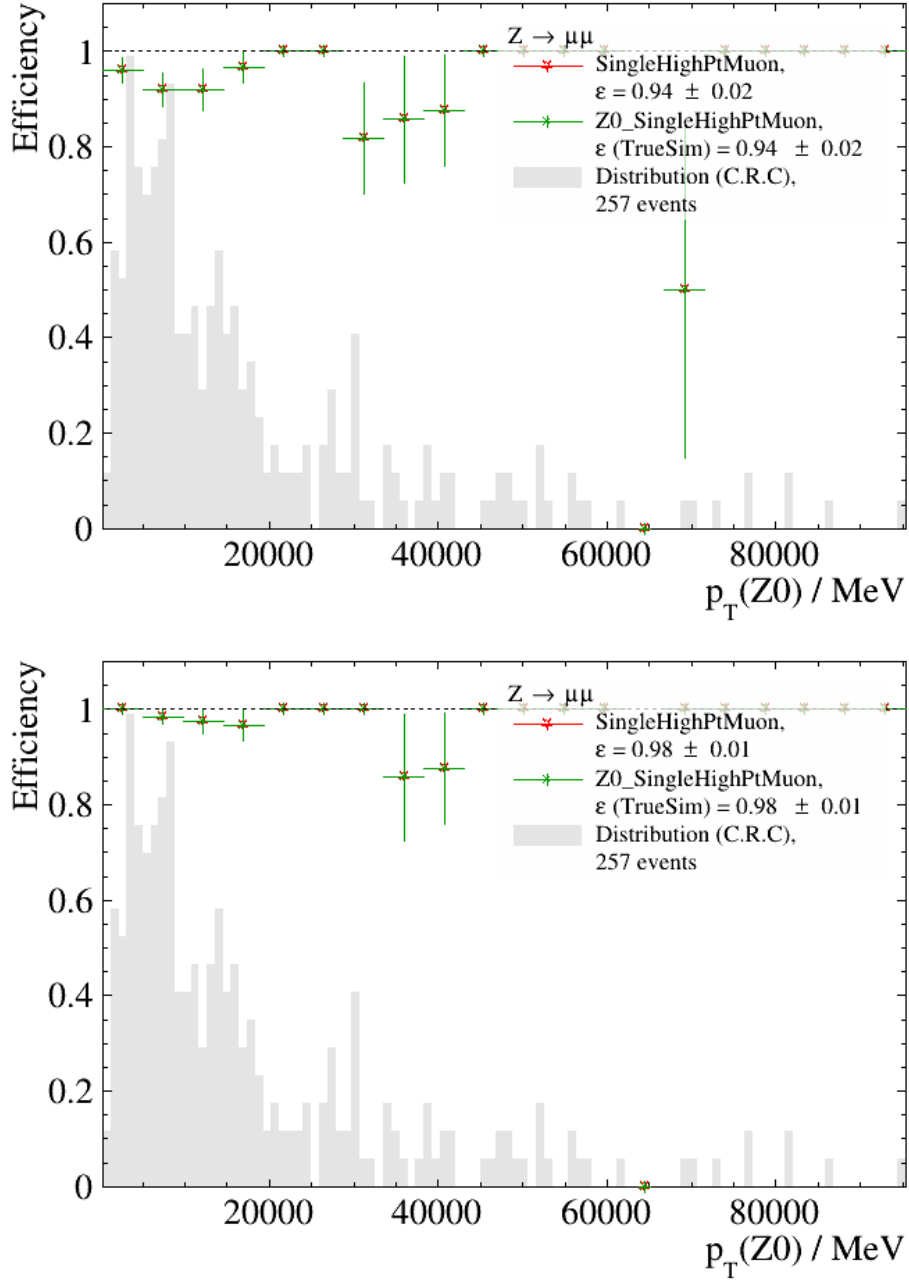


FIGURE 4.2: Trigger efficiencies for simulated $Z \rightarrow \mu^+\mu^-$ events (top) before and (bottom) after removing GECs. Produced using **HltEfficiencyChecker**. Efficiency increases from 94% to 98%. C.R.C is the CanRecoChildren denominator distribution.

merge requests I created underwent numerous revisions and updates, requiring me to adjust the changes to meet new specifications and settings.

Chapter 5

Analysis Setup

“The best that most of us can hope to achieve in physics is simply to misunderstand at a deeper level.”

– Wolfgang Pauli

In this section, we explore the dataset used in our analysis, providing an overview of its sources, main characteristics, and importance to the research. In addition, we outline the corrections applied to the simulated data, which are essential to ensure a precise comparison with experimental results. These corrections encompass adjustments for detector effects, reweighting for theoretical models, and any necessary calibrations to address systematic uncertainties. Corrections follow closely the strategy described in the previous LHCb m_W measurement paper[21]. By considering these aspects, we aim to build a solid foundation for the subsequent analysis, ensuring that our results are dependable and scientifically valid.

It is important to reiterate that since this is a proof-of-principle measurement, we are using a small dataset from the 5.02 TeV Run 2 in 2017, which is statistically limited, and as will be discussed, has some unique challenges.

We introduce the dataset and event selections in Section 5.1, providing a detailed account of the trigger mechanisms, datasets, and simulation samples used in the analysis. We then discuss the momentum calibration applied to the data in Section 5.2, focussing on the curvature bias corrections determined using the pseudomass method. In Section 5.3, we explore the corrections applied to the simulation, including momentum smearing, efficiency corrections, and isolation variable calibrations. Moving on to Section 5.4, we validate the corrected simulation by comparing it with the data, ensuring that it accurately models the experimental results. Finally, in Section 5.5, we provide a comprehensive overview of the software projects involved in the analysis. This section outlines the various components of the analytical framework, identifies the contributors to this analysis, and emphasises the key contributions made by the author.

5.1 Dataset and event selections

5.1.1 Datasets and simulation samples

The purpose of this study is to measure the differential cross-section in bins of muon transverse momentum and the mass of the W boson (m_W) based on the LHCb 2017 5.02 TeV Run 2 dataset, using an integrated luminosity of $99.86 \pm 2.00 \text{ pb}^{-1}$. The uncertainty on the luminosity is fixed to 2% [102]¹. The analysis focusses on the p_T range from 28 to 52 GeV and the pseudorapidity range from 2.2 to 4.4, and the isolation range from 0 to 8 GeV. We selected $W \rightarrow \mu\nu_\mu$ candidates by identifying data events in this kinematic region, resulting in approximately 27,000 $W^+ \rightarrow \mu^+\nu_\mu$ candidates and 21,000 $W^- \rightarrow \mu^-\bar{\nu}_\mu$ candidates. Furthermore, we have approximately 3,500 $Z \rightarrow \mu^+\mu^-$ candidates in the same dataset, which are used to calibrate the simulation samples and to determine efficiency corrections. They are also used to determine the curvature bias corrections to the data with the pseudomass method and to normalise the simulation backgrounds to the data, which is crucial for the background subtraction fit, as it is an extremely clean sample. It is important to mention that we filter our data candidates to retain only the “good runs”, where we have accurate measurements of luminosity, as indicated in the database provided by the LHCb luminosity team in Ref. [102].

We used simulated signal and background data, as well as data files obtained from the LHCb experiment. Our list of backgrounds includes various sources: QCD (e.g., pions and kaons), $Z \rightarrow \mu^+\mu^-$, $W \rightarrow \tau\nu_\mu$, $q\bar{q} \rightarrow t\bar{t}$, $gg \rightarrow t\bar{t}$, single-top, $c\bar{c}$, and $b\bar{b}$ processes. It is important to mention that the $Z \rightarrow \mu^+\mu^-$ process incorporates masses $> 20 \text{ GeV}$; however, the range between 20 and 40 GeV is actually simulated as a separate process. Additionally, we will use Z to refer to both the Z boson and its contributions with virtual photons, instead of explicitly saying Z/γ .

Electroweak backgrounds, such as single-top and $W \rightarrow \tau\nu_\mu$, are generated using simulations and are fixed and scaled to the $Z \rightarrow \mu^+\mu^-$ distribution. The cross-sections are normalised to Z and are constrained by the $Z \rightarrow \mu^+\mu^-$ data to align the simulated backgrounds to the observed data in terms of shape and rate, enabling precise background modelling. This can be summarised in this formula:

$$\text{Scale Factor}_{\text{comp}} = \frac{N_{Z \text{ data}}}{N_{Z \text{ MC}}} \times \frac{N_{Z \text{ MC}}^{\text{gen}}}{N_{\text{comp MC}}^{\text{gen}}} \times \frac{\sigma_{\text{comp}}}{\sigma_Z^{\text{fiducial}}}, \quad (5.1)$$

where $N_{Z \text{ data}}$ is the number of $Z \rightarrow \mu^+\mu^-$ candidates in the data, $N_{Z \text{ MC}}$ is the number of the $Z \rightarrow \mu^+\mu^-$ candidate in the simulation after applying the same selection criteria as the data, $N_{Z \text{ MC}}^{\text{gen}}$ is the number of generated $Z \rightarrow \mu^+\mu^-$ candidates in the simulation, $N_{\text{comp MC}}^{\text{gen}}$ is the number of generated candidates for the other component or process, $\sigma_{\text{comp}}^{\text{fiducial}}$ is the fiducial cross-section of the component, and $\sigma_Z^{\text{fiducial}}$ is the fiducial cross-section of the $Z \rightarrow \mu^+\mu^-$ process. These fiducial cross-sections are obtained from the MCFM program [103] using the NNPDF4.0 PDF set and the fiducial region that is defined as $2.0 < \eta < 4.5$ and $20 \text{ GeV} < p_T$.

Unlike the earlier LHCb m_W measurement [21], which modelled the QCD background using data samples and high- p_T tracks chosen by specialised triggers without incorporating muon

¹Note that this citation is only accessible to LHCb users.

identification, our approach focused on modelling misidentified hadrons through the imposition of specific trigger and muon identification criteria. This allows us to switch to simulated samples of hard-QCD background processes and apply a parameterisation of selection biases as weights to the fit templates. It is important to understand the differences between simulation and real data as well as the impact of charge asymmetries.

The simulation events used in the analysis are shown in Table 5.1. The selection criteria for the signal, QCD background, and $Z \rightarrow \mu^+\mu^-$ candidates are described in Section 5.1.3.

Additionally, samples from the 2017 5.02 TeV Run 2 dataset only have a downward magnet polarity, compared to the 13 TeV Run 2 dataset, which has upward and downward magnet polarities.

TABLE 5.1: Simulation events used in the analysis.

Process	Purpose
$W \rightarrow \mu\nu_\mu$	Signal
Inclusive hadrons, $p_T > 18$ GeV	Background
$Z \rightarrow \mu^+\mu^-, M > 40$ GeV	Background and calibrations
$W \rightarrow \tau\nu_\mu$	Background
$Z \rightarrow \tau^+\tau^-$	Background
$\gamma^* \rightarrow \mu^+\mu^-, 20 < M < 40$ GeV	Background
$q\bar{q} \rightarrow t\bar{t}$	Background
$gg \rightarrow t\bar{t}$	Background
Single top quark	Background
$c\bar{c} \rightarrow \mu X$	Background
$b\bar{b} \rightarrow \mu X$	Background

Quarkonia decays (e.g. $\Upsilon(1S) \rightarrow \mu^+\mu^-$ and $J/\psi \rightarrow \mu^+\mu^-$) are typically used in similar electroweak analyses in LHCb. For our analysis, they could have been used to calibrate the simulation, as they are clean samples of muons. For example, we can use them to calibrate the muon-reconstruction efficiencies and the momentum measurement. However, at present, there are no samples available at 5.02 TeV for $\Upsilon(1S) \rightarrow \mu^+\mu^-$ and $J/\psi \rightarrow \mu^+\mu^-$ as they were not included in the centralised processing of the 5.02 TeV dataset. We note that they could be recovered through a new stripping², but this is not a trivial task and would require significant human and computing resources from the LHCb computing project. For a reminder of the processing of the data in LHCb, consult Section 3.6. Thus, we mainly rely on the $Z \rightarrow \mu^+\mu^-$ process for our calibrations, which suffices for our needs, as will be shown shortly.

5.1.2 Trigger

It is important to require that the events have been selected by a specific pattern of trigger decisions, so that any inefficiencies or biases can be carefully understood, calibrated, and controlled. The trigger system at LHCb is designed to select events that are likely to contain interesting physics processes, such as $W \rightarrow \mu\nu_\mu$ decays. For a detailed explanation of the trigger system in LHCb, consult Section 3.5. In Run 2, a hardware trigger (L0) called **L0MuonEW** was used, which was fired when there was a muon in the event with $p_T > 4.8$ GeV. The first

²See 3.6.3 for a definition of stripping productions.

software trigger (HLT1) utilised the `Hlt1SingleMuonHighPt` line, which filtered events containing a muon with $p_T > 4.8$ GeV. Similarly, the second software trigger (HLT2) employed the `Hlt2SingleMuonHighPt` line to select events featuring a muon with $p_T > 10$ GeV. These triggers were used to select the $W \rightarrow \mu\nu_\mu$ candidates in the 5.02 TeV Run 2 dataset [67].

Additionally, we concluded that the other physics triggers for the 5.02 TeV dataset are too biasing in the selection of a hadron-enriched sample, which is necessary for the QCD background calibration. The use of “no-bias” data is still under investigation, but this is inconclusive. No-bias data are data that are not selected by any trigger and can be used to model the QCD background.

As a consequence, we allowed the QCD background normalisation to float in our fit of the differential cross-section, which will be described in Section 6.3. This approach was deemed acceptable, as the QCD background is not the primary source of uncertainty of the analysis, and normalisation is not expected to have a significant impact on the final results. We also tested various different approaches to this floating normalisation, and all gave consistent results, further supporting the validity of this approach.

5.1.3 Selections

The cuts in the stripping were replaced by the cuts at the analysis level. In this section, we review the selection methods used for the various samples. Our selections are categorised into three groups: signal, charged-hadrons (QCD) background, and Z cuts. With respect to the listed requirements, the high-level trigger and stripping are 100% efficient. For a revision of the meanings of the variables, consult the section on definitions and conventions in Section 1.2. The selection criteria for the different categories are as follows:

Signal selection cuts:

General selection of $W \rightarrow \mu\nu_\mu$ candidates:

- **Muon Kinematics:** The pseudorapidity range $2.2 < \eta < 4.4$ is chosen to match the coverage of the LHCb detector, ensuring that the selected muons are within the optimal acceptance of the detector. The transverse momentum range $28 < p_T < 52$ GeV is selected to focus on high- p_T muons, which are essential to reduce background contamination.
- **Impact Parameter:** An impact parameter $\chi^2 < 9$ is required to ensure that the selected muons originate from the primary vertex, minimizing contributions from secondary decays and improving the signal purity.
- **Energy Deposit:** The energy deposit in the hadronic calorimeter (Hcal) is required to be less than 5.0 GeV. This cut is applied to reduce contamination from hadrons misidentified as muons, ensuring that the selected muons are well-identified.
- **Relative Momentum Error:** A relative momentum error of less than 0.06 is imposed to ensure precise momentum measurement, which is crucial for accurate kinematic reconstruction and analysis.

- **Track Fit Quality:** The track fit quality, measured by the reduced χ^2/ndf , is required to be less than 1.8 to select well-reconstructed tracks and improve the accuracy of the muon momentum measurements.
- **Muon Momentum:** The momentum of chosen muons must be below 2000 GeV to prevent biases in the analysis caused by extremely high momentum muons, which are infrequent and might experience systematic uncertainties. This also serves as a sanity check for our units.
- **Trigger Criteria:** The muon must satisfy the **TOS**³ condition on all trigger lines [105], including **Hlt1SingleMuonVHighPT** and **Hlt2SingleMuonHighPT**. This requirement simplifies the corrections for trigger efficiencies and ensures that the event is selected by the high-level triggers designed to capture high- p_T muons.
- **Candidate Multiplicity:** Events are required to have no additional candidates. This ensures a clean sample by reducing combinatorial backgrounds and avoiding events with multiple muon candidates.
- **Z Veto:** Events are rejected if they contain another muon with transverse momentum $p_T > 25$ GeV and the invariant mass of the two muons $M(\mu^+, \mu^-) > 40$ GeV. This veto is applied to exclude potential background events from Z boson decays, which could mimic the signal and introduce contamination into the analysis.

Charged-hadrons selection cuts:

Special selection to make the charged-hadron model (QCD background):

- **Muon Kinematics:** The pseudorapidity range $2.2 < \eta < 4.4$ and the transverse momentum range $28 < p_T < 52$ GeV is chosen to mimic the same fiducial cuts for the signal.
- **Impact Parameter:** The impact parameter $\chi^2 < 9$ is applied to ensure that the selected daughter muons originate from the primary interaction vertex, minimising the likelihood of selecting muons from secondary decays or cosmic rays.
- **Isolation:** An isolation cut of $\mathcal{I} < 10.0$ GeV is imposed to ensure that the muons are not surrounded by significant energy deposits from other particles within a cone around the muon. This helps in reducing background from particles or other sources of non-isolated muons.
- **Track Fit Quality:** The requirement that the track fit $\chi^2/\text{ndf} < 1.8$ similar to what is done with the signal.
- **Muon Momentum:** The muon momentum is required to be less than 2000 GeV similar to what is done with the signal.
- **Candidate Multiplicity:** Events with multiple muon candidates are rejected, similar to what is done with the signal.

³This abbreviates “Trigger On Signal” [104].

Z candidate selection cuts:

Auxiliary selection of $Z \rightarrow \mu^+\mu^-$ events for normalisation equations that scale the simulated components to the data. These Z selections are applied to both muon charges.

- Muon Kinematics: The pseudorapidity range $1.7 < \eta < 5.0$ is selected to cover the full acceptance of the LHCb detector. The transverse momentum range $28 < p_T < 52$ GeV focuses on muons with high momentum, similar to what is done for the signal, to enhance signal purity and reduce background.
- Impact Parameter: The impact parameter $\chi^2 < 100$ is applied to ensure that the selected muons are likely to originate from the primary interaction vertex, but with a looser cut compared to the signal to accommodate potential differences in background distributions.
- Isolation: An isolation cut of $\mathcal{I} < 10.0$ GeV is used, similar to what is done for the signal, to reduce background by ensuring that the selected muons are not surrounded by significant energy from other particles.
- Invariant Mass Requirement: The invariant mass window $|m_Z - 91.0| < 14$ GeV is applied to select events where the invariant mass of the muon pair is close to the known Z boson mass, ensuring that the selected events are consistent with Z boson decays.
- Relative Momentum Error: A relative momentum error of less than 0.06 is imposed to ensure precise momentum measurement, similar to what is done for the signal.
- Track Fit Quality: The track fit quality, requiring $\chi^2/\text{ndf} < 1.8$, is used to select well-reconstructed tracks, similar to what is done for the signal.
- Muon Momentum: The momentum is required to be less than 2000 GeV, similar to what is done for the signal,
- Trigger Requirements: The event must trigger on the specified lines (`mu_LoMuonEWTOS`, `mu_H1t1TOS`, `mu_H1t2TOS`) to ensure that the event was selected by the relevant triggers designed to capture the desired signal, similar to what is done for the signal candidates.
- Candidate Multiplicity: Events with multiple muon candidates are rejected, similar to what is done for the signal, to maintain a clean dataset and avoid complications from multiple candidates.

5.2 Calibrations applied to the data

In this section, the correction of the momentum in the data is described.

The leading biases in the measurement of the momentum of a particle, with charge q , follow

$$\frac{q}{p} \rightarrow \frac{q}{p'} = \frac{q}{\alpha p} + \delta, \quad (5.2)$$

where p' is the “biased” momentum, and p is unbiased. Here, α is a momentum scale factor and δ is a charge-dependent curvature bias.

These calibrations are performed in parallel rather than in an arbitrary sequence; they are complementary in that one is multiplicative in p and the other is additive in q/p . However, it is important to mention that the scaling α is not necessary for this analysis. These calibrations improve the resolution and uniformity of the momentum response in the data, so that the momentum smearing⁴ model, described in Section 5.3.1, can be relatively simple. We do not apply a run–number–dependent calibration since our data at 5.02 TeV does not have a significant time dependence (some other datasets do). In fact, the entire dataset was obtained over a short period of time that did not exceed two weeks. The q/p corrections do have an associated systematic uncertainty because the momentum smearing makes no further correction for curvature biases.⁵

5.2.1 Curvature corrections with the pseudomass method

We apply charge–dependent curvature bias corrections determined with $Z \rightarrow \mu^+\mu^-$ decays and the pseudomass method [106]. Muons from Z decays tend to have very small curvature values, so this analysis is particularly sensitive to the alignment of the tracking system. The bulk of the strategy is summarised in [106]. The pseudomass is an estimate of the mass of two–particle final states, in which the magnitude of one of the momenta is ignored. Considering the decay $Z \rightarrow \mu^+\mu^-$, we define the two pseudomasses:

$$\mathcal{M}^\pm \equiv \sqrt{2p^\pm p_T^\pm \frac{p^\mp}{p_T^\mp} (1 - \cos \theta)}, \quad (5.3)$$

where p^+ and p^- denote the magnitudes of the μ^+ and μ^- momenta (and similarly for the transverse momenta $p_T^\pm$), and θ is the opening angle between the two muons. Effectively, the pseudomass estimates the mass under the assumption that Z bosons are produced with transverse momenta that are small compared to the Z boson mass. Crucially, where $\langle p_T^+ \rangle = \langle p_T^- \rangle$ we can write $\mathcal{M}^+ = \mathcal{M}^-$. Unlike the mass, in which biases of the form of Eq. 5.2 strongly cancel between the μ^+ and μ^- , the pseudomasses have first–order sensitivity to these biases. The pseudomass is a powerful tool for determining the curvature biases in the data, as it is sensitive to them in both the data and simulation and insensitive to the momentum scale and resolution. A simultaneous fit of the positive and negative pseudomass is performed at intervals of η and ϕ .⁶ This is directly translated into q/p bias corrections.

The data are corrected by accounting for the difference in q/p biases between the data and simulation, while the simulation remains uncorrected. This approach removes a small bias that arises from the vector and axial vector couplings inherent in the physics process [106]. The statistical uncertainties on the pseudomass asymmetries are propagated through systematic uncertainties on the differential cross–section measurement.

Curvature biases are determined in intervals of η and ϕ and are shown in Figure 5.1. As a reminder, our dataset is only for the 5.02 TeV run from 2017 with magnet down polarity, so the

⁴Smearing is used to denote both stochastic/resolution smearing and systematic biases.

⁵Technically it can, and this feature was used in the previous LHCb m_W analysis [21], but the new treatment is considered as a significant optimisation of the procedure.

⁶Our dataset is limited to a single year and one polarity, which limits us to a single fit. If additional data were available, with multiple years and polarities, we would perform a fit for each year and polarity.

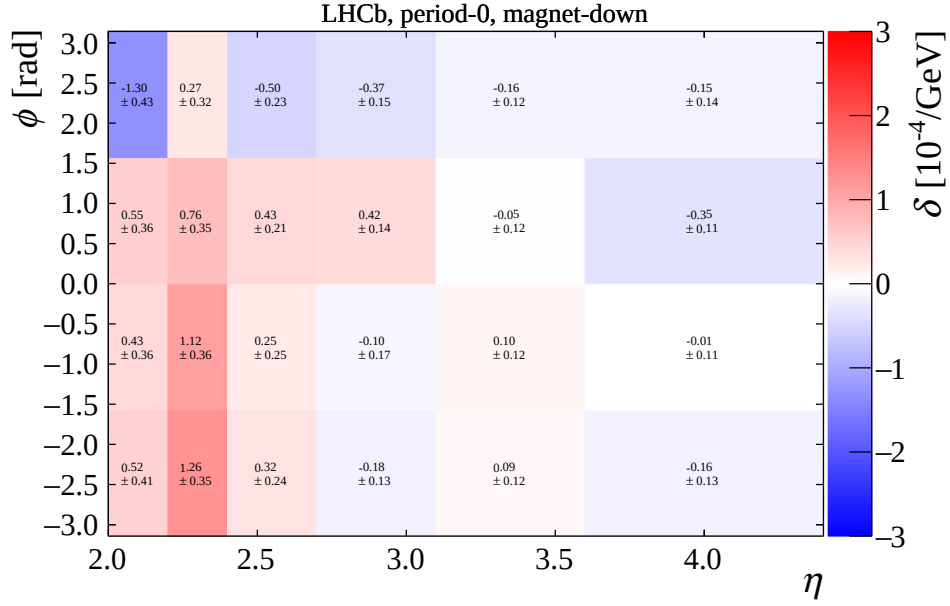


FIGURE 5.1: Curvature biases determined with the pseudomass in intervals of η and ϕ .

statistics are low. We binned the data into nonuniform bins of η and ϕ to ensure that the pseudomass asymmetry is well determined in each bin. We expect to have more statistics in the lower η bins, so we can afford to have smaller bins there, but we need to have larger bins in higher η regions to have enough statistics to determine the pseudomass asymmetry.

Figure 5.2 shows that the Z mass resolution is improved by applying the pseudomass corrections. One can observe that the pseudomass correction is unable to work very well with these low statistics of the 5.02 TeV sample, but at least the mass resolution is not worse compared to before. The pseudomass correction is expected to contribute greatly to the systematic uncertainties in the final result, but it is still a necessary step to ensure that the data are as unbiased as possible. A larger dataset is expected to enhance the precision of the pseudomass correction due to an increase in events, thereby reducing statistical uncertainties and producing more representative distributions for ϕ and η . This expectation is based on comparisons with other pseudomass applications at 13 TeV, which show demonstrably better results [21].

