

Improvement of the muon spectrometer and precision W-boson physics in ATLAS

*Amélioration du spectromètre à muons et physique de
précision du boson W dans ATLAS*

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Title: Improvement of the muon spectrometer and precision W-boson physics in ATLAS

Key words: *MPGD, Micromegas, ATLAS, NSW, Cosmic bench, Electroweak interactions, Standard Model, Large Hadron Collider, W-boson, Statistics*

Abstract: This thesis presents the work on the production and the qualification of the ATLAS NSW Micromegas Large Modules 1 (LM1), as well as the measurements of the W-boson transverse momentum and the W-boson mass using low pile-up data recorded by ATLAS at center-of-mass energies of $\sqrt{s} = 5.02$ and 13 TeV. Detailed protocols and results of the cosmic bench characterization for the LM1 Micromegas modules are documented, including the preliminary tests of the improved working gas mixture which was in the end adopted for the Run 3 data taking. The physics modelling and detector calibration for the W-boson analyses using the low pile-up data are described. The data-driven estimation of the multijet background is presented, emphasizing the improvements with respect to the procedures used in the previous analyses. The measurement of the W-boson transverse momentum achieves a remarkably high precision and will become a valuable input to the W mass analysis. Based on a profile likelihood fit, the framework of statistical analysis for the low pile-up W mass measurement has been studied and optimized. Benefiting from the detector calibration derived from the measurement of the W-boson transverse momentum, the low pile-up W mass analysis is expected to become a proof-of-concept measurement for the precise determination of the W-boson mass using future low pile-up datasets in ATLAS. Also, a coherent approach of uncertainty decomposition for profile likelihood fits is developed and illustrated for the first time.

Titre : Amélioration du spectromètre à muons et physique de précision du boson W dans ATLAS

Mots clés : *MPGD, Micromegas, ATLAS, NSW, Banc cosmique, Interactions électrofaibles, Modèle standard, Grand collisionneur de hadrons, boson W, Statistiques*

Résumé : Cette thèse présente les travaux sur la production et la qualification des modules NSW Micromegas "Large Modules 1" (LM1) de l'expérience ATLAS, ainsi que les mesures de l'impulsion transverse du boson W et de la masse du boson W à l'aide de données à faible empiement enregistrées par ATLAS à des énergies dans le centre de masse de $\sqrt{s} = 5.02$ et 13 TeV. Les protocoles et les résultats détaillés de la caractérisation du banc cosmique pour les modules Micromegas LM1 sont documentés, y compris les tests préliminaires avec un mélange gazeux qui aboutit à des performances améliorées et qui a finalement été adopté pour la prise de données du Run 3. La modélisation physique et l'étalonnage du détecteur pour les analyses du boson W à l'aide des données à faible empiement sont décrits. L'estimation basée sur les données du fond multijet est présentée, en soulignant les améliorations par rapport aux procédures utilisées dans les analyses précédentes. La mesure de l'impulsion transverse du boson W atteint une précision remarquablement élevée et deviendra une donnée précieuse pour l'analyse de la masse du boson W. Sur la base de l'ajustement par maximum de vraisemblance profilé, le cadre d'analyse statistique pour la mesure de la masse du boson W à faible empiement a été étudié et optimisé. Bénéficiant de l'étalonnage du détecteur obtenu pour la mesure de l'impulsion transverse du boson W, l'analyse de masse du boson W à faible empiement devrait devenir une preuve de concept pour la détermination précise de la masse du boson W à l'aide de futures données à faible empiement dans ATLAS. En outre, une approche cohérente de la décomposition de l'incertitude pour l'ajustement par maximum de vraisemblance profilé est développée et illustrée pour la première fois.

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¹¹² Introduction

Particle physics, as the center of our understanding of the fundamental laws of nature, studies the basic constituents and interactions of our universe, which concern the elementary particles and the forces between them. As a discipline of experimental science, the development of particle physics relies on the progress on both the theoretical side and the experimental side. While efforts have been constantly made to explore the possibility of extending our current understanding of particle physics, the experimental research of particle physics are also carried out based on various foundations, including the measurement of neutrino properties [1], the detection of cosmic ray [2] and the experiments at colliders [3], [4].

Ever since the discovery of the Higgs boson [5], [6] in 2012 declaring the completion of the Standard Model of particle physics (SM), no smoking gun proof has yet been found regarding the possible existence of new elementary particles and interactions beyond the SM. However, the lack of new discovery hardly marks the end of the particle physics, but on the contrary, as eminent physicist Albert Abraham Michelson, Nobel Prize laureate known for his work on the measurement of the speed of light and especially for the famous Michelson-Morley experiment, once remarked:

"The future truths of physical science are to be looked for in the sixth place of decimals."

The future of particle physics counts on precision measurement, indeed. One important parameter in the SM that needs to be precisely measured at the moment is the mass of the W-boson (m_W), since any discrepancy between the precisely calculable SM prediction and the results of the experimental measurements will be a strong indication of new physics. The W-boson is an elementary particle in the SM as the mediator of the electrically-charged weak interaction. The measurement of m_W started from the discovery of the W-boson in 1983 [7], and has been performed in various experimental conditions. However, considering that the most up-to-date world average result of direct m_W measurements is still less precise than the state-of-the-art SM prediction and that there is a discrepancy of around one standard deviation between the two quantities, more efforts have to be continually spent on the determination of m_W from the experimental side. Furthermore, the recently released CDF II result has achieved the best precision so far as a single measurement, but revealed a significant tension with the SM prediction [8], which calls for urgent needs for more and better m_W measurements at other experiments as crucial validations. As one frontier of experimental particle physics, the ATLAS experiment [9] located at the Large Hadron Collider (LHC) [10] benefits from the stably controlled collisions of the proton beams happening at high energies and high rates, as well as its advanced particle detection system and full coverage of the azimuthal angle, making it a highly competitive platform for carrying out precision measurements of m_W and other W-boson properties. With the increasing luminosity at LHC, the ATLAS detector has to be upgraded accordingly so as to better exploit its physics potential for both precision measurement and searches of new physics. During the last long shut-down of LHC, the inner end-cap of the ATLAS muon system has been replaced with the New Small Wheels (NSWs) [11]. The ATLAS NSWs are equipped with detectors made of improved technologies, specifically sTGC [12] and Micromegas [13]. In ATLAS, the muons not only carry important kinematic information that helps reveal the interactions happening during the proton-proton collisions, but also provides clean signature

for event triggering. The NSW upgrade is expected to enhance the background rejection for the triggering of the muon system and keep the precision tracking of the muons.

The work of this thesis is dedicated to the detector upgrade of the ATLAS experiment and the precision measurement of the W-boson physics at ATLAS using specialized datasets. The thesis is structured in the following way: The basic theoretical background is reviewed in Chapter 1, followed by an introduction to the LHC and ATLAS experiment in Chapter 2. Chapter 3 documents the cosmic bench qualification of the LM1 Micromegas modules for the ATLAS NSW upgrade. The discussions about precision W-boson physics start from an overview of m_W measurements in Chapter 4, then covered by Chapters 5, 6, 7 and 8 with more details. In Chapter 9, the work related to the development of the statistical method is presented, which benefits to the precision measurements at high energy physics experiments. Finally, Chapter 10 concludes the thesis.

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²³⁷ Chapter 1

²³⁸ Theoretical background

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This chapter briefly illustrates the main aspects of the theoretical background of the work related to the thesis. Particularly, the theory prediction of the W-boson mass and the production mechanism of W-bosons in proton-proton colliders will be discussed.

1.1 The Standard model of particle physics

The state-of-the-art of the understanding of particle physics is condensed in the Standard Model of particle physics (SM), which is a quantum field theory that comprises three generations of quarks and leptons as well as three fundamental forces between them: the electromagnetic interaction, the weak interaction and the strong interaction. The SM follows the gauge symmetries of the unitary product group $SU(3) \times SU(2)_L \times U(1)$, while the gravity is neglected in the SM because of its tiny magnitude on the microscopic scale and consequently its negligible effect on the predictions of the SM.

An overview of the elementary particles in the SM is summarized in Fig. 1.1, where the fermions of spin 1/2 contain six quarks: up, down, charm, strange, top and bottom, while the six leptons are the electron, muon, tau, electron neutrino, muon neutrino and tau neutrino. Each of the above fermions comes with its anti-particle that are also included in the SM. The gauge bosons carry a spin of 1 and mediate interactions via the bosonic fields. The photon is the mediator of the electromagnetic interaction, the $W^\pm$ and Z bosons mediate the weak interaction, and the eight gluons are the mediators of the strong interaction. One scalar particle in the SM, the Higgs boson, is introduced by the mechanism of spontaneous symmetry breaking (Section 1.1.3) to explain the origin of the masses of the $W^\pm$ and Z bosons.

Assuming neutrinos are normal Dirac fermions, the SM has 25 free parameters that have to be set by hand [2]. They are:

- The masses of the twelve fermions.
- The three coupling constants corresponding to the strengths of the gauge interactions: α , G_F and α_S .
- The vacuum expectation value of the Higgs potential v , and the mass of the Higgs m_H .
- The eight mixing angles of the PMNS and CKM matrices.

In addition, if Charge Parity (CP) violation is allowed in the strong interaction, the strong phase θ_{CP} can be counted as the 26th free parameter of the SM, even though its value is experimentally known to be extremely small and is usually considered to be zero.

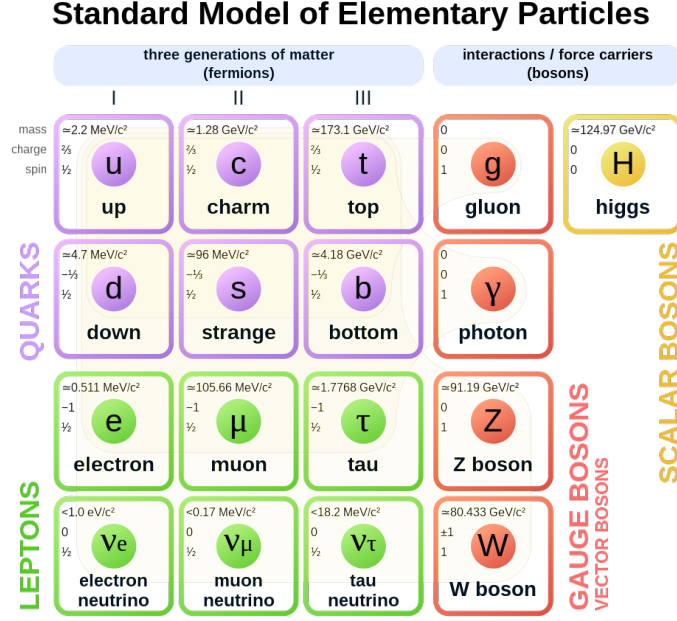


Figure 1.1 – Elementary particles of the Standard Model. (Figure from [1].)

The SM is by far the most successful theory in the domain of particle physics and represents one the triumphs of modern physics, achieving good descriptions of almost all the experimental data and the relevant phenomena. The SM has been tested by a variety of different experiments for decades. At TeV scale, the tests have been performed to a precision of typically 1 %, as shown in the example of Figure 1.2, where the SM total and fiducial production cross section measurements at ATLAS are found to be in good agreement with the theoretical calculations [3]. At lower energies, the SM has been tested to the precision of even better than 1 ppb for certain processes, such as the muon magnetic anomaly [4]. As an overview of the theoretical foundation of the work covered in this thesis, the electroweak sector of the SM (Section 1.1.1), the strong interaction sector (Section 1.1.2) of the SM and the BEH mechanism (Section 1.1.3) will be briefly summarized in the sections below, followed by an introduction to the global EW fit (Section 1.1.4).

1.1.1 The electroweak sector

Quantum field theory (QFT) is a set of ideas and tools that comprise three major themes of modern physics: quantum theory, concept of field and relativity [5]. QFT underlies particle physics and is the bridge between physics and mathematics.

Each of the fundamental forces covered in the SM can be described by QFT. In the case of electromagnetism, the theory of Quantum Electrodynamics (QED), from which QFT is established, dictates that the interactions between charged particles are mediated by the exchange of virtual photons. A virtual particle is a mathematical construct that represents the effect of summing over all possible time-ordered diagrams, and summing over the possible polarization states of the exchanged particle where it applies. The long range interaction of electromagnetic force can be described in terms of a photon exchange. Similar to the exchange of photon in QED, the weak charged-current interaction is mediated by the charged W^+ and W^- bosons, while

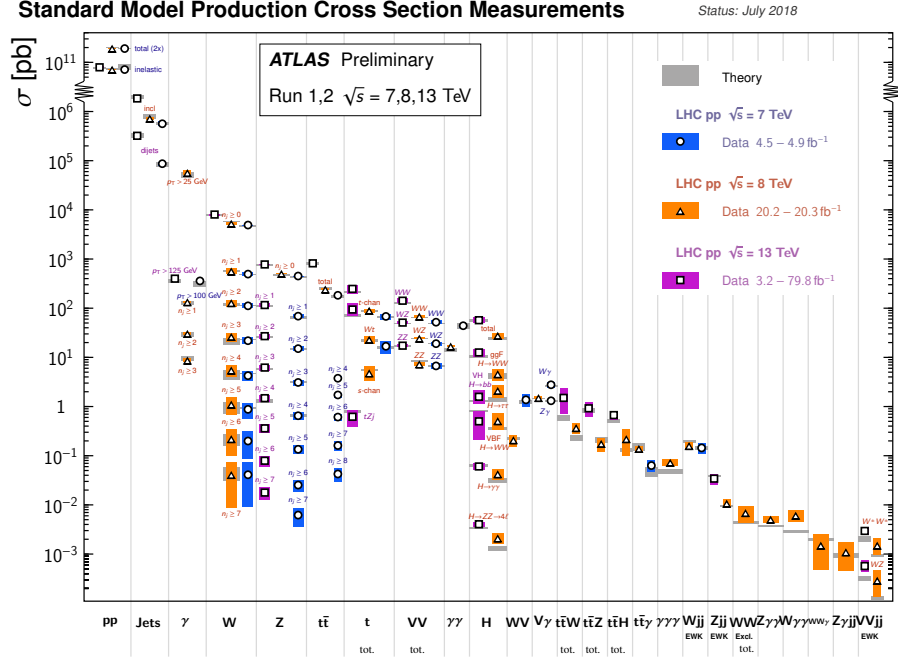


Figure 1.2 – Summary of several Standard Model total and fiducial production cross section measurements at ATLAS compared with the corresponding theoretical expectations. (Figure from [3].)

the weak neutral-current interaction is mediated by the electrically neutral Z boson.

QED is considered to be one of the most precise physical theories, of which the theoretical predictions match the experimental results up to $\mathcal{O}(10^{-11})$. QED follows the U(1) Abelian gauge symmetry and describes the electromagnetic interactions. The electromagnetic interaction is unified with the weak interaction by the electroweak (EW) theory [6]–[8], where a $SU(2)_L \times U(1)$ gauge symmetry is established, coming with four spin-1 vector bosons as the mediators: photon, W^+ , W^- and Z. The generator of the U(1) group is denoted as the weak hypercharge Y, while the $SU(2)_L$ group raises three generators τ^a ($a=1,2,3$) built upon the Pauli matrices. The two sets of generators are linked by the relation $Y = 2(Q - \tau_3)$, where Q is the electrical charge. The Lagrangian of the EW theory can be written down as following, using the two sets of generators mentioned above:

$$\mathcal{L}_{EW} = -\frac{1}{4}W_{\mu\nu}^a W_a^{\mu\nu} - \frac{1}{4}B_{\mu\nu} B^{\mu\nu} + \sum \bar{\psi}_i \gamma^\mu (i\partial_\mu - \frac{1}{2}g'YB_\mu - \frac{1}{2}g\tau_L^a W_\mu^a)\psi \quad (1.1)$$

Here g and g' represent the weak and electromagnetic couplings, respectively. B_μ and W_μ^a are the U(1) and SU(2) gauge fields. $B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu$ and $W_{\mu\nu}^a = \partial_\mu W_\nu^a - \partial_\nu W_\mu^a + g\epsilon^{abc}W_{b\mu}W_{c\nu}$ are the field strength tensors. The gauge boson for the electromagnetic interaction (photon) and those three for weak interactions are formulated as the linear combinations of the gauge fields B_μ and W_μ^a :

$$A_\mu = B_\mu \cos\theta_W + W_\mu^3 \sin\theta_W \quad (1.2)$$

$$Z_\mu = -B_\mu \sin\theta_W + W_\mu^3 \cos\theta_W \quad (1.3)$$

$$W_\mu^\pm = \frac{1}{\sqrt{2}}(W_\mu^1 \mp W_\mu^2) \quad (1.4)$$

θ_W is known as the weak mixing angle, which is associated to the masses of weak gauge bosons via $\cos\theta_W = m_W/m_Z$. The U(1) symmetry dictates the photon (A_μ) to be exactly massless, leading to the equation of $e = g\sin\theta_W = g'\cos\theta_W$. e is the unit charge.

1.1.2 Strong interaction

In the SM, the strong interaction is described by the Quantum Chromodynamics (QCD) theory. The SU(3) gauge symmetry of QCD leads to the eight generators, corresponding to the eight massless gluons carrying a quantum number of color charge. The three types of color charges, denoted as red (r), green (g) and blue (b) also apply to each flavor of the quarks, such that the strong interaction is interpreted to be the exchange of gluon between quarks. The Lagrangian of QCD is written as:

$$\mathcal{L}_{QCD} = -\frac{1}{4}F_{\mu\nu}^a F_a^{\mu\nu} + \sum_{flavor} \bar{\psi}_i (i\gamma^\mu (\partial_\mu \delta_{ij} - ig_s F_\mu^a T_{ij}^a) - m_\psi \delta_{ij}) \psi_j \quad (1.5)$$

ψ_i are the Dirac spinors of a quark with a color index and mass m . The F_μ^a are the eight gluon fields, T_{ij}^a are the generators of SU(3) group. Using the Gell-Mann matrices, $F_{\mu\nu}^a = \partial_\mu F_\nu^a - \partial_\nu F_\mu^a + g_s f^{abc} F_\mu^b F_\nu^c$ are the gluon field strength tensors, where the coefficients f^{abc} are the structure constants of the SU(3) group. g_s is the strong coupling constant and is universally shared by all the eight gluons. The first term of the Lagrangian describes the gluon self-interaction, due to the presence of which the strong interaction cannot become a long range interaction. The second term describes the quark-gluon interaction. The coupling constant of QCD is expressed in the form of running coupling constant $\alpha_S = g_s^2/4\pi$:

$$\alpha_S(Q^2) = \frac{12\pi}{(11n_c - 2n_f) \ln(Q^2/\Lambda_{QCD}^2)} \quad (1.6)$$

Q indicates the energy scale, or in another word, the energy transferred in the interaction. $n_c=3$ corresponds to the number of colors involved in the interaction. n_f is the number of flavors of quarks involved in the interaction, which varies between 3 to 6 depending on the energy scale. $\Lambda_{QCD} \approx 200 \text{ MeV}$ is the QCD energy scale, below which the perturbative calculation of the QCD is no longer valid, bringing a transition from the regime of perturbative QCD to quark confinement. Based on this equation, two important features of the QCD can be derived [2]:

- Color confinement: $\alpha_S(Q^2) \rightarrow \infty$ when $Q^2 \rightarrow 0$. For $|Q| \approx 1 \text{ GeV}$, the interaction between quarks increases rapidly when they are separated in space, such that the quarks are limited to a bound state or the latter stages in the hadronization process instead of being free. Consequently, quarks can only be found in hadrons or quark-gluon-plasma (QGP). Free quarks have never been observed.
- Asymptotic freedom: $\alpha_S(Q^2) \rightarrow 0$ when $Q^2 \rightarrow \infty$. This property suggests that as long as the energy scale is high enough, i.e. $|Q| > 100 \text{ GeV}$ is the typical scale for modern high-energy colliders, the strong coupling constant drops to a sufficiently small value $\alpha_S \approx 0.1$, allowing the application of the perturbation theory. Quarks turn to quasi-free states rather than being bound in the proton. Even though the perturbative expansion of QCD (pQCD) becomes a valid

approach for the calculation of the strong interaction with asymptotic freedom, the contributions beyond the lowest order cannot be neglected, since $\alpha_s \approx 0.1$ is still not so small as the coupling constant of QED.

1.1.3 BEH mechanism and spontaneous symmetry breaking

In order to preserve the $SU(2)_L \times U(1)$ gauge symmetry in the Lagrangian of the EW interaction, all fermions and bosons are supposed to remain massless, which is contradicted by experimental facts. For this reason, the spontaneous symmetry breaking hypothesis of the BEH mechanism is introduced to help explain how the particles in EW processes gain their masses [9]–[12]. In brief, the BEH mechanism combines the EW theory with an additional Higgs field that interacts with all particles, leading to a spontaneous symmetry breaking. Therefore, one extra degree of freedom is given to each gauge boson as a finite mass. The Higgs field ϕ_H itself, is a four-component scalar field that forms a complex doublet of the weak isospin $SU(2)$ symmetry. The Lagrangian of the Higgs part is:

$$\mathcal{L}_H = |(\partial_\mu - igW_\mu^a \tau^a - \frac{1}{2}ig'B_\mu)\phi|^2 - \mu^2 \phi^\dagger \phi - \lambda(\phi^\dagger \phi)^2 \quad (1.7)$$

τ^a are the three generators of the $SU(2)$ symmetry. Once assuming $\lambda > 0$ and $\mu^2 < 0$, the ground state of the Higgs potential becomes non-zero but $|\phi| = \sqrt{-\mu^2/2\lambda}$. Defined by $v^2 = \frac{-\mu^2}{\lambda}$. The vacuum expectation value v sets the scales for the masses of the Higgs boson and the weak gauge bosons. The mass of the Higgs boson is given by:

$$m_H = \sqrt{2\lambda}v \quad (1.8)$$

The coupling terms of ϕ to quadratic W_μ and B_μ give masses to the W and Z bosons:

$$m_W = \frac{1}{2}gv \quad (1.9)$$

$$m_Z = \frac{v}{2}\sqrt{g^2 + g'^2} = \frac{m_W}{\cos\theta_W} \quad (1.10)$$

The masses of the W and Z bosons are linked to the parameters of the Higgs field and the weak mixing angle $\cos\theta_W$, while the fermions gain their masses from the Yukawa couplings between fermion fields and Higgs field according to the following Lagrangian:

$$\mathcal{L}_Y = -g\bar{\psi}\psi\phi = -g\bar{\psi}\psi\frac{v}{\sqrt{2}} \quad (1.11)$$

According to the calculation, v is approximately 246 GeV, which gives the predictions of W-boson mass and Z-boson mass within the SM. However, the mass of the Higgs boson has to be measured experimentally due to unknown value of the free parameter λ . Since the discovery of the Higgs boson in 2012 by ATLAS [13] and CMS [14] experiments, efforts have been made to the measurement of the Higgs boson mass. The current world average of Higgs boson mass, measured by ATLAS and CMS at center-of-mass energies of 7, 8 and 13 TeV, is 125.25 ± 0.17 GeV [15].

1.1.4 Global electroweak fit and W mass

While the SM describes physics phenomena up to high energies at TeV scale and demonstrates impressively the predictive power of the electroweak unification as well as the quantum loop corrections, its overall consistency is also examined by the global electroweak fit. In particular, the mass of the W-boson is a very sensitive quantity for the consistency check of the SM, since its value is tightly constrained by the SM. Therefore, it is possible to obtain a predicted value of W mass from the global electroweak fit as an indirect measurement, which can be compared with the result of direct measurements, providing not only the probe of the SM, but also the excluding power of physics beyond the SM. A precise comparison relies on high precision of both the experiments and the theory predictions. For the latter, the global electroweak fit can nowadays be handled by the gFitter software package [16]–[18] that features the minimisation of a test statistics and its interpretation using frequentist statistics.

The gFitter package adopts a least-square-like notation. The corresponding test statistics is defined to be:

$$\chi^2(y_{mod}) = -2\ln\mathcal{L}(y_{mod}) \quad (1.12)$$

Where the y_{mod} are the theoretical expressions as functions of model parameters. The y_{mod} consists of three subsets: One corresponds to the unconstrained parameters of the theory. The second subset are theoretical parameters for which the fit uses the prior knowledge from direct measurements or from available calculations. The last subset includes those hard-to-quantify theoretical uncertainties, such as higher-order QCD corrections to a truncated perturbative series.

The likelihood function $\mathcal{L}(y_{mod})$ is essentially the product of two contributions

$$\mathcal{L}(y_{mod}) = \mathcal{L}(x_{theo}(y_{mod}) - x_{exp}) \cdot \mathcal{L}_{theo}(y_{mod}) = \prod_{i,j=1}^{N_{exp}} \mathcal{L}_{exp}(i,j) \cdot \prod_{k=1}^{N_{mod}} \mathcal{L}_{theo}(k) \quad (1.13)$$

x_{exp} are the involved measurements considered in the analysis, while x_{theo} are the theoretical expressions used to describe those x_{exp} . The first term on the right-hand side of the Equation 1.13 shows the experimental likelihood, which measures the agreement between the theoretical calculation and the measurements. N_{exp} individual likelihood components for observables are accounted for, of which the model predictions depend on a subset of y_{mod} parameters that come with the constraint power from the prior knowledge. The correlations between measurements as well as the experimental and systematical uncertainties are also taken into account. The second term indicates the theoretical likelihood, containing the prior knowledge of certain y_{mod} parameters. The $\mathcal{L}_{theo}(k)$ terms may be constant if there is no prior information, be bound, or may express a probability function whenever the information is reliable.

The SM tree-level relations and radiative corrections related to W and Z bosons are accounted for in the formalism adopted by gFitter. In particular, the electroweak unification leads to a tree-level relation between the Fermi constant G_F (precisely determined by the measurements of muon lifetime) and the electromagnetic coupling (α):

$$G_F = \frac{\pi\alpha}{\sqrt{2}\left(m_W^{(0)}\right)^2\left(1 - \frac{(m_W^{(0)})^2}{m_Z^2}\right)} \quad (1.14)$$

The superscript "(0)" indicates tree-level quantities. The radiative corrections can be parameterized by multiplying a form factor $(1 - \Delta r)^{-1}$ to the right-hand side of the Equation 1.14. Therefore, the m_W is resolved to be

$$m_W^2 = \frac{m_Z^2}{2} \left(1 + \sqrt{1 - \frac{\sqrt{8}\pi\alpha}{G_F m_Z^2 (1 - \Delta r)}} \right) \quad (1.15)$$

In the SM, the form factor Δr depends nearly quadratically on top-quark mass m_t and logarithmically on m_H due to the virtual quantum loop corrections illustrated in Figure 1.3. Consequently, the physical mass of the W-boson is different from the tree level calculation in Equation 1.14 according to

$$m_W = m_W^{(0)} + a m_t^2 + b \ln \left(\frac{m_H}{m_W} \right) + \dots \quad (1.16)$$

Where a and b are constants determined by calculation, m_t and m_H for the masses of the top quark and Higgs boson [2]. In gFitter, besides the one-loop corrections from m_t and m_H , the effect of Δr also takes into account the full two-loop results [19] and the four-loop QCD calculation. Therefore, according to the SM, the prediction of the W mass is parameterized as a function of Higgs mass, top-quark mass, Z-boson mass and the running QED coupling strength at m_Z^2 [20].

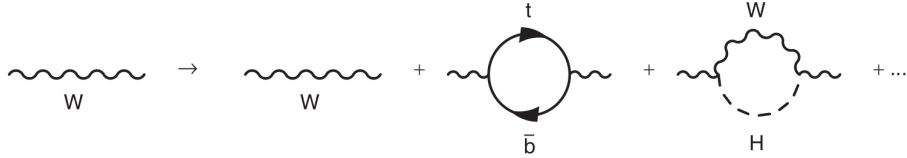


Figure 1.3 – Two largest loop corrections to the W mass. (Figure from [2].)

Five classes of observables are used in the global EW fit:

- Z resonance parameters: Z mass and width, total $e^+e^- \rightarrow Z \rightarrow \text{hadron}$ production cross-sections.
- Partial Z cross sections: Ratios of leptonic to hadronic cross-sections, and heavy-flavour hadronic to total hadronic cross-sections.
- Neutral current couplings: Effective weak mixing angle, as well as the left-right and the forward-backward asymmetries for universal leptons and heavy quarks.
- W-boson parameters: W mass and width.
- Higgs boson parameter: Higgs mass.
- Additional input parameters: Heavy-flavour quark masses (m_c , m_b and m_t . The masses of lighter quarks along with leptons are fixed to the world average values), QED and QCD coupling strengths at the Z-mass scale.

The most up-to-date global EW fit with gFitter uses the data of electroweak precision measurements at the Z pole by the LEP and SLD collaborations [21], the world averages of running quark masses, of W mass ($m_W = 80379 \pm 13 \text{ MeV}$) and W width and the latest determination of the five quark hadronic vacuum polarisation contribution to $\alpha(m_Z^2)$, known as $\Delta\alpha_{had}^{(5)}(m_Z^2)$. The mass of the top quark $m_t = 172.47 \pm 0.46 \text{ GeV}$, is taken from the average of direct measurements by the ATLAS and CMS experiments based on 7 and 8 TeV data, assigned with an additional theoretical uncertainty of 0.5 GeV due to the ambiguity in the kinematic top-mass definition, the

colour structure of the fragmentation process and the finite perturbative order of the relation between the pole and the $\bar{M}S$ mass currently known. The experimental input of the LHC Run 1 Higgs mass $m_H = 125.09 \pm 0.21 \pm 0.11$ GeV, is the weighted average of ATLAS and CMS results.

On the theoretical side, in addition to the aforementioned parameterized prediction of the W mass, the width of the W-boson is known up to one electroweak loop order. Dominant contributions from final-state QED and QCD radiations are also included in the calculations [20].

The indirect determination of W mass by the fit is $m_W = 80.354 \pm 0.007$ GeV, receiving uncertainty contributions from the measurements and the theoretical uncertainty of top-quark mass, the uncertainties of Z-boson mass, strong coupling constant, Higgs mass, and theoretical uncertainty of the unknown higher-order corrections to W mass $\delta_{th} m_W \approx 4$ MeV.

One crucial consistency check of the SM is the simultaneous indirect determination of m_W and m_t , with direct measurements of m_H included or excluded. The corresponding plot is shown in Figure 1.4, in which it can be appreciated that the indirect determination of m_W outperforms the world average of direct measurements. The additional symmetries and fields in the extensions to the SM, such as the Minimal Supersymmetric Standard Model (MSSM) [22], the Two-Higgs-Doublet Model (2HDM) [23] and the dark photons [24], would possibly modify the W mass with respect to the SM prediction. However, no significant deviation from the SM has so far been fully confirmed from the determination of the W mass.

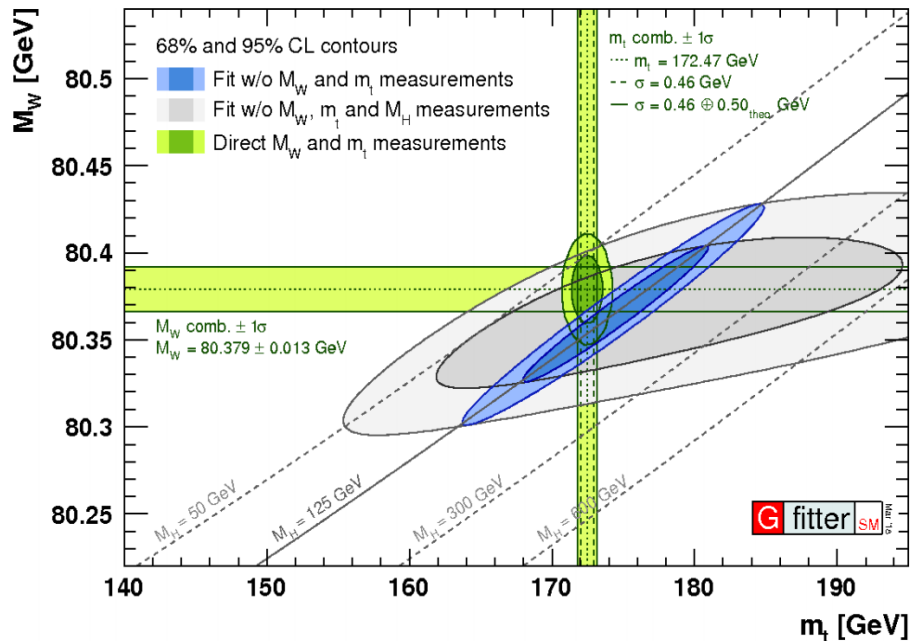


Figure 1.4 – The consistency of the SM is probed by the scanning m_W versus m_t for the global EW fit including or excluding m_H measurement, while excluding m_W and m_t from the fits. 68% and 95% CL contours are shown. (Figure from [20].)

1.2 Weak boson production in proton-proton collisions

Weak bosons can be produced from high-energy proton-proton collisions. By studying the quantum numbers, it appears as if each proton is made of two up quarks (u) and one down quark (d). This composition of "uud" forms the valence quarks in a proton. However, in the scope of inelastic scattering of the collision, the proton turns into a complicated dynamic structure, where the gluons that link the valence quarks through strong interaction can also split into virtual quark-antiquark pairs, known as sea quarks, or produce more gluons via gluon self-interaction. The quarks and gluons in this dynamic structure of the proton are considered to be point-like constituents. The high-energy proton-proton collisions are essentially the interactions between partons from the two protons colliding against each other.

The scattering processes can be divided into soft processes and hard processes by the amount of transferred transverse momentum following the QCD. In the soft processes with low values of Q^2 (i.e. $Q^2 < \Lambda_{QCD}^2$), such as single diffraction, double diffraction, multi-partonic interactions, soft initial state radiation (ISR) and final state radiation (FSR), and beam remnants at a hadron collider. In contrast, the productions of vector bosons are the typical hard processes, where the annihilation of a quark and an anti-quark is involved, coming with a high value of Q^2 . The hard processes can be computed perturbatively in the context of QCD, while the prediction of soft processes has to take into account the non-perturbative effects. In a collision event at hadron collider, a hard process is accompanied by soft processes. The collection of all these soft processes in that event is noted as the underlying event (UE).

The production of vector boson cross-section in a proton-proton collision can be calculated according to the QCD factorization theorem. As constructed by Drell and Yan [25], the process is factorized into the hard scattering of the partons multiplied by the momentum distribution of partons inside the proton $f_i(x, Q^2)$ (a.k.a. Parton Distribution Functions or PDFs, functions of the fraction of the momentum of the proton carried by the struck quark x at a given energy scale Q^2). The PDFs describe the non-perturbative QCD of the soft interactions and have to be extracted by experiments. Once evolved to a large enough energy scale, denoted as $Q^2 = \mu^2$, the cross-section of the $q\bar{q} \rightarrow V$ is expressed as:

$$\sigma_{pp \rightarrow V} = \sum_{a,b} \int dx_a dx_b f_{a/A}(x_a, \mu_F^2) f_{b/B}(x_b, \mu_F^2) \times \sigma_{ab \rightarrow V}(x_a p_1, x_b p_2, \mu_F^2) \quad (1.17)$$

where a and b denote the quark and the anti-quark involved in the scattering, A and B stand for the two colliding protons. The spurious factorization scale μ_F helps separate the non-perturbative factors from the total cross-section, curing the infrared divergences caused by a particle reaching zero momentum or a massless particle radiating another massless particle. At high energy scale, the strong coupling constant $\alpha_S \ll 1$, which allows the expansion of the hard QCD process, for instance:

$$\sigma_{ab \rightarrow V} = \sigma_0 + \alpha_S(\mu_R^2)\sigma + \alpha_S^2(\mu_R^2)\sigma^2 + \mathcal{O}(\alpha_S^3) \quad (1.18)$$

where the μ_R is the renormalization scale of the running QCD coupling, another spurious parameter introduced to resolve the ultraviolet divergence caused by large

momentum in the loops of the Feynman diagrams representing the amplitude. The first term σ_0 stands for the cross-section calculated at Leading Order (LO). The $\alpha_S(\mu_R^2)\sigma$ term and the $\alpha_S^2(\mu_R^2)\sigma^2$ term are the Next-to-Leading Order (NLO) and Next-to-Next-to-Leading Order (NNLO) calculations. In the perturbative theory, as long as the calculation of the cross-section is expanded to all orders of α_S , the result will no longer depend on the scale parameters μ_F and μ_R . However, with a finite computing power in reality, the calculation of the cross-section can only be carried out up to a certain level of accuracy, such as NLO or NNLO. Therefore, the choice of μ_F and μ_R has an impact on the outcome of the perturbative QCD calculation. The typical choice of the scale parameters for a Drell-Yan process is $\mu_F = \mu_R = m_V$, where m_V is the mass of the vector boson of interest, either W or Z.

In the center-of-mass frame for the hard scattering, neglecting the mass of the partons, the four-momenta of the two incoming partons can be written as:

$$p_1 = \frac{\sqrt{s}}{2}(x_1, 0, 0, x_1), p_2 = \frac{\sqrt{s}}{2}(x_2, 0, 0, x_2) \quad (1.19)$$

x_1 and x_2 are related to $\hat{s}$ by $\hat{s} = (p_1 + p_2)^2 = x_1 x_2 s$. The rapidity of the produced vector boson is $y = 1/2(\ln(x_1/x_2))$, which leads to:

$$x_1 = \frac{m_V}{\sqrt{s}}e^y, x_2 = \frac{m_V}{\sqrt{s}}e^{-y} \quad (1.20)$$

Such that the cross-section at leading order is:

$$\frac{d\sigma}{dy} = \frac{1}{s} \sum \sigma_0(m_V) f_1(x_1, m_V^2) f_2(x_2, m_V^2) \quad (1.21)$$

The partonic cross-section of W and Z production can be derived from:

$$\begin{aligned} \sigma^{q\bar{q}' \rightarrow W} &= \frac{\sqrt{2}\pi G_F m_W^2}{3} |V_{q\bar{q}'}|^2 \delta(\hat{s} - m_W^2), \\ \sigma^{q\bar{q}' \rightarrow Z} &= \frac{\sqrt{2}\pi G_F m_Z^2}{3} (v_q^2 + a_q^2) \delta(\hat{s} - m_Z^2) \end{aligned} \quad (1.22)$$

Here $\hat{s}$ is the square of the parton center-of-mass energy, $V_{q\bar{q}'}$ is the appropriate CKM matrix element, corresponding to the strength of the flavour-changing weak interaction. v_q and a_q are the vector coupling and axial vector coupling of the Z-boson to the quarks.

At leading order, the transverse momentum of the produced vector boson is zero. At higher orders, additional quarks and gluons come into play in the production of vector bosons, specifically the soft processes between partons that lead to the gluon radiations and the emissions or virtual loops of gluons and quarks in the hard process. The involvement of additional quarks and gluons brings a non-zero transverse momentum distribution to the produced vector boson. The hard process part of the higher order effect can be predicted by pQCD calculations. For instance, the NLO approximation of the differential cross-section at high p_T^W is:

$$\left(\frac{d\sigma}{d(p_T^W)^2} \right)_{\mathcal{O}(\alpha_S)} = \frac{4\alpha_S\sigma^0}{3\pi^2} \frac{1}{(p_T^W)^2} \ln \left(\frac{m_W^2}{(p_T^W)^2} \right) \quad (1.23)$$

However, it is obvious that this formulation of cross-section does not apply to the low p_T^W region. i.e. the logarithmic term expands rapidly for $p_T^W < m_W$ and compensates for the small α_S , causing the divergence of the higher-order terms. Particularly, at order N , there is a series of terms proportional to $\alpha_S^N \ln^M(m_W^2/(p_T^W)^2)$ in the differential cross-section, where $M \leq 2N - 1$. As for $p_T^W \ll m_W$, the logarithmic term becomes the leading contribution to the calculation, which makes the perturbative calculation become inapplicable. The contribution from soft process can be accounted for by modelling the gluon radiations (parton shower, or PS) based on the knowledge of soft QCD effects accumulated from the fixed target proton-neutron collisions. In an analytic approach, following the argument that the cross-section integrated over all $(p_T^W)^2$ should be finite [26], the divergence can be resolved by resumming all the logarithms to the leading logarithm by the following form:

$$\left(\frac{d\sigma}{d(p_T^W)^2} \right)_{\mathcal{O}(\alpha_S)} = \frac{d}{d(p_T^W)^2} \left\{ \sigma^0 \exp \left[-\frac{2\alpha_S}{3\pi} \ln^2 \left(\frac{m_W^2}{(p_T^W)^2} \right) \right] \right\} \quad (1.24)$$

Such that the full expression of the differential cross-section becomes:

$$\frac{d\sigma}{d(p_T^W)^2} = \left(\frac{d\sigma}{d(p_T^W)^2} \right)_{pQCD} + \left(\frac{d\sigma}{d(p_T^W)^2} \right)_{resumm} \quad (1.25)$$

In low p_T^W region, the contribution from soft effects are corrected by the resummation techniques. Moving to the high p_T^W region, the hard emissions can be predicted by pQCD at a given fixed order.

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⁶⁷¹ Chapter 2

⁶⁷² The LHC and the ATLAS experiment

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2.1 The Large Hadron Collider

The Large Hadron Collider (LHC) at CERN is by far the most powerful tool for particle physics research [1]. Mainly designed to collide proton beams with a centre-of-mass energy of 14 TeV at maximum, it can reach an unprecedented luminosity of $10^{34} \text{ cm}^{-2} \text{ s}^{-1}$. The LHC is a two-ring-superconducting-hadron accelerator and collider installed in the tunnel of its predecessor: Large Electron Positron Collider (LEP), of which the circumference is 26.7 km (Fig 2.1). The existing LEP tunnel is composed of 8 straight sections and 8 arcs buried between 45 m and 170 m underground. Linked by two transfer tunnels, the tunnel of LHC is connected to the CERN accelerator complex that functions as an injector.

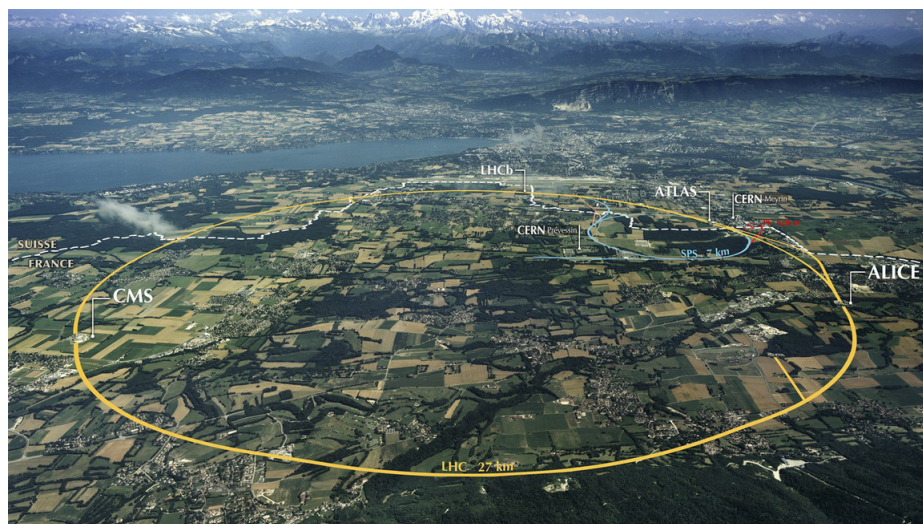


Figure 2.1 – A photo illustrating the LHC main tunnel and four major experiments.

The so-called "lattice magnets" on the LHC bend and tighten the trajectory of particles and are responsible for keeping the beams stable and precisely aligned. The LHC is equipped with 1232 main dipoles, each 15 m long and as heavy as 35 t, for

handling tighter turns of the beam. The dipoles make use of superconducting magnet technology and generate a peak dipole field of 8.33 T at the maximum proton beam energy of 7 TeV. These dipoles cooperate with multipoles correctors such as sextupole, octupole and decapole magnets, so that the small imperfections in the magnetic field at the extremities of the dipoles can be corrected. In addition, the quadrupoles made of four magnetic poles around the beam pipe squeeze the beam either vertically or horizontally, keeping the particles in a tight beam as a net effect, such that the RMS of the beam size in the arcs is kept at the level of 0.3 mm during collisions. The "insertion magnets" take over the beam in front of the LHC detectors. These insertion magnets are deployed in the form of inner triplet, which is a system based on three quadrupoles. There are eight inner triplets at LHC in total, every two of which are located at each of the four large LHC detectors: ALICE, ATLAS, CMS and LHCb. Thanks to the inner triplets, particles are squeezed closer together before they reach a detector, which increases the chance of colliding with particles coming from the opposite direction.

All the work of this thesis, covering both the Quality Assurance and Quality Control (QAQC) of particle detectors and the study of W-boson properties, is based on the ATLAS experiment. A Toroidal LHC Apparatus (ATLAS), is a multi-purpose particle detector of the Large Hadron Collider designed to collect data from high-energy proton-proton or heavy ion collisions. This chapter will first give a brief introduction to the principle of particle acceleration at LHC, as well as the key parameters to describe the particle beams at collisions. The second part of the chapter will review all the sub-systems of ATLAS, followed by a detailed description of the specific datasets used for W-boson physics analysis.

2.1.1 The acceleration system of LHC

The injected protons supplied to the LHC come from the CERN accelerator complex. The protons are initially accelerated by the Linac2 to reach 50 MeV, then transferred to the Proton Synchrotron Booster (PSB) where their energy increases to 1.4 GeV. As the last part of the pre-injector chain, the Proton Synchrotron (PS) receives beams from the PSB and further accelerate them to 25 GeV. After that, the Super Proton Synchrotron (SPS) takes over the acceleration and injects the proton beam to LHC at 450 GeV. Following the injection of the beam, the protons are accelerated at LHC up to an energy of 7 TeV via the Radio Frequency (RF) systems.

Inside a proton beam of the LHC, instead of spreading uniformly along the circumference, the particles are grouped in structures called bunches. The acceleration system operates coordinately, meeting the requirement of holding more than 2800 proton bunches per LHC ring with a nominal bunch spacing of 25 ns while achieving small transverse and well defined longitudinal emittances. The high center-of-mass energy in the collisions increases the cross-sections of hard-scattering processes in general and opens a wider kinematic window for rare physics processes.

2.1.2 Proton beam and luminosity

Two proton beams circulate in opposite clock-wise directions in the tunnel of the LHC, manipulated by separate magnet dipole fields and vacuum chambers in the main arcs. The two equally charged particles beams enter common sections only at

the insertions regions reserved for detectors. The number of events per unit of time generated in proton-proton collisions at LHC is described by the following relation [2]:

$$N_{event} = L\sigma_{event} \quad (2.1)$$

Here σ_{event} is the cross section of the process of interest, and L is the machine luminosity determined by the beam parameters. Assuming a Gaussian beam distribution, we have:

$$L = \frac{N_b^2 n_b f_{rev} \gamma_r}{4\pi \epsilon_n \beta_*} F \quad (2.2)$$

where N_b stands for the number of protons per bunch, n_b is the number of bunches per beam, f_{rev} for the revolution frequency, γ_r for the relativistic factor, ϵ_n for the normalized transverse beam emittance and β_* is the beta function at the interaction point (IP). The beta function describes the envelope of the single particle trajectory, which has to be squeezed to a minimum at the IP. The emittance is an invariant determined by the optics of the machine. It describes the space occupied by the particle in the transverse two-dimensional phase space. Due to the non-zero crossing angle of the two colliding beams at the IP, the luminosity is reduced according to the factor

$$F = \frac{1}{\sqrt{1 + \left(\frac{\theta_c \sigma_z}{2\sigma^*}\right)^2}} \quad (2.3)$$

in which θ_c , σ_z and σ^* are the full crossing angle at IP, the RMS bunch length and the transverse RMS beam size at IP, respectively.

As high luminosity experiments, ATLAS and CMS require a peak luminosity of $L = 10^{34} \text{ cm}^{-2} \text{ s}^{-1}$. The limit of the LHC machine performance depends on multiple factors, such as the beam-beam interactions, the aperture of the focusing magnets on the two sides of the IP, the maximum dipole field as well as the magnet quench limits, the capability of the beam dumping system and so on.

Continuous data taking at high instantaneous luminosity leads to a high integrated luminosity, therefore helping accumulate more statistics of interactions at a given center-of-mass energy. Together with the high beam energy, the high beam intensity also facilitates the exploration of rare processes.

A precise measurement of the integrated luminosity is a key part of the physics measurements at LHC. In particular, the uncertainty on the luminosity is an important concern in precision measurements.

2.2 The ATLAS detector

ATLAS [3] is a cylindrical multi-purpose particle detector designed to be forward-backward symmetry. The detector sub-systems are comprised of the inner tracking detector, the electromagnetic and hadronic calorimeters and the muon spectrometer, as shown in Fig 2.2. In addition to the detectors, there is a thin superconducting solenoid surrounding the inner detector, as well as three large superconducting toroid magnets (one for the barrel and two for the end-caps) distributed in an eight-fold azimuthal symmetric pattern around the calorimeters.

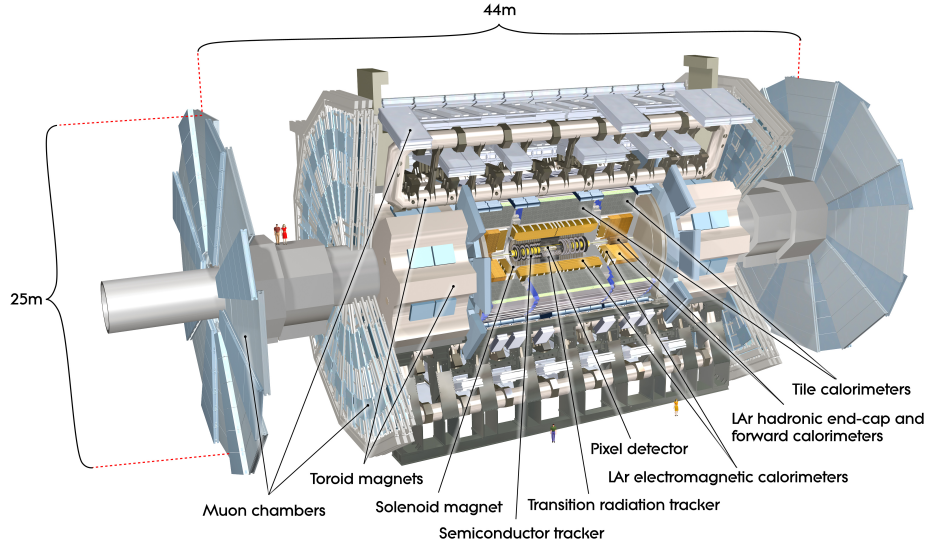


Figure 2.2 – An overview of the ATLAS magnet and detector sub-systems. (Figure from [3].)

2.2.1 The coordinate system in ATLAS

Based on its cylindrical geometry, ATLAS uses a right-handed coordinate system centering at the Interaction Point (IP). The z -axis is parallel to the beam pipe while the x - y plane is perpendicular to the beam. The positive direction of the x -axis points towards the center of LHC. The positive direction of the y -axis is defined to be pointing upwards. The cylindrical coordinates are introduced to the x - y plane to simplify the description of particle kinematics, with ϕ denoted to be the azimuthal angle and θ as the polar angle.

Benefiting from the cylindrical coordinate system of ATLAS, the rapidity of the particle is frequently used to study the kinematics, since the calculation of rapidity is additive under Lorentz transformations. For a particle with energy E and momentum p_z along the z -axis, the definition of rapidity is as following:

$$y = \frac{1}{2} \ln \frac{E + p_z}{E - p_z} \quad (2.4)$$

When the mass of the particle is negligible to its momentum, pseudorapidity is used for approximation:

$$\eta = -\ln [\tan(\theta/2)] \quad (2.5)$$

The pseudorapidity-azimuthal angular distance between two particles is defined as:

$$\Delta R = \sqrt{\Delta\eta^2 + \Delta\phi^2} \quad (2.6)$$

2.2.2 Magnet system

The magnet system of ATLAS is a hybrid of four superconducting magnets, stretching 22m in diameter and 26m in length. The ATLAS magnet system consists of one solenoid, one barrel toroid and two end-cap toroids. A volume of about 12 000 m³ is covered by the magnetic field provided by the magnet system.

The central solenoid is aligned on the beam axis, producing a 2 T axial magnetic field for the inner detector. In order to minimise the impact on the calorimeters outside, the material budget of the solenoid is kept as low as possible: the inner and outer diameters are 2.46 m and 2.56 m, respectively. The axial length of the solenoid is 5.8 m.

The barrel toroids comprises eight coils enclosed by individual vacuum-vessels. They generate a magnetic field that fills the cylindrical volume surrounding the calorimeters and two end-cap toroids. The total size of the barrel toroid system after installation is 25.3 m in length, 9.4 m and 20.1 m in inner and outer diameters.

Each of the two end-cap toroid, with 5 m in axial length and 10.7 m in outer diameter, contains eight flat, squared coil units in a rigid structure. The end-cap toroid coil system is rotated by an offset of 22.5° with respect to the barrel toroid coil system so as to provide radial overlap and to optimize the bending power at the interface between the two coil systems.

The toroid system provides the muon detectors a toroidal magnetic field, of which the magnitude is around 0.5 T in the central region and approximately 1 T in the end-cap regions.

2.2.3 Inner detectors

As the innermost detector sub-system of ATLAS, the Inner Detector (ID) is capable of providing pattern recognition, measuring momentum, determining both primary and secondary vertex for charged tracks within $|\eta| < 2.5$. Additional electron identification is also provided by ID within $|\eta| < 2$. The ID is located within a cylindrical coverage with 1150 mm in radius and 3512 mm in length (Fig 2.3). Thanks to the enclosing thin solenoid, there is an axial magnetic field of 2 T throughout the entire ID. As illustrated in Fig 2.4, the ID is made of three different types of sub-detectors: the silicon-pixels and the silicon micro-strip sensors (SCT) at inner radii, as well as the transition radiation tracker modules (TRT) at larger radii. The three sub-detectors operate independently but are complementary to each other (Fig 2.5).

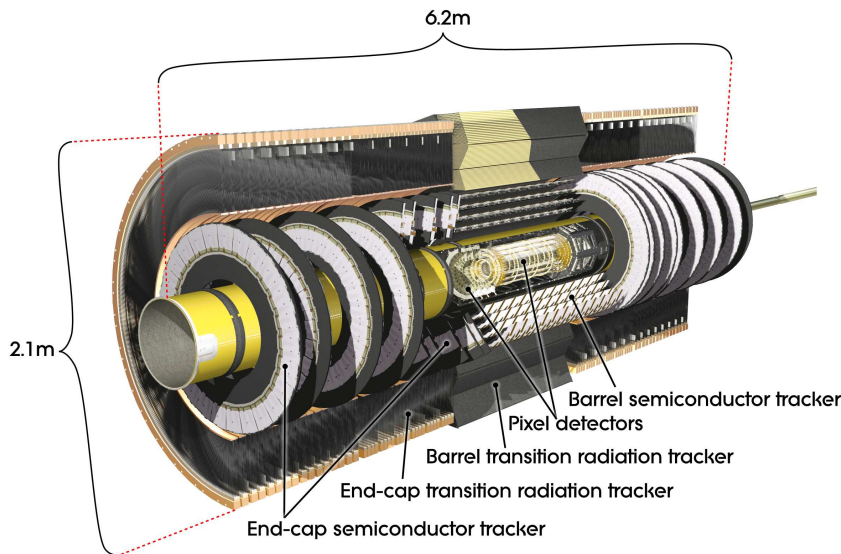


Figure 2.3 – Cut-away view of ATLAS Inner Detector. (Figure from [3].)

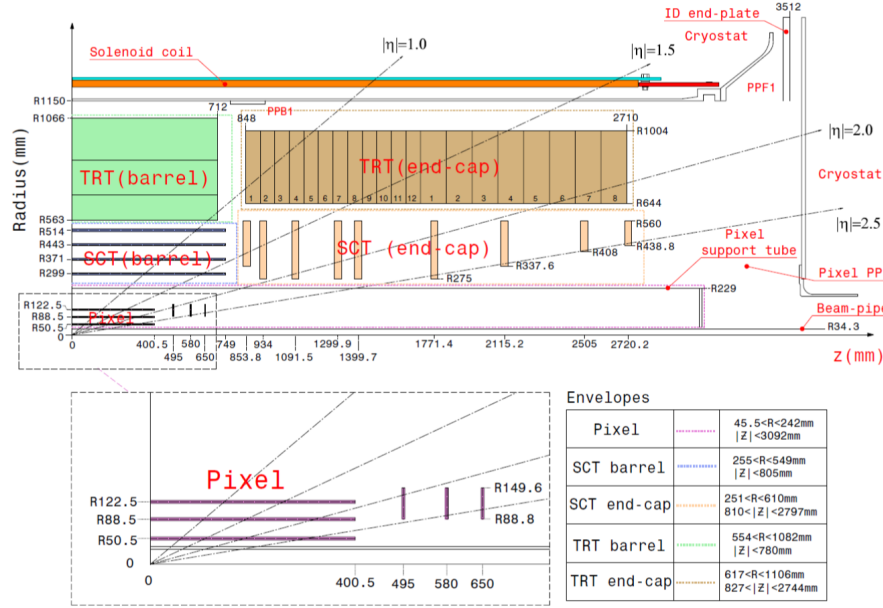


Figure 2.4 – Sectional view of ATLAS ID. (Figure from [3].)

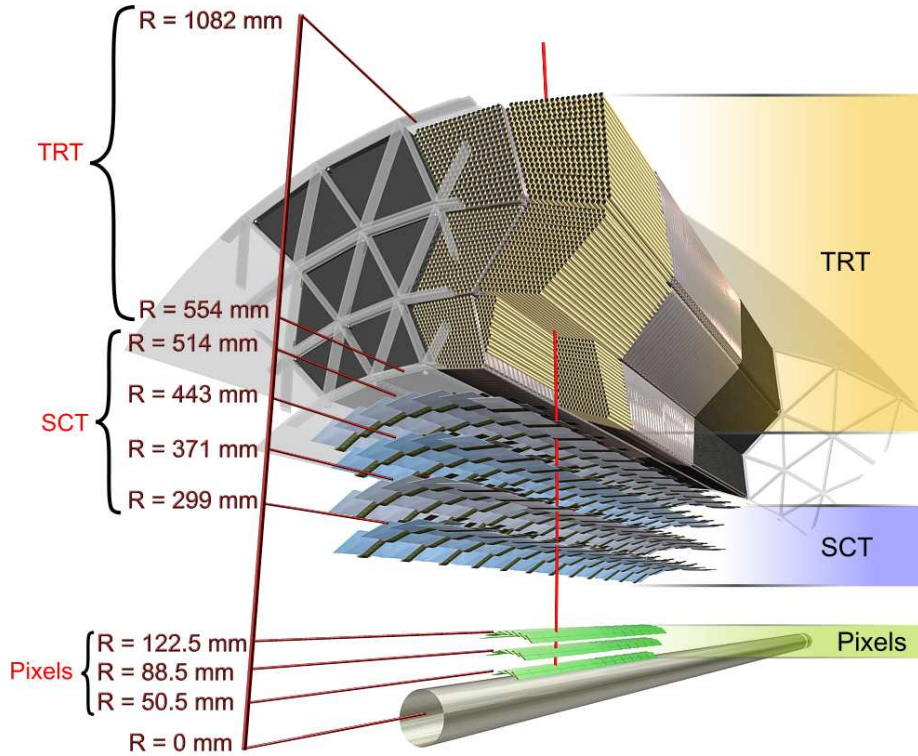


Figure 2.5 – ATLAS ID transversed by a charged particle. (Figure from [3].)

832 The working principle of silicon detectors is that when the charged particles pass
 833 through the depleted region in a PN-junction, the ionized free electrons and holes will
 834 quickly drift towards the electrodes under the external electric field, inducing signals
 835 at the same time [4]. The silicon pixel layers and the SCT layers closer to the beam
 836 pipe are designed to realise high-resolution pattern recognition abilities. The 1744
 837 modules of pixel detectors are arranged in three barrel layers at $R=50.5$ mm, 88.5 mm,
 838 122.5 mm and in three disk layers on each of the two end-caps at $|z|=495$ mm, 580 mm,

650 mm. The intrinsic accuracy of pixel detectors is $10\text{ }\mu\text{m}$ in $R - \phi$ direction and $115\text{ }\mu\text{m}$ in z (R) direction for barrel layers (disks). An extra pixel layer, called the Insertable B-Layer (IBL) [5], was inserted between the original pixel detectors and a new beam pipe during the shutdown of LHC in 2013-2014, aiming at improving the track reconstruction performance of the ATLAS ID system at a higher luminosity.

Covering a range of radius from 275 mm to 560 mm, the 4088 SCT modules are placed on four tilting coaxial cylindrical layers in the barrel regions and on nine disk layers in the end-cap regions. The SCTs provide an intrinsic tracking accuracy of $17\text{ }\mu\text{m}$ in $R - \phi$ and $580\text{ }\mu\text{m}$ in z (R) for barrel layers (disks).

The TRT is installed in the outermost regions of the ID. The basic TRT detector elements are polyimide drift tubes (straw tubes) with 4 mm in diameter. The low energy transition radiation photons can be absorbed by the Xe-based working gas of straw tubes, inducing a much larger signal than what is caused by the ionization of charged particle. As a result, the identification of electron is enhanced by distinguishing transition radiation from tracking signals. 73 layers of straw tubes are axially arranged in the barrel region while 160 straw tube planes are installed radially in each of the end-cap regions. Fibres or foils are inserted among the straw tubes in the barrel or end-caps, rendering the emissions of transition radiation from traversing electrons. The inner radius and the outer radius of the TRT are 563 mm and 1066 mm, respectively. The intrinsic accuracy of the TRT is $130\text{ }\mu\text{m}$.

2.2.4 Calorimetry

In order to measure the energies of particles that pass through the ID, the electromagnetic calorimeter and the hadronic calorimeters of ATLAS are designed to absorb most high-energy particles through electromagnetic showering and hadronic showering. The layout of ATLAS calorimeters are shown in Fig 2.6.

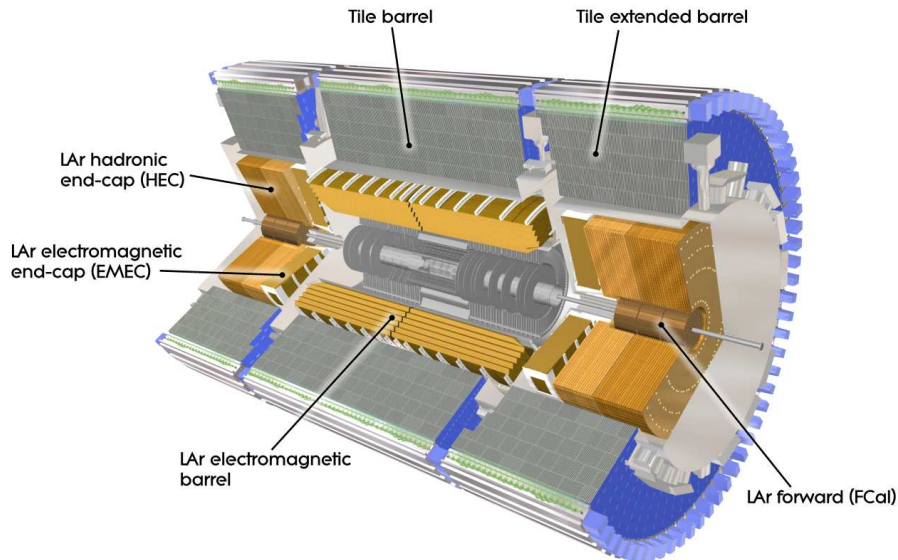


Figure 2.6 – The layout of ATLAS calorimeter systems. The LAr electromagnetic calorimeter in the barrel and on the end-caps measure the energy of electrons and photons through electromagnetic showering. The tile calorimeter, LAr hadronic end-caps and LAr forward calorimeters are designed to measure the energies of hadrons. (Figure from [3].)

Typically, electromagnetic showers happen when a high-energy electron, positron or photon enters a material and interacts with its atomic nucleus. The high-energy electron emits a photon due to bremsstrahlung while the high-energy photon produces a pair of electron-positron in the Coulomb field of the atomic nucleus. Then the electron and positron produced in the pair production further annihilate into a pair of photons. These processes develop a cascade of electron and photons, resulting in an electromagnetic shower. The electromagnetic showers come to an end if the energy of the photons is insufficient for pair production and the bremsstrahlung of secondary electrons is not strong anymore. After that, electrons lose energy via ionization while photons give up energy through Compton scattering and photoelectric effect until finally being absorbed by the material. The radiation length X_0 , depending on the material, is commonly used to describe the longitudinal development of electromagnetic showers. It is defined to be the average thickness of the material that reduces the average energy of high-energy electrons to $1/e$ of its initial value because of bremsstrahlung.

The hadronic showers begin when high-energy hadrons interact with atomic nuclei in the material through strong interaction and produce multiple secondary hadrons. The secondary hadrons, still with high-energy, will carry on the interaction with atomic nuclei and produce even more secondary particles. Even though most secondary particles are pions and nuclear fragments in hadronic showers, protons, neutrons, electrons, photons, muons and neutrinos can also be produced as secondary particles, which leads to the complexity of hadronic showers. The characteristic of longitudinal development of a hadronic shower is the nuclear interaction length:

$$\lambda = \frac{A}{N_A \rho \sigma} \quad (2.7)$$

where A is the atomic number, N_A is the Avogadro constant, ρ is the matter density and σ is the cross-section of inelastic scattering.

2.2.5 LAr electromagnetic calorimeter (EM)

The EM calorimeter of ATLAS is a sampling calorimeter made of lead-Liquid Argon (LAr) detectors with accordion-shaped kapton electrodes and lead absorber plates [6]. The lead absorbers are glued to two stainless-steel sheets with resin-impregnated glass-fibre fabric to improve the mechanical strength. The gaps between absorbers plates are filled by flowing liquid argon. The readout electrodes are also situated in the gaps between the absorbers and are made of three conductive copper layers insulated and separated by polyimide sheets. The two copper layers on the outside provide high voltage potential while the layer in the middle reads out signal through capacitive coupling.

As high energy particles pass through the EM calorimeter, electromagnetic showers take place in the lead absorber. Between the absorbers plates, the shower particles ionize liquid argon, creating ions and free electrons. Since an external electric field is applied to the gap filled by LAr, the ions and the electrons will drift towards the absorber and the electrode, therefore inducing a current signal to the readout electrodes. After converting the signals to energy deposit in individual cells, the energies and positions of the electromagnetic showers are recorded, which determines the interaction point and allows electron/photon identification as well as lepton isolation.

The accordion geometry of the calorimeter guarantees a fast extraction of the signal as well as a full coverage in ϕ without azimuthal cracks. Supported by independent cryostats, the EM calorimeter includes a barrel part covering $|\eta| < 1.475$ and two end-cap parts covering $1.375 < |\eta| < 3.2$. The LAr calorimeter shares the vacuum vessel with the adjacent central solenoid.

The barrel calorimeter is composed of two half-barrels, each of length 3.2 m and with inner and outer diameter of 2.8 m and 4 m, respectively. Between the two half-barrels there is only a small gap at $z = 0$. A half-barrel is divided into 16 modules of which the coverage of azimuthal angle is 22.5° per module. The thickness of a module is $22X_0$ at minimum, then increased from $22X_0$ to $30X_0$ between $0 < |\eta| < 0.8$ and from $24X_0$ to $33X_0$ between $0.8 < |\eta| < 1.3$. One barrel region module has three layers in depth, and contains 3424 readout cells in total including the presampler cells. The first calorimeter sampling is based on the strip cells that have fine granularity in η to measure the pointing of the photon and to facilitate the π_0 rejection. The main sampling for energy measurement comes from the square towers in the second layer of the calorimeter with corresponding sampling size of $\Delta\eta \times \Delta\phi = 0.025 \times 0.025$. The third sampling collects the tail of the electromagnetic shower and has a granularity of $\Delta\eta \times \Delta\phi = 0.05 \times 0.025$. The trigger tower of the barrel module is of the size $\Delta\eta \times \Delta\phi = 0.1 \times 0.1$. The 3-layer sampling structure of a module in the barrel calorimeter as well as the granularity of trigger tower is shown in Fig 2.7.

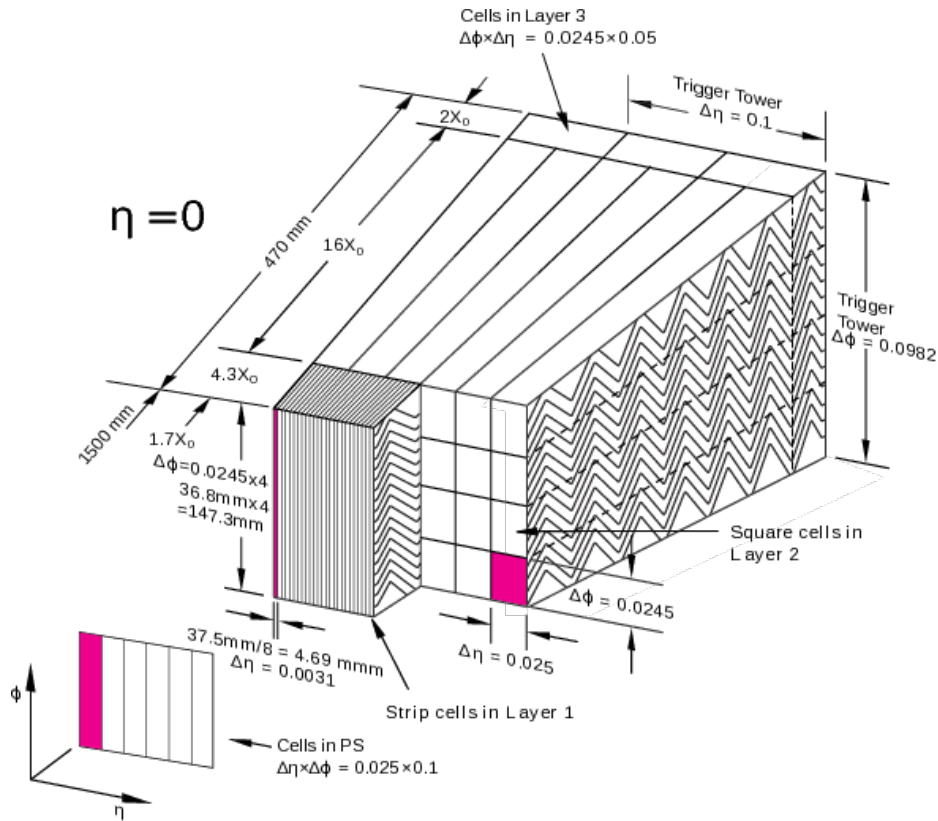


Figure 2.7 – Sketch of a module installed in the barrel EM calorimeter of ATLAS. (Figure from [3].)

The end-cap calorimeter on each side is divided into an outer wheel covering $1.375 < |\eta| < 2.5$ and an inner wheel covering $2.5 < |\eta| < 3.2$. Each end-cap wheel is 63 cm thick and stretches from 330 mm to 2098 mm along its radius. A LAr

presampler is installed in front of the end-cap wheel ranging $1.5 < |\eta| < 1.8$ in order to improve the energy measurement. Above $|\eta| > 1.475$, the total active thickness of the end-cap is more than $24X_0$. Specifically, the thickness increases from 24 to $38X_0$ for η going from 1.475 to 2.5 and increases from 26 to $36X_0$ for η going from 2.5 to 3.2. Similar to the barrel electromagnetic calorimeter, the precision region in the end-cap calorimeter is also divided into 3 layers in depth. The front layer with $4.4X_0$ in thickness is built of strips in η direction. The transverse size of the projective cell in the middle layer matches the precision of its barrel electromagnetic calorimeter counterpart. The $|\eta| < 1.5$ region of the outer wheel as well as the inner wheel is segmented into 2 longitudinal layers, which brings a coarser transverse granularity.

2.2.6 Hadronic calorimeters

The hadronic calorimeters locate beyond the EM calorimeters and consist of three components: the tile calorimeter, the liquid-argon hadronic end-cap calorimeter (HEC) and liquid-argon forward calorimeter (FCal).

The tile calorimeter functions as a sampling calorimeter of which the absorber is steel and the active material is scintillator. The scintillator tiles equipped with wavelength-shifting fibre readout on the edge are aligned radially and normal to the beam line, which achieves an almost full azimuthal coverage. The scintillation light is converted to current and multiplied by photomultiplier tubes (PMTs). The tile calorimeter is placed in $|\eta| < 1.7$ and divided into a 5.8m long central barrel for $|\eta| < 1.0$ along with two 2.6m long extended barrels for $0.8 < |\eta| < 1.7$. Each subdivision has an inner radius of 2.28m and an outer radius of 4.25m, contributing to a radial length of approximately 7.4λ . The segmentation is $1.5, 4.1$ and 1.8λ for the barrel region and $1.5, 2.6, 3.3\lambda$ for the extended barrel region. The dimensions of the sampling cells are $\Delta\eta \times \Delta\phi = 0.1 \times 0.1$ in the first two layers, followed by $\Delta\eta \times \Delta\phi = 0.2 \times 0.1$ in the last layer.

The HEC is made of two separate wheels on each end-cap behind the electromagnetic calorimeter. The HEC extends out to $|\eta| = 3.2$ and overlaps the forward calorimeter to compensate for the drop in material density in the transitional region (between the end-cap and forward calorimeter, around $|\eta| = 3.1$). Inwards, the HEC overlaps the η range of the tile calorimeter by extending to $|\eta| = 1.5$. One HEC wheel is divided into 2 segments along the depth, which makes 4 layers in total on each end-cap. The outer radius of the HEC wheels is 203mm. The inner radius is 47.5mm in general but reduced to 37.2mm in the overlap region with the forward calorimeter. The first segment of the wheel closer to the interaction point are built of 25 mm copper plates while the second segment uses 50 mm copper plates. Between the copper plates, a 8.5 mm gap is filled with LAr as the active medium for the sampling calorimeter. The readout cells correspond to the pads etched on the central foil in these gaps. Below $|\eta| < 2.5$ the size of the readout cells is $\Delta\eta \times \Delta\phi = 0.1 \times 0.1$, while in the region of higher η the granularity of the readout is $\Delta\eta \times \Delta\phi = 0.2 \times 0.2$.

Installed in the same cryostats as the end-cap calorimeters and providing a coverage over $3.1 < |\eta| < 4.9$, the FCal is about 10λ in depth comprised of three layers of modules in each end-cap. The first layer is made of copper and optimized for EM measurements. Meanwhile, the other two layers are built with tungsten, since they mainly focus on the measurement of the energy of hadronic interactions. Every single FCal module contains a metal matrix, where the regularly spaced channels are filled

with electrode structure made of concentric rods and tubes parallel to the beam. Between the rod and the tube, liquid argon functions as the active medium in the gap. Given that the FCal modules are all located at high η only around 4.7 m away from the IP, the LAr gaps are designed to be very small not only to avoid ion build-up issues when exposed to high particle fluxes, but also to maintain the highest possible density at the same time.

2.2.7 Muon spectrometer

The muon spectrometer makes up the outermost part of the ATLAS detector (Fig 2.8). It is dedicated to the detection of charged particles that penetrate the barrel and end-cap calorimeters. The momenta of the particles can be measured in the range $|\eta| < 2.7$, while the trigger on these particles covers the region $|\eta| < 2.4$. The most important performance goal is to achieve a stand-alone transverse momentum resolution of about 10 % for 1 TeV tracks, or equivalently a sagitta along the beam axis of around $500\ \mu\text{m}$ to be measured within a resolution of $50\ \mu\text{m}$.

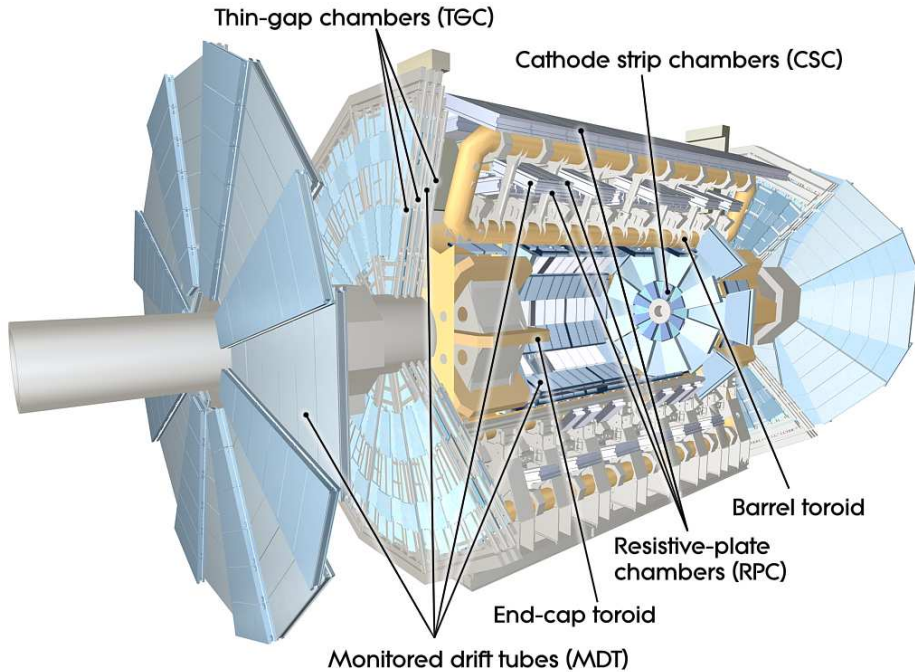


Figure 2.8 – Cut-away view of ATLAS muon system. The toroids are also drawn in the schematics (Figure from [3]).

The operation of the muon spectrometer is based on the magnetic deflection of muon tracks and high-precision tracking chambers combined with separate trigger chambers. Below $|\eta| < 1.4$, the magnetic bending is brought by the large barrel toroid, while between $1.6 < |\eta| < 2.7$ the muon tracks are bent by the end-cap toroids. In the transitional region $1.4 < |\eta| < 1.6$, both the barrel field and the end-cap fields contribute to the magnetic deflection. As for particle detection, in the barrel region, the tracks are measured by three concentric cylindrical shells of chamber around the beam at radii of approximately 5 m, 7.5 m and 10 m. In the transitional and end-cap regions, the muon chambers are also installed in three layers, but in planes perpendicular to the beam axis. The Muon chambers form large wheels in the

end-cap region, located at $|z| \approx 7.4$ m, 10.8 m, 14 m and 21.5 m with respect to the IP. In the central region of the detector around $|\eta| = 0$, a gap is left open to accommodate the services to the solenoid magnet, the calorimeters and the inner detectors.

The purpose of the precision-tracking chamber is to determine the coordinate of the track in the bending (η) plane. The precision momentum measurement relies on the Monitored Drift Tube chambers (MDT's). The MDTs cover the pseudorapidity range of $|\eta| < 2.7$ except for the innermost end-cap layer where its coverage is limited to $|\eta| < 2.0$. Each MDT chamber consists of three to eight layers of drift tubes with a diameter of 29.970 mm. The working gas of drift tube is 93 %Ar + 7 %CO₂ under an absolute pressure of 3 bar, allowing to achieve an average spatial resolution of 80 μ m per tube. A high-precision optical alignment system is also installed to monitor the positions and internal deformations of the MDT chambers, which is accompanied by track-based alignment algorithms.

Because of requirements of high rate capability and time resolution, Cathode-Strip Chambers are used in the innermost tracking layer in the forward region $2 < |\eta| < 2.7$. The CSC detectors are multiwire proportional chambers in which the cathode planes are segmented into strips in orthogonal directions, such that both coordinates can be measured from the induced-charge distributions. The spatial resolution of a CSC chamber of 40 μ m in the bending plane and about 5 mm in the transverse plane.

The capability to trigger on muon tracks is a crucial design criterion for the muon system, according to which the precision tracking chambers have to cooperate with a system of fast trigger chambers that delivers track information in tens of nanoseconds after the passage of the muon. In the barrel region $|\eta| < 1.05$, Resistive Plate Chambers (RPC) is chosen as the fast trigger, while for the end-cap region $1.05 < |\eta| < 2.4$ the choice is Thin Gap Chambers (TGC). Both types of trigger chambers can deliver signals with a spread of 15-25 ns, enabling the tagging of the beam-crossing. Besides the fast triggering, the trigger chambers also measure both η (bending) and ϕ (non-bending) coordinates of the track. Once the matching of the MDT and trigger chamber is complete, the trigger chamber's coordinate in the ϕ plane is used as the second coordinate of the MDT measurement.

In Run 3, the inner end-cap of the muon spectrometer, the previous Small Wheels, are replaced by the New Small Wheels (NSW) that are equipped with sTGC and Micromegas detectors. The upgrade of the inner end-cap ensures a high tracking resolution in the high luminosity environment of LHC. More details about the NSW upgrade is covered in Section 3.1.

2.2.8 The Data Acquisition (DAQ) and trigger system

The bunch-crossing interval as short as 25 ns at LHC makes the trigger a challenging task for ATLAS. It is essentially impossible to fully record all of the data for every proton-proton collision. In addition, the majority of the interactions are the soft processes and therefore not concerned by the physics interest in most analysis. For this reason, the Level 1 hardware trigger (L1) and the High Level Trigger (HLT) make up for the trigger system of ATLAS during Run 2 data taking, which aims at eventually reducing the event rate down to around 1.5 kHz [7]. The ATLAS DAQ system delivered high performance and availability during the Run 2 data taking. A detailed work-flow of ATLAS Run 2 trigger and data acquisition system is presented in Fig 2.9.

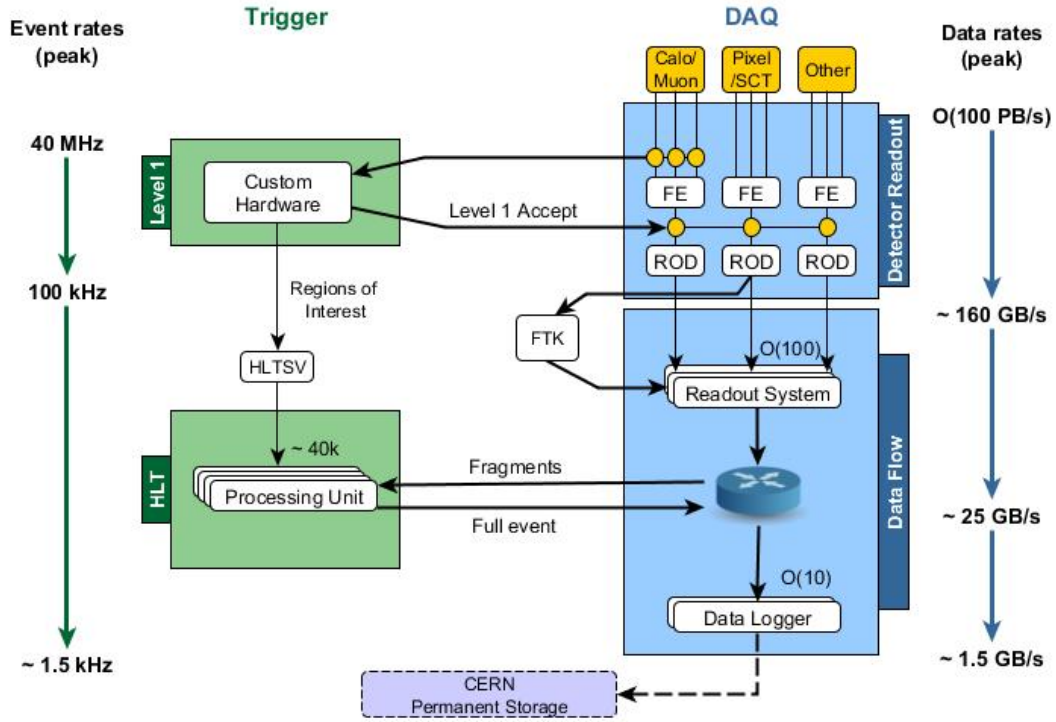


Figure 2.9 – Block diagram showing the work-flow of ATLAS trigger and data acquisition system in Run 2. (Figure from [7].)

The L1 trigger looks for specific signatures from high- p_T muons, electrons/photons, jets as well as τ decaying into hadrons. Events with large missing transverse energy (E_T^{miss}) and large total transverse energy are also selected by L1. The L1 trigger makes use of the reduced-granularity information from a subset of detectors, such as the RPC and TGC for high- p_T muons, the calorimeter sub-systems for electromagnetic clusters, jets, τ -leptons, large E_T^{miss} as well as large total transverse energy. The detector Regions of Interest (RoI) are flagged by the L1 trigger for further event processing performed by the HLT.

The possible trigger objects within an event identified by the L1 trigger is passed down to the HLT, which is seeded by RoI defined by these objects. The HLT will first exploit the information on the coordinates, energy and the type of signatures to limit the amount of data transferred from the detector readout. After that, the HLT continues to select the events based on the offline analysis on fully-built events, therefore reducing the rate to a level that can be recorded for the offline analysis. Full granularity and precision of the calorimeters and the muon chamber data is used by the HLT algorithms to further improve the trigger selections.

The event data from the detector front-end electronics systems are sent to the Readout System (ROS) in response to a L1 trigger acceptance signal. The ROS buffers the event data passing L1 selection. The HLT requests data from the ROS during its processing of each event, then decides either to accept this event and send it to the Data Logger for permanent storage, or to request the deletion of the event from the ROS buffers. The ROS in Run 2 is capable of buffering an input data rate

of 100 kHz while sending 50 % of the data to the HLT without loss of performance,
which comfortably satisfies the ATLAS requirement for Run 2.

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1098 Chapter 3

1099 Cosmic Bench Characterization of
1100 ATLAS NSW LM1 Micromegas
1101 quadruplets

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3.1 NSW upgrade of ATLAS experiment

The steadily increasing luminosity of LHC requires an upgrade to high rate and high-resolution capable detector technology for the inner end cap of the muon spectrometer of the ATLAS experiment. To fulfill this goal, the original two Small Wheels were replaced with New Small Wheels that are more powerful in terms of triggering and precision tracking [1]. In the high background of the upgraded LHC, the maximum hit rate predicted for the inner end-caps of the muon spectrometer will be approximately 15 kHz cm^{-2} [2]. In order to maintain the high precision of muon momentum resolution (ranges from 1.7 % to 4 % depending on transverse momentum as well as rapidity [3]), the spatial resolution of around $100 \mu\text{m}$ for track reconstruction is required.

The NSW sub-detector system is composed of two wheel-like structures of 10 m in diameter, as shown in Figure 3.1. Each wheel includes 8 small sectors and 8 large sectors. While the small sectors are settled towards the ATLAS Interaction Point (IP), the large sectors are installed beyond the small sectors, further away from the IP. The coverage area in radius is from 90 to 445 cm for small sectors and from 92 to 465 cm for large sectors [4]. Besides being compatible with the existing tracking detectors and the end-cap alignment system of ATLAS as well as having high-precision trigger and tracking capability, the layout of NSW also requires a certain degree of redundancy for tracking and triggering. The small and large sectors on a wheel are installed in the projective geometry of EM and EO Monitored Drift Tube (MDT) chambers [5]. The redundancy is achieved by placing 8 detection layers of small-strip Thin Gap

Chambers (sTGC) [6] and 8 detection layers of Micromegas (MM) [7] in each sector. The lever arm between primary trigger chambers are maximized by covering two Micromegas quadruplets with a sTGC quadruplet on each side, in a sTGC - MM - spacer frame - MM - sTGC pattern. With the overlap area between adjacent small and large sectors, the NSW is competent to obtain a full coverage of $|\eta|$ from 1.3 to 2.7 for precision tracking and from 1.3 to 2.5 for trigger acceptance.

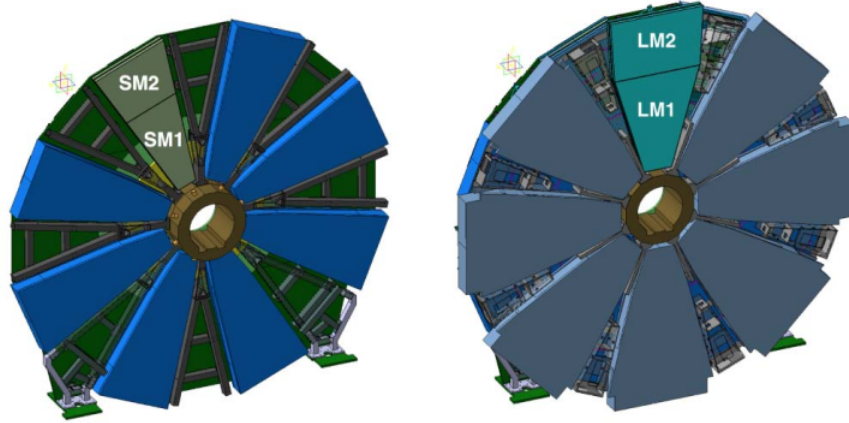


Figure 3.1 – View of the ALTAS NSW with 8 small sectors (left) covered by 8 large sectors(right). Each sector has a trapezoidal shape. (The mechanical structure is also presented. Figure from [4].)

Four different shapes of planar detectors are produced to fill the small and sectors for both sTGC and Micromegas, noted as LM1, LM2, SM1 and SM2. Since one large or small double-wedge sector consists of two quadruplets, there are 16 quadruplets of each planar detector on a wheel and 32 quadruplets in total.

3.2 Micromegas for NSW upgrade

3.2.1 The Micromegas technology

Micromegas [8] is one type of Micro-Pattern Gaseous Detectors (MPGD). The name Micromegas stands for MICRO-Mesh Gaseous Structure, which indicates the main characterization of the detector: the gas volume is asymmetrically divided by metallic mesh into two gaps. The primary ionization occurs in the thicker conversion gap by design, where the electric field supports a uniform drift of free electrons. After passing the mesh, the strength of electric field is higher so that the multiplication of electrons can be concentrated within the thinner amplification gap. This ratio between the electric field in the amplification gap and the one in the conversion gap needs to be tuned to a properly large value in order to achieve an optimal performance of the detector. A schematic view of the electric field in a Micromegas detector is shown in Fig 3.2 (Image taken from [9]). The high electric field field in the amplification gap also provides an advantage of collecting the ion cloud quickly, such that only a small part of the ions pass through the micromesh and escape to the conversion gap [10]. The amount of escaping ions is in fact inversely proportional to the electric field ratio described above. The electron multiplication process is restricted locally to the holes

on the mesh, providing a great potential for high rate capability as well as high spatial resolution.

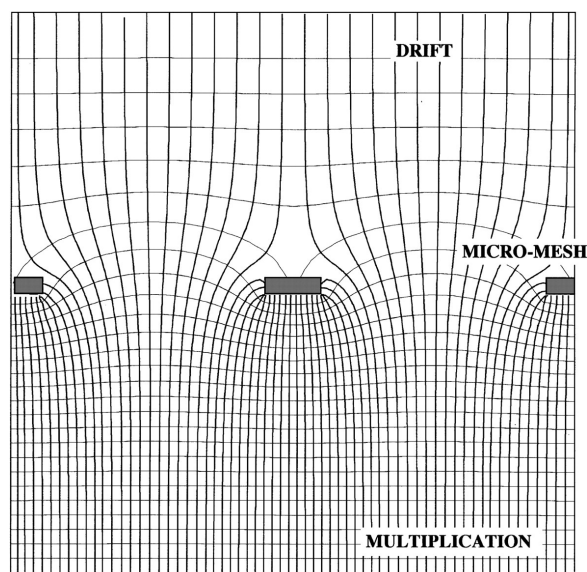


Figure 3.2 – The electric field map in Micromegas detector. (Figure from [9].)

A typical Micromegas detector is composed of the following components:

(1) Anode electrode. Anode strips can be directly placed at the end of gas gap for signal collection in non-resistive Micromegas. For Micromegas with resistive readout design, such as NSW Micromegas modules, the readout strips have to be separated and isolated from the resistive anode. The resistive readout scheme suppresses the global drop of voltage due to sparks at high rate, preventing the occurrence of a long dead time.

(2) Supporting pillars. Pillars are glued on the anode to support the micro-mesh structure during the operation. When high voltage is applied to the detector, the micro-mesh tends to be pressed towards the anode due to the electric force. Hence, the thickness of the amplification gap is determined by the height of supporting pillars, which is in the order of 100 μm . A photo of a supporting pillar used in NSW Micromegas is shown in Fig 3.3.

(3) The micro-mesh. The micro mesh tunes the electric field in the gas volume. Following the non-uniform local electric field, free electrons are induced to drift through the holes on the mesh. Although the mesh is usually fabricated in a dense pattern and shows a low optical transparency, the electric transparency is mainly influenced by the electric field ratio between two gaps and can easily become higher than 95 % for most designs.

(4) The cathode. The cathode, together with the micro mesh, defines the volume of conversion gap.

(5) The working gas. Working gas is filled in the gas volume to facilitate the primary ionization, drift as well as multiplication of electrons with quenching capabilities.

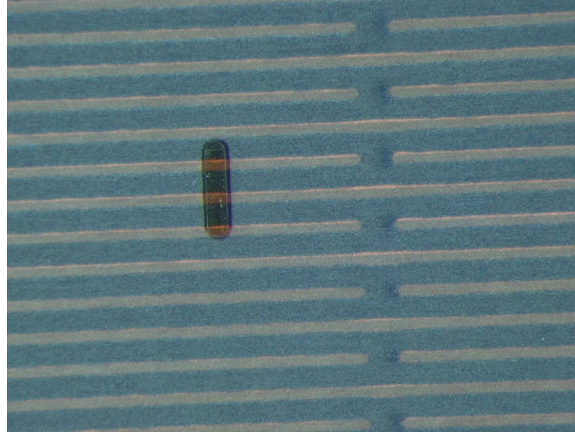


Figure 3.3 – Photo of a pillar on top of the readout panel of NSW Micromegas. The surface of this supporting pillar is about $1.2\text{ mm} \times 200\text{ }\mu\text{m}$. The strips in shallow color are the resistive strips.

3.2.2 An overview on NSW Micromegas

The Micromegas detectors used for NSW upgrade are constructed in parallel plate structures and working in proportional mode. While a negative voltage is applied to the cathode at the beginning of the 5 mm-thick drift gap, the mesh is connected to the ground. The positive voltage on the other side of amplification gap ensures a sufficient amount of electron multiplication. The NSW Micromegas adopted the resistive readout scheme, according to which the positive voltage is applied to a layer of resistive material above the readout PCB panel. Although the resistive strips have the same shape of readout strips and are aligned in the same position and direction to the readout strip, the two types of strips are separated by a glued Kapton foil as insulator to prevent any direct contact. The "ladders" between resistive strips are designed to realise a uniform resistivity on the anode surface. Figure 3.4 shows the appearance of a PCB readout panel, along with a zoomed view of the details on the PCB. The sensitive area of each detecting layer is surrounded by metal frame that defines the thickness of the drift gap. Above the PCB panel, an O-ring is inserted in the groove of the metal frame to ensure the gas tightness in the peripheral region, while in the central region of the panel far away from the O-ring, additional interconnections provide mechanical support of the structure. These interconnections are hollow cylinders appearing in all the gas gaps of the module. During the assembly in the clean room, a long screw is perpendicularly penetrated through the 4 layers of the detector, tightening the detector and fixing the position of each interconnection. The internal components of the Micromegas are shown in Figure 3.5.

At the start of the project, the working gas of NSW Micromegas was chosen to be the mixture of 93 % Ar and 7 % CO₂ at atmospheric pressure. While argon can be easily ionized by the high energy muons, the CO₂ plays the role of a quencher to maintain a stable operation.

The NSW Micromegas quadruplets are produced by different institutes (SM1: INFN, Italy, SM2: Germany, LM1: Saclay, France, LM2: Dubna, Russia - Thessaloniki, Greece - CERN) and then transported to CERN for construction of double-wedge sector as well as final installation. The structure of double-wedge sectors is shown in Figure 3.6.

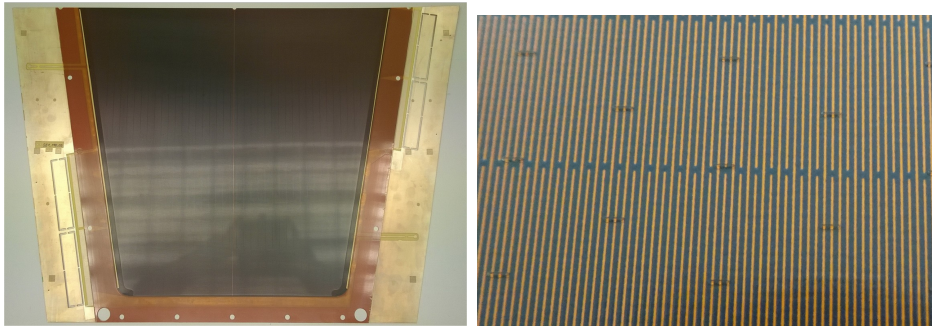


Figure 3.4 – Left: Picture of a trapezoidal PCB1 readout panel. Right: The resistive strips and pillars on the readout PCB. (Figures from [4].)

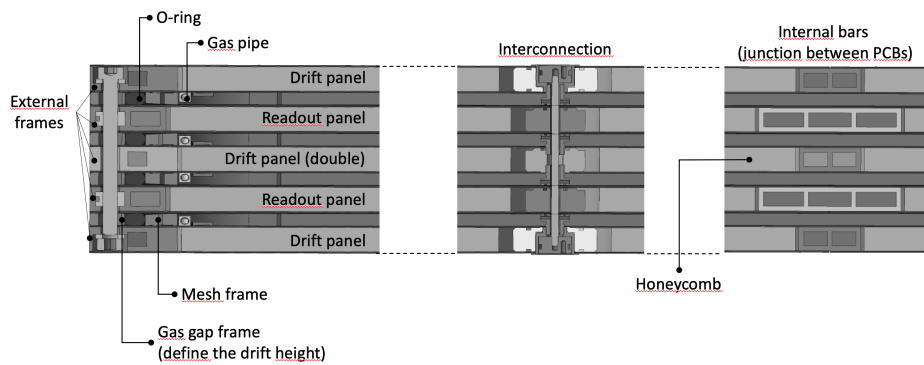


Figure 3.5 – Sectional view of a NSW Micromegas quadruplet, showing the internal components such as O-ring, the interconnections and so on. (Figure from [4].)

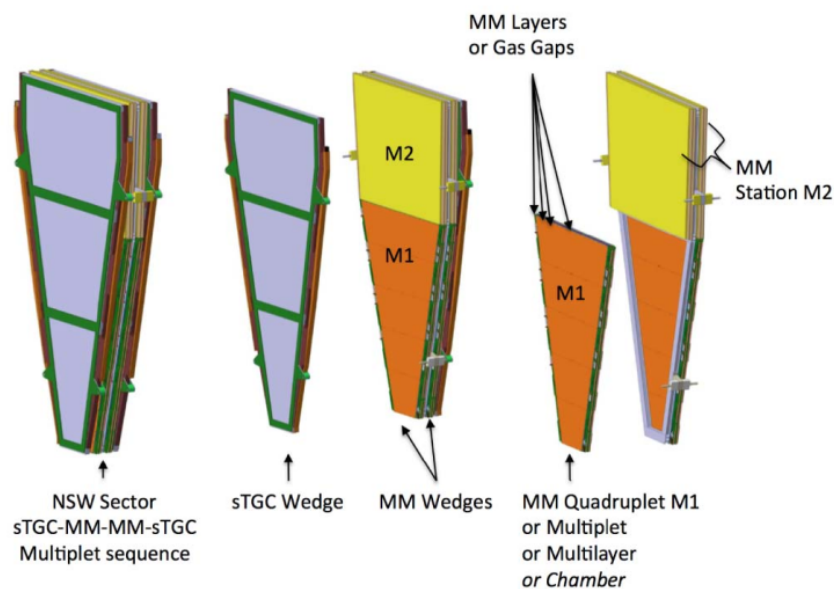


Figure 3.6 – A NSW double-wedge sector. (Figure from [4].)

3.2.3 NSW LM1 Micromegas quadruplets

The 32 LM1 Micromegas quadruplets are all produced at CEA Saclay. The production of the quadruplets comprises of the construction of readout and drift panels,

the resistivity measurement and passivation, planarity scan of the panels, as well as the assembling of the quadruplets [11]. After the production, the LM1 modules are validated both at Saclay and at CERN but via different sets of test procedures to ensure the detectors meet the requirement of NSW. Geometrically, the NSW LM1 Micromegas quadruplet is made of 4 trapezoidal detection layers with each layer containing 5 PCBs: starting from PCB1, the index increases from the small base to the large base. (Figure 3.7, Figure 3.8)(A picture to show the dimensions of LM1 modules, emphasis that LM1 is the largest type of Micromegas in NSW upgrade) The detection layers are arranged in a sandwich pattern. The readout strips on the first two readout panels (noted as Layer 1 and Layer 2) are perpendicular to the height of the panels and are also called eta panels since the measurement of position from those strips can offer the eta information of muon in the ATLAS experiment. Meanwhile, the other two panels (Layer 3 and Layer 4) are composed of readout strips carrying ± 1.5 degrees inclinations to reconstruct the second coordinate while allowing the measurement of the eta coordinates (Figure 3.9). With the 93%Ar + 7%CO₂ gas mixtrue, the drift voltage of NSW LM1 Micromegas quadruplets is -300 V, while the nominal high voltage on the resistive layer is 570 V. One PCB panel is divided into left and right sectors with separate high voltage supply. For each gas gap, the mesh is supported by $128\text{ }\mu\text{m}$ -high Pyralux pillars above the resistive layer [12], as the structure presented in Figure 3.10. To compensate for the insufficient surface resistance of the resistive layer, a technique called "passivation" is applied to the areas that are prone to sparks, including the borders of PCBs and the vicinity of mechanical interconnections. Passivation involves applying glue above the vulnerable sensitive areas of the readout panel to prevent the amplification of electrons from ever happening. Despite the loss of acceptance, this technique is proved to be an indispensable step in the production to ensure a good HV performance of the module.

In order to achieve the $100\text{ }\mu\text{m}$ spatial resolution required by NSW, nominally one PCB panel on a single layer of LM1 Micromegas should have 1024 readout strips, each $300\text{ }\mu\text{m}$ in width and $150\text{ }\mu\text{m}$ gap between neighbor strips. As for the signal readout, the 1024 strips are divided from the middle so that first 512 strips and the other 512 strips are readout from two sides of the PCB. However, contrary to traditional method, no soldered connector is used for LM1 modules. Instead, the electronics is always connected using Zebra elastomeric connector [13] on front-end electronics board, which increases the flexibility of connection. The metal brushes on the Zebra connectors are designed to stay contact with the footprints of the readout strips, enabling the transmission of electrical signals. The mounting position of the front-end electronics board is fixed by precision pins and the compression bars installed on the module.

3.3 Cosmic ray test bench at CEA Saclay

The cosmic ray test bench at Saclay occupies an area of 20 m^2 and is able to scan the surface of NSW LM1 Micromegas modules by sliding a supporting tray (Figure 3.11). On the bench, there are 3 external trackers called Mcube detectors that are Micromegas telescopes initially designed for muon tomography [14]. The layers of external trackers are denoted as HH, HB and BH from top to bottom. Each reference Mcube detector covers an area of 1 m^2 and consists of four $0.5\times 0.5\text{ m}^2$ bulk-

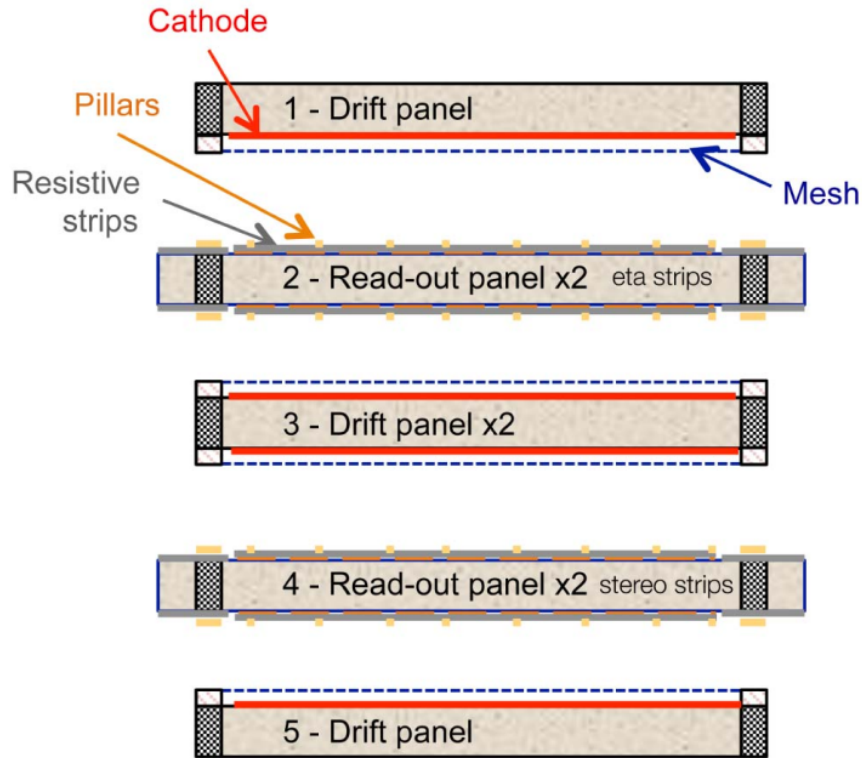


Figure 3.7 – Cross section view of NSW Micromegas illustrating the structure of 4 detecting layers. (Figure from [4].)

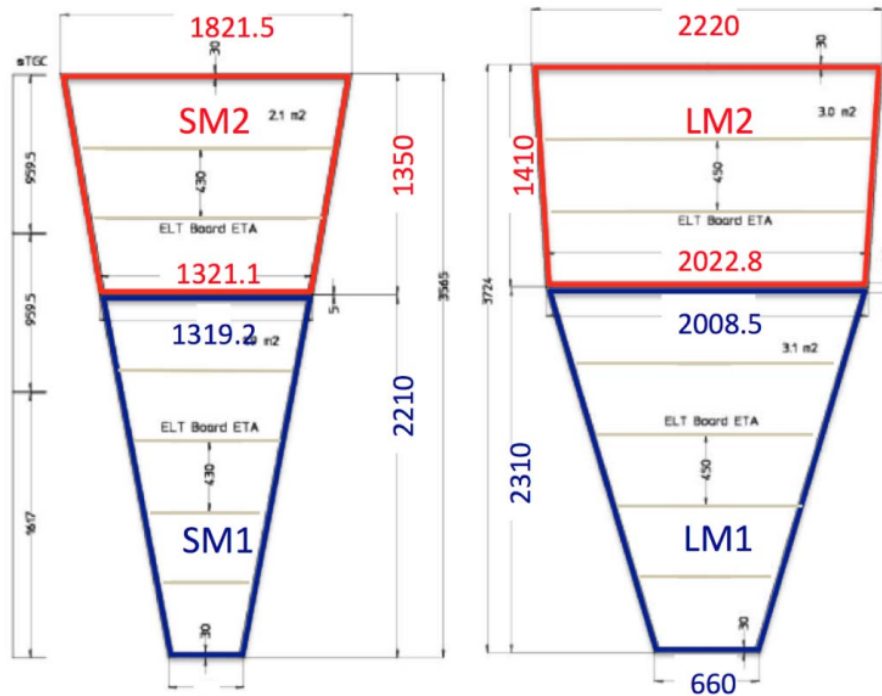


Figure 3.8 – Dimensions of different NSW Micromegas modules including non-sensistive area. LM1 quadruplet is the largest one among four types of detectors. (Figure from [4].)

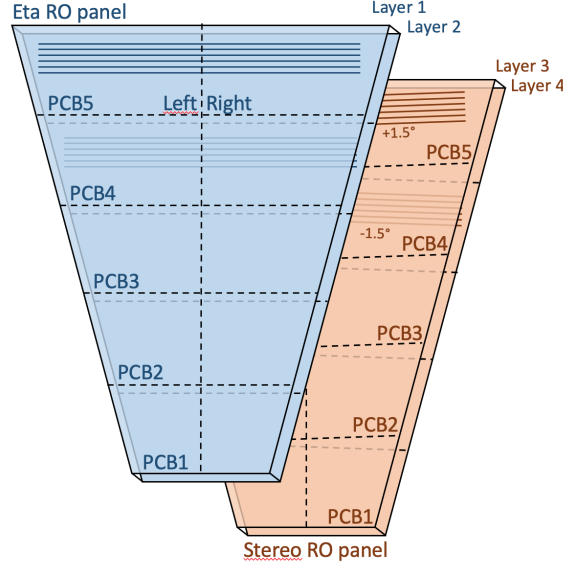


Figure 3.9 – Directions of readout strips on four readout planes in a quadruplet. (Figure from [11].)

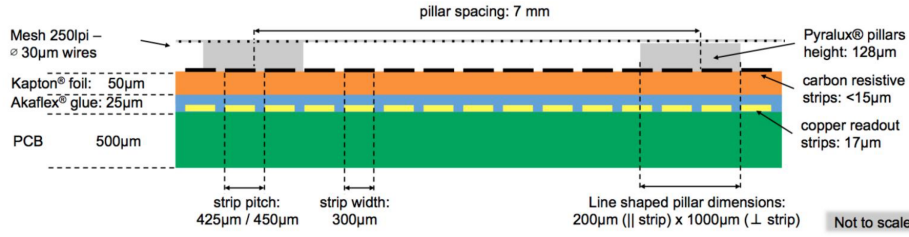


Figure 3.10 – Structure of the readout panel together with the mesh. (Figure from [11].)

1280 Micromegas (Multigen2D) with a drift gap of 1 cm [7]. The Mcube system has a
 1281 2-dimensional readout. Therefore, a track of the cosmic muon, of which the average
 1282 energy is 4 GeV [15], can be reconstructed by the three layers together. In order to save
 1283 electronics channels, the 1024 strips of one Multigen2D detector are multiplexed into
 1284 only 61 channels [16]. These external trackers work with 95 % Argon+5 % Isobutane
 1285 mixture provided by an independent gas system.

1286 Given that the development for LM1 cosmic ray characterization took place before
 1287 the readiness of VMM electronics [17] which is currently used on NSW, the cosmic
 1288 test bench at Saclay chose the available Deadtimeless Readout Electronics ASIC for
 1289 Micromegas [18] as the readout electronics. DREAM, developed at CEA for CLAS12
 1290 experiment [19], is able to cope with Micromegas detectors with high strip capacitance
 1291 (200 pF for 2 strips on PCB5 in a measurement) while keeping the signal to noise ratio
 1292 larger than 10. This setup of electronics for cosmic test bench has been validated in
 1293 both CLAS12 experiment and muon tomography [20].

1294 The signal readout of both the NSW modules and Mcube are achieved via a data
 1295 acquisition system based on DREAM electronics. A Front End Unit (FEU) board
 1296 consists of 8 DREAM ASIC chips and is placed inside the crate. Since one DREAM
 1297 corresponds to 64 channels and there are 8 FEU for the readout of LM1 module, the
 1298 test bench is then competent to take data of 4096 strips at the same time. However,
 1299 the LM1 module contains about 20,000 strips in total and thus around 4,000 strips

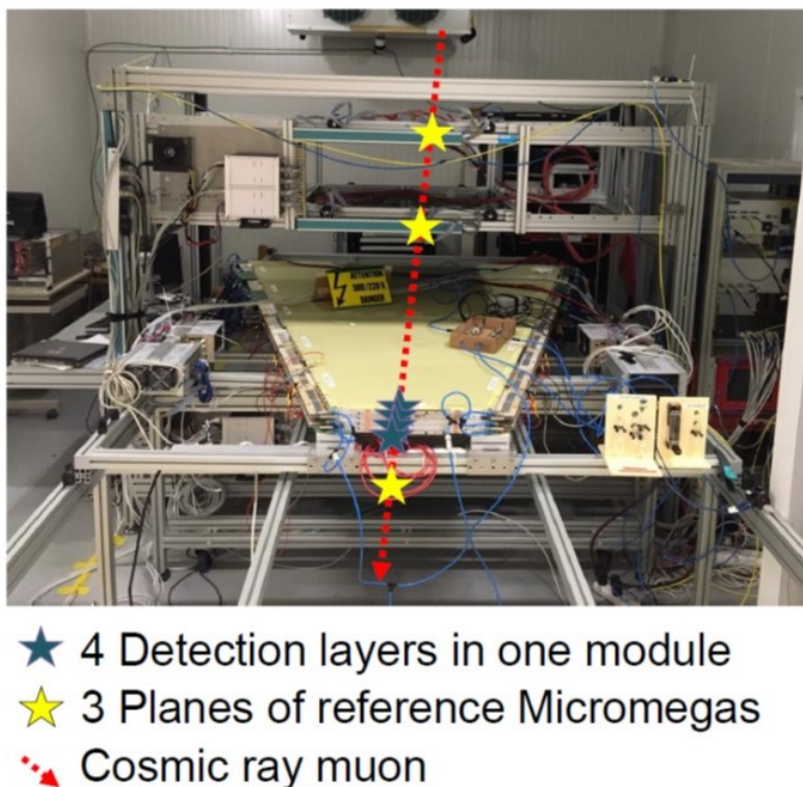


Figure 3.11 – A photo showing the setup of Saclay cosmic bench. 3 layers of Mcube system is included together with the LM1 module ready for test.

1300 per PCB. Connecting DREAMs directly to the zebra line on LM1 means the cosmic
 1301 bench is able to scan only one PCB of LM1 module at a time, which is not only time
 1302 consuming but also fail to profit of the full sensitive area of the external trackers.
 1303 As a solution to this scheme, electronics channels are adapted to zebra lines via a
 1304 multiplexing strategy. The idea is to use the adapter to associate one channels to two
 1305 strips. Considering the low cosmic ray multiplicity of LM1 module (only 3-4 strips for
 1306 a muon event), it is possible to design a multiplexing circuit with a special pattern of
 1307 connection to save electronics channels. A multiplexing factor of 2 has been decided
 1308 as the final option, as it maintains a good signal to noise ratio while fully making use
 1309 of the coverage of Mcube system.

1310 3.4 Test protocol

1311 Every NSW LM1 quadruplet is installed on the cosmic test bench after assembly
 1312 and quality checks described in [11]. The module can be sent to CERN for NSW
 1313 installation only if it has passed through the validation procedures on the cosmic
 1314 bench.

1315 The humidity is among the many factors (dust, panel planarity and so on) that
 1316 affect the high voltage stability of Micromegas. If the humidity exceeds a certain limit
 1317 in the gas gap, high voltage instabilities will show up upon the Micromegas modules.
 1318 It has been observed that water may remain inside the PCB panel after washing and

cleaning, which can slowly evaporate into the gas gap, causing high leakage current or even sparks during operation of the detector. To speed up the drying-up, while flushing the working gas, modules are covered by plastic foils and heated from the bottom to get rid of humidity. Once the relative humidity (RH) reached as low as 20 % at 45 °C, the heaters are turn off. The module will cool down to room temperature around 20 °C after 12 hours. By monitoring the relative humidity of going-out gas, one is able to compare the humidity of the module before and after the heating. The RH is reduced from 37 % to 8 % through a 3-day heating process. The monitoring of this process is shown in Figure 3.12.

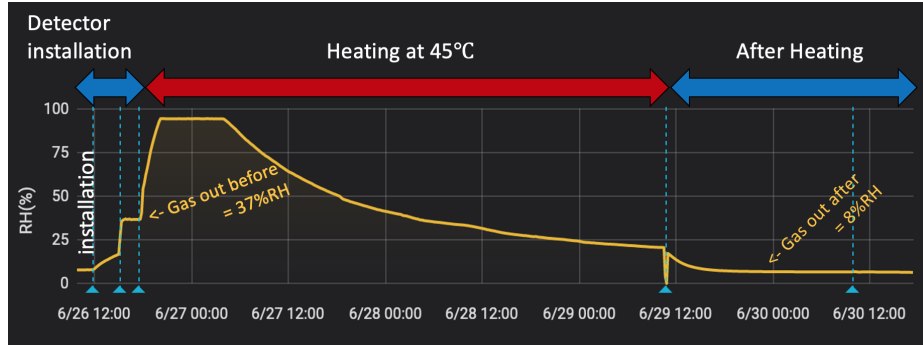


Figure 3.12 – The reduction of humidity by heating the LM1 module. The orange curve stands to the relative humidity of going-out gas, which represents the humidity inside the detector. Although the value of relative humidity depends on the temperature, in the end we still observed a great drop of the humidity after the module cooled down back to room temperature.

Once the module is verified to be dry enough, specifically when the relative humidity (RH) of going-out gas is lower than 10 %, the high voltage will be applied to the resistive strips. While the drift voltage is always 300 V for the 4 layers and can be easily achieved, the increase of the amplification voltage to 570 V requires a slow ramp up process. It turns out that the best results can be achieved only via a relatively slow and careful ramp-up of high voltage above all the 40 HV sectors. On the cosmic bench, the high voltage and corresponding current consumption are monitored by an automatic HV control program. To avoid tripping, a sector can rise up to the next HV step only when its current is low and stable after a certain time interval. In case the current fluctuates too much, the corresponding HV will be automatically reduced to a lower level for protection.

One more step before the actual data taking is to check the connection of Mux2 cards. As the Mux2 cards are always pressed to the module by metal compression bars, a careful cleaning is necessary before the mechanical installation of Mux2 cards on the module. This cleaning procedure makes sure that any dust or remaining glue from the tapes used in the production are removed from both the footprints on the module and the zebra lines on Mux2 cards. The first check of connection is by taking pedestal RMS of each electronics channel. A non-uniform pedestal may suggest some missing connection since the noise of a electronics channel is proportional to the capacitance. The known bad connections can then be improved by re-plugging of Mux2 card or adjusting the connectors of cables. The second check comes with a short data taking after the module is slid to the Mcube window. A short data taking is necessary because the multiplexing hides the missing connections in the noise run

of electronics pedestal. If the number of event counts in a channel is half of the others, it is very likely that the electronics channel in concern has a bad connection to one of its readout strips. The bad connections will be restored as good as possible before the HV scan and the long runs.

Due to the limited length of Mcube and the number of available electronics channels, the data taking has to be repeated for three times in order to scan all the strips in LM1 module. The test starts from PCB45, then moves to PCB23 and ends at PCB1 (Figure 3.13), which usually takes one week. For PCB3, 4 and 5, the surface beyond the Mcube acceptance will not be scanned because of the tight production schedule.

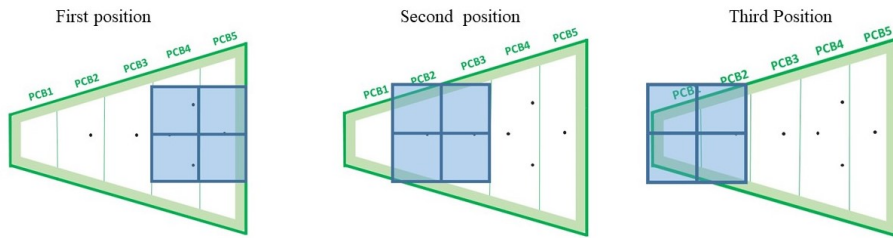


Figure 3.13 – The scheme of data taking for 5 PCBs. In order to scan all the 5 PCBs, the module needs to be at least posed at 3 positions. The blue square refers to the reference Mcubes, which are always fixed to the frame of cosmic bench. The offset between Mcube and the LM1 module along X direction can be obtained by data analysis after data taking and will be corrected before evaluating the performance of the LM1 module. (Figure from [11].)

3.5 Cosmic bench data analysis

The data analysis code of LM1 cosmic muon characterization [21] has been adapted from muon tomography studies and has been evolving to the last version of the analysis that was finished in 2020. It is divided into 4 steps that will be further illustrated:

1. Subtraction of pedestal: The amplitudes of electronics channels go through the common noise suppression, as well as the subtraction of mean amplitude value in the absence of a cosmic muon signal (the pedestal is subtracted to have equalisation). The suppression of a remaining common mode noise follows the pedestal subtraction: While the 50 samples are collected, the baseline amplitudes of all strips of a given Front-End-Unit (FEU) electronics are interpreted to be the common mode and are moved coherently.

2. Hit finding: After the subtraction of the pedestal and the suppression of the common mode noise, only the signals that have amplitudes 5 times than the standard deviation of channel pedestal are considered to be "hits".

3. Clustering: The neighboring hits gathered to form clusters. Since signals are readout in a multiplexing pattern, the conversion from electronics channels to strips requires a dedicated demultiplexing algorithm.

4. Track to cluster matching: The cluster matching algorithm associates muon tracks to the reconstructed clusters on each layer by an interpolation. Requirements

on the quality of the track are applied. A count of detected or undetected events gives the efficiency of each strip.

The above mentioned stages are all configured by a general configuration file. This configuration file is uniquely linked to the analysis of a specific data taking, where the geometric setup of the cosmic bench, the path to the pedestal as well as signal files, the cluster reconstruction algorithm, the arrangement of readout electronics and so on, are defined for the programs at each stage.

3.5.1 Subtraction of pedestal

The pedestals are calculated as the mean values of amplitudes without response to cosmic muon. In the pedestal taking, the amplitudes are recorded according to random triggers. The pedestals for all channels in one FEU are shown in Figure 3.14. Here Figure 3.15 shows an event display of 8 DREAMs in a FEU when muon passes through the LM1 module. 32 samples are collected for every event, with a sampling period of 60 ns. While pedestal subtraction is applied to the signals, the baseline of strips are also moving coherently for a given FEU. Their remaining common mode will also be suppressed. By the end of this stage, the information of pedestals is recorded in a text file while the processed signals are stored into root files.

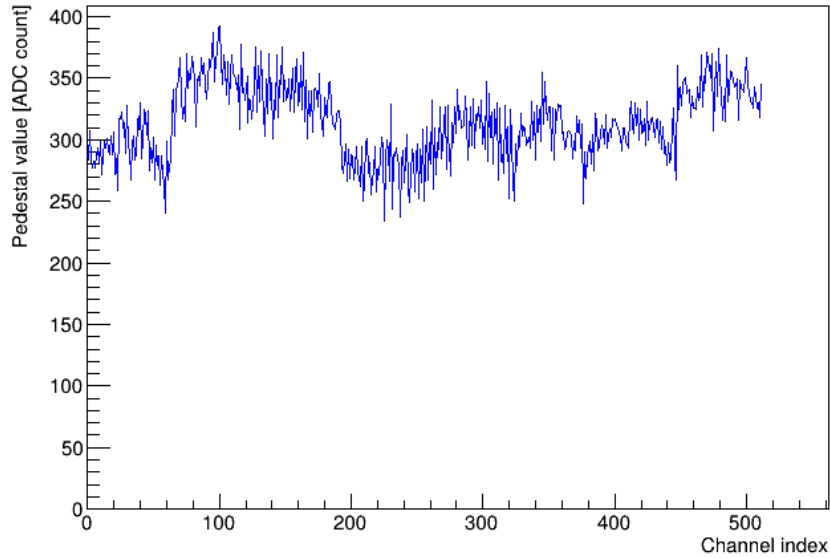


Figure 3.14 – The pedestal values of all channels on a FEU board during a pedestal taking.

3.5.2 Hit finding

The standard deviations of pedestals are obtained along with their mean amplitudes and are usually about 10-15 ADC counts. The hits are then selected according to the 5 times pedestal RMS criteria. Therefore, the amplitude of hits is supposed to be larger than 50-75 ADC units in most cases. The selected hits are saved as another root file, ready for the process of clustering. Considering that the connectors to the zebra line on the Mux2 card can be worn out over time, replacement of the zebra line connectors is necessary, which involves unplugging and re-plugging the readout cables

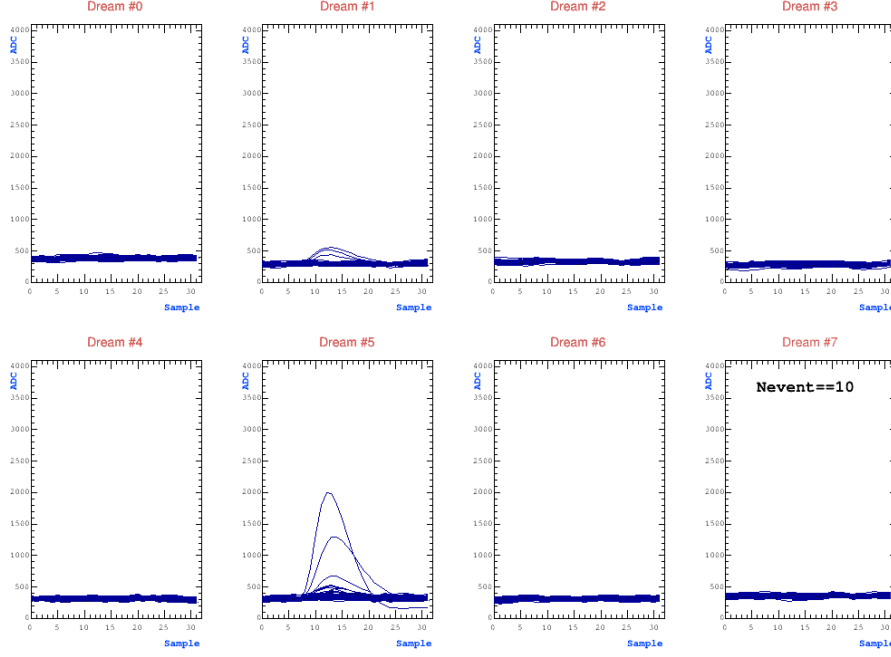


Figure 3.15 – An event display extracted from the raw data for 8 DREAMs on a FEU board. Signals from two DREAMs are observed, since these two DREAMs are fixed in the same X and Y locations of the module but on different layers. Pedestals around 256 are visible, as no correction is applied.

to the zebra line connectors. However, the readout cable can be inserted to the socket of the zebra line connector either way round, bringing the possibility of connecting them in an inverted way by mistake. The hit finding algorithm therefore includes the flexibility of handling inverted zebra line connection, such that the data analysis can be carried out regardless of this type of mistake. When an inverted zebra line connection happens, as the example presented in Fig 3.16, the mistake can be easily identified by comparing the physical position of the hit on the LM1 module with the cosmic muon predicted by the external trackers. Once identified, the mistake can be fixed by the hit finding algorithm as long as the identified inverted connection is manually labelled out in the corresponding configuration file.

3.5.3 Clustering

Clustering is performed separately on each layer, since the amplification is independent in the four gas gaps. The objective of clustering is to convert the "hit map" (Figure 3.17(a)) of electronics created by "hit finding", to a "cluster map" (Figure 3.17(b)) that records the positions of reconstructed clusters per detecting layer.

Due to the fact that one electronics channel is linked to 2 strips via the Mux2 card in the way indicated in Fig 3.18, referred to as "multiplexing", it is impossible to determine the position of a single hit. The demultiplexing is thus necessary in order to perform the reconstruction of clusters. To begin with, the hit with highest amplitude is chosen and matched with other hits according to a channel-to-strips map created in configuration file for this run. For every matching, both combinations are considered, assuming the hit comes from either of the two strips associated to each

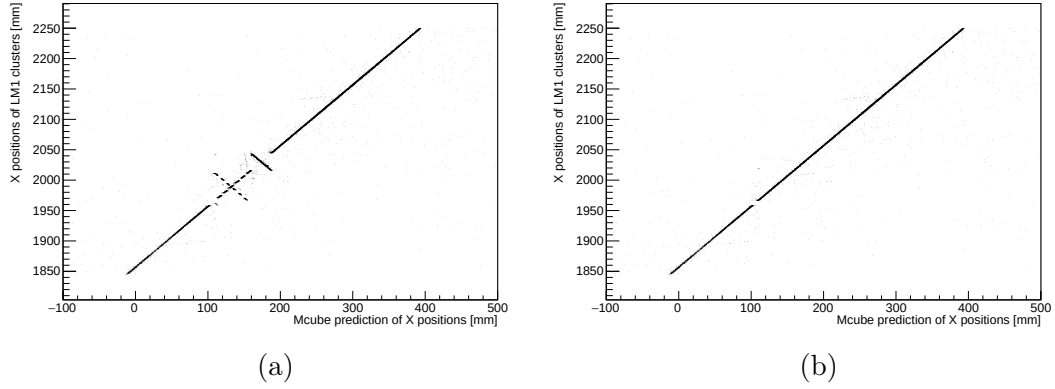


Figure 3.16 – The correlation of X position between Mcube prediction and LM1 clusters for M35 PCB5 layer2. The offset of X position is not set to zero in this plot. (a) The correlation in the presence of an inverted cable connector. (b) The correlation after taking into account the inversion of cable connector in the configuration of data analysis.

channel. Two strips are considered to be neighbors if they are within the interval of 2 strips, or in another word, 0.9 mm from each other. Once the neighbors are found, the responding strips of both hits are fixed and the two hits are merged into a cluster. The demultiplexing algorithm then loops over the rest of the hits, trying to form a cluster as large as possible, provided that one possible strip of the next hit turns out to be the neighbor of a determined strip in the growing cluster. The clustering of this cluster ends when no more neighbor strip can be found. After that, the clustering of another cluster begins in the same method, gathering remaining clusters into a new cluster. The clustering finishes until all hits are uniquely assigned to a cluster. In case that a hit can never find any neighbor for both of its associated strips, two unity size clusters are generated in the position of each strip. Given that time window of the reference Mcubes is $32 \text{ samples} \times 40 \text{ ns}$, the small flux of cosmic muon ensures a low rate of signal for the reconstruction of a single cosmic ray muon most of the time. The trigger rate of Mcubes is below 20 Hz due to its efficiency, the acceptance angle and the high threshold on the electronics side to work in auto trigger. Even if sometimes more than one correlated particles are produced in the cosmic ray shower (of which the size is in the scale of meter), the muon tracks may become difficult to reconstruct in the Mcubes, therefore making the event being rejected due to the low quality of the reconstructed track. As a result, mainly the events with one cosmic muon are used to study the performance of the LM1 Micromegas, so that the duplication of multiplexing has a negligible effect to cluster and will not bring any major bias due to the mixture of multiple muon tracks. It has been checked by comparing to the data taking without multiplexing that the deficit in efficiency due to multiplexing is small, as long as the detectors work in the nominal HV.

The charge of a cluster is defined to be the sum over the amplitudes of all the constituent hits. The position of the cluster is the average weighted strip number of its hits where the strip number is weighted by the amplitude. While a hit must come from a physical strip, the position of a cluster refers to the centroid position except cluster of unity size. By the end of this step, the properties of clusters are saved in a format easy for the analysis of track matching.

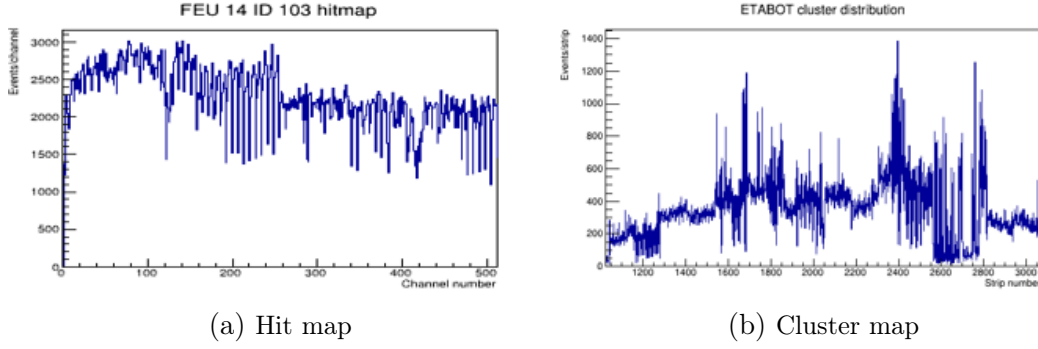


Figure 3.17 – Examples of (a) The "hit map" of electronics and (b) The "cluster map" of one detecting layer of the LM1 module during the cosmic bench characterization.

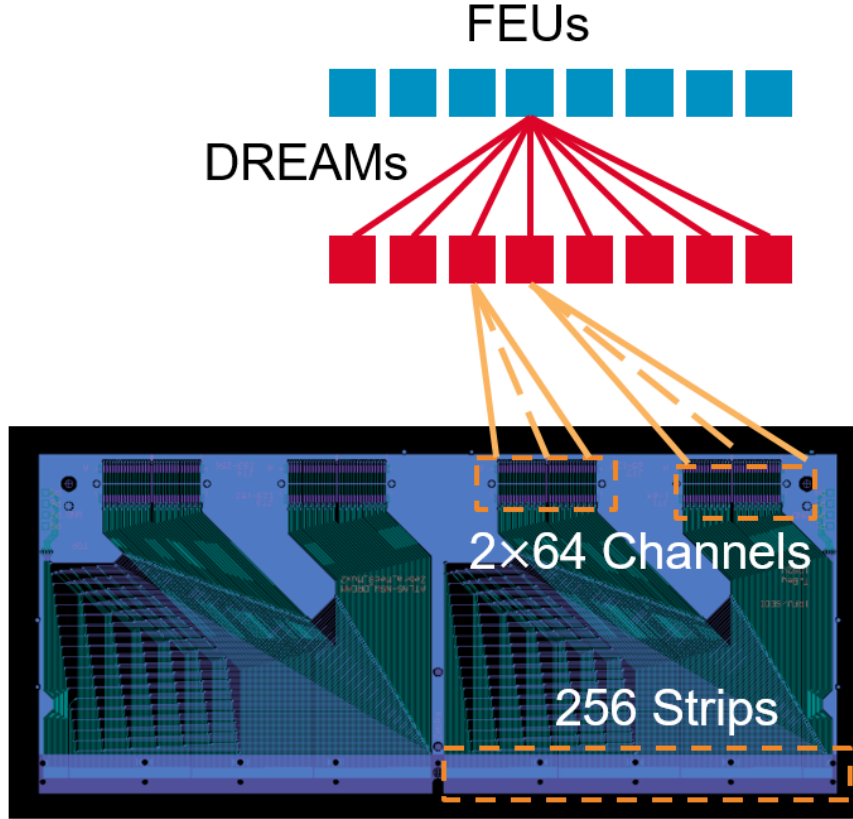


Figure 3.18 – Schematic view of the multiplexing readout. Each 128 electronics channels of 2 adjacent DREAM ASIC chips are associated to 256 readout strips of the LM1 module through the Mux2 card. The channel-to-strip projection is not always continuous, such that the granularity of the strips can be retained in cosmic ray tests.

3.5.4 Track to cluster matching

A Mcube track matches a cluster if the residual along X direction is smaller than the residual cut, which is set to 10 mm. However, selection is applied to both Mcube tracks and clusters before matching of a cosmic muon event.

Based on the signal readout of 3 Mcubes, the cosmic muon rays are reconstructed by a tomography program of which the essential idea is similar to the algorithm described in [20]. Considering the geometry of the cosmic bench, the selection of

good tracks requires muon rays with χ_X^2 less than 2.5, χ_Y^2 less than 10 and incident angle no more than 0.4 rad. Here X and Y are orthogonal axes of Mcubes. The cut in the Y direction is looser than that in the X direction because of worse resolution of the Mcubes along the Y direction. The Y is also the direction of readout strips stretching on the eta layers of LM1 module on cosmic bench. In some rare cases, there might be more than one possible muon rays that meet the requirement of good track, hence the selection of best track is introduced to keep the muon ray that has the minimum χ_X^2 from those match the most clusters on 4 layers of LM1 module. Matching is only performed between the best Mcube track and reconstructed clusters. As for a zero match event, the best track is simply the good ray with minimal χ_X^2 .

Once the Mcube track is determined, it is the time to select the clusters. For each layer, only the cluster with minimum residual in X with respect to the position predicted by Mcube track is chosen to be the response of the module on that layer. When applying the residual cut, this cluster with best reconstructed position could be removed if it is too far away from one predicted X position, yielding an inefficient event for the layer. A typical cut flow for cosmic muon events is listed in Table 3.1. The number of events used for matching is reduced by a factor of ~ 10 from the initial data after the selection criteria.

Cut	Entries			
Mcube events	499949			
At least 1 track	334791			
At least 1 good track ($\chi_X^2 < 2.5$, $\chi_Y^2 < 10.0$ and incident angle < 0.4 rad)	54583			
	Layer 1	Layer 2	Layer 3	Layer 4
Geometric cut	44447	44142	43319	45297
Residual cut (10 mm)	38831	38186	36818	39884

Table 3.1 – Cut flow of cosmic muon analysis for Module 13 PCB4 and PCB5

The distribution of residuals that reveals the difference between the X position of best cluster and the prediction of Mcube track, is shown in Figure 3.19 left. Due to a translation along X direction, this distribution is not necessarily centered at 0 before offsetting the module. The offset has to be found for each data taking in order to make sure the matching is based on proper muon rays and clusters. In order to find the center of the residual distribution for the adjustment of module offset, the cluster residual is always fitted by a double Gaussian function per PCB per layer as a proxy of effect of the LM1 spatial resolution convoluted over muon angles, smeared by the imperfect alignment as well as the resolution of the external Mcube trackers. The fraction of the wider Gaussian in this fit is artificially restricted to be within 5 %, yielding a corresponding standard deviation around 3 mm. Although neither standard deviation of the double Gaussian fit can be interpreted to be the intrinsic spatial resolution of the detector, that of the wider component offers a reliable reference for setting the residual cut. After studying the dependence of efficiency on the choice of residual cut, it was decided to set the residual cut to 10 mm for all the measurements. This universal residual cut is from 3 to 5 times larger than the standard deviation of the wider Gaussian, thus ensures a good coverage of the real detector responses to the cosmic muon while excluding the clusters that are too far away from their predictions. From Figure 3.20, it can be seen that the cut of 10 mm locates on the plateau of efficiency, which means changing the criterion of residual cut by a reasonable amount

will hardly influence the measurement of efficiency. The residual cut should not be set to the turning point of the curve without any margin, otherwise the efficiency would be too sensitive to the performance of each individual PCB as well as the subtly of Mcube alignment during data taking. As a cross-check in Figure 3.19 right, the correlation between the best cluster position and the Mcube tracker prediction along X follows an expected linear relationship. The sparse random dots spread beyond the linear correlation are considered to be the noise and are therefore effectively excluded from the efficient events, thanks to the residual cut.

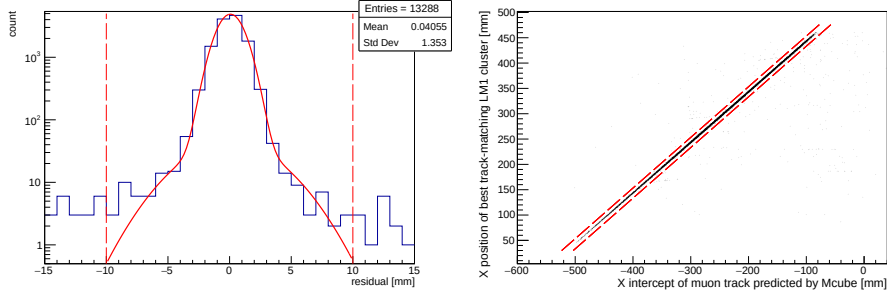


Figure 3.19 – Left: The cluster residual distribution of Module 11 PCB1 Layer 1 before applying the residual cut. The red curve represents the double Gaussian fit. Right: The corresponding correlation of X positions between best reconstructed clusters and their predictions by the external trackers. The dashed lines in both plots indicates the 10 mm residual cut.

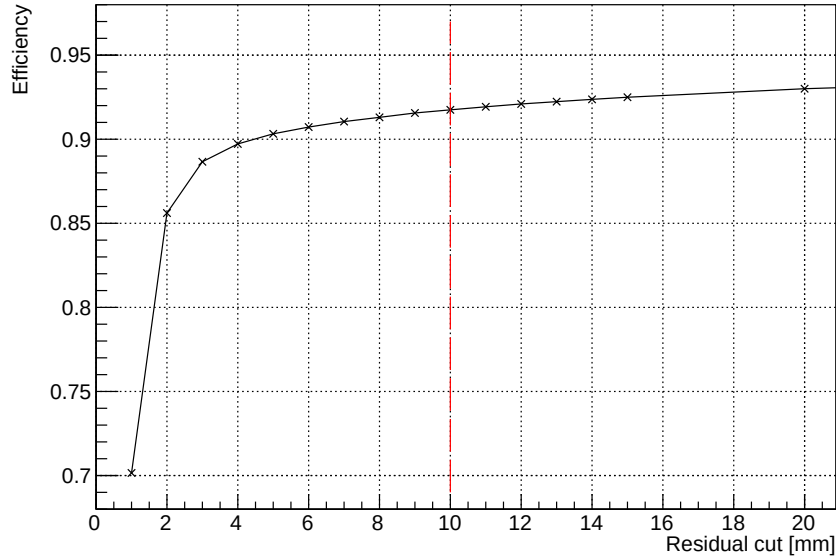


Figure 3.20 – The efficiency of LM1 module with different criteria of residual cut. The 10 mm residual cut is indicated by the red dashed line, corresponding to an efficiency of about 92%.