5.3 Corrections to the simulation

This section provides an overview of the corrections applied to our samples of simulated events. These corrections are essential to ensure the accuracy and reliability of our results and to account for discrepancies between the simulated and actual data. Since we intend to compare the simulation to the data, it is crucial to address any systematic uncertainties that may arise from the simulation's imperfections. For instance, the simulation of electroweak processes has limited perturbative accuracy, and the characterisation of the detector material is merely an approximation. These various inaccuracies must be corrected to ensure that our measurement achieves

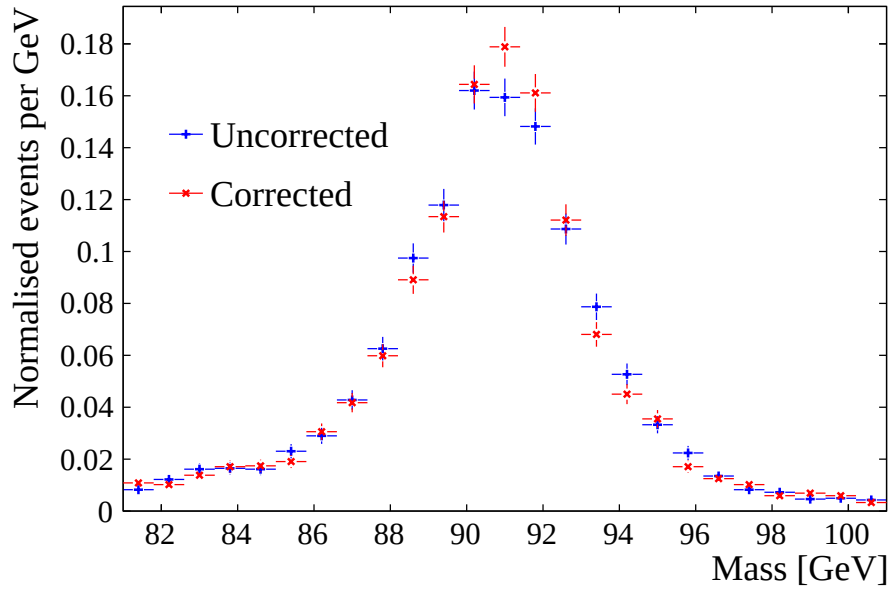


FIGURE 5.2: Mass distribution of Z candidates with and without the pseudomass corrections.

the desired level of precision. These corrections encompass a range of factors, from momentum smearing to isolation calibrations, each of which can significantly impact the outcomes of the simulation.

5.3.1 Momentum smearing

Having applied the calibrations to the momenta in the data, the resolution and uniformity are much closer to the idealised case in the simulation. However, smearing is still required. The master smearing formula is:⁷

$$p \rightarrow (p + \beta)(1 + \alpha)(1 + f(\eta)\mathcal{R}_1(\eta)\sigma_1)(1 + \mathcal{R}_2 p \sigma_2^a)(1 + \mathcal{R}_2 p \sigma_2^b / \cosh \eta), \quad (5.4)$$

where

$$f(\eta) = \begin{cases} 1 & \eta < 3.3 \\ \zeta & \eta \geq 3.3 \end{cases}, \quad (5.5)$$

and where the $\mathcal{R}$ are random numbers sampled from a standard normal distribution, β is an energy offset parameter, α is an momentum scale offset, and σ_1 and σ_2 are momentum and curvature smearing parameters, respectively [21]. A fit of the $Z \rightarrow \mu^+ \mu^-$ data and simulated events is made to determine the best values of the α and σ_2 parameters. All other parameters have fixed values. We tested the effect of changing the parameter ζ , which is a factor that scales the resolution of the momentum in the forward region, but found that it had no significant

⁷This is a generalised formula, where there is a choice of either using the σ_2^a or σ_2^b parameter, while fixing the other to zero.

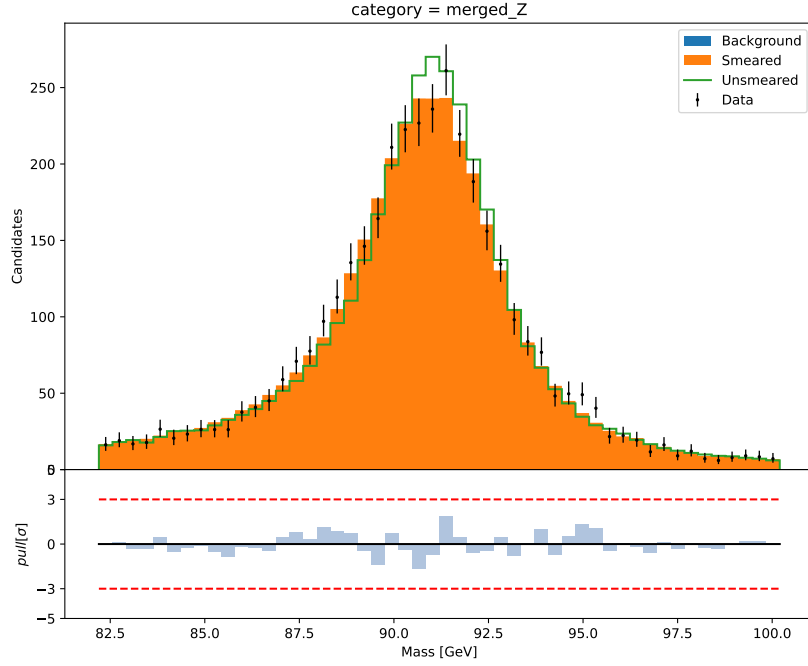


FIGURE 5.3: Mass distribution of Z candidates. The momentum scale fit results are overlaid.

effect on the fit results. As such, we do not associate any systematic uncertainty with the fixed parameter ζ .

Table 5.2 shows the results of the momentum smearing fit, with $\chi^2/\text{ndf} = 106.49/146$. Figure 5.3 shows the Z mass distribution, after combining all η categories, compared to the fit results.

TABLE 5.2: Results of the momentum smearing fit, where the parameters correspond to the terms of Eq. 5.4.

Parameter	Value
$\alpha: i_{\eta}=0$	-0.1765 ± 0.9986
$\sigma_2^b: i_{\eta}=0$	16.6855 ± 3.4400
$\alpha: i_{\eta}=1$	-1.2356 ± 0.5879
$\sigma_2^b: i_{\eta}=1$	22.5124 ± 3.1867
σ_1	2.0000 (fixed)
σ_2^a	0.0000 (fixed)
δ	0.0000 (fixed)
β	0.0000 (fixed)
ζ	1.5000 (fixed)

The systematic uncertainties on the momentum smearing model include the propagation of statistical uncertainties from the momentum smearing fit results, computed by generating variations through random sampling and then calculating their covariance matrix. The root mean

square (RMS) of the differential cross-section values from these variations determines the uncertainty.

5.3.2 Efficiency corrections

The overall efficiency includes the requirement that the muons are reconstructed as long tracks and identified as muons, and that the $Z \rightarrow \mu^+\mu^-$ candidates pass the selection requirements described in Sec. 5.1.3.⁸

The imperfect simulation-based efficiency modelling is corrected with weights.

- **Trigger:** The combined TOS (`L0MuonEW_TOS` and `Hlt1SingleMuonHighPT_TOS`) efficiencies are measured with the usual “tag-and-probe” method with $Z \rightarrow \mu^+\mu^-$ candidates. Tag-and-probe is simply a method to measure the efficiency of a selection requirement where one muon is required to pass a selection requirement, and the other is not. One “tag” muon is required to be TOS, so that the trigger decision is independent of the other “probe” muon. The fraction of events in which the probe is also TOS defines the estimate of the efficiency. The ratio of efficiency measured in the data and the simulation can be used to define a set of weights to be applied to the simulation.
- **Identification:** This is treated similarly to the trigger efficiency, except that it only includes $Z \rightarrow \mu^+\mu^-$ decays, specially selected without a muon identification requirement on one muon.
- **Tracking:** A dedicated selection of $Z \rightarrow \mu^+\mu^-$ candidates is used, in which the probe muon is reconstructed with a special algorithm [107] that only uses information from the muon stations and the TT. The fraction of events in which the probe track is in a match with a long track defines the tracking efficiency estimate.

Altogether, these three efficiencies account for the total muon reconstruction efficiency:

$$\varepsilon = \varepsilon_{\text{trk}} \times \varepsilon_{\text{L0, Hlt}} | \text{ID} \times \varepsilon_{\text{ID}}, \quad (5.6)$$

To properly model the variations in efficiency across the phase space that signal muons occupy, we must group the probes and derive efficiency corrections depending on the kinematics of a given signal muon. For all three efficiencies described above, we first split into four bins of azimuthal angle ϕ and eleven bins of muon η , covering the fiducial region of the detector. The kinematics of the signal muon determine the parameterisation picked up and the point in $1/p_T$ at which ε_{MC} and $\varepsilon_{\text{D}}/\varepsilon_{\text{MC}}$ are evaluated. These parameterisations can be uniform, where the efficiency is assumed to be constant across the kinematic range, or linear, where the efficiency is modelled as a linear function of the kinematic variable. Within this analysis, we rely only on uniform parameterisations. These parameterisations are derived from simulations and experimental data to capture the efficiency change in bins of eta and phi.

⁸The signal efficiency of the track fit χ^2/ndf and impact parameter χ^2 cuts are about 99%, so we do not need to worry about the determination of corrections or uncertainties associated with these cuts

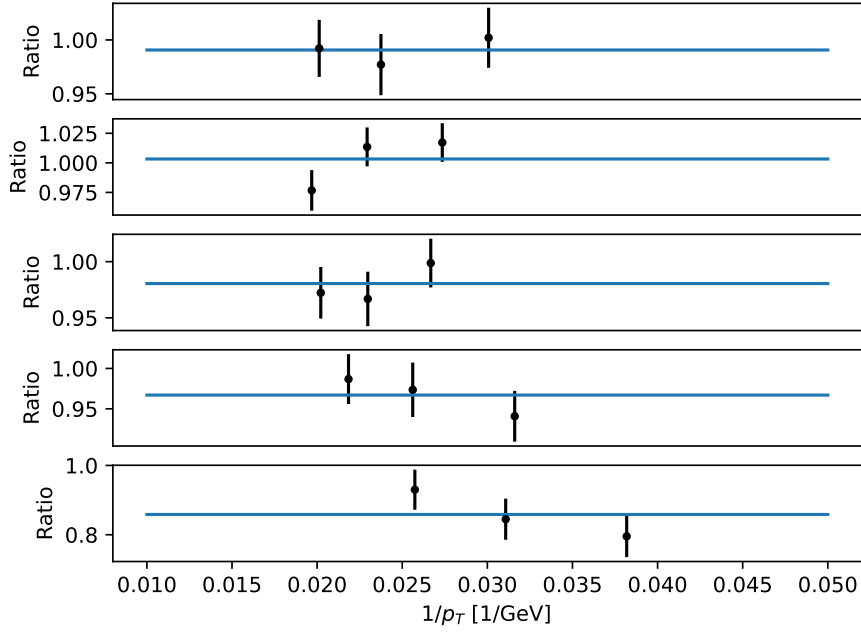


FIGURE 5.4: Ratio of the trigger efficiency between data and simulation as a function of $1/p_T$, with the fit results overlaid. The rows span intervals in η with boundaries at [1.65, 2.2, 2.75, 3.3, 3.85, 4.4]. The blue line shows the fitted uniform $1/p_T$ dependence. We observe that the efficiencies are high.

Each muon gets a weight according to the efficiency in equation 5.6, which are then multiplied together to give a total weight due to the mismodeling of muon reconstruction. For an event with a single muon in the final state, the appropriate weight is the scale factor $\varepsilon_D/\varepsilon_{MC}$. For events with two muons in the final state, the weight assigned is based on the number of muons that pass the selection requirement. For example, if the μ^+ in an event passes while the μ^- fails, the appropriate weight is

$$w = \frac{\varepsilon_{\text{data}}^+}{\varepsilon_{\text{MC}}^+} \times \frac{1 - \varepsilon_{\text{data}}^-}{1 - \varepsilon_{\text{MC}}^-}, \quad (5.7)$$

thereby correcting for the mismodelled *efficiency* of the requirement on μ^+ , and the mismodelled *inefficiency* of the μ^- based on where it is in the phase space. Since these events are rejected based on the p_T of the second positively identified muon, they are weighted to correct for mismodelled tracking and ID (in)efficiencies based on whether the second muon passes or fails the tracking and ID requirements. The second muon's truth information gives the kinematic dependence of these weight factors in the cases where the true muon falls within the LHCb acceptance but is not properly reconstructed or positively identified as a muon.

Figure 5.4 shows the ratio of the trigger efficiency between the data and the simulation, in intervals of η , with the uniform $1/p_T$ dependence overlaid. The splitting in intervals of ϕ is included in the investigation of systematic uncertainties in Section 6.5. Figures 5.5 and 5.6 show the corresponding ratios for identification and tracking efficiencies, respectively.

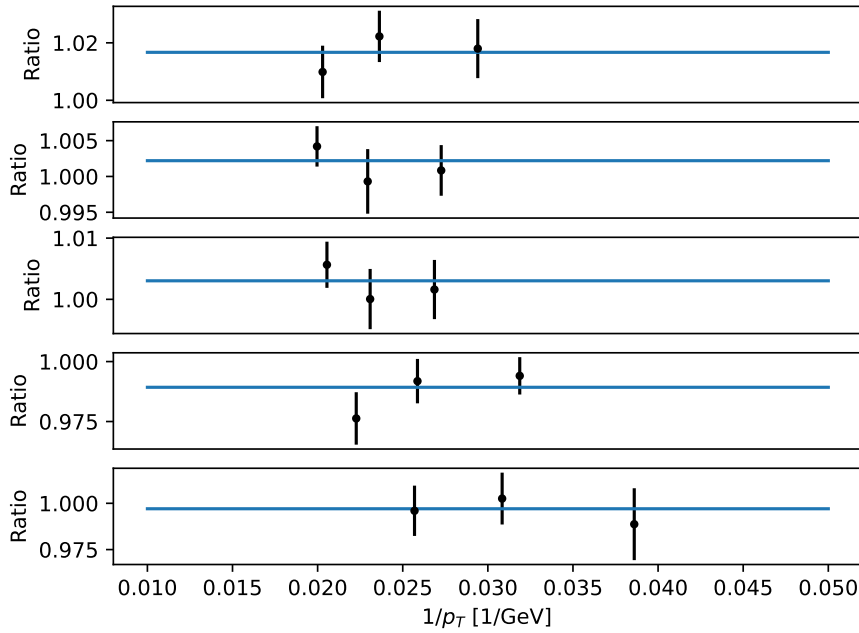


FIGURE 5.5: Ratio of the muon identification efficiency between data and simulation as a function of $1/p_T$, with the fit results overlaid. The rows span intervals in η with boundaries at [1.65, 2.2, 2.75, 3.3, 3.85, 4.4]. The blue line shows the fitted uniform $1/p_T$ dependence. We observe that the efficiencies are high.

The systematic uncertainties in efficiency corrections include the propagation of statistical uncertainties by randomly varying efficiency values within their uncertainties, followed by fitting the p_T dependence.

The systematic uncertainties will be discussed in more detail in the systematic uncertainty section 6.5.

5.3.3 Corrections to the isolation variable

QCD background isolation calibration

NoMuID data, which refer to events where no muons are identified, are particularly useful for describing charged-hadron (QCD) backgrounds in analyses involving leptons (like muons) because QCD events often produce particles that can mimic the signatures of isolated leptons. In QCD processes, particularly in multi-jet events, hadrons can decay into particles that may appear similar to muons in the detector but do not pass the stringent identification criteria. Hence, the **NoMuID** data serve as a proxy for the QCD background. We use a dual-calibration strategy that incorporates both **NoMuID** data and the $Z \rightarrow \mu^+ \mu^-$ data to derive a robust calibration factor.

We used 13 TeV **NoMuID** data to calibrate the isolation variable ($\mathcal{I}$) for the simulated QCD background sample. This is because we do not have the equivalent 5.02 TeV **NoMuID** data, and the

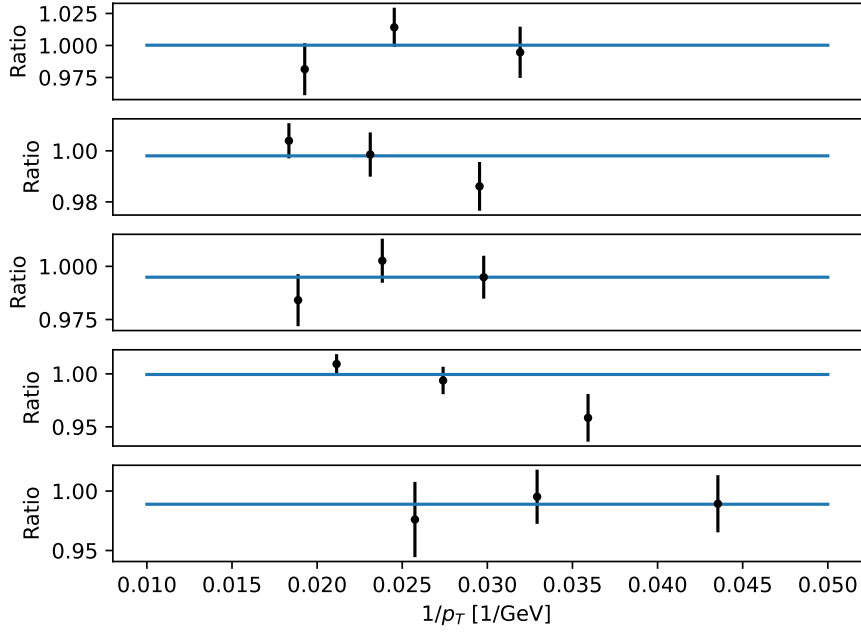


FIGURE 5.6: Ratio of the tracking efficiency between data and simulation as a function of $1/p_T$, with the fit results overlaid. The rows span intervals in η with boundaries at [1.65, 2.2, 2.75, 3.3, 3.85, 4.4]. The blue line shows the fitted uniform $1/p_T$ dependence. We observe that the efficiencies are high.

corrections are expected to be similar at both energies.⁹ This choice allows the $\mathcal{I}$ calibration factor to be derived by comparing the binned NoMuID $\mathcal{I}$ to the binned QCD background simulated $\mathcal{I}$. Furthermore, to improve the accuracy of the calibration, an additional constraint is introduced by incorporating $Z \rightarrow \mu^+ \mu^-$ data. This constraint is crucial for accounting for potential missing electroweak processes at low isolation levels, where the contribution from electroweak interactions may not be well represented in the original QCD background simulations. We fit the isolation distribution both muon charges simultaneously; hence, a delta factor is floated to account for the difference in the calibration factor between the two charges, where the calibration factor is now defined as

$$C' = C + q\delta. \quad (5.8)$$

This charge asymmetry factor, δ , is observed in the distributions and fit results, as shown in Figure 5.7 and Table 5.3, and is considered significant. This dual calibration strategy ensures a more robust and accurate $\mathcal{I}$ calibration, considering both the specific characteristics of the QCD background and potential deficiencies in the original simulations, ultimately enhancing the precision of the analysis.

We set the $\mathcal{I}$ range to be from 0 until the same $\mathcal{I}$ cut-off value that we implement later on in the background subtraction fit (8 GeV). A Gaussian kernel smoothing technique is also used to

⁹This is a simplification, but a reasonable one that was tested comparing the 5.02 TeV and 13 TeV Z data and simulation.

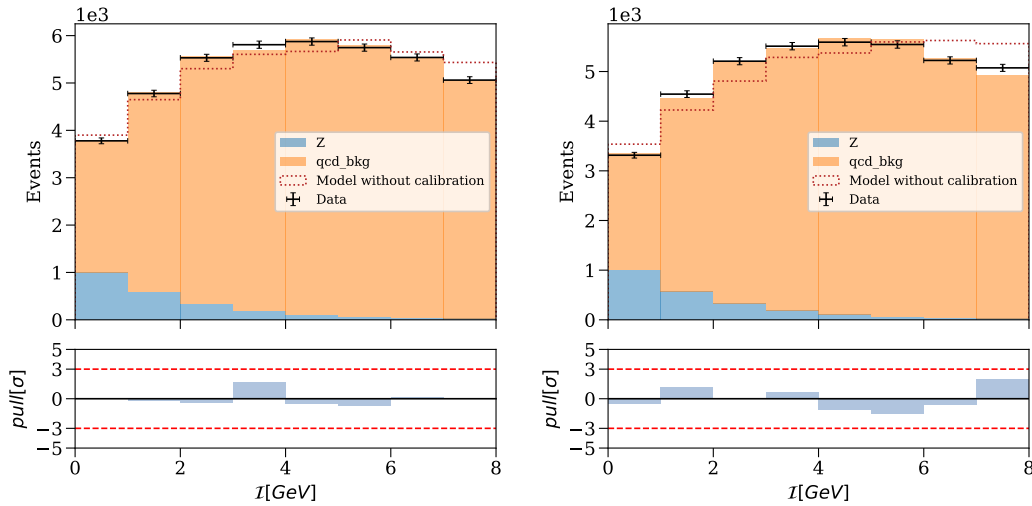


FIGURE 5.7: Distributions of $\mathcal{I}$ for the (left) positively charged and (right) negatively charged candidates in the inverted muon–identification samples. The results of the background isolation calibration fit are overlaid.

create smoothed histograms from the unbinned $\mathcal{I}$ sample data [108]. This smoothing process is useful for mitigating statistical fluctuations during the minimisation process and preventing abrupt changes in the minimisation function.

We then use a simple χ^2 minimisation to find the best calibration factor, where the fit model is the sum of the QCD background and the smoothed histogram of the data Z , each scaled by a fit fraction. We compare this fit model to the smoothed **NoMuID** data histogram, determining the optimal calibration factor as the one that minimises the value of χ^2 . This calibration factor is then used to scale the QCD $\mathcal{I}$ parameter during the background subtraction process. The fit results can be seen in Table 5.3, where we get $\chi^2/\text{ndf} = 7.444/10$.

parameter	value	error
δ	0.055	0.012367
C	0.821	0.012378
Hadron fraction (positive charge)	0.945	0.005119
Hadron fraction (negative charge)	0.943	0.004772

TABLE 5.3: Results of the isolation calibration fit for the hadronic background.

Signal Isolation Calibration

The Signal Isolation Calibration, on the other hand, is performed using $Z \rightarrow \mu^+\mu^-$ data. This calibration is applied to the W^+ and W^- simulated samples, which represent our signal. We obviously cannot use the W data to calibrate the isolation variable, as the W data are what we are trying to measure.

We do a similar process for the signal as we did for the QCD background, but we use the $Z \rightarrow \mu^+\mu^-$ data to calibrate the isolation variable for the $Z \rightarrow \mu^+\mu^-$ simulated sample, and apply

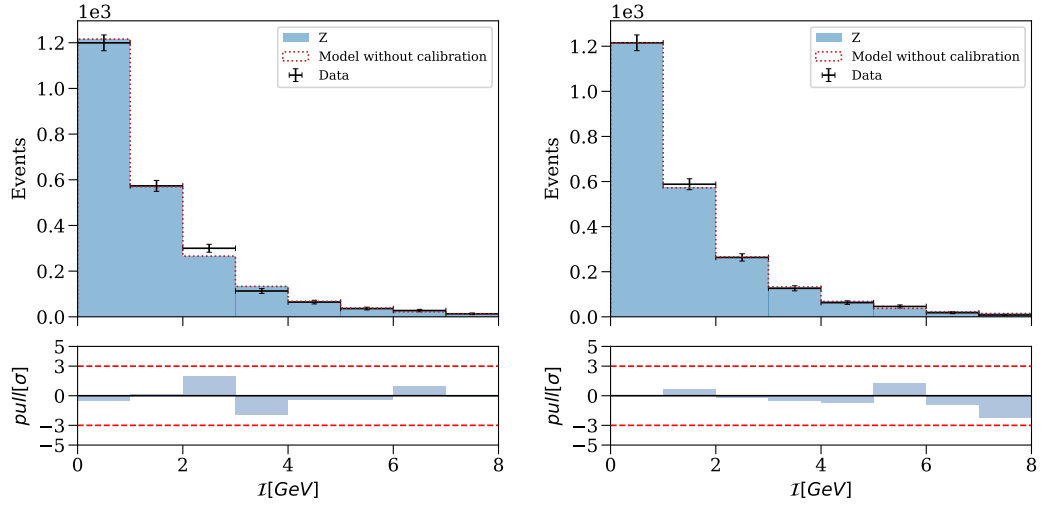


FIGURE 5.8: Distributions of $\mathcal{I}$ for the (left) positively charged and (right) negatively charged candidates in the $Z \rightarrow \mu^+\mu^-$ samples. The results of the signal isolation calibration fit are overlaid.

that calibration to the W^+ and W^- simulated samples (signal). The delta factor to account for the charge asymmetry has been found to be insignificant and therefore is omitted here. As such, we fix the delta factor to zero in the fit. The fit results can be observed in Table 5.4, where we get $\chi^2/\text{ndf} = 15.207/13$. The applied calibration results can be seen in Figure 5.8. We can also

parameter	value	error
C	0.938	0.01653

TABLE 5.4: Results of the isolation calibration fit for the electroweak signals.

check that the corrected $Z \rightarrow \mu^+\mu^-$ isolation distributions look fine in smaller intervals of recoil projection than can be used for the background subtraction fits. We can also plot in bins of η . These plots are shown in Appendix A.