3.6 Performance of modules

3.6.1 Efficiency

Each matching cluster that passes the residual cut will be assigned a position in the module. Besides the layer and the X position is known from the clustering, the Y position of the cluster is also given by the Mcube track thanks to the 2D readout of Mcube detectors. Although the resolution along Y is not as good as that in X, this method is still powerful enough to perform the 2D characterization for LM1 modules. The efficiency is studied at the base of every HV sector and summarized per PCB per layer. The overall efficiency for each layer or each PCB is the average weighted efficiency over the corresponding HV sectors. The weight used for the average is the inverse of the squared error of efficiency in the HV sector, in order to ensure the sensitive area of each sector is properly taken into account in the average. The choice of weight is based on the assumption of a constant efficiency, following which the number of accepted events in a given amount of time for each sector is proportional to its area, so that effectively the sensitive area is proportional to the inverse of the squared relative statistical error. i.e. a HV sector with larger sensitive area reconstructs more cosmic muon events (which also means a smaller error of efficiency) and therefore should have a greater contribution to the average of efficiency. For this reason, the data taking is supposed to have more or less same amount of cosmic muon events recorded by Mcubes for different coverage of a LM1 module, otherwise the weight of PCBs that accept more events will be overestimated because of lower statistical fluctuation. Here the errors of efficiency are obtained along with the efficiency itself via the TEfficiency class of Root [22] on 1σ confidence level. Despite the asymmetric errors of Poisson statistics, the up and down errors are symmetrized to a single number, so as to simplify the analysis workflow. The symmetrization of up and down errors is justified as their absolute values are generally very close to each other for a given HV sector.

However, considering the fact that the acceptance of Mcube detectors is smaller than the Y-width of PCB 3,4 and 5, the efficiency on the borders of these PCBs cannot be studied and will not be included in the final results. The count of cluster matching on the rest of the area gives a so-called raw efficiency, which does not take into account the effect of bad connections or the passivation on edges. Even though the connection is always carefully checked via inspecting "hit map" before data taking, remaining glues, degrading of Mux2 cards or even broken zebra lines may still lead to a loss of efficiency that is hidden behind multiplexing. In order to correct this underestimation of efficiency caused by bad connection of electronics, the strips with efficiency less than 10% in a data taking will be excluded from the calculation of efficiency. Meanwhile, the edge passivation can be observed on PCB1, 2 and 3 by Mcubes on the cosmic bench, which will produce a large inefficient area if we look at the 2D distribution of efficiency. This impact of edge passivation is corrected by introducing a global margin factor to the border of the module instead of simply using the geometric edges. As an example, the efficiency of a module is summarized to each layer in Tab 3.2 and to each PCB in Tab 3.3. The efficiency significantly goes up for PCB1 after excluding the effect of edge passivation. Yet the difference is smaller for PCB 2 and 3 since their proportion of passivated area is lower compared to that of PCB1.

Layer	Raw Efficiency	Efficiency excluding bad connections	Efficiency excluding passivated areas
L1/Eta top	0.86 ± 0.11	0.88 ± 0.11	0.89 ± 0.03
L2/Eta bottom	0.85 ± 0.06	0.87 ± 0.05	0.88 ± 0.03
L3/Stereo top	0.84 ± 0.06	0.86 ± 0.05	0.88 ± 0.03
L4/Stereo bottom	0.87 ± 0.08	0.88 ± 0.07	0.89 ± 0.03

Table 3.2 – Mean efficiencies of the four layers of Module 13.

Layer	Raw Efficiency	Efficiency excluding bad connections	Efficiency excluding passivated areas
PCB1	0.72 ± 0.06	0.74 ± 0.5	0.91 ± 0.04
PCB2	0.82 ± 0.03	0.86 ± 0.01	0.92 ± 0.03
PCB3	0.86 ± 0.03	0.91 ± 0.02	0.91 ± 0.02
PCB4	0.87 ± 0.05	0.88 ± 0.03	0.88 ± 0.03
PCB5	0.85 ± 0.04	0.88 ± 0.01	0.88 ± 0.01

Table 3.3 – Mean efficiencies per PCB for Module 13.

When looking at the final efficiency values after two corrections, the module achieved not only the relative high efficiency per individual layer (90 %), but also an overall good homogeneity for all the 4 layers (<4 %). An explicit view of the efficiency comes from the efficiency map (Figure 3.21), which is the combined 2D distribution of efficiency from the data taking at 3 locations on cosmic bench. The blue area near the small base is due to the edge passivation while some circular inefficient spots come from the interconnections that support the mechanical structure of the quadruplet. Figure 3.22 shows the efficiency per layer of every module validated on the cosmic bench. Generally speaking, they were all fulfilling the requirements during the cosmic muon characterization period at Saclay.

3.6.2 Charge homogeneity

In addition to the efficiency, the charge of matching clusters are also studied to validate the properties of LM1 modules in terms of gain. The charge distribution of different HV sectors are fitted with Landau function and the Most Probable Value (MPV) obtained from this fit is chosen to represent the charge in ADC unit for the sector (Figure 3.24). Although the actual number of gain relies on the transfer function of electronics which heavily depends on the connection scheme, especially the capacitance of the connected readout strip, it is still possible to estimate the level of gain by some simulations.

With the setup of electronics on the cosmic bench, the total input capacitance including capacitance of the strip, the cable, the cross-talk as well as the parasitic capacitance of the Mux2 card, is about 280 pF on PCB1. As longer strips bring larger capacitance and stronger cross-talk, the total input capacitance increases to 520 pF

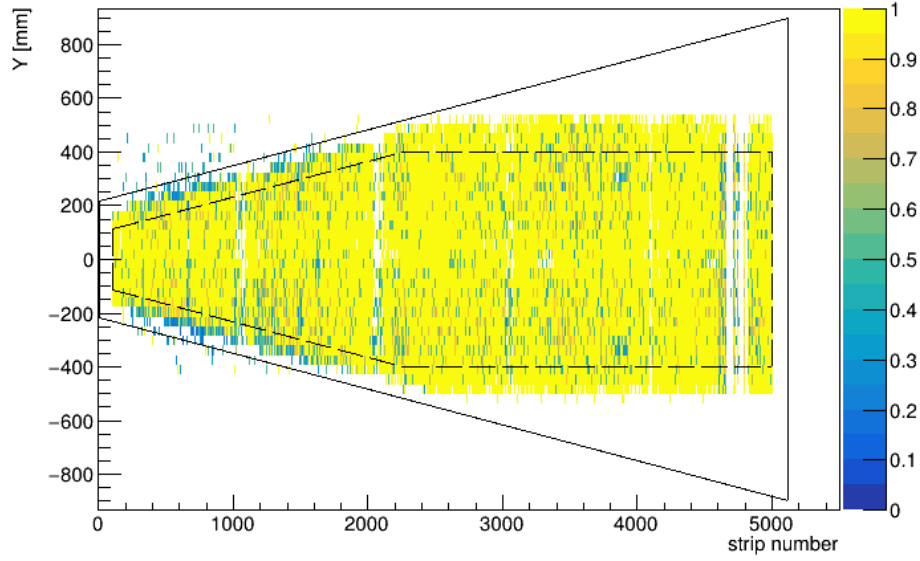


Figure 3.21 – 2D efficiency map of Module 11 layer 2, where the bins in yellow stand for efficiency higher than 95 %. The solid lines outline the nominal sensitive area of this layer. But the area taken into account for efficiency calculation is indicated by dashed lines, which excludes both the surface inactivated by edge passivation and the area beyond Mcube acceptance. Low efficiency caused by passivation can be seen near the small base of the trapezoid between the solid lines and dashed lines.

on PCB5. At the average strip length, the corresponding capacitance is inferred to be about 400 pF. By contrast, the transfer function of DREAM electronics decreases with the input capacitance. At 280 pF, 400 pF and 520 pF, the values of transfer function according to the simulation are 8.4, 6.7 and 5.8, in the unit of bin/fC.

The estimation of the gain uses the value of transfer function at average strip length, which is 6.7 bin/fC. As simulated by Heed [23] and presented in Figure 3.23, the primary ionization occurs in a 5 mm gas gap gives 47.4 primary electrons (N_{PI}) in average. The momentum and the angular distributions of cosmic ray is accounted for by adopting a phenomenological model [24]. By multiplying the amplitude factor to the MPV value of charge distribution in ADC unit and then dividing the number by the scale of primary ionization, one gets the absolute gain about 7,000. For illustration, the explicit calculation of this estimation is done in Equation 3.1.

$$Gain = \frac{ADC}{e \times f_{Mux2}^{Response} \times N_{PI}} = \frac{350}{1.6 \times 10^{-19} \text{ C} \times 6.7 \text{ fC}^{-1} \times 47.4} \approx 6888 \quad (3.1)$$

It should be noted that the distribution of charge as well as the Landau fit need to be studied individually for each HV sector, since even a slightly different voltage in another sector is going to change the gain exponentially, which in the end blurs the total distribution because of two displaced Landau peaks. The overall charge for each layer or each PCB still comes from the averaged weighted charge of constitutional HV sectors. But this time the weight is the inversed square error of MPV of the fit, which in principle still keeps a weight proportional to the sensitive area. Similar to the data processing of efficiency, there is also a charge map for the 2D graphic view of the performance, as shown in Figure 3.25. The 2D distribution of cluster charge has been

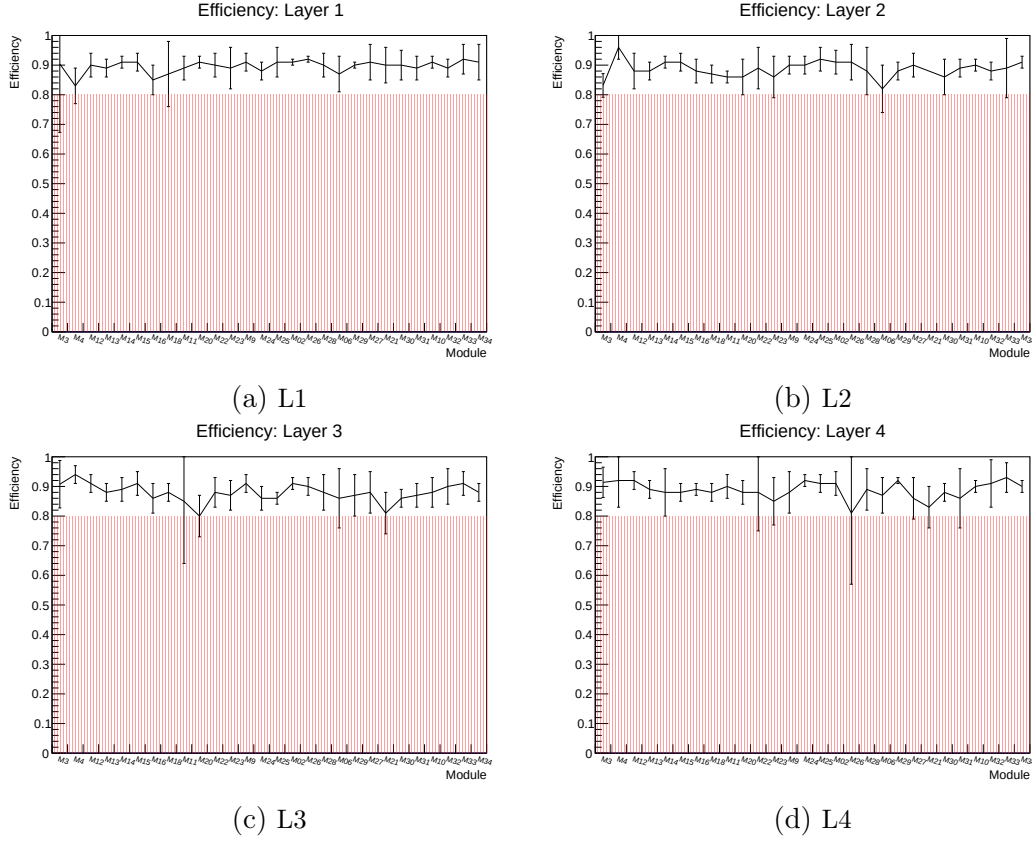


Figure 3.22 – Average efficiency per layer of LM1 modules. The uncertainty bars are the RMS evaluated from efficiencies of 10 HV sectors on the layer. Points above the red band fulfill the requirement of efficiency.

re-binned so that the gap of bad connections and the blank due to interconnection are even clearer on this map.

The average charge per layer is recorded for all the modules, of which the results are shown in Figure 3.26. The first data point here stands for module 3. The values are systematically higher than other modules because the signals are read out without a multiplexing scheme for this module. The Mux2 card enables one electronics channel to the associated two strip in the module, which almost doubles the capacitance attached to the channel and thus brings a lower amplitude. At nominal HV of 570 V, this lower amplitude caused by multiplexing introduces a slightly loss of efficiency.

Another attribute of charge in ADC unit is its reduction with the increase of connected impedance, which reflects the influence of transfer property of DREAM electronics. In an LM1 module, from the small base to the large base of the trapezoid, the strips become longer. When electronics are attached to longer strips, the corresponding capacitance also increases, which introduces a reciprocal relationship to the amplification factor. As a result, there is a tendency of charge going down with larger PCB numbers. Though the connection to strips is not exactly continuous on each PCB panel due to the multiplexing pattern, this trend turns out to be more or less matching the reciprocal of averaged strip lengths per PCB as shown in Figure 3.27, confirming the responding of the transfer function.

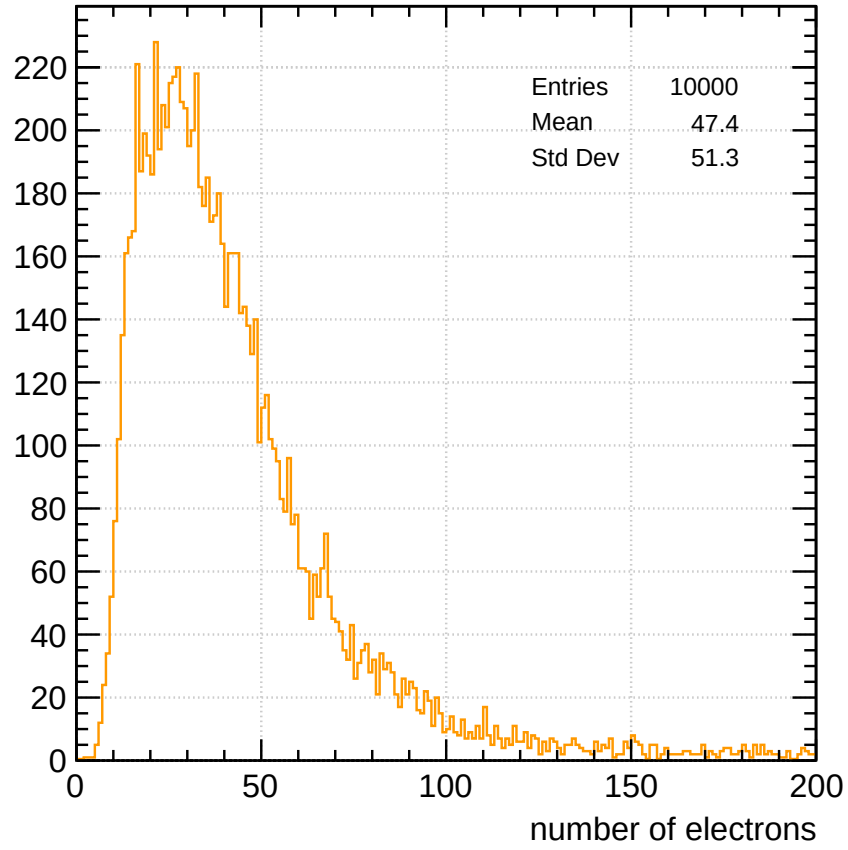


Figure 3.23 – The number of free electrons ionized by cosmic ray muons in a gas gap of 5 mm obtained from simulation. The gas medium is set to 93 %Ar + 7 %CO₂ in the simulation.

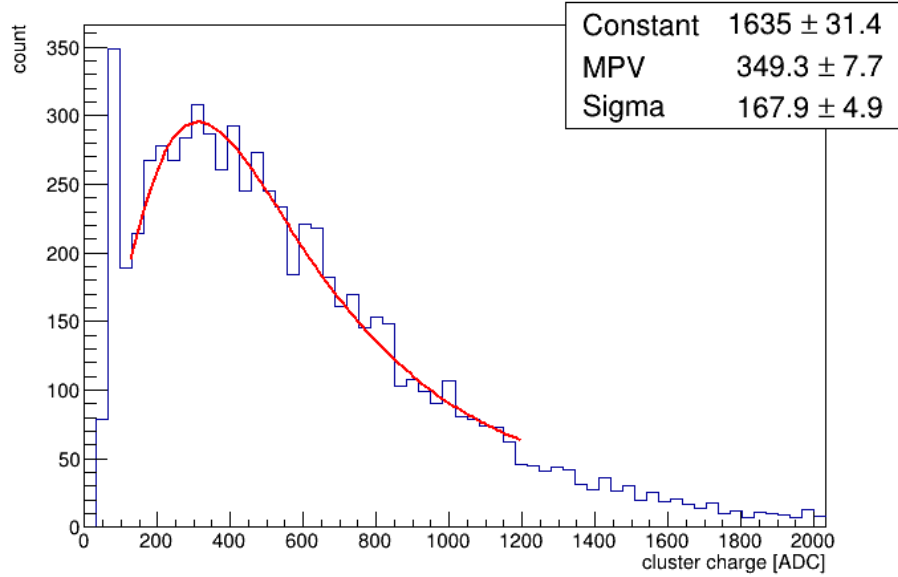


Figure 3.24 – The distribution of charge in ADC unit. This plot shows the charge of M13 PCB4 Layer 2, right HV sector. The small bump at low ADC is mixture of noise and partially reconstructed clusters.

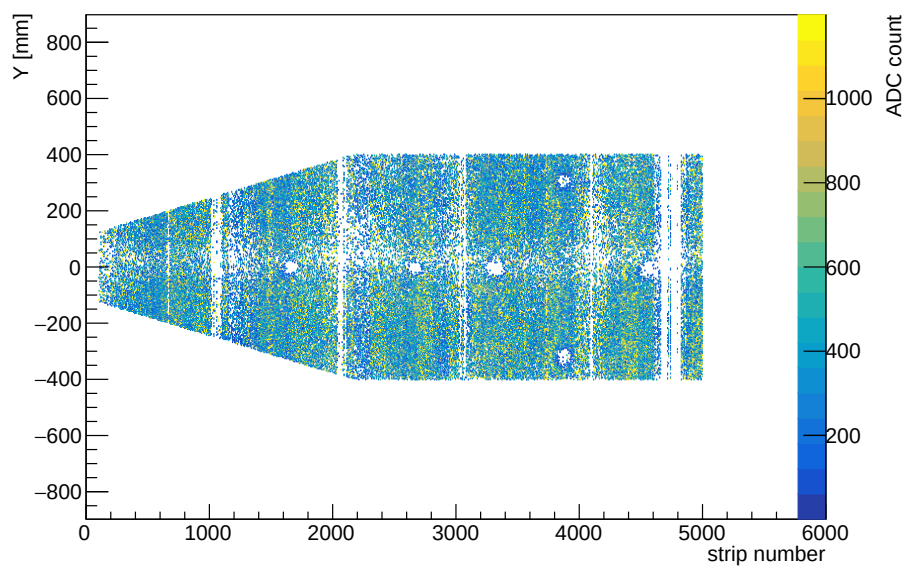


Figure 3.25 – The 2D charge map for Module 11 layer 2, all PCBs. The empty gaps are due to bad connections and passivation between PCBs. The 6 interconnections are clearly revealed on the map.

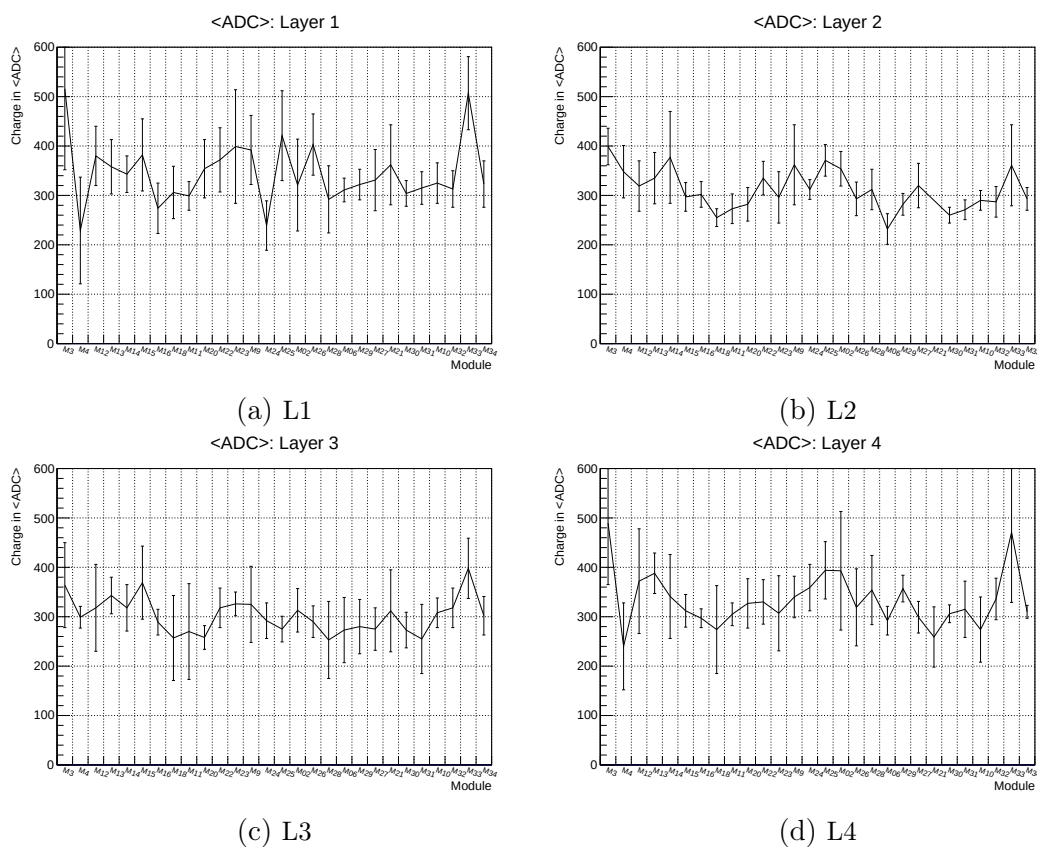


Figure 3.26 – Average charge in ADC unit per layer of LM1 modules. The uncertainty bars are the RMS evaluated from the charge of 10 HV sectors on the layer.

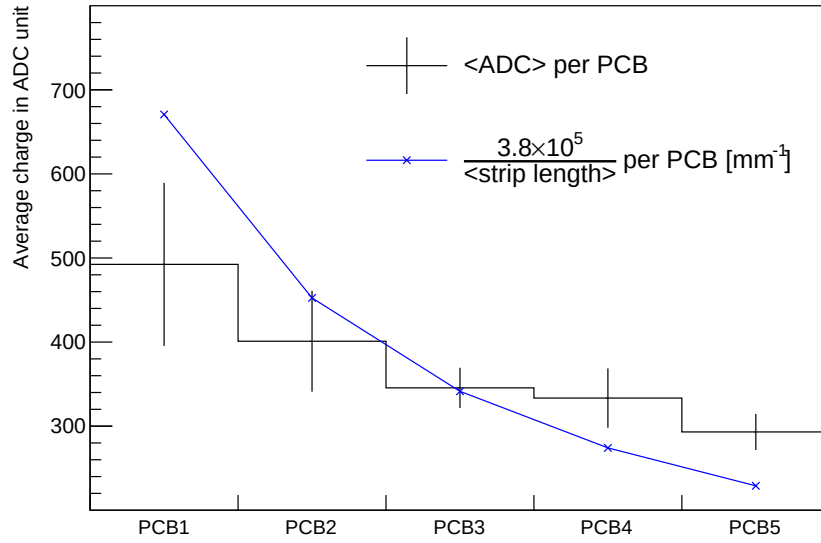


Figure 3.27 – Example taken from Module02: The dependence of charge in ADC unit on PCB. The reciprocal of average strip length per PCB are scaled for a better visualization. The regression of the gain roughly follows the reciprocal of average strip length.

3.6.3 The impact of interconnections

As mentioned in 3.2.2, interconnections are the mechanical points introduced to the design of NSW LM1 Micromegas to support the structure of the detector by well defining the gas gap, thus preventing the inflation of the detector. A close look of the interconnection is shown in Fig 3.28. However, the interface of an interconnection between resistive layer and metal mesh turns out to be susceptible to sparks, causing HV instability of the section or even serious damage in the adjacent area. Therefore, passivation was also applied to the area around interconnections.

The passivation of an interconnection consists of two parts, both on the resistive layer and on the metal mesh. In principle, the two parts of the passivation should well overlap each other. But a perfect overlap can hardly be realized in practice, which brings a even larger dead area than the passivated glue on the resistive layer. Moreover, the additional thickness of passivation around interconnections may deform the metallic mesh and thus locally increase the thickness of amplification gap, further reducing the efficiency of the detector near interconnections. Given that the LM1 modules are designed to be high precision trackers with good efficiency, the loss of efficiency due to the impact of interconnections should be carefully estimated.

This impact was studied for PCB4 over 10 different LM1 modules. The basic idea is to compare the efficiency between an interconnection-included area and a reference area that does not contain any interconnection. A graphic illustration can be seen in Fig 3.29, where 3 rectangles are outlined to cover each interconnection on PCB4. In order to minimize the effect of bad connections, the corresponding reference boxes are chosen to be parallel to the areas of interest. The reference boxes are also designed to have the same surface of the interconnection-included rectangles to ensure a reasonable comparison in term of statistics.

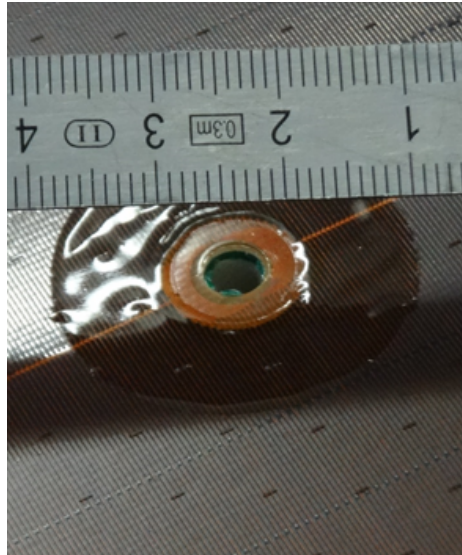


Figure 3.28 – Photo showing a real interconnection with passivation. The total diameter is 30 mm.

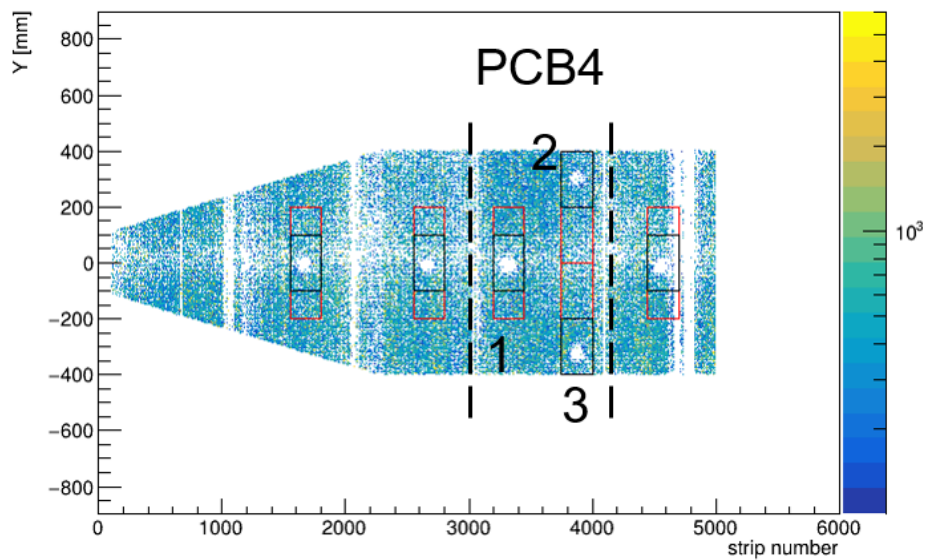


Figure 3.29 – A charge map showing the loss of efficiency due to interconnections. The impact of the three interconnections on PCB4 was estimated. Each box surrounded by black lines include one interconnection. The corresponding reference boxes are indicated in red.

$$\epsilon_{IC} = \epsilon_{Ref} \times (S_{Box} - S_{IC,geo} \times r_{actual/geo}) / S_{Box} \quad (3.2)$$

1645 Equation 3.2 illustrates the relation between the efficiency of box including an
 1646 interconnection and efficiency of the corresponding reference box, where ϵ_{IC} and ϵ_{Ref}
 1647 are the efficiencies in the two boxes, S_{Box} is the surface of the boxes, $S_{IC,geo}$ is the ge-
 1648 ometric surface of an interconnection and $r_{actual/geo}$ is the ratio to be measured. Once
 1649 the efficiencies are measured within the area with interconnection and the reference
 1650 area, it is possible to know what is the actual dead area caused by an interconnection

compared to its geometric size.

After a series of surveys for 10 different modules, it turns out that the loss of efficiency due to interconnections is much larger than what one may expect based on the surface of the glue. The measured surface ratios $r_{actual/geo}$ for the three interconnections on PCB4 are shown in Figure 3.30. Generally speaking, the surface ratio between the consequent dead area on efficiency map and the geometric size of the interconnection is more than 3 according to the measurements. Assuming that the impact of interconnection is at the same level for different PCBs, an optimistic estimation using the factor of 3 would yield a loss of efficiency for the entire sensitive area, which is unfortunately a non-negligible number for Micromegas. Therefore, the passivation of interconnection has to be treated properly in the simulation of the detector to account for the loss of efficiency. The results of this estimation are also given to the construction team of Saclay who then makes efforts to minimize the surface of passivation around interconnections.

3.6.4 Intrinsic spatial resolution

Although the efficiency and gain are without doubt the two most crucial features of the NSW Micromegas modules, the spatial resolution is equally important when the detectors are operated during the actual data taking in ATLAS. As pointed out by the Technical Design Report [2], the precision trackers on the NSW have to measure the transverse momentum of passing muons within a precision of 10 % for 1 TeV muons in the full pseudo-rapidity of its coverage. This can be translated to a position resolution better than 100 μm per detecting layer for the Micromegas modules. The test beam is always a good approach to study the intrinsic spatial resolution of gaseous detectors, provided that the errors such as the imperfection of beam profile are well understood. Nonetheless, the unique setup of the cosmic test bench at Saclay also allows a simple estimation of intrinsic spatial resolution for LM1 Micromegas modules.

For the study of intrinsic spatial resolution using cosmic test data, the readout scheme of the electronics has to be specially prepared: the Mux2 card is no longer used, instead each electronics channel is connected to only one readout strip (Mux1). This choice of readout scheme is mostly due to the inevitable ambiguity of hit-to-strip conversion and the compromised signal-to-noise ratio with Mux2 readout. Even if the measurement of efficiency is not heavily affected by these effects, the clustering is prone to an incomplete reconstruction of charge, which smears determination of the charge center of the cluster. In contrast, data taking without multiplexing not only enhances the signal-to-noise ratio, but also avoids the potential mistake of demultiplexing in the analysis, both of which help improve the quality of cluster reconstruction. Given that the crossing point of cosmic muon track on the detecting layers is inferred from the charge centers of LM1 clusters, the distribution of residual of this inference, or in another word the spatial resolution, will significantly benefit from the Mux1 readout. Since there is no need for performing a scan of intrinsic spatial resolution, this measurement can focus on a single PCB. Therefore, the number of available electronics channels on the bench are still sufficient even without using multiplexing.

Considering the precision of the quantity to be measured, a strategy has to be worked out to reduce the dependence on the millimeter level predictions of muon tracks that are provided by the external trackers, as shown in Fig 3.31. The basic

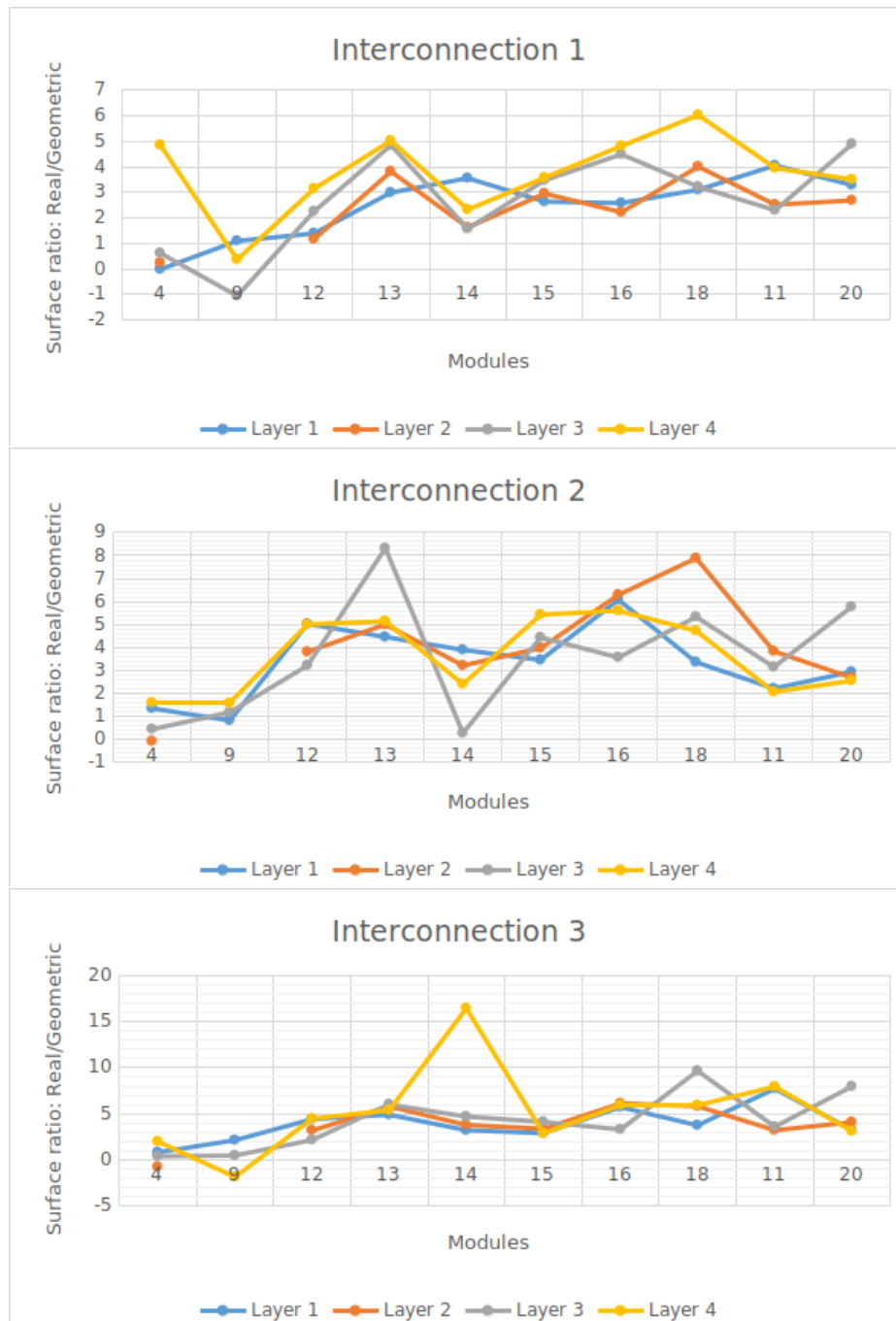


Figure 3.30 – The surface ratios $r_{actual/geo}$ indicating the actual dead area due to the three interconnections on PCB4 compared to their geometric sizes. If ϵ_{IC} occasionally happens to be slightly higher than ϵ_{Ref} in the measurement, negative $r_{actual/geo}$ can be reported.

1697 principle is hence: for each cosmic muon reconstructed by the external trackers, the
 1698 charge center of the LM1 clusters from the Layer1 and Layer2 are recorded if they
 1699 both pass the track matching mentioned in 3.5.4. The charge center of LM1 cluster
 1700 on each layer represents the position where the cosmic muon passes through the
 1701 corresponding detecting layer. Given that the first two layers of LM1 module have
 1702 parallel readout strips along the Y-direction, the difference in the X-position of the
 1703 two clusters between Layer1 and Layer2 ($X_1 - X_2$) helps reveal the spatial resolution
 1704 of LM1 Micromegas (the X, Y directions have been defined in 3.5.4). The residual of

the position difference in X-direction is the result of double smearing by two detecting layers, which suggests that in the vicinity of zero $angleX$ (the projection of the incident angle on X-axis), the standard deviation of the residual divided by $\sqrt{2}$ can be interpreted as the intrinsic spatial resolution of a single detecting layer.

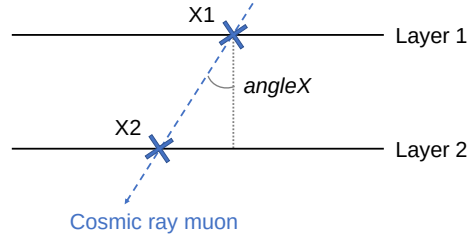


Figure 3.31 – Schematics showing the definitions of $X1 - X2$ and $angleX$ for the measurement of intrinsic spatial resolution.

However, two major issues were discovered as the study unfolded. First, without any available external alignment, there is a limitation of the mechanical alignment for the external trackers at the level of several millimeters, which propagates a bias to the incident angle, adding an unknown offset to $angleX = 0$. More importantly, a correlation exists between $X1 - X2$ and $angleX$, which blurs the spatial resolution measured from any finite interval of $angleX$ even if the interval is close to 0. Solutions to these two difficulties are described in below.

In order to correct for the mechanic mis-alignment of the external trackers, each quarter of the 3-layer Mcube system is studied separately to achieve a small adjustment of the alignment parameters. The corrections include two parts: the optimization of the alignment parameters inside each quarter and offsetting different quarters along the X-axis. The directions of the readout panels of Multigen2D detectors are denoted as X and Y following the coordinate system of the cosmic bench and are each assigned with a number from 0 to 23 (Fig 3.32(a)). For the first part of the correction, the top left quarter is taken as an example for detail illustration: The bottom layer is used as reference such that the X, Y positions of 16, 17 readout panels are never changed in this stage of the correction. Then instead of reconstructing the tracks from all the three layers of Multigen2D (as it is done for the characterization of LM1 modules), the cosmic muon tracks are only reconstructed by 0, 1 together with 16, 17. On the horizontal plane, the mismatch of the reconstructed track with regards to the response of 8,9 is added to the alignment parameters of 8, 9 as correction. After that, muon tracks are reconstructed by 8, 9 and 16, 17 to check its consistency with the response of 0, 1. Similarly, the horizontal offsets in X and Y are added to the alignment parameters of 0, 1. This two procedures are repeated for a few iterations until the muon tracks reconstructed by 16, 17 and 0, 1 (8, 9) match the response of 8, 9 (0, 1). The configuration of the Multigen2D detectors and the distributions of residual for one of them after the alignment is presented in Fig 3.32. The same adjustment is carried out on the other three quarters as well. In the second part, the offset along X-axis between different quarters is observed by the LM1 module. If nothing is done after the correction within individual quarters of the Mcube, the Left and Right halves of a PCB panel of the LM1 module will show different central values of cluster residuals. To compensate for the inter-quarter misalignment, the top left

1741 quarter of the Mcube is chosen to be the baseline while the discrepancy is added to
 1742 2, 11 and 19 as a global offset. The effect of adjusting the X-offset between Mcube
 1743 quarters is shown in Fig 3.33.

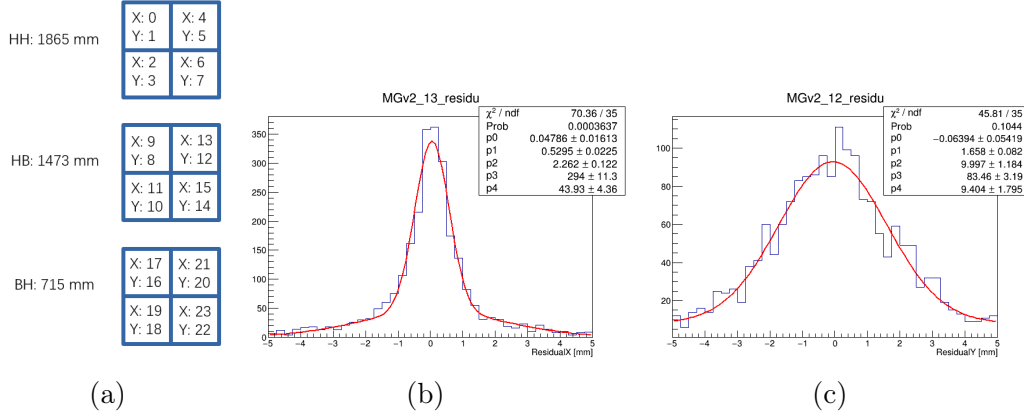


Figure 3.32 – The illustration of Mcube alignment within individual quarters. (a) The indices and directions of Multigen2D readout panels, as well as the heights of 3 layers. (b) The residual of one Multigen2D in X direction centers at zero after the alignment in the corresponding quarter of Mcube. (c) The residual of one Multigen2D in Y direction centers at zero after the alignment in the corresponding quarter of Mcube.

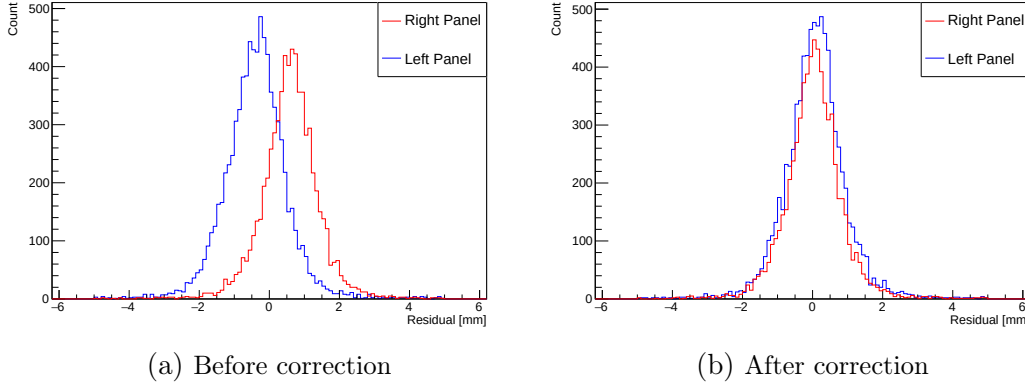


Figure 3.33 – The distributions of cluster residual for M12 Layer1 PCB1 before and after the second part of the alignment. The distributions are separated into left and right halves of the PCB panel to emphasis the improvement brought by the alignment.

1744 Regarding the angle correction, in Fig 3.34(a) a 2-D histogram of " $X1 - X2$ vs
 1745 $angleX$ " is drawn for cosmic muons with $angleX$ smaller than 0.35 rad (or 20°). A
 1746 linear fit is performed to determine the specific amount of correction required for a
 1747 given $angleX$ (Fig 3.34(b)). The result of the linear fit is added to the variable $X1 - X2$
 1748 as a function of $angleX$, so that the average of $X1 - X2$ no longer depends on $angleX$.
 1749 From the same 2-D histogram after angle correction (Figure 3.34(c)), the $X1 - X2$ in
 1750 the central bin of $angleX$ ($-0.05 < angleX < 0.05$) is projected to a 1-D histogram
 1751 in Fig 3.34(d) and fitted by a Gaussian function between $[-0.45 \text{ mm}, 0.45 \text{ mm}]$. Since
 1752 the σ of the Gaussian is the standard deviation of the residual distribution being

measured twice, the intrinsic spatial resolution should be this standard deviation divided by $\sqrt{2}$, which is $110\text{ }\mu\text{m}$.

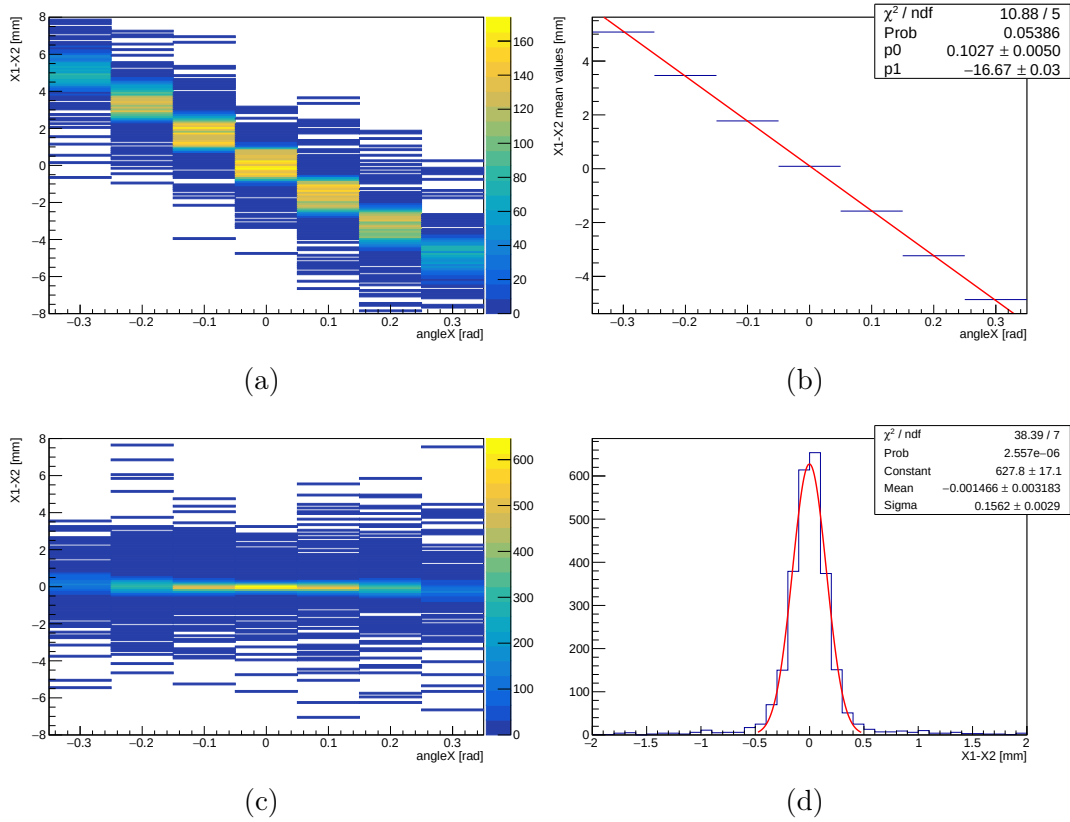


Figure 3.34 – The distributions of cluster residual for M12 Layer1 PCB1 before and after the second part of the alignment. The distributions are separated into left and right halves of the PCB panel to emphasize the effect of the alignment. (a) $X1 - X2$ vs $angleX$ before angle correction. (b) The linear fit for angle correction. (c) $X1 - X2$ vs $angleX$ after angle correction. (d) Gaussian fit of $X1 - X2$ for the bin $-0.05 < angleX < 0.05$.

The estimation made with cosmic bench data proves that the level of intrinsic spatial resolution for a LM1 Micromegas module is within the specifications of the NSW Technical Design Report.

3.6.5 Performance of the detector using isobutane gas mixture

Poor HV performance of NSW Micromegas modules is a common issue frequently reported in all consortia of production when the detectors are operated with 93 %Ar + 7 %CO₂ as working gas. The suppression of sparking can be achieved to some extent by the passivation technique mentioned in 3.2.3. Meanwhile, an alternative solution has also been tested, which involves the replacement of the initial gas mixture with 93 %Ar + 5 %CO₂ + 2 %iso-C₄H₁₀ (isobutane). The presence of isobutane greatly reduces the amplitude of amplification HV needed for reaching the efficiency plateau.

This gas mixture with isobutane is also available on the cosmic test bench at Saclay. Some Mux2 data was taken for PCB1 of Module25 while circulating 93 %Ar + 7 %CO₂ in the gas gaps. As it is shown in Fig 3.35, an efficiency over 90 % on Layer2

1769 can be achieved using the isobutane gas mixture under 515 V. Comparing with
 1770 the HV performance and the efficiency of the 93 %Ar + 7 %CO₂ data in Fig 3.36, it
 1771 is evident that the LM1 module is able to maintain the efficiency at lower HV in
 1772 isobutane gas mixture while keeping more working sectors.

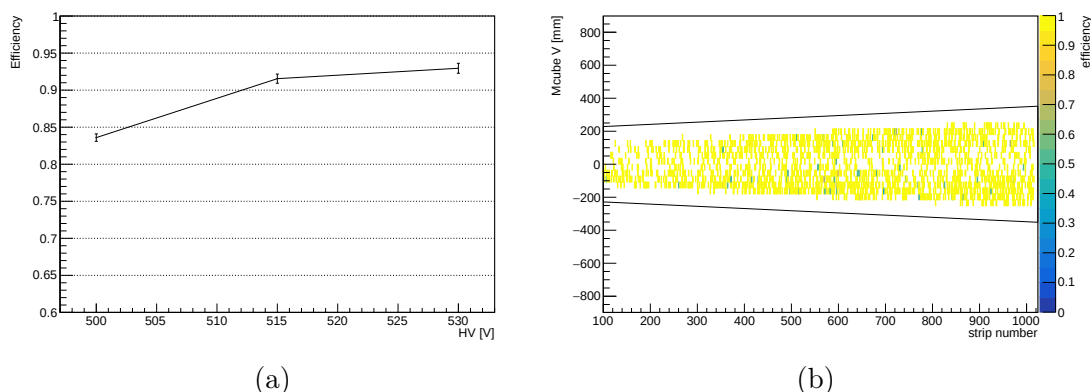


Figure 3.35 – The efficiency of M25 PCB1 Layer2 using 93 %Ar + 5 %CO₂ + 2 %iso-C₄H₁₀ gas mixture. (a) HV scan of efficiency. (b) Efficiency map of M25 PCB1 Layer2 at 515 V.

	Left	Right		Left	Right
Layer 1	0.86	0.88	Layer 1	0.86	0.87
Layer 2	0.94	0.90	Layer 2	0.93	0.90
Layer 3	0.24	0.14	Layer 3	0.87	0.87
Layer 4	0.93	0.89	Layer 4	0.94	0.89

(a) 93 %Ar + 7 %CO₂ (b) 93 %Ar + 5 %CO₂ + 2 %isobutane

Figure 3.36 – The HV status and efficiency of 8 sectors on M25 PCB1 using two different gas mixtures. For the 93 %Ar + 7 %CO₂ mixture, the nominal HV is 570 V. Those sectors unable to sustain nominal HV are labelled in red. For the isobutane gas mixture, the efficiencies are measured at 515 V.

1773 Even though there are signs showing that the isobutane gas mixture is very likely
 1774 to outperform the initial choice of working gas, a major concern with the isobutane is
 1775 its aging effect on the detector. For this reason, dedicated beam test was carried out in
 1776 the Gamma Irradiation Facility (GIF++) at CERN, studying long term degradation
 1777 of the performance for NSW Micromegas modules. No significant evidence shows
 1778 that the 93 %Ar + 7 %CO₂ gas mixture is better than the isobutane gas mixture
 1779 in terms of aging. In February 2022, in view of the improvement of stability and
 1780 experience in aging, ATLAS made the final decision to change the nominal gas to
 1781 93 %Ar + 5 %CO₂ + 2 %iso-C₄H₁₀.

3.7 Summary

Cosmic ray muons provide a powerful tool for the scan and characterization of large area detectors. The cosmic test bench at Saclay is efficient to characterize the performance of ATLAS NSW LM1 Micromegas in terms of efficiency, charge homogeneity, and even intrinsic spatial resolution. Serving as both trigger and external tracker, the Mcube system records cosmic rays and reconstructs their tracks within the acceptance area. While the acceptance of Mcubes is unable to cover the full surface of LM1 Micromegas modules with the given scheme and the pressing time schedule, the test of all readout strips is still possible thanks to the sliding tray and multiplexing readout strategy. Data taking at 3 different locations along the tray rail is enough for the scan of all 5 PCBs in a LM1 Micromegas. Meanwhile, the Mux2 cards adapt DREAM electronics on FEU boards to LM1 module by associating 2 strips to one electronics channel in a carefully-designed pattern so that the total number of channels is reduced by a factor of 2, while the granularity of strips is still retained.

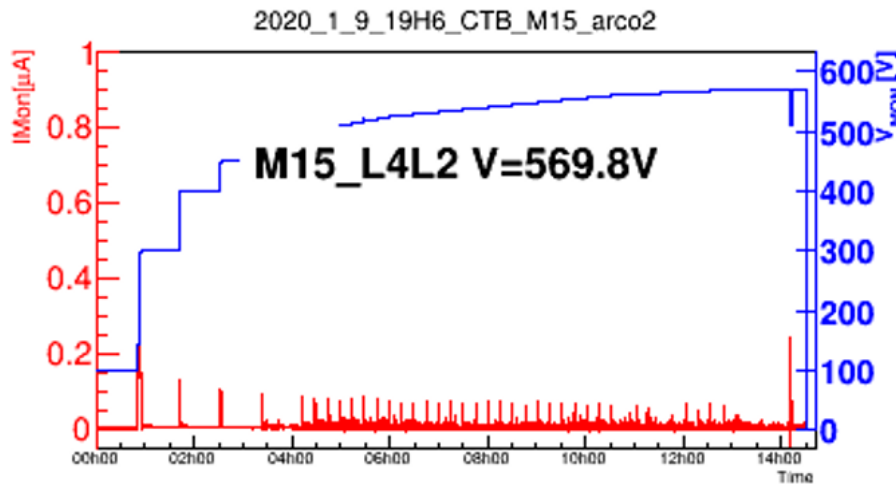


Figure 3.37 – The ramp-up of a sector on M15. The high voltage steadily rises up to the nominal HV.

A complete test protocol has been established for the cosmic muon characterization of NSW LM1 Micromegas at Saclay. Covering the LM1 Micromegas with a plastic tent, the heating on cosmic bench effectively accelerates the drying of modules. The high voltage ramp-up takes place only after the relative humidity of going-out gas reaches below 10 %. An example of HV ramp-up is presented in Fig 3.37. Considering the fragility of HV status for a Micromegas as large as a NSW LM1 quadruplet, an automatic program is hereby introduced to controls the HV ramping up of the module through monitoring and reacting to the leak current in real time. The program prevents rising up the HV of a sector to next preset level unless the current is stable enough according to a certain criteria. All sectors in the module can smoothly achieve the nominal HV of 570 V after 12 hours if everything goes well and without the presence of bad sector. A check of pedestals is always necessary after the installation of the module but before the actual data taking. The immediate output of pedestal

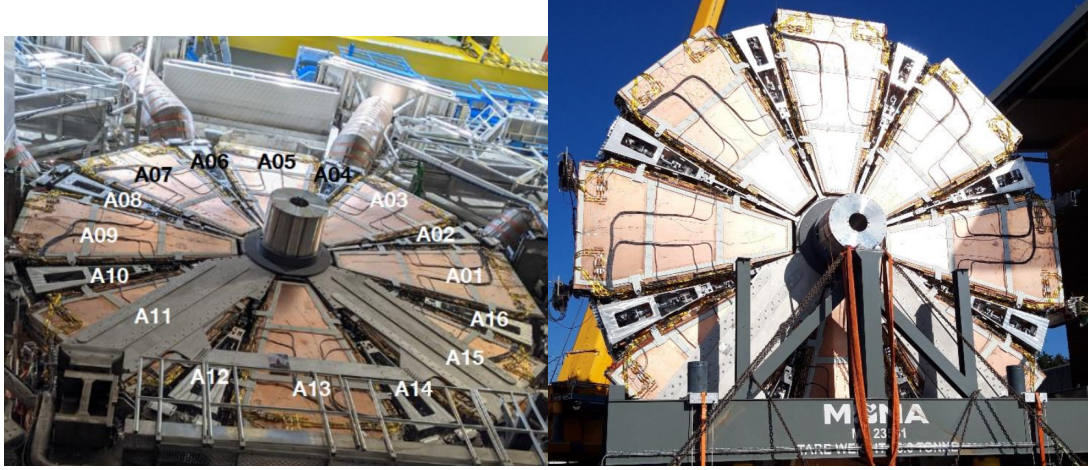
taking offers important clues to some possible connection issues. Once the LM1 module is pushed to the data taking position between Mcube layers HH and HB, a short run will be launched. The "hit map" of this short run can further indicate some other suspicious connection problems so that they can be fixed without going through the entire data analysis process. The entire data taking contains scans of the LM1 module at 3 locations. Given that the change of Mux2 card is necessary when switching to the next data taking location, the connection of electronics has to be done at least three times for a module.

The data analysis of cosmic bench at Saclay is another crucial part of cosmic test, from which the statistics summary will be produced as a basic reference of validation and the input of database. Each step of the analysis program uses root files to pass the processed information in a format easier for the next step to handle. The pedestal and signal samples collected by DREAM electronics go through the pedestal subtraction, hit finding, clustering and are finally matched with cosmic muon rays reconstructed by Mcubes, yielding output such as efficiency maps and charge maps. A demultiplexing algorithm is applied to the clustering stage, for the signals of two strips are readout by only one electronics channel with the help of Mux2 card. In addition to the analysis of LM1 module, there is also another program taking care of the track reconstruction of cosmic muon rays for Mcubes. The possible muon rays reconstructed by Mcubes are selected and matched to the LM1 clusters at the final phase of the analysis.

The data analysis mostly focuses on two properties of LM1 modules: efficiency and charge. Both efficiency and charge are studied in a basis of HV sector. While efficiency of a sector can be obtained by simply counting cluster-track matching events, the Landau function is chosen to fit MPV of the charge distribution. The summary of efficiency and charge per PCB and per layer requires a weighted average over the constitutional HV sectors, where the weight refers to the reverse of squared error in order to represent the surface of sensitive area for each sector. Excluding the effect of bad connections as well as edge passivation, the final efficiencies are able to reveal the intrinsic performance of the detector and the homogeneity for different layers. Based on the same exclusion as final efficiencies, analysis of charge also monitors the charge homogeneity of the tested module. The modules validated on cosmic bench meet the validation requirement and show a good performance in terms of efficiency and charge. The study of intrinsic spatial resolution of LM1 modules with cosmic bench data confirms that the produced detectors satisfy the designed tracking precision. An alternative working gas mixture for NSW Micromegas containing isobutane has been tested on cosmic bench at Saclay. Preliminary results show a positive effect on the detector performance and that the HV stability is improved without compromising the efficiency. Finally, this gas mixture has been chosen for the operation of NSW Micromegas during Run 3.

After the production and validation at Saclay, the LM1 modules are transported to CERN, where they are qualified and assembled into large Micromegas double-wedge chambers together with the LM2 quadruplets. A first HV test is performed for each chamber to further validate the HV performance after assembly. Furthermore, around 70 % of the Micromegas chambers are also put into a HV test under high irradiation at GIF++, so that the behaviour of the detectors can be measured at different incident fluxes. Current instability as well as sparking effects are monitored in the tests at GIF++, since they are likely to affect the performance in the expected 15 years of operation of NSW. Then the Micromegas chambers receive mechanical and services

1858 integration before being equipped with electronics and assembled with sTGC wedges
 1859 to form sectors. During the commissioning, the sectors are installed on the NSWs,
 1860 with services and electronics boxes connected. Both of the NSWs have been lowered
 to the ATLAS cavern at P1 of CERN (Fig 3.38).



(a) NSW A in ATLAS cavern

(b) NSW C in transportation

Figure 3.38 – Photos of ATLAS NSWs

1861 In Run 3, the NSWs are operational and have already become part of the ATLAS
 1862 muon spectrometer for the data taking. More work is still going on to fully exploit the
 1863 benefit of the NSW upgrade, so as to further enhance the performance of the muon
 1864 spectrometer in terms of both trigger and track reconstruction. As an indispensable
 1865 step of the detector production, the cosmic bench test carried out at Saclay is a key
 1866 to the success of the NSW upgrade.
 1867

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1961 Chapter 4

1962 Overview of the W mass 1963 measurements

Contents

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As a quantity tightly linked to the BEH mechanism and spontaneous symmetry breaking described in Section 1.1.3, the mass of the W-boson is a fundamental parameter of the SM. While the W mass can be determined from the precisely known values of the fine-structure constant, the Fermi constant and the mass of the Z-boson at the lowest order, the higher order loop corrections, denoted as Δr , are also predicted by the SM, bringing dependencies on the gauge boson coupling and the masses of the heavy particles in the SM. In particular, Δr is sensitive to the top-quark mass and to the Higgs-boson mass. However, in extended theories, the correction terms also receive contributions from additional particles and interactions of new physics. Therefore, comparing the W mass measured from the experiments and the value predicted by theoretical calculations not only allows the consistency check of the SM, but may also reveal a hint of the physics beyond the SM.

4.1 The history of W-boson measurement

Since the W mass is such an important probe for the consistency check of the SM as well as a possible indication of new physics, countless efforts have been devoted to the precise determination of W mass from experiments. It all started in 1983, when the mass of the W-boson was firstly measured during its discovery at CERN SPS $p\bar{p}$ (proton-antiproton) collider by UA1 at a center-of-mass energy of 546 GeV and UA2 at 630 GeV [1]. Given the six and four W candidates found by the two experiments, the value of W mass was measured to be 81 ± 5 GeV and 80_{-6}^{+10} GeV. The measurements were followed by those of the Tevatron $p\bar{p}$ collider with the CDF and D0 detectors at a center-of-mass energy of 1.8 TeV in Run 1 (1987-1996) and at 1.96 TeV in Run 2 (2002-2011) [2]–[5]. The combined value of the measurements from the two experiments gives a W mass of 80.387 ± 0.016 GeV. In the LEP electron-positron collider, the ALEPH, DELPHI, L3 and OPAL collaborations performed the measurement at center-of-mass energies between 161-209 GeV using $e^+e^- \rightarrow W^+W^-$ process. The combined value of W mass at LEP is 80.376 ± 0.033 GeV [6]–[9]. In 2018, ATLAS published its first W mass measurement at 7 TeV center-of-mass energy: 80.370 ± 0.019 GeV [10]. Before 2021, the average of the direct measurements achieves a result of 80.379 ± 0.013 GeV, an extraordinary result approaching the frontier of the experimental techniques. However, when compared with the SM predictions using the latest top mass and Higgs mass results, the prediction of 80.354 ± 0.007 GeV can hardly set tight constraint to new physics, as the comparison is limited by the experimental precision.

LHCb recently released its proof-of-concept measurement of W mass at 13 TeV using 1.7 fb^{-1} luminosity, giving 80.354 ± 0.032 GeV, which is largely dominated by statistics [11]. A future measurement of the W mass at LHCb using the full LHC

Run 2 data is expected to bring partial cancellation of PDF uncertainties with the ATLAS measurement. A more recent publication of CDF II data, with an integrated luminosity of 8.8 fb^{-1} $p\bar{p}$ collision at 1.96 TeV, presents by far the most precise single measurement of W mass: $80.433 \pm 0.009 \text{ GeV}$ [12]. The precision of the CDF II result has exceeded the combination of all the previous measurements excluding LHCb, approaching the accuracy of the state-of-the-art theoretical prediction. However, a 7σ tension with respect to the SM is also observed by CDF II, which requires confirmation from other experiments, emphasizing the needs for new W mass measurements at LHC with improved systematic uncertainties as well as the reprocessing the 7 TeV ATLAS data using improved method of statistical analysis. A new on-going measurement of W mass at ATLAS with the specialized datasets (Section 4.3.2) will be described in Chapter 8.

4.2 The methodologies of the precision W mass measurements

The precision measurement of the W mass started from the LEP era. As an electron positron collider, LEP produced the W-bosons in pairs via the annihilation of an electron and a positron. After the production, the W^+W^- pair decays subsequently in three modes: The fully hadronic decay results in four jets, of which the total energy adds up to the center-of-mass energy; the semi-leptonic decay yields a dijet product, a lepton signature and the missing energy in the detector brought by the neutrino; the fully leptonic decay comes with low statistics and significant missing energy in the detector. The reconstruction of the W-boson from jets at LEP benefited from the low event rate and the clean environment of the collisions. Also, thanks to the known total energy in the electron-positron collisions, in the case of the leptonic decays, the neutrino energy can be inferred from the conservation of the total energy and of the total momentum in three dimensions. The direct mass reconstruction method was used at LEP at 172, 183 and 189 GeV center-of-mass energies based on the fully hadronic decay and the semi-leptonic decay. Despite the lower branch ratio, the measurement based on the semi-leptonic decay mode had clean signal and no contribution from the dijet pairing problem compared with the fully hadronic decay model. As a result, the two decay modes yielded similar precision. In addition to the direct reconstruction of the W-boson invariant mass, LEP also determined the W mass by measuring the W^+W^- production cross-section at the threshold center-of-mass energy around 161 GeV [13]–[16]. The precision of the threshold analysis is limited by the small size of the dataset taken at 161 GeV.

At hadron colliders, the measurement of the W mass relies on the production of W-boson via the Drell-Yan process, as described in Section 1.2. In contrast to LEP, the direct reconstruction of the W-boson invariant mass is extremely challenging regardless of the W-boson decay mode. In the hadronic decay, the QCD background is orders of magnitude higher than the electroweak production of the W-boson. In the leptonic decay, the conservation of the four momentum can no longer be used at hadron colliders to infer the momentum and the direction of the neutrino. For the above reasons, all W mass measurement at hadron colliders are based on the study of the leptonic decays of the W-boson with the help of the template fit method for statistics inference. The basic principle of the template fit method, as illustrated

in Figure 4.1, is to generate the kinematic distributions associated to the leptonic decay of the W-boson, but with different hypothetical values of the W mass in the simulation. And these simulated kinematic distributions with different hypotheses of the W mass are referred to as templates. The W mass hypothesis of the template that is closest to the corresponding kinematic distribution of the data is taken to be the most probable value of the W mass from the measurement. In order to avoid the bias brought by the experimenter-expectancy effect, when using the template fit, an unknown offset has to be added to the value of the W mass for the production of the templates and is only removed after the measurements are compatible in all channels as well as categories, and the analysis procedure is finalized. In the simplest test statistic of the template fit method, for instance, measurements [10]–[12], the most probable value of the W mass is inferred from statistics without taking into account the bias introduced by the systematic effects. This method of statistical inference is usually referred to as the χ^2 offset method. The systematic uncertainties in the χ^2 offset method are taken into consideration later on by fitting the systematic variations as pseudo data and studying the offsets in the central value of the fit brought by the variations. A more advanced technique of statistical inference based on the template fit method, known as profile likelihood fit, has been recently introduced to the W mass analysis [17], which corrects for the small systematic biases and allows to use the power of data to constrain the systematic uncertainties. Consequently, the profile likelihood fit improves both the accuracy (the central value) and the precision (the total uncertainty) of the measurement, with respect to the χ^2 offset method. More detailed discussions about the template fit method will be covered in Section 8.1.

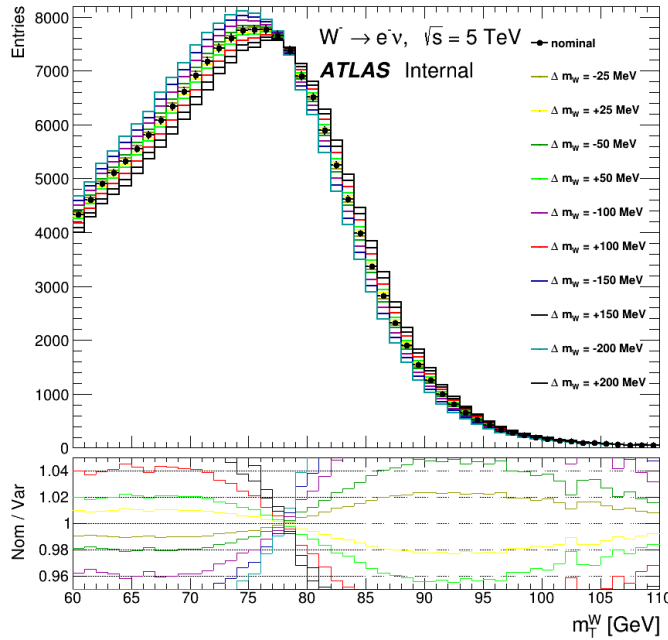


Figure 4.1 – An illustration of the template fit method for the measurement of the W mass. Templates of the reconstructed level kinematic distribution assuming different values of W mass are compared with the nominal simulation. The sensitivity to W mass can be appreciated from the changes in the land-shape.

Despite its success in the measurement of the W mass, one major limitation of the

template fit method arises from the systematic uncertainty in the modelling of the underlying transverse momentum distribution of the W-boson (p_T^W) in the simulation. Conventionally, the modelling of p_T^W is extrapolated from the well measured transverse momentum distribution of the Z-boson (p_T^Z). However, the extrapolation from p_T^Z to p_T^W is inclined to introduce large modelling uncertainty to the measurement. For example, in the ATLAS measurement of the W mass at 7 TeV center-of-mass energy [10], the systematic uncertainty on the p_T^W modelling amounts to as high as 6 MeV out of the 19 MeV total uncertainty. One possible solution to constrain the p_T^W modelling uncertainty is to perform a direct measurement of the p_T^W spectrum with fine granularity and use it as the input to the W mass analysis. Given that the neutrino in the leptonic decay of the W-boson escapes the detection, this p_T^W measurement relies on the precise measurement of the hadronic system in a clean collision environment, which would only be feasible at ATLAS by analyzing the specialized low pile-up dataset that will be introduced in Section 4.3. The measurement of the p_T^W spectrum using the low pile-up data and the propagation of the p_T^W uncertainty from the direct measurement to the W mass analysis will be discussed in Chapter 7 and Section 8.3.2, respectively. The signal modelling and the background estimation of the W-boson analyses will be covered in Chapter 5. Chapter 6 will give a detailed description of the physics objects used in the analyses.

4.3 ATLAS low pile-up data

For pp collisions at LHC, no more than 1 hard inelastic process happens in every bunch crossing. However, additional inelastic pp collision may occur in the same bunch crossing along with the hard process. The number of those additional inelastic pp collisions per bunch crossing other than the one containing the hard process, is referred to as pile-up. The high instantaneous luminosity reached by the LHC implies a high pile-up environment. For instance, the average pile-up ($\langle\mu\rangle$) of the 13 TeV Run 2 dataset is 33.7 [18]. And it is the combination of all the individual collisions that yields the signals seen by the detectors. As a result, deteriorated detector resolution and enhanced soft QCD background largely affect the measurements of W-boson production, imposing significant challenges to the precision measurement of W-boson mass performed in ATLAS. Fortunately, even though low pile-up requires low instantaneous luminosity, the cross-sections for W and Z-boson production are large enough at LHC that a considerable amount of samples can still be collected within a few days of low instantaneous luminosity [19]. For these reasons, dedicated low pile-up data was taken during LHC Run 2 to help improve the relevant precision measurements.

4.3.1 General information of the datasets

The average pile-up is as high as 15-60 for typical LHC Run 2 runs. But for the low pile-up datasets that serve the needs of electroweak precision measurements, such as the measurement of W-boson mass, $\langle\mu\rangle$ is set to a much lower value. Putting aside the performance of the LHC machine, the optimization of $\langle\mu\rangle$ is essentially a trade-off between fast accumulation of samples and high proportion of triggered events with only one collision vertex [20]. Higher instantaneous luminosity comes with higher $\langle\mu\rangle$, which means that the events pass the trigger much more frequently but the

fraction of events with only one interaction is reduced. Toward another extremity, if $\langle\mu\rangle$ approaches zero, almost all the triggered events have exactly one interaction, which enhances the experimental precision. This will be discussed in Section 4.3.2 with more details. However, the low instantaneous luminosity in this data taking condition will result in a longer time for the acquisition of sufficient statistics. Considering the above two aspects, the final decision of $\langle\mu\rangle$ is made to be around 2 for the Run 2 low pile-up datasets as a middle ground.

The Run 2 low pile-up pp collision data used in the work of this thesis covers two center-of-mass energies: 5.02 TeV and 13 TeV [21], [22]. The 5.02 TeV pp collision data was recorded in 2017, with an integrated luminosity of 255 pb^{-1} . The corresponding number of W -boson events after selection is estimated to be about 1.45 million. The total luminosity uncertainty of 5.02 TeV data is 1 %. In 2017, ATLAS collected 147 pb^{-1} of low pile-up data at 13 TeV. Another round of low pile-up data taking was performed at 13 TeV in 2018, of which the integrated luminosity is 191 pb^{-1} . The total integrated luminosity of 13 TeV low pile-up data amounts to 338 pb^{-1} , providing 4.36 million W -boson events after selection. The combined 2017 and 2018 low pile-up data at 13 TeV has a total luminosity uncertainty of 0.92 %. The uncertainties in the integrated luminosity of the low pile-up datasets are determined following the method discussed in [19], using the LUCID-2 detector [23] for the primary luminosity measurement and complemented by the measurements using the inner detector and calorimeters. The luminosity-weighted distribution of $\langle\mu\rangle$ for full Run 2 pp collision data at 13 TeV is shown in Fig 4.2, where the low pile-up data is contrasted by the other datasets from 2015 to 2018.

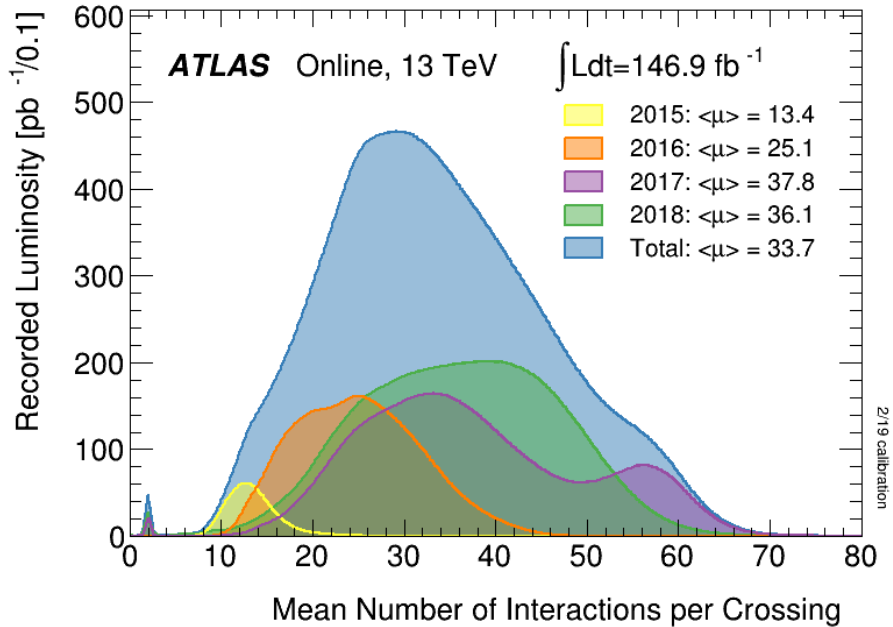


Figure 4.2 – Luminosity-weighted distribution of average interactions per bunch crossing for the full Run 2 pp collision data at 13 TeV. The little bump on the very left side of the plot indicates the 13 TeV low pile-up data recorded in 2017 and 2018 [18].

4.3.2 The physics interest of low pile-up conditions

The existing W mass measurement at ATLAS using 7 TeV data [10] reveals clear difficulties related to the impact of pile-up on the study of W-boson production. The multiple proton interactions as high as $\langle \mu \rangle \approx 9$ degrades the detector response to the hadronic system recoiling against the W decay products in the transverse plane (a.k.a. the recoil or u_T). The poor resolution of the recoil weakens the ability of the data to constrain the modelling of the W transverse momentum (p_T^W) distribution, leading to a large systematic uncertainty in the p_T^W modelling. Furthermore, since some kinematic distributions used for the extraction of W mass, such as the neutrino missing transverse momentum p_T^{miss} and the W transverse mass m_T , are based on the u_T distribution, the statistical sensitivity of these observables to W mass are hence reduced. On top of that, soft QCD background (multijet) in high pile-up conditions is also a limiting factor of the total uncertainty. These three disadvantages are the main limitations of further improving the precision of m_W measurements using high pile-up datasets in ATLAS.

Given that no good modelling of p_T^W to the desired precision of a resolution better than $\sim 6\text{--}7\text{ GeV}$ exists to describe the data of u_T [24], p_T^W has to be extrapolated from p_T^Z after parton shower tuning in most cases. Both the experimental measurement of p_T^Z and the extrapolation contribute to the modelling uncertainty of p_T^W . However, the theoretical uncertainty of the Z to W extrapolation is not only debatable but also crucial to the total uncertainty of W mass. Following this strategy, the p_T^W uncertainty is estimated to be more than 6 MeV in the 7 TeV W mass measurement of ATLAS, according to a quadratic sum of uncertainties related to parton shower tuning, charm quark mass and perturbative calculation of parton shower as well as its corresponding missing higher order uncertainties. With the low pile-up data, it will be possible to directly measure the p_T^W spectrum instead of performing the theoretical extrapolation from p_T^Z . An uncertainty about 1-2% with bin width close to 6-7 GeV is desired for the direct measurement, which is expected to halve the related QCD modelling uncertainty in W mass. The ratio of W^+ over W^- p_T distributions will also achieve a good precision thanks to the cancellation of experimental systematics, rendering better constraints on the theoretical modelling of W-boson production. More detailed discussion of p_T^W and m_W measurement using Run 2 low pile-up data will be covered in Chapter 7 and Chapter 8.

From the experimental side, the Run 2 low pile-up datasets taken at two center-of-mass energies have slightly different features. The data recorded at $\sqrt{s}=13\text{ TeV}$ comes with a better detector calibration and more signal candidates in total, while the $\sqrt{s}=5.02\text{ TeV}$ dataset has in principle smaller underlying event and cleaner signals, including a lower fraction of the multijet background that cannot be reliably simulated. Consequently, the measurements of m_W at two center-of-mass energies are complimentary for their own advantages in different experimental aspects. From the theoretical side, it is very interesting to study the energy dependence of p_T^W spectrum, since the distribution is expected to vary by 20-30% between 5.02 and 13 TeV. Another effect is that the fraction of charm-induced W-boson production increases rapidly from 5.02 to 13 TeV, which means that the precise measurements at both centre-of-mass energies will help understand the impact of heavy flavor quarks on p_T^W spectrum. These studies relevant to the measurement of the p_T^W spectrum will benefit to the constraint on the physics modelling in the m_W measurement.

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2286 Chapter 5

2287 W signal modelling and background 2288 estimation

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5.1 MC samples in ATLAS Run 2 low pile-up data analyses

As mentioned earlier in Section 4.3, dedicated data-taking was carried out during Run 2 at low average interactions per bunch crossing, so as to help the precision electroweak measurements, such as the W transverse momentum (p_T^W) and the W mass (m_W).

The signal process considered in the p_T^W and the m_W analyses is the W-boson production via the proton-proton collision described in Section 1.2, then followed by the leptonic decay of the given W-boson. Only electron and muon are referred to as leptons in the signal process, while W-boson decaying into tau and neutrino is one major electroweak background in the analyses. A few Feynman diagrams of this process at leading-order (LO) and next-to-leading-order (NLO) are shown in Figure 5.1 for illustrative purpose. Since the pile-up distributions are adjusted to those of the datasets, no further pile-up reweighting needs to be applied to the MC samples.

In order to model the signal process, the top-quark related background as well as the electroweak background, MC16 production campaigns are launched to fully simulate and reconstruct the ATLAS Monte Carlo (MC) samples. These MC samples are set to match the special data taking conditions at low pile-up, specifically the pile-up overlay, the topo-cluster noise threshold of the calorimeter and the corresponding trigger menu.

The main signal event samples for W and Z production are generated using the Powheg [1]–[4] event generator using the CT10 global PDF set [5]. For the simulation of the parton shower, Powheg is interfaced to Pythia8 [6] using the AZNLO tune [7]. The effect of QED final state radiation in the Powheg+Pythia8 samples is

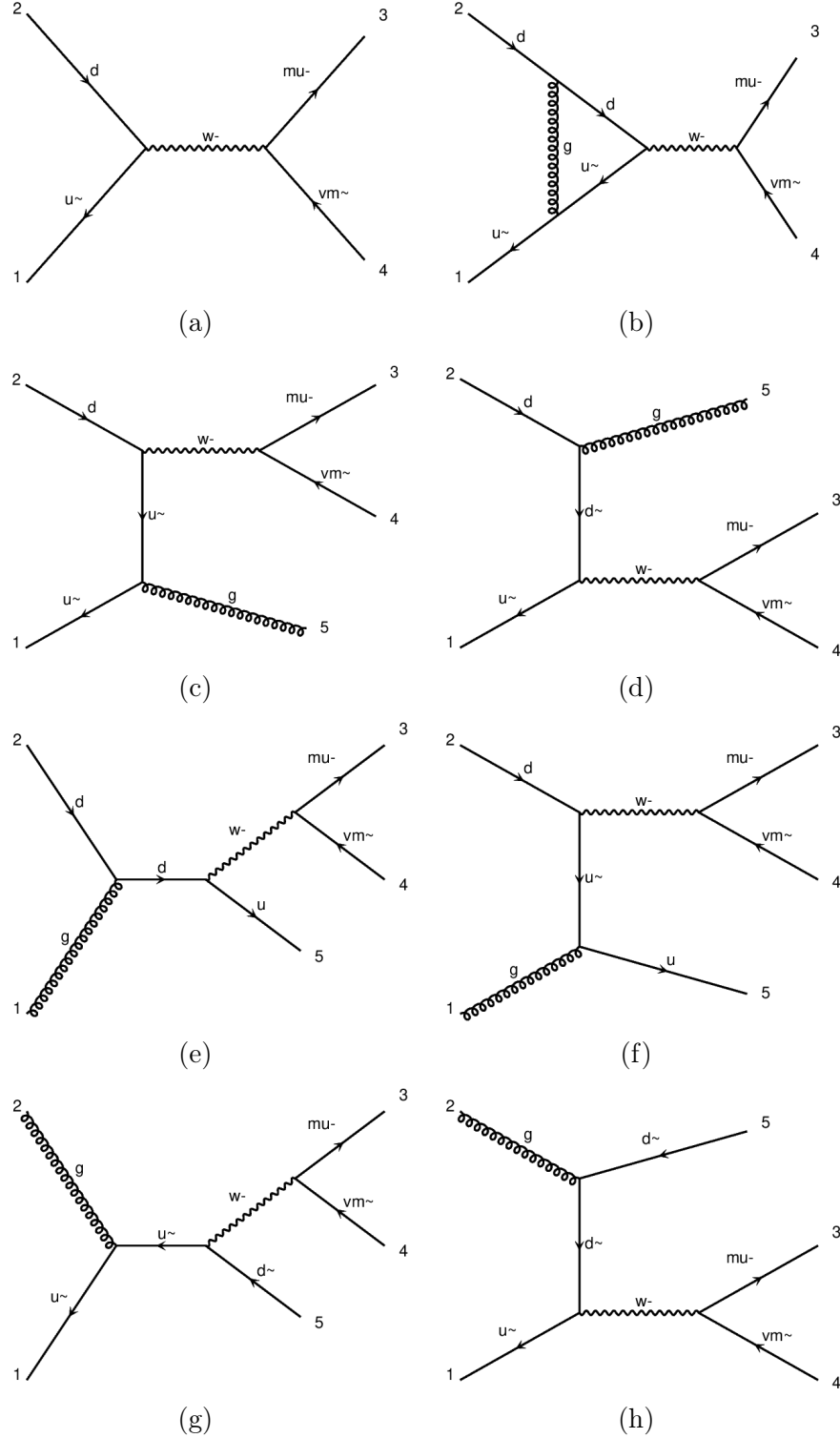


Figure 5.1 – The LO and NLO processes of W^- production and decaying into a pair of $\mu^- \nu$, where the d -quark, $\bar{u}$ -quark and gluon are involved as initial partons. (a) LO process. (b) Virtual loop. (c)(d) Gluon emission. (e)(f)(g)(h) Quark emission.

2334 simulated by Photos++ [8]. The cross-sections of W and Z processes are normal-
 2335 ized to NNLO calculation performed using the DYTURBO [9], [10] program with the
 2336 MMHT2014nnlo PDF set [11]. Given that the uncertainties on the choice of the PDF

set plus the PDF set internal uncertainties ($\sim 3 - 4\%$), the dependence of factorization and renormalization scales ($< 1\%$), as well as the uncertainty of strong coupling constant α_S ($\sim 1 - 2\%$), a conservative total uncertainty of 5% is assigned to the event count predictions normalized using these cross-sections.

Single top-quark production (Wt, t-channel and s-channel) and top-quark pair-production $t\bar{t}$ are generated with Powheg+Pythia8 as well, among which the cross-section of $t\bar{t}$ at 5.02 TeV is taken from the Top++ prediction [12] reported by CMS [13]. Various combinations of di-bosons background ($VV, V = W, Z$) are produced with Sherpa [14] in all the decay channels for at least one real lepton in the final state. Cross-section uncertainties of 10% are assigned to the top-quark related background and the di-boson background, except for $t\bar{t}$ process, of which the cross-section uncertainty is 7%.

The effect of multiple interactions per bunch crossing (pile-up) is modelled by overlying the simulation of minimum bias events over the hard-scattering event. The minimum bias events are generated with Pythia8 using NNPDF2.3LO [15] as the global PDF set and A3 tune [16]. The soft QCD background (multijet or MJ) is derived via a data-driven method described in Section 4.4.

5.2 Object definitions and event selection

The reconstruction of primary vertices requires at least two charged-particle tracks. The primary vertex with the highest $\sum p_T^2$ of all associated tracks is considered to be the hard-scattering vertex. The distribution of z-position of the primary vertex in the simulation is corrected to the data using a data-driven method in Z events. The 5.02 TeV simulation is adapted to perfectly match the data out-of-box. However, at 13 TeV, the data and simulation do not match very well, which mainly originates from the slightly different beam spots in 2017 and 2018 datasets. The residual acceptance effects due to different z-position distribution of the primary vertex can be corrected by reweighting the MC by the data/MC ratio taken from $Z \rightarrow ee$ and $Z \rightarrow \mu\mu$ selections, such that the data and MC distributions by construction match each other well after the reweighting. Despite marginal impact on p_T^W measurement, further studies might be required in the future to understand whether or not the effect of this correction is also negligible in the m_W analysis.

An electron is reconstructed from the clusters of energy deposition in the calorimeter that can be associated to a track in the ID, while a muon is reconstructed as a track crossing the muon spectrometer and the ID. Both types of object have to pass the kinematic cut $p_T^\ell > 25$ GeV. A cut at 25 GeV, above the p_T cut in the lepton trigger, is chosen so as to avoid the migration of events from outside the kinematic region brought by the relatively low resolution of p_T used in the trigger. Considering the acceptance of the detector, the electron is limited to $|\eta| < 2.47$, excluding the transitional region of the calorimeter $1.37 < |\eta| < 1.52$. The muon is required to satisfy $|\eta| < 2.4$. As for lepton identification, the electron candidates must reach the Medium LH identification criteria [17]. Track isolation is applied to the lepton isolation selection, for which the nearby activity is measured by tracks in a cone of size $\Delta R < 0.2$ around the lepton candidate (ptcone20). The electron candidates have to be isolated according to ptcone20/Min(p_T , 50 GeV). The muon candidates are required to pass the Medium identification criteria [18] and be isolated by the ptcone20/Min(p_T ,

2382 50 GeV) track isolation criteria.

2383 The lepton candidates of the $W \rightarrow \ell\nu$ and $Z \rightarrow \ell\ell$ analyses must originate from
 2384 at least one primary vertex that is associated with at least two ID tracks. A trigger
 2385 is supposed to be fired for at least one muon candidate with p_T above 14 GeV or
 2386 loose-likelihood identified electron with p_T above 15 GeV. Two impact parameters are
 2387 introduced in order to select the primary vertex. The longitudinal impact parameter
 2388 of a lepton track is the distance between the track and the primary vertex along
 2389 the beam line (Δz_0) times the $\sin\theta$ of the track. This quantity is required to fulfill
 2390 $|\Delta z_0 \times \sin\theta| < 0.5$ mm. The second parameter describing the impact in the transverse
 2391 plane, known as the significance of the transverse impact parameter, is defined to be
 2392 the transverse impact parameter ($d0$) of the lepton track with respect to the nominal
 2393 beam line divided by its estimated uncertainty ($\sigma(d0)$). The significance has to satisfy
 2394 $|d0|/\sigma(d0) < 5$ for electron candidates and $|d0|/\sigma(d0) < 3$ for muon candidates.

2395 Regarding W-boson production, the neutrino kinematics needs to be calculated
 2396 at a high precision, which can only be achieved by the measurement of the hadronic
 2397 recoil (u_T). In a proton-proton collision, due to the QCD initial state radiation of
 2398 gluon or quark, the vector boson gains a non-zero momentum in the transverse plane,
 2399 formulated by:

$$\vec{p}_T(W/Z) = \vec{p}_T^{\text{lepton1}} + \vec{p}_T^{\text{lepton2}} = -\Sigma \vec{p}_T^{\text{ISR}q,g} = -\vec{u}_T \quad (5.1)$$

2400 the left hand side $\vec{p}_T(W/Z)$ denotes the transverse momentum of the vector boson
 2401 and $\vec{p}_T^{\text{lepton}}$ is the transverse momentum of the leptons. The hadronic recoil, $\vec{u}_T$ is
 2402 therefore the transverse momenta of the partons from the initial state radiation at
 2403 detector level. The hadronic recoil turns out to be an important quantity in the m_W
 2404 measurement. With the use of particle flow objects (PFOs [19]) as input ingredients to
 2405 the reconstruction of u_T in Run 2, instead of reconstructing the vector sum of all topo-
 2406 clusters [20], u_T has become an even more powerful tool for the precise determination
 2407 of neutrino transverse momentum, according to:

$$\vec{E}_T^{\text{miss}} = \vec{p}_T^\nu = -(\vec{u}_T + \vec{p}_T^\ell) \quad (5.2)$$

2408 $\vec{E}_T^{\text{miss}}$ is the missing transverse energy of the neutrino. The W-boson transverse mass
 2409 m_T is then derived from the kinematic information of the charged lepton and of the
 2410 neutrino as follows:

$$m_T = \sqrt{2p_T^\ell E_T^{\text{miss}}(1 - \cos\Delta\phi)} \quad (5.3)$$

2411 here $\Delta\phi$ is the azimuthal opening angle between $\vec{p}_T^\ell$ of the charged lepton and neutrino
 2412 $\vec{E}_T^{\text{miss}}$. Another critical quantity ΣE_T , defined to be the scalar sum of the p_T of all
 2413 PFOs, represents the total event activity. ΣE_T is related to the resolution of the
 2414 u_T measurement and grows with $p_T^V = u_T$. In order to represent the activity from
 2415 underlying event, pile-up as well as the emissions beyond the order of $V + 1$ jet
 2416 process,

$$\Sigma \bar{E}_T = \Sigma E_T - u_T \quad (5.4)$$

2417 is introduced to correct the event activity by the directed recoil activity, disentangling
 2418 the dependence of ΣE_T on boson dynamics.

5.2.1 W analysis event selection

Candidates of $W \rightarrow \ell\nu$ ($\ell = \mu, e$) events are recorded in the data with single-lepton triggers. The primary, unscaled triggers are used in both 5.02 and 13 TeV datasets. *HLT_e15_lhloose_nod0_L1EM12* single-lepton trigger is required for electron, where at least one reconstructed electron with E_T larger than 15 GeV and passing loose identification is required. *HLT_mu14* is applied to muons, requiring p_T of the muon to be higher than 14 GeV.

W events are selected to contain exactly one identified and isolated lepton candidate satisfying the criteria of object definitions. In order to suppress the Z background, any event with additional leptons of the same flavour with p_T higher than 20 GeV meeting certain identification criteria is vetoed. Background fraction, in particular the multijet background, can be effectively reduced once the following two kinematic cuts are applied to the events in addition to the lepton kinematic cut: $E_T^{\text{miss}} > 25$ GeV and $m_T > 50$ GeV. The three cuts of lepton p_T , E_T^{miss} and m_T are kept coherent, such that an optimal trade-off can be achieved in balancing background rejection and signal selection efficiency.

An overview of pre-selection, object selection and event selection for W candidates is summarized in Table 5.1.

Pre-selection		
GRL Vertex	Pass Good Run List Number of primary vertex ≥ 1	
Object selection		
p_T	Electron $p_T > 25 \text{ GeV}$	Muon $p_T > 25 \text{ GeV}$
η	$ \eta < 1.37$ or $1.52 < \eta < 2.47$	$ \eta < 2.4$
Identification	MediumLH	Medium
Isolation	$ptcone20/Min(p_T, 50 \text{ GeV}) < 0.1$	$ptcone20/Min(p_T, 50 \text{ GeV}) < 0.1$
Trigger	<i>HLT_e15_lhloose_nod0_L1EM12</i>	<i>HLT_mu14</i>
Event selection		
N leptons	Exactly one electron or muon	
E_T^{miss}	$E_T^{\text{miss}} = \vec{u}_T + \vec{p}_T^\ell > 25 \text{ GeV}$	
m_T	$m_T > 50 \text{ GeV}$	

Table 5.1 – Summary of pre-selection, object selection and event selection for W-boson candidates in low pile-up data

5.2.2 Calibration of vertex selection efficiency

In the leptonic decay of a W-boson, due to the neutrino in the decay product, only the association of the single lepton to the primary vertex can be performed via the impact parameter cuts. The second track used in the search of vertex has to rely on the hadronic activities associated to the W-boson production, either jets or the underlying event. What is more complicated is that the single lepton used in the vertex association may actually originate from pile-up rather than the hard-scattering,

especially in the case of electron. The reconstructed vertices associated to the tracks of these leptons will not be selected as the event primary vertex due to the leptons not passing the impact parameter cuts. Although Z-boson selected events are generally used to calibrate the W-boson analyses, the two reconstructed leptons in a Z event help the vertex selection become almost independent of the hadronic activities and hardly influenced by the effects that lead to vertex selection inefficiency in W events.

It has been found that the efficiency of lepton matching the primary vertex is not well reproduced by the simulation of W events with Powheg+Pythia8 in the low p_T region [21], mainly due to a mis-modelling in the underlying events. The effect of this mis-modelling is studied by comparing between the data and the MC the fraction of $Z \rightarrow \mu\mu$ events with 0 or 1 additional tracks (excluding the two leptons in the decay product) matching the z-position of the hard interaction. A significant discrepancy is found between data and MC in the low p_T region where the majority of the W-bosons is produced. Therefore, a dedicated correction of vertex selection efficiency is derived for W events.

The correction of vertex selection efficiency relies on the extraction of the efficiency of lepton matching the primary vertex as a function of boson p_T , in three categories of events: with 0, 1, or more than 1 additional tracks matching the z-position of the hard interaction. The inefficiency in each category of events ($1 - \epsilon(p_T)_i$) is obtained from W simulation. The fraction of number of events falling into each category ($frac(p_T)_i$) is extracted from Z events as a function of $p_T^{\ell\ell}$, separately for data and simulation. The overall efficiency for a given sample S, to be either data or a MC simulation, can be expressed in:

$$\begin{aligned} \epsilon(p_T)_S = & frac(p_T)_{(0,S)} \cdot \epsilon(p_T)_{(0,Powheg+Pythia8)} \\ & + frac(p_T)_{(1,S)} \cdot \epsilon(p_T)_{(1,Powheg+Pythia8)} \\ & + frac(p_T)_{(>1,S)} \cdot \epsilon(p_T)_{(>1,Powheg+Pythia8 \rightarrow S)} \end{aligned} \quad (5.5)$$

While $1 - \epsilon(p_T)_{(0,Powheg+Pythia8)}$ and $1 - \epsilon(p_T)_{(1,Powheg+Pythia8)}$ are simply the inefficiencies of lepton matching primary vertex for 0 or 1 additional tracks from Powheg+Pythia8 simulation, a 2D-reweighting is applied to the Powheg+Pythia8 W simulation to match the ($\sum E_T^{charged} \equiv \sum E_T$ of tracks close to the interaction point, $p_T^{\ell\ell}$) distribution in the target sample S before determining $\epsilon(p_T)_{(>1,Powheg+Pythia8 \rightarrow S)}$.

In the nominal case for p_T^W analysis where the data is studied by unfolding, the Powheg+Pythia8 simulation is corrected by a function of boson p_T :

$$\epsilon(p_T)_{data} / \epsilon(p_T)_{Powheg+Pythia8} \quad (5.6)$$

In the p_T^W analysis, the uncertainty on the vertex selection efficiency due to mis-modelling in the underlying events is estimated by comparing two generators that are known to model underlying events differently. So, in order to estimate only the residual difference between the two MC setups generated in different ways, Sherpa is first reweighted to Powheg in $(p_T^{true,V}, y)$, then calibrated to Powheg+Pythia8 nominal signal samples (following the same procedure as how Powheg is calibrated to data). After that, the difference in the results obtained with calibrated Sherpa and the nominal Powheg+Pythia8 is taken as the generator systematic uncertainty and propagated to the analysis chain.

5.3 Modelling of boson kinematics and physics corrections

The inclusive $Z \rightarrow \ell\ell$ and $W \rightarrow \ell\nu$ cross-sections are six-dimensional due to the two final-state leptons coming with fixed masses. A factorization of the differential cross-section has been demonstrated and successfully implemented in the 7 TeV ATLAS m_W measurement [20]:

$$\begin{aligned} \frac{d\sigma}{dp_1 dp_2} = & \left[\frac{d\sigma(m)}{dm} \right] \left[\frac{d\sigma(y)}{dy} \right] \left[\frac{d\sigma(p_T)}{dp_T} \Big|_y \frac{d\sigma(y)^{-1}}{dy} \right] \\ & \times \left[(1 + \cos^2\theta) + \sum_{i=0}^7 A_i(p_T, y, m) P_i(\cos\theta, \phi) \right] \end{aligned} \quad (5.7)$$

The four factorized quantities are related to the di-lepton (or boson) system. p_1 and p_2 are the lepton and anti-lepton four momenta. m , p_T and y are the invariant mass, transverse momentum and rapidity of the di-lepton system. If lepton refers to the negatively charged lepton from W^- or Z -boson as well as the neutrino from a W^+ boson, then θ and ϕ are the polar angle and the azimuth of the lepton in the rest frame of the di-lepton system. A_i 's are the numerical angular coefficients representing the fraction of each helicity-state cross-section of the boson production, depending on boson transverse momentum, rapidity and mass. However, an approximation is introduced as discussed in the 7 TeV measurement that the dependence of A_i 's on the final state invariant mass can be neglected. Each of the angular coefficient is multiplied by a spherical harmonic of order zero, one or two, noted as P_i in the formula.

The above factorization facilitates a simplified modelling of the mass-distribution $d\sigma(m)/dm$, the boson rapidity distribution y as well as the eight angular coefficients A_i predicted by fixed-order calculation, and the boson- p_T distribution at a given rapidity $d\sigma(p_T)/dp_T|_y$, which can be described by resummation or tuned parton shower as discussed in Section 1.2.

5.3.1 Parameterization of boson resonance

The boson resonance ($d\sigma(m)/dm$) in Equation 5.7 is parameterized by a Breit-Wigner distribution. To start with, the gauge interactions and couplings are expressed as:

$$\begin{aligned} \mathcal{L}_{NC} &= \sum_{i=\gamma, Z} \alpha_i \bar{f} \gamma_\mu (v_{f_i} - a_{f_i} \gamma^5) f V_i^\mu \\ \mathcal{L}_{CC} &= \alpha_W V_{f\bar{f}'} \gamma_\mu (1 - \gamma^5) f W^\mu \end{aligned} \quad (5.8)$$

where the gauge coupling constants α_i , the generation-mixing terms $V_{f\bar{f}'}$ and the vector and axial couplings v_i , a_i for the incoming and outgoing fermions are summarized for the gauge bosons in Table 5.2.

Using the same set of notation, the parton-level $q\bar{q} \rightarrow \ell\ell$ and $q\bar{q}' \rightarrow \ell\nu$ cross-

Vertex	$\gamma f \bar{f}$	$Z f \bar{f}$	$W f \bar{f}'$
α_i	α_{em}	$G_\mu m_Z^2 / (2\sqrt{2}\pi)$	$G_\mu m_W^2 / (\sqrt{2}\pi)$
v_i	Q_f	$I_{3f} - 2Q_f \sin^2 \theta_W$	I_{3f}
a_i	0	I_{3f}	I_{3f}
V_i	1	1	$V_{ff'}$

Table 5.2 – The SM gauge coupling constants, fermion vector and axial-vector coupling factors as well as the generation mixing terms. $V_{ff'}$ is known from the CKM matrix.

sections at leading order are written as:

$$\begin{aligned}
\hat{\sigma}(\hat{s}) &\propto \sum_{i,j} \alpha_i \alpha_j V_i V_j B_{ij} P_{ij}(\hat{s}) \\
B_{ij} &= (v_i v_j + a_i a_j)_{in} (v_i v_j + a_i a_j)_{out} \\
P_{ij}(\hat{s}) &= \hat{s} \frac{(\hat{s} - m_i^2)(\hat{s} - m_j^2) + m_i m_j \Gamma_i \Gamma_j}{\left[(\hat{s} - m_i^2)^2 + (m_i \Gamma_i)^2 \right] \left[(\hat{s} - m_j^2)^2 + (m_j \Gamma_j)^2 \right]}
\end{aligned} \tag{5.9}$$

The summation runs over the exchanged gauge bosons in the s-channel. In the case of charged-current interactions mediated by W-boson, the indices i, j only take one value, while in the neutral-current interactions mediated by γ and Z-boson they take two values. This is interpreted as $i=j$ representing the squared amplitude corresponding to the exchange of a given boson in the diagonal and the cross terms when $i \neq j$ describing the contribution of γ -Z interference. The mass and width of a gauge boson i are denoted to be m_i and Γ_i . $\hat{s}$ stands for the available energy for the parton-level process.

The formula above is known as the fixed-width parameterization and was commonly used in the MC generators. An alternative parameterization, as used in the Z peak studies in LEP [22] and in the earlier m_W measurement [20], is the running-width form, where every factor of $m\Gamma$ is substituted by $\frac{\hat{s}}{m}\Gamma$, such that

$$P'_{ij}(\hat{s}) = \hat{s} \frac{(\hat{s} - m_i'^2)(\hat{s} - m_j'^2) + \frac{\hat{s}^2}{m_i' m_j'} \Gamma_i' \Gamma_j'}{\left[(\hat{s} - m_i'^2)^2 + (\frac{\hat{s}}{m_i'} \Gamma_i')^2 \right] \left[(\hat{s} - m_j'^2)^2 + (\frac{\hat{s}}{m_j'} \Gamma_j')^2 \right]} \tag{5.10}$$

When the running-width Breit-Wigner parameterization is used, the physical masses and widths of the gauge bosons should be taken from the PDG values [23] for the nominal simulation.

The two parameterizations are equivalent up to a redefinition of boson resonance parameters [24]:

$$\begin{aligned}
m_i &= m_i' / \sqrt{1 + (\Gamma_i' / m_i')^2} \\
\Gamma_i &= \Gamma_i' / \sqrt{1 + (\Gamma_i' / m_i')^2}
\end{aligned} \tag{5.11}$$

The W and Z samples produced by Powheg+Pythia8 are reweighted using the photon couplings in the vicinity of the Z peak defined as $\alpha_{em}(s)$, taking into account theoretical inputs for lepton loops as well as a mixture of theory and measurements for the

quark contributions:

$$\alpha_{em}(s) = \frac{\alpha_{em}(0)}{1 - \Delta\alpha(s)} \quad (5.12)$$

$$\Delta\alpha(s) = \frac{\alpha_{em}(0)}{3\pi} (13.4955 + 3 \ln s) + 0.00165 + 0.00299 \ln(1 + s)$$

The experimental value of G_μ is taken for the weak boson couplings as a constant.

So far, the formalism of boson resonance doesn't consider the real photon emission corrections. In ATLAS, those corrections are applied to the simulation with Photos as an afterburner, such that photon emissions are generated based on the event record returned by the Powheg+Pythia8, modifying the lepton kinematics in the final state. In Photos [8], multiple photon radiation is generated in the full available phase space starting from the Born-level final state without QED radiation, where the emission of each photon is calculated according to the corresponding QED matrix element. After each emission, the available phase space is updated by excluding the emitted real photon in the vertex, as illustrated in Figure 5.2, so that another iteration is carried out using the same perturbative accuracy. The iteration comes to an end when the energy of the emitted photon is 10^{-7} times smaller than the parent lepton in the decay rest frame.

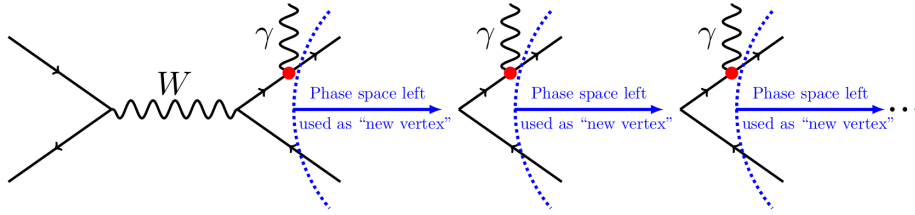


Figure 5.2 – The iterative calculation of QED FSR applied in ATLAS using Photos. Figure taken from [25].

5.3.2 Boson rapidity and transverse momentum

For the m_W measurement, the differential cross-section as a function of boson rapidity, $d\sigma(y)/dy$, is modelled with fixed-order perturbative QCD predictions at $\mathcal{O}(\alpha_S^2)$ in the perturbative expansion of the strong coupling constant using the CT18NNLO PDF set [26].

Given that most gauge bosons generated for the analyses come with a low transverse momentum in the region $p_T^W < 30$ GeV, the predictions of the boson transverse momentum cannot rely solely on fixed-order perturbative QCD calculations. The non-perturbative effects of QCD are accounted for by parton showers to predict the transverse momentum spectrum of vector boson at a given boson rapidity $d\sigma(p_T)/dp_T|_y$ in equation 5.7. The AZNLO tune of Powheg+Pythia8 [7] is chosen to be the parameter setup of parton shower. The CTEQ6L1 PDF set [27] is used for the parton shower.

As it is revealed in the measurement of W transverse momentum with low pile-up data, the default prediction of rapidity-inclusive boson transverse momentum significantly deviates from the data. In order to avoid starting off the analysis with a clearly known non-optimal modelling that is due to bring bias to the physics results, a

2564 data-driven correction is applied to the boson transverse momentum at particle-level,
 2565 such that a good data-MC agreement is observed on boson transverse momentum
 2566 at reco-level. The particle-level reweighting functions are derived based on center-
 2567 of-mass energy, boson charge and the choice of reco-level kinematic distribution as
 2568 observable. More detailed description of this reweighting technique will be covered in
 2569 Section 7.3.

2570 5.3.3 Angular coefficients

2571 At leading order without QCD ISR, the vector bosons are only transversely po-
 2572 larized during the production. The angular distributions of the decay leptons only
 2573 depend on the polar angle of the lepton in the boson rest frame. At higher orders, the
 2574 non-zero transverse momentum brought by QCD ISR affects the polarization states
 2575 of the bosons. Consequently, azimuthal asymmetries arise in the angular distribution
 2576 of the decay leptons [28]–[31]. The polarization of the W and Z boson is completely
 2577 embodied in the eight A_i 's of equation 5.7. The dependence of the differential cross-
 2578 section on the polar and the azimuthal angles of the lepton in the boson rest frame is
 2579 written down explicitly by:

$$\begin{aligned}
 \frac{d\sigma}{dp_T^2 dy dm d\cos\theta d\phi} &= \frac{3}{16\pi} \frac{d\sigma}{dp_T^2 dy dm} \times [(1 + \cos^2\theta) + A_0 \frac{1}{2}(1 - 3\cos^2\theta) \\
 &\quad + A_1 \sin 2\theta \cos \phi \\
 &\quad + A_2 \frac{1}{2} \sin^2 \theta \cos 2\phi \\
 &\quad + A_3 \sin \theta \cos \phi \\
 &\quad + A_4 \cos \theta \\
 &\quad + A_5 \sin^2 \theta \cos \phi \\
 &\quad + A_6 \sin 2\theta \sin \phi \\
 &\quad + A_7 \sin \theta \sin \phi]
 \end{aligned} \tag{5.13}$$

2580 Making use of the completeness of the spherical harmonics, the A_i 's can be extracted
 2581 via the projector defined as follows [32]:

$$\langle P_i(\theta, \phi) \rangle = \frac{\int_{-1}^1 d\cos\theta \int_0^{2\pi} d\phi d\sigma(\theta, \phi) P_i(\theta, \phi)}{\int_{-1}^1 d\cos\theta \int_0^{2\pi} d\phi d\sigma(\theta, \phi)} \tag{5.14}$$

2582 Such that

$$\begin{aligned}
 A_0 &= 4 - 10 \langle \cos^2 \theta \rangle, A_1 = 5 \langle \sin(2\theta) \cos \phi \rangle, \\
 A_2 &= 10 \langle \sin^2 \theta \cos(2\phi) \rangle, A_3 = 4 \langle \sin \theta \cos \phi \rangle, \\
 A_4 &= 4 \langle \cos \theta \rangle, A_5 = 5 \langle \sin^2 \theta \sin(2\phi) \rangle, \\
 A_6 &= 5 \langle \sin(2\theta) \sin \phi \rangle, A_7 = 4 \langle \sin \theta \sin \phi \rangle
 \end{aligned} \tag{5.15}$$

2583 represent the ratios of helicity cross-sections with respect to the unpolarized one. The
 2584 angles θ and ϕ are defined in a boson rest frame with arbitrary momentum along the
 2585 z-axis, known as the Collins-Soper (CS) frame [33].

2586 According to the measurement of Z A_i 's at 8 TeV [34], the fixed-order NNLO
 2587 perturbative QCD predictions of the angular coefficients describe well the data. A_i 's

depend on boson mass, p_T and rapidity. Since the dependence on boson mass is relatively weak, the predictions are made with a fixed boson mass. However, it is understood that the angular correlations between the two decay leptons due to the polarization of the mother vector boson, are not properly modelled in the baseline simulation of Powheg+Pythia8 that produces W and Z samples. The generator only reaches NLO accuracy in QCD and predicts negative A_0 as $p_T^{true,V}$ approaches 0. For this reason, the original angular coefficients are reweighted to NNLO predictions as what is done in the 7 TeV m_W analysis.

With CT10NNLO PDF set [35], DYTURBO predicts the angular coefficients as functions of boson p_T and rapidity, individually for W^+ , W^- and Z. Noting A_i the initial prediction of an angular coefficient and A'_i an alternative prediction, the $\cos(\theta)$ and ϕ of the events are reweighted by the following weight:

$$w = \frac{1 + \cos^2 \theta + \sum_{i=0}^7 A'_i(p_T, y) P_i(\cos \theta, \phi)}{1 + \cos^2 \theta + \sum_{i=0}^7 A_i(p_T, y) P_i(\cos \theta, \phi)} \quad (5.16)$$

where the angular coefficients of Powheg+Pythia8 appear in the denominator and those of DYTURBO at $\mathcal{O}(\alpha_S^2)$ appear in the numerator.

5.4 Estimation of the MJ background

In addition to the electroweak and top-quark related background described in Section 5.1, multijet (MJ) background also contaminates the events with single W or Z-boson final states. The MJ background originates from semi-leptonic decays of heavy quarks, in-flight pion decays. And in the case of electron channels, MJ may come from pion faking an electron or photon conversion as well. In W events, although the MJ background is in principle well rejected by the lepton isolation selection and kinematic cuts on m_T and E_T^{miss} , a non-negligible contribution is still attributed to both the large jet production cross-section and the limited resolution of E_T^{miss} . In contrast, the MJ contribution is much smaller in Z channels due to the requirement of two opposite-sign, well-identified and isolated leptons. MJ background can hardly be simulated precisely at detector-level because its effect after signal selection is essentially a large production cross-section multiplied by a tiny efficiency. Good control is demanded for both quantities in order to obtain the accurate shape and the yield of the MJ background, which is not easily achievable. Instead, based on what has been proposed in [36], a data-driven method is used for the estimation of MJ background and the relevant uncertainties in the low pile-up p_T^W and m_W measurements [37].

The W-boson phase-space in the signal region (SR) that is used for the main measurement of p_T^W spectrum, is defined by the following selections:

- $p_T^\ell > 25 \text{ GeV}$ and $|\eta_\ell| < 2.4$.
- $E_T^{\text{miss}} > 25 \text{ GeV}$.
- $m_T > 50 \text{ GeV}$.
- Lepton isolation as $\text{ptcone20}/\text{Min}(\text{pT}, 50 \text{ GeV}) < 0.1$.

It is shown in [38] that $\text{ptcone20}/\text{Min}(\text{pT}, 50 \text{ GeV})$ tightens the lepton isolation consistently at high u_T region and ensures a good compatibility between electron and muon channels.

Four regions are defined by relaxing the kinematic cuts and inverting the lepton isolation selection (Figure 5.3):

- 2630 — Signal region (SR): Isolated leptons and signal phase-space for p_T^ℓ , E_T^{miss} and
- 2631 m_T .
- 2632 — Fit region (FR): Isolated leptons, but without cut on E_T^{miss} or on m_T .
- 2633 — Control region 1 (CR1): Anti-isolated leptons ($\text{ptcone20}/\text{Min}(p_T, 50 \text{ GeV}) > 0.1$)
- 2634 and relaxed kinematic cuts like FR.
- 2635 — Control region 2 (CR2): Anti-isolated leptons with full kinematic cuts like SR.

Fit Region (FR)	Signal Region (SR)
• Lepton $p_T > 25 \text{ GeV}$ • $m_T > 0 \text{ GeV}$ • $\text{met} > 0 \text{ GeV}$ • Lepton isolation < 0.1	• Lepton $p_T > 25 \text{ GeV}$ • $m_T > 50 \text{ GeV}$ • $\text{met} > 25 \text{ GeV}$ • Lepton isolation < 0.1
Control Region 1 (CR1)	Control Region 2 (CR2)
• Lepton $p_T > 25 \text{ GeV}$ • $m_T > 0 \text{ GeV}$ • $\text{met} > 0 \text{ GeV}$ • Lepton isolation > 0.1	• Lepton $p_T > 25 \text{ GeV}$ • $m_T > 50 \text{ GeV}$ • $\text{met} > 25 \text{ GeV}$ • Lepton isolation > 0.1

Figure 5.3 – The definition of four regions used in MJ study, according to the division of fiducial space and lepton isolation.

2636 The basic idea of MJ estimation is to use the two MJ enriched control regions, CR1
 2637 and CR2, to infer the fraction and the shape of MJ background in SR. Considering
 2638 that the MJ production is concentrated at lower values of p_T^ℓ , E_T^{miss} and m_T , it is
 2639 more practical to first determine the fraction of MJ in FR where m_T and E_T^{miss} cuts
 2640 are relaxed, such that the peak of the MJ distribution facilitates the convergence of
 2641 the fraction fit:

$$N_{data}^{FR} = \alpha \cdot N_{EW+top}^{FR} + T \cdot N_{MJ}^{CR1} \quad (5.17)$$

2642 N_{data}^{FR} , N_{EW+top}^{FR} and N_{MJ}^{CR1} stand for the total number of data events in FR, expected
 2643 contribution from electroweak and top-quark related process in FR and the number
 2644 of MJ events in CR1 ($N_{MJ}^{CR1} = N_{data}^{CR1} - N_{EW}^{CR1}$). T is the scale of MJ template in FR to
 2645 be determined by the fit, while α represents an overall normalization of contribution
 2646 from electroweak and top-quark related process approximately equal unity, within the
 2647 uncertainties in luminosity and production cross-sections. Then the MJ yield in FR
 2648 is given by:

$$N_{MJ}^{FR} \approx T \cdot N_{MJ}^{CR1} \quad (5.18)$$

2649 While SR and CR2 share the same set of kinematics cuts, FR and CR1 also share
 2650 another set of kinematics cuts. The only differences between these two sets of kine-
 2651 matics cuts are the additional selections on m_T and E_T^{miss} in SR and CR2. Therefore,

the yield in FR can be extrapolated to the yield in SR with a transfer factor ϵ derived from the two control regions:

$$N_{MJ}^{SR} = \epsilon N_{MJ}^{FR} = \frac{N_{MJ}^{CR2}}{N_{MJ}^{CR1}} \times N_{MJ}^{FR} = \frac{N_{data}^{CR2} - N_{EW+top}^{CR2}}{N_{data}^{CR1} - N_{EW+top}^{CR1}} \times N_{MJ}^{FR} \quad (5.19)$$

The shape of MJ template used in the fraction fit in FR, as well as the expected MJ shape in SR, has to be derived from CR1 and CR2, respectively.

Section 5.4.1 will describe a dedicated correction applied to the MJ estimation so as to eliminate a bias in the recoil construction in CR1 and CR2. Section 5.4.2 and 5.4.3 explain how the yield and the shape of MJ in SR are obtained. Section 5.4.4 is about the uncertainty sources of MJ estimation for p_T^W analysis. Section 5.4.5 summarizes the results of MJ for p_T^W analysis.

5.4.1 Recoil isolation correction

During the recoil reconstruction for a W event, all the PFOs are used for the given event, but excluding the cone of $\Delta R = 0.2$ around the selected lepton. The energy of pile-up and underlying event removed in this cone is subsequently compensated by looking at another cone with the same size centered at the same $|\eta|$ but in an azimuth away from any lepton or hard activity. The energy in this selected cone is added to the recoil, with its orientation aligned to the direction of the removed lepton. In SR and FR, the leptons are isolated, meaning they are typically away from jets and are surrounded by soft activities. The removal and replacement of the cone does not introduce a bias to the recoil. However, in anti-isolated regions, the lepton in interest is on average close to a jet. Replacing the cone around the lepton with another cone of soft activities removes the hard contribution and biases the recoil reconstruction. The m_T and E_T^{miss} are affected by this bias, since the computation of their values uses the recoil as input.

The projection of the recoil along the direction of the lepton momentum $u_{||}^\ell = \vec{u}_T \cdot \frac{\vec{p}_T^\ell}{|\vec{p}_T^\ell|}$, is pertinent to the bias, since the magnitude in its direction is affected by this cone replacement on average. In the absence of bias, $u_{||}^\ell$ is mostly independent on lepton isolation, whereas the dependence of $u_{||}^\ell$ on `ptcone20` indicates the bias. The quantity `ptcone20`, sum of transverse momentum of tracks around a lepton, is correlated to the energy to be removed around the lepton, even though the neutral components are not accounted for in `ptcone20`. So, a correction to recoil u_T as a function of `ptcone20` is proposed:

$$\vec{u}_T^{\text{corr}} = \vec{u}_T^{\text{baseline}} + (k_{iso} \times \text{ptcone20}) \frac{\vec{p}_T^\ell}{|\vec{p}_T^\ell|} \quad (5.20)$$

where u_T^{baseline} is the recoil before correction, and u_T^{corr} the recoil after the correction. The linear coefficient k_{iso} takes into account the possible difference in the scale between `ptcone20` and the total energy removed in the cone around the lepton. Considering the correlation between `ptcone20` (contributions from only the charged components) and the total energy of hard activity (both neutral and charged components) within the cone remains unknown, the coefficient k_{iso} has to be measured from the data by performing a linear fit:

$$\langle u_{||}^{\ell, \text{baseline}} \rangle = -k_{iso} \langle \text{ptcone20} \rangle + b \quad (5.21)$$

Channel	k_{iso}
13 TeV $W^- \rightarrow e^- \nu$	1.179 ± 0.009
13 TeV $W^+ \rightarrow e^+ \nu$	1.178 ± 0.009
13 TeV $W^- \rightarrow \mu^- \nu$	1.426 ± 0.004
13 TeV $W^+ \rightarrow \mu^+ \nu$	1.429 ± 0.004
5.02 TeV $W^- \rightarrow e^- \nu$	1.141 ± 0.019
5.02 TeV $W^+ \rightarrow e^+ \nu$	1.160 ± 0.020
5.02 TeV $W^- \rightarrow \mu^- \nu$	1.383 ± 0.009
5.02 TeV $W^+ \rightarrow \mu^+ \nu$	1.422 ± 0.007

Table 5.3 – The fitted values of k_{iso} for 13 and 5.02 TeV, the uncertainty are the statistical uncertainty from the fitting.

The term b is the constant in the linear fit, representing the average $u_{\parallel}^{\ell}$ in the anti-isolated regions. The correction procedure is shown Figure 5.4, where the large dependence of $u_{\parallel}^{\ell}$ on ptcone20 is reduced to almost zero after the correction. The value of k_{iso} is fitted for each individual channel as the opposite of the slope in the linear fit performed between 5 and 25 GeV. The results of k_{iso} are listed in Table 5.3, in which the values are consistent for each lepton flavor at a given center-of-mass energy. The corrections needed for muons are larger than electrons, which is interpreted as the electron Bremsstrahlung energy leaking outside the cone is not fully removed from the recoil before the correction. The fraction fits in FR performed before and after the correction are compared in the left and middle columns of Figure 5.7, where the MJ template is taken from an anti-isolated slice in CR1. As an example, Figure 5.5 show the effect of the recoil isolation correction on the $\text{ptcone20}/\text{Min}(\text{pT}, 50 \text{ GeV})$ scan in the electron and muon channels at 13 TeV. A strong impact is seen for the dependence of the fitted yield on the isolation range.

The bias discussed above is corrected for in the hadronic recoil reconstruction only for the MJ background estimation. In isolated regions, recoil reconstruction is consistent with applying no correction, as the value of ptcone20 is in general close to zero if the lepton isolation cut is satisfied.

5.4.2 Yield extrapolation

The determination of MJ yield in SR starts from the MJ fraction fit in FR. Given that the shape of MJ distribution in the FR is not known in the first place, it has to be inferred from the CR1. MJ background is the dominant contribution in CR1 due to the inverted lepton isolation and relaxed kinematic cuts. The data in CR1 subtracted by the small predicted contribution from the electroweak and top-quark related processes is expected to well represents the MJ distribution in CR1. However, the shape of MJ distribution is not independent of lepton isolation, which means the shape of MJ in CR1 is different from that in FR and should not be taken directly as the template for the fraction fit. In order to catch the dependence of MJ shape on lepton isolation in CR1, four consecutive anti-isolated slices are chosen, with isolation variable lays between $[0.1, 0.2]$, $[0.2, 0.3]$, $[0.3, 0.4]$ and $[0.4, 0.5]$. In each slice $A < \text{ptcone20}/\text{Min}(\text{pT}, 50 \text{ GeV}) < B$, the MJ shape for the normalized distribution of X in CR1, after the

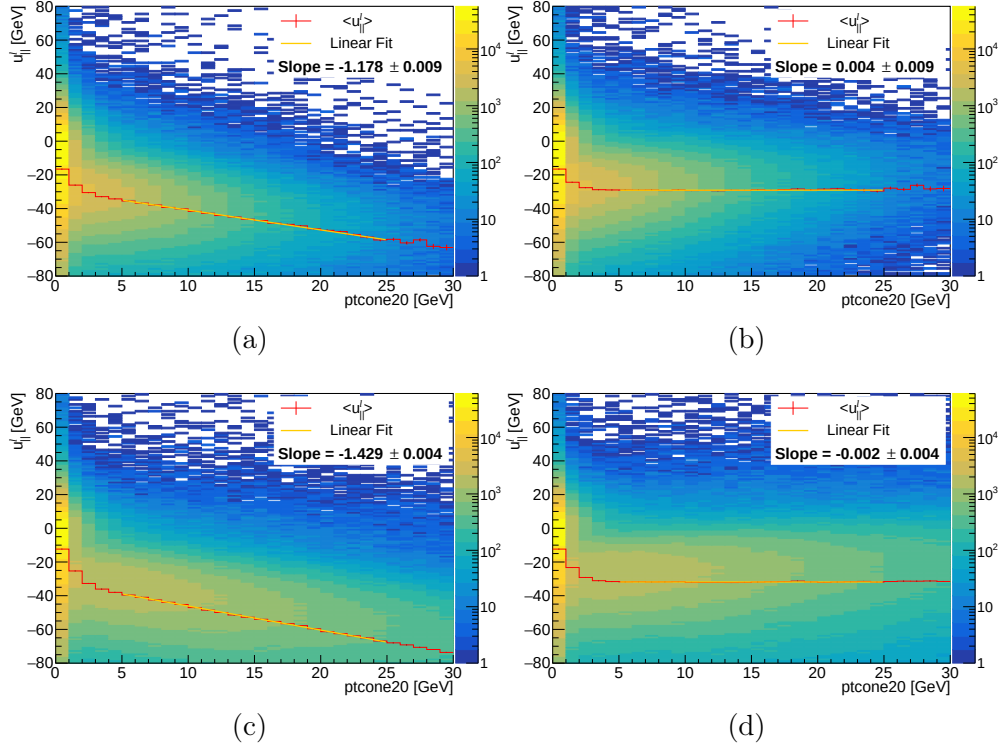


Figure 5.4 – $\langle u_{||}^{\ell} \rangle$ as a function of ptcone20 before (left) and after (right) recoil isolation correction in CR1+FR of data for 13 TeV $W^+ \rightarrow e^+ \nu$ channel (a)(b) and $W^+ \rightarrow \mu^+ \nu$ channel (c)(d).

2721 subtraction of EW and top-quark related processes, is defined to be:

$$H_{MJ}^{[A,B]}[X] = H_{data}^{[A,B]}[X] - H_{EW+top}^{[A,B]}[X] \quad (5.22)$$

2722 The assumption that the MJ shape changes linearly with respect to the isolation
 2723 variable is justified in Figure 5.6, so that the degree of change in MJ shape by every
 2724 0.1 of ptcone20/Min(pT, 50 GeV) is approximately:

$$\Delta^{CR1}[X] = \frac{1}{2} \left\{ \frac{H_{MJ}^{[0.1,0.2]}[X] - H_{MJ}^{[0.3,0.4]}[X]}{2} + \frac{H_{MJ}^{[0.2,0.3]}[X] - H_{MJ}^{[0.4,0.5]}[X]}{2} \right\} \quad (5.23)$$

2725 with which the MJ shape in FR is uniquely determined via:

$$H_{MJ}^{FR}[X] = H_{MJ}^{[0.1,0.2]}[X] + \Delta^{CR1}[X] \quad (5.24)$$

2726 The correction term $\Delta[X]$ used in the above formula is purely built upon the MJ
 2727 shapes in the chosen four anti-isolated slices of CR1. For MJ template taken from
 2728 each of the four slices, its shape is corrected to $H_{MJ}^{FR}[X]$ in 5.24 then normalized to
 2729 the initial integral of MJ distribution in that slice. For this reason, once the shape
 2730 correction is implemented in CR1, no matter which anti-isolated slice the MJ template
 2731 comes from, the fraction of MJ in FR determined by the fraction fit is always the
 2732 same for a given observable X (Figure 5.8(b)(e) in contrast to (a)(c)). Significant
 2733 improvement can be seen from the middle and right columns of Figure 5.7 after the
 2734 shape correction in CR1.

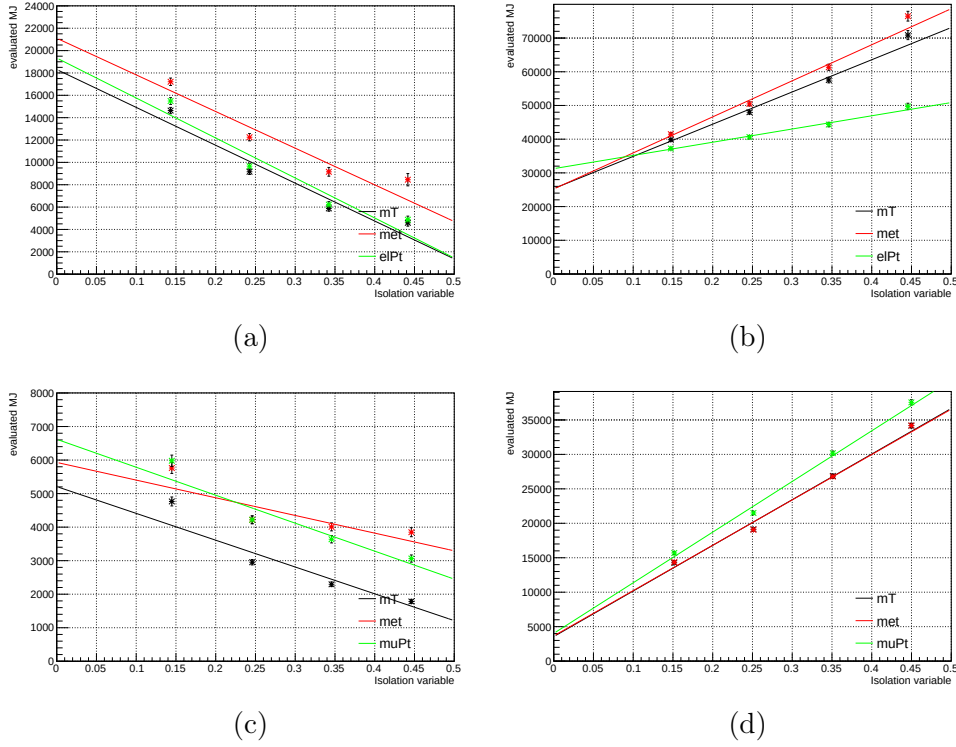


Figure 5.5 – The estimated total number of the MJ background in SR scanned as a function of the isolation. Each point represents an interval of anti-isolated region, of which the horizontal coordinate is the average $p_{Tcone20}/\text{Min}(p_T, 50 \text{ GeV})$ in the corresponding slice of CR2. The scans are shown for 13 TeV $W^+ \rightarrow e^+ \nu$ (top) and $W^+ \rightarrow \mu^+ \nu$ (bottom), without (left) and with (right) recoil isolation correction. No correction has ever been applied to the shape of MJ templates derived in CR1.

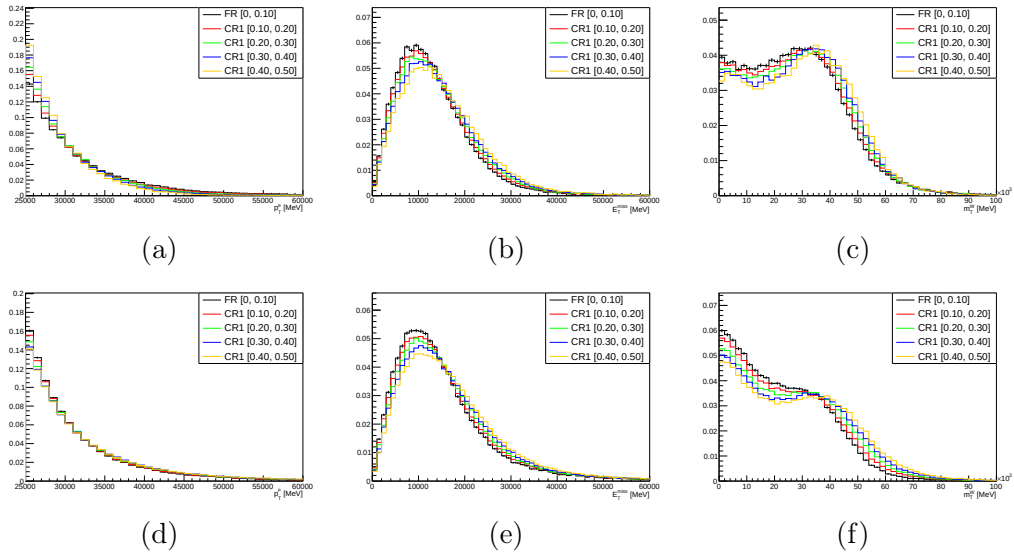


Figure 5.6 – Normalized MJ shapes in anti-isolated slices of CR1 and the shape corrected to FR for 13 TeV $W^+ \rightarrow e^+ \nu$ (top) and $W^+ \rightarrow \mu^+ \nu$ (bottom). (a)(d) p_T^ℓ , (b)(e) E_T^{miss} , (c)(f) m_T .

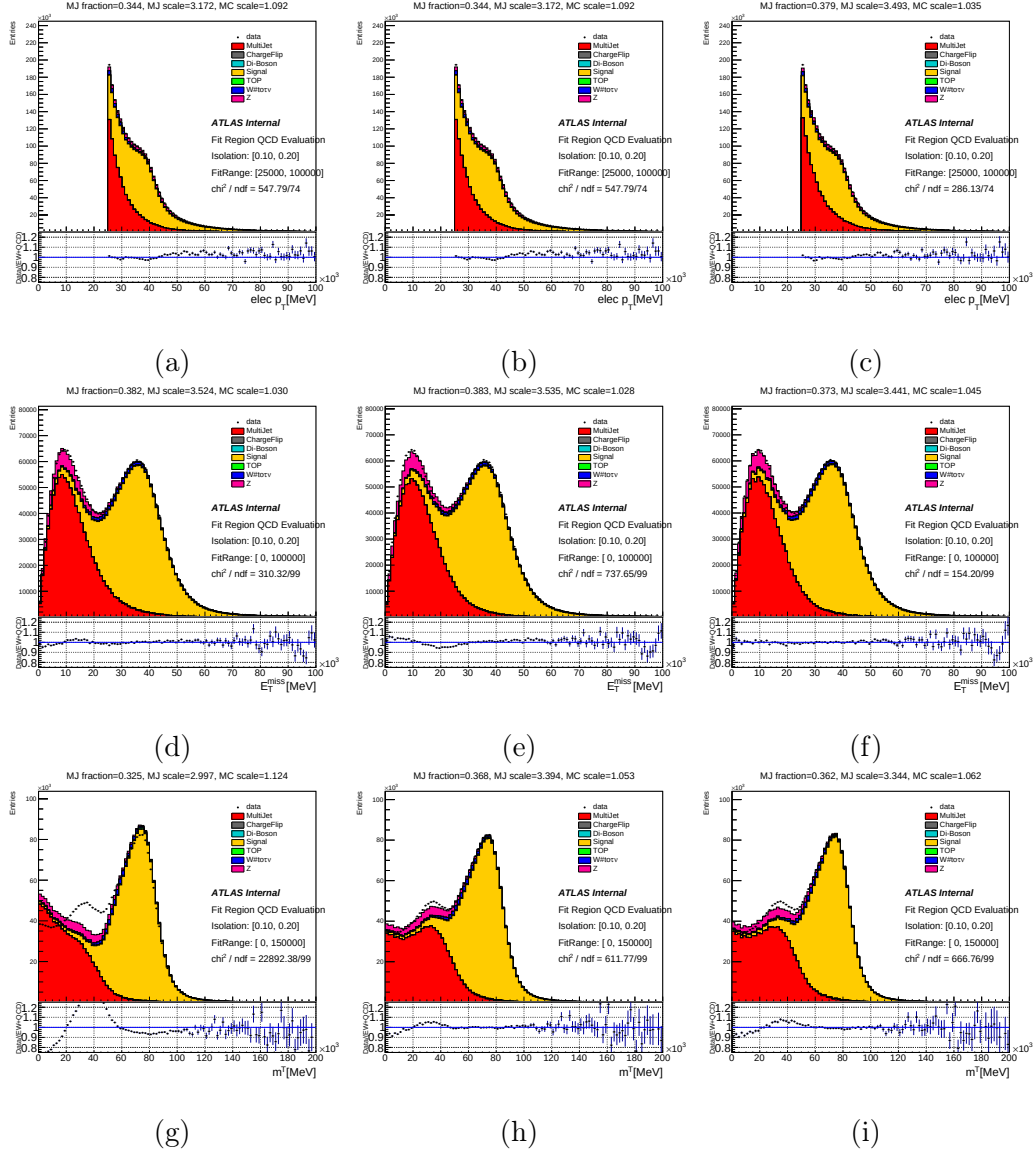


Figure 5.7 – $W^+ \rightarrow e^+ \nu$ channel at 13 TeV. The FR post-fit control plots showing MJ fraction fit using p_T^ℓ (top), E_T^{miss} (middle) and m_T (bottom) distributions. (a)(d)(g) Without recoil isolation correction or shape correction in CR1. (b)(e)(h) After recoil isolation correction but before shape correction in CR1. (c)(f)(i) With both recoil isolation correction and shape correction in CR1.

in SR. Yet the ratio of $\epsilon = N_{MJ}^{CR2}/N_{MJ}^{CR1}$ is not a constant because the efficiencies of the tighter E_T^{miss} and m_T cuts in the control regions depend on lepton isolation. This dependence is visualized in Figure 5.8, where the value of isolation for each point is taken from the average value of isolation in individual anti-isolated slice of CR2. The anti-isolated slices in CR2 are defined consistently with CR1. For now, a simple linear relationship is assumed. A fair good match between the yield of data-driven method and the yield of $b\bar{b}$ and $c\bar{c}$ samples in Section 5.4.5 validates the assumption of the linear extrapolation.

With the known N_{MJ}^{FR} and the linear extrapolation of ϵ , the yield in SR becomes:

$$\begin{aligned}
 N_{MJ}^{SR}(ptcone20/Min(pT, 50 \text{ GeV}) = x) &= \epsilon(x) N_{MJ}^{FR} \\
 &\approx k \cdot x + N_{MJ}^{SR}(ptcone20/Min(pT, 50 \text{ GeV}) = 0)
 \end{aligned}
 \tag{5.25}$$

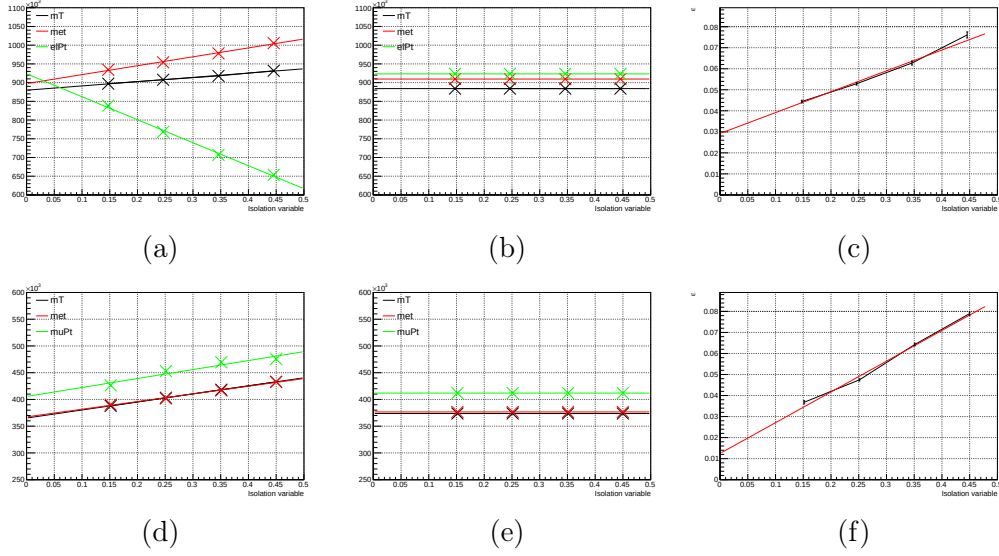


Figure 5.8 – Dependence of the predicted MJ yield in FR on the anti-isolated slices in CR1 before shape correction (left), after shape correction (middle) and the relation between ϵ and isolation variable (right) for $W^+ \rightarrow e^+ \nu$ channel (a)(b)(c) and $W^+ \rightarrow \mu^+ \nu$ channel (d)(e)(f) at 13 TeV.