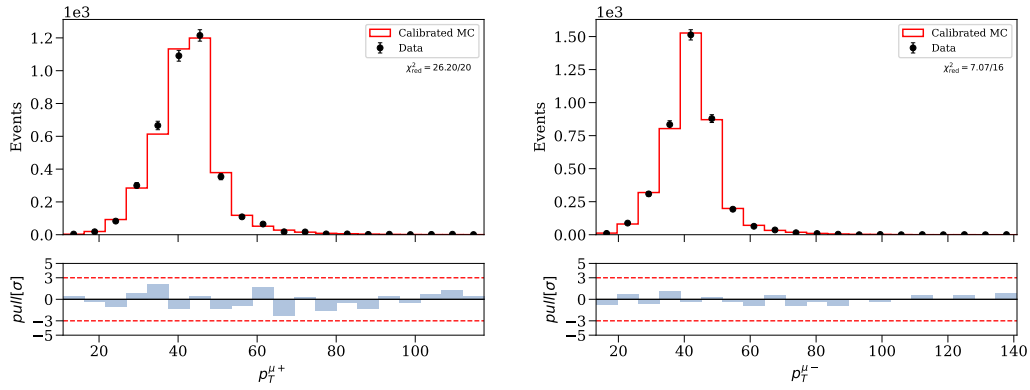


FIGURE 5.9: Distributions for μ^+ transverse momentum (left) and μ^- transverse momentum (right) comparing the data and the calibrated simulation.

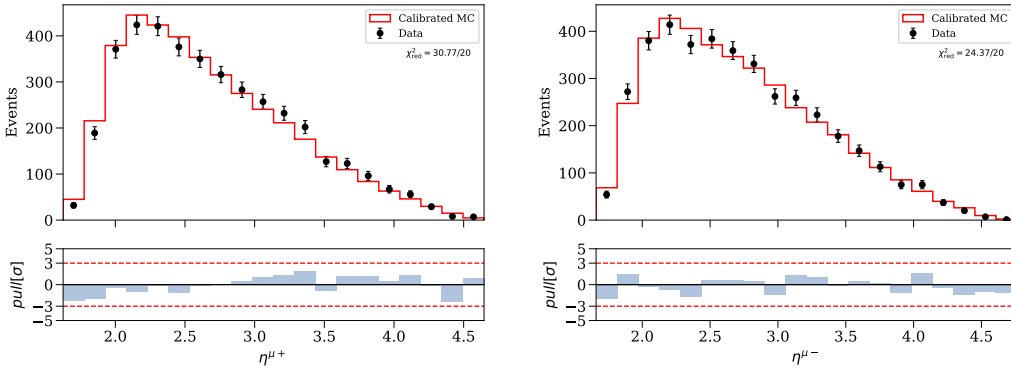


FIGURE 5.10: Distributions for μ^+ pseudorapidity (left) and μ^- pseudorapidity (right) comparing the data and the calibrated simulation.

5.4 Validation of the corrected simulation

Having applied the smearing and efficiency weights to the simulation, it is important to validate that it adequately models the data. The distributions shown in Figures [5.9 – 5.13] provide a comparison between the data and the simulation. All selection requirements are applied, except for requirements for the variable whose distribution is presented. Note that the specific variables shown are not of particular importance and were selected somewhat arbitrarily. The important thing to observe is the agreement between simulation and data.

Observing these plots, we can see that the simulation is well calibrated and accurately models the data. This is also evidenced by the reduced χ^2 values displayed on the plot, which are close to 1, signifying a strong agreement between the simulation and the data. Again, this is critical for the success of the analysis, as the simulation is used to model the background and signal processes, and any discrepancies could lead to inaccurate results. We can then proceed with the background subtraction fit and unfolding process to obtain the differential cross-section measurement. All plots use 30 bins to ensure consistency across all visualisations and display the reduced χ^2 values

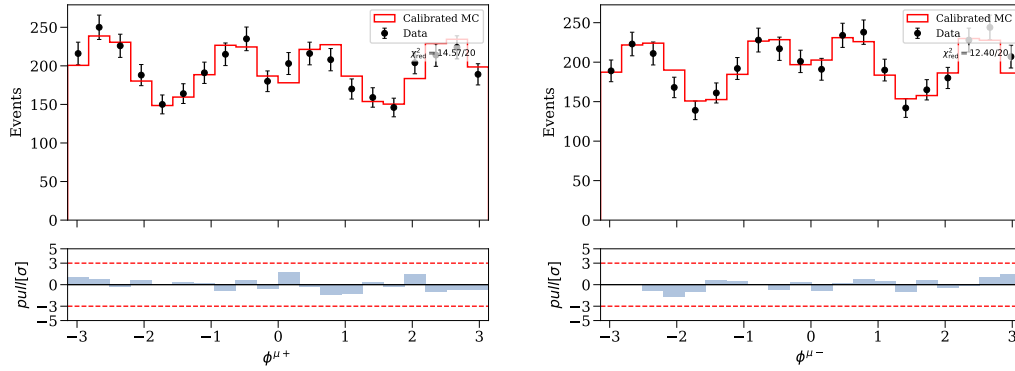


FIGURE 5.11: Distributions for μ^+ ϕ angle (left) and μ^- ϕ angle (right) comparing the data and the calibrated simulation.

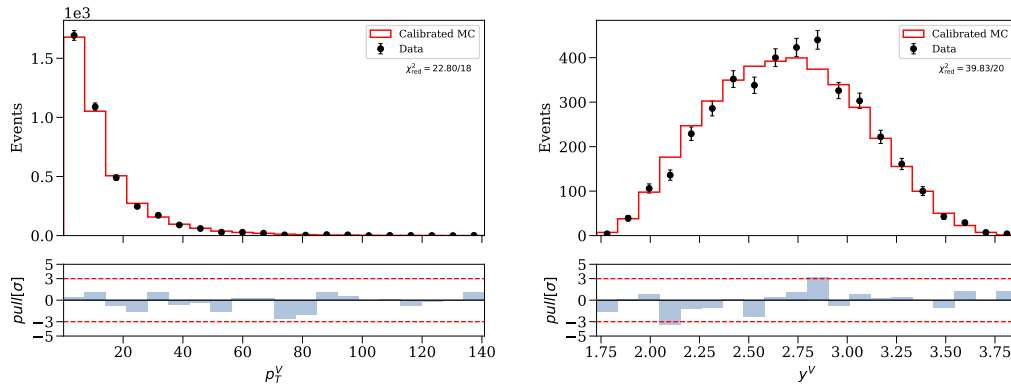


FIGURE 5.12: Distributions for vector boson transverse momentum (left) and vector boson rapidity (right) comparing the data and the calibrated simulation.

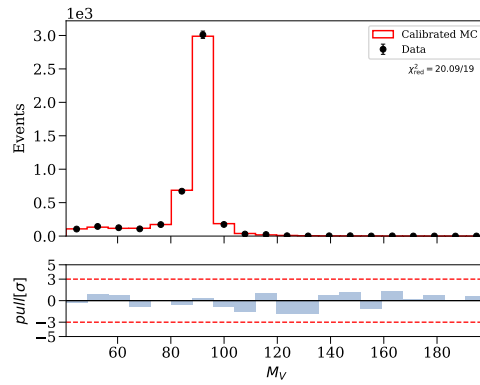


FIGURE 5.13: Distributions for vector boson mass comparing the data and the calibrated simulation.

along with the pulls. Some plots contain empty bins, so the number of degrees of freedom may be less than 30. The reduced χ^2 values are near 1, showing a strong agreement between simulation and data.

5.5 Analysis software projects

The analysis configuration is based on collaborative efforts of the LHCb “ew-analyses” group, which originated from the initial m_W analysis software project and now involves more than 20 collaborators around the world. We are part of the LHCb QCD, Electroweak, and Exotica (QEE) working group. We use Python, C++, Bash, ROOT [109], Snakemake [110], Gitlab [90], and LHCb software for physics data analysis and enhancement of our joint Gitlab repository, which contains numerous shared codes for various analyses¹⁰. We employ CERN’s batch system and software to create simulation templates, accessible via a project repository called “ew-tuples”. We have used a variety of event generators, simulators, and precision measurement tools (including Pythia, POWHEG, Geant4, DYTURBO, MCFM, etc.) to simulate events and the detector response, as referenced in the “ew-generators” project repository [44, 111, 76, 43, 103]. Furthermore, we maintain a joint Gitlab repository for document preparation “ew-docs”, where we draft our publications and reports using L^AT_EX.

My individual work, which I completed independently with feedback from my supervisors and colleagues, includes handling 5.02 TeV data, calibrating the isolation variable, performing background subtraction and unfolding processes, and conducting the final differential cross-section measurement. In addition, I created the simulation templates for the 5.02 TeV data, which were incorporated into the analysis, and managed the fiducial cross-sections necessary for the background subtraction fit. I also played a role in improving our approach to handling DYTURBO simulations, which are used to apply higher-order corrections to the templates to construct the fit model used in the determination of the mass of the W boson. This analysis bases the W boson mass determination on previous measurements by the group as referenced in [21]. However, the approach used in the new analysis presented in this thesis is distinct as it uses unfolded data to gauge the mass of the W boson, contrasting with the other measurement of the mass of the W boson at LHCb, which incorporates backgrounds and efficiencies into the fit model.

¹⁰The links are currently only accessible to members of the LHCb collaboration.

Chapter 6

W boson differential cross-section measurement

“Nature uses only the longest threads to weave her patterns, so that each small piece of her fabric reveals the organization of the entire tapestry.”

– Richard P. Feynman

As a reminder, the main goal of the analysis is to measure the differential cross-section in bins of muon transverse momentum ($\frac{d\sigma}{dp_T}$) for $W^+ \rightarrow \mu^+ \nu_\mu$ and $W^- \rightarrow \mu^- \bar{\nu}_\mu$ production, which involves correcting the raw p_T distribution of W candidates for background, detection efficiency, and resolution.

In this chapter, we will start with an overview of how we account for the uncertainties in the data and simulation using the Beeston–Barlow lite method in Section 6.1 [112]. Then, we will discuss the unfolding process and how to apply it “on-the-fly” in Section 6.2 and 6.3 respectively. We will present the fit results in Section 6.4. We will then discuss the systematic uncertainties in Section 6.5. Finally, we perform a comparison of the cross-section results with the theoretical predictions in Section 6.6.

6.1 Accounting for the uncertainties on data and simulation

This analysis depends on template fits. A template fit involves comparing experimental data to a set of simulated templates to determine the presence and characteristics of a particular signal within the data. Performing a template fit implies being affected by uncertainties arising from the limited number of events in the simulated samples, as well as from other factors such as model assumptions, imperfect simulation of the detector response, and systematic variations in the background and signal shapes. These sources of uncertainty can influence the accuracy of the fit, potentially leading to biases in the derived parameters and limiting the precision of the measurement. The Beeston–Barlow method [113] allows one to take into account both the accuracy of the distributions obtained from the simulated samples as well as the statistical precision derived from the data. In principle, this method introduces as many nuisance parameters

as bins are considered in the fit for each component, leading to a high-dimensional parameter space in the minimisation function. The curse of dimensionality is reduced by condensing the fluctuations in a single parameter by assuming that all the components can be stacked together (they are Poissonian or Gaussian in nature) and considering that the only important magnitude is the overall number of simulated candidates in each bin. This method is proposed in [112], and constitutes an improvement with respect to a standard fit χ^2 . The derivation starts with the definition of the probability of observing a given number of events given the expectation

$$P(k, \lambda) = \frac{1}{k!} \lambda^k e^{-\lambda}. \quad (6.1)$$

When fitting a model to data we typically use the $-\ln \mathcal{L}$ as the minimisation function which, expressed as the sum of contributions over the bin index i , becomes

$$-\ln \mathcal{L} = \sum_i^n [-k_i \ln \lambda_i + \lambda_i], \quad (6.2)$$

where we have dropped the term $\ln k!$ since it does not have any impact in the minimisation. In the asymptotic limit, Eq. 6.1 translates into a probability-density function that follows the normal distribution

$$G(k, \lambda) = \frac{1}{\sqrt{2\pi\lambda}} e^{-\frac{(k-\lambda)^2}{2\lambda}}, \quad (6.3)$$

and the minimised quantity resembles a χ^2 distribution [114]

$$\chi^2 = \sum_i^n \frac{(k_i - \lambda_i)^2}{\lambda_i}. \quad (6.4)$$

In Eq. 6.4, the parameters k_i and λ_i correspond to the number of observed data candidates and predicted candidates in bin i , respectively. The simplified Beeston–Barlow likelihood [112], also referred to as Beeston–Barlow lite (BBLite), is defined as

$$-\ln \mathcal{L} = \sum_i^n \left[-k_i \ln \beta_i \lambda_i + \beta_i \lambda_i + \frac{(\beta_i - 1)^2}{2\sigma_{\beta_i}^2} \right], \quad (6.5)$$

which has a lot of similarities to Eq. 6.2, but here the predicted events are multiplied by a factor β_i that is constrained to unity using the uncertainty on the prediction with a penalty term coming from a normal distribution (Gaussian constraint). The values of β_i are obtained from the consideration that the derivative of the minimised function must be zero for each β_i :

$$\beta_i^2 + (\lambda_i \sigma_{\beta_i}^2 - 1) \beta_i - k_i \sigma_{\beta_i}^2 = 0, \quad (6.6)$$

which is a quadratic equation that can be solved for β_i . It is important to note that Eq. 6.5 and Eq. 6.6 only stand as long as the number of overall events (both in simulation and in data) is sufficiently large to be in the asymptotic limit. If not, the constraint should be modified accordingly. We consider our existing dataset size to be adequately large. The BBLite method is hence a powerful tool to account for the uncertainties in the data and simulation, and it is used

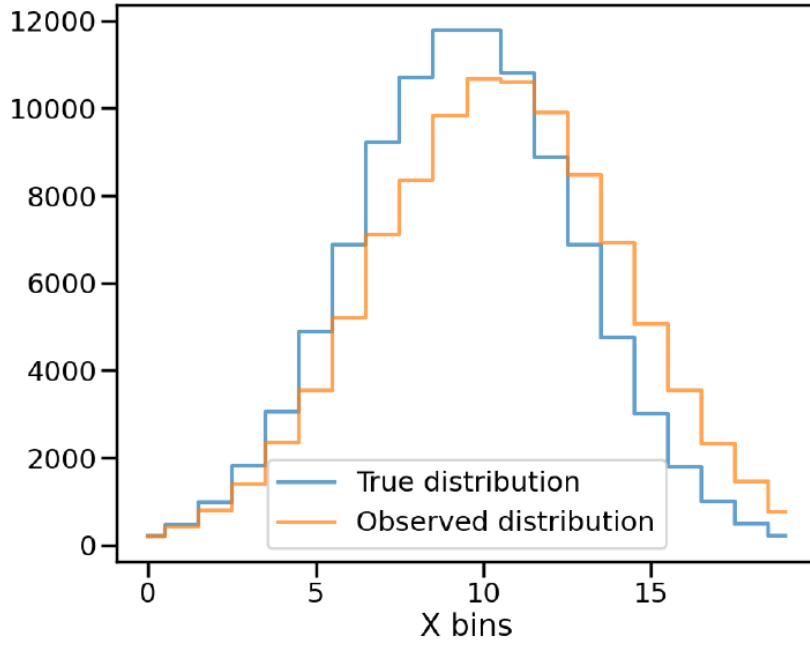


FIGURE 6.1: A diagram showing an example of the differences between the measured/observed distribution using a non-perfect detector vs the true distribution. For this example we plot the distributions as histograms with 20 uniform bins. The true distribution is derived from sampling a normal (Gaussian) distribution, while the observed distribution is the same but includes some added random Gaussian noise.

in this analysis to ensure that the fit results are robust and reliable.

6.2 Unfolding

In an ideal world, we would have a perfect instrument or detector that could measure our desired quantities with absolute accuracy and without errors. Unfortunately, in reality, detectors have finite precision, limited acceptance, ineradicable biases, and statistical uncertainties. Therefore, the value measured or observed by the detector is usually not the same as the true value. An example of this is shown in Figure 6.1.

To address the shortcomings mentioned above, we can use an unfolding process. Unfolding is a method of correcting a measured value, usually the distribution of a measurable quantity, for distortions caused by detector effects such as acceptance and resolution. We can use this process to infer the true distribution/spectrum from the measured/observed one [115].

We can mathematically link the measured quantity to the true quantity using equation 6.7.

$$\nu = R\mu, \quad (6.7)$$

where ν is the vector of measured events, μ is the vector of true events, and R is the response matrix that maps the true distribution to the measured distribution. Typically, the goal is to

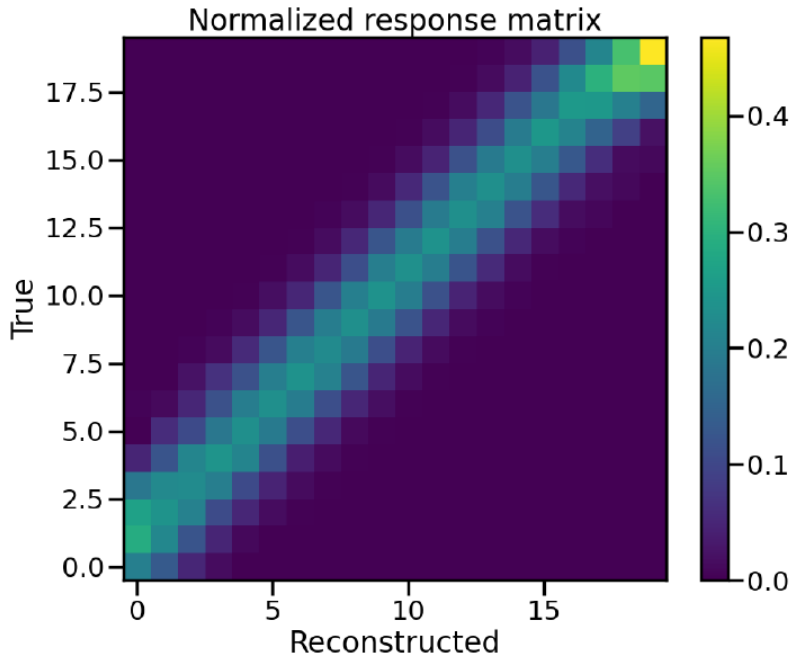


FIGURE 6.2: A plot of the normalised response matrix that encapsulates a detector's bias and resolution that cause us to measure a smeared version of the true distribution. This is based on the distributions illustrated in Figure 6.1.

obtain the true and unsmeared vector from the measured vector. To do this, equation 6.7 can be inverted to get

$$\mu = R^{-1}\nu, \quad (6.8)$$

where R^{-1} is the inverse of R . This process of inverting the response matrix and applying it to the measured vector to obtain the truth vector is referred to as unfolding.

The two-dimensional matrix R can typically be obtained using simulation events where both the generator-level and the reconstructed-level information are available. We then fill a matrix with all events for which there are both true and corresponding measured/reconstructed values. This matrix is filled with events for which both true and measured/reconstructed values are known, encoding the migration of events between bins.

Note that this matrix is dependent on the shape of the spectrum inside the true bins, which are estimated using a Monte Carlo approach. The response matrix values are subject to fluctuations due to the binomial nature of the simulated events. This must be taken into account as an additional source of systematic uncertainty. For example, one can fluctuate the matrix according to a multinomial distribution. An example of a normalised response matrix is shown in Figure 6.2.

After obtaining this response matrix, we can try to invert it and apply it to our measured data vector or histogram. However, we must be aware that if the response matrix is ill-conditioned, that makes the unfolding process unstable. As a result, even slight modifications to the measured

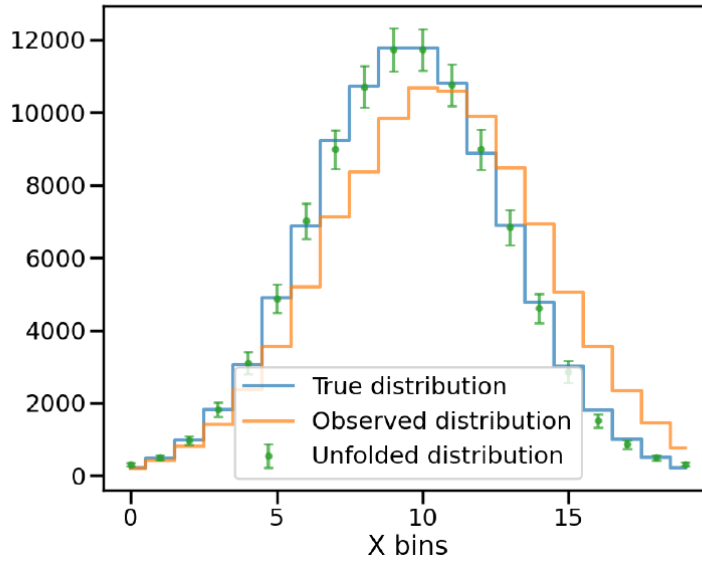


FIGURE 6.3: A plot showing a typical result of a successful unfolding procedure applied to measured data. We can observe that the unfolded result is in agreement with the true data.

data can lead to significant changes in the inferred true distribution. Therefore, a regularisation procedure is usually necessary to prevent high variance, except in cases where the response matrix is well-behaved, with little migration, and the measured data are highly precise.

Unfolding a data sample can be accomplished via many different methods, all of them subject to different advantages and disadvantages. For example, one could simply invert the matrix with an inversion tool, but this approach usually yields an incorrect result since it disregards the need for regularisation techniques that are essential for stabilising the unfolding process and reducing the impact of noise and fluctuations in the data. Another popular approach is bin-by-bin inversion, which, as the name implies, involves inverting each bin of the response matrix individually, treating each bin as an individual problem without taking into account the correlations or interdependencies between adjacent bins. This also lacks regularisation and ignores systematic effects. However, this method may be adequate if the migration effects are minor and can be compensated by a multiplicative factor (generalised efficiency).

Some of the well-known techniques include the iterative Bayesian unfolding method [116], the unfolding approach using singular value decomposition (SVD) [117], and methods employing machine learning algorithms such as neural networks [118].

A key drawback of the Bayesian unfolding method is its reliance on strong assumptions regarding the unfolded distribution, which can lead to biased results [116]. Various regularisation techniques have been proposed to mitigate these biases, including those presented in [119], but they often prove to be either computationally inefficient or difficult to apply in this analysis.

What we do in this analysis is, in fact, not unfolding, but folding. We will not be unfolding

the measured data to obtain the true distribution but rather folding the true distribution to obtain the measured data. Inversion of the response matrix is not necessary, and we can use the measured data directly in the fit. This is a much simpler and more straightforward approach and is the one we will be using in this analysis. However, since ‘unfolding’ is commonly used to describe such procedures, we will refer to this process as unfolding for consistency.

6.3 Unfolding on-the-fly

The goal of providing a measurement of the *W* cross-section as a function of the bare-level p_T becomes challenging due to the necessity of unfolding the measured p_T distribution for resolution effects. The approach adopted in this analysis is to perform the unfolding while fitting the data, which has already been proposed in other studies, such as [120]. The translation to reconstructed-level p_T occurs through the response matrix R_{ij} , which is obtained from simulation, and defines the probability of an event that with bare-level p_T in bin j ends up in a reconstructed bin i . The bin index i denotes reconstructed-level bins, while the index j denotes the bare-level p_T bin. The reconstructed-level distribution $\lambda_i^{\text{sig.}}$ is obtained from the bare-level distribution p_T^{bare} through the transformation

$$\lambda_i^{\text{sig.}} = \sum_j R_{ij} p_{T,j}^{\text{bare}} \quad (6.9)$$

Performing unfolding on-the-fly requires including the bare-level p_T distribution of the signal as parameters of interest in the fit, one parameter for each bin.

The definition of reconstructed-level bins is general in such a way that we can be referring to different p_T , η or $\mathcal{I}$ bins, and it could be represented as a function of different indices (one per variable) $R_{i_1, i_2, \dots, j}$ or flattened R_{ij} . In this explanation, the latter is chosen for simplicity.

However, the response matrix is subject to inaccuracies that arise from the limited number of simulated signal candidates. To account for the associated uncertainty, the response matrix can be sampled using a multinomial distribution, where for a given bare-level p_T bin, we define the probabilities of falling in each reconstructed bin as the elements of the response matrix. Note that, by definition, the sum of the elements of the response matrix in all reconstructed bins evaluates to the efficiency $\sum_i R_{ij} = \varepsilon_j$. In addition, the efficiency can be defined as a fraction, where the numerator is the count of selected events in bins of the bare muon transverse momentum, whereas the denominator is the count of simulated events in bins of p_T^{bare} , with the bare η constrained by the fiducial domain.

The number of Bernoulli trials is considered as the number of bare-level candidates in each p_T bin from simulation. With this information, it is possible to sample several expected signal events in the reconstructed bins, and therefore obtain a set of response matrices. For a given prediction of the p_T^{bare} distribution, we can obtain all the reconstructed-level predictions for the signal coming from each sampled response matrix. The covariance matrix associated with the

obtained values will therefore contain the information on the statistical fluctuations of the simulated sample size, as well as correlations among the different reconstructed-level bins. When processing the signal predictions, we therefore obtain the central prediction from the main response matrix and the covariance matrix from the sampled response matrices. This constitutes our signal model. Keep in mind that for the initial values of our p_T^{bare} differential cross-section data, we include additional underflow and overflow bins to accommodate further migration. However, these overflows are fixed because our tests showed that they have a negligible impact on the results. We perform a cross-check to verify this, with the findings detailed in Section 7.5.