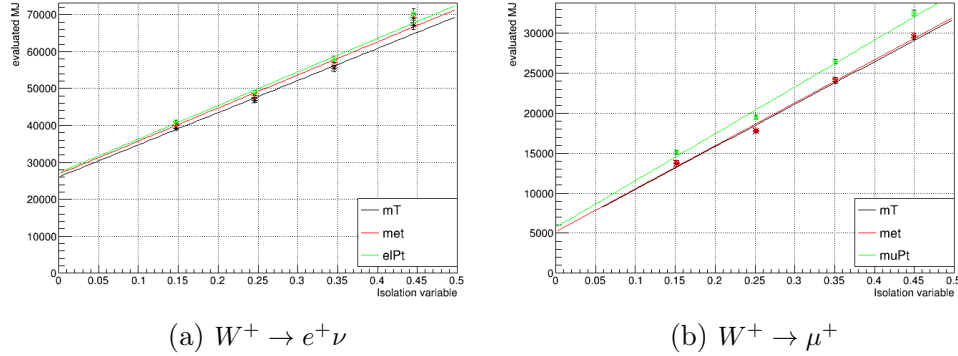


Figure 5.9 – The extrapolation of MJ yield to SR for $W^+ \rightarrow e^+ \nu$ channel (a) and $W^+ \rightarrow \mu^+ \nu$ channel (b) at 13 TeV. Each point is located at the average isolation of the corresponding anti-isolated slice of CR2.

where x is the target of the extrapolation and k is the slope in Figure 5.9. The average value of isolation of the MJ background in SR is not zero and can only be inferred from appropriate MC samples. For this work, the average value of isolation of the MJ background in SR is investigated with $b\bar{b}$ and $c\bar{c}$ samples in the muon and anti-muon channels as shown in Table 5.11, which favors choosing a baseline value of $x=0.025$ as the extrapolation target. A similar study can hardly be carried out in the electron channels because of the significant contribution from mis-identified hadrons. It is assumed that the isolation distributions for pions are not very different than those for the leptons from heavy-flavor quark decays. Hence, the baseline extrapolation target being checked in the muon channels applies to the electron channels as well.

While the baseline $x=0.025$ is considered the central value of the yield extrapolation, the difference of yield between $x=0.025$ and $x=0$ is taken as one source of MJ yield uncertainty due to the ambiguity of the extrapolation target.

It is observed that in the post-fit control distributions in FR there is a discrepancy between observation and total prediction, which can be possibly traced back to a mis-modelled dependence of the fraction fit on the jet activity. The effect of this mis-modelled dependence is studied by categorizing events as a function of recoil u_T , a quantity capable of representing the overall jet activity. The inclusive u_T distribution is therefore split into 5 intervals: $[0, 20, 40, 60, 80, 1000]$ GeV. Once the fraction fit is performed in each slice of u_T , the post-fit agreement is in general better than the outcome of inclusive u_T , which confirms the dependence of the MJ template on u_T . Following the baseline method MJ yield estimation, the fraction fits are carried out in the 5 slices of u_T using E_T^{miss} and m_T distributions. After that, the MJ yield is extrapolated to the SR in the corresponding u_T slice. The procedures are shown in Figure 5.10. Given that the cut $p_T^\ell > 25$ GeV is never relaxed in FR and CR1, a fraction fit based on p_T^ℓ does not benefit from the kinematic peak of MJ. Especially in the high u_T region where the Jacobian peak of W-boson production is no longer prominent either, the MJ distribution in p_T^ℓ becomes indistinguishable from the EW and top-quark related processes, causing convergence issue in the fraction fit. Therefore, p_T^ℓ distribution is not used in the estimation of MJ yield in slices of u_T . The estimation of MJ yield in slices of u_T offers not only a cross-check of the MJ shape in the SR (Figure 5.11), but also helps define a source of MJ yield uncertainty that is the difference between the nominal u_T inclusive yield and the sum of yields in u_T slices.

The correlation of u_T -dependence uncertainty across channels is understood by examining the response of the sum of yields towards the division of u_T spectrum in different channels. The MJ yield estimation is repeated upon a few alternative division of the entire u_T region, along with the inclusive one and the 5-slice split:

- 1 slice (u_T -inclusive): $[0, 1000]$ GeV for u_T .
- 2 slices: $[0, 50, 1000]$ for u_T .
- 4 slices: $[0, 25, 50, 1000]$ for u_T .
- 5 slices (the systematic variation): $[0, 20, 40, 60, 80, 1000]$ GeV for u_T .
- 10 slices: $[0, 10, 20, 30, 40, 50, 60, 70, 80, 90, 1000]$ GeV for u_T .

The estimation of overall yield tends to inflate when more u_T slices are used (more degrees of freedom are given). In Figures 5.12 and 5.13, with 5 slices of u_T , the sum of yields approaches the plateau, which justifies the choice of using the sum of yields in 5 u_T slices as the systematic variation of this uncertainty source. Also strong correlation is observed both between two lepton flavours and between the two center-of-mass energies. It is hence concluded that the u_T -dependence yield uncertainty originates from the method of assessment itself and should be fully correlated across both charges, both lepton flavours and both center-of-mass energies.

5.4.3 Extraction of shape

The shape of MJ background in SR is determined from four anti-isolated slices in CR2, each one with $\text{ptcone20}/\text{Min}(\text{pT}, 50 \text{ GeV})$ ranging within $[0.1, 0.2, 0.3, 0.4, 0.5]$. For any observable X such as p_T^ℓ , E_T^{miss} and m_T , the shape extrapolation uses a formula similar to equation 5.24, with MJ shapes in CR1 replaced by those in CR2.

$$H_{MJ}^{SR}[X] = H_{MJ}^{[0.1, 0.2]}[X] + \Delta^{CR2}[X] \quad (5.26)$$

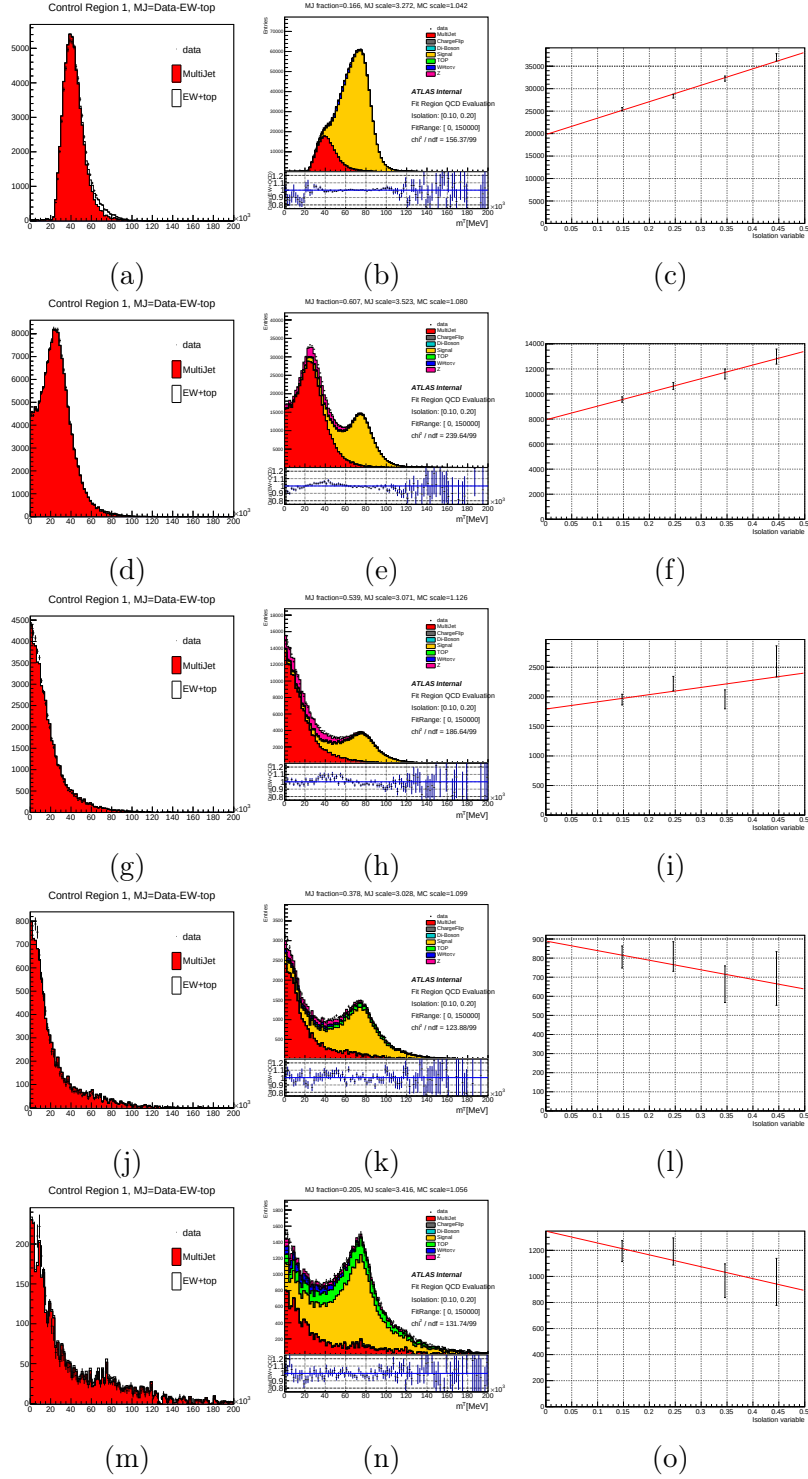


Figure 5.10 – MJ yield estimated in slices of u_T . Left: CR1 MJ shape template with shape correction implemented, middle: MJ fraction fit in FR, right: yield extrapolated to SR. (a)(b)(c): $u_T \in [0,20]$ GeV, (d)(e)(f): $u_T \in [20,40]$ GeV, (g)(h)(i): $u_T \in [40,60]$ GeV, (j)(k)(l): $u_T \in [60,80]$ GeV, (m)(n)(o): $u_T \in [80,1000]$ GeV.

2801 The shape extrapolation is shown in Figure 5.14. The statistical uncertainty in the
 2802 MJ distribution in the SR follows the error propagation of equation 5.26, where in
 2803 each anti-isolated slice of CR2 the corresponding statistical uncertainty comes from
 2804 the data-EW-top subtraction per bin. The statistical uncertainty is bin-by-bin uncor-

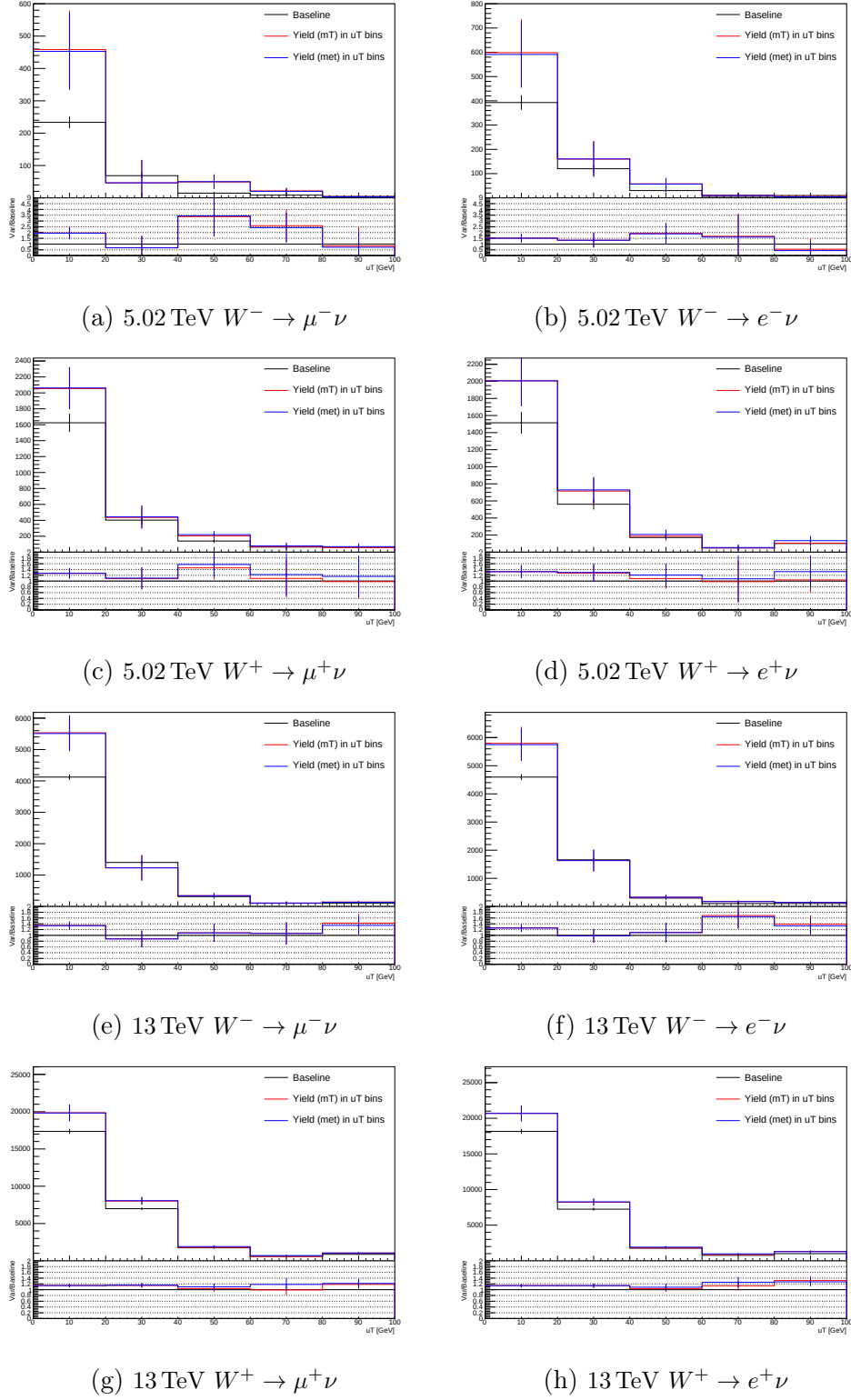


Figure 5.11 – MJ background in SR compared in broad bins of u_T between u_T -sliced estimations and the baseline shape extrapolation. The last bin includes overflow.

2805 related in this way. But in p_T^W analysis, a dedicated smoothing is applied to the region
 2806 $u_T > 100$ GeV in certain channels to improve the compatibility between electron and
 2807 muon unfolded p_T^W spectra [38]. In this case, the shape variations of the smoothing
 2808 parameters are correlated in the full u_T distribution, representing the MJ statistical

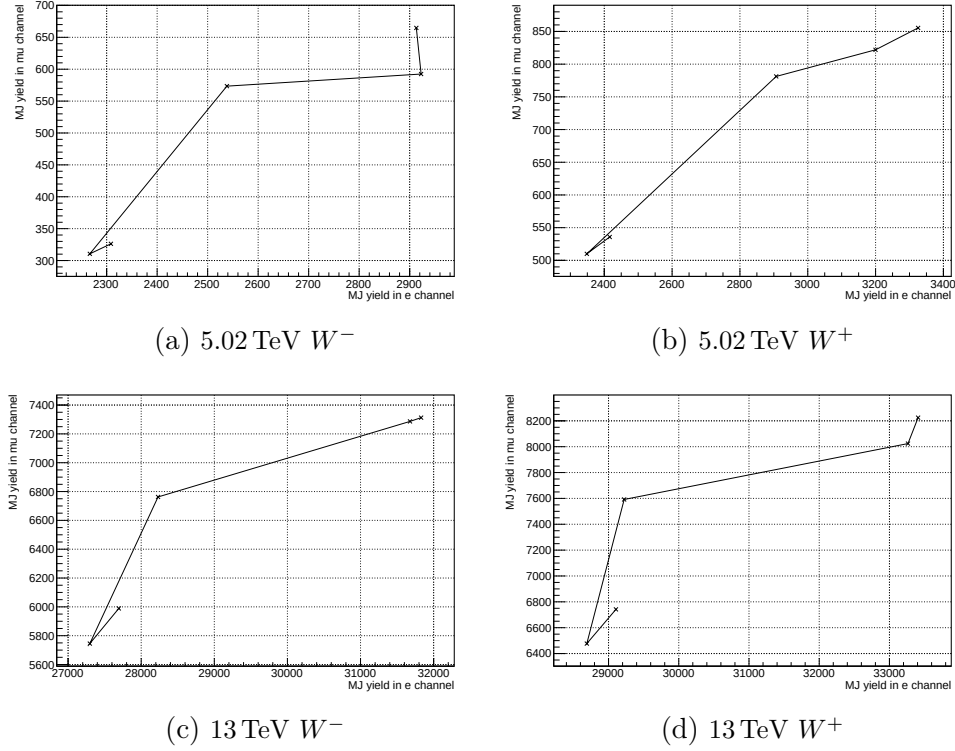


Figure 5.12 – The MJ yield in $W \rightarrow e\nu$ channel (horizontal) and $W \rightarrow \mu\nu$ channel (vertical) depending on the number of u_T slices (1, 2, 4, 5, 10, the overall trend goes from bottom left to top right) used for the estimation of yield. A strong correlation is found in all cases.

uncertainty for $u_T > 100$ GeV.

5.4.4 Sources of uncertainties

For the p_T^W analysis, the uncertainty in the MJ background estimation arises dominantly from three aspects: systematic uncertainties in the estimation of MJ yield in the SR, systematics in the determination of the shape in the SR and statistical uncertainty of the MJ distribution. Among them, the systematic uncertainties in yield and shape turn out to be correlated across the channels, while statistical uncertainty is uncorrelated between channels because of its statistical nature. Besides the above three major types of uncertainties, all sources of uncertainty have been investigated throughout the study of MJ background estimation and are listed below. These two sources of uncertainties are assumed to be uncorrelated between channels:

- The data and MC statistics: The statistical uncertainty from the limited data and MC samples is obtained from the error propagation of equation 5.26. The data and MC statistics is bin-by-bin uncorrelated in most cases. Only in the electron channels at 13 TeV, the statistical uncertainty in $u_T > 100$ GeV is parameterized into shape variations by smoothing and hence becomes bin-by-bin correlated.
- Uncertainty in the yield caused by the linear extrapolation as a function of isolation. The linear extrapolation of the yield accounts for the statistical uncertainty in the control regions and the uncertainty of the MJ fraction in FR obtained from the fraction fit. For the nominal estimation, the impact of

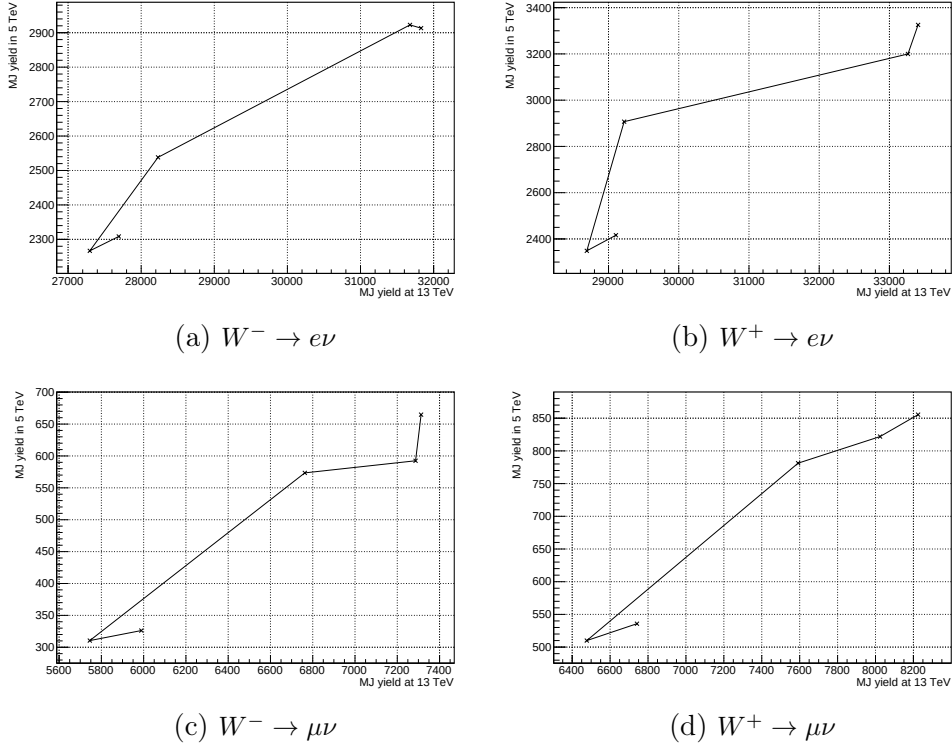


Figure 5.13 – The MJ yield in 13 (horizontal) and 5.02 TeV (vertical) center-of-mass energies depending on the number of u_T slices (1, 2, 4, 5, 10, the overall trend goes from bottom left to top right) used for the estimation of yield. A strong correlation is found in all cases.

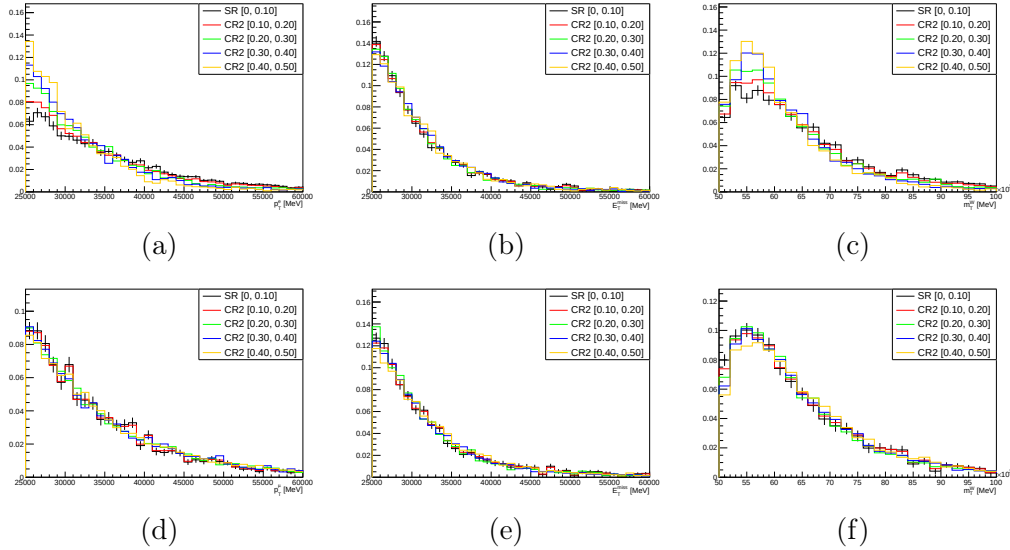


Figure 5.14 – Normalized MJ shapes in anti-isolated slices of CR2 and the shape corrected to SR for 13 TeV $W^+ \rightarrow e^+\nu$ (top) and $W^+ \rightarrow \mu^+\nu$ (bottom). (a)(d) p_T^ℓ , (b)(e) E_T^{miss} , (c)(f) m_T .

this uncertainty source turns out to be well covered by the leading systematic uncertainties of the yield. It is only included when presenting the uncertainty on the yield in slices of u_T .

The following uncertainties are correlated across channels, and across bins for a given

2834 kinematic distribution:

- 2835 — Uncertainty in the shape of MJ background in SR. This source of uncertainty is
- 2836 represented by the difference between the extrapolated shape in equation 5.26
- 2837 and the shape of the $[0.1, 0.2]$ anti-isolated slice in CR2. This source of uncer-
- 2838 tainty is propagated to the final result of MJ estimation.
- 2839 — The dependence of the fitted and extrapolated MJ yields on u_T . This uncer-
- 2840 tainty affects the total yield of MJ in the SR. It is taken to be the difference
- 2841 between the central value of yield based on u_T -inclusive estimation and the
- 2842 sum of yields from u_T -sliced estimations.
- 2843 — Uncertainties in the luminosity and in the cross-sections of the simulated MC
- 2844 samples. They appear in the subtraction of the EW and top processes from
- 2845 CR1 and CR2.
- 2846 — Uncertainty in the choice of the extrapolation target for the isolation scan.
- 2847 This is assessed by changing the isolation target from 0.025 to 0. As long as
- 2848 the shape correction is applied to the CR1 templates, the impact of changing
- 2849 the isolation target on the yield is consistent among p_T^ℓ , m_T and E_T^{miss} . So,
- 2850 the arithmetic average of the impact on the yield over the three kinematic
- 2851 distributions is propagated to the final result of MJ estimation.
- 2852 — The uncertainty in the yield due to the recoil isolation correction k_{iso} is es-
- 2853 timated by changing the isolation correction by 10%, a number way more
- 2854 conservative than the fit uncertainty presented in Table 5.3.

2855 For p_T^W analysis, the u_T -dependence of the fitted yield is important in all the W-
 2856 boson decay channels. In the muon channels, the uncertainty in the choice of the ex-
 2857 trapolation target is comparable to the u_T -dependence uncertainty. In terms of yield,
 2858 other sources are found to be sub-dominant in comparison. Since the u_T -dependence
 2859 uncertainty and the extrapolation target uncertainty are the two independent lead-
 2860 ing uncertainty sources, a quadratic sum of them is taken to be the total MJ yield
 2861 uncertainty. The shape uncertainty is propagated to the analysis chain as varia-
 2862 tions. The data and MC statistical uncertainty in MJ estimation is only considered
 2863 for $u_T > 100$ GeV, where it is no longer negligible compared with the data statistical
 2864 uncertainty of each u_T bin.

2865 5.4.5 Summary of results

2866 The final results of the MJ background estimation for p_T^W analysis are presented
 2867 in this section for the eight individual channels. The yield in the SR is obtained
 2868 by determining the yield in FR with the MJ template extracted from CR1, then
 2869 extrapolated to the SR via the ratio of numbers of MJ events in CR2 over CR1. The
 2870 yield is estimated using three kinematic distributions: p_T^ℓ , E_T^{miss} and m_T . The total
 2871 yield of MJ background based from the three distributions are summarized in Table 5.4
 2872 $\sim$ Table 5.7 for u_T inclusive estimation and u_T -sliced estimations. The central value
 2873 of yield is taken to be the arithmetic average of the u_T inclusive estimation based on
 2874 p_T^ℓ , E_T^{miss} and m_T , since the three numbers are compatible with each other. The MJ
 2875 yields between the two charges are within the tolerance of uncertainties as well. The
 2876 absolute and relative uncertainties of MJ yields are shown in Tables 5.8 and 5.9, where
 2877 the statistical uncertainty of the linear extrapolation of yield to the SR is included
 2878 but only for reference.

2879 As for the cross check of the shape extrapolation, the MJ shape extrapolated from

Distribution	u_T bin (GeV)	Electron Channel	Positron Channel	Δ
Lepton p_T	Inclusive	28041 ± 852	29362 ± 883	1321 ± 1227
m_T	[0,20]	19789 ± 615	20584 ± 641	795 ± 888
m_T	[20,40]	8023 ± 403	8176 ± 409	153 ± 574
m_T	[40,60]	1826 ± 168	1847 ± 173	21 ± 241
m_T	[60,80]	586 ± 92	919 ± 107	333 ± 141
m_T	[80,1000]	1086 ± 127	1415 ± 147	329 ± 194
m_T	Inclusive	26692 ± 805	28078 ± 834	1386 ± 1159
E_T^{miss}	[0,20]	19767 ± 614	20563 ± 640	796 ± 887
E_T^{miss}	[20,40]	8075 ± 406	8241 ± 412	166 ± 578
E_T^{miss}	[40,60]	1929 ± 178	1940 ± 182	11 ± 255
E_T^{miss}	[60,80]	696 ± 110	1011 ± 118	315 ± 161
E_T^{miss}	[80,1000]	1118 ± 134	1325 ± 139	207 ± 193
E_T^{miss}	Inclusive	27399 ± 826	28858 ± 857	1459 ± 1190

Table 5.4 – MJ background yield estimated from the lepton p_T , m_T and E_T^{miss} distributions in the 13 TeV $W \rightarrow e\nu$ channel. Statistical uncertainties are quoted.

CR2 to SR is merged to broad bins and compared to the yields estimated in slices of u_T in Figure 5.11. Good agreement is found in these comparisons.

A further check of the data-driven MJ background estimation is achieved in muon channels by comparing with heavy-flavour quark decay MC samples, since the MJ background in the muon channels is supposed to be dominated by $b\bar{b}$ and $c\bar{c}$ decays. More information of the samples are provided in Tables 5.10. Even though the production cross-sections are greater than those of the signal and other backgrounds, the calculation of these processes carry significant uncertainties from the dijet production cross-section, $b/c \rightarrow \mu$ decay fractions and fragmentation functions. After signal event selection and normalization to data luminosity, the number of $b\bar{b}$ and $c\bar{c}$ candidates are reported in Table 5.11, along with the average isolation in the SR. The shape comparison between the data-driven MJ background and the $b\bar{b}$ and $c\bar{c}$ samples is shown in Figures 5.15 and 5.16, where a certain degree of discrepancy is seen for the m_T distribution at 13 TeV.

The difference between the data-driven results and the MC samples can be attributed to a couple of reasons:

- The uncertainty in the generated cross-sections of MC samples.
- Background sources other than $b\bar{b}$ and $c\bar{c}$.
- Possible mis-modelling of $b\bar{b}$ and $c\bar{c}$ distributions bringing bias to cut efficiencies.

The estimated MJ yields in SR for p_T^W analysis are reminded in Table 5.12 for all channels, with central value and that of the yield uncertainty converted to a fraction of number of total candidates in data.

5.5 Total yields and control plots

The total expected yield passing the event selections of Section 5.2.1 is the sum of MC samples mentioned in Section 5.1 and the estimation of MJ background in

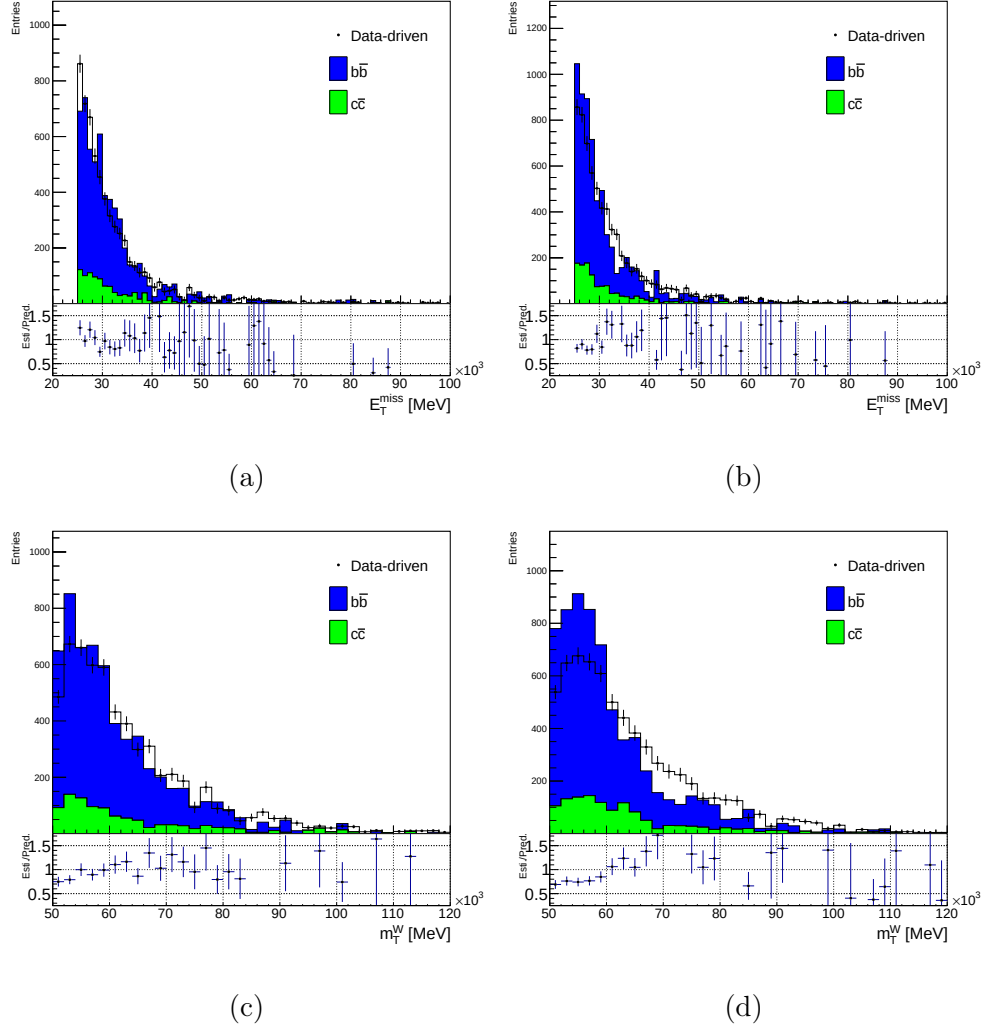


Figure 5.15 – Comparison of SR distributions of E_T^{miss} (top) as well as m_T (bottom) in muon channels using $b\bar{b} + c\bar{c}$ MC samples and data-driven estimation at 13 TeV. (a)(c) $W^- \rightarrow \mu^- \nu$, (b)(d) $W^+ \rightarrow \mu^+ \nu$. The MC samples are normalized to the yield of data-driven MJ background.

2906 Section 5.4. The observed yields are compared with the expected yields in Table 5.13
 2907 for each W decay signal channel.

2908 Control plots for 5.02 and 13 TeV low pile-up datasets are produced after apply-
 2909 ing all the selection in Table 5.1, as well as all the object corrections and physics
 2910 corrections described in Sections 5.2 and 5.3. The total predictions including MJ are
 2911 normalized to the integral of data to offer a better visualization of the agreement in
 2912 the shapes. In the ratio panel, the shallow grey band represent the total systematic
 2913 uncertainty that includes all the experimental uncertainties. Luminosity uncertainty
 2914 is excluded in the error band because of the normalization. No modelling uncertainty
 2915 is taken into account in the control plots. The brown band adds in quadrature the
 2916 signal MC statistical uncertainty on top of the total systematics, coming with a tiny
 2917 contribution everywhere over the spectra. The χ^2/ndf is calculated based on the full
 2918 covariance matrix of statistical uncertainties and the above systematics accounting
 2919 for bin-by-bin correlations.

2920 The hadronic recoil u_T in Figure 5.17 is the reco-level observable for the measure-

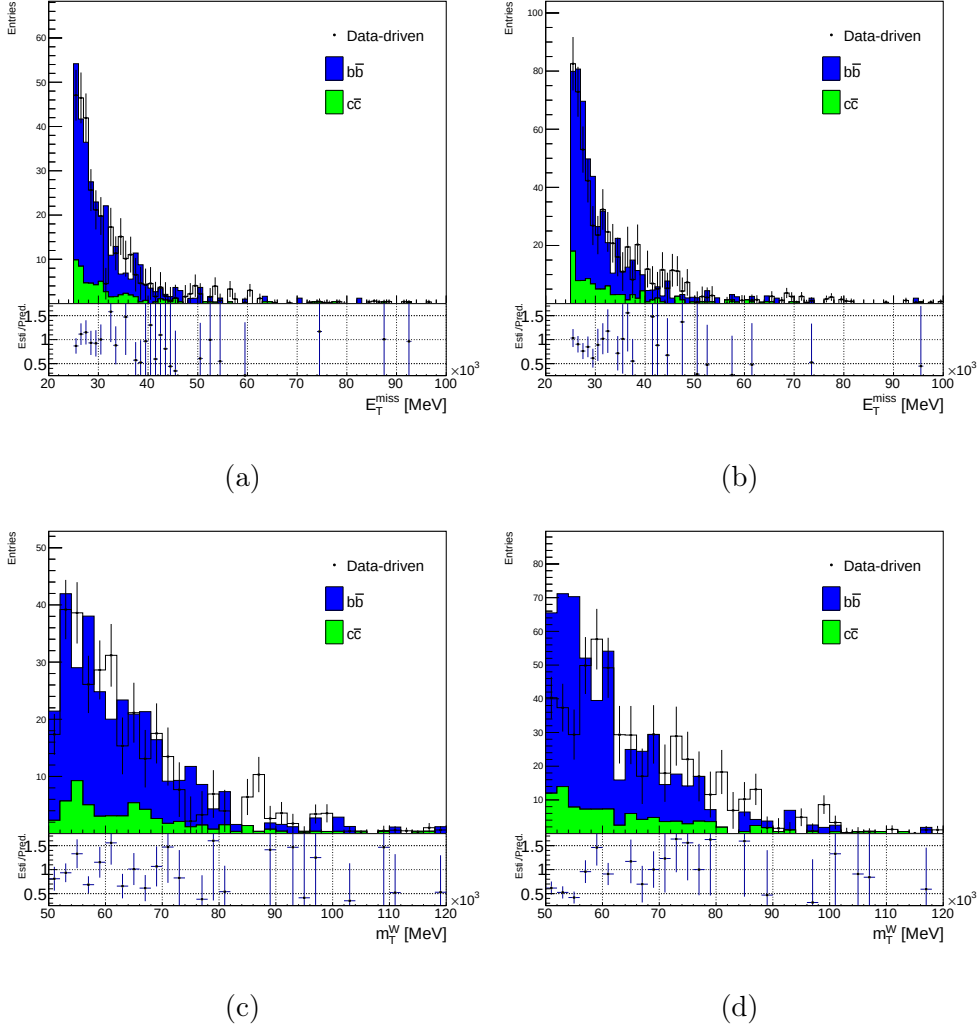


Figure 5.16 – Comparison of SR distributions of E_T^{miss} (top) as well as m_T (bottom) in muon channels using $b\bar{b} + c\bar{c}$ MC samples and data-driven estimation at 5.02 TeV. (a)(c) $W^- \rightarrow \mu^- \nu$, (b)(d) $W^+ \rightarrow \mu^+ \nu$. The MC samples are normalized to the yield of data-driven MJ background.

2921 ment of p_T^W , while p_T^ℓ (Figure 5.18), m_T (Figure 5.19) and E_T^{miss} (Figure 5.20) are
 2922 commonly used observables for the m_W measurements.

Distribution	u_T bin (GeV)	Muon Channel	Anti-muon Channel	Δ
Lepton p_T	Inclusive	6123 ± 331	6924 ± 358	801 ± 488
m_T	[0,20]	5529 ± 233	5740 ± 247	211 ± 340
m_T	[20,40]	1115 ± 154	1512 ± 162	397 ± 224
m_T	[40,60]	291 ± 67	299 ± 69	8 ± 96
m_T	[60,80]	96 ± 34	167 ± 36	71 ± 50
m_T	[80,1000]	136 ± 29	160 ± 30	24 ± 42
m_T	Inclusive	5617 ± 296	6247 ± 312	630 ± 430
E_T^{miss}	[0,20]	5504 ± 233	5685 ± 246	181 ± 339
E_T^{miss}	[20,40]	1108 ± 153	1507 ± 162	339 ± 222
E_T^{miss}	[40,60]	290 ± 67	299 ± 69	9 ± 96
E_T^{miss}	[60,80]	95 ± 34	162 ± 36	67 ± 50
E_T^{miss}	[80,1000]	125 ± 28	153 ± 30	28 ± 41
E_T^{miss}	Inclusive	5638 ± 298	6304 ± 316	666 ± 434

Table 5.5 – MJ background yield estimated from the lepton p_T , m_T and E_T^{miss} distributions in the 13 TeV $W \rightarrow \mu\nu$ channel. Statistical uncertainties are quoted.

Distribution	u_T Cut (GeV)	Electron Channel	Positron Channel	Δ
Lepton p_T	Inclusive	2248 ± 311	2405 ± 348	157 ± 467
m_T	[0,20]	2044 ± 242	2005 ± 263	-39 ± 357
m_T	[20,40]	422 ± 136	753 ± 151	331 ± 203
m_T	[40,60]	216 ± 42	188 ± 54	-28 ± 68
m_T	[60,80]	78 ± 37	61 ± 33	-17 ± 50
m_T	[80,1000]	58 ± 35	119 ± 47	61 ± 59
m_T	Inclusive	2181 ± 301	2269 ± 327	88 ± 444
E_T^{miss}	[0,20]	2060 ± 244	2010 ± 264	-50 ± 359
E_T^{miss}	[20,40]	428 ± 138	766 ± 153	338 ± 206
E_T^{miss}	[40,60]	232 ± 45	209 ± 60	-23 ± 75
E_T^{miss}	[60,80]	85 ± 41	70 ± 38	-15 ± 56
E_T^{miss}	[80,1000]	64 ± 38	129 ± 51	65 ± 64
E_T^{miss}	Inclusive	2258 ± 312	2356 ± 339	98 ± 461

Table 5.6 – MJ background yield estimated from the lepton p_T , m_T and E_T^{miss} distributions in the 5.02 TeV $W \rightarrow e\nu$ channel. Statistical uncertainties are quoted.

Distribution	u_T Cut (GeV)	Muon Channel	Anti-muon Channel	Δ
Lepton p_T	Inclusive	321 ± 123	586 ± 148	265 ± 192
m_T	[0,20]	480 ± 89	614 ± 103	134 ± 136
m_T	[20,40]	34 ± 54	165 ± 63	131 ± 83
m_T	[40,60]	46 ± 22	57 ± 26	11 ± 34
m_T	[60,80]	24 ± 10	7 ± 12	-17 ± 16
m_T	[80,1000]	5 ± 7	7 ± 10	2 ± 12
m_T	Inclusive	293 ± 112	528 ± 132	235 ± 173
E_T^{miss}	[0,20]	474 ± 88	606 ± 102	132 ± 135
E_T^{miss}	[20,40]	34 ± 54	165 ± 63	131 ± 83
E_T^{miss}	[40,60]	47 ± 23	57 ± 25	10 ± 34
E_T^{miss}	[60,80]	22 ± 10	6 ± 12	-16 ± 16
E_T^{miss}	[80,1000]	4 ± 6	6 ± 7	2 ± 9
E_T^{miss}	Inclusive	298 ± 114	532 ± 133	234 ± 175

Table 5.7 – MJ background yield estimated from the lepton p_T , m_T and E_T^{miss} distributions in the 5.02 TeV $W \rightarrow \mu\nu$ channel. Statistical uncertainties are quoted.

Channel	Electron	Positron	Muon	Anti-muon
MJ background yield	27377	28766	5792	6492
Stat. unc. (p_T^ℓ)	852 (3.1%)	883 (3.1%)	331 (5.7%)	358 (5.5%)
Stat. unc. (E_T^{miss})	826 (3.0%)	857 (3.0%)	298 (5.1%)	316 (4.9%)
Stat. unc. (m_T)	805 (2.9%)	834 (2.9%)	296 (5.1%)	312 (4.8%)
Stat. arithmetic average	828 (3.0%)	858 (3.0%)	308 (5.3%)	329 (5.1%)
Extrapo. unc. (p_T^ℓ)	2498 (9.1%)	2414 (8.4%)	1518 (26.2%)	1556 (24.0%)
Extrapo. unc. (E_T^{miss})	2440 (8.9%)	2370 (8.2%)	1404 (24.2%)	1426 (22.0%)
Extrapo. unc. (m_T)	2377 (8.7%)	2306 (8.0%)	1400 (24.2%)	1414 (21.8%)
Extrapo. arithmetic average	2438 (8.9%)	2364 (8.2%)	1440 (24.9%)	1465 (22.6%)
u_T -dependence unc.	4208 (15.4%)	4315 (15.0%)	1375 (23.7%)	1386 (21.3%)
Total unc.	4863 (17.8%)	4920 (17.1%)	1991 (34.4%)	2017 (31.1%)

Table 5.8 – The absolute and relative statistical uncertainty of MJ yield in the $W \rightarrow e\nu$ and $W \rightarrow \mu\nu$ channel at $\sqrt{s}=13$ TeV.

Channel	Electron	Positron	Muon	Anti-muon
MJ background yield	2229	2344	304	549
Stat. unc. (p_T^ℓ)	311 (14.0%)	348 (14.8%)	123 (40.5%)	148 (27.0%)
Stat. unc. (E_T^{miss})	312 (14.0%)	339 (14.5%)	114 (37.5%)	133 (24.2%)
Stat. unc. (m_T)	301 (13.5%)	327 (14.0%)	112 (36.8%)	132 (24.0%)
Stat. arithmetic average	308 (13.8%)	338 (14.4%)	116 (38.3%)	138 (25.1%)
Extrapo. unc. (p_T^ℓ)	293 (13.1%)	352 (15.0%)	222 (73.0%)	240 (43.7%)
Extrapo. unc. (E_T^{miss})	295 (13.2%)	345 (14.7%)	209 (68.8%)	221 (40.3%)
Extrapo. unc. (m_T)	284 (12.7%)	332 (14.2%)	206 (67.8%)	220 (40.1%)
Extrapo. arithmetic average	291 (13.1%)	343 (14.6%)	212 (69.7%)	227 (41.3%)
u_T -dependence unc.	642 (28.8%)	841 (35.9%)	268 (88.2%)	295 (53.7%)
Total unc.	705 (31.6%)	909 (38.8%)	342 (112.5%)	372 (67.8%)

Table 5.9 – The absolute and relative statistical uncertainty of MJ yield in the $W \rightarrow e\nu$ and $W \rightarrow \mu\nu$ channel at $\sqrt{s}=5.02$ TeV.

Process	$\sqrt{s}$	Data set	Generator	$\sigma \cdot \text{BR} \cdot \epsilon_{\text{filter}}$ [nb]
$b\bar{b} \rightarrow \mu$	13 TeV	361250	Pythia8B+A14+NNPDF23LO	186
$c\bar{c} \rightarrow \mu$	13 TeV	361251	Pythia8B+A14+NNPDF23LO	58
$b\bar{b} \rightarrow \mu$	5.02 TeV	361250	Pythia8B+A14+NNPDF23LO	46
$c\bar{c} \rightarrow \mu$	5.02 TeV	361251	Pythia8B+A14+NNPDF23LO	14

Table 5.10 – MC simulation of $b\bar{b}$ and $c\bar{c}$ decay to muon.

Channel	$b\bar{b}$		$c\bar{c}$		$b\bar{b} + c\bar{c}$		Data-driven Yield
	Yield	$< \text{ptcone20}/\text{Min}(\text{pT}, 50 \text{ GeV}) >$	Yield	$< \text{ptcone20}/\text{Min}(\text{pT}, 50 \text{ GeV}) >$	Yield	$< \text{ptcone20}/\text{Min}(\text{pT}, 50 \text{ GeV}) >$	
13 TeV $W^+ \rightarrow \mu^+\nu$	4365	0.0246 ± 0.0018	1067	0.0228 ± 0.0021	5433	0.0243 ± 0.0015	6492
13 TeV $W^- \rightarrow \mu^-\nu$	4436	0.0306 ± 0.0019	941	0.0243 ± 0.0023	5377	0.0295 ± 0.0016	5792
5.02 TeV $W^+ \rightarrow \mu^+\nu$	428	0.0282 ± 0.0025	111	0.0225 ± 0.0026	538	0.0270 ± 0.0020	549
5.02 TeV $W^- \rightarrow \mu^-\nu$	393	0.0292 ± 0.0026	100	0.0185 ± 0.0025	493	0.0270 ± 0.0021	304

Table 5.11 – The numbers of $b\bar{b}$ and $c\bar{c}$ events in $W^+ \rightarrow \mu^+\nu$ and $W^- \rightarrow \mu^-\nu$ signal channel at $\sqrt{s}=13$ TeV and 5.02 TeV, and the corresponding average value of $\text{ptcone20}/\text{Min}(\text{pT}, 50 \text{ GeV})$. The data-driven MJ background yield is included for comparison; all numbers are normalized to the integrated luminosity of the data.

Channel	13 TeV		5.02 TeV	
	Yield	Fraction	Yield	Fraction
$W^- \rightarrow e^-\nu$	27377 ± 4863	$(2.88 \pm 0.51)\%$	2229 ± 705	$(0.81 \pm 0.26)\%$
$W^+ \rightarrow e^+\nu$	28766 ± 4920	$(2.38 \pm 0.41)\%$	2344 ± 909	$(0.54 \pm 0.21)\%$
$W^- \rightarrow \mu^-\nu$	5792 ± 1991	$(0.60 \pm 0.21)\%$	304 ± 342	$(0.11 \pm 0.12)\%$
$W^+ \rightarrow \mu^+\nu$	6492 ± 2017	$(0.52 \pm 0.16)\%$	549 ± 372	$(0.12 \pm 0.08)\%$

Table 5.12 – Evaluation of MJ background yields at 13 TeV and 5.02 TeV.

Channel	Observed	Expected
5.02 TeV $W^+ \rightarrow e^+ \nu$	430662	430930 ± 933
5.02 TeV $W^+ \rightarrow \mu^+ \nu$	457223	457473 ± 429
5.02 TeV $W^- \rightarrow e^- \nu$	274375	276272 ± 726
5.02 TeV $W^- \rightarrow \mu^- \nu$	288026	289016 ± 384
13 TeV $W^+ \rightarrow e^+ \nu$	1207652	1192055 ± 4930
13 TeV $W^+ \rightarrow \mu^+ \nu$	1245755	1230252 ± 2042
13 TeV $W^- \rightarrow e^- \nu$	949297	947264 ± 4871
13 TeV $W^- \rightarrow \mu^- \nu$	964514	965341 ± 2012

Table 5.13 – Observed and expected event yields for each W decay channel after event selection. The uncertainty of the expectation is the quadratic sum of MC background statistics and multijet yield uncertainty.

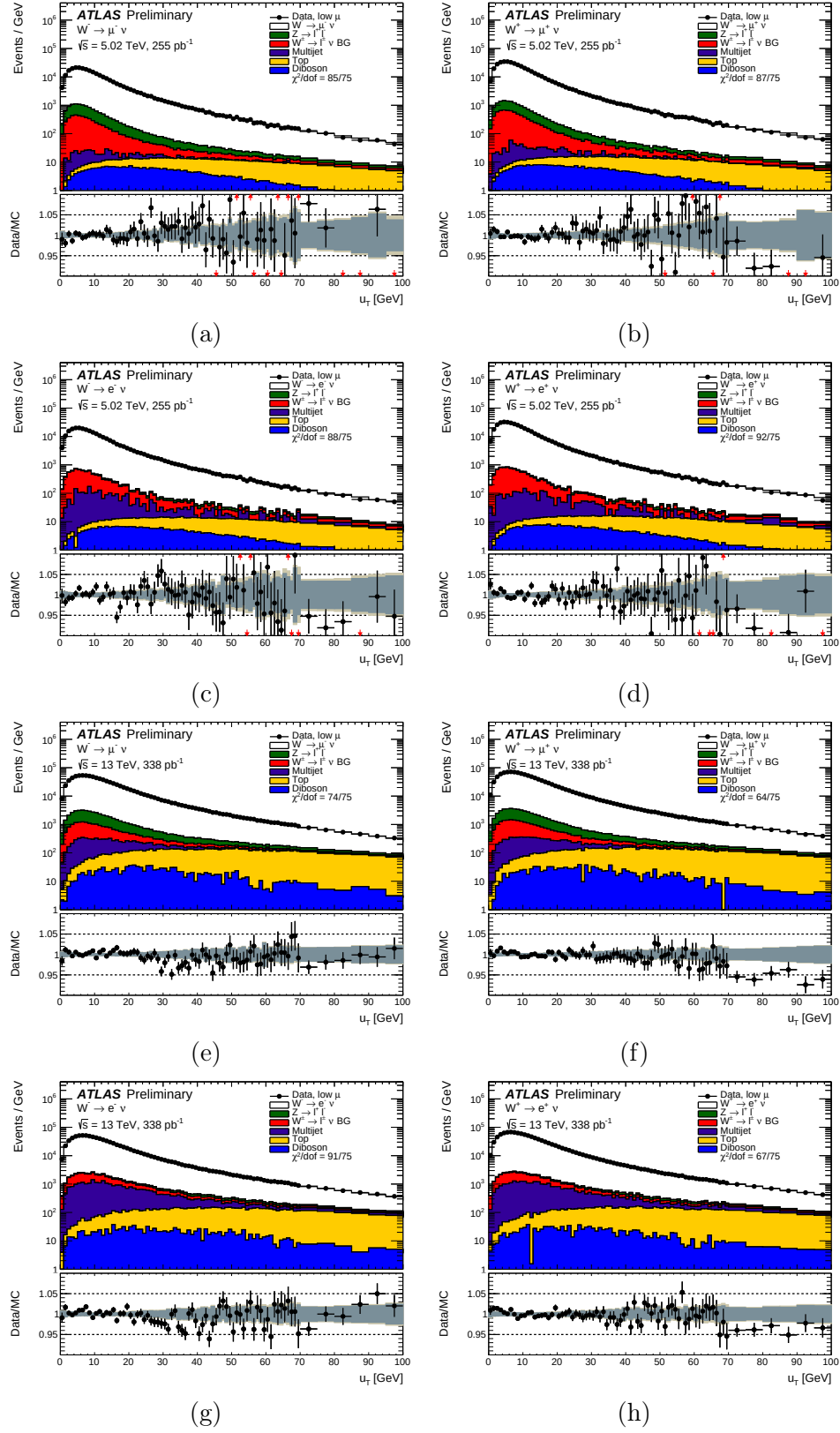


Figure 5.17 – u_T distribution in the muon and electron channels for the $\sqrt{s} = 5.02$ and 13 TeV datasets.

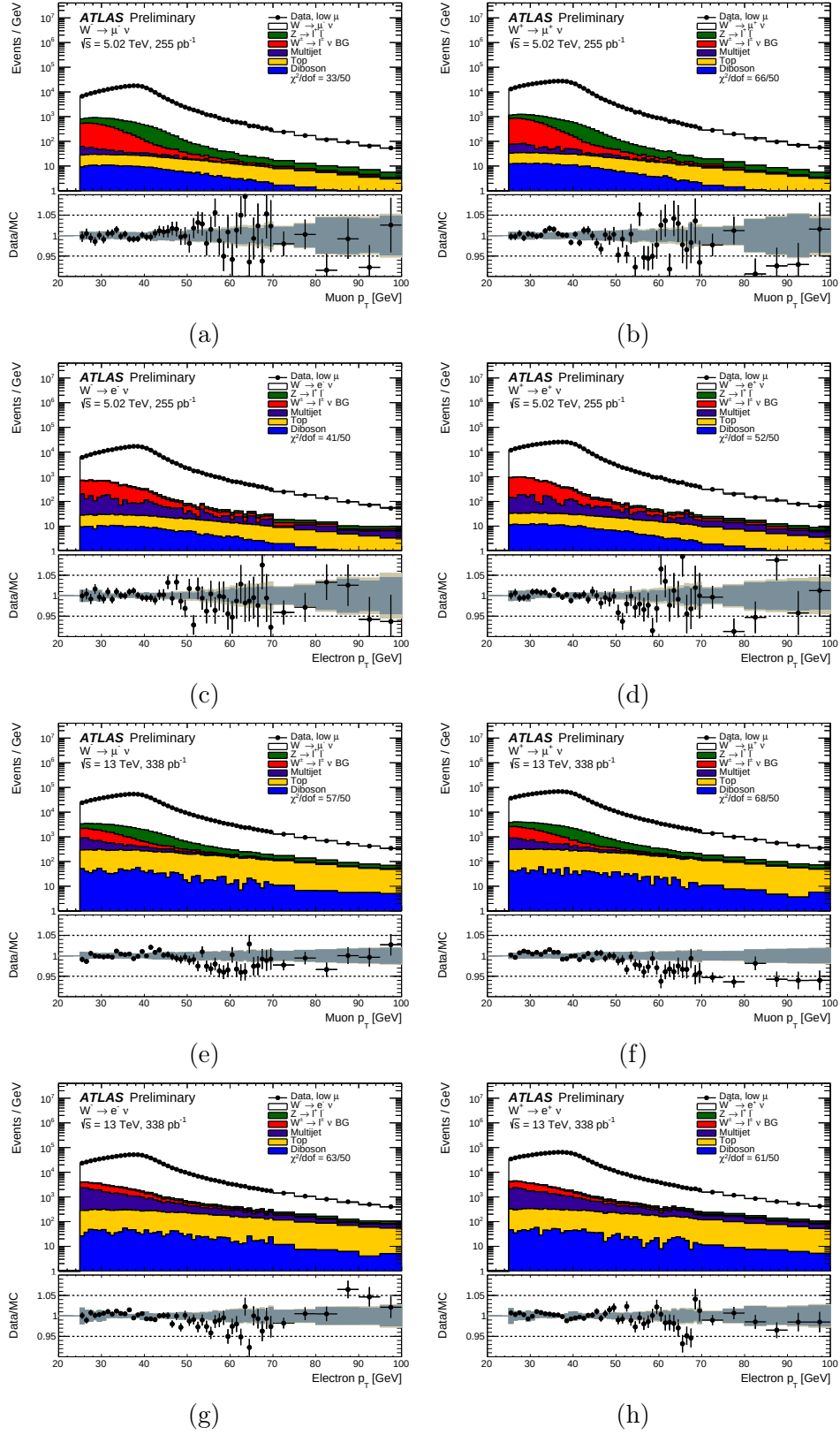


Figure 5.18 – Lepton p_T distribution in the muon and electron channels for the $\sqrt{s} = 5.02$ and 13 TeV datasets.

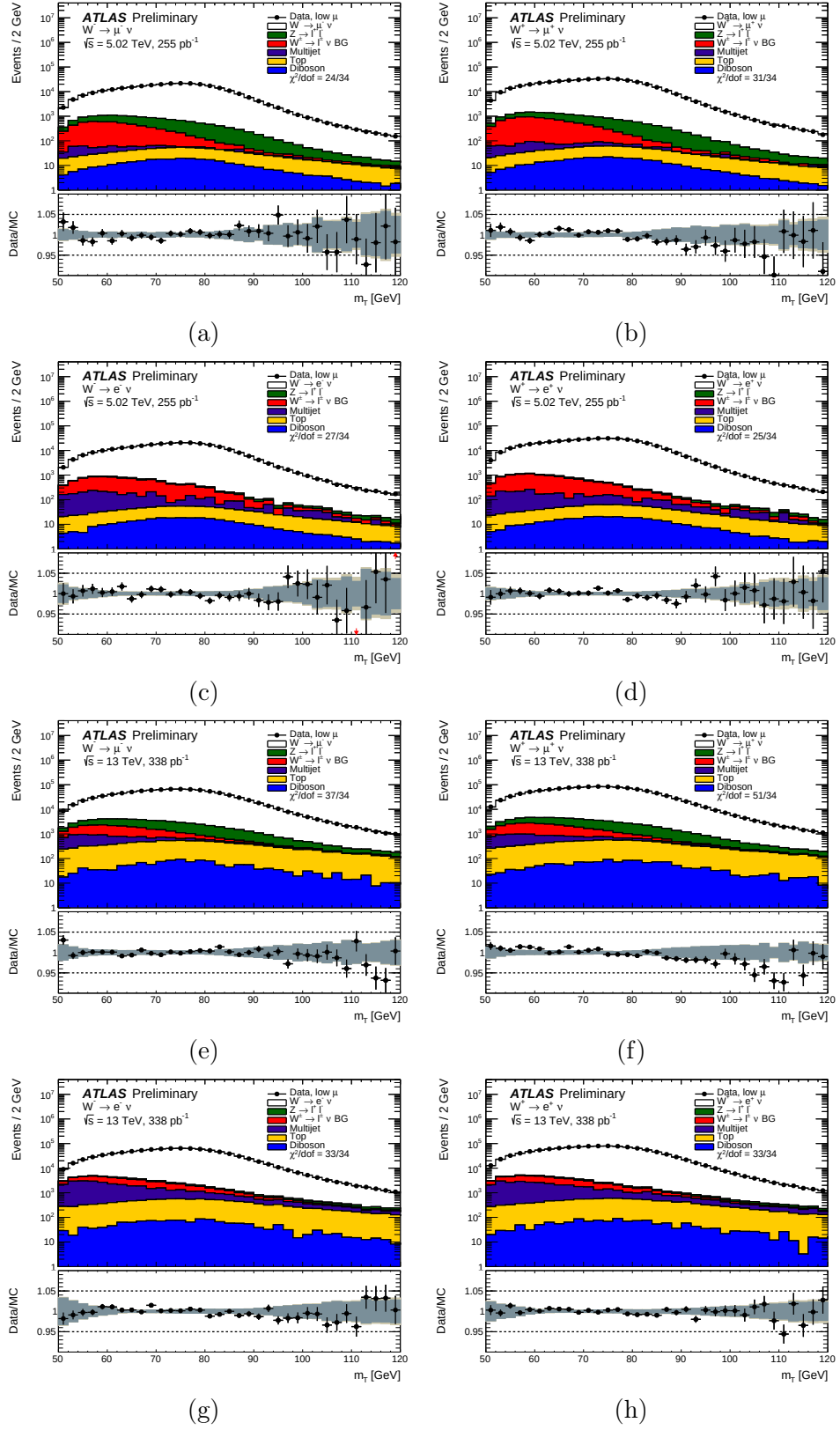


Figure 5.19 – m_T distribution in the muon and electron channels for the $\sqrt{s} = 5.02$ and 13 TeV datasets.

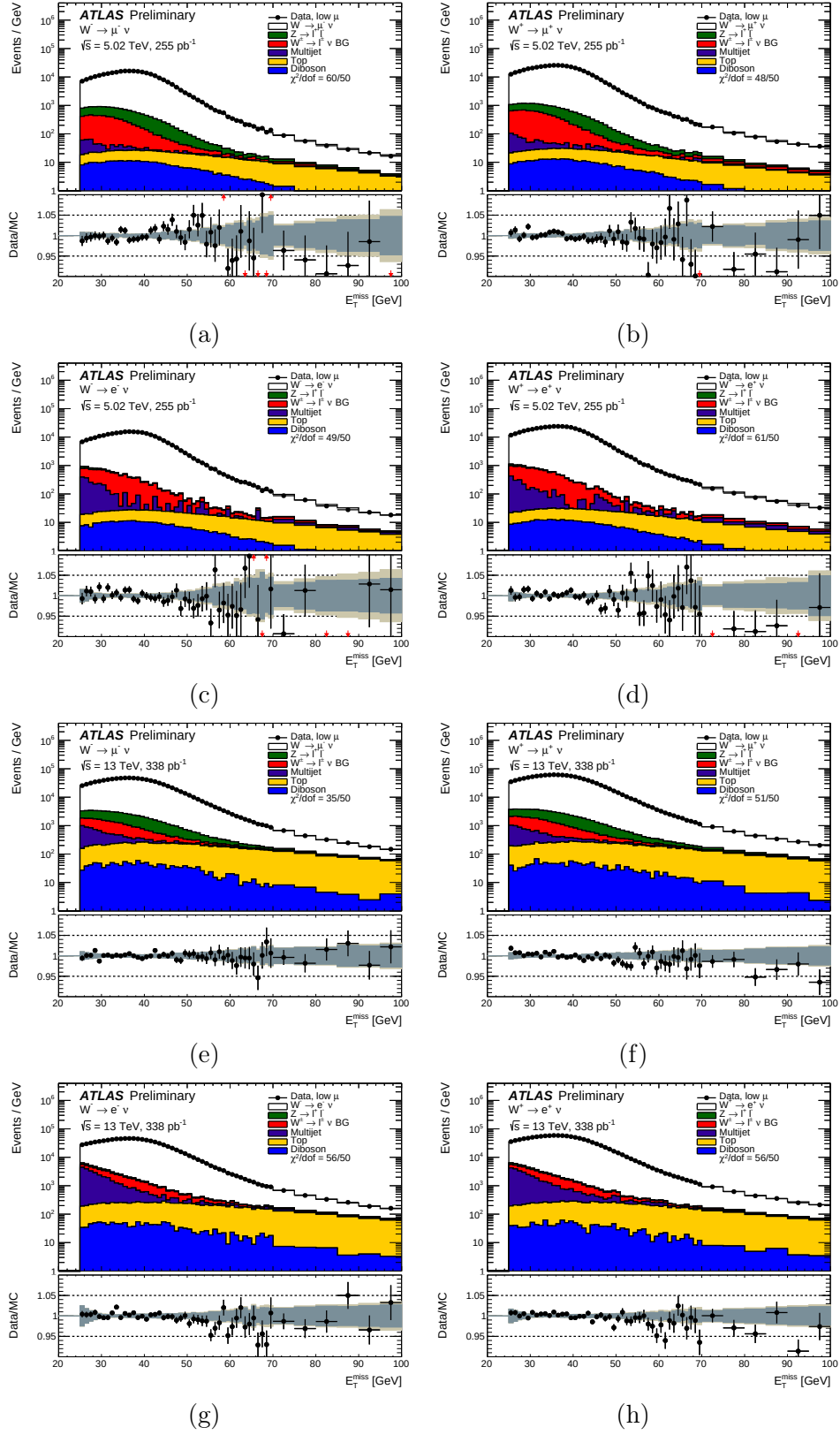


Figure 5.20 – E_T^{miss} distribution in the muon and electron channels for the $\sqrt{s} = 5.02$ and 13 TeV datasets.

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3085 Chapter 6

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6.1 Electron measurements

6.1.1 Electron reconstruction and selection

The offline electron and photon reconstruction in ATLAS in Run 2 is based on dynamic, variable-size clusters known as superclusters [1], so as to facilitate the energy recovery from bremsstrahlung photons or from photon-converted electrons. The dynamic clustering algorithm coordinates with dedicated calibration techniques to maintain the linear energy response and good stability as a function of pile-up.

An electron is defined to be an object comprising a cluster built from energy deposits in the calorimeter and a matched track or multiple tracks in the ID. An electron candidate is required to originate from the primary vertex, satisfying $|z_0| \cdot \sin(\theta) < 0.5 \text{ mm}$ and $|d_0|/\sigma_{d_0} < 5$. A converted photon is defined to be a cluster matching to a conversion vertex or several vertices, while an unconverted photon is a cluster matching to no electron track or conversion vertex.

The electron selection in the W-boson analyses consists of four requirements [2].

- Reconstruction: The object has to be reconstructed as an electron, judging from its energy deposit in the EM calorimeter and the track matching to the ID.
- Identification: The candidate for a prompt electron is required to pass a given working point (WP) of likelihood, where multiple quantities related to the calorimeter and the track are used to distinguish the prompt electron from light flavor jets, converted photons, or non-prompt electrons from the semi-leptonic decay of hadrons containing b or c quarks. The "MediumLLH" working point is chosen for the analyses, since it is found to render a good balance between efficiency and background rejection.
- Isolation: The reconstructed and identified electron needs to pass the isolation working point $\text{ptcone20}/\text{Min}(\text{pT}, 50 \text{ GeV}) < 0.1$.
- Trigger: The reconstructed and identified electron passing the isolation selection has to pass the trigger selection. For the single lepton trigger used in W-boson analyses, the trigger fires while at least one selected electron matches to the objects that fire the trigger.

The analyses use the combination of the above four requirements to select electrons, in order to improve the purity of electrons and reject backgrounds. The efficiencies of the selections may differ between data and simulation due to the complex data-taking and detecting environment, non-realistic simulation of tracks and clusters, mis-modelling of backgrounds and jet activities. This means that the electron distributions from the simulations have to be corrected for, in order to match instrumental inefficiencies in the data. The discrepancy between data and MC simulation are accounted for by scale factors (SF): the ratio of a given selection efficiency in the data over the predicted efficiency in MC. The total correction of selection efficiency applied to a single electron in W events is factorized into a product of the SFs for the four selection requirements:

$$W_{event}^{W \rightarrow e\nu} = SF_{reco} \cdot SF_{ID} \cdot SF_{Isolation} \cdot SF_{Trigger} \quad (6.1)$$

The efficiencies are measured with the tag-and-probe method using $Z \rightarrow ee$ events [3]. First, adequate kinematic cuts are applied to select $Z \rightarrow ee$ events, ensuring the purity and quality of the electron-positron pairs. Then, a tight selection is applied to one object in each pair to get the 'tag' while keeping the single electron trigger still matched, so that the trigger bias is removed. For events passing both the ordinary Z selection and the tight selection, the second object in the electron-positron pair is highly probable to be a prompt and isolated electron. And it is taken as 'probe' to measure the efficiency of a selection. The relevant systematic uncertainties are assessed by varying the requirements in the two levels of selection, which primarily evaluates the stability of the background subtraction.

For the low pile-up data taking, the *HLLT_e15_lhloose_nod0_L1EM12* unprescaled trigger is used for the electrons, which requires $p_T^e > 15$ GeV, loose ID criteria and no isolation selection. The trigger efficiency is defined to be the fraction of reconstructed, MediumLLH-identified and isolated electrons passing the trigger selection. Given that both ID and isolation are applied to both electrons, background becomes negligible for the measurement of electron trigger efficiency.

The central values and uncertainties of electron reconstruction SFs are obtained by extrapolating the standard high pile-up SFs to the low pile-up conditions, for both the 5.02 and 13 TeV datasets, as a function of p_T in bins of η . The identification SFs are measured in-situ using $Z \rightarrow ee$ events of the low pile-up data, but separately for the two center-of-mass energies. The isolation and trigger SFs are measured using $Z \rightarrow ee$ events in the low pile-up data, combining 5.02 and 13 TeV datasets.

6.1.2 Electron calibration

The extraction of electron energy scale and resolution uses the same method as the high pile-up data [4]. For $Z \rightarrow ee$ events, the energy scale factors α_i is applied to the data to correct the measured energy E^{data} to the scale of the simulation E_{MC} , such that

$$E^{data} = E^{MC} / (1 + \alpha_i) \quad (6.2)$$

The correction is implemented in η bins of the calorimeter cluster with the index i . The data are calibrated by multiplying the factor $1 + \alpha_i$ after the derivation of the α_i factors.

The relative resolution of the measured electron energy is a quadratic sum of three contributions:

$$\frac{\sigma(E)}{E} = \frac{a}{\sqrt{E}} \oplus \frac{b}{E} \oplus c \quad (6.3)$$

The first term a is the stochastic fluctuations related to the development of the electromagnetic shower in the calorimeter. Since the development of the shower is linear to the energy E , the stochastic fluctuations of the development of the shower is therefore proportional to $1/\sqrt{E}$. The second term with b describes the electronics and pile-up noise, followed by a constant term c accounting for the non-uniformity of the detector. An additional smearing c'_i is applied to the simulation in each calorimeter bin of η , making the width of the MC mass distribution match that of the data:

$$\left(\frac{\sigma(E)}{E}\right)_i^{data} = \left(\frac{\sigma(E)}{E}\right)_i^{MC} \oplus c'_i \quad (6.4)$$

24 calorimeter bins of η are used to extract α and c' correction factors in-situ for the low pile-up runs [3]. Uncertainties are assigned to electron calibration according to the statistical uncertainties in α and c'_i factors, along with the systematics related to imperfect modelling of the calorimeter.

6.2 Muon measurements

6.2.1 Muon reconstruction and selection

The muon reconstruction and identification are based on the information of three sub-detectors: MS, ID and calorimeter. A muon candidate can be classified according to a few types, depending on the available information, and consequently the algorithm used for the reconstruction. The types of muon candidates relevant to the low pile-up data are listed below [5]:

- MuonStandAlone (SA): Muon candidates are identified by track reconstruction only in the MS. The MS-only track is extrapolated to the IP after accounting for the energy-loss in the calorimeter. The information of the muon candidate is determined by this Muon-Extrapolated (ME) track.
- Combined (CB): The reconstruction of the tracks comes from a combined fit using both ID and MS information. An ID track and a ME track are needed for the final fit, where the ID track is extrapolated outward to the MS to find additional segments.
- CaloTagged (CT): Muon candidates are identified via an ID track with $p_T > 5$ GeV matching an energy deposit of a minimum ionization particle (MIP) in the calorimeter and being isolated.

For the muon reconstruction and identification of the Run 2 low pile-up datasets, the Medium working point (WP) is chosen as the muon selection criterion [6]. The tracks in both ID and MS are combined to fit the final reconstructed track of the muon (CB muon). The Medium WP in general provides very good acceptance, permille level fake rate, as well as small systematic uncertainties. For $0.1 < |\eta| < 2.5$, the Medium WP relies on CB muons with hits in no less than 2 MDT stations. For $|\eta| < 0.1$, CB muons with hits in 1 MDT station are allowed.

Given that the tracks are reconstructed separately in the ID and MS, the muon reconstruction efficiency can be factorized into a product of reconstruction efficiency in the ID, the reconstruction efficiency in the MS and the matching efficiency between the ID and MC measurements. Similar to the measurement of electron selection efficiency, a tag-and-probe method is also used for muons [7]. The efficiency to obtain a Medium muon is computed as:

$$\epsilon(\text{Medium}) \simeq \epsilon(\text{ID}|\text{ME}) \times \epsilon(\text{ME}|\text{CT}) \times \epsilon(\text{Medium}|\text{ID} \wedge \text{ME}) \quad (6.5)$$

Here $\epsilon(\text{ID}|\text{ME})$ stands for the efficiency of an ID track measured with respect to a ME probe. $\epsilon(\text{ME}|\text{CT})$ is the efficiency of an MS-only track measured with respect to a CT probe. The last term $\epsilon(\text{Medium}|\text{ID} \wedge \text{ME})$ is a measurement of the combination and the medium identification criteria with respect to a two-track-probe, in which the ID and ME tracks are required to stay $\Delta R < 0.05$ with each other.

The muon objects have to satisfy $|z_0| \cdot \sin(\theta) < 0.5$ mm and $|d_0|/\sigma_{d_0} < 3$ (track-to-vertex association requirements, or TTVA) to ensure they originate from the primary vertex (PV). Since the two requirements are not included in the muon identification criteria, their selection efficiency has to be studied independently. Muons passing the Loose identification criterion as well as $|\eta| < 2.5$ and $p_T > 10$ GeV are used as probes. Considering that the TTVA efficiency doesn't rely on the choice of identification WP, the corrections can be directly applied to the Medium WP. The matches are realized by adding the TTVA selection.

The muon trigger efficiency is determined by the tag-and-probe approach. The tag is an isolated and triggered muon, while the probe is an isolated muon. The invariant mass of the muon pair falls in the kinematic window near the Z resonance peak. The main unprescaled trigger is HLT_mu14, which comes with a significantly lower threshold and no isolation requirement in contrast to the single-muon triggers used for high pile-up Run 2 data taking. The trigger efficiency is studied in three dimensions: p_T , η and ϕ , so as to meet the requirements of the analyses in which the physics results are highly dependent on p_T but are inclusive in ϕ . To estimate the impact of systematic uncertainties on muon trigger efficiency, the selection requirements of the probe and the matching window to the trigger object are varied, including the kinematics of the Z mass window, the tightness of isolation, detector azimuthal symmetry and the inclusion of interaction point (IP) cuts.

Muon isolation in the low pile-up W analyses uses the track based isolation $\text{pt-cone20}/\text{Min}(p_T, 50 \text{ GeV}) < 0.1$. The isolation selection efficiencies are studied with the tag-and-probe method in $Z \rightarrow \mu\mu$ candidates in both data and MC simulation. Given that the background level is estimated to be negligible, no background subtraction is done in the determination of the isolation efficiency in both the data and the MC. Muon isolation SFs, measured as a function of the probe muon η and p_T , are used in the analyses to correct the MC events to match the isolation selection efficiency in the data. The muon isolation SFs are derived using combined 5 and 13 TeV low pile-up datasets. The per mille level isolation SF uncertainties are dominated by statistical uncertainties. Additional systematic sources include the range of pair invariant mass required in the selection, the probe quality, the tag-probe separation, the jet separation and the tag isolation.

6.2.2 Muon calibration

Muon momentum calibration involves the corrections to the momentum scale and resolution, as well as the sagitta bias. The corrections to momentum scale and resolution are applied to the muons in the simulation, bringing the scales and resolutions closer to the data, though at the expense of optimal performance.

The reconstructed muon tracks in the inner detector (ID tracks) and those in the muon spectrometer (ME tracks) receive separate corrections. The corrections are propagated from ID tracks and ME tracks to combined muon tracks (CB tracks) thanks to a reweighting of the two contributions. The muon calibration in terms of momentum scale and resolution is obtained from the high pile-up datasets [5] and cross-checked with low pile-up datasets using J/ψ and Z-boson samples. Generally speaking, after the implementation of the standard corrections, the agreement of both the scales and resolutions in MC with data is improved. The data-MC discrepancy for scale is at permille level most of the time. The residual difference can be mostly covered by systematic uncertainties related to calibration corrections. This implies that at least for p_T^W analysis, which is not very sensitive to muon momentum calibration, the high pile-up calibration can be applied to the low pile-up data. An improved muon calibration might be needed for the W mass analysis, using dedicated in-situ derivation of muon momentum corrections for the low pile-up data.

The other part of muon calibration is the correction for charge-dependent momentum biases (sagitta bias). The misalignments within the ID, within the MS or between the ID and the MS can introduce sagitta bias to the muon momentum reconstruction. The effect of sagitta bias on the reconstructed muon transverse momentum is parameterized as:

$$p_T^{meas} = \frac{p_T^{reco}}{1 + q \times \delta_{sagitta} \times p_T^{reco}} \quad (6.6)$$

p_T^{meas} is the measured (and biased) transverse momentum. p_T^{reco} represents the unbiased reconstructed transverse momentum. The sagitta bias, quantified by $\delta_{sagitta}$, is studied in bins of lepton η while ignoring the distribution in ϕ to suppress the statistical fluctuations. The sagitta bias is evaluated using a combination of three methods [8]:

- ID momentum measurement with electrons. From $Z \rightarrow ee$ or $W \rightarrow e\nu$ events, the charge bias of the electron track is determined by:

$$\delta_{sagitta} = \frac{\langle E/p_{track} \rangle^+ - \langle E/p_{track} \rangle^-}{2 \langle p_T \rangle} \quad (6.7)$$

Here p_{track} is the track momentum recorded by the inner detector and E is the calibrated electron energy measured by the EM calorimeter. $p_T = E \sin \theta$ is the transverse momentum of the electron derived from the charge-blind EM calorimeter measurement and the track direction. All the quantities are measured in η bins on average.

- $p_T(\mu)$ method. The muon momentum bias is determined by comparing the p_T distribution of muons with that of the anti-muons in $Z \rightarrow \mu\mu$ events. The difference caused by electroweak effects is known from the simulation and turns out to be marginal. Therefore, the major difference between the p_T spectra is expected to arise from the sagitta bias.

— Z-mass method. The reconstructed Z-boson invariant mass peak from $Z \rightarrow \mu\mu$ events may have an offset from the expected value, which can be assigned to the sum of the sagitta biases of the two muon tracks. The measured biases are corrected iteratively, until the sagitta bias of each individual muon track converges and Z mass is independent of η and ϕ bins.

It is observed that the sagitta bias in the low pile-up data is around 10 times larger than in the MC simulation. A dedicated sagitta bias correction is derived by using the results differential in η from the Z-mass method corrected by a global offset $< \delta_{sagitta} >$. The overall offset of $< \delta_{sagitta} >$ is evaluated by adjusting its value until the ratio of positive and negative muon p_T spectra in $Z \rightarrow \mu\mu$ events becomes flat [6]. The data-MC difference is applied to the data to achieve a good agreement in data/MC comparison.

6.3 Hadronic recoil measurements

In the decay of the $W \rightarrow \ell\nu$ process, where the neutrino escapes the detection, the transverse momentum of the neutrino can only be inferred from the transverse momenta of all other particles produced in a p-p collision. The transverse momentum of the neutrino is not only required in the event selection for W-boson measurements, but also enters the reconstruction of the W transverse mass, a quantity of particular importance in the W mass analysis. In order to determine neutrino kinematics from W decay in high-precision, the ATLAS analyses use the general methodology known as the hadronic recoil, according to which the vectorial and scalar sums of all particles from the initial state gluon and quark radiation are measured. In this way, the transverse momentum of the W-boson can be directly measured by $\vec{u}_T + \vec{p}_T^W = 0$. One step further, by combining the measurement of the charged lepton with the hadronic recoil, the neutrino transverse momentum can be inferred by making use of the momentum balance in the transverse plane.

The hadronic recoil has been implemented successfully in previous ATLAS analyses, for instance, the measurement of W-boson transverse momentum distribution at $\sqrt{s} = 7$ TeV [9] and the measurement of W-boson mass at $\sqrt{s} = 7$ TeV [10]. In contrast, the ATLAS standard E_T^{miss} algorithm used for physics in high pile-up environment with track-based soft term (TST) is not suitable for the W analyses at low p_T^V , mostly because of the main E_T^{miss} algorithm on reconstructed jets with a p_T threshold of about 20 GeV and the very limited response of the tracking term.

6.3.1 PFOs and hadronic recoil reconstruction

The hadronic recoil is calculated using an algorithm based on Particle Flow Objects (PFOs) [11], which helps the algorithm stay fully transparent regarding its application to $W \rightarrow \ell\nu$ events and the calibration with $Z \rightarrow \ell\ell$ events.

In the PFO algorithm, first the inner detector (ID) tracks are selected with the tight criteria, requiring at least nine hits in the silicon detectors and tracks not missing hits in the pixel detector when those hits are expected from the trajectory under the given operating conditions. The best recoil resolution is obtained by using the constituent PFO in the largest possible phase space, such that no extra cut is applied apart from the $p_T > 0.5$ GeV requirement of the track reconstruction. The maximal

pseudo-rapidity corresponds to the coverage of the ATLAS FCAL up to $|\eta| = 4.9$ [12]. Tracks matching to either an electron or muon with $p_T > 10 \text{ GeV}$ identified by the medium criteria are removed. The rest of the tracks are then matched to a single calorimeter cluster constructed with the usual topo-cluster 4/2/0 algorithm calibrated to the electromagnetic (EM) scale [13]: Based on the cluster position and the track momentum, the expected energy deposit of the track of the particle in the calorimeter is computed. Topo-clusters of calorimeter cells start from cells with absolute energy measurements exceeding the expected energy by four times of the standard deviation of the expected noise (both electronics noise and contribution from the pile-up). After that, the adjacent cells with absolute energy two times above the noise are added iteratively to the topo-clusters, followed by a final addition of all the neighbouring cells. In case a single particle deposits energy in multiple clusters, the additional clusters within the geometric cuts are added to the system. The recovery of split showers is handled by a specific procedure. In the end, the difference of the expected calorimeter energy deposited by the track and the sum of the associated clusters is compared with the expected calorimeter noise. If the difference is compatible with zero within 1.5σ , the clusters are removed, otherwise the remnant topo-cluster is retained.

The PFOs are divided into two separate types. The neutral PFOs (nPFOs) correspond to topo-clusters not associated to ID tracks of charged particles. The charged PFOs (cPFOs) correspond to the selected ID tracks. The cPFOs originating from vertices other than the primary vertex are removed by requiring $|z_0 \times \sin(\theta)| < 2 \text{ mm}$ from the primary vertex.

The reconstruction of the hadronic recoil needs to be transparent between W and Z, or between electron and muon decay channels for a precise calibration and well-understood correlations. Although both electrons and muons lose energy in the calorimeter, the nature of energy loss as well as the profile of the energy deposit is very different: while electrons develop showering in the material, muons lose energy mainly because of isolation. For this reason, the lepton-related difference between electron and muon have to be removed in the reconstruction algorithm. Empirically in ATLAS, removing a cone of radius $\Delta R < 0.2$ around the lepton axis is sufficient to eliminate those effects, while bringing a minor impact on the rest of the event reconstruction. In practice, this is achieved by excluding the PFOs inside the cone around the identified leptons from the calculation of the recoil. Subsequently, the bias caused by removing energy deposits from underlying event and pile-up in that cone is compensated for by looking for another cone of the same size and same $|\eta|$ but at a random ϕ position. The chosen replacement cone also needs to be $\Delta R > 0.4$ away from the recoil prior to this correction or from any lepton. After that, the activities in replacement cone are rotated in ϕ to the removed lepton direction and added vectorially to the recoil. The leptons considered in the removal procedure are required to be well measured and pass the criteria of $p_T > 10 \text{ GeV}$, medium ID and $|\eta| < 2.4$ for muons, or LooseLH+B-layer hit and $|\eta| < 2.47$ excluding the cracks for electrons.

The hadronic recoil in the transverse plane $\vec{u}_T$ is the vectorial sum of all cPFOs and nPFOs:

$$\vec{u}_T = \sum_i \vec{E}_{T,i}^{cPFO} + \sum_i \vec{E}_{T,i}^{nPFO} \quad (6.8)$$

The quantity u_T is the magnitude of the vector $\vec{u}_T$.

6.3.2 Variables for the study of hadronic recoil

The characterization and calibration of the hadronic recoil are based on variables independent of the azimuthal ϕ angles of the event activities but directly correlated with the vector boson kinematics. In Z events, $\vec{p}_T^{\ell\ell}$ is used to represent the transverse momentum of Z-boson. Two scalar quantities are built using the projection of $\vec{u}_T$:

The component of the hadronic recoil perpendicular to the axis of Z-boson is defined as:

$$u_{\perp} = \frac{|\vec{p}_T^{\ell\ell} \times \vec{u}_T|}{p_T^{\ell\ell}} \quad (6.9)$$

The component of the hadronic recoil parallel to the axis of Z-boson is:

$$u_{\parallel} = \frac{\vec{p}_T^{\ell\ell} \cdot \vec{u}_T}{p_T^{\ell\ell}} \quad (6.10)$$

$u_{\perp}$ is zero on average, and its spread reflects the hadronic recoil resolution in the perpendicular direction. Ideally, $u_{\parallel}$ balances the Z-boson transverse momentum according to $u_{\parallel} = -p_T^{\ell\ell}$. However, this balance can hardly be fulfilled because of effects such as particles escaping detection, energy losses in dead material and the non-compensating nature of the calorimeter. Therefore, a non-zero bias of the momentum balance in the transverse plane almost always exists and is introduced to the study of the hadronic recoil, as a more convenient notation in many contexts:

$$b = u_{\parallel} + p_T^{\ell\ell} \quad (6.11)$$

The average of b is usually found to be positive. And the spread of b distribution indicates the hadronic recoil resolution in the parallel direction.

In W events, the useful projections of the hadronic recoil rely on the direction of charged lepton because of the neutrino in the decay product:

$$\begin{aligned} u_{\perp}^{\ell} &= \frac{|\vec{p}_T^{\ell} \times \vec{u}_T|}{p_T^{\ell}} \\ u_{\parallel}^{\ell} &= \frac{\vec{p}_T^{\ell} \cdot \vec{u}_T}{p_T^{\ell}} \end{aligned} \quad (6.12)$$

In this way $u_{\perp}^{\ell}$ and $u_{\parallel}^{\ell}$ provides the information of u_T in the direction related to event kinematics, similar to the quantities $u_{\perp}$, $u_{\parallel}$ and b in Z events.

6.3.3 Hadronic recoil calibration and correction

The calibration procedure for the hadronic recoil is adapted from the method developed in the analysis of the 2011 7 TeV data [14]. The calibration uses Z-boson events as a standard candle, since the transverse kinematics of the final state lepton pair can be fully reconstructed and precisely measured, of which the resolution is one order of magnitude better than the hadronic recoil. Three steps are carried out in a sequence:

- Reweighting of MC to improve the modelling of $\sum \vec{E}_T$ (defined in Equation 5.4).
- The azimuthal bias is corrected for.

— The resolution and the response of the two projections of $\vec{u}_T$ in MC are corrected to match the data.

$\sum \bar{E}_T$ reweighting

Despite the similarities between W and Z events that help the transfer of calibration at good accuracy, differences in the energy scale and the partons involved in the hard scattering lead to both different underlying events and the boson- p_T spectra. Therefore, the transfer of the calibration correction from Z to W relies on a calibration obtained as a function of $\sum \bar{E}_T$ and p_T^V . Both the modelling of $\sum \bar{E}_T$ and the correlation between $\sum \bar{E}_T$ and p_T^V have to be corrected for in the MC, such that in each bin of p_T^V the MC distribution of $\sum \bar{E}_T$ matches that of the data.

Three steps of event reweighting are implemented to obtain a correct modelling of event activity $\sum \bar{E}_T$ and its correlation with p_T^V for the MC samples. A ($\sum \bar{E}_T, p_T^{\ell\ell}$) 2D reweighting is derived from the ratio of the data to Z simulation, where electron and muon channels are added together:

$$w_{2D}^Z(\sum \bar{E}_T, p_T^{\ell\ell}) = \frac{h_{2D}^{data,Z}(\sum \bar{E}_T, p_T^{\ell\ell})}{h_{2D}^{MC,Z}(\sum \bar{E}_T, p_T^{\ell\ell})} \quad (6.13)$$

Here h_{2D} stands for the normalized 2D distribution of ($\sum \bar{E}_T, p_T^{\ell\ell}$).

The second step of reweighting is applied to the W processes, where the residual of a few percent for $\sum \bar{E}_T$ is observed between data and MC, if the above Z-based procedure is implemented without further correction (i.e. simply taking $w_{2D}^Z(\sum \bar{E}_T, p_T^{true,V})$). After the 2D reweighting, an improved modelling of $\sum \bar{E}_T$ in W events is obtained by reweighting $\sum \bar{E}_T$ in slices of u_T :

$$w_{j,sliced}^{W^\pm}(\sum \bar{E}_T) = \frac{h_j^{data,W^\pm}(\sum \bar{E}_T)}{h_j^{MC,W^\pm,Z2D}(\sum \bar{E}_T)} \quad (6.14)$$

h_j denotes the normalized $\sum \bar{E}_T$ for a given slice j of u_T distribution in the data (the sum of signal and background predictions) for the numerator (denominator) after the W event selection. The index 'Z2D' suggests that the histogram in the denominator has gone through the 2D reweighting.

Since the above two steps of reweightings together modify the spectrum of boson p_T in an unintended way, a third reweighting is necessary to recover the initial spectrum of $p_T^{V,true}$ in the full phase-space, so that the critical modelling of boson p_T in the MC is not affected:

$$w_{1D}^{W^\pm}(p_T^{true,V}) = \frac{h^{MC,W^\pm,mod}(p_T^{true,V})}{h^{MC,W^\pm,orig}(p_T^{true,V})} \quad (6.15)$$

The 1D histogram h is the modified (numerator) or original (denominator) $p_T^{true,V}$ of simulated W events without selection cut. The final weight applied to W events in the simulation with the value of u_T in the corresponding slice j is given by $w_{2D}^Z(\sum \bar{E}_T, p_T^{\ell\ell}) \times w_{j,sliced}^{W^\pm}(\sum \bar{E}_T) \times w_{1D}^{W^\pm}(p_T^{true,V})$.

Correction of recoil direction

Due to effects such as beam crossing angle, calorimeter non-uniformity and beam displacement, the direction of the recoil is not perfectly modelled in the simulation. Empirically, the correction is carried out by an additive offset as the difference between the mean values of the x, y components ($\langle u_x \rangle$, $\langle u_y \rangle$) of the recoil in data and

simulation. A linear dependence on $\sum \bar{E}_T$ is assigned to the correction:

$$\begin{aligned} u_x^{MC,corr} &= u_x^{MC} + [(\langle u_x^{data} \rangle - \langle u_x^{MC} \rangle)(\sum \bar{E}_T)] \\ u_y^{MC,corr} &= u_y^{MC} + [(\langle u_y^{data} \rangle - \langle u_y^{MC} \rangle)(\sum \bar{E}_T)] \end{aligned} \quad (6.16)$$

The remaining discrepancy between data and MC for the direction of the recoil is below the data statistics of the Z events and is therefore considered as a fluctuation. The effect of the correction of the recoil direction has been observed having a negligible impact on u_T , confirming a weak correlation between the direction and the magnitude of the recoil.

Resolution and response corrections

The recoil resolution is derived by comparing the width of $u_\perp$ as well as the mean and width of the bias b in the data and in the simulation of Z events. The differences in both perpendicular and parallel directions are corrected individually.

The data-to-MC resolution correction for $u_\perp$ derived as the ratio $r(\sum \bar{E}_T, p_T^{\ell\ell}) = \sigma(u_\perp)^{data}/\sigma(u_\perp)^{MC}$. The ratio of $\sum \bar{E}_T$ and $p_T^{\ell\ell}$ is fitted in separate bins of $p_T^{\ell\ell}$ as a function of $\sum \bar{E}_T$:

$$\sigma(u_\perp)(\sum \bar{E}_T) = c + d \cdot \sqrt{\sum \bar{E}_T} \quad (6.17)$$

The corrected value of $u_\perp$ in the simulation of W events is scaled to make the width match to that of the data, with $p_T^{\ell\ell}$ replaced by $p_T^{true,V}$:

$$u_\perp^{MC,corr} = u_\perp^{MC} \times r(\sum \bar{E}_T, p_T^{true,V}) \quad (6.18)$$

The correction of $u_\parallel$ in the simulation of W events comprises a response and a resolution correction:

$$u_\parallel^{MC,corr} = \langle u_\parallel^{data} \rangle + (\langle b^{data} \rangle - \langle b^{MC} \rangle) \cdot r_\parallel + (u_\parallel^{MC} - \langle u_\parallel^{data} \rangle) \cdot r_\parallel \quad (6.19)$$

Here $\langle u_\parallel^{data} \rangle(\sum \bar{E}_T, p_T^{\ell\ell})$ averages over data in each bin of $(\sum \bar{E}_T, p_T^{\ell\ell})$. And the dependence on $p_T^{\ell\ell}$ is fitted in each $\sum \bar{E}_T$ bin linearly:

$$\langle u_\parallel^{data} \rangle(p_T^{\ell\ell}) = e + f \cdot p_T^{\ell\ell} \quad (6.20)$$

Similarly, $(\langle b^{data} \rangle - \langle b^{MC} \rangle)(\sum \bar{E}_T, p_T^{\ell\ell})$ is derived as the average of bias in bins of $(\sum \bar{E}_T, p_T^{\ell\ell})$ and fitted in each bin of $\sum \bar{E}_T$ by a linear function of $p_T^{\ell\ell}$. When the correction is applied to the simulation of W events, $p_T^{\ell\ell}$ is replaced by $p_T^{true,V}$.

6.3.4 Recoil calibration uncertainties

The uncertainties on the recoil calibration stem from the data statistics in the resolution and response correction, as well as the non-closure in the extrapolation from Z to W [12].

$\sum \bar{E}_T$ reweighting uncertainty

The Sherpa MC samples are used as pseudo-data to perform a closure test of $\sum \bar{E}_T$ reweighting procedure. The reweighting coefficients are derived from Sherpa W and Z events and applied to Powheg samples following the $\sum \bar{E}_T$ reweighting procedures in Section 6.3.3. The reweighted Powheg is then compared with Sherpa in 2D

($\sum \bar{E}_T, p_T^{true,V}$) distribution. The non-closure for the $\sum \bar{E}_T$ distribution in each bin of $p_T^{true,V}$ defines a systematic variation as a 2-point systematic uncertainty.

Resolution and response uncertainties

The MC statistical uncertainty are neglected, since the size of simulated MC samples is one order of magnitude larger than the size of data samples. In a given bin of $p_T^{\ell\ell}$, the uncertainties of two parameters of the fitted resolution function $\sigma(u_\perp)(\sum \bar{E}_T)$ stem from limited data statistics. The variations of these two parameters are transformed into 2 up and 2 down uncorrelated variations as the linear combinations of the initial uncertainty variations.

The diagonalization of the statistical uncertainty variations is also done for the two parameters in the fit of $\langle u_{||}^{data} \rangle (p_T^{\ell\ell})$ in bins of $\sum \bar{E}_T$. The non-linearity of the response is accounted for by an alternative 2nd polynomial fit to $\langle u_{||}^{data} \rangle$ as a function of $p_T^{\ell\ell}$ in bins of $\sum \bar{E}_T$. The variation with respect to the nominal linear fit is taken to be an additional source of uncertainty correlated over the $\sum \bar{E}_T$ bins.

The extrapolation uncertainties are also considered, accounting for the potential detector response and resolution differences to the hadronic recoil between W and Z events. First, the W samples are reweighted in 2D ($\sum \bar{E}_T, p_T^V$) to match the distributions of the Z process. Then, the neutrino of the W is treated to be a charged lepton in the recoil reconstruction, with the cone replacement implemented in the exact way as what is done for an electron or a muon. In the end, the events with truth FSR photons are discarded in both W and Z samples to get rid of the difference between W and Z due to their different patterns of final state photon radiation. After that, $u_{||}$ and $u_\perp$ are computed using $\bar{p}_T^{V,true}$, and compared between W and Z process. The extrapolation uncertainties are treated as correlated in the analyses.

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3555 Chapter 7

3556 Measurement of the W-boson 3557 transverse momentum

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7.1 The motivation of the p_T^W measurement

The measurement of the spectrum of the W-boson transverse momentum (p_T^W) is mainly driven by the needs for the W mass measurement. The lepton p_T spectrum, as one important observable used for the extraction of W-boson mass, is heavily dependent on the modelling of p_T^W . A well-controlled modelling of p_T^W distribution with a precision as good as around 1 % is required in the region of low p_T^W (below 20-30 GeV) [1]. The cross-section in this low p_T region, known as the Sudakov peak, cannot be predicted using only fixed-order perturbative QCD calculation, and has to be handled with either analytical resummation [2] or the description of parton shower [3].

In the previous ATLAS measurement of W mass [4], the modelling of p_T^W relies on the experimentally measured the Z-boson transverse momentum spectrum (p_T^Z) to tune the parton shower [5]. The tuning parameters are passed to the prediction of p_T^W spectrum, with an uncertainty due to the experimental precision of p_T^Z measurement. In addition, the parton shower PDF uncertainty also contributes to the uncertainty on the p_T^W modelling. Other uncertainty sources include the uncertainty in the masses of charm and bottom quarks, of which the impact can be estimated in Pythia 8 [3], as well as the uncertainty from the higher-order QCD corrections. While the variations of masses of charm and bottom quarks turn out to have a relatively small impact on W mass measurement, the uncertainty from the higher-order QCD corrections is much larger and has to be assessed carefully. The missing higher-order QCD corrections to the parton shower are estimated by the variations of the factorization scale (μ_F) for QCD ISR. Although these corrections are mostly correlated between the W and Z-boson productions initiated by light quarks (u, d and s), therefore cancelled out in the ratio of p_T^W/p_T^Z , there is a certain level of decorrelation due to the different fractions of heavy quarks involved in the production. In general, the QCD scales variations have been considered to be an unreliable way to estimate the missing higher-order QCD uncertainty, as they only provide an approximate size of the missing higher-order effect without an insured coverage. On top of that, these variations offer no indication about how the uncertainties are correlated across a distribution.

To overcome the difficulty brought by the extrapolation of p_T^Z to the p_T^W modelling, it is worth measuring the p_T^W spectrum directly from the data, so that the estimation of uncertainty using QCD scale variations will be avoided. In the low- μ data taking conditions, the improved recoil resolution is expected to bring a <2 % uncertainty in

3604 p_T^W with a granularity of around 6 GeV in the Sudakov peak, which will greatly reduce
 3605 the modelling uncertainty in m_W .

3606 7.2 Methodology

3607 The measurement of the p_T^W spectrum corresponds to a measurement of the dif-
 3608 ferential Drell-Yan cross-section with a direct decay of the W-boson: $W \rightarrow \ell\nu$, where
 3609 $\ell = e$ or μ . The measurement in the e and μ channels as well as in the combination
 3610 are reported in a common fiducial region close to the definition of the event selection
 3611 in Table 5.1, such that the extrapolation stays minimal:

3612 — $W \rightarrow \ell\nu$: $p_T^\ell > 25$ GeV, $|\eta^\ell| < 2.5$, $p_T^\nu > 25$ GeV and $m_T > 50$ GeV.

3613 The lepton kinematics in this definition of this cross-section measurement is taken
 3614 from the state before QED FSR (a.k.a at Born-level).

3615 Due to the neutrino in the decay product of the leptonic decay of the W-boson, the
 3616 kinematics of the boson cannot be fully reconstructed in the transverse plane using
 3617 the decay product. Therefore, the p_T^W has to be inferred from the hadronic recoil u_T ,
 3618 which originates from the QCD ISR that gives the none-zero transverse momentum
 3619 distribution to the boson.

3620 The detector effects bias and smear the measurement of p_T^W via the hadronic
 3621 recoil. In order to extract the true underlying p_T^W spectrum from the response of
 3622 the detector, all the detector effects, such as the limited detector acceptance, the
 3623 efficiency of the object reconstruction and the finite detector resolution, are accounted
 3624 for by an unfolding procedure. For events passing both fiducial volume selection and
 3625 reconstructed-level selection ($T\&R$), the migration from p_T^W at truth-level to u_T at
 3626 reco-level, normalized to unity in each truth bin j , is described by the migration
 3627 matrix in the simulation:

$$M_{ij} = \frac{N_{ij}}{\sum_k N_{kj}} \quad (7.1)$$

3628 N_{ij} is the number of events in u_T bin i and p_T^W bin j . Weights at truth level and
 3629 reconstructed level are kept separately for these events. The baseline binnings of the
 3630 u_T and of the p_T^W spectrum after the optimization (Section 7.4) are chosen to be:

3631 — u_T : 1 GeV bin size up to $u_T=100$ GeV, 5 GeV bin size for $u_T \in [100,200]$ GeV,
 3632 10 GeV bin size for $u_T \in [200,600]$ GeV.

3633 — p_T at 5.02 TeV: [0, 7, 14, 21, 28, 35, 42, 49, 56, 63, 77, 92, 115, 145, 175,
 3634 220] GeV.

3635 — p_T at 13 TeV: [0, 7, 14, 21, 28, 35, 42, 49, 56, 63, 77, 92, 115, 145, 175, 220,
 3636 310, 600] GeV.

3637 Examples of migration matrices are shown in Figure 7.1. The non-diagonal fea-
 3638 ture of the migration is attributed to the large detector resolution in u_T as well as
 3639 the detector response of $\langle u_T/p_T^W \rangle$ is below one and dependent on p_T^W , which can be
 3640 understood as the recoil measurement suffers from the signal loss due to noise cuts,
 3641 particles traversing beyond the detector acceptance as well as the non-compensating
 3642 nature of the calorimeter. The conditional probability to observe a reconstructed
 3643 value of u_T in bin i and a given truth value of p_T^W in bin j is defined by correcting for
 3644 the efficiency in the migration matrix M_{ij} , which yields the response matrix R_{ij} :

$$R_{ij} = M_{ij}\epsilon_j \quad (7.2)$$

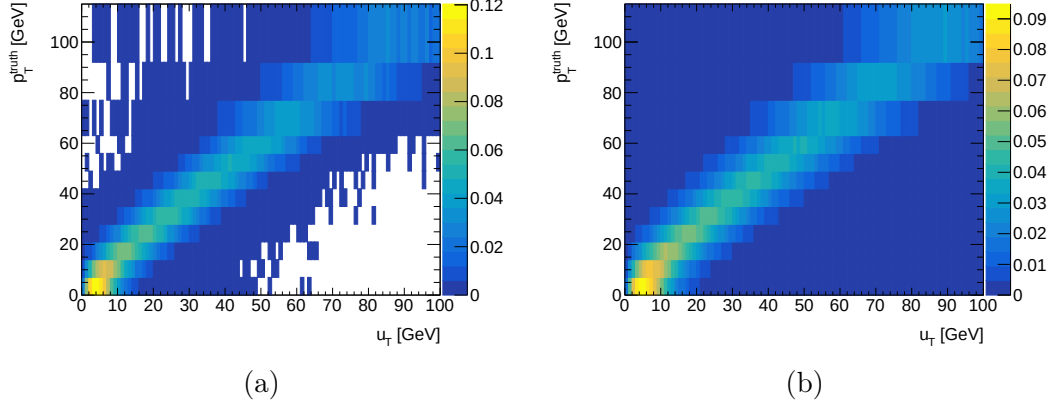


Figure 7.1 – Examples of migration matrices for $W^- \rightarrow e\nu$ at (a) 5.02 TeV and at (b) 13 TeV.

And the efficiency of reconstructing a truth level signal, in bin j of p_T^W is:

$$\epsilon_j = \frac{N_j^{reco\&gen}}{N_j^{gen}} = \frac{\sum_i N_{ij}}{N_j^{gen}} \quad (7.3)$$

Here $N_j^{reco\&gen}$ refers to the number of events in a p_T^W bin j passing $T\&R$ selection cuts, and N_j^{gen} the number of events in the corresponding bin passing only the fiducial volume selection.

Starting from the underlying true distribution T after the fiducial cuts, the detector level distribution D in bins of u_T can be expressed in:

$$D_i = \sum_j R_{ij} T_j + B_i \quad (7.4)$$

The term B_i accounts for the presence of background after the reconstructed level cuts. Considering the fact that events beyond the fiducial volume can be reconstructed, an additional correction of purity has to be applied to the reconstructed level distribution in the unfolding by:

$$p_i = \frac{N_i^{reco\&gen}}{N_i^{reco}} = \frac{\sum_j N_{ij}}{N_i^{reco}} \quad (7.5)$$

The numerator of the purity correction in this formula covers the entire fiducial volume. The level of efficiency and purity corrections are shown in Figures 7.2.

Including the effect of purity correction, the best estimate of the underlying true p_T^W distribution determined by the unfolding procedure is given by:

$$\tilde{U}_j = \sum_i U_{ij} (D_i - B_i) p_i \quad (7.6)$$

where the unfolded spectrum estimating the underlying distribution T is presented by $\tilde{U}_j$. The unfolding transformation is symbolized by U_{ij} . Technically, the unfolding matrix U_{ij} in the analysis is determined by the iterative Bayesian regularized unfolding approach with the implementation of the D'Agostini iterative scheme [6][7]. For a given resolution of the unfolded spectrum, the number of unfolding iterations has to be chosen for the solution to minimize the total uncertainty as well as the possible bias. The optimization of the unfolding will be discussed in Section 7.4.

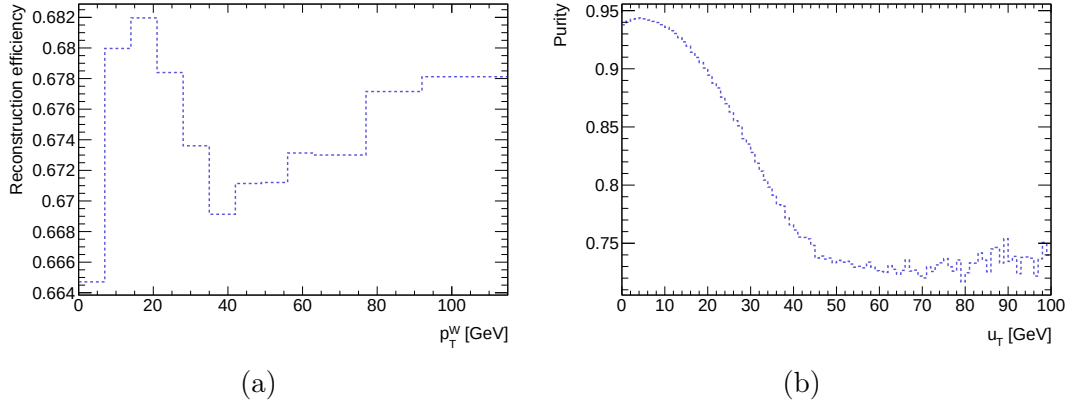


Figure 7.2 – Example of (a) efficiency correction and (b) purity correction in 13 TeV $W^- \rightarrow e\nu$ channel.

Based on Equation 7.6, the uncertainties sources are divided to four types and propagated to the unfolded level following the corresponding procedure:

- Data statistics. The data distribution D_i is fluctuated at reconstructed level by bootstrap. The spread of toys after unfolding is taken to be the data statistical uncertainty at unfolded level.
- MC statistics. The unfolding transformation U_{ij} is varied by fluctuating the migration matrix, the efficiency correction and the purity correction via bootstrap. The spread of toys is considered to be the MC statistical uncertainty at unfolded level.
- Experimental systematics. The effect of the systematic variation is applied to the unfolding transformation U_{ij} , including migration matrix, efficiency correction and purity correction. After that, the change in the unfolded spectrum is taken to be the experimental systematic of the given source.
- Background uncertainty. The background estimation B_i at reconstructed level is varied according to the relevant systematics. The change in the unfolded spectrum is then the background systematic of source under study.

On the technical side, the bootstrap method [8] is implemented by assigning a random weight to each event for each (pseudo) data set, defined as follows:

$$w = \mathcal{P}(n, 1) \quad (7.7)$$

where n is an integer randomly generated according to the Poisson distribution of unit mean, and $w = \mathcal{P}(n, 1)$ is the value of the Poisson probability to observe n events with an expectation of 1. The covariance matrices of data statistics and MC statistics are therefore the standard deviation of the N_{bs} bootstrap toys:

$$C_{kl}^{stat} = \frac{1}{N_{bs} - 1} \sum_{\alpha=1}^{N_{bs}} (\tilde{U}_k^\alpha - \langle \tilde{U} \rangle_k) (\tilde{U}_l^\alpha - \langle \tilde{U} \rangle_l) \quad (7.8)$$

with k and l looping all the bins. The systematic uncertainty is fully correlated across all the bins in the unfolded spectrum. Therefore, the elements of the covariance matrix for a given source of systematic a are:

$$C_{kl}^a = \delta \tilde{U}_k^a \delta \tilde{U}_l^a = (\tilde{U}_k^a - \tilde{U}_{Nom}^a) (\tilde{U}_l^a - \tilde{U}_{Nom}^a) \quad (7.9)$$

Then, the total covariance matrix for a given channel is expressed as a sum of all the uncertainty components:

$$C_{kl}^{tot} = C_{kl}^{stat,Data} + C_{kl}^{stat,MC} + \sum_a C_{kl}^a \quad (7.10)$$

with a running over all source of systematic uncertainties.

One major caveat of the unfolding procedure is that the hypothetical underlying distribution of a given physics variable in the simulation is not guaranteed to exactly match its true value, therefore causing a bias in the simulation of the migration matrix, as well as the efficiency and the purity corrections. While using the regularized iterative unfolding, this prior hypothesis bias introduces an uncertainty on the unfolded result. In ATLAS, it is recommended to estimate this source of uncertainty by following a procedure of data-driven closure test, of which the guiding principle is:

- The MC events passing $T\&R$ selections are reweighted at truth level to achieve the optimal agreement in shape at reconstructed level between data and MC for the physics observable under study. It has been checked that whether or not applying the reweighting to those events passing truth level selection but failing reconstruction selection doesn't change the result.
- For events passing $T\&R$ selections, the reconstructed-level MC distribution after the above truth level reweighting is unfolded as pseudo data using the migration matrix from the initial unreweighted MC.
- The unfolding bias is estimated by comparing the unfolded result with the reweighted truth distribution for events passing $T\&R$ selections.

Conventionally, the reweighting of the truth-level distribution is taken from the data-to-MC ratio at reconstructed level. But in the p_T^W measurement, due to the fact that the possible discrepancy (between the modelling of p_T^W and true physics) at truth-level may be smeared out at reconstructed level due to the large resolution, and that the proportionality between u_T and p_T^W is a function dependent on p_T^W , the reweighting has to use parameterized functions in p_T^W determined from a data-driven method described in Section 7.3, otherwise a simple reweighting of the data-to-MC ratio at reconstructed level will lead to large unfolding bias.

In parallel to the p_T^W measurement, the p_T^Z is carried using the same low pile-up datasets to validate the analysis strategy of unfolding the hadronic recoil to obtain the boson- p_T spectrum, since the p_T^Z can be measured by unfolding the u_T or by unfolding the $p_T^{\ell\ell}$ distribution. The latter offers a precision of about one order of magnitude better than the former. In the end, it is shown that the results of p_T^Z measurements based on u_T and based on $p_T^{\ell\ell}$ are compatible with each other, which validates method of unfolding the u_T distribution for the p_T^W measurement. The fiducial cross-section of Z production, along with the cross-section ratios, will be reported in Section 7.5.2.

7.3 p_T^W reweighting and bias uncertainties

The mis-modelling of p_T^W in the simulation brings large discrepancy of transverse momentum distribution at detector level between data and MC. It has been observed that the unfolding starting from this non-optimal initial distribution leads to a percent level unfolding bias as a source of uncertainty at unfolded level.

The analysis results can be optimized by reweighting the MC simulation according to a function of p_T^W , such that the best possible data to MC agreement on the

reconstructed u_T distribution is achieved, ensuring an unbiased unfolding result. This correction procedure is referred to as the p_T correction, of which the relevant uncertainties are derived from the possible variations of this correction. The form of p_T correction is derived from the following χ^2 minimization:

$$\begin{aligned}\chi^2 &= \sum_{ij} \Delta_i^T C_{ij}^{-1} \Delta_j \\ \Delta_i &= (D_i - B_i) - \sum_{ij} R_{ij} \times (w_T(p_T^W))_j\end{aligned}\tag{7.11}$$

Here Δ is the difference between the background subtracted data ($D_i - B_i$) and the reconstruction level response (R_{ij}) of the simulation after the particle level reweighting $w(p_T^W)$. The χ^2 is computed using the total covariance matrix at reconstructed level (C_{ij}). It was found that a satisfactory data-to-MC agreement can be realized by choosing the following form for the baseline reweighting function:

$$w_T(p_T^W) = N[(1 + ap_T^W + b(p_T^W)^2) \cdot (1 - c + c \cdot r_{NNPDF/CT10}(p_T^W))]\tag{7.12}$$

The ratio $r_{NNPDF/CT10}(p_T^W)$ is calculated by DYTURBO. It represents the correction in the full phase space, from the default PDF set (CT10 [9]) used in the Powheg [10]–[13] samples to a more modern PDF set: NNPDF3.0 [14]. The additional terms in the function adds more flexibility to the fit. The form of the baseline reweighting function is optimized separately for each centre-of-mass energy and for each charge of the boson. The correction is only derived over $0 < p_T^W < 100$ GeV. The value of the correction is frozen at 100 GeV for any larger p_T^W . The average of correction functions in the e and μ channels are used because of the good compatibility between them.

The bias uncertainty related to this data driven p_T correction comprises three parts:

- Fit uncertainty: The uncertainty in the fit parameters.
- Parameterization uncertainty: The arbitrariness in the form of the reweighting function.
- Initial (p_T^W, y) distribution uncertainty: The potential mis-modelling in the correlation between boson p_T and rapidity.

The fit uncertainty is derived by repeating the fit of the p_T -correction, but using p_T -corrected migration matrix, as well as the efficiency and purity corrections in the MC. The new correction factor of p_T -reweighting given by the fit is compatible with 1, showing a small non-closure within the uncertainty due to the effect of averaging the p_T reweighting between the electron and muon channels. After the fit, based on the covariance matrix, the function parameters are transformed to be orthogonal and are parameterized by the corresponding nuisance parameters. Given that there are three parameters in the baseline fit of Equation 7.12, three variations of the fit uncertainty are produced following this prescription.

The parameterization uncertainty is determined by repeating the reweighting procedure of Equation 7.11, but replacing the baseline $w_T(p_T^W)$ with other functions:

$$Expo. \times Pol. : w_T(p_T^W) = N[(1 + ap_T^W + b p_T^{W^2})(1 - c + \frac{c}{\tau} e^{-p_T^W \cdot \tau})]\tag{7.13}$$

3769

$$2nd - Order Pol. : w_T(p_T^W) = N[1 + ap_T^W + b p_T^{W^2}]\tag{7.14}$$

$$\text{Gaus.} + \text{Pol.} : w_T(p_T^W) = N[(1 + a p_T^W + b p_T^{W^2} + c G(p_T^W; \mu, d)] \quad (7.15)$$

$$\text{Gaus.} \times \text{Pol.} : w_T(p_T^W) = N[(1 + a p_T^W + b p_T^{W^2})(1 - c + c G(p_T^W; \mu, \sigma))] \quad (7.16)$$

$$2 \text{Gaus} + \text{Pol.} : w_T(p_T^W) = N[1 + a p_T^W + b p_T^{W^2} + c G(p_T^W; \mu, d) + e G(p_T^W; \mu, f)] \quad (7.17)$$

Only the functions that bring a reasonable data-to-MC compatibility of u_T at reconstructed level are chosen to be the variations for the parameterization uncertainty. After rejecting the functions leading to large χ^2 values of data-to-MC, Equation 7.13 is chosen to be the alternative function form at 13 TeV, while Equations 7.14 and 7.15 are taken to be the variations at 5.02 TeV, for W^+ and W^- respectively.

The uncertainty arising from the initial (p_T^W, y) distribution in the simulation is assessed by once again performing the reweighting procedure of p_T -correction, but starting from one of the following predictions of (p_T^W, y) in the simulation: DYTURBO [15][16] using CT10 PDF, DYTURBO using NNPDF3.0 PDF, Herwig7 [17] and Pythia [3]. The nominal MC is reweighted in 2D to the possible alternative prediction of (p_T^W, y) . For each alternative prediction of (p_T^W, y) , the form of reweighting function is chosen to be the one from 7.12 to 7.17 that gives the best description of data for u_T at reconstructed level. In the end, DYTURBOCT10 is chosen to be the proxy of initial (p_T^W, y) uncertainty variation at 13 TeV while DYTURBONNPDF3.0 is chosen for all 5.02 TeV channels. These two choices give a reasonable data-MC agreement at reconstructed level and are the most conservative variations in the corresponding center-of-mass energy.

7.4 Optimization of Iterative Bayesian Unfolding

The guideline of measuring the p_T^W spectrum is to minimize the total uncertainty while reaching a fine resolution after unfolding. In particular, a low contribution from the unfolding bias uncertainty (approximately at 1 % level) will always be favoured, as a precise estimation of the bias uncertainty is extremely challenging in the context of this analysis. In order to achieve this goal, the two major parameters for the refinement of the Iterative Bayesian Unfolding, namely the bin width at truth level and the number of unfolding iterations, are optimized simultaneously in a dedicated study.

In each of the eight decay channels, the uncertainty on the unfolded spectrum is studied in bin widths of 5, 6, 7, 8 and 9 GeV in the low p_T^W region and with 1, 3, 5, 7, ..., 29 iterations of unfolding. In total the 75 versions of unfolding in a 2D mapping of bin width and number of iterations enable an optimization of both quantities. The optimization accounts for statistical uncertainty, experimental and background uncertainties, as well as bias uncertainties. Considering that the scheme of the optimization should be consistent with the combination of electron and muon channels and that the statistics is roughly the same in the two channels for a given charge at the same center-of-mass energy, the statistical uncertainties are always divided by a factor of $\sqrt{2}$, even though the study is performed per individual channel. Meanwhile, the experimental systematics, background uncertainties and bias uncertainties are considered correlated between the electron and the muon channels, since the dominating components among them are the recoil calibration uncertainty and the bias

uncertainty that are correlated between the two channels. The total uncertainty in each channel is the quadratic sum of the above components.

The total uncertainties in various options of bin width and with different number of iterations are presented in Figures 7.3 and 7.4, where a linear interpolation between the centers of the two nearest bins is performed at a given value of p_T^W , so as to facilitate the comparison between the choices of bin width. It turns out that both the 5.02 TeV bin width and the 6 GeV bin width lead to a significantly larger uncertainty at both center-of-mass energies, mainly due to the large unfolding bias uncertainty. A wider binning tends to help reduce the total uncertainty in the p_T^W region thanks to not only the lower statistical uncertainty in each bin, but also a smaller bias uncertainty. Consequently, a bin width of 7 GeV is chosen for both 5 and 13 TeV, as a trade-off between the total uncertainty and the resolution. In this way, the 7 GeV bin width keeps a sufficiently low total uncertainty while preserving a reasonable resolution of the unfolded spectrum.

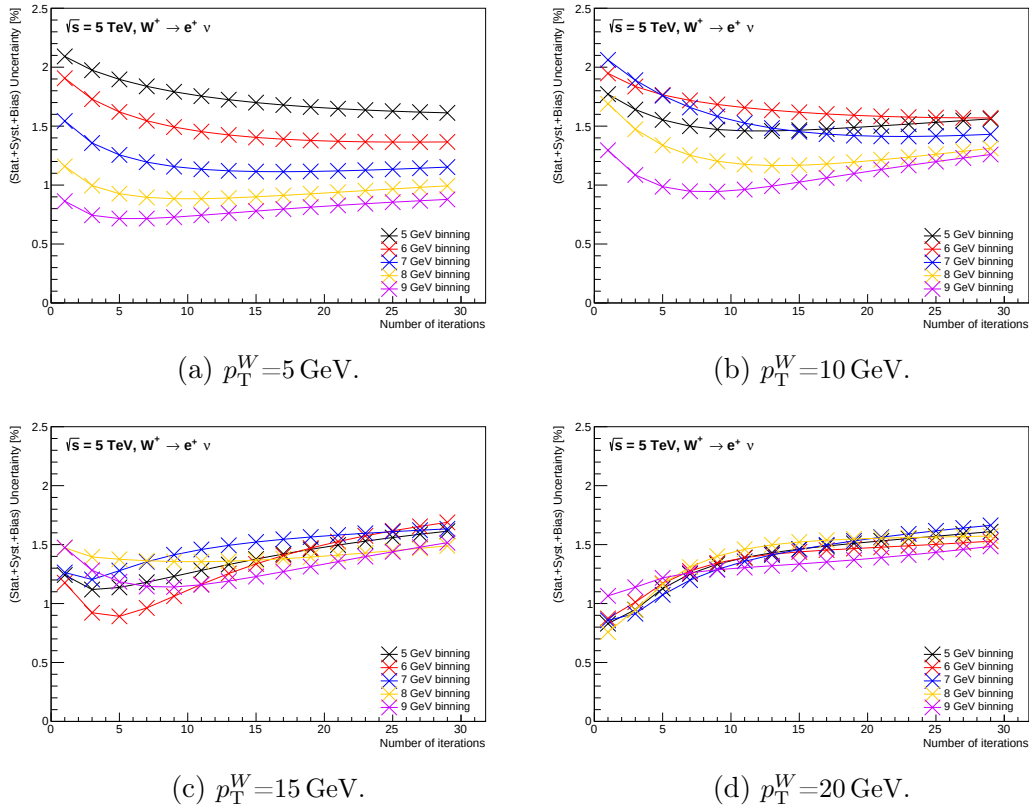


Figure 7.3 – Total uncertainty in various bin widths over different iterations for 5.02 TeV $W^+ \rightarrow e \nu$ channel. The total uncertainty is interpolated to $p_T^W = 5, 10, 15, 20$ GeV.

Once the bin width is chosen, the optimization is focused on the number of iterations, which is investigated separately for 5.02 and 13 TeV, as shown in Figures 7.5 and 7.6. The general tendency is that the increased number of iterations reduces the bias uncertainty but comes with a price of inflated both statistical and the other systematic uncertainties. The main guideline of choosing the number of iterations for a given bin width is to keep the bias uncertainty as low as possible in the $p_T^W < 10$ GeV region, since the bias uncertainty is the least well understood uncertainty source in

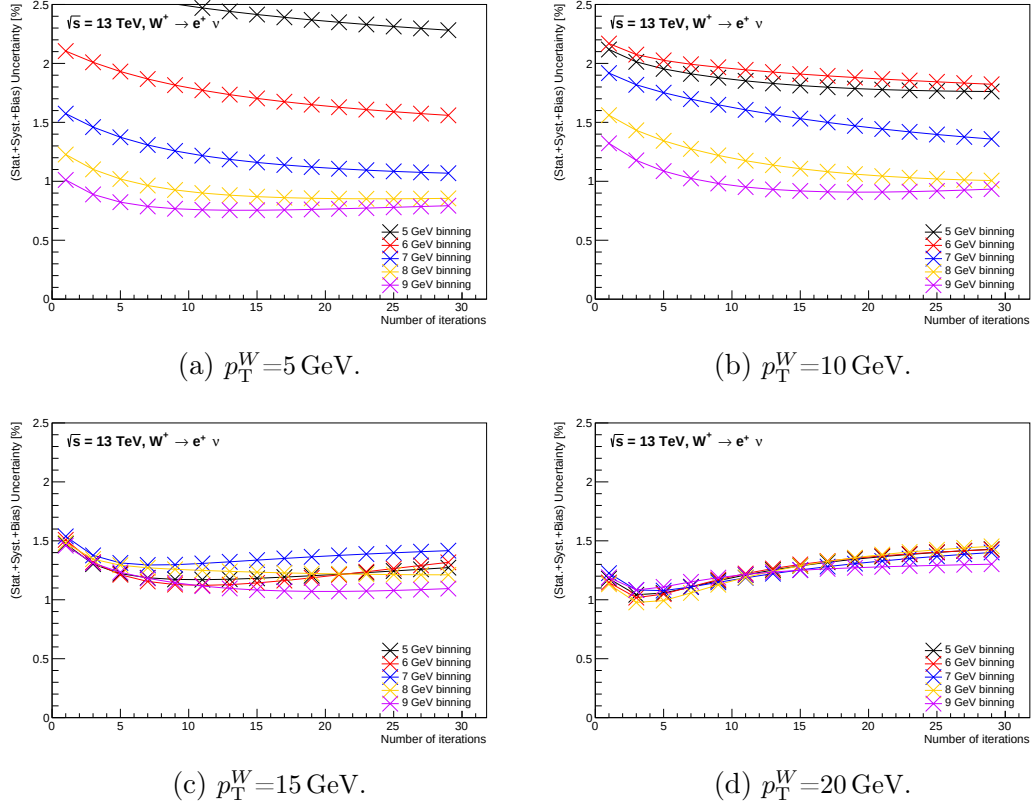


Figure 7.4 – Total uncertainty in various bin widths over different iterations for 13 TeV $W^+ \rightarrow e\nu$ channel. The total uncertainty is interpolated to $p_T^W = 5, 10, 15, 20 \text{ GeV}$.

this region. If possible, the total uncertainty needs to be kept low as well, in addition to the low bias uncertainty. While reaching the best uncertainty in the region $p_T^W < 10 \text{ GeV}$, it is also necessary not to worsen the result at higher p_T^W unreasonably. Following all these criteria, at 13 TeV, a large number of iterations is preferred in the end, so as to keep the bias uncertainty better under control, especially in the $p_T^W < 10 \text{ GeV}$ region. In this low p_T region, the other uncertainty components only increase very slowly after 20 iterations. For the 5.02 TeV channels, the effective drop of the bias uncertainties along with the relative slow increase of the statistical and the other systematic uncertainties brings a plateau of total uncertainty in the low p_T^W region, starting from around 9 iterations, even though the total uncertainty always increases with the number of iterations at above 15 GeV. The final decision of the unfolding optimization is to use a bin width of 7 GeV in the low p_T^W region for both center-of-mass energies, with 25 iterations of unfolding for 13 TeV and 9 iterations for 5.02 TeV.

7.5 Results

7.5.1 Unfolded data spectra and uncertainty break-down

The unfolded absolute different cross-sections are compared to the prediction of the Powheg+Pythia 8 samples with p_T^W reweighting correction in Figures 7.7 and 7.8.

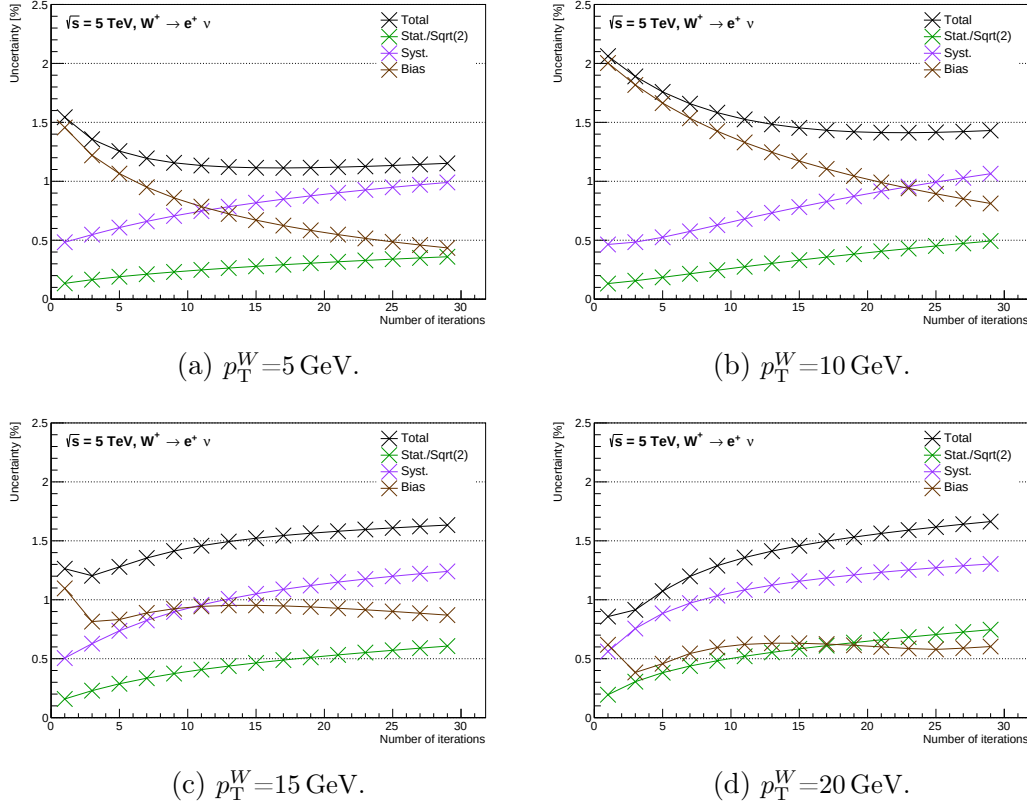


Figure 7.5 – Total uncertainty in various binnings over different iterations for 5.02 TeV $W^+ \rightarrow e\nu$ channel. The total uncertainty is interpolated to $p_T^W = 5, 10, 15, 20 \text{ GeV}$.

3851 The ratios of the unfolded normalized differential cross-sections in different predictions
 3852 to the unfolded data are shown in Figures 7.9 and 7.10. Putting aside the luminosity
 3853 uncertainty, the experimental uncertainties vary from about 2 % in the low p_T^W region
 3854 to about 5 % (4 %) when $p_T^W = 100 \text{ GeV}$ at 5.02 TeV (13 TeV). On top of that, the
 3855 luminosity contributes to 1 % at 5.02 TeV and 0.92 % at 13 TeV. The predictions
 3856 for comparison include Powheg+Pythia 8 AZNLO [5] without p_T^W correction, Pythia
 3857 8 AZ, Sherpa and DYRes [18][19]. For both center-of-mass energies, both lepton
 3858 flavors and both W boson charges, a similar level of agreement with data is found for
 3859 Powheg+Pythia8 AZNLO and Pythia 8 AZ. Sherpa predicts a softer spectrum while
 3860 DYRes predicts a harder one.

3861 The breakdown of the uncertainty at the reconstructed level is shown in Fig-
 3862 ures 7.11, 7.12 for the region $0 < p_T^W < 100 \text{ GeV}$. The absolute uncertainties at the
 3863 unfolded level are shown in Figures 7.13, 7.14. The corresponding shape uncertainties
 3864 at unfolded level are presented in Figures 7.15, 7.16, in which the central value and
 3865 the uncertainty variations of the p_T^W distributions are normalized to unity when the
 3866 uncertainty are calculated. The region of interest for p_T^W shown in the uncertainty
 3867 breakdown plots at unfolded level is restricted to $0 < p_T^W < 63 \text{ GeV}$, although the
 3868 measurement of the p_T^W spectra higher than 63 GeV is also performed in the analysis.

3869 At the unfolded level, the total experimental uncertainty starts from 1 % at the low-
 3870 est p_T bin and remains below around 2 % up to $p_T = 40 \text{ GeV}$ in general. For $p_T > 40 \text{ GeV}$
 3871 the uncertainty grows moderately at 13 TeV and increases steeply at 5.02 TeV. The

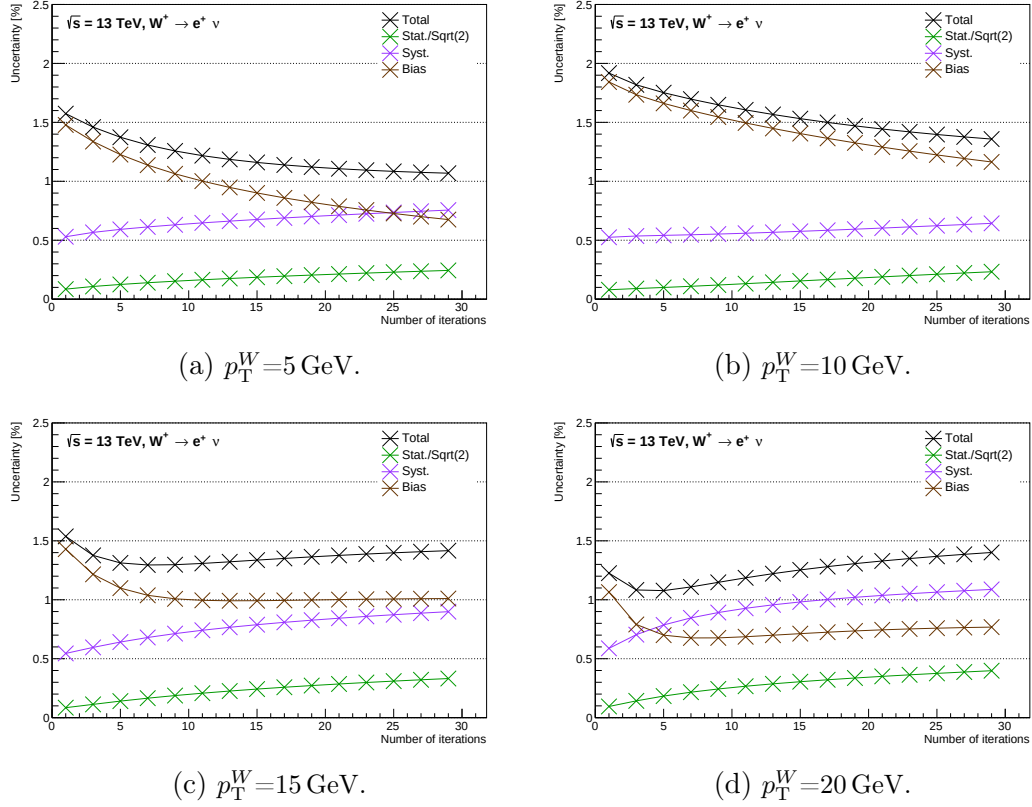


Figure 7.6 – Total uncertainty in various binnings over different iterations for 13 TeV $W^+ \rightarrow e\nu$ channel. The total uncertainty is interpolated to $p_T^W = 5, 10, 15, 20 \text{ GeV}$.

data statistics, recoil calibration systematics, the signal generator systematic and bias uncertainties are the largest uncertainty components. The three first p_T bins receive a very important impact from the bias uncertainties. All the other remaining uncertainties are small in comparison, except for the background uncertainties at 13 TeV because of the larger multijet and $t\bar{t}$ contamination.

7.5.2 Channel combination

The combination of electron and muon channel cross-sections follows the standard prescription of the Best Linear Unbiased Estimator method (BLUE) [20][21] using the following χ^2 minimization:

$$\chi^2 = (X - \bar{X})^T C^{-1} (X - \bar{X}) \quad (7.18)$$

where X is the joint $2N$ -bin distribution in the electron and muon channels $X = \{X_1^e, \dots, X_n^e, X_1^\mu, \dots, X_n^\mu\}$ and $\bar{X} = \{\bar{X}_1, \dots, \bar{X}_n; \bar{X}_1, \dots, \bar{X}_n\}$ the vector of the average to be determined. The covariance matrix is of the size $2N \times 2N$:

$$C = \begin{pmatrix} C^e & C^{e\mu} \\ C^{e\mu} & C^\mu \end{pmatrix} \quad (7.19)$$

In this notation, the $N \times N$ matrices C^e and C^μ are the total covariance matrices derived in the electron and the muon channels according to Equation 7.10. And $C^{e\mu}$

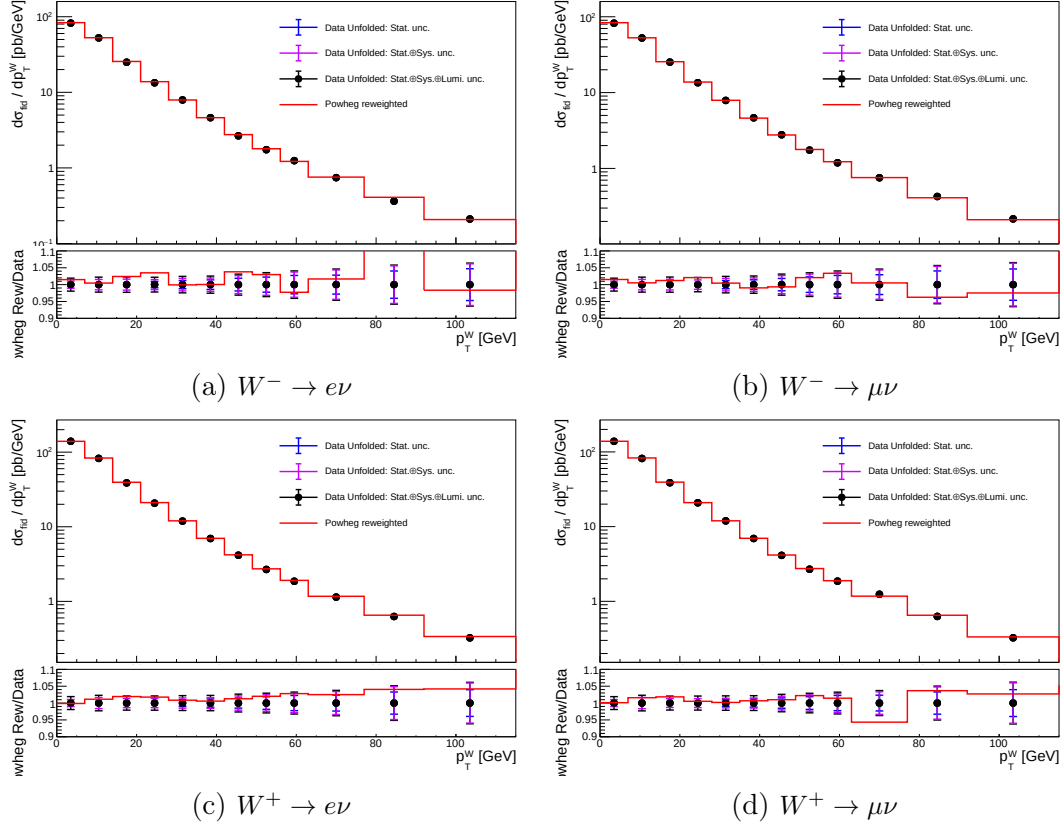


Figure 7.7 – Unfolded measurement results at $\sqrt{s}=5.02$ TeV compared with the nominal p_T -reweighted Powheg+Pythia 8 MC prediction. The bottom panel shows the prediction-to-data ratio.

reflects the sources of systematic uncertainty correlated between the two channels.
The analytical solution to the minimization of the χ^2 is given by:

$$\bar{X} = (H^T C^{-1} H)^{-1} H^T C^{-1} X \quad (7.20)$$

with the $2N \times N$ matrix indicating the structure of the linear system:

$$H = \begin{pmatrix} 1 & & & 0 \\ & \ddots & & \\ 0 & & 1 & \\ 1 & & & 0 \\ & \ddots & & \\ 0 & & & 1 \end{pmatrix} \quad (7.21)$$

where the $2N$ lines correspond to the measured electron and muon spectra and the N columns correspond to the combined spectrum. In the end, the total covariance of the combination is solved by:

$$\bar{C} = H^T C^{-1} H^{-1} \quad (7.22)$$

This combination procedure is valid in the Gaussian limit and under the assumption that the uncertainties are constant in the covariance matrix, not depending on the

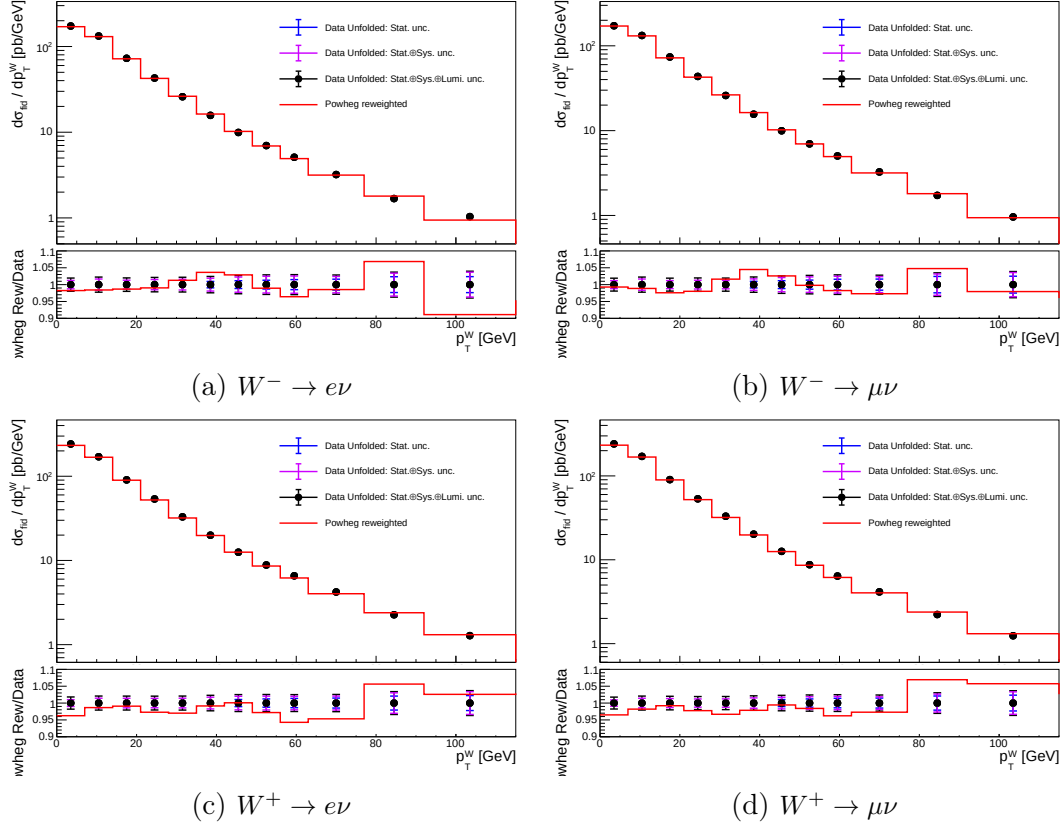


Figure 7.8 – Unfolded measurement results at $\sqrt{s}=13$ TeV compared with the nominal p_T -reweighted Powheg+Pythia 8 MC prediction. The bottom panel shows the prediction-to-data ratio.

value of the measured quantity. When it comes to luminosity or efficiency uncertainties that are the fixed fractions of the measured cross-section, a given relative source of uncertainty has a smaller impact for lower measured values, biasing the combination towards the lowest measured cross-section. A remedy of this issue would be to iterate the combination while rescaling the elements of the covariance matrix according to the combined uncertainties of the previous iteration [22]. In the p_T^W analysis, 4 iterations of combination are carried out, achieving stable results in the end.