The QCD background deserves special treatment as well. Given the observation of discrepancies in the p_T^{reco} between real data and simulated candidates in the 13 TeV data samples, and due to the lack of suitable real data samples to compute simulation corrections, such discrepancies are taken into account in the fit by scaling the p_T^{reco} distribution with different formulas. The associated contribution is constituted from the template rescaled according to a set of nuisance parameters, multiplied by an overall normalisation factor to the data. The scaling formula is chosen to be similar to the correction that is performed for other electroweak measurements at 13 TeV, and corresponds to a factor

$$\alpha = 1 + \frac{1}{\omega}. \quad (6.10)$$

For the remaining backgrounds, the predictions are determined from simulation and scaled to the Z distribution using the known cross-sections as highlighted in 5.1. The uncertainties in the background predictions are obtained using the number of events or the sum of squared weights in each bin as a proxy for the variance. This variance is scaled to the data, and the uncertainty is propagated accordingly taking into account the statistical power of the number of data candidates we are processing, such as

$$\sigma^2(\lambda_{i,\text{bkg.}}) = f_{\text{bkg.}}^2 N_{\text{data}}^2 \sigma^2(t_{i,\text{bkg.}}) + N_{\text{data}} f_{\text{bkg.}} t_{i,\text{bkg.}}, \quad (6.11)$$

where $\lambda_{i,\text{bkg.}}$ is the prediction of one background in bin i , $f_{\text{bkg.}}$ is the total fraction of data candidates for that particular background, $t_{i,\text{bkg.}}$ is the fraction of events for that particular background in the bin and N_{data} corresponds to the overall number of data candidates. Note that in the limit of high statistics, the variance can be approximated by the template and we can write

$$\sigma^2(t_{i,\text{bkg.}}) = \sigma^2(\alpha_{\text{bkg.}} t'_{i,\text{bkg.}}) = \alpha_{\text{bkg.}}^2 \sigma^2(t'_{i,\text{bkg.}}) = \alpha_{\text{bkg.}}^2 t'_{i,\text{bkg.}}, \quad (6.12)$$

where $t'_{i,\text{bkg.}}$ corresponds to the un-normalized template and $\alpha_{\text{bkg.}}$ is its normalization factor

$$\alpha_{\text{bkg.}} = \frac{1}{\sum_i t'_{i,\text{bkg.}}}. \quad (6.13)$$

As a consequence, the first element in Eq. 6.11 will scale with the square of $\alpha_{\text{bkg.}}$ while the second will be linear, so in the limit of a high number of template events we get

$$\lim_{\alpha_{\text{bkg.}} \rightarrow 0} \sigma^2(t_{i,\text{bkg.}}) = N_{\text{data}} f_{\text{bkg.}} t_{i,\text{bkg.}}, \quad (6.14)$$

which evaluates to the template of the model scaled according to the number of data candidates. Therefore, in the limit of having infinite simulations to build the templates, the corresponding uncertainties become negligible, and we recover the classic definition of the χ^2 in binned fits [112].

The overall covariance matrix is obtained from the sum of the different contributions

$$C = \sum_b C_b + C_{\text{sig.}}, \quad (6.15)$$

where $\sum_b C_b$ is the sum of the covariance matrices of the backgrounds, which are all diagonal; and $C_{\text{sig.}}$ is the covariance matrix of the signal, which has off-diagonal elements to account for the bin-to-bin correlations, and is obtained through the procedure discussed above.

The overall covariance matrix can be used to either do a χ^2 based fit, following Eq. 6.4, which can be redefined as

$$\chi^2 = \sum_{i,k}^n (k_i - \lambda_i) C_{ik}^{-1} (k_k - \lambda_k), \quad (6.16)$$

where n is the number of reconstructed-level bins; or following Eq. 6.5 to obtain

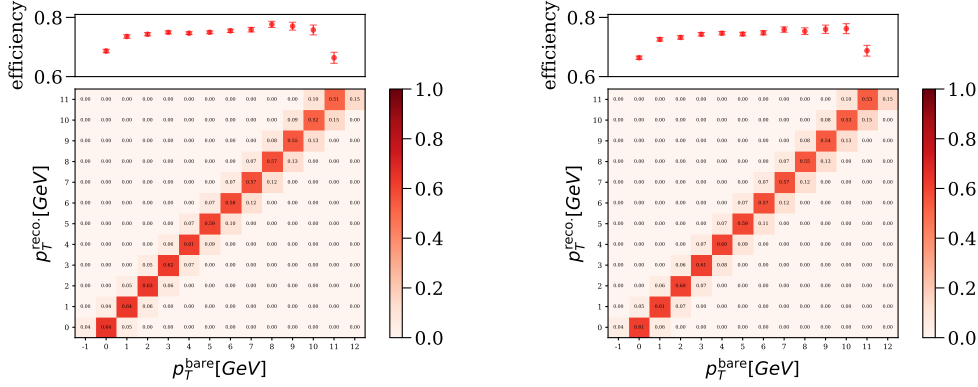
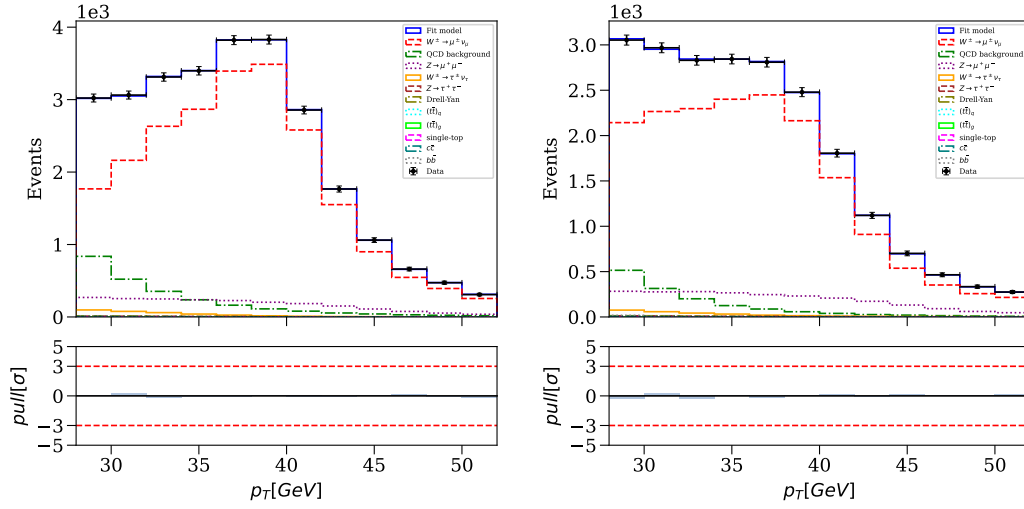
$$-\ln \mathcal{L} = \sum_i^n [-k_i \ln \beta_i \lambda_i + \beta_i \lambda_i] + \frac{1}{2} \sum_{i,k}^n (\beta_i - 1) C_{\beta_i, \beta_k}^{-1} (\beta_k - 1), \quad (6.17)$$

where the covariance matrix C_{β_i, β_k} can be obtained by dividing the original covariance by the outer product of λ_i :

$$C_{\beta_i, \beta_k} = \frac{C_{i,k}}{\lambda_i \lambda_k}. \quad (6.18)$$

The fit model depends on the bare-level p_T cross-sections for the signal, considered the only parameters of interest; and the background fractions (which are all fixed except for the QCD background) and the QCD background scaling parameter ω , which are considered nuisance parameters. The response matrix for positive and negative candidates is shown in Figure 6.4.

As mentioned earlier, it can be observed that for this particular analysis, the amount of bin migration is relatively low, which implies that the unfolding problem is relatively easy to solve. As a result, we anticipate that the unfolding process would have a minimal effect on the final results. This is why we have chosen to use the Beeston-Barlow method to account for the uncertainties in the data and simulation, alongside this simple unfolding method, given that we have a relatively low amount of bin migration.

FIGURE 6.4: Normalised response matrix for (left) W^+ and (right) W^- .FIGURE 6.5: Distributions of p_T for the (left) W^+ and (right) W^- candidates, integrated over the muon isolation range $0 < \mathcal{I} < 8$ GeV and pseudorapidity range $2.2 < \eta < 4.4$. The results of the differential cross-section fit are overlaid.

6.4 Fit results

After performing the background subtraction and unfolding process. We can verify that the fit model can describe the data. The fit results in the p_T projection are shown in Figure 6.5 and in logarithmic scale in Figure 6.6. The fit results for the $\mathcal{I}$ projection are shown in Figure 6.7 and in logarithmic scale in Figure 6.8. The isolation plots per p_T bin are shown in Appendix B.

The differential cross-section can be calculated with respect to the transverse momentum, as shown in equation 6.19. At present, we have used equal bin widths.

$$\frac{d\sigma}{dp_T} = \frac{n^{\text{obs.}} - n^{\text{bkg.}}}{\mathcal{L}\varepsilon\Delta p_T}, \quad (6.19)$$

where $n^{\text{obs.}}$ and $n^{\text{bkg.}}$ are the number of counts/events that correct for resolution for the signal

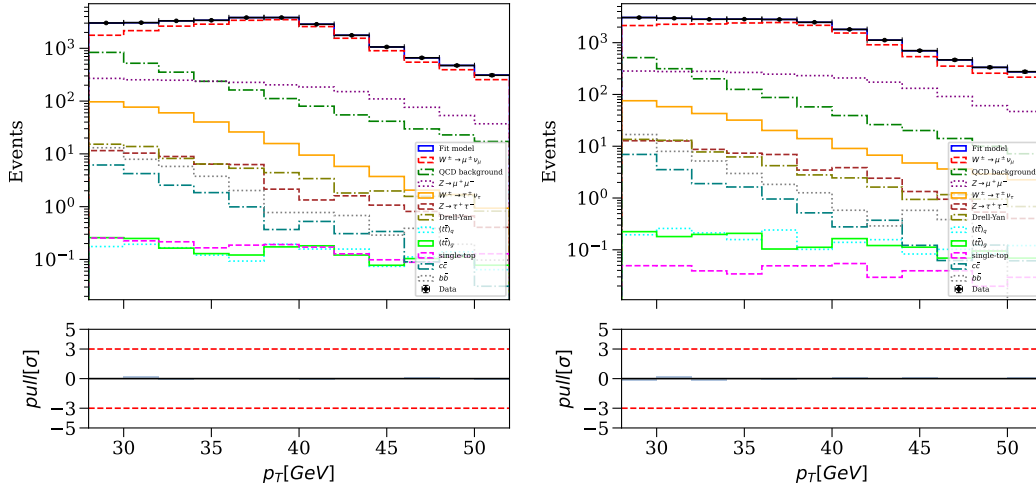


FIGURE 6.6: Distributions of p_T for the (left) W^+ and (right) W^- candidates in logarithmic scale, integrated over the muon isolation range $0 < \mathcal{I} < 8$ GeV and pseudorapidity range $2.2 < \eta < 4.4$. The results of the differential cross-section fit are overlaid.

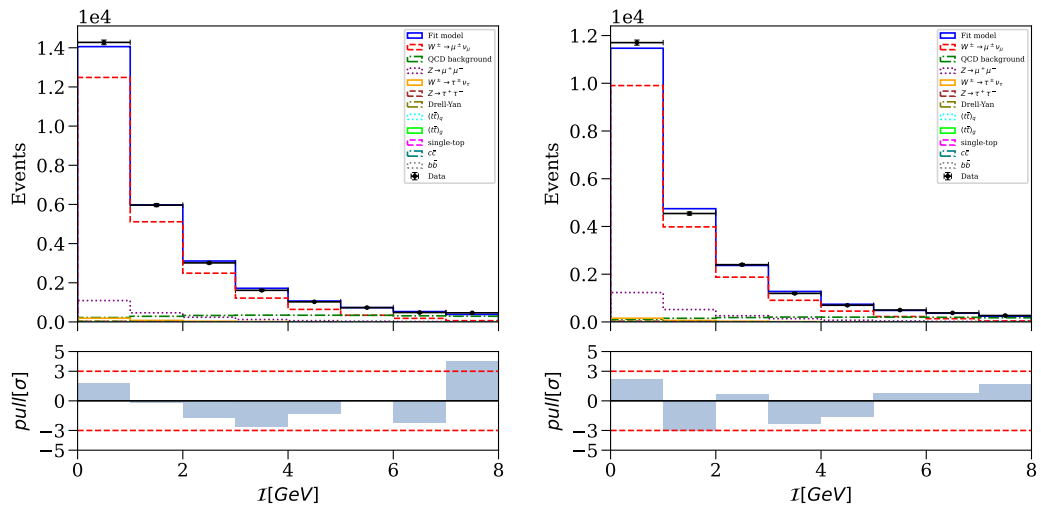


FIGURE 6.7: Distributions of $\mathcal{I}$ for the (left) W^+ and (right) W^- candidates, integrated over the muon transverse momentum range $28 < p_T < 52$ GeV and pseudorapidity range $2.2 < \eta < 4.4$. The results of the differential cross-section fit are overlaid.

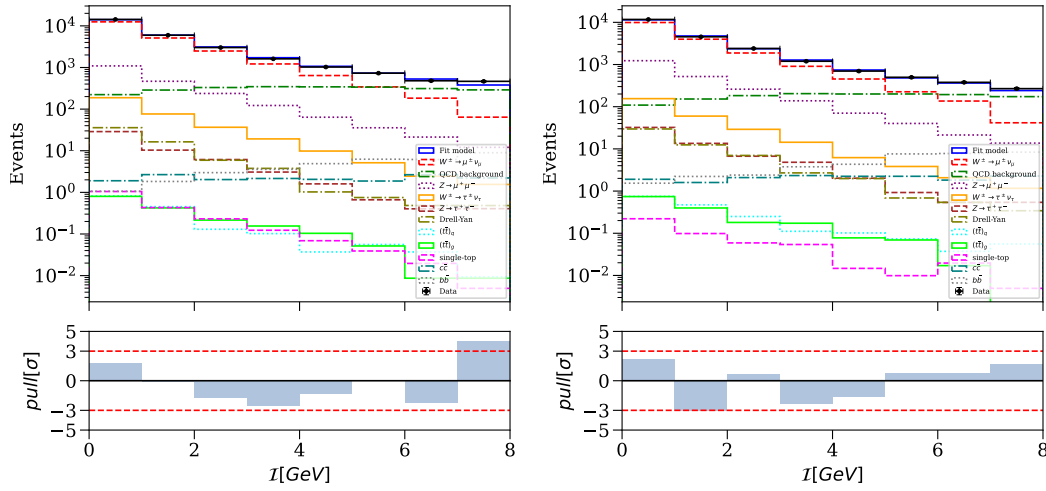


FIGURE 6.8: Distributions of $\mathcal{I}$ for the (left) W^+ and (right) W^- candidates in logarithmic scale, integrated over the muon transverse momentum range $28 < p_T < 52$ GeV and pseudorapidity range $2.2 < \eta < 4.4$. The results of the differential cross-section fit are overlaid.

and the backgrounds, respectively, Δp_T is the p_T bin width, ε represents the efficiencies per p_T bin.

The unfolded differential cross-section values in each bin and their corresponding covariance matrix are obtained from the Minuit fit [121]. In this step, the resulting covariance matrix only reflects the statistical fluctuations of the data and simulation. Statistical uncertainties are obtained by taking the square root of the diagonal elements of the covariance matrix. The total uncertainties also include systematic uncertainties and uncertainties from the integrated luminosity. The correlation matrix can be obtained by dividing the covariance matrix by the product of the square roots of the diagonal elements. The correlation matrix is useful for understanding the correlation between the differential cross-sections in different p_T bins. The correlation matrix can be viewed in Appendix C.

In addition to the statistical and systematic uncertainties, we also need to account for the uncertainty in the integrated luminosity. To estimate this uncertainty, we used an error ratio of 2%, as provided by the LHCb Luminosity group [102]. We obtain a fiducial cross-section value of

$$\begin{aligned}\sigma_{W^+ \rightarrow \mu^+ \nu_\mu} &= 301.62 \pm 3.00_{(\text{stat})} \pm 3.74_{(\text{syst})} \pm 6.03_{(\text{lumi})} \text{ pb}, \\ \sigma_{W^- \rightarrow \mu^- \bar{\nu}_\mu} &= 237.80 \pm 2.72_{(\text{stat})} \pm 2.67_{(\text{syst})} \pm 4.76_{(\text{lumi})} \text{ pb},\end{aligned}\tag{6.20}$$

where the first uncertainty represents the statistical uncertainty, the second is the systematic uncertainty, and the third comes from the integrated luminosity.

The differential cross-section values for both W^+ and W^- , along with their corresponding statistical errors per p_T bin, are presented in Table 6.1. The table is organised in such a way that the W^+ results are listed first, followed by the W^- results for easy comparison.

Bin index	$\frac{d\sigma}{dp_T}$ [pb/ GeV]	Error [pb/ GeV]	Bin index	$\frac{d\sigma}{dp_T}$ [pb/ GeV]	Error [pb/ GeV]
W^+			W^-		
0	11.8616	0.4375	12	14.8938	0.4734
1	14.6391	0.4554	13	15.7066	0.4776
2	17.7772	0.4811	14	15.6723	0.4732
3	18.9066	0.4998	15	16.0925	0.4764
4	22.5991	0.5503	16	16.5701	0.4830
5	23.8211	0.5745	17	14.5386	0.4675
6	17.1502	0.5149	18	10.5255	0.4171
7	10.4171	0.4224	19	6.0249	0.3358
8	5.8581	0.3330	20	3.5004	0.2758
9	3.4495	0.2797	21	2.3237	0.2316
10	2.6776	0.2530	22	1.6304	0.2053
11	1.6549	0.1974	23	1.4216	0.1758

TABLE 6.1: Differential cross-section results. The first 12 intervals cover the W^+ and the last 12 the W^- . The uncertainties are statistical only.

6.5 Systematic uncertainties

To account for systematic uncertainties in this analysis, variations are applied to the data by varying the default configuration settings and observing the effect on the differential cross-section values. This allows us to calculate the covariance matrices for each source of uncertainty.

The systematic uncertainties arise from various sources, such as the detector response, the theoretical models used in the analysis, and the calibration of the experimental setup. By calculating the covariance matrices, we can quantify the correlation between these sources of uncertainty and determine their impact on the final measurement. For this analysis, we classify our systematic uncertainties into two categories; theoretical and experimental.

Once the covariance matrices are obtained, they are propagated through the measurement process using appropriate statistical methods. In one approach, random sampling is performed by varying the parameters according to their uncertainties and the root mean square of the resulting variations is used to estimate the total uncertainty. This method effectively accounts for the contribution of each individual uncertainty to the overall uncertainty in the measurement. For example, this method can be used to quantify the impact of changing the number of isolation bins and the distribution of events between these bins. Alternatively, we can systematically vary key parameters by one standard deviation (1σ) or perform discrete variations to assess the impact of systematic uncertainties. For example, this approach can be used to determine the effect of a systematic variation, such as modifying the efficiency modelling function. The total uncertainty is then estimated by taking the maximum and minimum shifts observed during these variations. Using these approaches appropriately, we ensure that the statistical and systematic uncertainties are properly accounted for in the final result [122, 123].

By summing all the covariance matrices, we can fully account for all the uncertainties and correlations arising from the different sources of systematic uncertainty. It also enables us to identify

the sources of uncertainty that contribute the most to the final result and prioritise efforts to reduce their impact in future analyses. We presume that every source is independent to add these matrices. We divide our systematic uncertainties into different categories. A list of the categories of the differential cross-section systematic uncertainties is as follows:

- **Hadronic Background:** This category addresses uncertainties related to the modelling of hadronic background contributions. It includes:
 - **QCD background misidentification rate model:** Variations in QCD background misidentification rates are considered by testing 20 randomised models. This accounts for potential discrepancies in the way background events are identified and classified.
 - **Hadron species abundance:** Our hadronic background is made up of kaons, pions, and protons. We can vary the proportions of these hadrons in our background model. The uncertainty in the abundance of the hadron species is explored by varying the fractions of different hadron species. This is done by adjusting the fractions to 0.8, 1.0, or 1.2 for each parameter and evaluating all possible combinations of these variations to assess their impact on the background estimation.
- **Background subtraction fit:** This category examines the uncertainties related to the subtraction of background contributions from the measured data. Variations include:
 - **Muon isolation bins:** The impact of changing the number of muon isolation bins is assessed. Variations include [5, 12] bins, excluding 8 bins (which is the default), to evaluate how different binning choices affect the background subtraction process.
- **Efficiencies:** This category covers uncertainties related to the efficiencies of various detection and selection processes. Variations include:
 - **Trigger efficiency seed:** Variations in the seed parameter for trigger efficiency are tested to assess its impact on the measured cross-section (default is 0).
 - **Muon identification seed:** Similar to trigger efficiency, the seed parameter for muon identification efficiency is varied to understand its effect on identification accuracy (default is 0).
 - **Trigger efficiency details:** The impact of using a flat parameterisation instead of linear efficiency functions for the trigger is evaluated.
 - **Muon identification efficiency details:** The difference between using flat and linear efficiency functions for muon identification is analysed.
 - **Track Efficiency Details:** The impact of using a flat track efficiency function instead of a linear function is assessed.
- **Momentum smearing:** This category considers the uncertainties related to the smearing of momentum measurements. Variations include:

- **Momentum smearing seed:** The uncertainty is analysed by varying the seed parameter for momentum smearing to gauge its impact on the smearing effects (default is 0).
- **Curvature biases:** This category addresses uncertainties related to biases in the measurement of particle curvature. Variations include:
 - **Pseudomass seed:** The effect of varying the seed parameter for pseudomass parameterisation is evaluated to understand how curvature biases impact the analysis.
- **Isolation calibration:** This category covers uncertainties in the calibration of isolation measurements. Variations include:
 - **Signal isolation calibration offset:** The effect of altering the calibration offset for signal isolation with {upward, downward} adjustments by one sigma is evaluated.
 - **Hadronic background isolation calibration offset:** Similar variations are considered for the calibration offset of QCD background isolation by {upward, downward} adjustments by one sigma.

The systematic uncertainties are taken into account by summing all the covariance matrices of different individual sources of uncertainty. The resulting covariance matrix is then added to the statistical uncertainty covariance matrix to obtain the “total” covariance matrix of the uncertainty in the measurement. The total correlation matrix of systematic uncertainty can be seen in Figure 6.9, where again the correlation is defined as the covariance divided by the product of the standard deviations of the two variables. The figure shows the correlation state between W^+ and W^- concerning the combined systematic uncertainties. Generally, W^+ and W^- display a high correlation, though certain regions show a lack of correlation or even anti-correlation. This is attributed to certain uncertainties, like curvature biases, due to the curvature biases having the opposite effect for μ^+ and μ^- , leading to the “bluer” areas in the figure. For detailed correlation matrices of individual systematics, see Appendix C.

We show the breakdown of the systematic uncertainties for each variation subcategory by plotting the relative uncertainties of the systematic uncertainties in Figure 6.10. The relative uncertainty is defined as the standard deviation (uncertainty) divided by the central value of the differential cross-section for each p_T bin. A plot of the individual systematic uncertainties is shown in Figure C.14.

We note that, as mentioned in Section 5.2, the main source of systematic uncertainty comes from corrections for the curvature bias of the data using the pseudomass method, and more so at higher p_T , where the true curvature is smaller, and the underlying true cross-section is smaller. Again, these curvature corrections are not perfect because of the lack of real data sample statistics to apply the corrections, especially in the high η region. Also, we observe that the background is more important at low p_T , which makes perfect sense given that the background is steeply falling versus p_T .

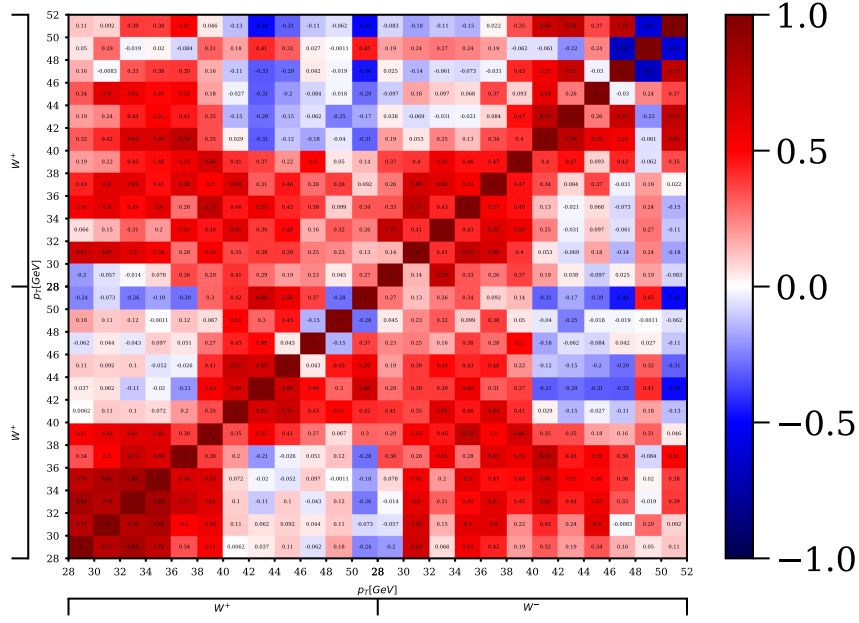


FIGURE 6.9: Correlation matrix corresponding to the total systematic uncertainty. The impact of pseudomass corrections is evident at higher p_T values, where the correlation pattern shifts from positive to negative.

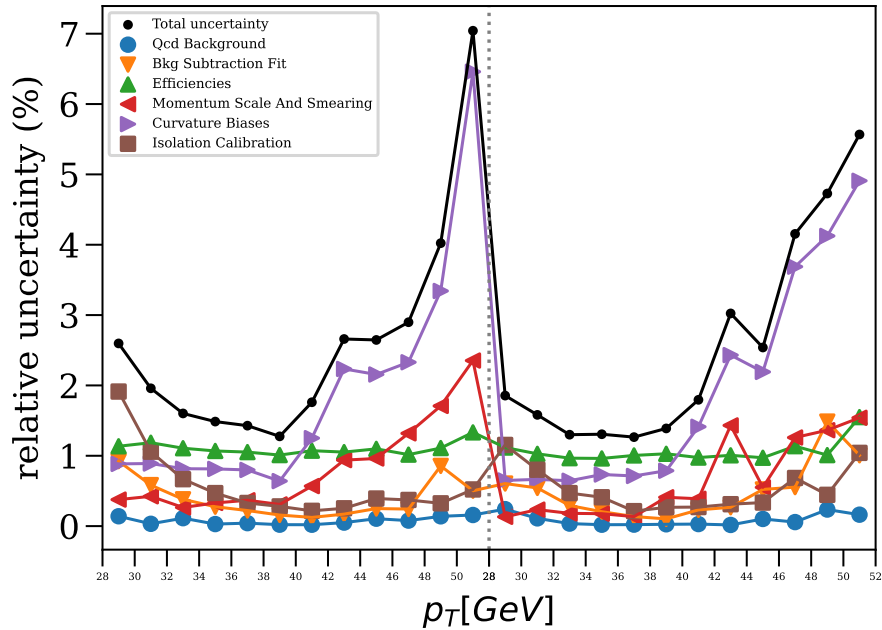


FIGURE 6.10: Relative systematic uncertainties on the differential cross-sections. The first and last 12 intervals correspond to W^+ and W^- , respectively. The primary source of uncertainty is evidently the curvature biases, particularly at higher p_T . The pseudomass approach seeks to mitigate these biases.

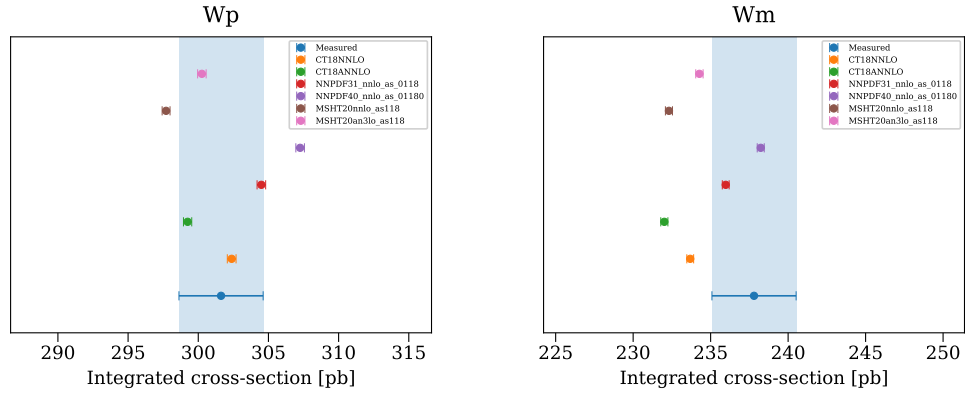


FIGURE 6.11: Integrated cross-sections of (left) W^+ and (right) W^- are compared with the theoretical predictions from the MCFM generator using various PDF sets.

The other sources of systematic uncertainty are clearly relatively smaller compared to the pseudomass curvature bias corrections, as seen in Figure C.14.

6.6 Comparison with predictions

We compare the measured cross-section results with the theoretical predictions of the generator MCFM [103] using various PDF sets. The comparison is shown in Figure 6.11. We can see that the measured cross-section values are in good agreement with the theoretical predictions from the MCFM generator using various PDF sets. Please note that the uncertainties presented in the predictions are solely statistical and do not encompass theoretical uncertainties. Addressing these will be a focus of future analysis. Furthermore, the measured cross-section values are presented at the bare level, which means that they include QED radiation effects. MCFM predictions are therefore reduced by 3% to adjust for QED radiation effects/corrections and align them with the measured cross-section values. This adjustment factor was determined by comparing the p_T distributions at the bare and born levels using PYTHIA simulations. The factor is simply the ratio of events with born kinematics in the fiducial region to those with bare kinematics in the same region. This approach for separating QED radiation effects/corrections is not flawless, and a more precise method would involve running PYTHIA with QED corrections enabled and disabled to acquire accurate cross-section values. We plan to implement this improvement in the near future. However, the current method, which also relies on simulations from PYTHIA, is a reasonable approximation. Additionally, it should be noted that we have not yet compared the differential cross-section with theoretical predictions, as this comparison is performed in Chapter 7.

Chapter 7

W boson mass measurement

“The most beautiful thing we can experience is the mysterious. It is the source of all true art and science.”

– Albert Einstein

The final step of this analysis involves determining the mass of the W boson by fitting it to the unfolded differential cross-sections obtained in Chapter 6. The fitting process is carried out by comparing the observed data with templates generated from simulations of the W processes. We use a single sample of W events generated using PYTHIA at tree-level, which lacks the accuracy required for our purposes¹. The hypothesised m_W value can be changed simply by assigning weights to this single sample. To improve template accuracy, reweighting techniques using predictions with higher formal perturbative accuracy can be used to adjust the simulated events to match the observed data. The templates are then fitted to the differential cross-section distribution to extract the optimal mass of the W boson, and the residual limitation in the accuracy of the model is mitigated by tuning the strong coupling and the model of non-perturbative effects. Hence, we have floating parameters such as the W boson mass and α_s .

We discuss the fit procedure in Section 7.1, which includes the reweighting of the simulated template events to match the distributions of the higher-accuracy predictions and to be comparable to the data. We will then present the results of the Z boson differential cross-section fit in Section 7.2, which is used to obtain strong coupling and non-perturbative parameters required for the W boson mass fit. Next, we will discuss the fitting of the W boson mass and the systematic uncertainties associated with the model in Section 7.3 and Section 7.4 respectively. Finally, we conclude by reviewing the various cross-checks implemented to validate our methodologies and measurements, as detailed in Section 7.5.

¹This follows the same configuration used in the full simulation by LHCb.

7.1 The fit procedure

The fit for m_W (and Z cross-section) is based on simulated events produced using PYTHIA 8 [124]. Different m_W hypotheses are emulated by assigning weights to events based on a relativistic Breit–Wigner function with a width dependent on mass [125]. The Breit–Wigner function acts as a probability distribution function representing the probability of a particle having a specified mass as given by

$$f(M) = ((M^2 - m_W^2)^2 + M^4 \Gamma^2 / m_W^2)^{-1}, \quad (7.1)$$

where M is the mass of the particle, m_W is the central mass or resonance mass of the particle (e.g. the W boson) and Γ is its natural width, which quantifies the spread of the mass distribution and is related to the particle’s lifetime.

As mentioned in Section 2.3.5, we use the DYTURBO software (version 1.1) to generate high-precision predictions for the Drell–Yan process [43] (at NNLO for the unpolarised cross-section and at NLO for the angular cross-section). This is a crucial step in the analysis, as the predictions from DYTURBO are used to reweight the simulated events to improve their accuracy. The simulated events are produced at the leading-order using PYTHIA, which lacks the precision required for the analysis [21].

We employ a five-dimensional weighting technique to align the prediction of the model with the results of the unfolded data. This 5-dimensional weighting technique enables us to adjust the model parameters to reduce the discrepancies between the model and the data [21]. The five variables used to refine the fit model are as follows:

- p_T ; the dilepton’s transverse momentum.
- $|y|$; the dilepton’s absolute rapidity.
- M ; the invariant mass of the lepton pair.
- $\cos \theta$; the cosine of the lepton’s polar angle in the Collins–Soper frame.
- ϕ ; the lepton’s azimuthal angle in the Collins–Soper frame,

where the Collins–Soper frame was introduced and illustrated in Figure 2.11.

The parameters of the model are fine-tuned by minimising the log-likelihood function. Switching to a log-likelihood function based on the Beeston–Barlow lite method [113] was an option, but for the current analysis, the simplicity of having a larger number of template events compared to the statistical precision of the differential cross-section measurement was sufficient, given the small dataset used. For a recap of the Beeston–Barlow lite method, refer to Section 6.1.

The fit procedure is blinded to minimise biases. The blinding process involves randomly adjusting the mass of the W boson within the range $[-0.5, 0.5]$ GeV. When the minimisation asks for a particular value of m_W , the code actually uses a shifted value in the fit model.

The full LHCb simulation sequence does not readily permit the substitution of the event generator. Therefore, a higher-order generator must be emulated by adjusting the weights of the template events to align with their predictions. Keep in mind that we are not truly obliged by the approach taken in LHCb; it is simply more practical to adopt that method, consistent with previous analyses such as those found in the m_W measurement at LHCb [21]. Fortunately, we can improve the accuracy by reweighting these events. This process is divided into two stages: QCD reweighting and QED final-state radiation (FSR).

7.1.1 QCD reweighting

Equation 2.21 demonstrated the factorisation of the Drell–Yan differential cross-section into an unpolarised cross-section (σ^{unpol}) and an angular factor, which is a sum of multiple angular terms. Consequently, the QCD-related weights split into an unpolarised cross-section weight and an angular cross-section weight.

The computation starts by running the precise version of PYTHIA used in the LHCb simulation to generate predictions at the generator level, which constitute the denominator of the weights. Equivalent predictions are generated using a preferred higher-order tool, which forms the numerator. Predictions at the generator level are considerably quicker than the complete LHCb simulation, where the primary hindrance is the detector simulation. Based on feedback from senior members of the team involved in other electroweak measurements in the LHCb experiment, DYTURBO was identified as the best program capable of accurately predicting angular coefficients [21]. Consequently, we employ it to generate predictions for both the unpolarised and angular cross-sections.

The cross-sections depend on the vector boson mass M , rapidity y , and transverse momentum p_T^V ; denoted as $\frac{d^3\sigma}{dM dy dp_T^V}$. The weighting factor for the unpolarised cross-section is the ratio between DYTURBO's $\frac{d^3\sigma}{dM dy dp_T^V}$ and the leading-order $\frac{d^3\sigma}{dM dy dp_T^V}$ from PYTHIA.

$$w_{\text{unpol}} = \frac{\left(\frac{1}{\sigma_{\text{DYTurbo}}}\right) \left(\frac{d^3\sigma_{\text{DYTurbo}}}{dM dy dp_T^V}\right)}{\left(\frac{1}{\sigma_{\text{Pythia}}}\right) \left(\frac{d^3\sigma_{\text{Pythia}}}{dM dy dp_T^V}\right)}. \quad (7.2)$$

Note that the denominator of the weight relies on the cross-section as well as the differential cross-section. This is because the cross-section is the quantity that is used to normalise the templates, and the differential cross-section is the quantity that is used to fit the templates to the data.

DYTURBO generates 3D histograms for p_T , $|y|$, and M . Thus, trilinear interpolation is employed to determine the weight for each event at the corresponding p_T , $|y|$, and M coordinates.

The angular weight, w_{ang} , is computed in a similar way, using the angular coefficients of the angular component of the differential cross-section equation 2.21.

$$w_{\text{ang}} = \frac{1 + \cos^2 \theta + \frac{1}{2} A_0^{\text{DyTurbo}} (1 - 3 \cos^2 \theta) + \dots}{1 + \cos^2 \theta + \frac{1}{2} A_0^{\text{Pythia}} (1 - 3 \cos^2 \theta) + \dots} \quad (7.3)$$

The angular coefficients also depend on the vector boson mass M , the rapidity y , and the transverse momentum p_T^V ; A_0^{DyTurbo} represents the first angular coefficient of DYTURBO, while A_0^{Pythia} denotes the first angular coefficient obtained from PYTHIA. There are eight such coefficients (from A_0 to A_7), and we compute the weight for each one. However, in practice, the last 3 coefficients (A_5, A_6, A_7) are equal to zero at NLO [126].

Event generators deliver differential cross-section predictions, from which angular coefficients must be derived using a projection method, as described in [127]. Essentially, this method involves averaging the differential cross-section over θ and ϕ , applying the corresponding spherical-harmonic projection function (leveraging the orthogonal nature of spherical harmonics to isolate a specific term), and then dividing by the unpolarised cross-section. We have many systematic uncertainties associated with QCD reweighting, which are discussed in Section 7.4.

7.1.2 QED final-state radiation

QED final-state radiation leads to a reduction in energy in the boson system, lowering the bare-level muon p_T spectrum in $W^+ \rightarrow \mu^+ \nu_\mu$ compared to the Born-level muon p_T . This effect can be corrected by applying a weight to the muon p_T spectrum, which is derived from three different generators: PHOTOS, HERWIG, and PYTHIA. As there was no significant disparity between these generators, we use the average of the three weights. The individual weights of PYTHIA, PHOTOS, and HERWIG are preserved to evaluate the systematic uncertainty arising from the selection of the model, which is then propagated to the parameter of interest. The uncertainties related to the QED FSR are discussed in Section 7.4.

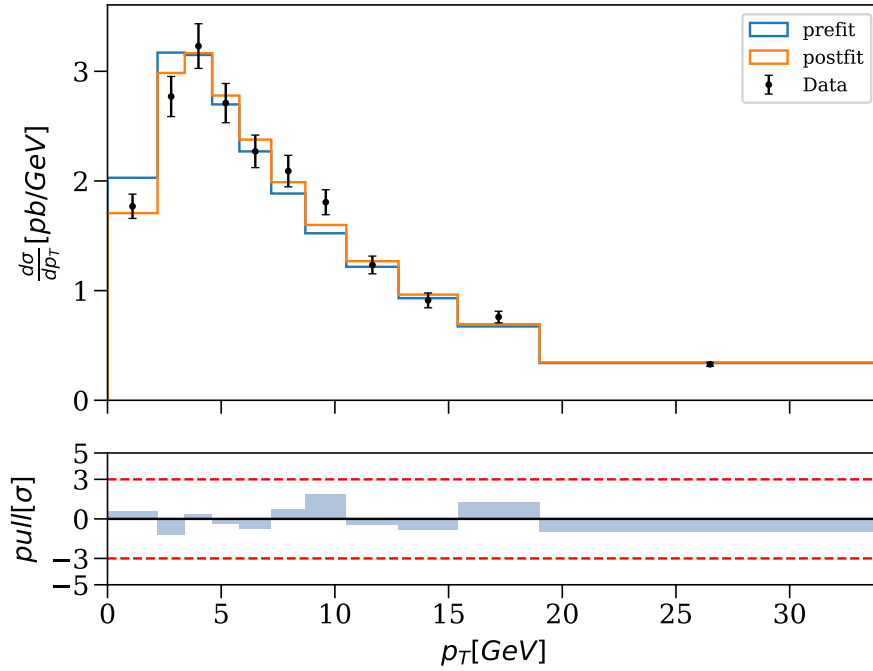
7.2 Strong coupling and non-perturbative parameter fit

In addition to the W boson mass fit, we fit the Z boson cross-section as well. The data of the Z boson cross-section are obtained from the LHCb experiment at 5.02 TeV [128], and the fit is performed using the same methodology as the W boson fit. The objective of this fit is to estimate α_s and g (which are floated and varied during this fit), as well as to confirm that the model appropriately represents the Z p_T data. The g -factor governs the non-perturbative segment of the cross-section and is used to adjust the cross-section to align with the data.

The results of the Z boson fit are shown in Table 7.1 and Figure 7.1, with $\chi^2/\text{ndf} = 11.3/9$. The starting values of α_s and g are 0.11 and 1.0 respectively. The starting g value corresponds to the DYTURBO default and the starting value of α_s is close to the PDG value cited in Ref. [129].

We can then use the results of the Z boson fit to constrain the W boson fit. We choose to fix the α_s value to the value obtained from the Z boson fit and float the g factor. This is because the simplified non-perturbative model does not include possible flavour-dependent effects, which are expected to be different for W and Z productions. Cross-checks were performed using

Parameter	Value	Error
α_s	0.1113	0.0050
g	1.9267	0.4148

TABLE 7.1: Results of the α_s and g fit to the Z data.FIGURE 7.1: Differential cross-section results for Z boson production as a function of p_T from Ref. [128]. The results of the α_s and g fit are overlaid.

different strategies for using the Z boson fit, as well as the choices of fixing or floating α_s and the g factor. The results of these cross-checks are shown in Section 7.5. A more direct check would be to compare the luminosity inferred from counting Z events and the NNLO cross-sections with the measured luminosity from the LHCb luminosity group [102]. We could also perform a combined measurement of the cross-sections W and Z at some point.

7.3 Fitting the W boson mass

The total uncertainty of the differential cross-section measurement, including statistical and systematic uncertainties, is represented using a covariance matrix that combines these uncertainties linearly. In addition to the mass of the W boson, the g -factor is allowed to freely float in the fit. As mentioned above, the resulting mass value is blinded to avoid any bias. We note that we fit both W boson charges simultaneously to obtain m_W . The measured differential cross-section, with the results of the W boson mass fit (with systematic uncertainties) overlaid, is shown in Figure 7.2.

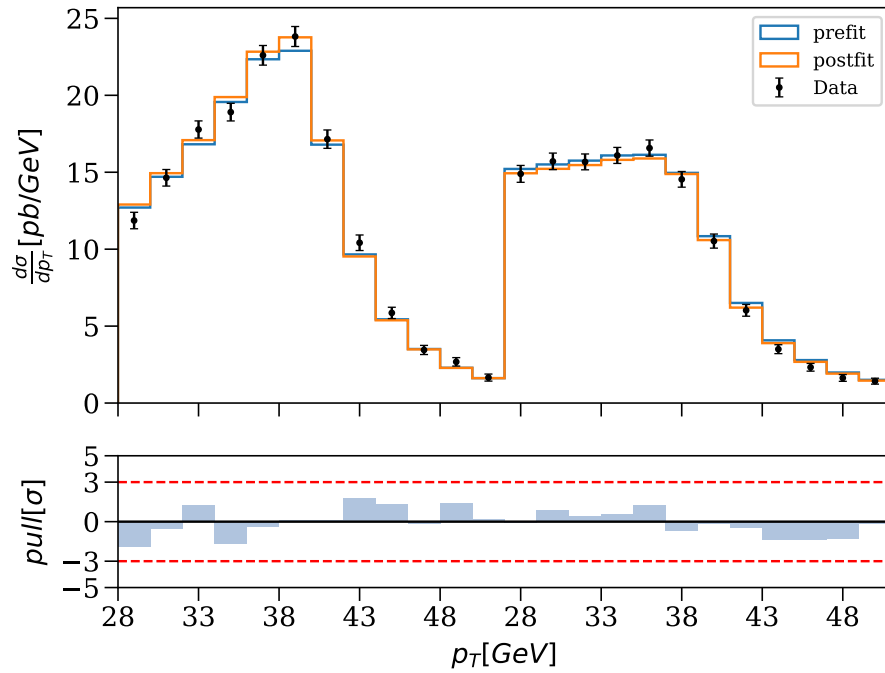


FIGURE 7.2: Measured differential cross-section, with the W boson fit results overlaid.

The m_W value we obtain for both muon charges is

$$m_W = 80496 \pm 100_{(\text{stat})} \text{ MeV (blinded)}. \quad (7.4)$$

where the uncertainty corresponds to statistical uncertainties only. The rationale behind merging both statistical and systematic uncertainties is that it allows us to perform a single fit for m_W , by summing the statistical covariance matrix with all the systematic covariance matrices of the various changes. We expect this to be the same method that theorists would typically use to compare their predictions to our measurement. We tried the more labour-intensive approach of fitting m_W for each systematic variation separately, but the study is still ongoing. The result of the W boson mass fit results with systematic uncertainties is shown in Table 7.2, with a $\chi^2/\text{ndf} = 29/20$. Statistical uncertainty dominates our measurement, which is not surprising

Parameter	Value	Error
g	0.9218	0.5807
m_W [MeV]	80355	128
Charge asymmetry	0.1181	0.0063
α_s	0.1113	0.0050

TABLE 7.2: Results of the m_W fit including the systematic uncertainties on the data.

given the limited size of our dataset.

7.4 Systematic uncertainties on the W mass fit model

As mentioned earlier, we are using DYTURBO to perform high-precision calculations of cross-sections and distributions for Drell-Yan processes, which involve the production of a lepton-antilepton pair via the annihilation of a quark and an antiquark. The software includes the computation of electroweak corrections, which are essential for accurate predictions and comparisons with experimental data.

The software requires us to provide an input file with many different control inputs. Some of the most important ones include:

- The process (W^+ , W^- , or Z)
- The center-of-mass energy of the collision (5.02 TeV for this analysis).
- The PDF set (choices of NNPDF31_nlo (default), NNPDF40_nlo, CT18NNLO or MSHT20nlo).
- The PDF set member number ([0,100], [0,100], [0, 58], [0, 65] for NNPDF31_nlo (default), NNPDF40_nlo, CT18NNLO and MSHT20nlo, respectively).
- The order of the perturbative calculation (N3LO, NNLO, NLO, etc.).
- The mass of the W boson (80.385 GeV by default).
- Inverse Bessel transform cutoff value (b_{lim}) (default is 3.0, and we vary it to 1.0 or 2.0)
- QCD scale factor (renormalisation, factorization, resummation scales) (default is 1.0, and we can vary them to 0.5 or 2.0 according to the 7-point variation strategy, which is explained later)

Some of the control inputs are fixed, such as the centre-of-mass energy of the collision. However, other inputs can be varied to study their impact on the results. For example, we can vary the renormalisation and factorisation scales, the PDF set, and the mass of the W boson to estimate the systematic uncertainties associated with these inputs. Some of the variations can be perceived as cross-checks as they have little impact on the results. However, other variations can have a more significant impact on the results, and it is essential to study them carefully to understand their impact on the final measurement. These then are the variations that we will focus on in this section, and handled as systematic uncertainties.

7.4.1 Perturbative accuracy

Theoretical uncertainties originate from multiple sources, including the renormalisation and factorisation scales of the QCD scale factor (denoted $k_{\mu,\text{ren}}$ and $k_{\mu,\text{fac}}$, respectively). These scales can take values of {0.5, 1.0, 2.0}. We adjust these values following the 7-point variation strategy [127], which excludes the combinations {0.5, 2.0} and {2.0, 0.5} for $\{k_{\mu,\text{ren}}, k_{\mu,\text{fac}}\}$. This reduces our selection from the 9 possible pairs in a Cartesian product to 7 valid combinations.

We also have uncertainties from QED FSR. By default, we use the average of 3 generators: PHOTOS, HERWIG, and PYTHIA. The systematic is based on the variation of these generators.

Sub-category	Error [MeV]	Number of additional fits
QED	11	3
QCD scales	13	6

TABLE 7.3: Results of all theory variations and corresponding systematic uncertainties.

7.4.2 PDF uncertainties

Each PDF set includes several variant members. For example, the NNPDF31_nnlo_as_0118 PDF set contains 101 members, labelled from 0 to 100. These members are Monte Carlo replicas designed to represent statistical uncertainties through random data point variations. PDF variations are calculated by varying the PDF set member number and calculating the difference between the central value and the varied value. The rms of these variations is then taken to estimate the systematic uncertainty associated with the PDF set. We currently deal with the default PDF sets of NNPDF31_nnlo_as_0118. The results of the PDF variations are shown in Table 7.4. We also deal with other PDF sets for a particular α_s value; however, we only use them as a cross-check, without varying the PDF set members.

Name	Method	Central value [MeV]	Number of additional fits	Error [MeV]
NNPDF 3.1	rms	80355	100	13

TABLE 7.4: Results of changing the default PDF set members.

We plan to vary the members for other PDF sets, not just NNPDF31_nnlo_as_0118, in the future. We also plan to get the m_W result by averaging the results from using all the PDF sets (again using member 0 as the default).

The m_W value we obtain using systematic uncertainties is

$$m_W = 80355 \pm 128_{(\text{stat})+(\text{syst})} \pm 21_{\text{theory}} \text{ MeV (blinded)}, \quad (7.5)$$

where the uncertainty arises from a combination of both statistical and systematic uncertainties, along with theoretical uncertainties derived from QED and QCD variations.

To give some context, the blinded result of this analysis can be seen alongside the results of other experiments (including the earlier LHCb measurement) in Figure 7.3. Note that this is just shown to try and put the currently blinded result in context. For cross-checks related to the m_W results, see the variations in Section 7.5.

7.5 Cross-checks

As Richard P. Feynman once said, “The first principle is that you must not fool yourself, and you are the easiest person to fool”. The cross-checks are essential to ensure that the results are

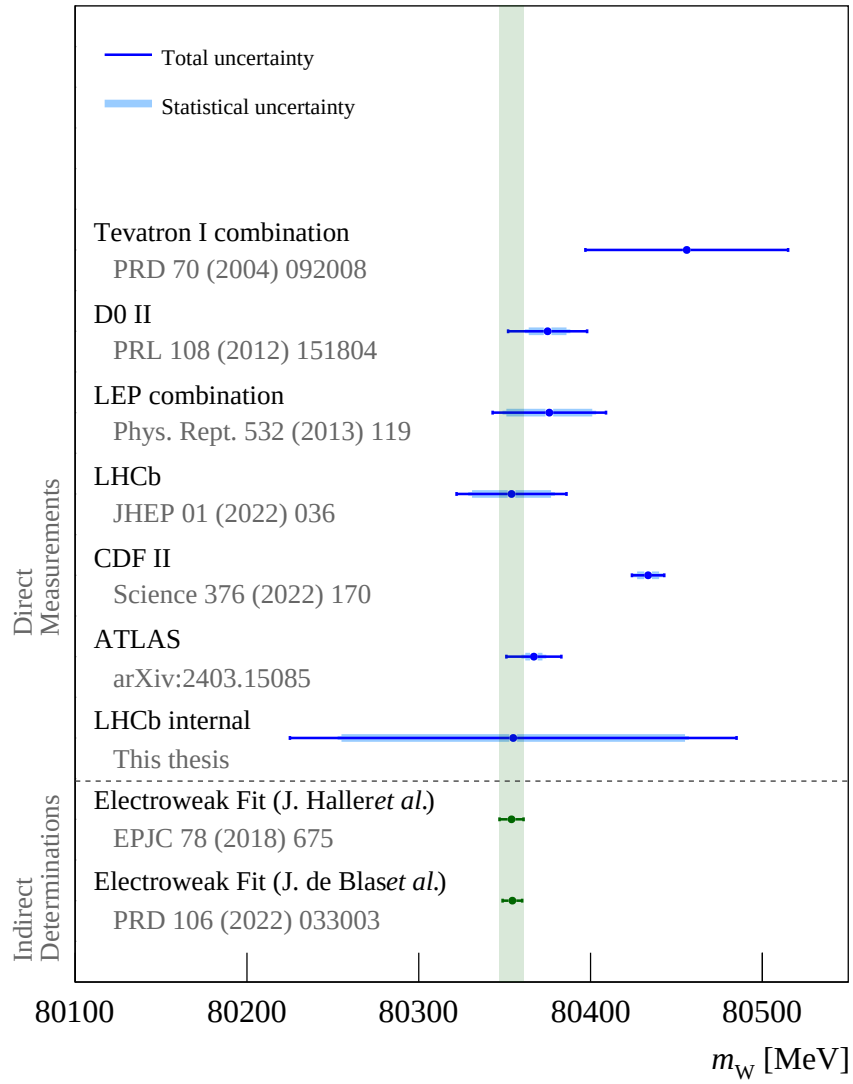


FIGURE 7.3: A chart illustrating the W mass measurements from different experiments alongside theoretical determinations. Note that the measurement presented in this thesis is blinded.

robust and reliable.