The number of systematic sources involved in the combination is summarized in Table 7.1. In addition to the cross-section uncertainties of simulated background mentioned in Section 5.1, the experimental uncertainties considered in the p_T^W analysis have been described in Sections 6.1, 6.2 and 6.3. The multijet statistical uncertainties are also propagated to the unfolded level only in the region $u_T > 100$ GeV, by fluctuating each bin separately according to its statistical uncertainty. The uncertainty variations are considered to be uncorrelated between channels. The multijet statistical uncertainties are accounted for in this way mainly because of their non-negligible contribution appearing only in the high p_T region. The multijet yield and shape uncertainties derived in Section 5.4 are, however, taken to be correlated across all the channels due to its methodological nature. The unfolding bias uncertainties represent the freedom in the underlying p_T^W distribution assumed in the MC for the simulation of the migration matrices as well as the selection efficiencies. All the 5 sources of unfolding bias uncertainty, including the 3 NPs of the fit uncertainty, the

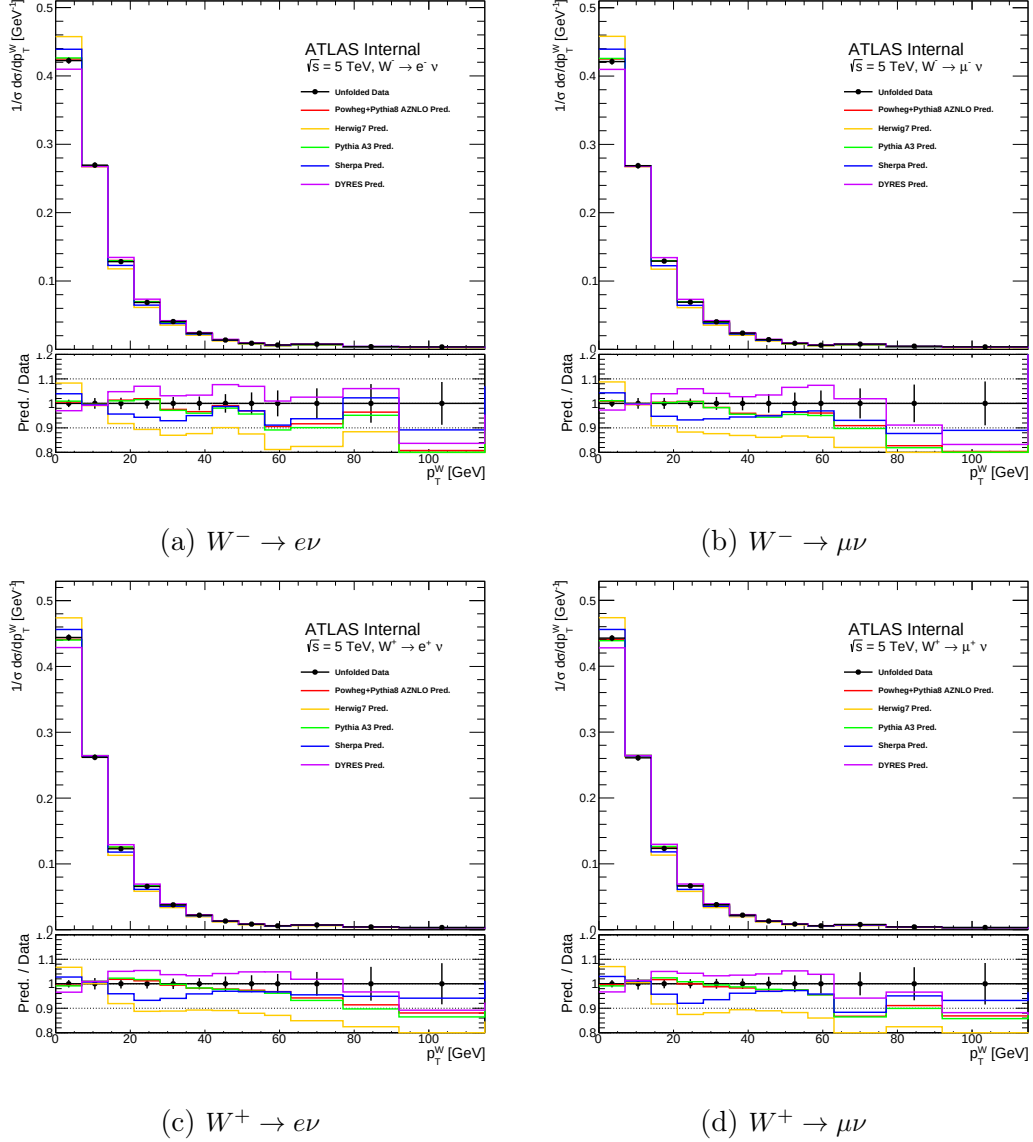


Figure 7.9 – Ratios of the normalized unfolded differential cross-sections by different predictions compared with the unfolded data at $\sqrt{s} = 5.02$ TeV.

parameterization uncertainty and the initial (p_T, y) distribution uncertainty, are considered to be correlated between the electron and the muon channels for a given boson charge and at a given center-of-mass energy. The signal generator systematic uncertainty, taken as the difference in the result when Sherpa is used instead of the nominal Powheg, reflects the possible impact of alternative perturbative and non-perturbative QCD modellings for a fixed p_T^W distribution, which is therefore fully correlated across all the eight channels. A simplified description of the correlation model is shown in Table 7.2 for all the uncertainty components.

The p_T^W spectra are combined for both absolute differential cross-sections and normalized distributions. The combination results are shown in Figure 7.17, along with the uncertainty breakdown after the combination in Figure 7.18. The p_T^Z spectra shown in Figure 7.19 are measured from the $p_T^{\ell\ell}$ distribution, combining $Z \rightarrow ee$ and $Z \rightarrow \mu\mu$ processes at each center-of-mass energy. The integrated fiducial cross-

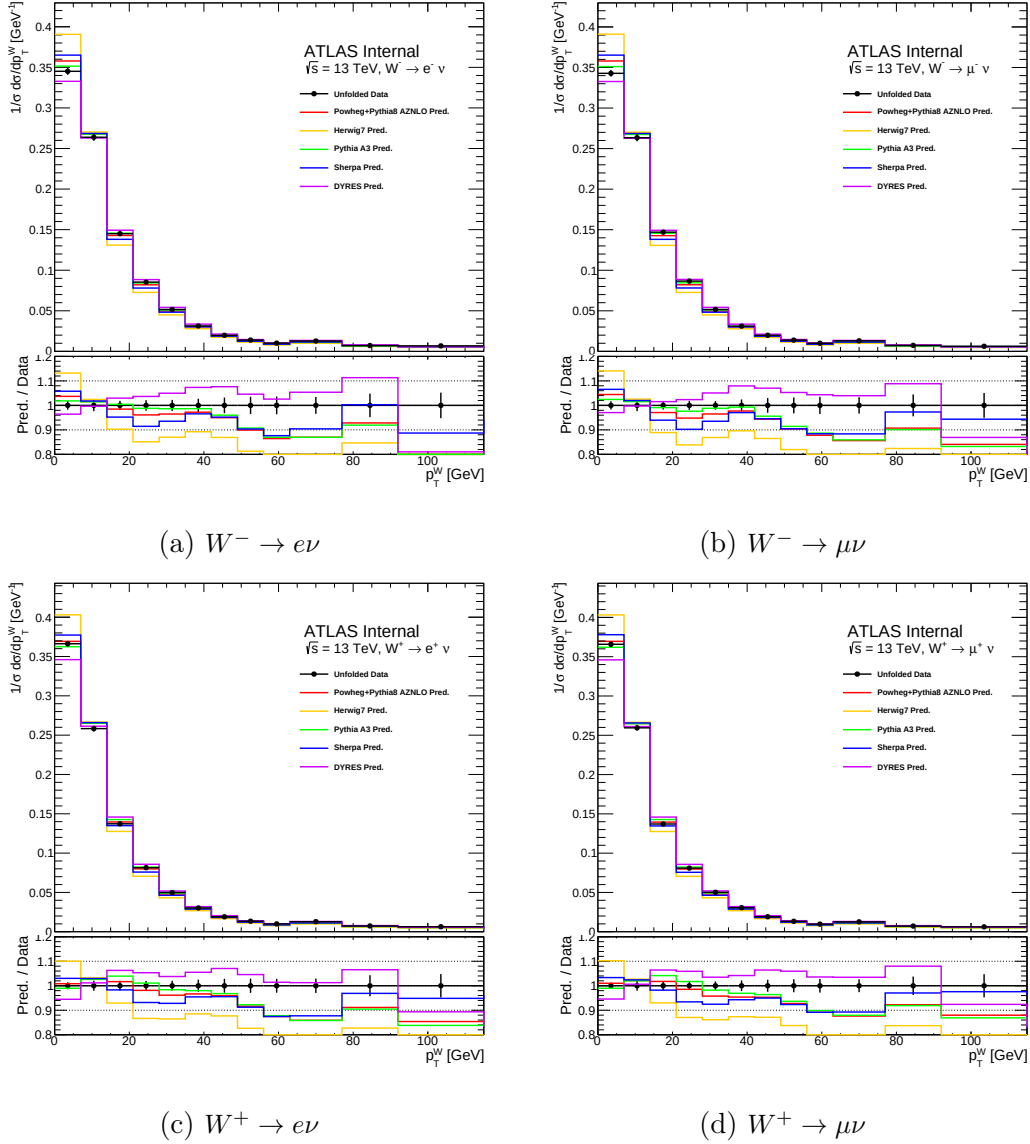


Figure 7.10 – Ratios of the normalized unfolded differential cross-sections by different predictions compared with the unfolded data at $\sqrt{s}=13$ TeV.

sections for W^+ , W^- and Z production at 5.02 and 13 TeV are listed in Table 7.3. The cross-section ratios are presented in Table 7.4.

From Figure 7.20 to 7.24, the measurements of p_T^W and p_T^Z are compared with a variety of ATLAS MCs and the NNLO+NNLL DYTURBO predictions using several global PDF sets. Thanks to the high precision of the measurement, the data is discriminating in the Sudakov peak. In particular, percent-level precision is achieved in the measurement of W^+/W^- ratios. Small discrepancy can be appreciated from the comparison between the measurement and the theory predictions.

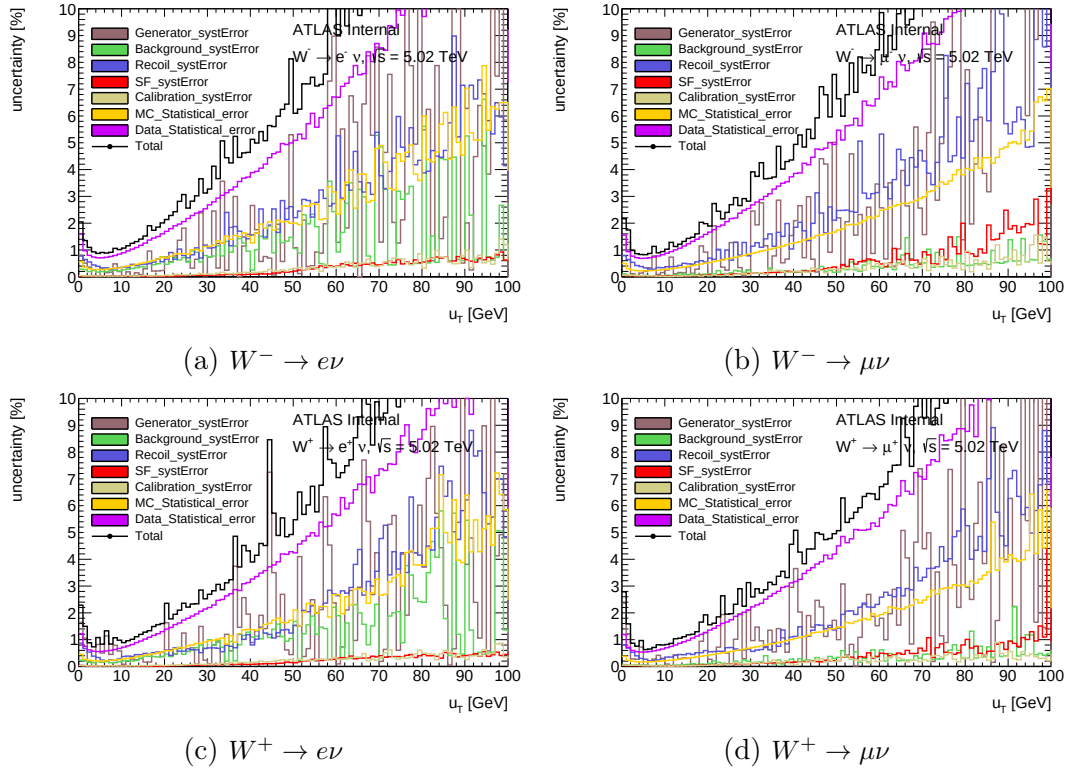


Figure 7.11 – Breakdown of uncertainties on the reconstructed level u_T distributions at $\sqrt{s} = 5.02$ TeV.

Uncertainty category	Nb. of sources at 5.02 TeV	Nb. of sources at 13 TeV
Electron calibration	108	108
Electron efficiencies	620	627
Muon calibration	48	48
Muon efficiencies	908	930
Recoil calibration	68	62
EW & Top backgrounds	24	22
Multijet background	11	11
Unfolding bias	5	5
Generator systematic	1	1
Total	1793	1814

Table 7.1 – Sources of systematic uncertainties in the p_T^W measurement at 5.02 and 13 TeV.

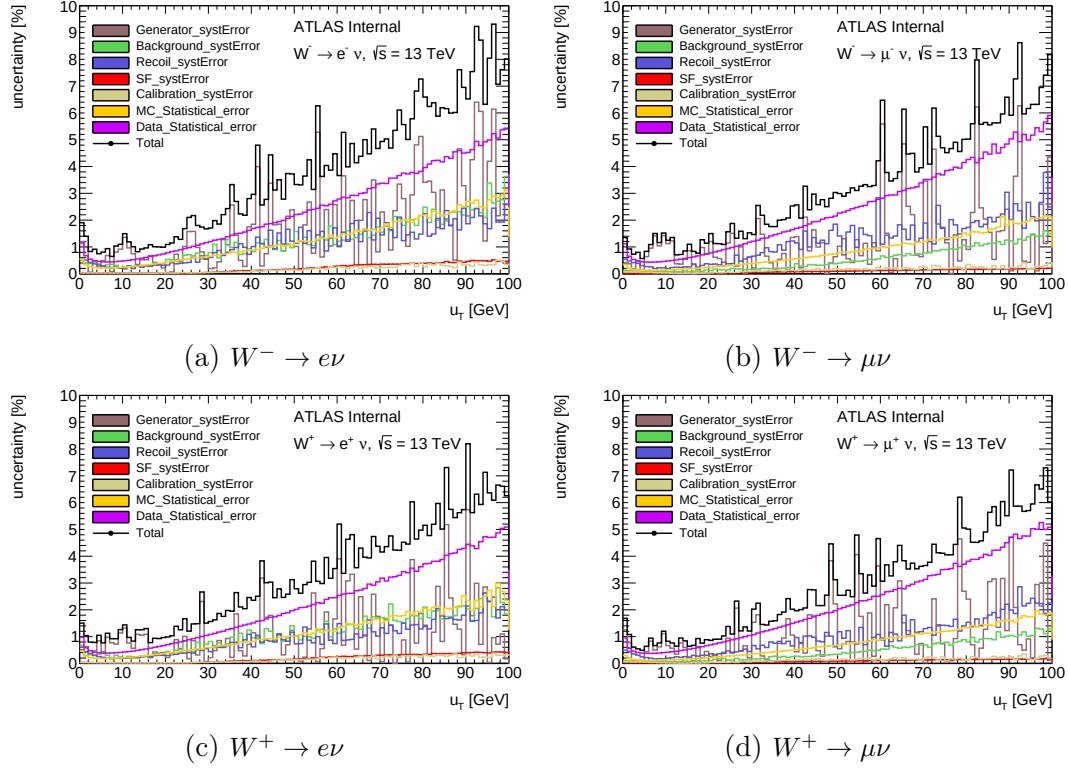


Figure 7.12 – Breakdown of uncertainties on the reconstructed level u_T distributions at $\sqrt{s}=13$ TeV.

Source of uncertainty	W^+ vs. W^-	e vs. μ	5.02 vs. 13 TeV
Data & MC statistics	No	No	No
Lepton calibration & efficiencies	Full	No	Partial
Recoil calibration (Z-based)	Full	Full	No
Recoil calibration ($\sum \bar{E}_T$)	Full	Full	Full
EW & Top backgrounds	Full	Full	Full
Multijet background (statistics)	No	No	No
Multijet background (systematics)	Full	Full	Full
Unfolding bias	No	Full	No
Generator systematic	Full	Full	Full
Luminosity	Full	Full	No

Table 7.2 – A simplified description of the correlation model taken into account when combining the p_T^W spectra and calculating the integrals of fiducial cross-sections.

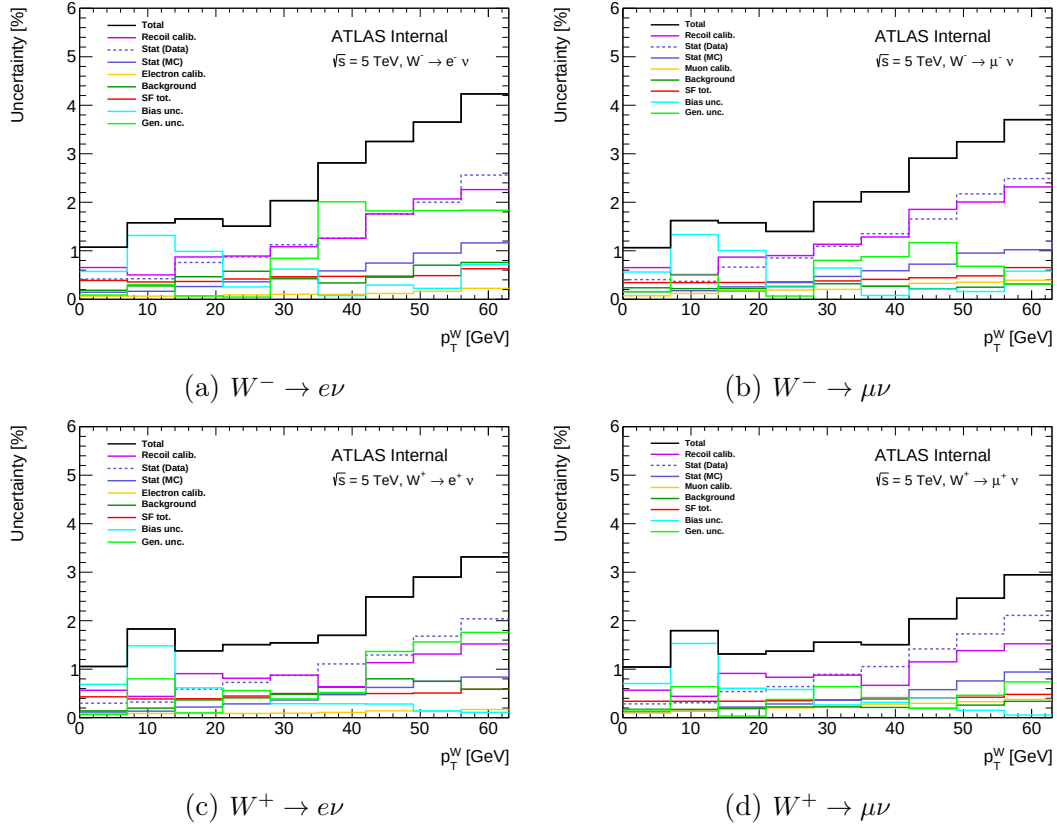


Figure 7.13 – Breakdown of absolute uncertainties on the unfolded level p_T^W distributions at $\sqrt{s} = 5.02$ TeV.

Process	Cross-section and uncertainties
5.02 TeV $W^+ \rightarrow l\nu$	2227.9 ± 2.6 (stat.) ± 7.8 (syst.) ± 23.1 (lumi.)
5.02 TeV $W^- \rightarrow l\nu$	1384.5 ± 2.1 (stat.) ± 5.4 (syst.) ± 14.5 (lumi.)
13 TeV $W^+ \rightarrow l\nu$	4571.4 ± 3.3 (stat.) ± 21.4 (syst.) ± 44.2 (lumi.)
13 TeV $W^- \rightarrow l\nu$	3485.8 ± 2.9 (stat.) ± 17.7 (syst.) ± 34.0 (lumi.)
5.02 TeV $Z \rightarrow ll$	333.0 ± 1.2 (stat.) ± 2.2 (syst.) ± 3.3 (lumi.)
13 TeV $Z \rightarrow ll$	780.3 ± 2.6 (stat.) ± 7.1 (syst.) ± 7.1 (lumi.)

Table 7.3 – Integrated fiducial cross-sections for W^+ , W^- and Z productions in pb at 5.02 and 13 TeV.

Process	Ratio at $\sqrt{s} = 5.02$ TeV [pb]	Ratio at $\sqrt{s} = 13$ TeV [pb]
W^+/W^-	1.611 ± 0.003 (stat.) ± 0.004 (syst.)	1.312 ± 0.001 (stat.) ± 0.003 (syst.)
W^-/Z	4.16 ± 0.01 (stat.) ± 0.05 (syst.)	4.46 ± 0.01 (stat.) ± 0.07 (syst.)
W^+/Z	6.69 ± 0.02 (stat.) ± 0.08 (syst.)	5.84 ± 0.01 (stat.) ± 0.09 (syst.)
$W^\pm/Z$	10.85 ± 0.04 (stat.) ± 0.11 (syst.)	10.31 ± 0.02 (stat.) ± 0.15 (syst.)

Table 7.4 – Integrated cross-sections ratios at 5.02 and 13 TeV.

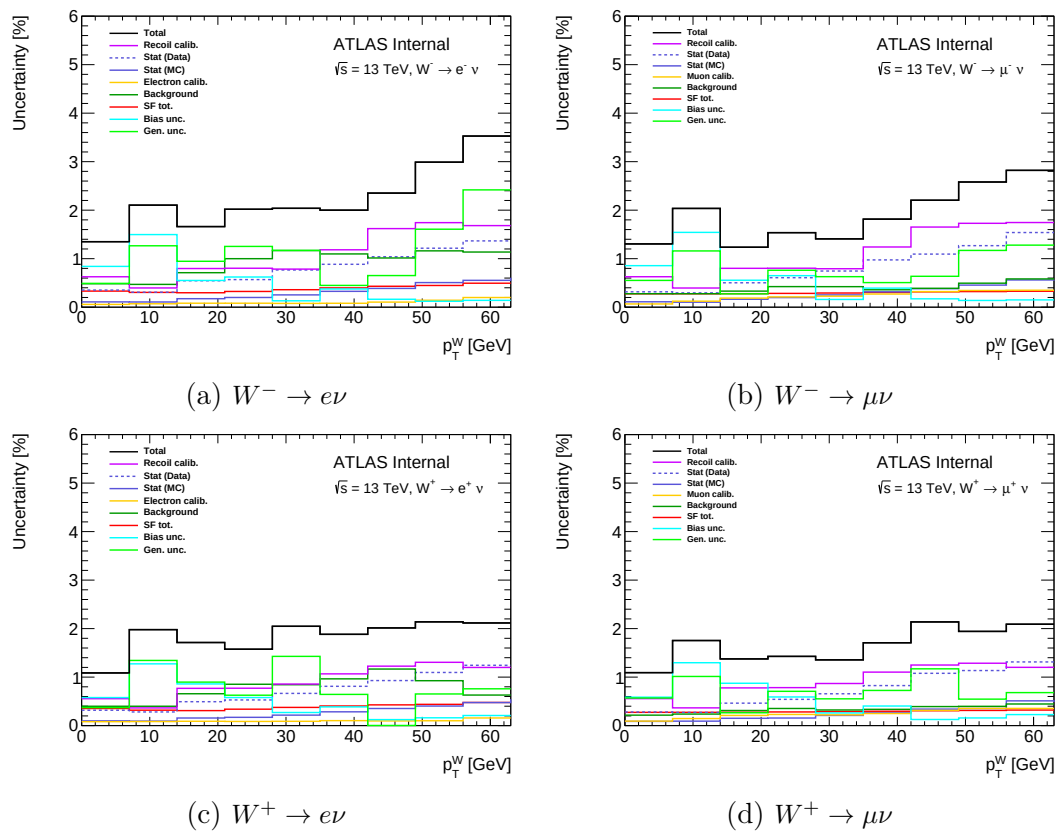


Figure 7.14 – Breakdown of absolute uncertainties on the unfolded level p_T^W distributions at $\sqrt{s} = 13$ TeV.

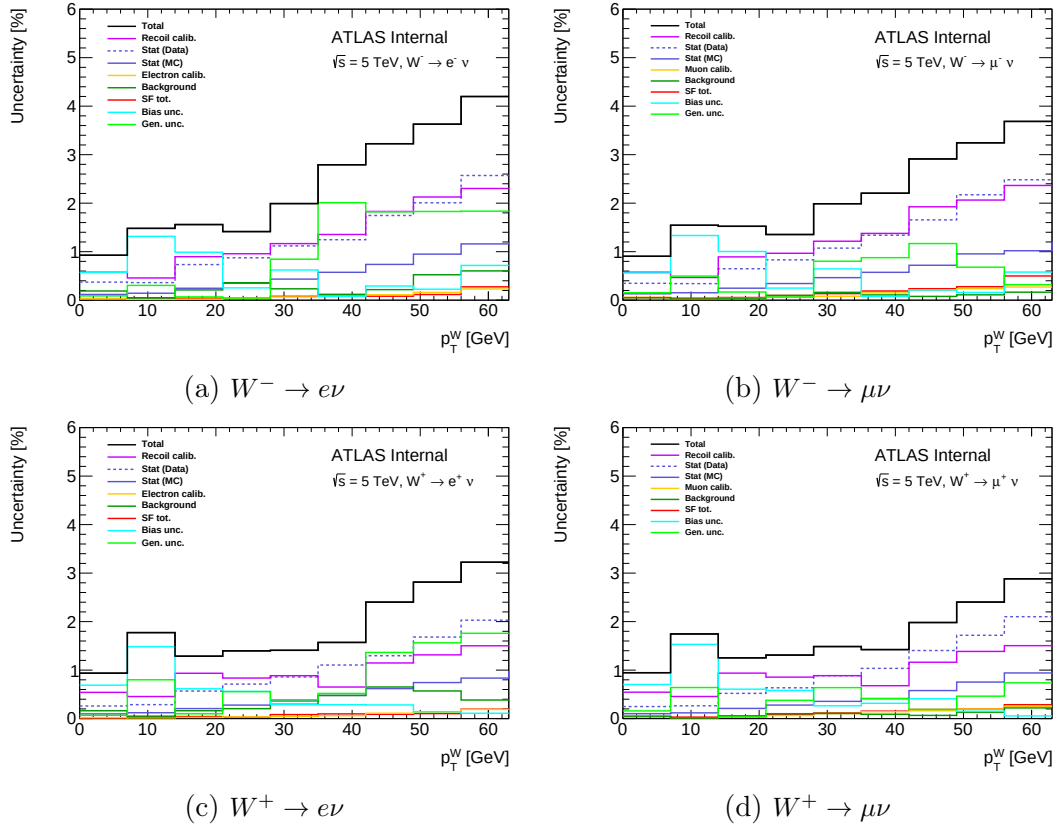


Figure 7.15 – Breakdown of shape uncertainties on the unfolded level p_T^W distributions at $\sqrt{s} = 5.02$ TeV.

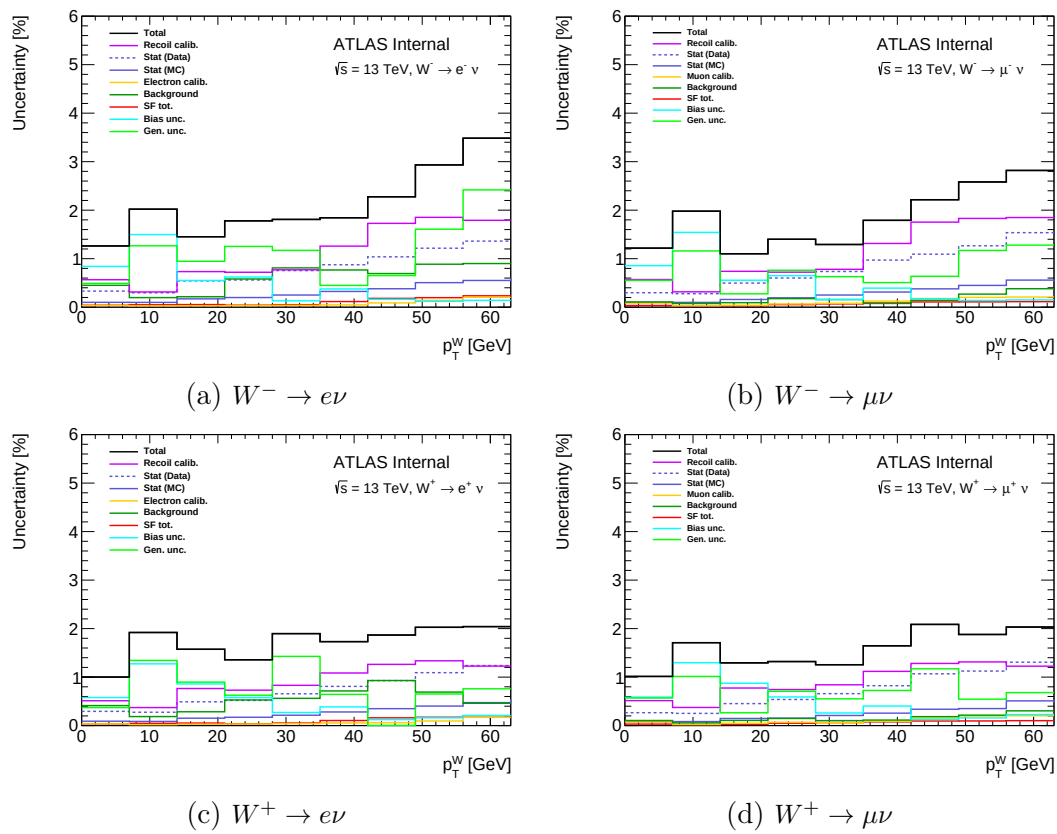
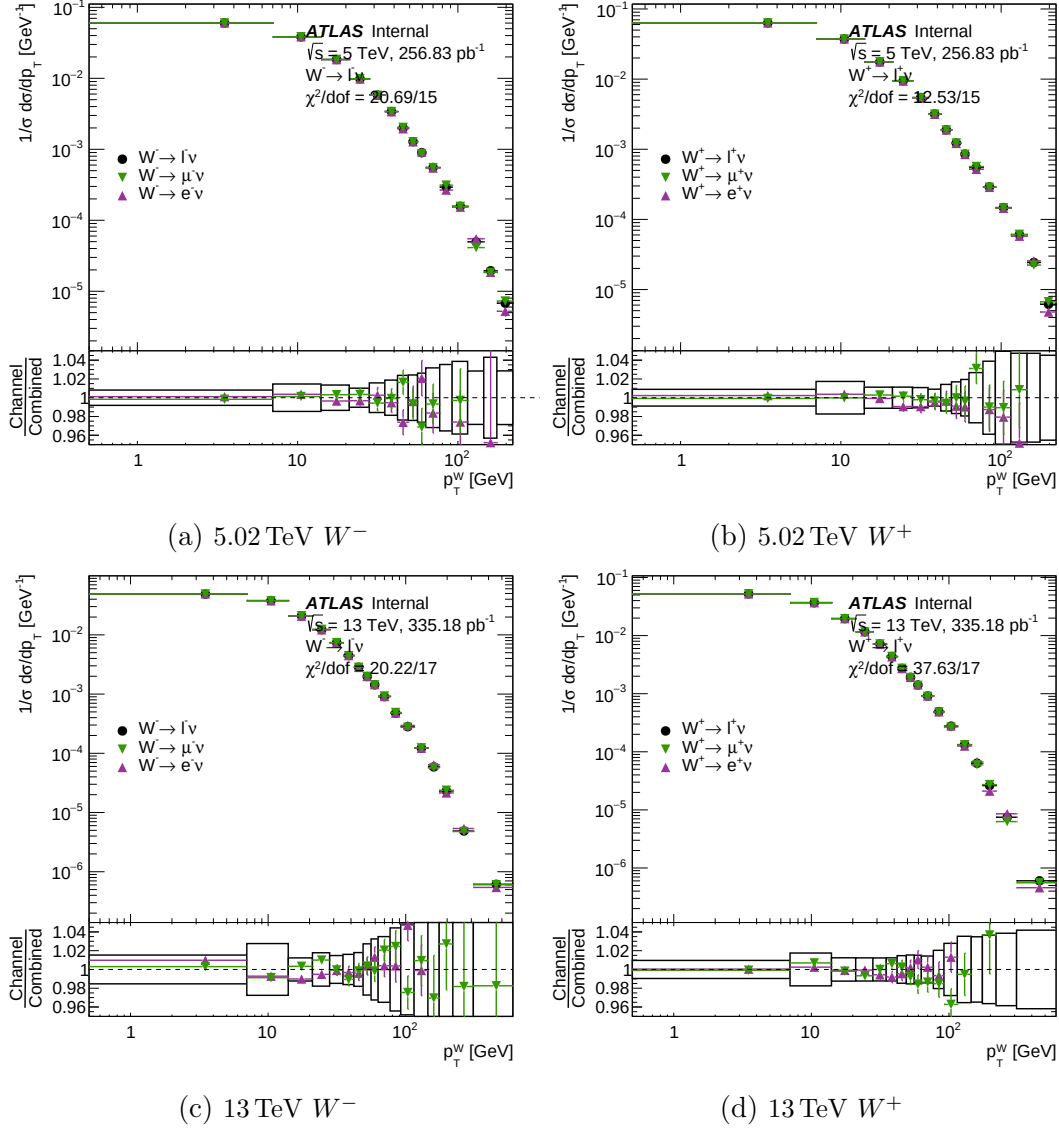
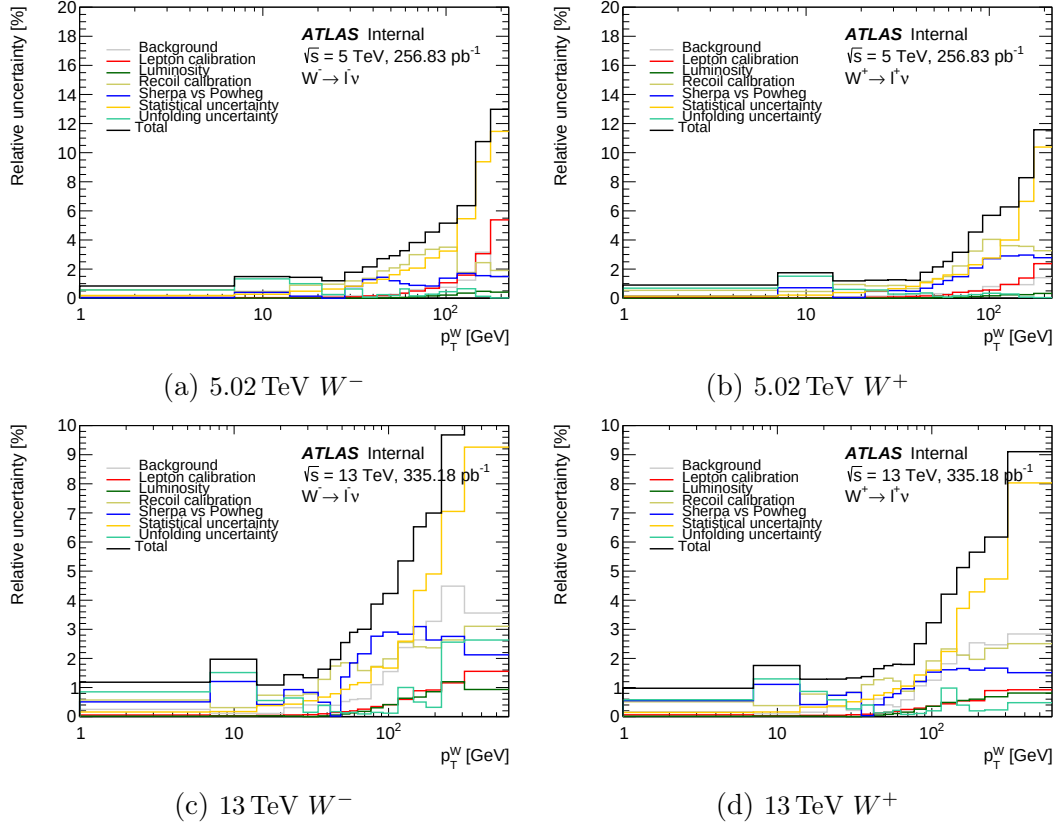
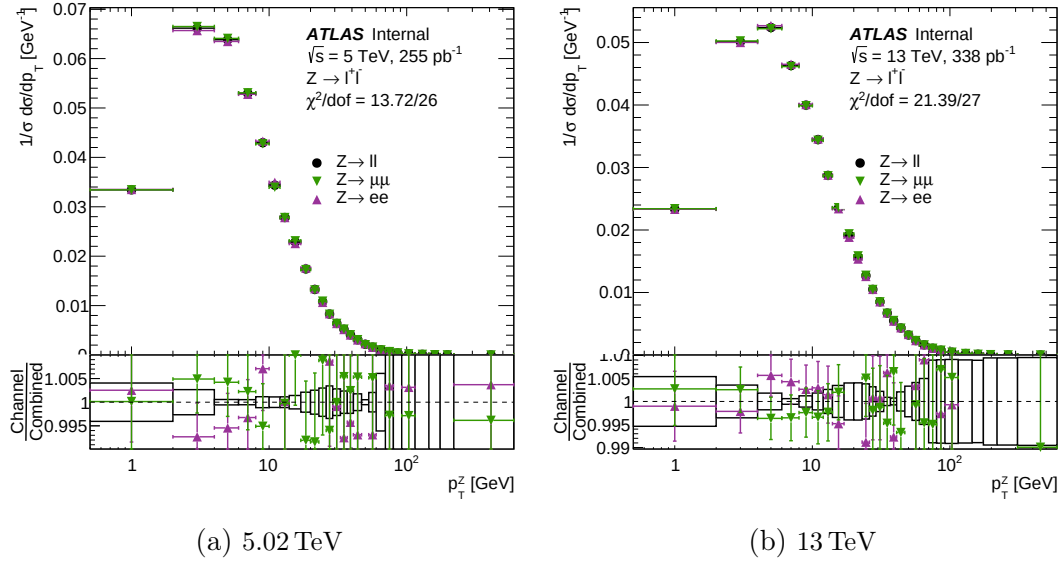


Figure 7.16 – Breakdown of shape uncertainties on the unfolded level p_T^W distributions at $\sqrt{s}=13$ TeV.

Figure 7.17 – Combination of the normalized p_T^W distributions.

Figure 7.18 – Breakdown of the uncertainties for the normalized p_T^W distributions.Figure 7.19 – Combination of the normalized p_T^Z distributions.

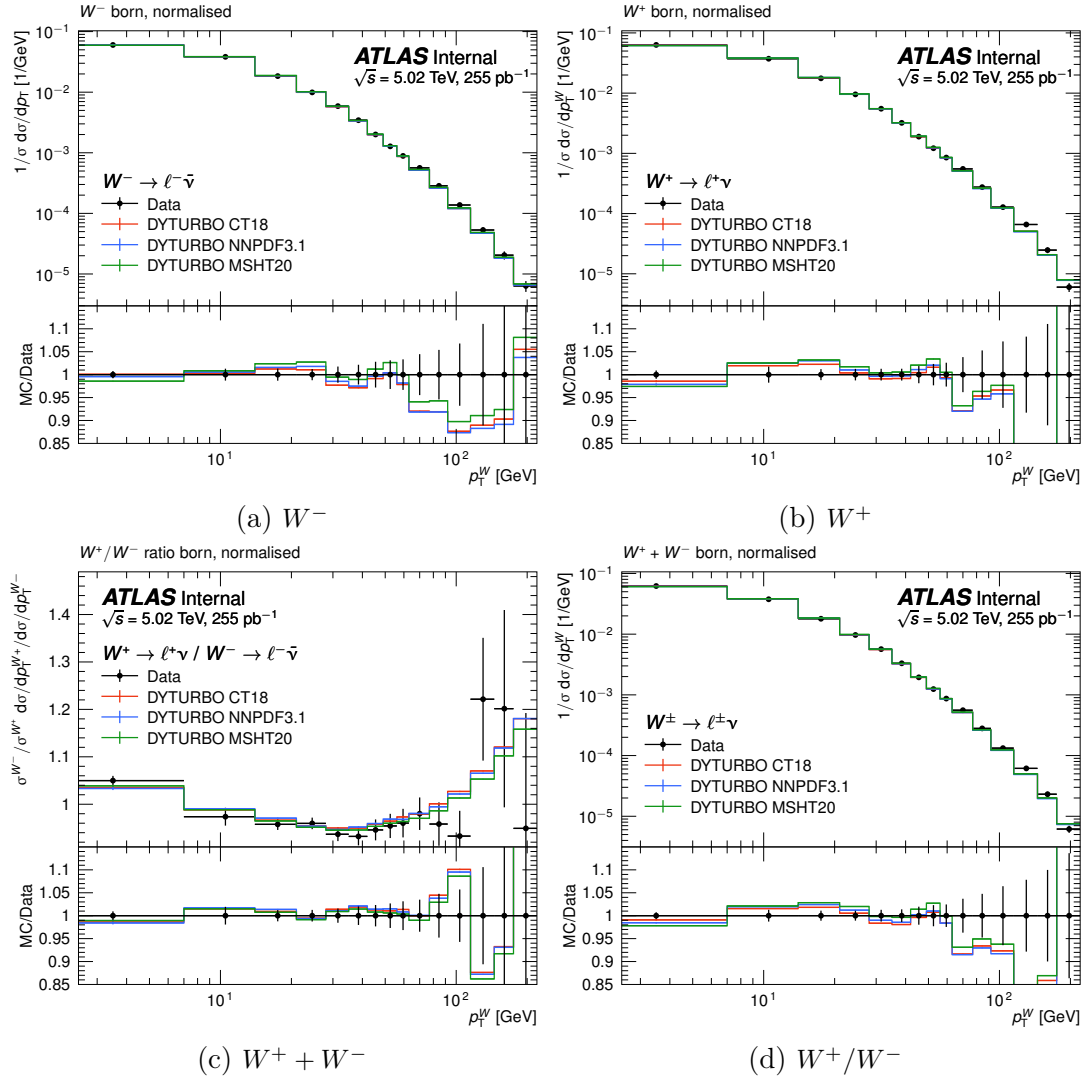


Figure 7.20 – Measurement of normalized p_T^W distributions at 5.02 TeV compared with DYTURBO predictions. The error bars of the data give the total measurement uncertainty.

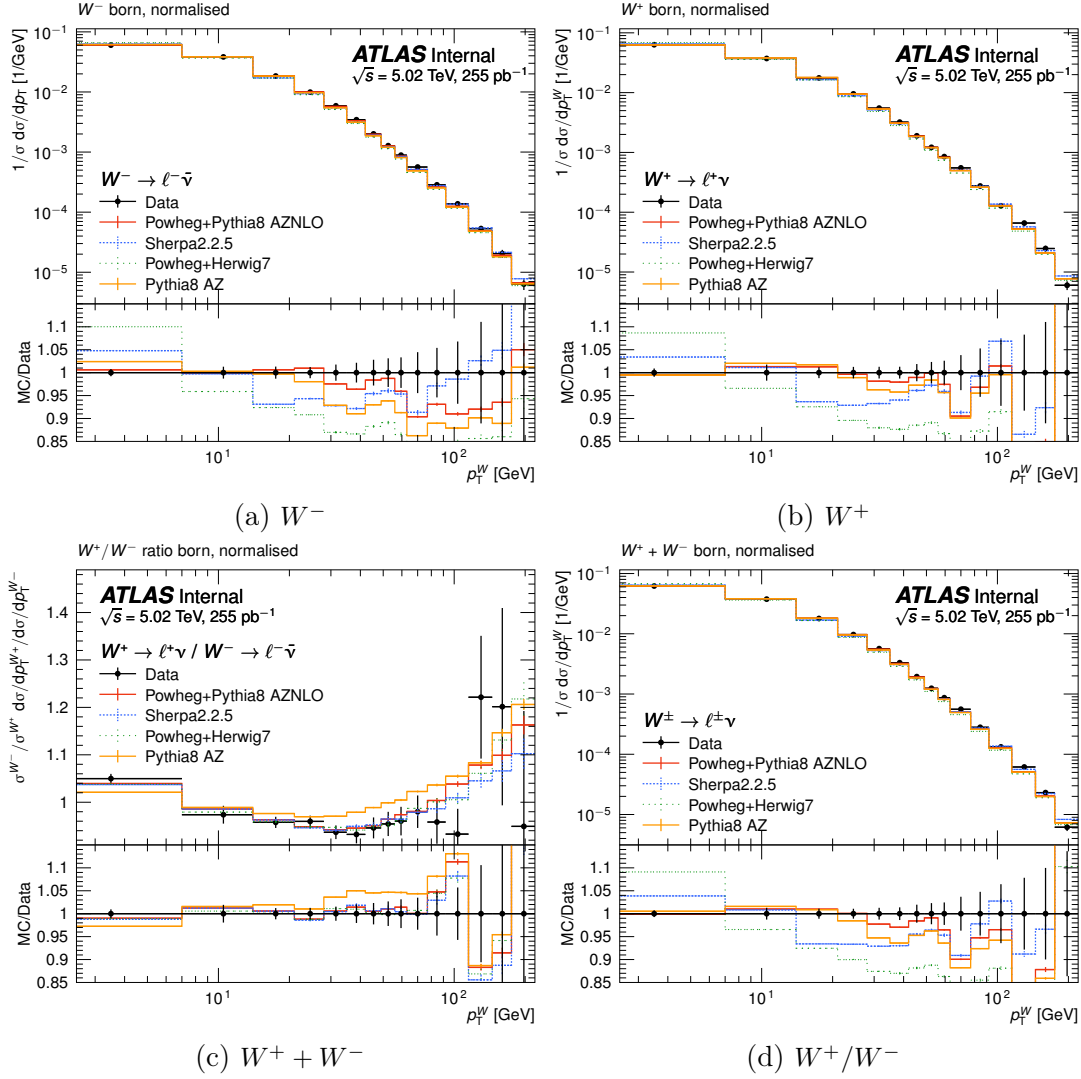


Figure 7.21 – Measurement of normalized p_T^W distributions at 5.02 TeV compared with MC predictions. The error bars of the data give the total measurement uncertainty.

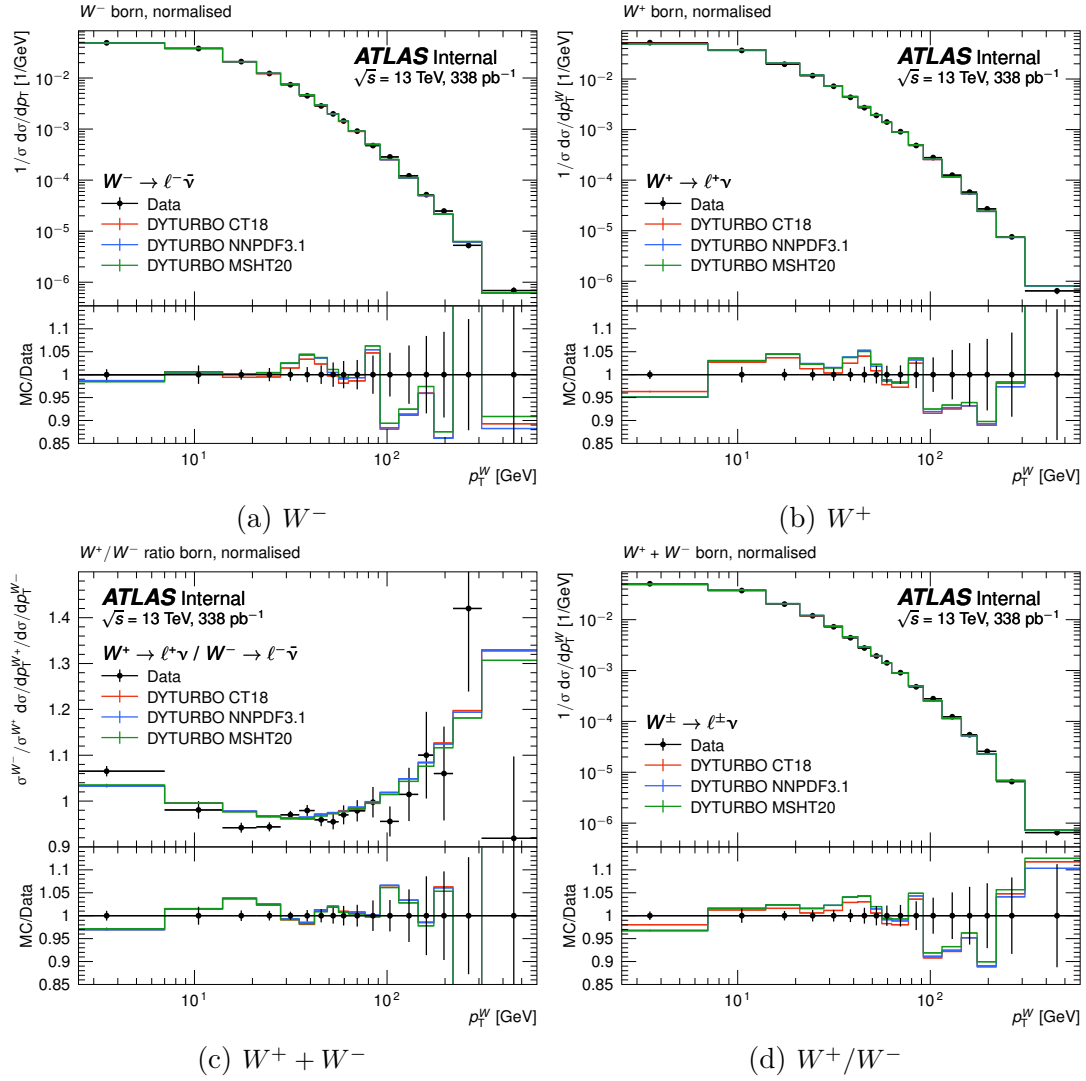


Figure 7.22 – Measurement of normalized p_T^W distributions at 13 TeV compared with DYTURBO predictions. The error bars of the data give the total measurement uncertainty.

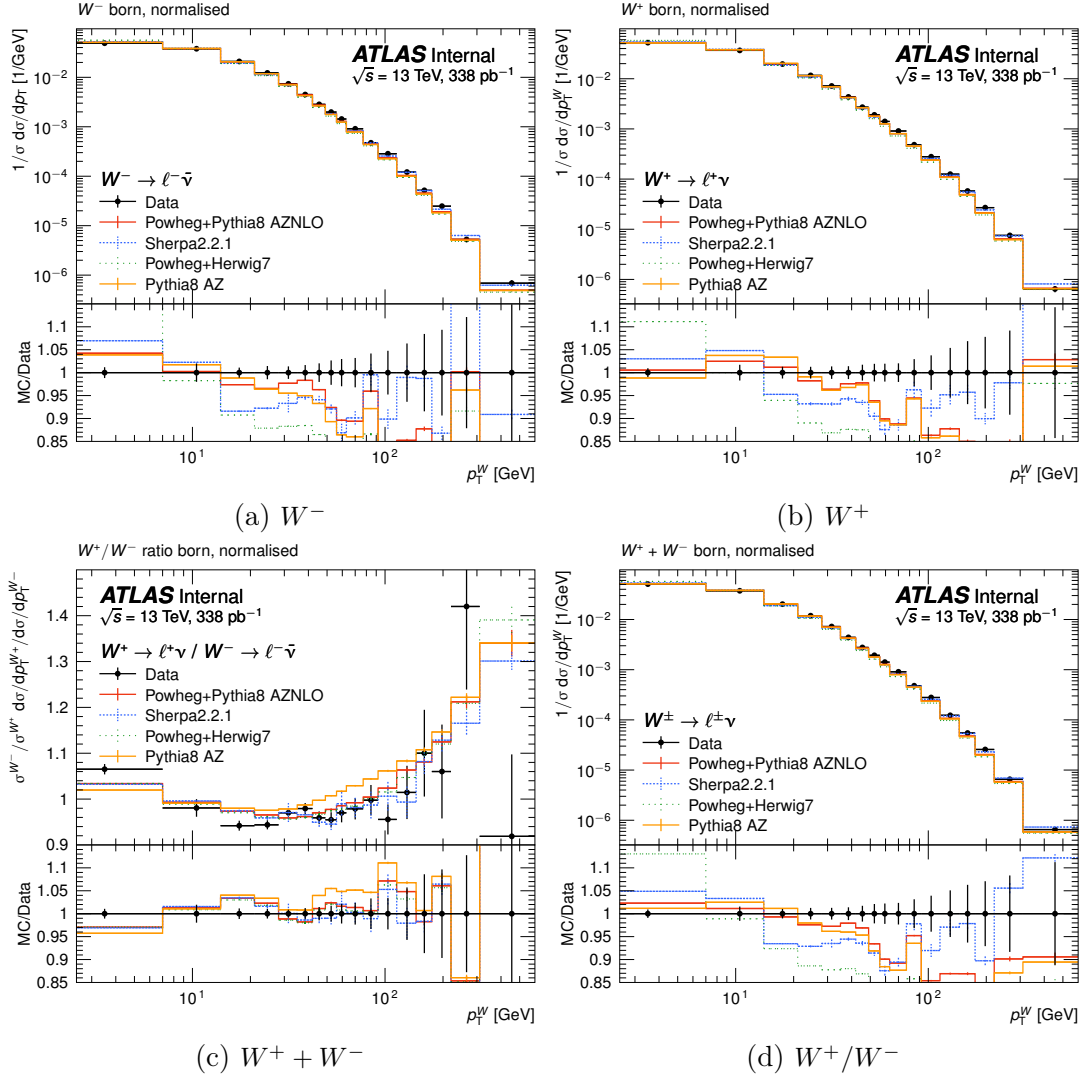


Figure 7.23 – Measurement of normalized p_T^W distributions at 13 TeV compared with MC predictions. The error bars of the data give the total measurement uncertainty.

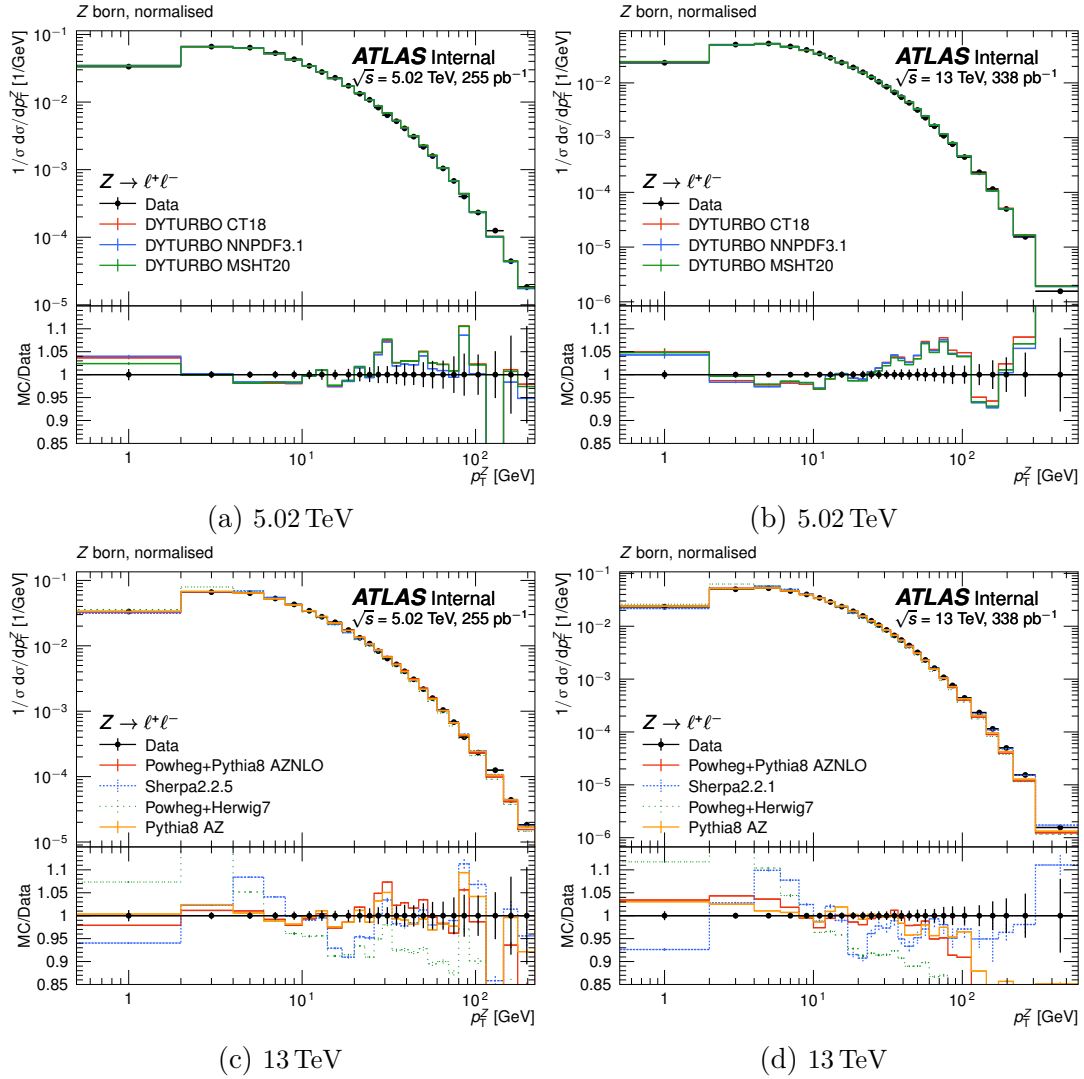


Figure 7.24 – Measurement of normalized p_T^Z distributions with DYTURBO (top) and MC (bottom) predictions. The error bars of the data give the total measurement uncertainty.

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4038 Chapter 8

4039 Measurement of the W-boson mass 4040 with low pile-up data

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The measurement of W-boson mass documented in this work uses the ATLAS Run 2 low pile-up data described in Section 4.3. A dedicated MC simulation campaign was launched to match the special data-taking conditions, as discussed in Section 5.1. The low pile-up W mass measurement shares the same physics modelling corrections as well as detector calibrations with the p_T^W measurement. Dedicated efforts have been spent on the improvement of the statistical analysis for the extraction of m_W from the data. The results presented in this chapter are derived from an old luminosity measurement [1]. The studies will soon be carried out again using the updated luminosity measurement [2], though no significant impact on the results is expected.

8.1 Basic methodology

The extraction of m_W in hadron colliders are based on a template fit, where the impact of hypothetically changing the mass of the boson on the kinematic spectrum is described by a reweighting of the resonance peak [3]:

$$\frac{d\sigma}{dm} \propto \frac{m^2}{(m^2 - m_V^2)^2 + m^4 \Gamma_V^2 / m_V^2} \quad (8.1)$$

with m_V and Γ_V as the assumed values of W-boson mass and width in the simulation, while m stands for the invariant mass of the W-boson. W mass templates with various assumed values of m_W are generated via reweighting the events in the baseline simulation according to Equation 8.1. The W mass template that best describes the data indicates the most probable value of W-boson mass inferred from the data. Considering the sensitivity of the kinematic distributions to m_W , the most common choices of observables used for the extraction of W mass in hadron colliders are lepton p_T and m_T . This can be understood as for W-bosons at rest, the Jacobian edges of these two distributions are $m_V/2$ and m_V , respectively [4].

The conventional approach for conducting the template fit for W mass, known as the offset method, will be briefly recapped in Section 8.1.1. The profile likelihood fit (pLH-fit), as an advanced fitting technique based on the template fit, is applied to the low pile-up W mass analysis, of which the basic concept and notations will be formulated in Section 8.1.2.

8.1.1 The χ^2 offset method

The χ^2 is a test statistic that helps evaluate the discrepancy between the expected frequencies and the observed frequencies. When applied to histograms, it takes the following binned format. For instance, the Neyman's χ^2 is expressed as:

$$\chi^2 = \sum_{i=1}^{N_{bin}} \frac{(n_i - v_i)^2}{\sigma_i^2} \quad (8.2)$$

for which n_i and v_i are the number of entries in the i -th bin of the observed (pseudo) data and of the expectation predicted by the simulation. The bin-by-bin uncorrelated uncertainty showing up in the denominator is a quadratic sum of the statistical uncertainties in both data and MC simulation in that bin, denoted as:

$$\sigma_i \equiv \sigma_i^{datastat.} \oplus \sigma_i^{MCstat.} \quad (8.3)$$

Considering that the kinematic distribution of physics observables, such as p_T^ℓ and m_T , are dependent on the value of W mass in the simulation (m_W), the number of entries in each bin of the kinematic distribution is therefore a function of m_W , and can be expanded into Taylor series in the vicinity of $m_W = m_W^0$, the resonance peak of the W-boson:

$$v_i = v_i \left(\frac{m_W - m_W^0}{m_W^0} \right) = \sum_{j=0}^{\infty} \frac{v_i^{(j)}(0)}{j!} \left(\frac{m_W - m_W^0}{m_W^0} \right)^j \approx v_i(0) + v_i'(0) \left(\frac{m_W - m_W^0}{m_W^0} \right) \quad (8.4)$$

The change of W mass in the template (in general, $m_W - m_W^0 \approx < 100$ MeV) is way smaller than the assumed central value of W mass ($m_W^0 \approx 80.400$ MeV). This allows to take the linear approximation from the Taylor series, such that the Neyman's χ^2 becomes a 2nd-order polynomial function of m_W :

$$\chi^2(m_W) = \sum_{i=1}^{N_{bin}} \frac{\left[n_i - v_i(0) - v_i'(0) \left(\frac{m_W - m_W^0}{m_W^0} \right) \right]^2}{\sigma_i^2} \quad (8.5)$$

The χ^2 between data and W mass templates are computed at different values of assumed W mass and fitted by a parabola. The minimum of the curve is the most probable value of W mass $\hat{m}_W$ ($\chi^2(\hat{m}_W) = \chi_{Min}^2$), while the difference between $\hat{m}_W$ and the value of m_W satisfying

$$\chi^2(m_W) = \chi^2(\hat{m}_W \pm \Delta_{m_W}^{stat.}) = \chi_{Min}^2 + 1 \quad (8.6)$$

is the statistical uncertainty of the measurement ($\Delta_{m_W}^{stat.}$). The systematic uncertainties are accounted for by replacing the (pseudo) data in the above procedure with the systematic variation and determining the new minimum of the parabola fit. The

offset between $\hat{m}_W$ and the new minimum is taken to be the uncertainty ($\Delta_{m_W}^{syst.p}$) arising from corresponding systematic source p .

In the end, the result of the measurement using the χ^2 offset method, in terms of central value and the associated uncertainties, is presented by:

$$\hat{m}_W \pm \Delta_{m_W}^{stat.} \oplus \sum_{\oplus p} \Delta_{m_W}^{syst.p} \quad (8.7)$$

The total uncertainty is a quadratic sum of statistical uncertainty and systematic uncertainties.

8.1.2 Profile likelihood fit

Despite the low CPU time consumption and the good numerical precision of the χ^2 offset method (typically at sub-MeV level for the W mass measurement), the biases caused by systematics are not corrected in the probability model, therefore leading to a less accurate central value. Moreover, even though the measurement of W mass has sensitivity to the systematics to some extent, none of the systematic uncertainties is constrained by the data in the offset method, which is a shortcoming especially when the measurement is dominated by systematics.

A more powerful tool for the statistical analysis, also following the principle of the template fit, would be the binned profile likelihood fit (pLH-fit), for which the probability density function is parameterized by the parameter of interest (POI) and a series of nuisance parameters (NPs) to account for the normalization as well as the effect of systematics [5]:

$$\mathcal{P}(n_i, a_p | \alpha_{mW}, \alpha_p, \phi_{CME}, \gamma_i) = \prod_{i \in bins} Pois(n_i | v_i) \prod_{p \in syst. + \Gamma} f_p(a_p | \alpha_p) \quad (8.8)$$

- n_i is the observed number of events in the i -th bin of the histogram known from the (pseudo) data.
- v_i is the total predicted number of events in the i th bin of the histogram.
- α_{mW} is the parameter of interest (POI) that can be translated to m_W .
- α_p is the nuisance parameter in the probability model parameterizing the effect of systematic uncertainty. Conventionally, $\alpha_p = \pm 1$ indicates the distortion of ± 1 standard deviation on the kinematic spectrum brought by the systematic uncertainty.
- ϕ_{CME} is the free floating normalization factor for the prediction of the signal. The normalization factors are introduced mainly to absorb the bias due to the effect of missing higher-order QCD calculation on the signal cross-section. The difference between the calculations of W^+ and W^- productions is considered to be weakly dependent on higher-order QCD correction. Hence one single ϕ_{CME} is shared by all the W-boson channels at each center-of-mass energy.
- γ_i is the parameter accounting for MC statistical uncertainty. The details of these parameters will be elaborated in Section 8.2.2. The collection of all γ_i is denoted as Γ .
- a_p is the central value of the auxiliary measurement for the systematic source p , known as the global observable. The prior knowledge of this systematic source is embodied by the form of the constraint term f_p . The constraint term is usually a normal distribution with the conventional definition of NP,

if nominal expectation for the main measurement has been corrected by the results of the auxiliary measurements (e.g. detector calibration). The overall likelihood $\mathcal{P}$ is a product of Poisson likelihood for each bin and the constraint terms for systematics. The prediction of the Poisson likelihood is parameterized by a template morphing that consists of a polynomial interpolation and linear extrapolation:

$$I_{poly|lin.}(\alpha; I^0, I^+, I^-) = \begin{cases} \alpha(I^0 - I^-) & \alpha \leq -1 \\ \sum_{i=1}^6 c_i \alpha^i & |\alpha| < 1 \\ \alpha(I^+ - I^0) & \alpha \geq 1 \end{cases} \quad (8.9)$$

with a set of coefficients c_i making the interpolation satisfy the continuity in the value $\eta(\alpha = \pm 1)$, the first derivative $d\eta/d\alpha|_{\alpha=\pm 1}$ and the second derivative $d^2\eta/d\alpha^2|_{\alpha=\pm 1}$ at the boundaries $\alpha = \pm 1$.

For multiplicative parameterization such as luminosity, the template morphing is written explicitly as:

$$\eta(\boldsymbol{\alpha}) = 1 + \sum_{p \in syst.} I_{poly|lin.}(\alpha_p; 1, \eta_p^+, \eta_p^-) \quad (8.10)$$

Here $\eta_p^\pm$ correspond to the deviation of η from unity brought by ± 1 standard deviation of systematic source p . And similarly for the additive parameterization, for example the shape of a distribution, the morphing takes the form:

$$h_i(\boldsymbol{\alpha}) = h_i^0 + \sum_{p \in syst.} I_{poly|lin.}(\alpha_p; h_{p,i}^0, h_{p,i}^+, h_{p,i}^-) \quad (8.11)$$

where h_i^0 is the content of the i -th bin at nominal value, α_p is the NP that parameterizes the effect of systematics, $h_{p,i}^\pm$ are the change in bin content brought by ± 1 standard deviation of the systematic source p . Conventionally, one sets $h_{p,i}^0 = h_i^0$.

Unlike fitting the background subtracted data in the χ^2 offset method, the Poisson likelihood for each bin adopts the form of an extended p.d.f., taking into account the contribution from background in the total prediction. Specifically, in the W mass fit one has:

$$v_i = \phi_{CME} \cdot v_i^{sig} + \alpha_{lumi.} \cdot v_i^{bkg} + v_i^{MJ} \quad (8.12)$$

where v_i^{sig} is the MC prediction of the signal process, v_i^{bkg} the prediction of the simulated background and v_i^{MJ} the data-driven estimation of the multijet background described in Section 5.4. v_i^{sig} is parameterized by the effect of systematic uncertainties as well as that of the hypothetical change in W-boson mass, the parameterization of both sources follows Equation 8.11. v_i^{bkg} can be varied by the cross-section uncertainty of each simulated background process, of which the uncertainty templates are also morphed according to Equation 8.11. The uncertainties arising from the data-driven method are assigned to v_i^{MJ} , with the parameterization of Equation 8.11. The normalization of the signal prediction is reflected by free floating parameter ϕ_{CME} , while the luminosity measurement parameterized by $\alpha_{lumi.}$ enters the normalization for the simulated background. The template morphing for the luminosity uncertainty takes the form of Equation 8.10. No normalization is applied to the result of data-driven estimation of the multijet background.