As mentioned in Sections 6.5 and 7.4, we have performed several cross-checks to ensure the robustness of the measurement. In this chapter, we present the results of the cross-checks for the scalar outcome, the W boson mass.

Most of these cross-checks are related to the DYTURBO calculations and involve varying the control inputs to study their impact on the results. The control inputs that we have varied include the PDF set, the order of the perturbative calculation, the b_{lim} value, and the electroweak scheme.

7.5.1 Changing the PDF set

We have tried using four different PDF sets for the W boson mass measurement:

- NNPDF31_nnlo_as_0118 (default)
- CT18NNLO
- MSHT20nnlo_as118
- NNPDF40_nnlo_as_0118

The default PDF set is NNPDF31_nnlo_as_0118, and we have used the other PDF sets as cross-checks. Choosing a PDF set is arbitrary since each one models parton distributions differently, using distinct assumptions and methodologies. We intend to use the average m_W value across all PDF sets as our final result, consistent with the other m_W analysis by LHCb [21].

Name	Central value [MeV]	Stat. unc.[MeV]	χ^2	Variation [MeV]
NNPDF 3.1 nnlo	80355	128	29	0
MSHT20nnlo	80364	127	31	9
CT18NNLO	80355	128	27	-0
NNPDF 4.0 nnlo	80350	128	29	-5

TABLE 7.5: Results of changing the PDF set.

More detailed results on these particular variations are shown in Tables 7.2, D.1, D.2, D.3.

7.5.2 Changing the scheme

These schemes refer to different choices for the renormalisation and factorisation scales (μ) used in the QCD calculations. The scale μ is a critical parameter in QCD because it determines at which energy scale the strong coupling constant, α_s is evaluated. The choice of scale can significantly impact the precision of theoretical predictions. The default scheme is 2, and we have used the other scheme 1 as a cross-check. As a reference, these are the schemes:

- **Scheme 1:** $\mu^2 = m_l^2$
- **Scheme 2:** $\mu^2 = m_l^2 + p_T^2$

Scheme 1 is suitable for cases where the transverse momentum of final-state particles is fairly low. In contrast, scheme 2 incorporates both the invariant mass and transverse momentum of the lepton pair, rendering it more appropriate for our dataset.

We do not vary the resummation scale, as using the scheme 2 in this case is equivalent to the default scheme 1.

Name	Central value [MeV]	Stat. unc. [MeV]	χ^2	Variation [MeV]
2	80355	128	29	0
1	80354	128	29	-2

TABLE 7.6: Results of changing the scheme.

7.5.3 Changing the inverse Bessel transform cut-off value

Since we are dealing with the inverse Bessel transform when using the DYTURBO calculations, we need to provide a cut-off value for the inverse Bessel transform to avoid the Landau pole. This Landau pole is a singularity in the QCD coupling constant at high energies, and it can cause the perturbative expansion to break down. This cutoff value is called the “ b_{lim} ” value and can be varied to study its impact on the results. The default value is 3.0 GeV. We have used 1.0 and 2.0 GeV as cross-checks.

Name	Central value [MeV]	Stat. unc. [MeV]	χ^2	Variation [MeV]
3.0	80355	128	29	0
1.0	80361	128	30	5
2.0	80356	128	29	1

TABLE 7.7: Results of changing the b_{lim} value.

7.5.4 Changing the order of the unpolarised component

By default, the unpolarised component is calculated at NNLO (next-to-next-to-leading order). We have used N3LO (next-to-next-to-next leading order) as a cross-check. This is only done at the level of the unpolarised component. As for the angular component, we only stick to using NLO, as NNLO is prohibitively expensive. Furthermore, as mentioned earlier, the angular component is already quite well modelled at fixed NLO [21].

Name	Central value [MeV]	Stat. unc. [MeV]	χ^2	Variation [MeV]
NNLO	80355	128	29	0
N3LO	80350	130	32	-6

TABLE 7.8: Results of changing the order of the unpolarised component.

7.5.5 Changing the W/Z fit strategy

Additionally, we investigated the impact of the W/Z fit strategy on the results, in which we can choose to fit the W/Z cross-section data separately or simultaneously. The Z cross-section data

is static and obtained from the recent LHCb measurement in Ref. [128] for Z production at 5.02 TeV.

Note that we could also choose to measure the Z cross-section as well, but we have chosen not to do so in this analysis, and we have only used the Z cross-section data to constrain the W boson mass fit.

We have 3 strategies relating to the W/Z fit strategy:

- Z fit off: Fit the W data only.
- Spectating: Fit the Z data but do not use the fit result parameters for the W fit (equivalent results to Z fit off).
- Sequential: Fit the Z data and use the fit result parameters for the W fit (default).

In addition, we can choose to float or fix the α_s value, as well as the g value. By default, we choose to use the sequential strategy, fix the value of α_s and float the value of g . Doing the Cartesian product of these combinations, we have 12 different strategies, which we have used as cross-checks. However, since technically the “off” and “spectating” strategies produce identical results, we actually have just 8 strategies/combinations. They are therefore presented in Table 7.9. As a reminder, the starting values of α_s and g are 0.11 and 1.0 respectively.

Name	Central value [MeV]	Stat. unc. [MeV]	χ^2	Variation [MeV]
sequential, Fixed α_s , Floating g ,	80355	128	29	0
sequential, Fixed α_s , Fixed g ,	80277	121	32	-78
sequential, Floating α_s , Floating g ,	80359	128	29	3
sequential, Floating α_s , Fixed g ,	80291	121	31	-64
off, Fixed α_s , Floating g ,	80358	128	29	3
off, Fixed α_s , Fixed g ,	80364	120	29	8
off, Floating α_s , Floating g ,	80359	128	29	3
off, Floating α_s , Fixed g ,	80366	120	29	10

TABLE 7.9: Impact of modifications to the W/Z fit strategy.

Additionally, we provide another table showing the effects of changing these strategies on the measurement of the W boson mass without including systematic uncertainties from the measurement of the cross-section in Table 7.10.

[htbp]				
Name	Central value [MeV]	Stat. unc. [MeV]	χ^2	Variation [MeV]
sequential, Fixed α_s , Floating g ,	80496	100	45	0
sequential, Fixed α_s , Fixed g ,	80419	94	49	-77
sequential, Floating α_s , Floating g ,	80496	100	44	-0
sequential, Floating α_s , Fixed g ,	80431	94	47	-66
off, Fixed α_s , Floating g ,	80496	100	44	-0
off, Fixed α_s , Fixed g ,	80494	93	44	-3
off, Floating α_s , Floating g ,	80496	100	44	-0
off, Floating α_s , Fixed g ,	80493	93	44	-4

TABLE 7.10: Impact of modifications to the W/Z fit strategy (without propagating the differential cross-section systematic uncertainties).

We observe notable deviations from the default values presented in Tables 7.9 and 7.10. Further discussion on these differences will be provided in sub-section 7.5.8.

7.5.6 Changing the overflows treatment

We shift our focus to the variations to determine the W cross section, as discussed in Chapter 6. One cross-check we performed is related to how we define the overflow and underflow of our truth cross-section histograms that we fit as part of our background subtraction and unfolding procedure. This was mentioned in Section 6.3. We propose four options:

- uu: Vary the underflow and overflow bins up.
- dd: Vary the underflow and overflow bins down.
- ud: Vary the underflow bin up and the overflow bin down.
- du: Vary the underflow bin down and the overflow bin up.

Name	Central value [MeV]	Stat. unc. [MeV]	χ^2	Variation [MeV]
uu	80356	128	29	0
dd	80355	128	29	-2
ud	80356	128	29	-1
du	80355	128	29	-1

TABLE 7.11: Results of changing the overflows treatment.

7.5.7 Changing the QCD Physics weights

Another cross-check we performed pertaining to the W cross-section is related to the QCD physics weights. We have three variables that we vary individually:

- The weight of the unpolarised component.
- The weight of the angular component.
- The mass of the W boson.

The weights for both unpolarised and angular components can be toggled on or off, and the mass of the W boson can be adjusted by ± 100 MeV to assess how the results respond to small variations in the assumed W boson simulation modelling mass.

These modifications are applied to the simulation data that generate our signal distributions utilised in the template fitting process. It is important to mention that the standard m_W mass used in these simulations is taken from PYTHIA, with a value of 80.385 MeV.

Name	Central value [MeV]	Stat. unc. [MeV]	χ^2	Variation [MeV]
Unpolarised Weights=False, Angular Weights=False, $m_W = 80.385$ GeV,	80355	128	29	0
Unpolarised Weights=True, Angular Weights=False, $m_W = 80.385$ GeV,	80347	128	29	-9
Unpolarised Weights=False, Angular Weights=True, $m_W = 80.385$ GeV,	80371	127	28	16
Unpolarised Weights=False, Angular Weights=False, $m_W = 80.485$ GeV	80352	128	29	-4

TABLE 7.12: Results of changing the physics weights.

7.5.8 Comments on the cross-checks

By comparing the results of the cross-checks with the default results, we can see that some variations have a significant impact on the results, while others have a negligible impact.

Effect	Impact
Changing the PDF set (subsection 7.5.1)	Single digits in units of MeV
Changing the scheme (subsection 7.5.2)	Single digits in units of MeV
Changing the b_{lim} value (subsection 7.5.3)	Single digits in units of MeV
Changing the perturbative order (subsection 7.5.4)	Single digits in units of MeV
Changing the W/Z fit strategy (subsection 7.5.5)	Up to ~ 80 MeV
Changing the overflows treatment (subsection 7.5.6)	Single digits in units of MeV
Changing the QCD physics weights (subsection 7.5.7)	Up to small double digits in units of MeV

TABLE 7.13: Summary of the effects of different cross-check variations on the measurement. The anomalous result, which deviates significantly from the expected values, is highlighted in light red for emphasis.

From Table 7.13, we can conclude that we have passed almost all cross-checks, apart from one; involving the changing of the W/Z fit strategy. We can observe in Table 7.9 that there is a discrepancy between the g -factor fit results from the Z and those from the W fit. This discrepancy is further supported by comparing the g values of the Z fit in Table 7.1 with those of the W fit presented in Table 7.2. This issue requires additional investigation and remains unresolved at the time this thesis was written.

Chapter 8

Summary and Outlook

“I can’t give you brains, but I can give you a diploma.”

– The Wizard of Oz

This thesis has covered the development and application of a novel technique for measuring the differential cross-section of the $W \rightarrow \mu\nu_\mu$ process using the LHCb detector, as well as the W boson mass. Using simulated template fits, background subtraction techniques, and unfolding methods that correct for detector biases, we managed to extract the differential cross-sections in bins of transverse momentum (p_T). Subsequently, these unfolded yields were used in the W boson mass (m_W) fit process, where they were fitted against reweighted Breit–Wigner simulation templates.

Uncertainties arising from different aspects of the analysis have been evaluated and propagated to the final result. Leveraging the unfolded distribution offers theorists a framework to test and compare their diverse models against the data we have measured. This research highlights the LHCb detector’s capability in executing high-precision electroweak-physics measurements, an area traditionally dominated by other experiments at the LHC, as well as LEP and the Tevatron.

For this analysis, a proof-of-principle, we used the small 5.02 TeV data sample of pp collisions from the LHCb Run 2 runs in 2017, with an integrated luminosity of $99.86 \pm 2.00 \text{ pb}^{-1}$. Despite the limited dataset, the methodology developed here can be seamlessly extended to larger datasets and higher energies, such as the full Run 2 dataset at 13 TeV, or future Run 3 datasets, enabling more precise and accurate measurements. Furthermore, it is anticipated that employing the complete Run 2 dataset at 13 TeV will enable us to achieve a total uncertainty of approximately 15 MeV. This is derived from an initial projection by our analysis team. This level of precision is highly competitive and could significantly affect the combined results with ATLAS and CMS, the other LHC experiments measuring m_W .

Integrating over our differential cross-section measurement, we obtain the following results for the integrated fiducial cross-sections for each W charge:

$$\begin{aligned}
\sigma_{W^+ \rightarrow \mu^+ \nu_\mu} &= 301.62 \pm 3.00_{(\text{stat})} \pm 3.74_{(\text{syst})} \pm 6.03_{(\text{lumi})} \text{ pb}, \\
\sigma_{W^- \rightarrow \mu^- \bar{\nu}_\mu} &= 237.80 \pm 2.72_{(\text{stat})} \pm 2.67_{(\text{syst})} \pm 4.76_{(\text{lumi})} \text{ pb}.
\end{aligned}
\tag{8.1}$$

The results have been compared against some theoretical predictions using the MCFM software, and were found to be compatible. More work needs to be done to further enhance our theoretical prediction models and to include QED FSR.

By fitting the W mass to the unfolded data, a result of

$$m_W = 80355 \pm 128_{(\text{stat})+(\text{syst})} \pm 21_{\text{theory}} \text{ MeV (blinded)} \tag{8.2}$$

has been obtained. The m_W result remains blinded, which means that we cannot yet compare our measurement directly with those of other experiments. However, at first glance, it seems that our analysis result and precision look promising, with a central value that is consistent with the theoretical and experimental average. The cross-checks performed suggest that this is a viable approach for measuring m_W using the 13 TeV data. It is important to note that this analysis is still ongoing, and further refinements and improvements are anticipated before its final release.

This work lays the foundation for several avenues of future research. As we move further into Run 3 of the LHCb experiment, the methods developed here can be applied to larger datasets, allowing for more precise measurements of the $W \rightarrow \mu \nu_\mu$ differential cross-section. Additionally, by extending this approach to other processes, such as $W \rightarrow e \nu_e$, we can continue to test the Standard Model and explore potential new physics.

There are also several potential improvements to this analysis. Future iterations could incorporate more sophisticated background modelling techniques, refine the unfolding process, improve the modelling and calibration of the isolation variable, and explore different strategies to reduce systematic uncertainties. Furthermore, improvements in detector calibration and alignment during Run 3 will likely reduce the uncertainties associated with the detector response, thus enhancing the precision of the final measurement.

Moreover, this thesis highlighted the importance of collaboration and teamwork in particle physics, showing the collective efforts of a diverse group of researchers around the world. The Real-Time Analysis initiative has achieved substantial improvements in data processing and online analysis at LHCb. Although the author's role in this project is limited, it highlights the collaborative nature of these improvements. With numerous small contributions and the ongoing commitment and ingenuity of the particle physics community, we can expect even more groundbreaking discoveries in the years to come.

In conclusion, this thesis demonstrates that the LHCb experiment and detector is a powerful tool for precision measurements in particle physics, with the capability to contribute significantly to the LHC program's goals in electroweak physics, and rival other experiments. The methods

developed here not only provide a path forward for measuring the $W \rightarrow \mu \nu_\mu$ differential cross-section with high precision, but also offer a blueprint for future analyses aimed at measuring fundamental quantities while ensuring the reproducibility and reusability of the results.

Appendix A

Isolation Calibration Verification

Plots comparing the calibrated Z MC isolation to the Z Data isolation are displayed across 5 bins of recoil projection, defined as

$$\text{Recoil} = \vec{u}_T \cdot \frac{\vec{p}_T^W}{|\vec{p}_T^W|}, \quad (\text{A.1})$$

along with 2 bins of pseudorapidity ranging from 2.2 to 4.4, consistent with our fiducial range.

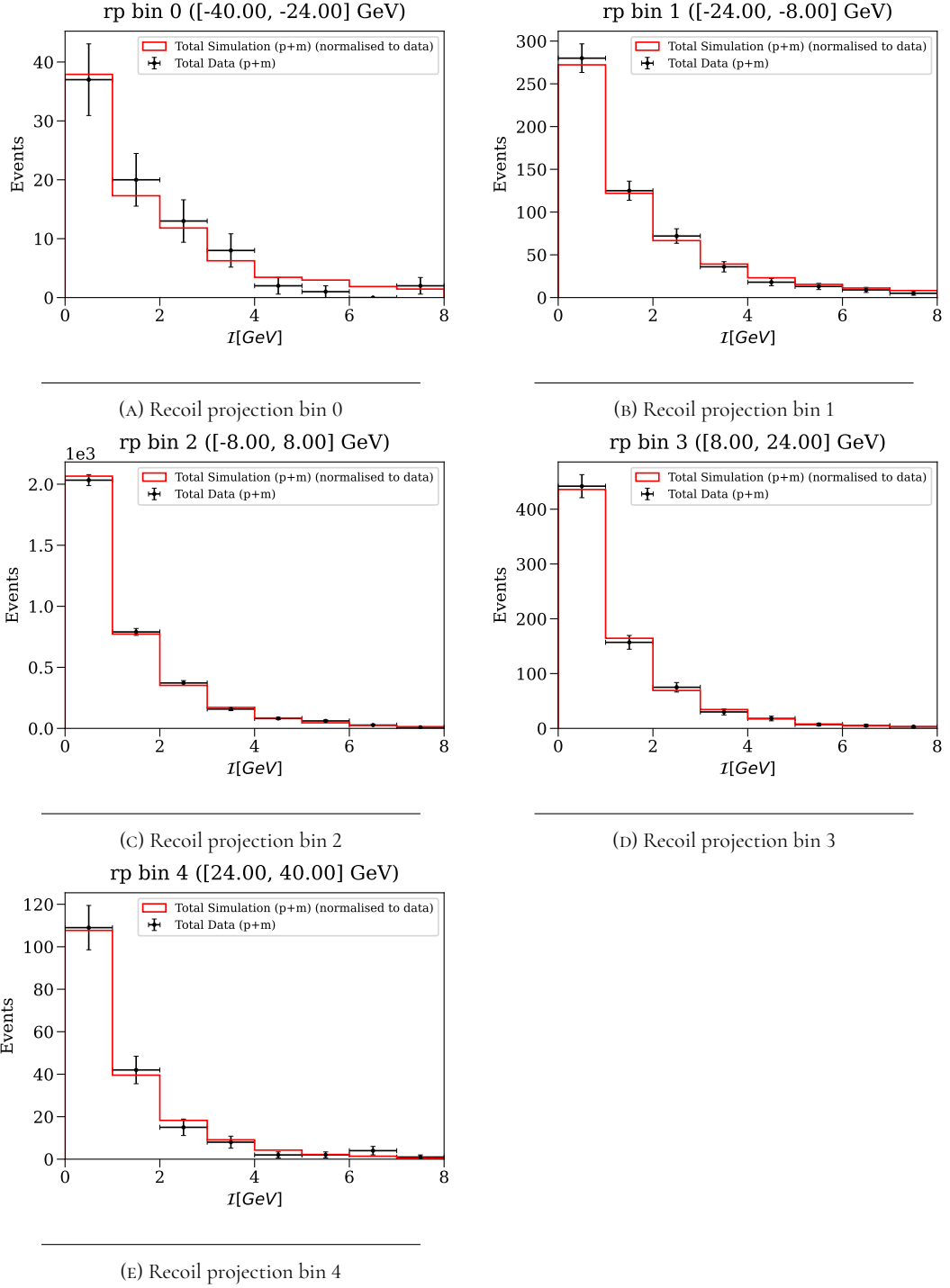


FIGURE A.1: Isolation calibration in bins of recoil projection from -40 to 40 GeV.

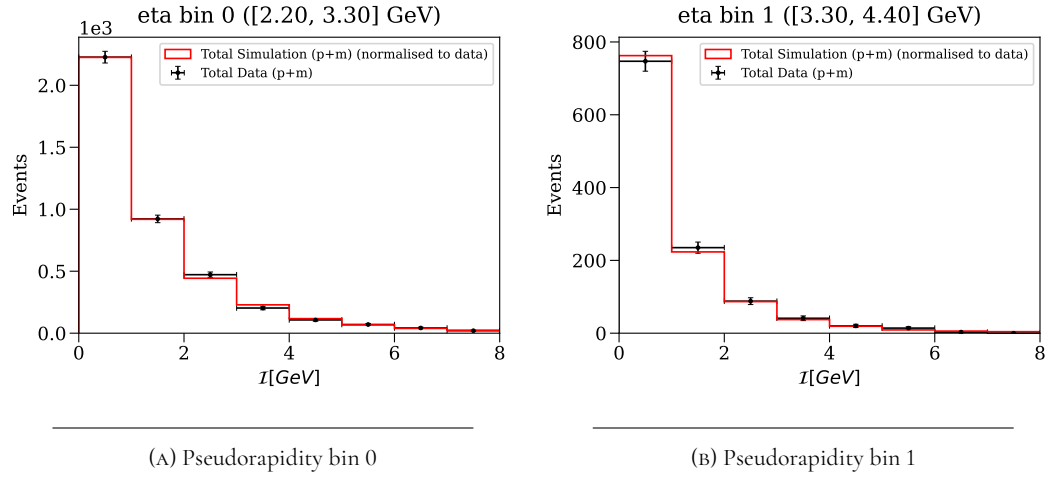


FIGURE A.2: Isolation calibration in bins of pseudorapidity.

Appendix B

Background Subtraction and Unfolding Fit Results

A summary of the background subtraction and unfolding fit outcomes for W^+ and W^- , showing $\mathcal{I}$ being used for background subtraction in bins of p_T . The results are validated with a pull plot at the bottom. The fit results seem visually acceptable, yet the pull plot shows some deviations, and the reduced χ^2 values are slightly higher than expected. These will certainly be tackled in due course.

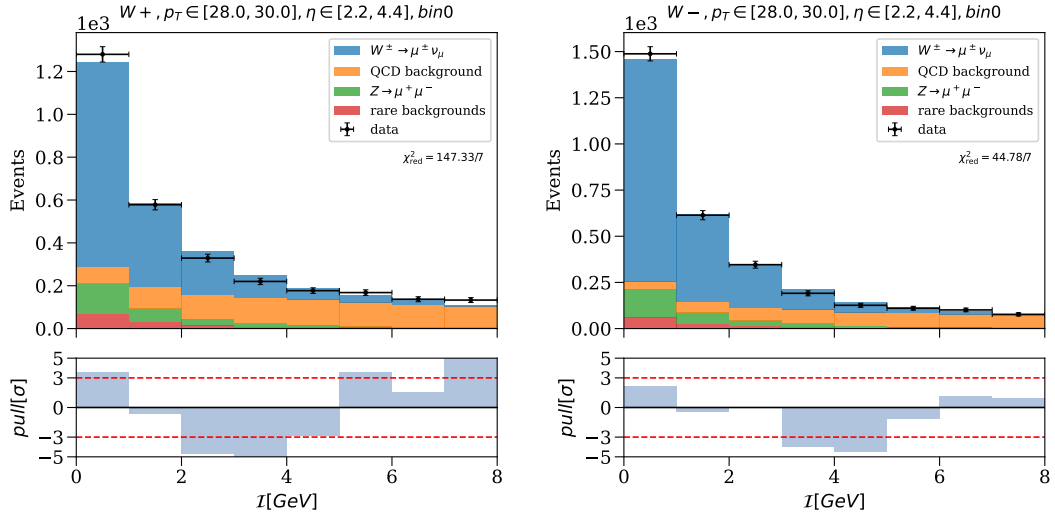


FIGURE B.1: Fit results histograms comparing fit model and grouped components to the data for p_T bin 0. The variable shown is the isolation variable. W^+ is on the left and W^- is on the right.

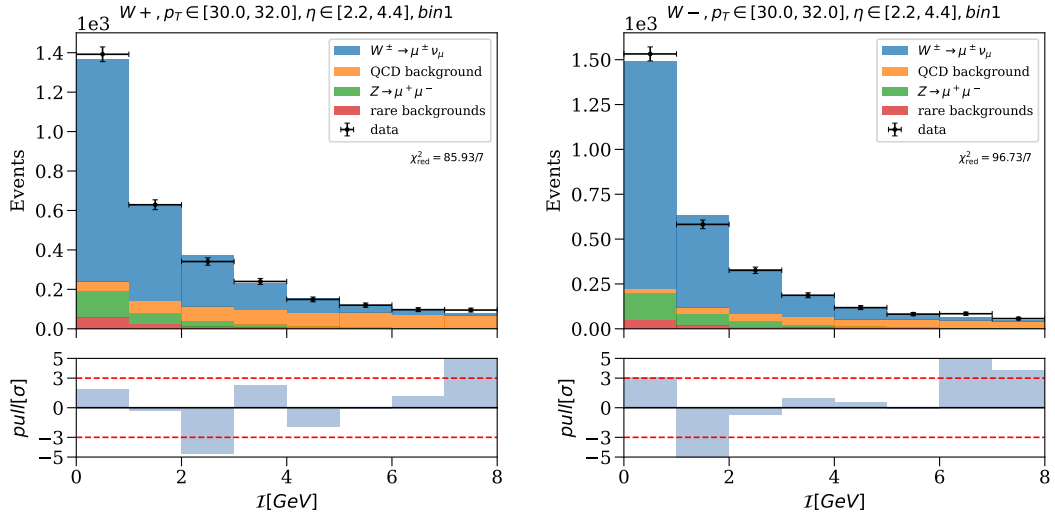


FIGURE B.2: Fit results histograms comparing fit model and grouped components to the data for p_T bin 1. The variable shown is the isolation variable. W^+ is on the left and W^- is on the right.

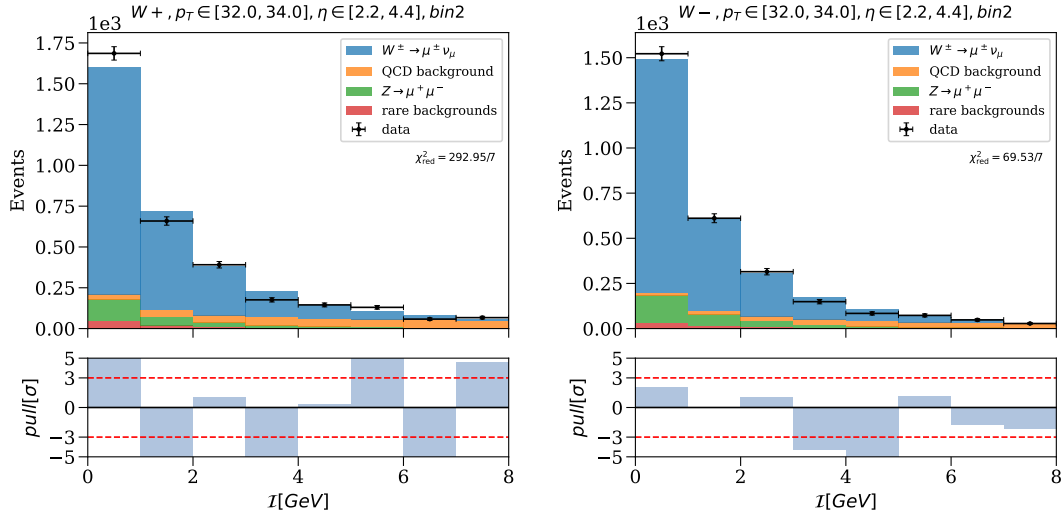


FIGURE B.3: Fit results histograms comparing fit model and grouped components to the data for p_T bin 2. The variable shown is the isolation variable. W^+ is on the left and W^- is on the right.

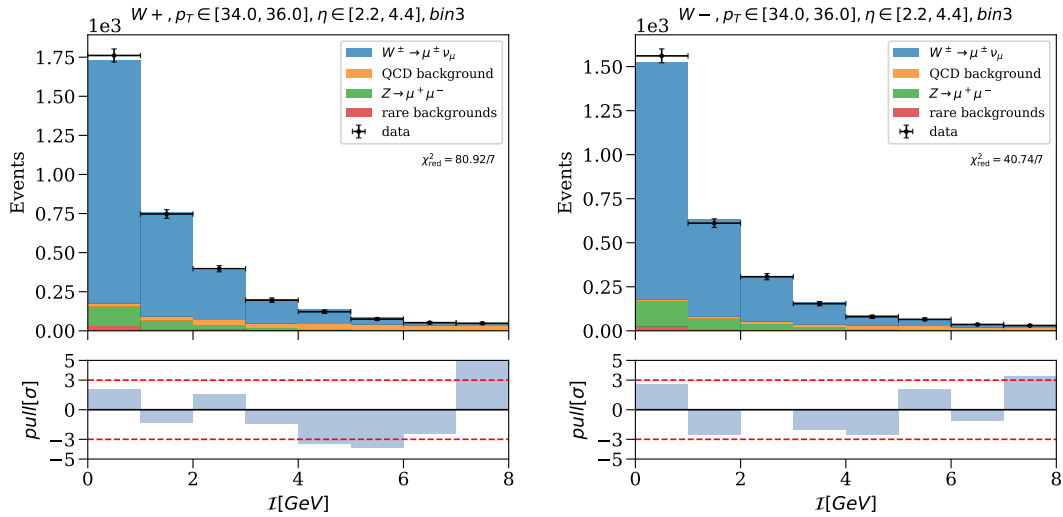


FIGURE B.4: Fit results histograms comparing fit model and grouped components to the data for p_T bin 3. The variable shown is the isolation variable. W^+ is on the left and W^- is on the right.

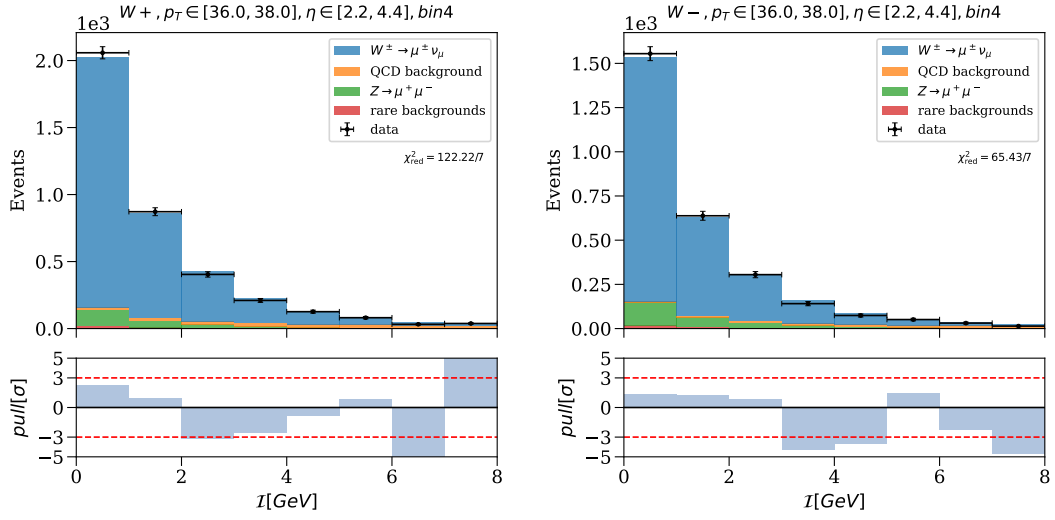


FIGURE B.5: Fit results histograms comparing fit model and grouped components to the data for p_T bin 4. The variable shown is the isolation variable. W^+ is on the left and W^- is on the right.

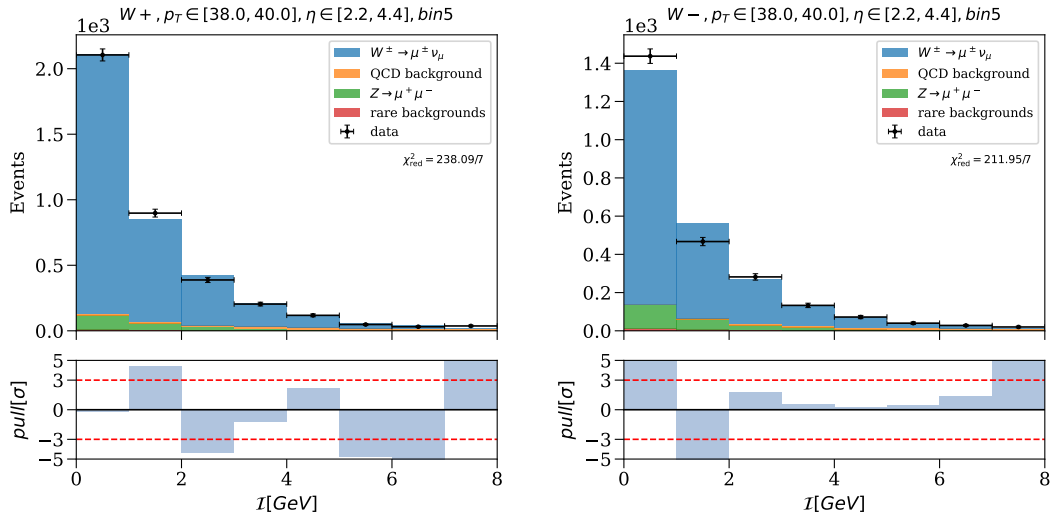


FIGURE B.6: Fit results histograms comparing fit model and grouped components to the data for p_T bin 5. The variable shown is the isolation variable. W^+ is on the left and W^- is on the right.

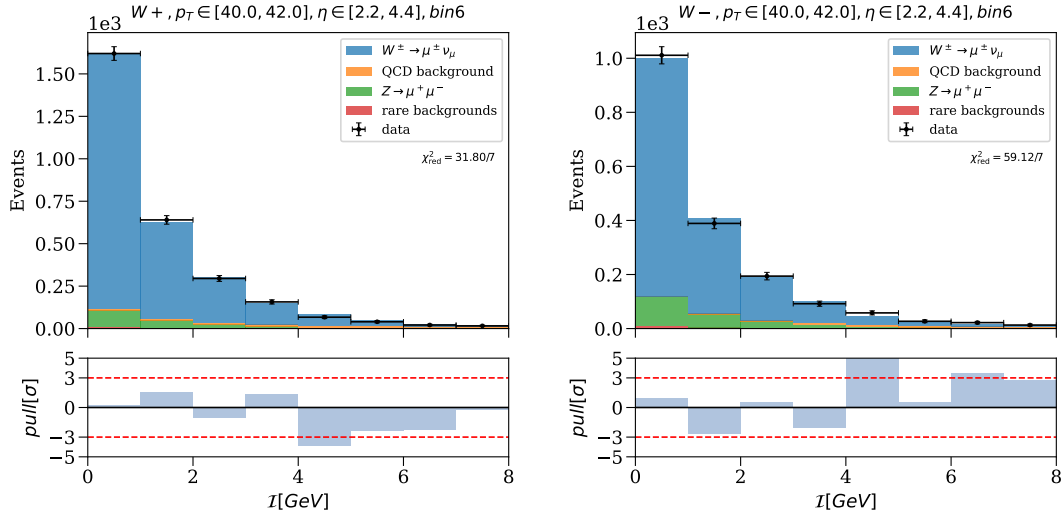


FIGURE B.7: Fit results histograms comparing fit model and grouped components to the data for p_T bin 6. The variable shown is the isolation variable. W^+ is on the left and W^- is on the right.

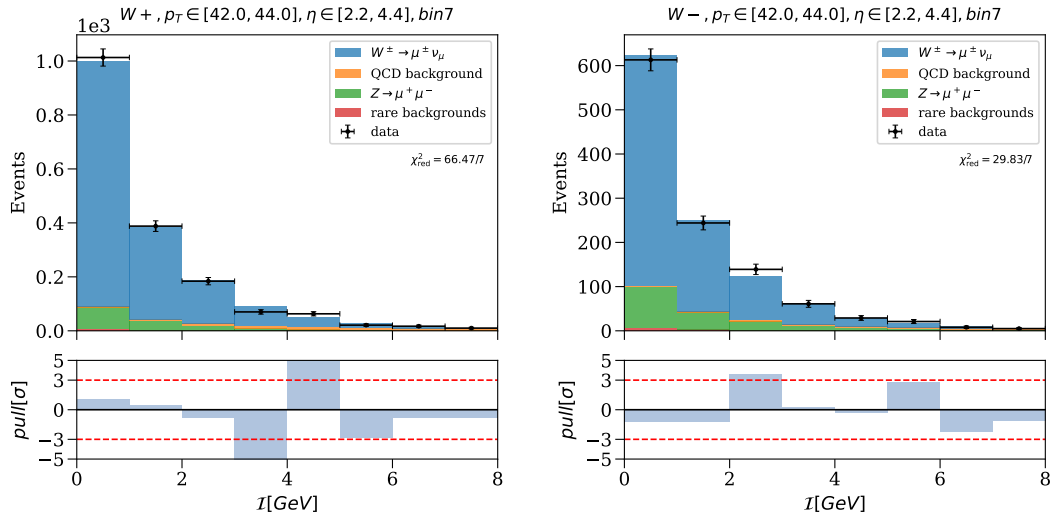


FIGURE B.8: Fit results histograms comparing fit model and grouped components to the data for p_T bin 7. The variable shown is the isolation variable. W^+ is on the left and W^- is on the right.

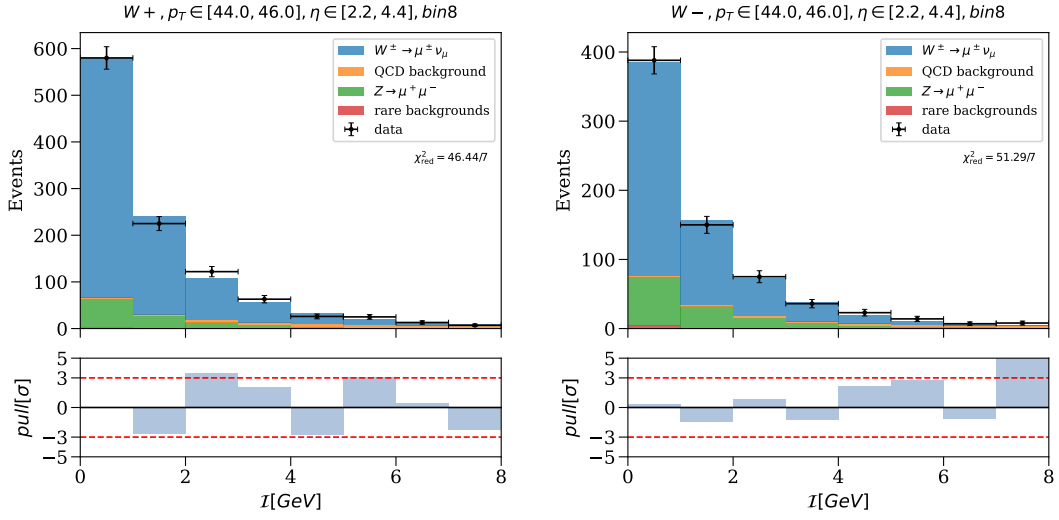


FIGURE B.9: Fit results histograms comparing fit model and grouped components to the data for p_T bin 8. The variable shown is the isolation variable. W^+ is on the left and W^- is on the right.

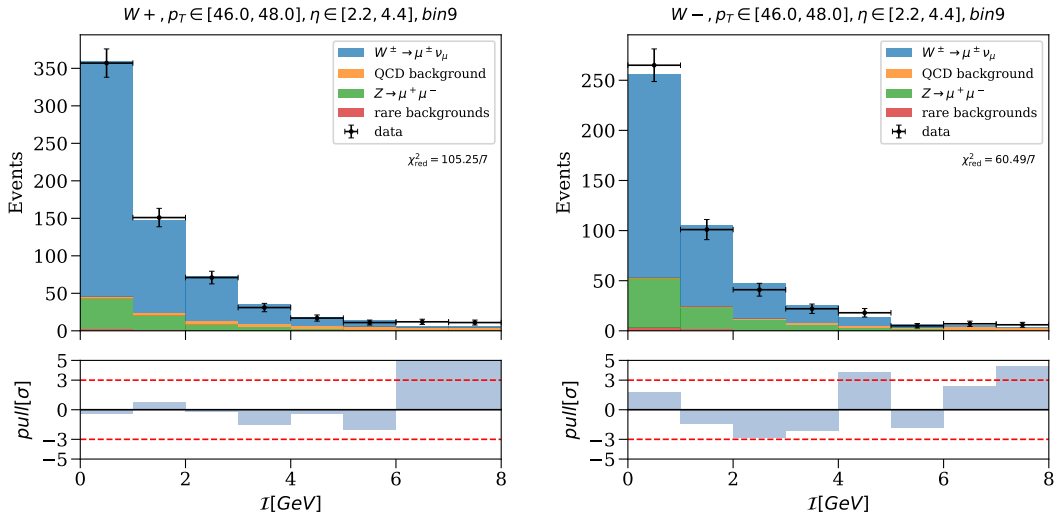


FIGURE B.10: Fit results histograms comparing fit model and grouped components to the data for p_T bin 9. The variable shown is the isolation variable. W^+ is on the left and W^- is on the right.

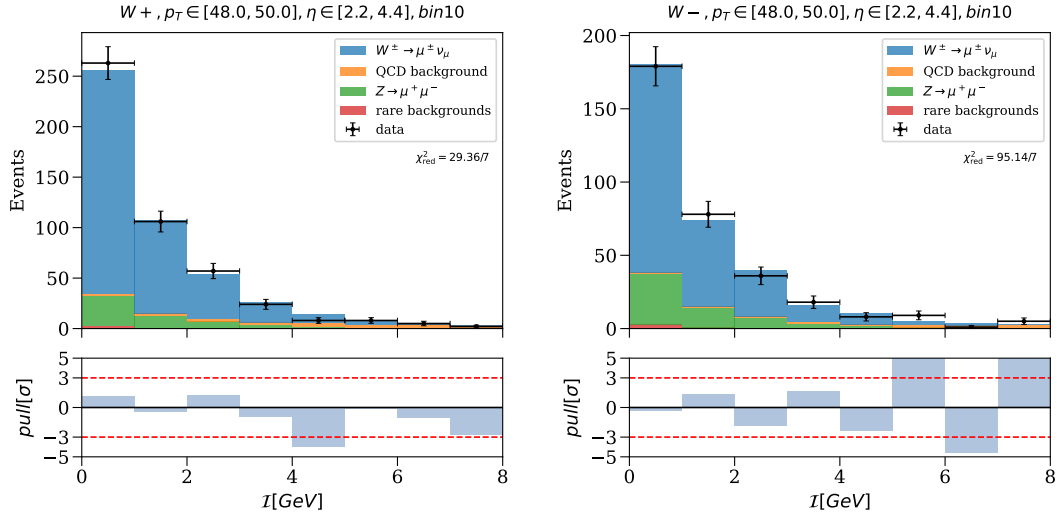


FIGURE B.11: Fit results histograms comparing fit model and grouped components to the data for p_T bin 10. The variable shown is the isolation variable. W^+ is on the left and W^- is on the right.

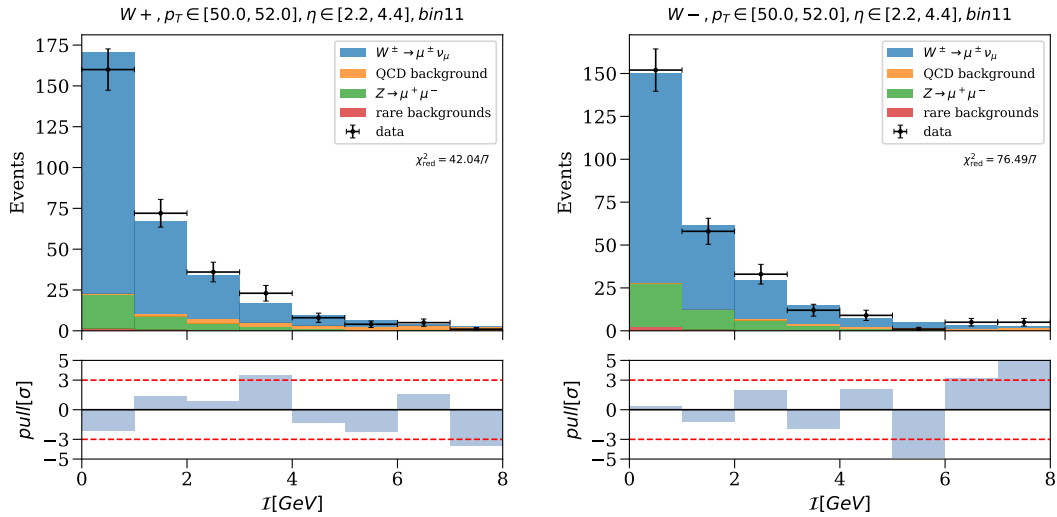


FIGURE B.12: Fit results histograms comparing fit model and grouped components to the data for p_T bin 11. The variable shown is the isolation variable. W^+ is on the left and W^- is on the right.

Appendix C

Systematic Uncertainties

The table of the systematic uncertainties per p_T bin for W^+ and W^- are shown in Table C.1, where the uncertainties are obtained from the diagonal of the covariance matrix. The table is organised in such a way that the W^+ results are listed first, followed by the W^- results for easy comparison. The total systematic uncertainty can be obtained by summing the individual systematic uncertainties in quadrature.

The covariance matrices of the individual systematic uncertainties are shown next.

bin index	Qcd Bkg Missd Rate Fit Model Syst	Qcd Bkg Abundance Syst	Iso Bins	Trigger Stats	Muid Stats	Trk Stats	Efficiency Phi Binning	Efficiency Eta Binning	Momentum Smearing	Pseudomass	Qcd Bkg Iso	Signal Iso
0	0.0101	0.0122	0.1096	0.0929	0.0194	0.0405	0.0637	0.0585	0.0447	0.1049	0.0264	0.2254
1	0.0007	0.0047	0.0853	0.1079	0.0208	0.0516	0.0862	0.0893	0.0620	0.1303	0.0165	0.1540
2	0.0022	0.0199	0.0678	0.1280	0.0291	0.0597	0.0933	0.0963	0.0471	0.1450	0.0118	0.1180
3	0.0028	0.0047	0.0521	0.1288	0.0271	0.0609	0.0982	0.1007	0.0623	0.1537	0.0095	0.0882
4	0.0068	0.0066	0.0500	0.1512	0.0312	0.0710	0.1211	0.1140	0.0830	0.1805	0.0051	0.0742
5	0.0026	0.0041	0.0375	0.1544	0.0322	0.0701	0.1241	0.1126	0.0734	0.1525	0.0076	0.0663
6	0.0002	0.0034	0.0212	0.1153	0.0259	0.0568	0.0879	0.0942	0.0976	0.2145	0.0031	0.0371
7	0.0033	0.0040	0.0178	0.0732	0.0133	0.0299	0.0511	0.0544	0.0977	0.2328	0.0017	0.0263
8	0.0042	0.0046	0.0146	0.0441	0.0074	0.0183	0.0228	0.0358	0.0563	0.1263	0.0026	0.0228
9	0.0010	0.0025	0.0084	0.0263	0.0044	0.0069	0.0080	0.0201	0.0455	0.0804	0.0013	0.0128
10	0.0002	0.0037	0.0229	0.0212	0.0033	0.0110	0.0013	0.0173	0.0457	0.0895	0.0014	0.0086
11	0.0016	0.0020	0.0083	0.0158	0.0053	0.0074	0.0101	0.0072	0.0389	0.1069	0.0004	0.0086
12	0.0030	0.0358	0.0901	0.1083	0.0200	0.0360	0.1183	0.0157	0.0196	0.0967	0.0154	0.1712
13	0.0009	0.0173	0.0852	0.1003	0.0187	0.0391	0.1128	0.0368	0.1043	0.0118	0.0188	0.1253
14	0.0013	0.0055	0.0438	0.0967	0.0210	0.0411	0.0980	0.0427	0.0284	0.1011	0.0059	0.0733
15	0.0028	0.0016	0.0330	0.0980	0.0198	0.0433	0.0875	0.0666	0.0284	0.1181	0.0024	0.0659
16	0.0001	0.0035	0.0225	0.1037	0.0192	0.0448	0.0974	0.0714	0.0223	0.1181	0.0061	0.0347
17	0.0026	0.0027	0.0149	0.0925	0.0177	0.0406	0.0857	0.0670	0.0596	0.1145	0.0009	0.0388
18	0.0006	0.0030	0.0243	0.0707	0.0151	0.0301	0.0414	0.0517	0.0409	0.1487	0.0043	0.0282
19	0.0003	0.0010	0.0161	0.0448	0.0093	0.0193	0.0105	0.0332	0.0863	0.1465	0.0014	0.0187
20	0.0002	0.0036	0.0183	0.0279	0.0040	0.0070	0.0056	0.0168	0.0193	0.0767	0.0025	0.0114
21	0.0008	0.0010	0.0129	0.0192	0.0045	0.0057	0.0008	0.0166	0.0293	0.0857	0.0017	0.0158
22	0.0005	0.0038	0.0242	0.0117	0.0054	0.0063	0.0079	0.0012	0.0222	0.0673	0.0024	0.0068
23	0.0012	0.0020	0.0143	0.0149	0.0045	0.0063	0.0027	0.0140	0.0219	0.0698	0.0006	0.0148

TABLE C.1: Systematic uncertainties per p_T bin for W^+ and W^- . The first 12 rows correspond to W^+ , and the last 12 rows correspond to W^- . The uncertainties are obtained from the diagonal of the covariance matrix.

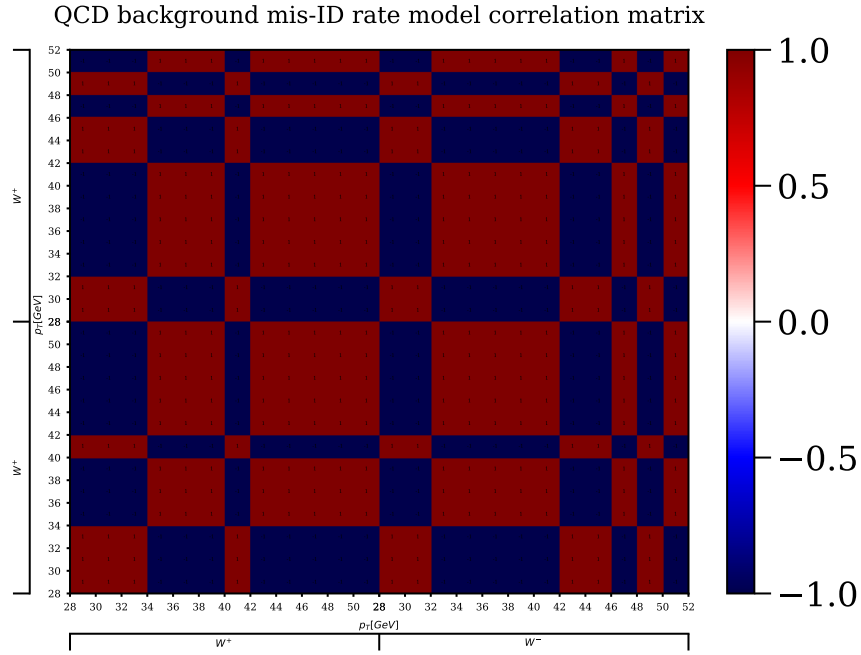


FIGURE C.1: The correlation matrix of the Hadronic background misidentification rate model systematic for the differential cross-section with respect to p_T for both W^+ and W^- .

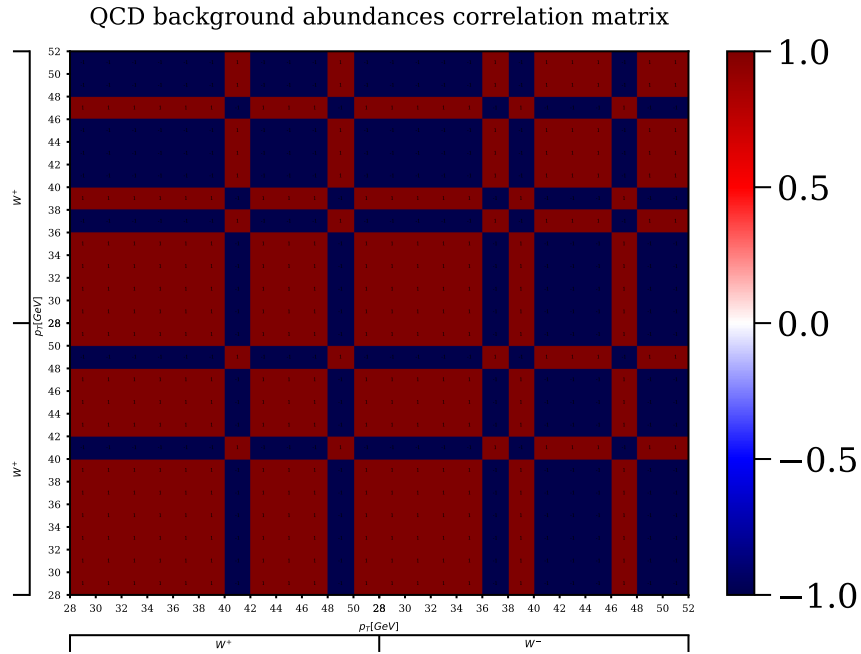


FIGURE C.2: The correlation matrix of the Hadronic background hadron species abundance systematic for the differential cross-section with respect to p_T for both W^+ and W^- .

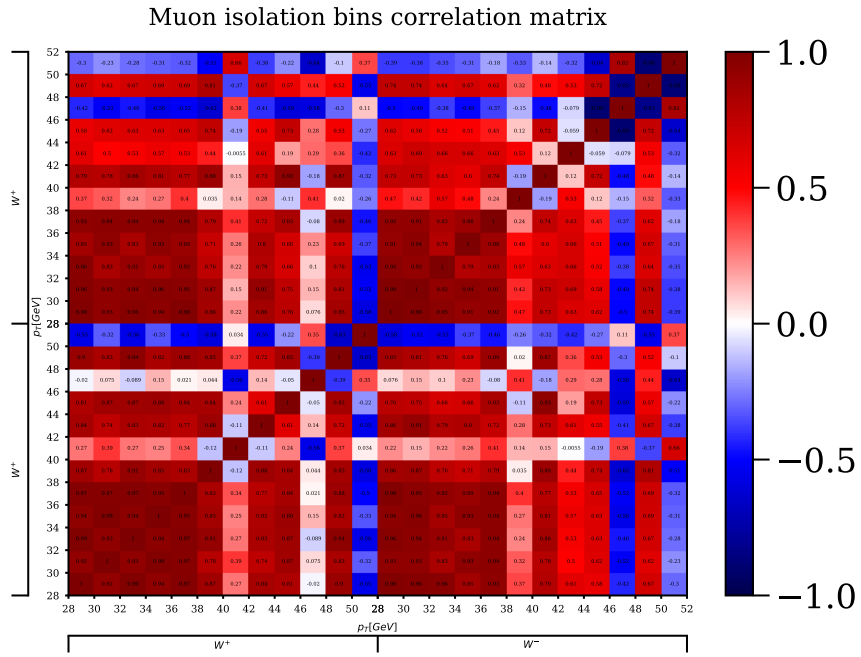


FIGURE C.3: The correlation matrix of the Isolation bins systematic for the differential cross-section with respect to p_T for both W^+ and W^- .

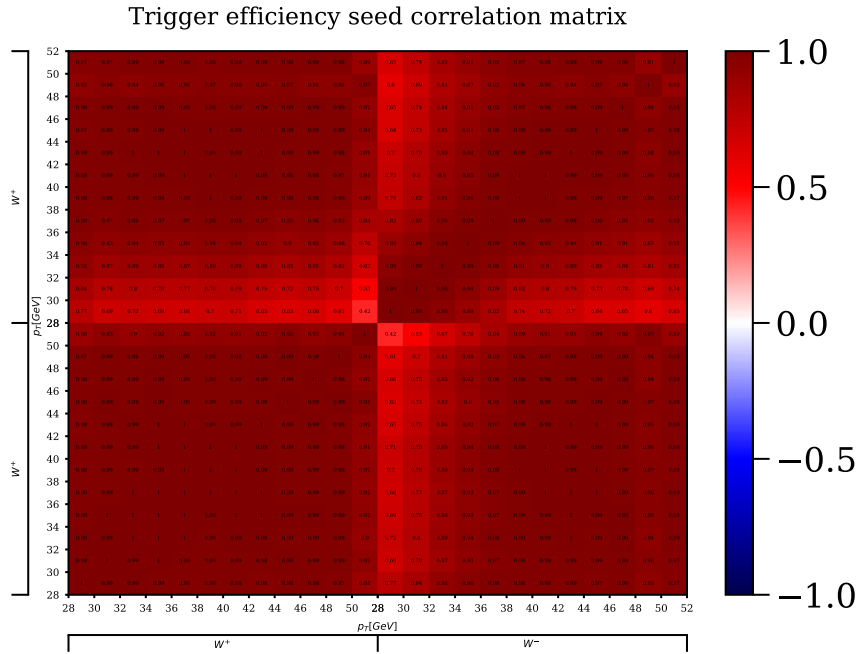


FIGURE C.4: The correlation matrix of the Trigger statistics systematic for the differential cross-section with respect to p_T for both W^+ and W^- .

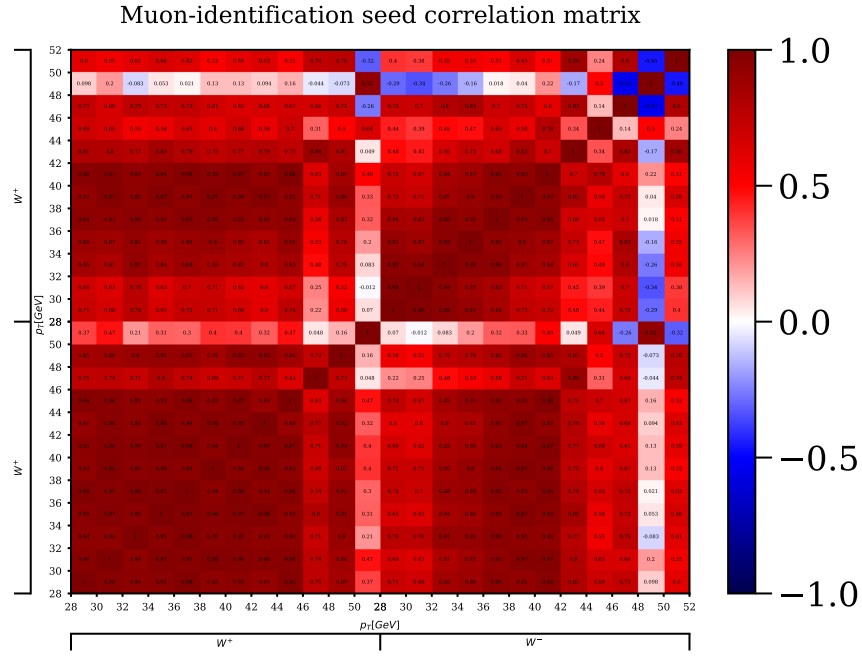


FIGURE C.5: The correlation matrix of the Muon identification statistics systematic for the differential cross-section with respect to p_T for both W^+ and W^- .

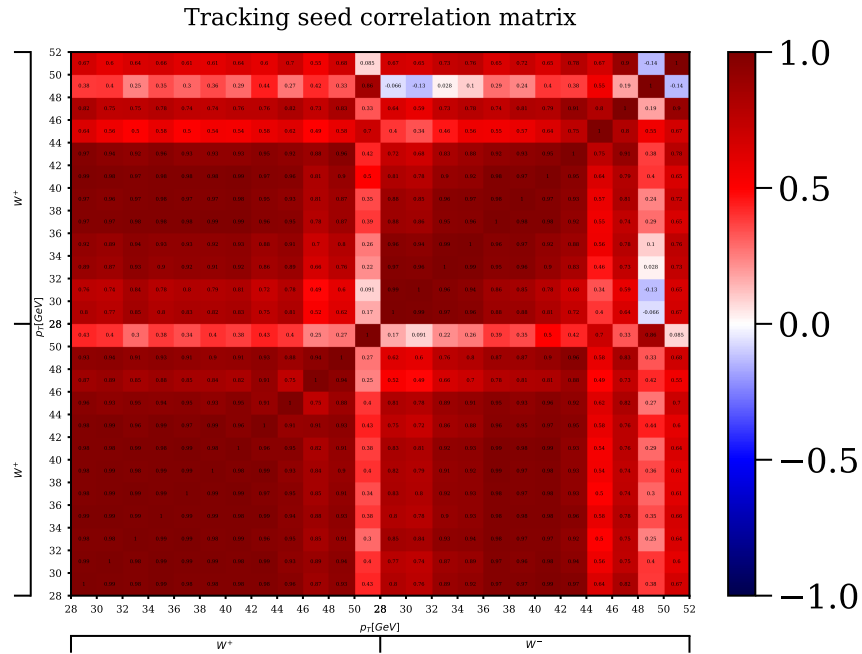


FIGURE C.6: The correlation matrix of the Tracking statistics systematic for the differential cross-section with respect to p_T for both W^+ and W^- .

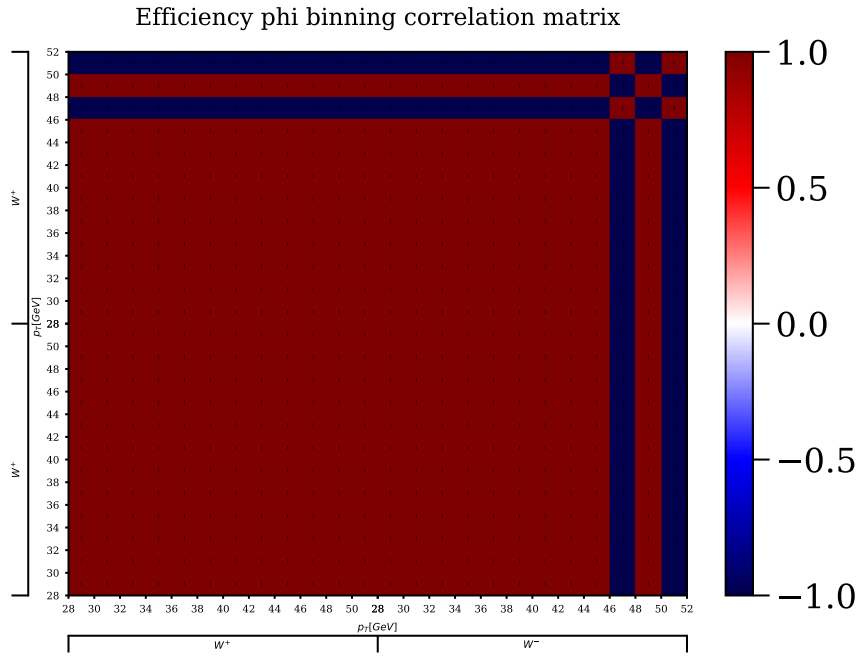


FIGURE C.7: The correlation matrix of the Efficiency ϕ binning systematic for the differential cross-section with respect to p_T for both W^+ and W^- .

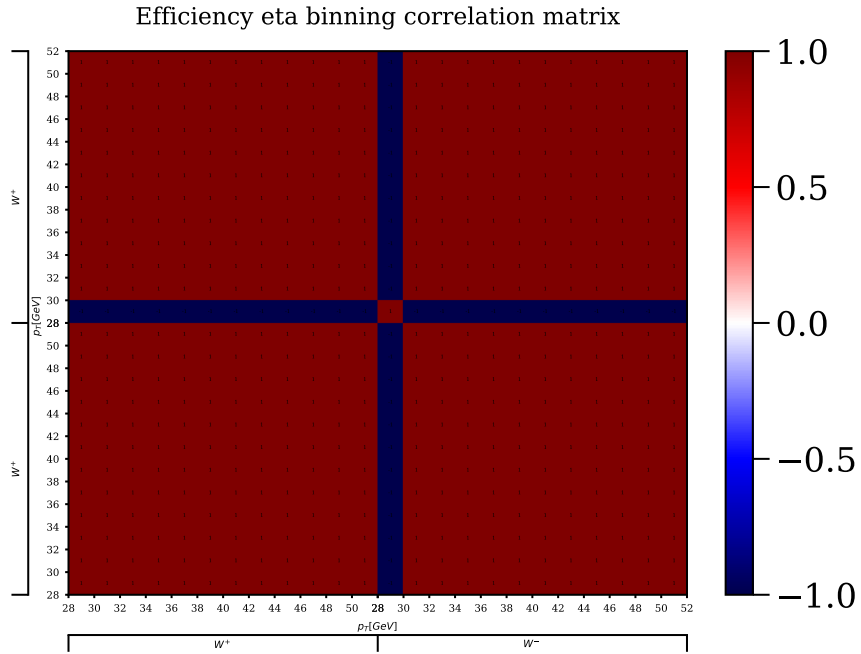


FIGURE C.8: The correlation matrix of the Efficiency η binning systematic for the differential cross-section with respect to p_T for both W^+ and W^- .

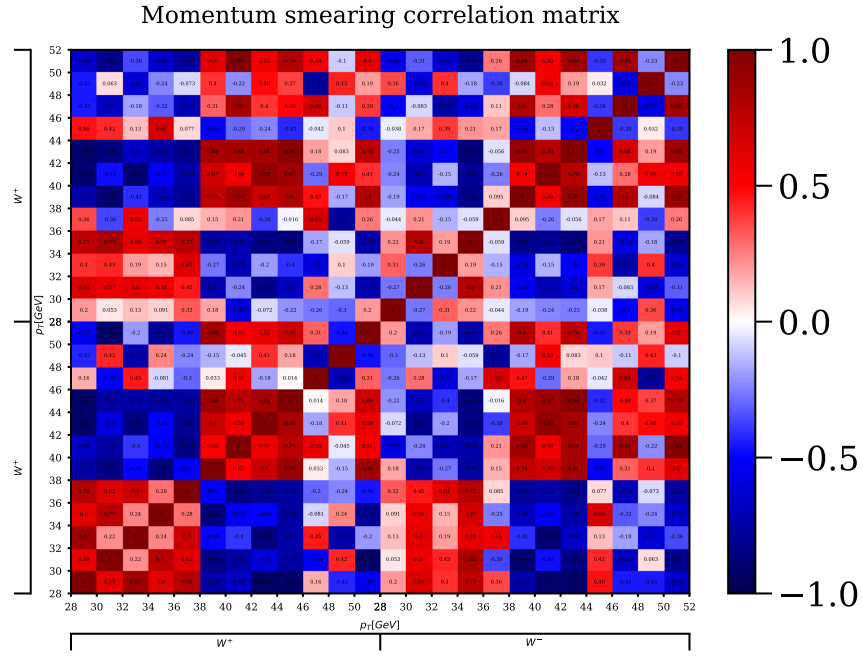


FIGURE C.9: The correlation matrix of the Momentum smearing systematic for the differential cross-section with respect to p_T for both W^+ and W^- .

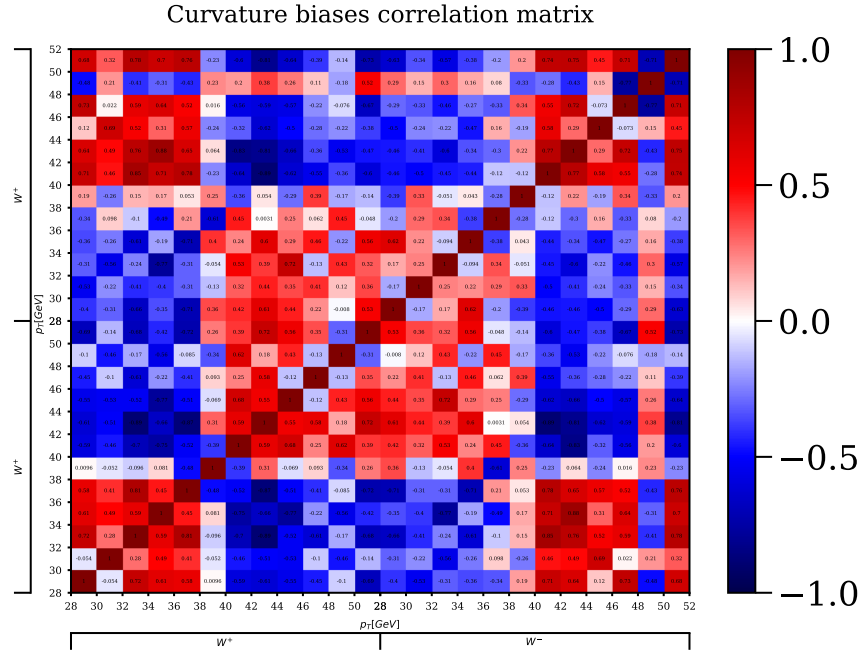


FIGURE C.10: The correlation matrix of the Curvature bias systematic for the differential cross-section with respect to p_T for both W^+ and W^- .

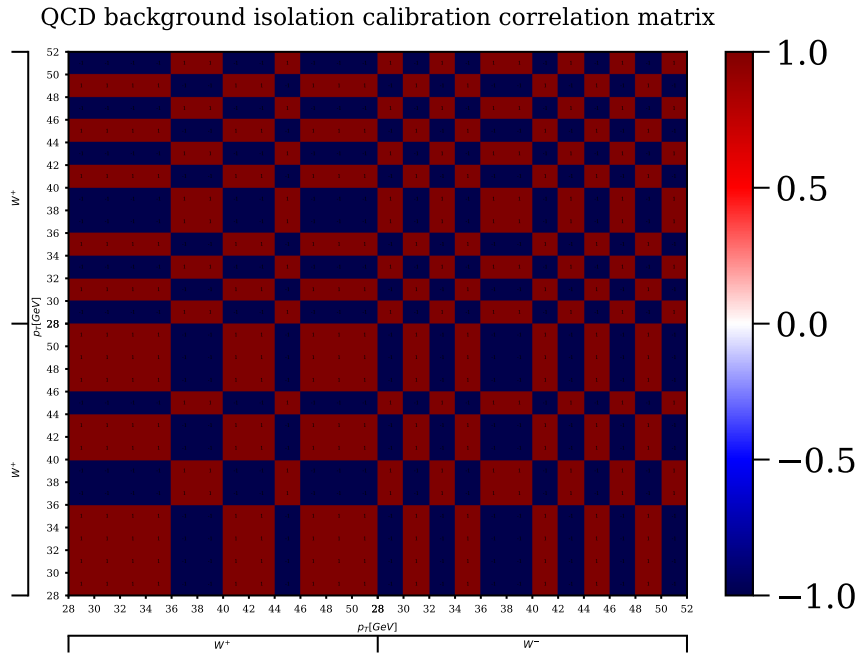


FIGURE C.11: The correlation matrix of the Hadronic background isolation calibration systematic for the differential cross-section with respect to p_T for both W^+ and W^- .

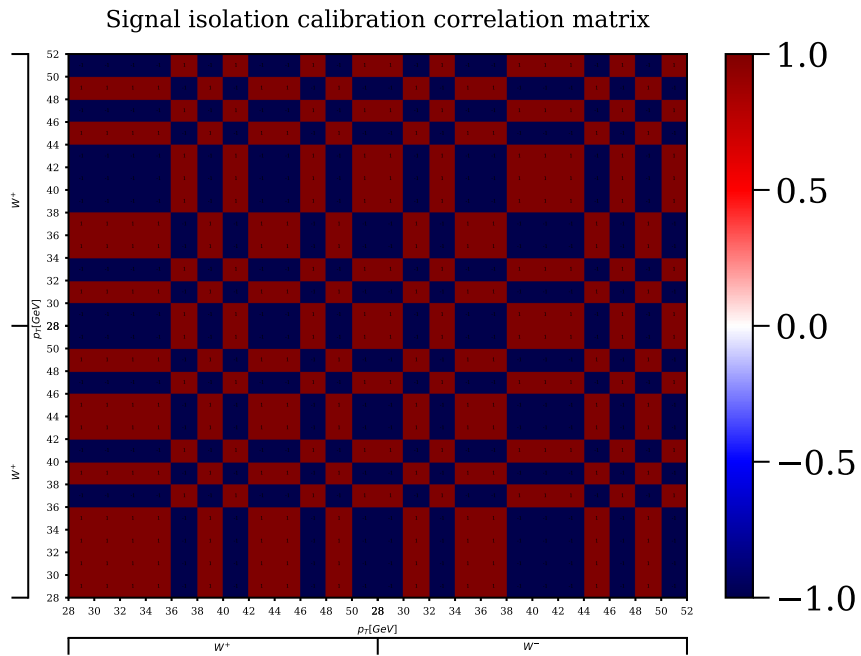


FIGURE C.12: The correlation matrix of the Signal isolation calibration systematic for the differential cross-section with respect to p_T for both W^+ and W^- .

The plots of the statistical correlation matrices are shown in Fig. C.13.

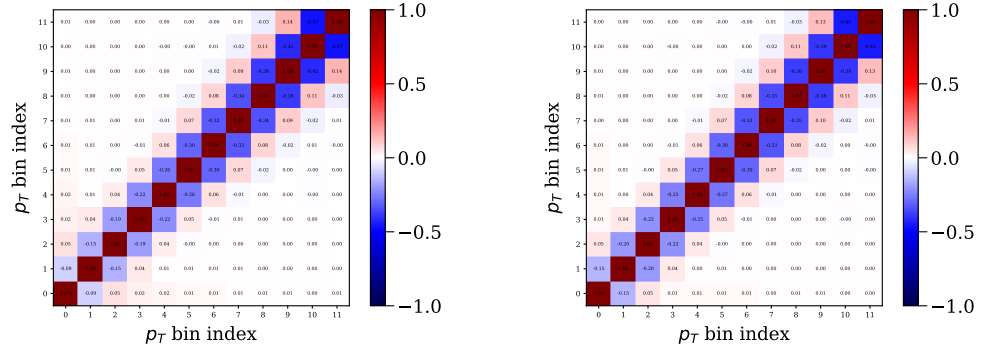


FIGURE C.13: The correlation matrices for the differential cross-sections with respect to muon transverse momentum (p_T) for (left) W^+ and (right) W^- .

A plot of the individual systematic uncertainties overlaid is shown in Fig. C.14.

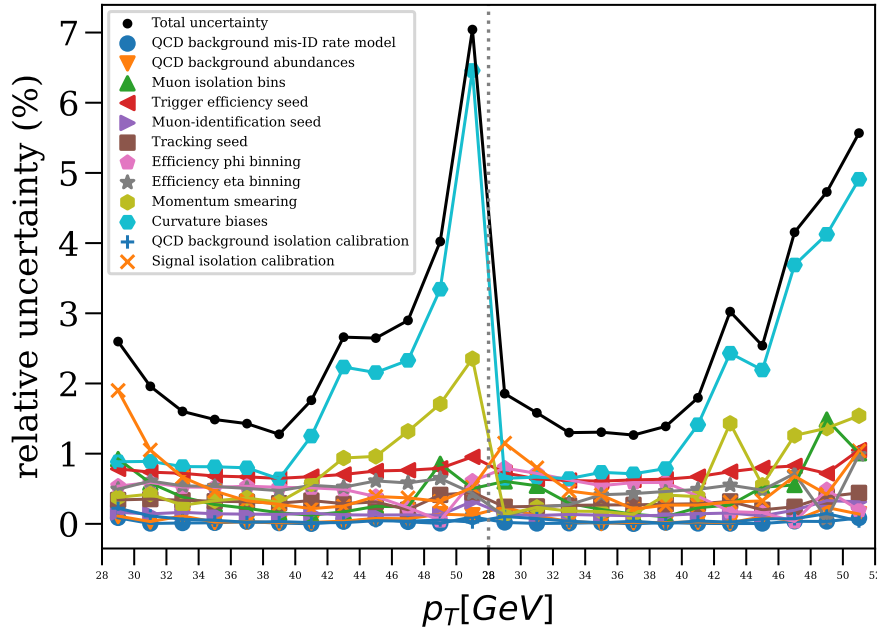


FIGURE C.14: The relative uncertainties of the systematic uncertainties for the differential cross-sections with respect to muon transverse momentum (p_T) for (left) W^+ and (right) W^- . The relative uncertainties are stacked to show the contribution of each source of uncertainty.

Viewing the plots above, it is evident that curvature biases significantly influence our systematic uncertainties. Moreover, the high degree of anticorrelation between both charges impacts the overall correlation matrix, as previously shown in Section 6.5.

Appendix D

Varying PDFs

These are the detailed fit results from varying the PDFs when performing the W boson mass fit as shown in Section 7.5.

Parameter	Value	Error
g	0.8475	0.5786
m_W [MeV]	80364	127
Charge asymmetry	0.1186	0.0063
α_s	0.1113	0.0050

TABLE D.1: Results of changing the PDF set to MSHT20nnlo.

Parameter	Value	Error
g	1.0282	0.5856
m_W [MeV]	80355	128
Charge asymmetry	0.1177	0.0063
α_s	0.1113	0.0050

TABLE D.2: Results of changing the PDF set to CT18NNLO.

Parameter	Value	Error
g	0.9604	0.5815
m_W [MeV]	80350	128
Charge asymmetry	0.1181	0.0063
α_s	0.1113	0.0050

TABLE D.3: Results of changing the PDF set to NNPDF40_nnlo_as_0118.

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