The setup and the execution of the profile likelihood fit used in this work are handled by HistFactory [5] interfaced to RooFit [6]. The most probable values of the fit parameters are inferred from the (pseudo) data by minimizing the Negative

Logarithm of the Likelihood (NLL), which is an approach mathematically equivalent to maximizing the likelihood but with a considerable numerical advantage. The calculation of the first derivative of the NLL leads to the most probable values of the fit parameters, requiring:

$$\frac{\partial NLL}{\partial x} = 0 \quad (8.13)$$

with x looping over all fit parameters.

The calculation of the errors on the fit parameters consists of two steps. Firstly, the Hessian matrix $\mathcal{H}$ (the matrix of secondary derivatives) of the NLL with respect to the fit parameters is computed:

$$\mathcal{H} = \left(\frac{\partial^2 NLL}{\partial x \partial y} \right) \quad (8.14)$$

with x and y independently looping over all fit parameters. The post-fit uncertainties of the POI and the NPs can be assessed by the retrieving the diagonal elements of the inverse of the Hessian matrix $\mathcal{H}^{-1}$, producing what is referred to as the Hesse Error in the context of RooFit. In the second step, after the derivation of the Hesse Error, the ± 1 standard deviation of a given fit parameter can be found by performing a scan of the likelihood contour along this parameter and solving the two points fulfilling $\Delta NLL = 0.5$. This is similar to solving the Equation 8.6 in the χ^2 offset method. But in a profile likelihood fit, this formula is applicable to all fit parameters. The uncertainty derived in the second step is denoted as the Minos Error. The computation of the Minos Error is more time consuming than the Hesse Error.

The Hesse Errors are symmetric and are only the approximate post-fit errors of the fit parameters, since they are computed from the secondary derivative at the minimum of the NLL and dependent on the minimization path being taken. In most cases the Hesse Errors are lower than the true errors on the fit parameters, meaning that some validations might be necessary for one to understand if the values are reliable. In contrast, the Minos Errors are obtained from the scan of the likelihood contour beyond the minimum. Therefore, the Minos Errors are allowed to be asymmetric and are expected to better reflect the true errors on the parameters. Both the Hesse Error and the Minos Error on the POI are estimates of the total uncertainty on POI, including statistical uncertainties and systematic uncertainties. In order to ensure a reliable total uncertainty in W mass with optimized usage of computational resources, the Minos Error is always used for the POI in the profile likelihood fit, while the Hesse Error can be considered to attain a reasonable precision for the errors on the NPs.

Compared to the χ^2 offset method, the profile likelihood fit not only enables corrections to the small biases in the measurement caused by systematics, but also allows to constrain the systematic uncertainties with the statistical power of the dataset. The downsides of the profile likelihood fit and the possible remedies to these known issues are discussed in Section 8.2.

8.1.3 Configuration of the fit for m_W measurement

The full probability model of the low pile-up W mass fit includes 8 decay channels. The main kinematic distribution is chosen to be m_T due to its better sensitivity to W mass in the low pile-up datasets. A 2D categorization in u_T and η_ℓ is implemented in the probability model to facilitate the constraint of systematic uncertainties. The

optimization of range and of the bin width of m_T and u_T has been studied in the 5 TeV $W^- \rightarrow e\nu$ channel, only taking into account data statistics. Only the categorization in u_T is involved in this study. The baseline is set to be: $60 < m_T < 100$ GeV with 2 GeV bin width, $u_T < 15$ GeV with 5 GeV bin width. Each time, one of the following four quantities is varied to see its impact on the total uncertainty, as shown in Figure 8.1: (a) Range of m_T , (b) Bin width of m_T , (c) Range of u_T and (d) Bin width of u_T . In the end, the combination of $60 < m_T < 100$ GeV with 2 GeV bin width and $u_T < 25$ GeV with 5 GeV bin width is chosen for the input of the fit. Finer granularity or extended fit range brings marginal improvement to the statistical uncertainty. The categorization in lepton η is kept consistent with the ATLAS 7 TeV measurement [7] in general, mainly driven by the acceptance and performance of the detector systems. Therefore, the alignment of η_e and η_μ bins is unnecessary:

- $|\eta_e|$: [0, 0.6, 1.2, 1.8, 2.47]
- $|\eta_\mu|$: [0, 0.8, 1.4, 2.0, 2.4]

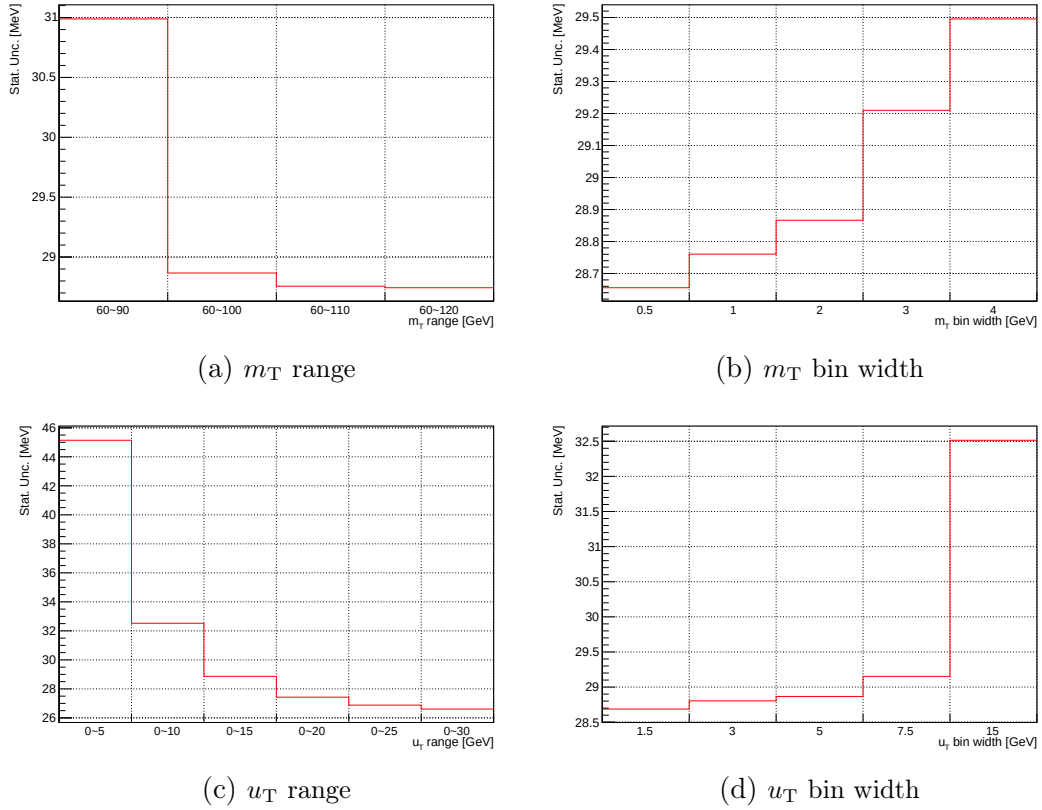


Figure 8.1 – The optimization of fit range and bin width for m_T and u_T using the signal MC in 5 TeV $W^- \rightarrow e\nu$ channel.

The 2D categorization effectively makes the fit a 3D measurement in (m_T, u_T, η_ℓ) . The full information on the kinematic distribution is presented by the unrolled m_T distribution, where the m_T spectra in different categories are concatenated into a long 1D histogram. The preliminary control plots of the unrolled m_T distribution are shown from Figure 8.2 and to Figure 8.5, where the total prediction is normalized to the data integral. Only the statistical uncertainties are presented in the control plots. The value of m_W is blinded in the simulation. The categorization in the plots first

goes through the 4 η_ℓ categories in the first u_T category, then the 4 η_ℓ categories in the next u_T categories, until the 4 η_ℓ categories in the last u_T categories. Neither the data scaling in Section 8.2.2 nor the reconstructed level PDF reweighting in Section 8.3.3 has been applied to the unrolled m_T distributions in the control plots.

Clear trends of discrepancy between data and simulation have been observed in the control plots. Some trends of discrepancy are repeated in different categories and shared by multiple channels, which might indicate that further improvement of physics modelling is necessary.

8.2 The improvement of profile likelihood fit

Despite clear advantages over the χ^2 offset method in terms of both accuracy (central value) and precision (total uncertainty), the use of profile likelihood fit is largely limited by two factors: CPU time consumption and the treatment of statistically correlated observables. The improvement for the first limitation is discussed in Sections 8.2.1 and 8.2.2. A few possible solutions related to the second limitation are mentioned in Section 8.2.3.

8.2.1 Reduction of nuisance parameters

Due to the calculation of the Hessian matrix during the minimization, the CPU time consumption goes up proportionally to the square of the number of fit parameters. Empirically, a fit with several hundreds of fit parameters takes about a few hours for a single CPU core to complete the minimization. An even larger number of fit parameters in the probability model will make the computation extremely expensive. Therefore, dedicated efforts need to be spent in studying the possible schemes of NP reduction in the probability model.

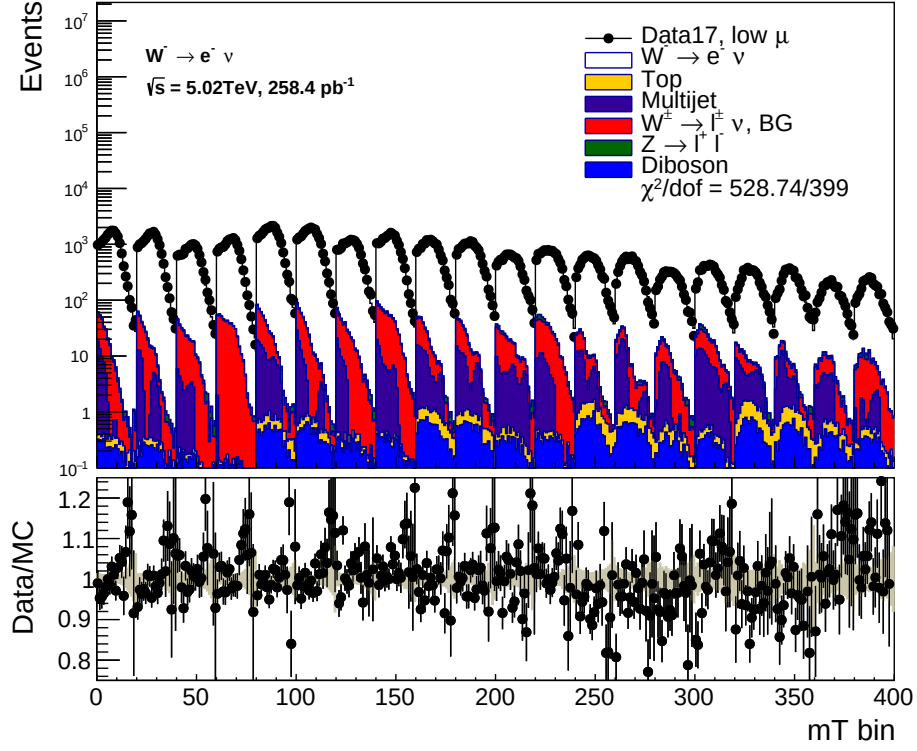
The large number of NP sources mainly originates from the experimental systematics such as the uncertainties related to the lepton selection efficiencies. However, instead of injecting each source of uncertainty one by one to the profile likelihood fit, a good approximation can be achieved by replacing the individual uncertainty variations with the error vectors of the corresponding covariance matrix. For the binned kinematic distribution of a given observable, the covariance matrix of systematics is written as:

$$C_{ij}^{syst} = \sum_p C_{ij}^p = \sum_p (H_i^p - H_i^{nom}) (H_j^p - H_j^{nom}) = \Sigma^T \Sigma \quad (8.15)$$

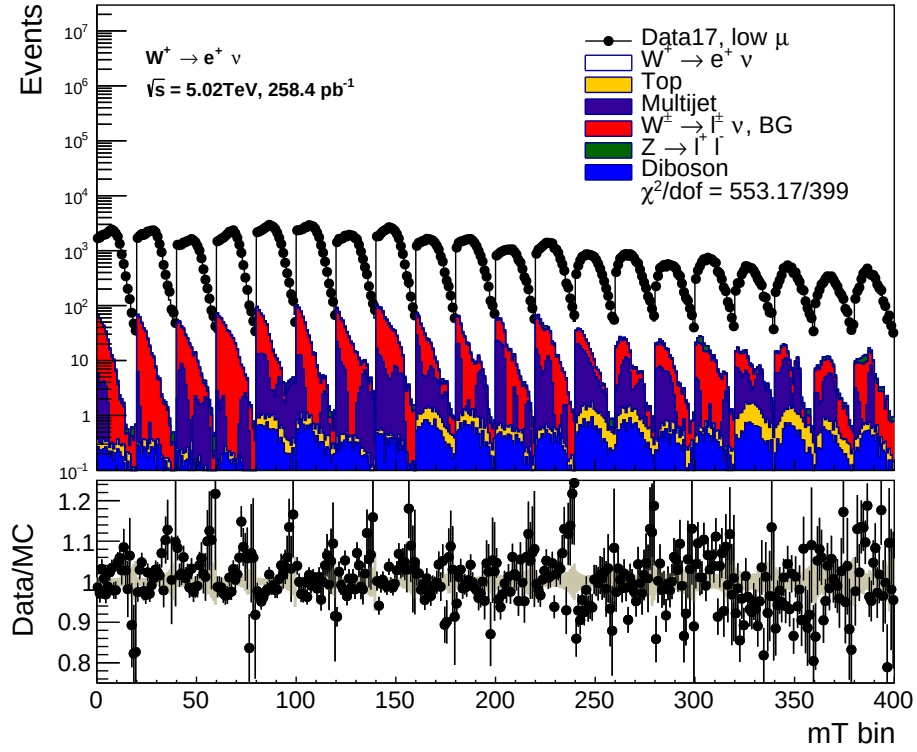
where i, j indicate the indices of bins ($i, j=1, 2, 3 \dots N_{bins}$) of the kinematic distribution, $(\Sigma)_i^p = H_i^p - H_i^{nom}$, H^{nom} is the nominal spectrum and H^p is the systematic variation. By performing the diagonalization of the matrix C^{syst} , one obtains a set of error vectors Γ . The number of the error vectors is equal to the rank of C^{syst} and is no more than the number of bins of the distribution:

$$\Sigma^T \Sigma = C^{syst} = \Gamma^T \Gamma = \left(\sqrt{D} X_L \right)^T \left(\sqrt{D} X_L \right) \quad (8.16)$$

D is a diagonal matrix filled with N_{bins} eigenvalues of C^{syst} . X_L is made of N_{bins} rows of eigenvectors. Denoting $Z = \Sigma \Gamma^{-1}$, the original NPs α_p are transformed to a new

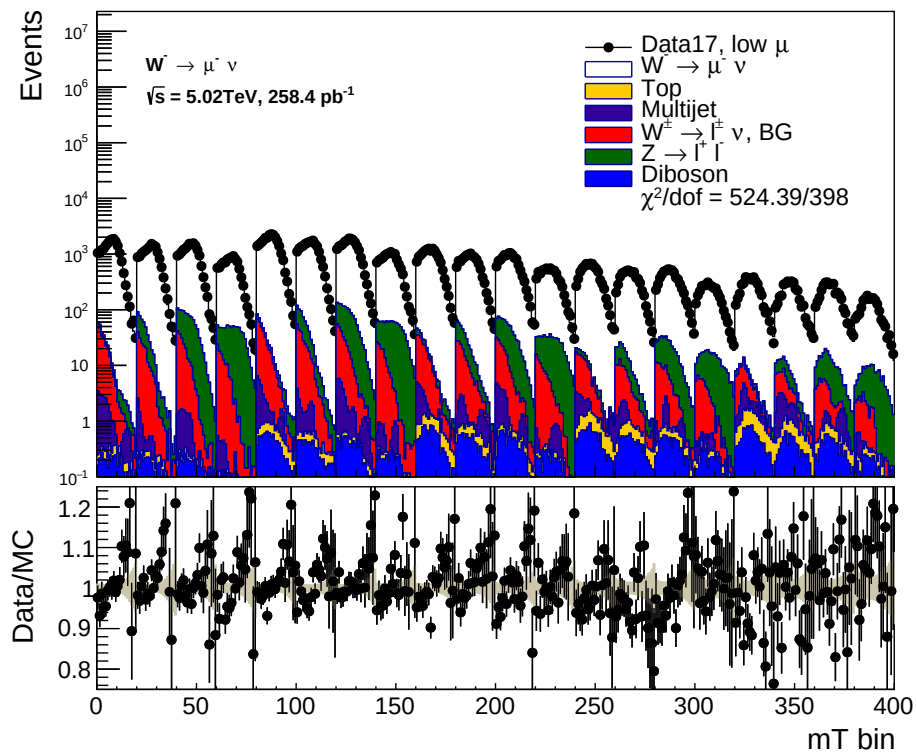


(a)

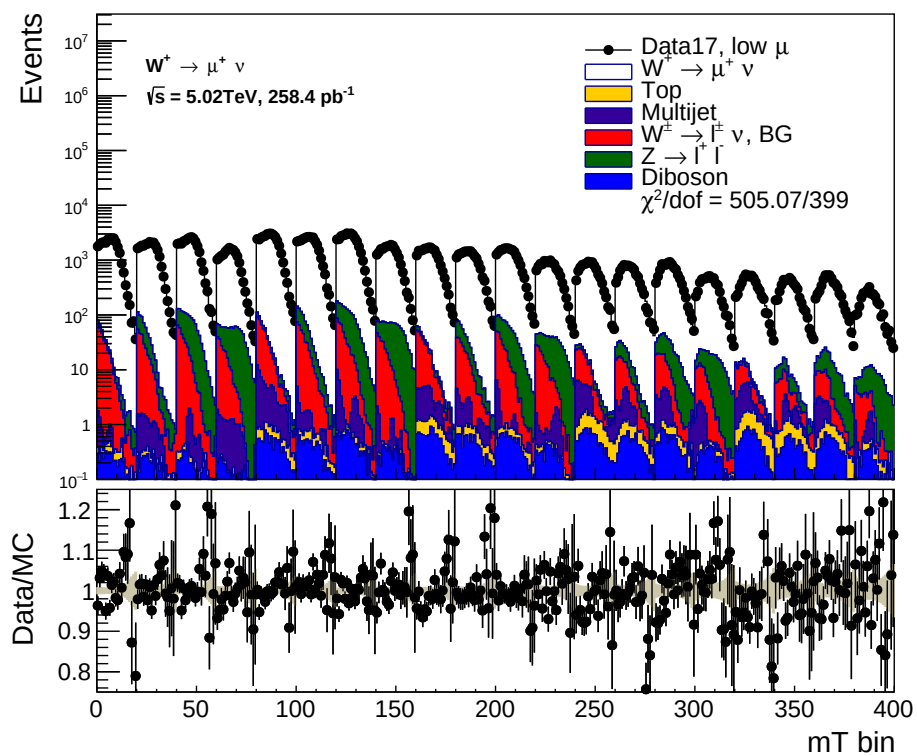


(b)

Figure 8.2 – The control plots of unrolled m_T spectra in 5 TeV electron channels.

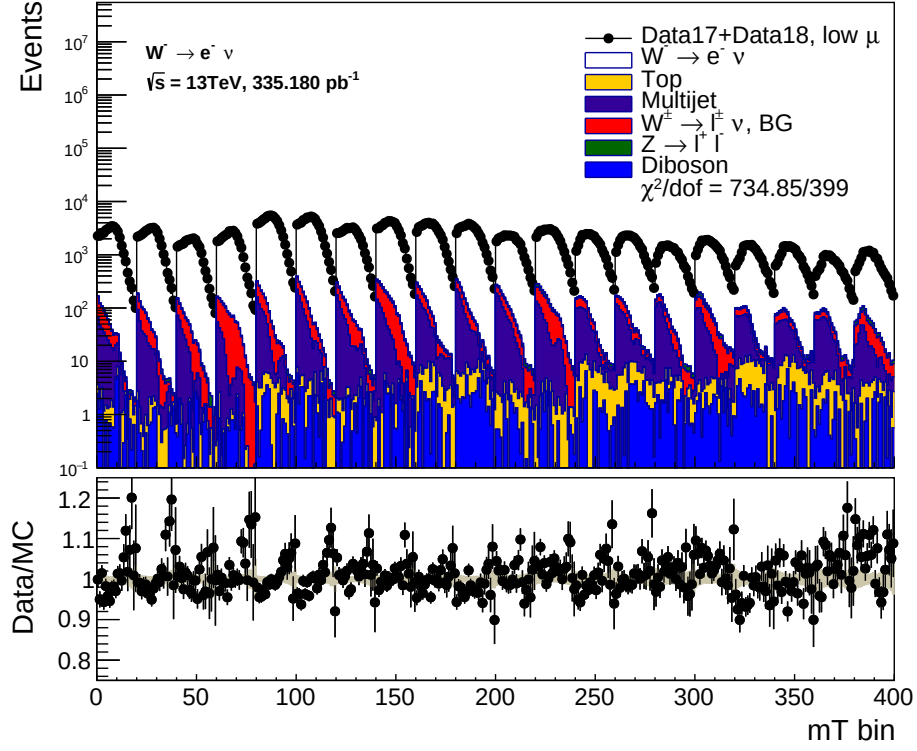


(a)

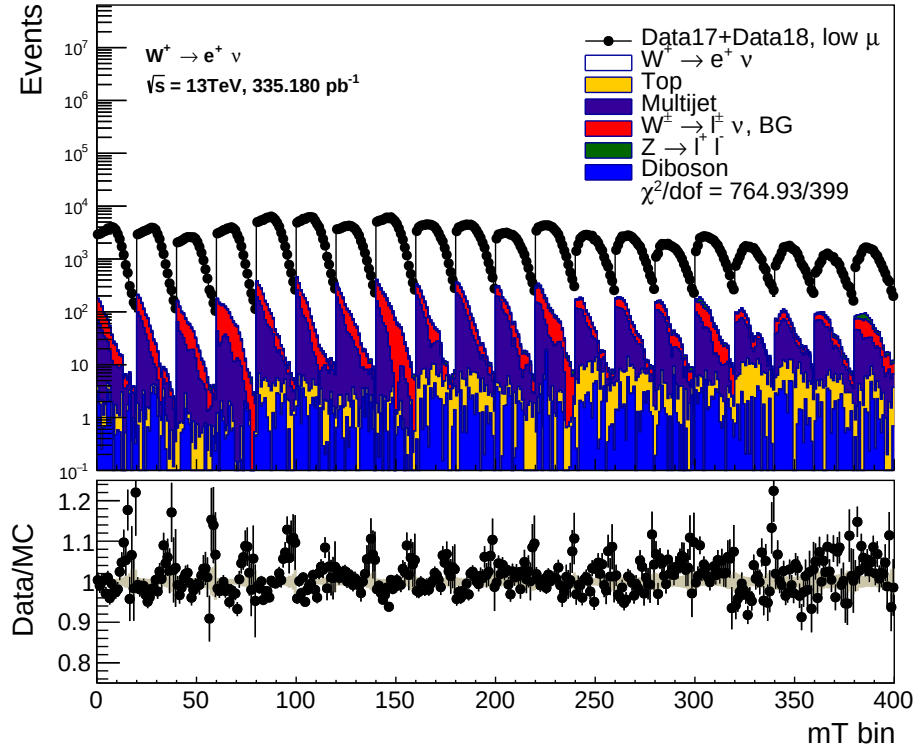


(b)

Figure 8.3 – The control plots of unrolled m_T spectra in 5 TeV muon channels.

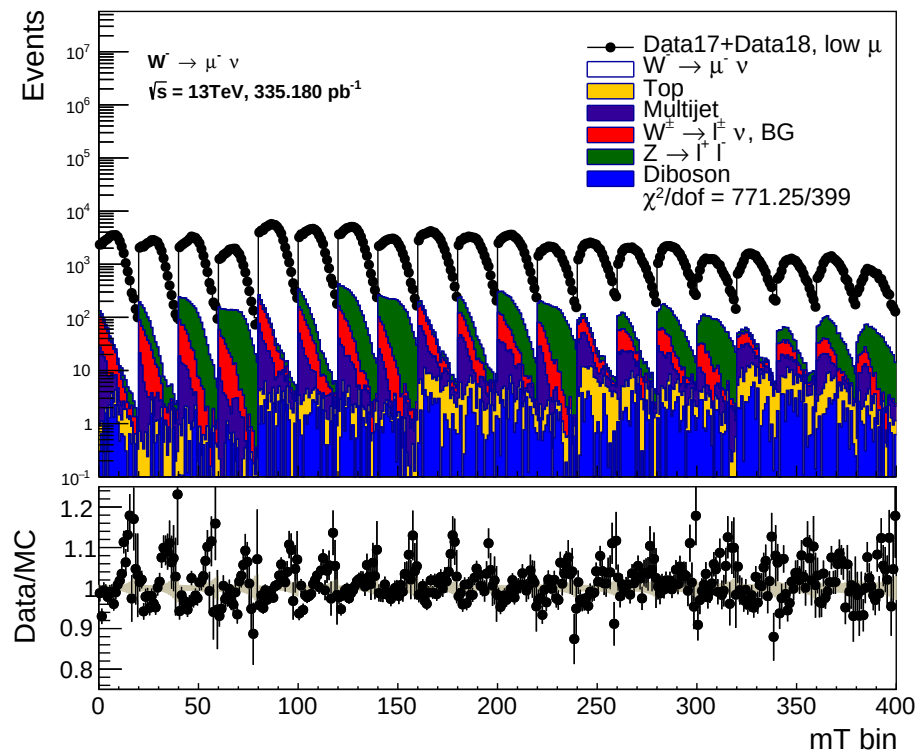


(a)

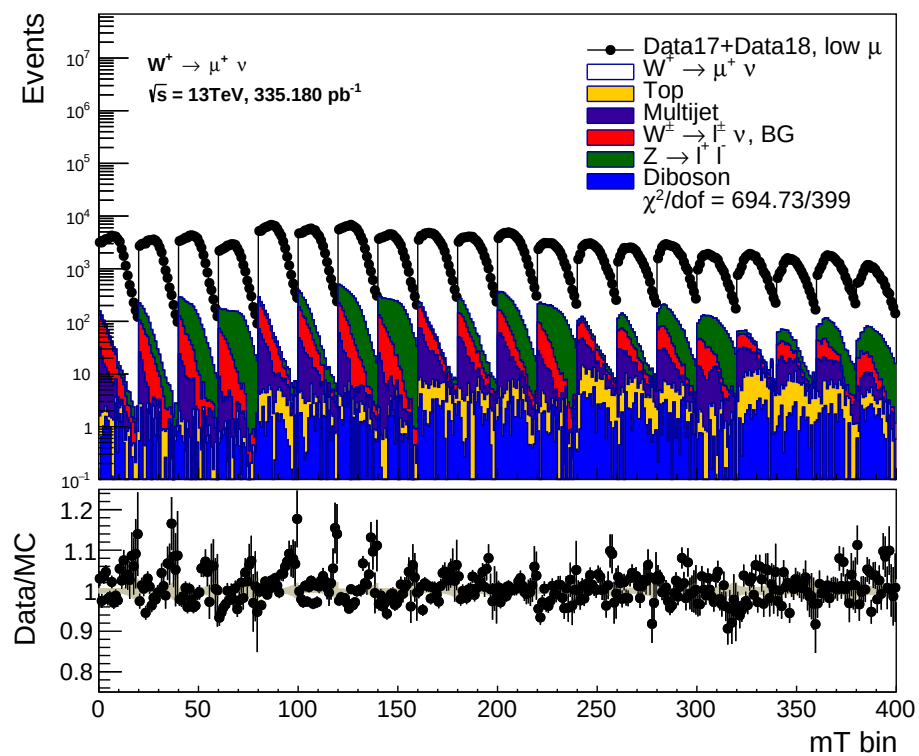


(b)

Figure 8.4 – The control plots of unrolled m_T spectra in 13 TeV electron channels.



(a)



(b)

Figure 8.5 – The control plots of unrolled m_T spectra in 13 TeV muon channels.

4289 set of NPs β_q according to:

$$\beta_q = \sum_p z_q^p \alpha_p \quad (8.17)$$

4290 It can be easily shown that neither the post-fit value of the POI nor the total uncer-
 4291 tainty of POI is affected by this transformation. In numerical computations, when
 4292 C^{syst} is not full rank, the eigenvalues beyond the rank are not exactly zero but tiny
 4293 values fluctuating around zero. For this reason, the absolute values are taken for the
 4294 eigenvalues, so as to avoid the square root of negative values. As a result, a set of
 4295 N_{bins} error vectors is always returned after the decomposition no matter C^{syst} is full
 4296 rank or not. The up and down for each of these error vectors can be symmetrized
 4297 with respect to 0 and then added to the nominal distribution to construct new paired
 4298 systematic variations.

4299 When the number of bins is smaller than the number of effective input systematic
 4300 sources, the above procedure is guaranteed to reduce the number of NPs. However, in
 4301 a real W mass measurement with combined channels and categorization, the number
 4302 of bins can easily exceed the number of NPs, especially when one would like to perform
 4303 the NP transformation within each type of systematics to preserve the basic physics
 4304 meanings of the NPs, i.e. in groups of electron SF, electron calibration, hadronic
 4305 recoil and so on. As introduced in [8], a further step of approximation can be brought
 4306 here to mitigate this issue by sorting the importance of the error vectors of C^{syst} from
 4307 the highest to the lowest and then imposing a truncation below a certain threshold.
 4308 The importance of the error vectors is assessed by the ranking of the corresponding
 4309 eigenvalues, following the steps below.

4310 First, all eigenvalues of the covariance matrix C^{syst} (taking the covariance matrix
 4311 in Figure 8.6(a) for instance) are sorted from the highest to the lowest. Then the
 4312 sum of absolute values of the eigenvalues are computed and normalized to unity. The
 4313 absolute values of the eigenvalues are scaled according to the normalization. After
 4314 that, a running sum is performed to sum up the absolute value of the eigenvalues from
 4315 the last one to the n -th one, as shown in Figure 8.6(b), meaning that the running
 4316 sum is equal to unity when $n = 1$ and decreases monotonically when n goes up.
 4317 When setting a truncation at 1% for the relative importance of the eigenvalues in this
 4318 example, the NPs for electron reconstruction SF systematics can be reduced from the
 4319 initial 272 to only 3. Consequently, the probability model is largely simplified without
 4320 losing much information of the systematics.

4321 Only the error vectors associated with the eigenvalues above the truncation will
 4322 be calculated and propagated to the probability model as new systematic variations.
 4323 That being said, those error vectors associated with the eigenvalues below the trun-
 4324 cation should in principle not be ignored either, otherwise one would expect an un-
 4325 derestimation of the total uncertainty from the fit due to less systematic effects being
 4326 accounted for in the probability model. One way to ensure a conservative total un-
 4327 certainty, or in other words, to approximately preserve the total covariance matrix,
 4328 is to neglect the bin-by-bin correlation of the error vectors of those eigenvalues below
 4329 the truncation, treat them as bin-by-bin uncorrelated uncertainties and add them in
 4330 quadrature. After that, these bin-by-bin uncorrelated systematic uncertainties are
 4331 denoted to be σ_i^{trunc} for the i -th bin, and will still be taken into account in the pro-
 4332 file likelihood fit in a way similar to the treatment of MC statistical uncertainty (see
 4333 Equation 8.24 for reference).

4334 The above procedure of NP transformation can be justified for most experimental

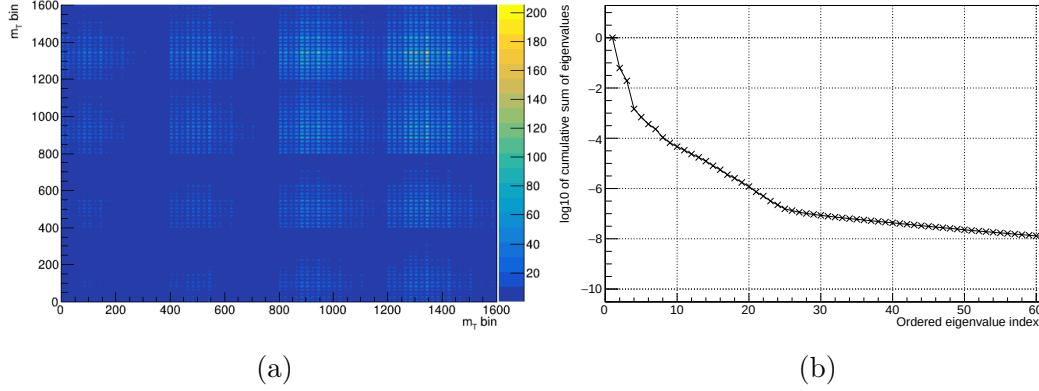


Figure 8.6 – NP reduction based on covariance matrix. (a) The covariance matrix of electron reconstruction SF uncertainties on the m_T distribution. (b) The reserved running sum of eigenvalues of the covariance matrix.

uncertainties because of their statistical nature. In contrast, all NPs accounting for theoretical uncertainties are kept unchanged, so as to avoid obscuring the physics meaning of the sources or mixing the possible non-Gaussian behaviour of some theory uncertainties. This strategy is expected to help the diagnostics of the fit and facilitate the interpretation of the result.

8.2.2 The handling of MC statistical uncertainty

Due to the finite number of recorded data events, fluctuations are expected in each bin of the observed kinematic distribution, bringing an uncertainty of statistical nature to the measurement when one extracts a physics quantity, such as m_W , from the kinematic distribution. Since the W mass is inferred by comparing the data with the MC templates, a similar uncertainty originates from the limited number of the MC samples as well. The fluctuations in the prediction of each bin due to the limited number of MC samples result in the MC statistical uncertainty. Unlike in the χ^2 offset method, where the MC statistical uncertainty enters the total uncertainty following Equation 8.2, the MC statistical uncertainty in the profile likelihood fit needs to be handled by the Barlow-Beeston prescription [9], according to which one additional NP is assigned per bin per sample to reflect the rate of statistical fluctuation. Once including these additional NPs into the probability model, the total uncertainty returned by the profile likelihood fit will correctly account for the effect of MC statistical uncertainty, as long as the difference of weights of MC events is relatively well under control. In practice, in order to reduce the total number of NPs in the probability model, a lighter version of this approach is implemented in HistFactory [5], which assigns one Barlow-Beeston NP to each bin for all the MC samples. Following the Barlow-Beeston-Lite prescription, the statistical model for each bin is modified by a factor due to the presence of MC statistics:

$$Pois(n_i | v_i(\boldsymbol{\alpha}) + \gamma_i v_i^{MC}(\boldsymbol{\alpha})) Poiss(m_i | \gamma_i \tau_i) \quad (8.18)$$

where $v_i(\boldsymbol{\alpha})$ stands for the MC predictions that do not need to include statistical uncertainty, and $v_i^{MC}(\boldsymbol{\alpha})$ are the MC predictions in which the statistical uncertainty

has to be accounted for. The NP per bin, denoted as γ_i , indicates the true rate that may differ from the original prediction $v_i^{MC}(\boldsymbol{\alpha})$ by a certain amount. If the total MC statistical uncertainty is δ_i for the i -th bin, then the relative statistical uncertainty becomes v_i^{MC}/δ_i . Therefore, assuming a reasonable spread of the MC weights, the effective size of the MC samples can be written as $m_i = (v_i^{MC}/\delta_i)^2$. The estimation of the MC bin content is deemed as an auxiliary measurement in this notation, so that the Poisson constraint term $Pois(m_i|\gamma_i\tau_i)$ is multiplied to the total likelihood, suggesting that m_i would fluctuate around $\gamma_i\tau_i$ when generating a new set of MC samples. While γ varies around unity in the fit, $\tau_i = (v_i^{MC}/\delta_i)^2$ is treated to be a fixed constant as if it does not fluctuate when new MC samples are generated by pseudo-experiments. One advantage of Barlow-Beeston-Lite prescription is that the conditional maximum likelihood estimate $\hat{\gamma}_i(\boldsymbol{\alpha})$ can be solved analytically as follows [5]:

$$\hat{\gamma}_i(\boldsymbol{\alpha}) = \frac{-B + \sqrt{B^2 - 4AC}}{2A} \quad (8.19)$$

4374

$$\begin{cases} A = v_i^{MC}(\boldsymbol{\alpha})^2 + \tau_i v_i^{MC}(\boldsymbol{\alpha}) \\ B = v_i(\boldsymbol{\alpha})\tau_i + v_i(\boldsymbol{\alpha})v_i^{MC}(\boldsymbol{\alpha}) - n_i v_i^{MC}(\boldsymbol{\alpha}) - m_i v_i^{MC}(\boldsymbol{\alpha}) \\ C = -m_i v_i(\boldsymbol{\alpha}) \end{cases} \quad (8.20)$$

Despite the fact that the Barlow-Beeston-Lite approach is widely used in the analyses involving MC statistics, it does not turn out to be suitable for the low pile-up W mass analysis where the joint fit of all the channels and all the categories involves a few thousand bins. However, neglecting completely the MC statistics in the fit will lead to an underestimated total uncertainty, which may affect the proper interpretation of the result of the measurement. To resolve this dilemma, an effective approximation is introduced to the low pile-up W mass analysis to completely remove the Barlow-Beeston NPs, as long as the size of the MC samples is at least a few times larger than that of the dataset. A similar method developed for the fit at Gaussian limit is described in [10]. The formulation of this approximation, referred to as 'data scaling' technique, will be illustrated in the paragraphs below.

Starting off without the presence of MC statistics, at the Gaussian limit of data statistics (more than a few tens of events in each bin), it is possible to describe the likelihood for each bin of the distribution by a Gaussian likelihood:

$$Pois(n_i|v_i) \approx Gaus(n_i|v_i, \sigma_i^{data}) \quad (8.21)$$

where σ_i^{data} is the data statistical uncertainty of the i -th bin. With the Gaussian likelihood, the MC statistics can be included in a straight-forward way:

$$Gaus(n_i|v_i, \sigma_i^{data}) \rightarrow Gaus(n_i|v_i, \sqrt{(\sigma_i^{data})^2 + (\sigma_i^{MC})^2}) \quad (8.22)$$

Due to the addition of MC statistics, the standard deviation enlarges, meaning the statistical power of the data diminishes accordingly, even though the observation stays at n_i . The Gaussian likelihood is capable of describing the mean and the standard deviation separately thanks to the two independent parameters in the function form. In a Poisson likelihood, considering that the mean and the variance are described by the same single parameter, the only way to reflect the reduced statistical power is hence to scale down the mean and the observation of the Poisson by the same factor

that catches the change in the variance:

$$Pois(n'_i|v'_i) = Pois\left(n_i \cdot \frac{(\sigma_i^{data})^2}{(\sigma_i^{data})^2 + (\sigma_i^{MC})^2} \middle| v_i \cdot \frac{(\sigma_i^{data})^2}{(\sigma_i^{data})^2 + (\sigma_i^{MC})^2}\right) \quad (8.23)$$

Denoting the data scaling factor to be $s_i = \frac{(\sigma_i^{data})^2}{(\sigma_i^{data})^2 + (\sigma_i^{MC})^2}$, the small bin-by-bin uncorrelated systematic uncertainties σ_i^{trunc} discussed in Section 8.2.1 can be treated in the same way as MC statistics and added back with an extended form of s_i :

$$s_i = \frac{(\sigma_i^{data})^2}{(\sigma_i^{data})^2 + (\sigma_i^{MC})^2 + (\sigma_i^{trunc})^2} \quad (8.24)$$

s_i is computed for each bin individually. The idea of the data scaling technique is consistent with the Barlow-Beeston-Lite prescription, which is to adjust the statistical power of the data by adding the effect of other sources of bin-by-bin uncorrelated uncertainties. It has to be pointed out that the MC statistical uncertainty is not parameterized in the formula of data scaling. As a result, the profile likelihood fit will gain significantly on the performance side. The approximation is anticipated to be acceptable in the low pile-up m_W analysis, since the result is hardly dominated by statistics and that the size of the signal MC samples is typically 5 to 10 times larger than the corresponding dataset. This approximate treatment of statistical uncertainty in the MC samples should only bring a marginal impact on the determination of the central value of m_W . A validation of this method, in terms of both the central value and the total uncertainty, has been performed with the following setup:

- Channel and kinematic distribution: 5 TeV $W^- \rightarrow e\nu$ channel. Unrolled m_T distribution in categories of u_T and η_ℓ .
- Statistical uncertainties: data, signal MC, background MC and multijet.
- Systematic uncertainties: PDF, electron calibration, luminosity, background cross-section uncertainty, multijet shape and yield uncertainties.

When applying the data scaling, the NP reduction is only applied to the electron calibration systematics in the signal part of the probability model. The total uncertainty is retrieved from the nominal fit, where the pseudo data is the sum of nominal MC, nominal background and nominal multijet. In addition, 400 bootstrap toys of the pseudo data are generated in order to quantify the non-closure between the Barlow-Beeston-Lite prescription (BBL) and the data scaling: the discrepancy between the central values of the POI between the two methods for the 400 toys is presented in 8.7. As it turns out, the central values of the fits using data scaling technique slightly deviate from those of the fit using the BBL prescription, but well within the level of MC statistical uncertainty (around 10 MeV in the channel under study). Also, as summarized in Table 8.1, the total uncertainty of the fit with the data scaling is compatible and more conservative with respect to the Barlow-Beeston-Lite method in this exercise.

The final data scaling for each bin takes into account the statistical uncertainty in signal MC, in the background simulation, in the multijet estimation, as well as the small systematic uncertainties that are considered to be bin-by-bin uncorrelated. The level of the scaling is shown in Figures 8.8 and 8.9 for 5 and 13 TeV, respectively.

The factor of data scaling is smooth in the muon channels. However, due to the larger fraction of multijet background and the associated larger statistical uncertainty,

NP for systematics	MC statistics	Total uncertainty [MeV]
Full NP	No MC stats.	40.76
Full NP	BBL	42.63
NP reduction	No MC stats.	40.76
NP reduction	Data scaling	42.64

Table 8.1 – The status of the systematic uncertainties in the W mass measurement. The MC statistics is taken into account by either the Barlow-Beeston-Lite prescription (BBL) or the data scaling.

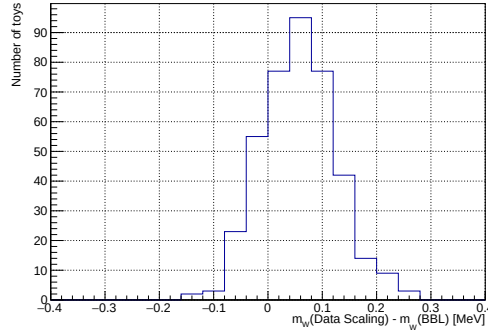


Figure 8.7 – The non-closure of the central value of the POI between data scaling and Barlow-Beeston-Lite prescription estimated from 400 bootstrap toys of the pseudo data. Even though there is a noticeable bias caused by the data scaling with respect to the Barlow-Beeston-Lite prescription, the bias itself is within the spread of the non-closure. On top of that, both the bias and the spread of non-closure are extremely small compared to the total uncertainty of 43 MeV.

the effect of data scaling is in general stronger in the electron channels. The noticeable spikes of the data scaling in the electron channels are also brought by the larger statistical uncertainty in the multijet estimation.

8.2.3 Combination of statistically correlated observables

Another limitation of the profile likelihood fit, compared to the χ^2 offset method, is the difficulty in combining the measurements from statistically correlated observables. While the results of the χ^2 offset method can be combined using the standard BLUE procedure [11][12], provided that the correlations between different observables are known, the Poisson terms in the profile likelihood fit make it impossible to introduce a simple correlation matrix to describe the statistical correlation between the bins of two kinematic distributions. Although one can always add one more dimension of the distribution to measure a new observable simultaneously, the events of the data in each measurement bin is drastically diluted when one more dimension is added to the probability model. Not only will there be a higher chance of finding an empty bin, but also the computation of the fit will become more time consuming with more bins. In order to mitigate or ultimately solve this drawback of the profile likelihood fit, several proposals have been worked out so far.

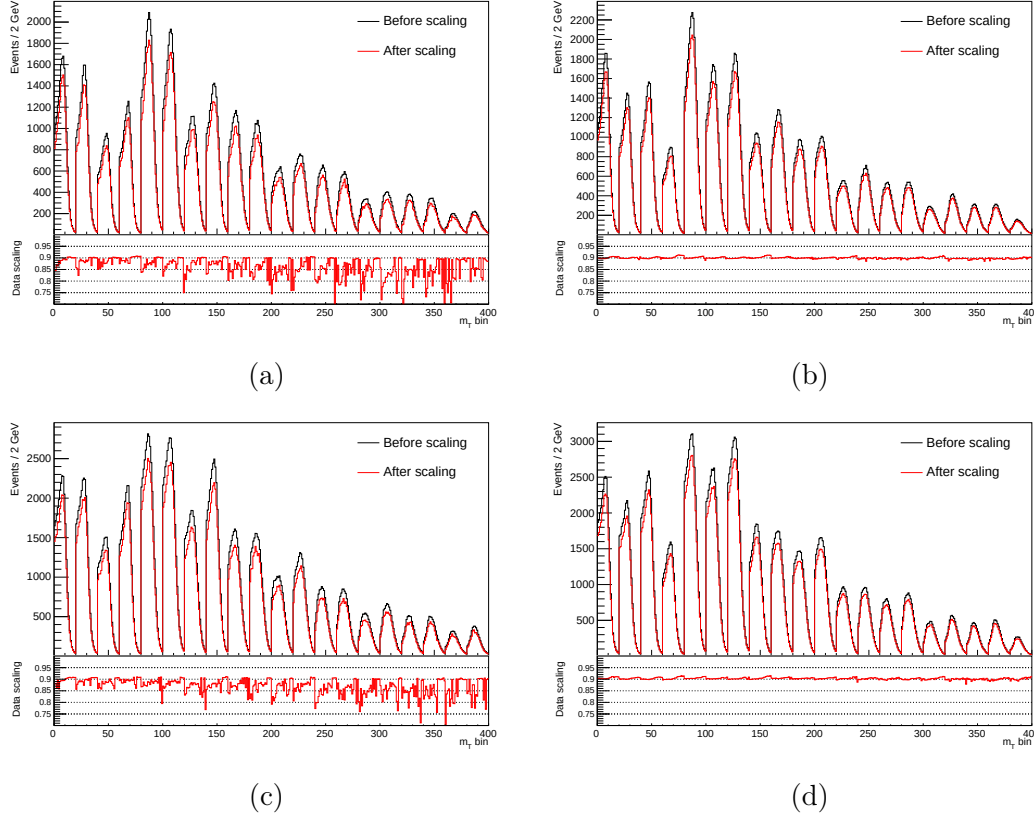


Figure 8.8 – The data scaling applied to the unrolled m_T spectra at 5 TeV. The distribution before the data scaling is pseudo data made of a sum of nominal signal MC, simulated background and multijet. (a) $W^- \rightarrow e\nu$ channel. (b) $W^- \rightarrow \mu\nu$ channel. (c) $W^+ \rightarrow e\nu$ channel. (d) $W^+ \rightarrow \mu\nu$ channel.

The first plan would be to fill a linear combination of the observables of interest and fit to the distribution of the new observable. For instance, the lepton p_T and m_T are statistically correlated in the W mass analysis. Both of these two quantities have good sensitivity to W mass. Instead of doing the measurement using each observable separately and then combining the two results, a linear combination, for example

$$Sm_T = 2a \cdot p_T^\ell + (1 - a) \cdot m_T \quad (8.25)$$

can be devised for each event in data and MC. The distribution of Sm_T is a valid input to the profile likelihood fit, since its bins are statistically uncorrelated, such that the information of both lepton p_T and m_T is used simultaneously with double counting. A factor of 2 is multiplied to p_T to make the scales of the two observables more compatible in the sum, as $m_T \approx 2p_T^\ell$ in the $u_T < \approx 25$ GeV region that is interesting for the W mass analysis. The combination coefficient a , which resembles the combination weight in the BLUE method, cannot be known from any prior knowledge and has to be optimized using (pseudo) data to achieve the smallest total uncertainty. The flexibility of this proposal is that the Equation 8.25 can be easily extended to include other observables such as the neutrino E_T^{miss} to better constrain recoil calibration systematics. In case there is categorization in the fit, the combination coefficient a needs to be determined separately in each category. The determination of the combination coefficients is shown in Figure 8.10 for an illustration, where only the signal

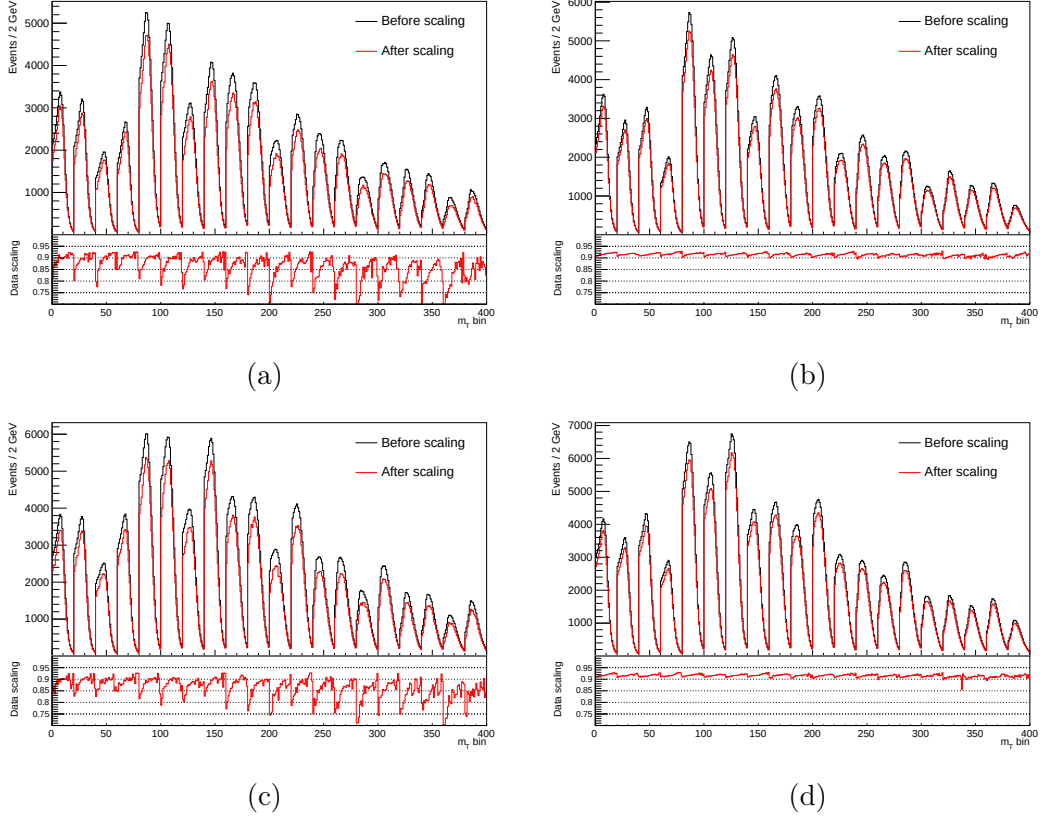


Figure 8.9 – The data scaling applied to the unrolled m_T spectra at 13 TeV. The distribution before the data scaling is pseudo data made of a sum of nominal signal MC, simulated background and multijet. (a) $W^- \rightarrow e\nu$ channel. (b) $W^- \rightarrow \mu\nu$ channel. (c) $W^+ \rightarrow e\nu$ channel. (d) $W^+ \rightarrow \mu\nu$ channel.

part of the probability model is fitted and only the data and MC statistical uncertainties are taken into account. The Sm_T is ultimately optimized against statistical uncertainty, recoil calibration systematics, lepton calibration and efficiency scale factor systematics, PDF uncertainty and p_T^W modelling uncertainty. The optimization of the coefficients is studied with an interval of 0.1. The corresponding result is summarized in Table 8.2. A joint Sm_T fit of all the 20 categories using the optimized combination coefficients yields an total uncertainty of 37.22 MeV. In comparison, the total uncertainty of the m_T fit is slightly worse: 37.84 MeV. A better performance of the Sm_T is anticipated, provided that the two observables in the combination share a moderate correlation but carry comparable BLUE weights.

$a(u_T, \eta_e)$	$ \eta_e < 0.6$	$0.6 < \eta_e < 1.2$	$1.2 < \eta_e < 1.8$	$1.8 < \eta_e < 2.47$
$u_T < 5$ GeV	0.4	0.3	0.2	0
$5 < u_T < 10$ GeV	0.2	0.3	0.2	0.3
$10 < u_T < 15$ GeV	0	0.1	0	0
$15 < u_T < 20$ GeV	0.1	0	0.1	0
$20 < u_T < 25$ GeV	0	0	0	0

Table 8.2 – The optimal combination coefficients for Sm_T in 5 TeV $W^- \rightarrow e\nu$ channel.

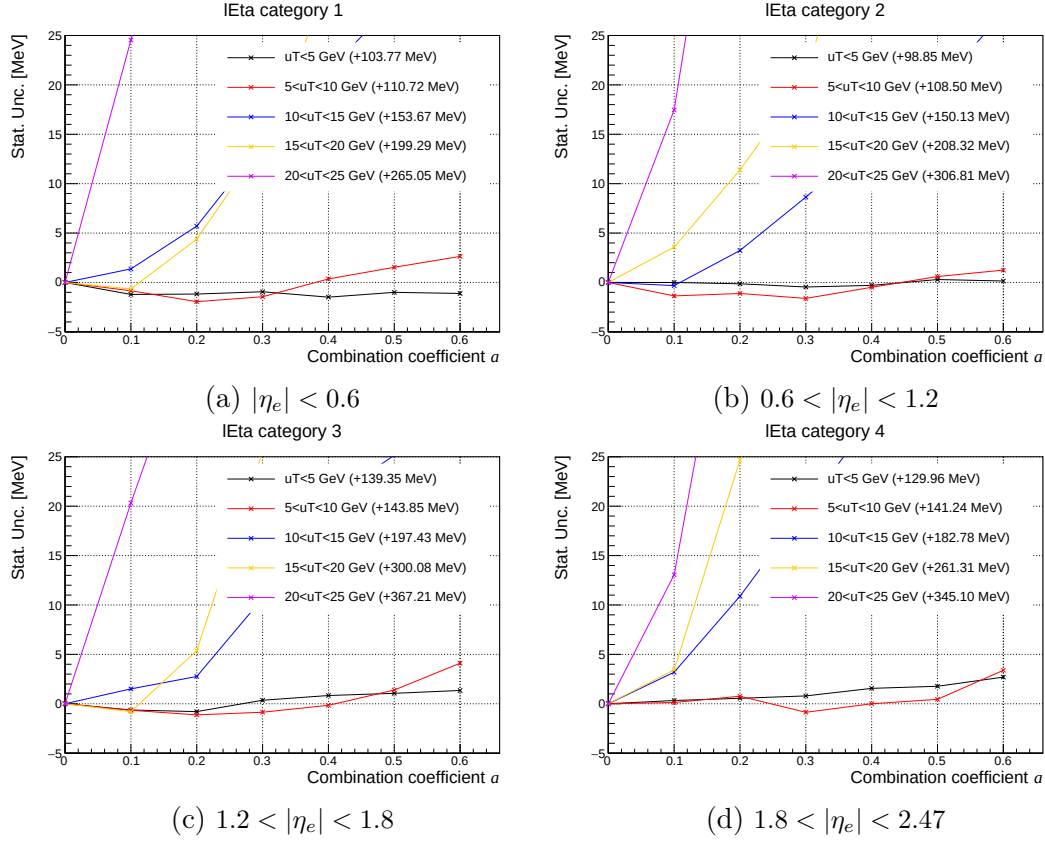


Figure 8.10 – The optimization of combination coefficients for Sm_T in 5 TeV $W^- \rightarrow e\nu$ channel taking into account the data and MC statistical uncertainties. The minima of the uncertainty curves for each channel correspond to the preferred value of the combination coefficients.

Despite the elegance of the formula, the optimization of the combination coefficients against the total uncertainty is time consuming. Besides, due to the bin-to-bin correlation in the systematic variations, the optimization in separate categories including systematics may not lead to the optimal total uncertainty when all categories are combined in a joint fit. Finally, with the Run 2 low pile-up data, the improved recoil resolution greatly helps m_T gain statistical sensitivity to W mass. Particularly, in the region $15 < u_T < 25$ GeV, m_T becomes way more sensitive to W mass than p_T^ℓ in comparison. The addition of p_T^ℓ is not expected to improve the result of m_T by a large degree, which justifies the setup of the profile likelihood fit focusing on only the m_T distribution at this stage of the analysis.

A possible alternative solution, as already implemented in the reanalysis of ATLAS measurement of W mass at 7 TeV [13], would be to simultaneously run bootstrap toys in m_T fit and p_T^ℓ fit, for both the data n_i and the global observables a_p in Equation 8.8. The correlation revealed by the bootstrap toys can be used to combine the uncertainties of the measurement from multiple observables.

8.3 Review of uncertainties in the W mass fit

The types of uncertainty sources entering the determination of W mass should include:

- Statistical uncertainty: bin-by-bin uncorrelated uncertainty stemming from data, signal MC, background MC and multijet estimation.
- Experimental uncertainty: electron SF, electron calibration, muon SF, muon calibration and recoil calibration.
- Background uncertainty: luminosity uncertainty, cross-section uncertainty of simulated background process, multijet yield and shape uncertainties.
- Uncertainty in electroweak corrections.
- QCD modelling uncertainty: fixed-order PDF uncertainty, spin-correlation and p_T^W modelling uncertainty.

All the systematic variations are symmetrized with respect to the nominal in the probability model for the W mass fit.

8.3.1 Experimental, background and luminosity uncertainties

All experimental uncertainties, as discussed in Sections 6.1 6.2 and 6.3, along with the background cross-section uncertainties and luminosity uncertainties reuse what has been applied to the p_T^W analysis.

The electron reconstruction SFs are extrapolated from the standard high pile-up SFs to the low pile-up regime and is applied to both the 5.02 and 13 TeV datasets. The electron identification SFs are measurement in-situ using Z samples, separately for 5.02 and 13 TeV data. The electron isolation and trigger efficiencies SF are measured in-situ using combined 5.02 and 13 TeV low pile-up datasets. The uncertainties of the electron SFs measurements include the statistical uncertainty due to the limited number of tag-and-probe pairs and the systematic uncertainties dominated by the photons and background control performance. The statistical uncertainties are uncorrelated in the SF bins, while the systematics are considered to be correlated. The electron energy scale and resolution corrections are determined from the Z events in the low pile-up datasets using the standard ATLAS procedures [14]. In addition to the statistical uncertainties in the electron energy calibration, the systematic uncertainties at high pile-up that are related to the calorimeter modelling are also included for the low pile-up data analysis, which contribute to large electron energy scale uncertainties.

The muon reconstruction SFs are extrapolated from the high pile-up results [15], while the trigger SFs, the isolation SFs and the TTVA SFs are measured in-situ using the low pile-up data. The muon SF uncertainties consist of both the statistical uncertainty of the tag-and-probe method and the systematic uncertainties mainly attributed to the background subtraction and the variation of the selection criteria used in the tag-and-probe method [16]. The muon momentum scale and resolution are derived from the high pile-up data. The muon sagitta bias corrections are derived from the 2017 low pile-up datasets. The muon calibration will have to be further improved for the m_W measurement, since large uncertainties have been found in muon energy scale and muon sagitta bias.

The hadronic recoil calibration corrects for the mis-modelling of the underlying events, the direction of the recoil, as well as the response and the resolution the recoil.

The recoil calibration uncertainties originate from the non-closure in the extrapolation of the calibration from Z events to W events, and from the statistical uncertainty in the response and resolution corrections.

A cross-section uncertainty of 10 % is assigned to the simulated top-quark related background processes and di-boson background, except for an uncertainty of 7 % on the $t\bar{t}$ process. The cross-section uncertainty of background contributed by Z-boson, charge flip of the lepton detection and W-boson decaying into a pair of tau and neutrino are conservatively estimated to be 5 %. The luminosity uncertainty, 1.6% for 5 and 1.5% for 13 TeV are determined from the luminosity calibration using a combination of the primary measurement with the LUCID-2 detector [17] and the complementary measurements with inner detector and calorimeters [1]. The result of updated luminosity calibration will be used for the measurement in the future. The multijet background is estimated using the data-driven method for the p_T^W analysis. However, small modifications are introduced to the determination of the yield and shape for both the central value and the associated uncertainties, in order to adapt to the different event selection as well as the categorization in the m_W analysis.

An additional selection $u_T < 25$ GeV is applied to all four regions for the data-driven multijet estimation defined in Section 5.4. The extrapolation of the yield keeps the same methodology, based on p_T^ℓ , m_T and E_T^{miss} spectra. For the moment, the multijet yield uncertainty is only given by the extrapolation uncertainty.

The linear extrapolation formula Equation 5.23 requires ample statistics in all anti-isolated slices. But in the W mass measurement, the estimation of the multijet shape relies on the unrolled m_T spectrum in categories of u_T and η_ℓ . The 2D categorization causes low statistics in each anti-isolated slice of CR2 with an interval of 0.1, making the linear extrapolation formula no longer reliable. Therefore, the central value of the multijet shape in SR ($0 < \text{isolation} < 0.1$) is simply taken from the first anti-isolated slice in CR2 ($0.1 < \text{isolation} < 0.2$), while the shape variation is defined to be the overall shape in CR2 for $0.1 < \text{isolation} < 0.5$. The histogram smoothing in ROOT based on the algorithm 353QH [18] is applied to both the central value and the shape variation. Any bins of the distribution below zero are set to zero manually afterwards. The statistical uncertainty is directly taken from the $0.1 < \text{isolation} < 0.2$ slice in CR2 regardless of the smoothing or the manual modification of negative bin content. As an example, the multijet background on m_T for the W mass fit in is shown in Figure 8.11 for the 13 TeV $W^+ \rightarrow \mu\nu$ channel. The central value with the statistical uncertainty, the yield variations and the shape variations are all displayed in the plot.

8.3.2 Propagation of the p_T^W uncertainty from the measurement

In the 7 TeV m_W measurement by ATLAS [7], the systematic uncertainty that stems from the p_T^W modelling is at the level of 6 MeV, including the uncertainty in parton shower tuning, the parton shower PDF uncertainty, the variations of the masses of the charm and bottom quarks, as well as the variations of the factorization scale in the calculation of QCD ISR. In order to constrain the p_T^W uncertainty, the m_W measurement using low pile-up data will make use of the result of the direct p_T^W measurement. As a preliminary study of the m_W fit, one important goal is to estimate the level of uncertainty originating from p_T^W in low pile-up data. This estimation can be achieved by propagating the unfolded level p_T^W uncertainty from the direct

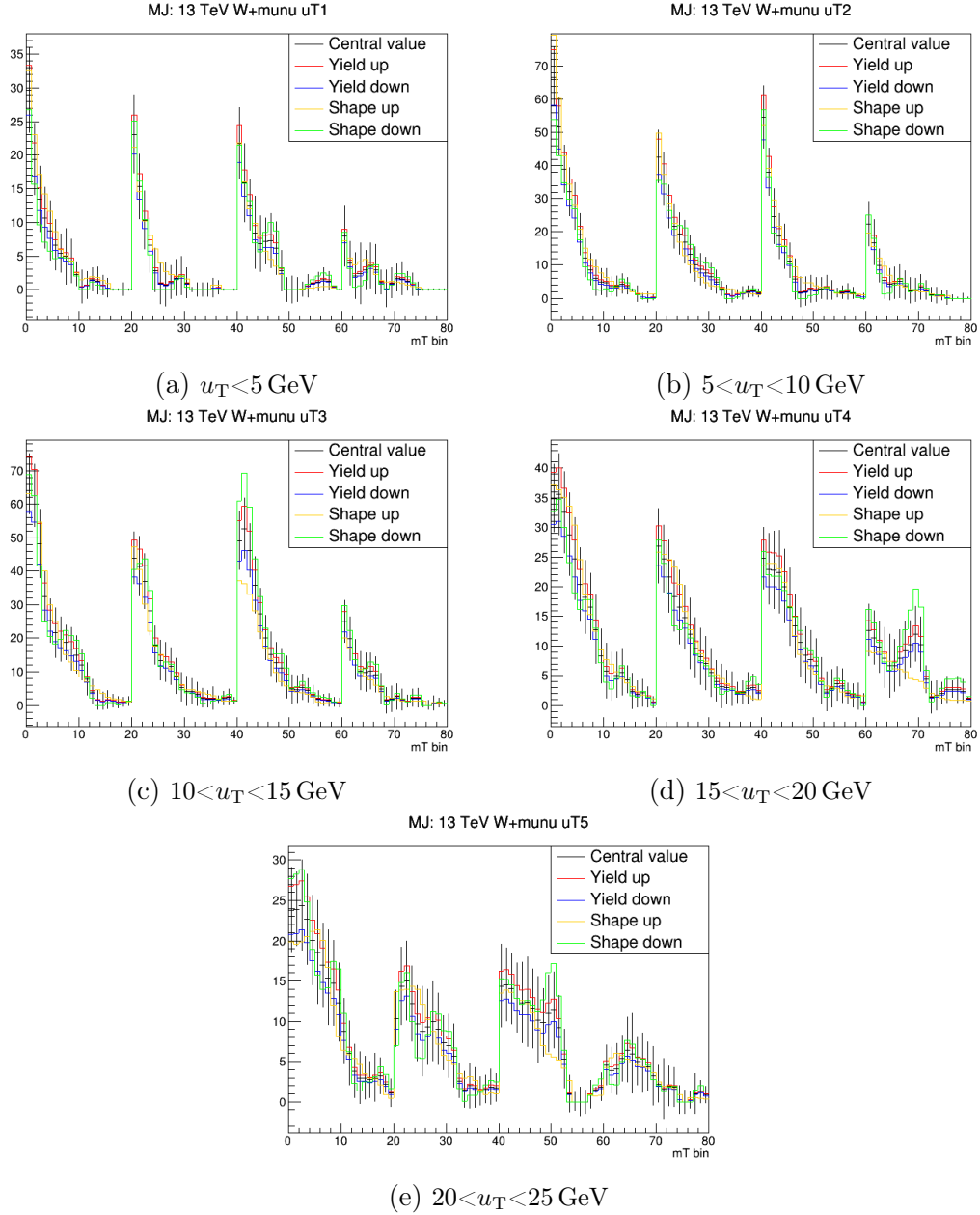


Figure 8.11 – The data-driven multijet background of m_T spectra in 13 TeV $W^+\mu\nu$ channel. The data scaling described in Section 8.2.2 is not applied to the distributions. Every 20 bins in the unrolled m_T distribution correspond to a category of η_ℓ .

4589 measurement reported in Section 7.5.1 to the probability model of W mass fit.

4590 The central value of the underlying p_T^W distribution in the simulation is reweighted

4591 in p_T^W as described in Section 7.3. The uncertainty variations of p_T^W are derived from

4592 the total covariance matrix at unfolded level, also provided by the p_T^W measurement.

4593 The total covariance matrix at unfolded level for $p_T^W < 63$ GeV (9 bins in each dimension)

4594 is shown in Figure 8.12 for the electron channels. An eigenvalue decomposition

4595 is performed on the covariance matrix to produce 9 pairs of error vectors of p_T^W , with

4596 each pair symmetrized with respect to the central value of the measurement. Similar

4597 to the reweighting of the central value, the underlying p_T^W distribution in the simulation

4598 can also be reweighted to the up and down variations of the 9 pairs of error

vectors, shown in the top rows of Figures 8.13 and 8.14. For the m_W measurement, all kinematic distributions of the signal process at the reconstructed level are produced with the underlying p_T^W reweighted to the central value of the measurement and to the 9 pairs of up and down variations. For example, the bottom rows of Figures 8.13 and 8.14 present the p_T^W variations on the m_T spectra. After that, the p_T^W variations of the reconstructed level distributions are injected to the probability model for template morphing, and 9 NPs are assigned to the p_T^W systematics per boson charge and per center-of-mass energy.

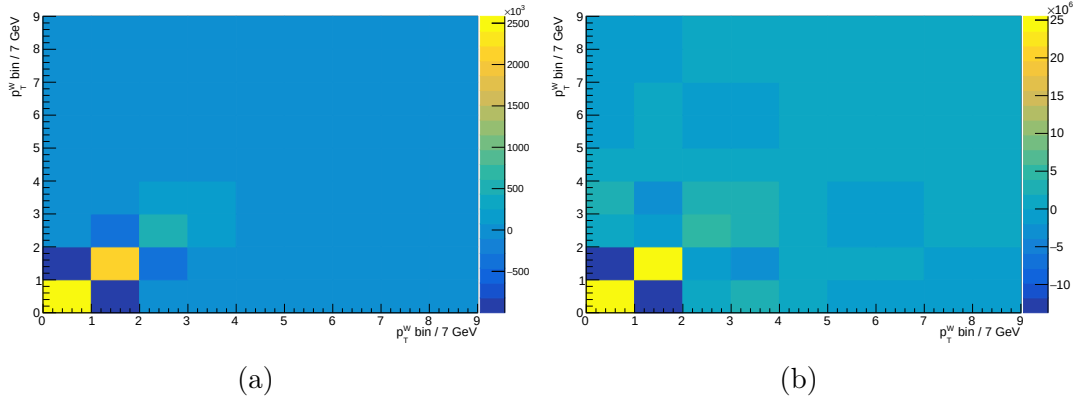


Figure 8.12 – Examples of the total covariance matrix of p_T^W at unfolded level. (a) 5 TeV $W^- \rightarrow e\nu$ channel. (b) 13 TeV $W^- \rightarrow e\nu$ channel.

Although the propagation of the p_T^W uncertainty described above provides a simplified treatment of p_T^W modelling uncertainty in the low pile-up W mass measurement, by no means it should become the final strategy for p_T^W uncertainty in the analysis. First, the correlation between boson p_T and rapidity are ignored in the 1D p_T^W reweighting. Second, the recoil calibration uncertainties are major sources of systematic uncertainty shared by p_T^W analysis and W mass analysis. The correlation of recoil calibration uncertainties at unfolded level for p_T^W analysis and at reconstructed level for W mass analysis is not properly accounted for in the current method. A possible solution to making use of the results of the p_T^W measurement in a more appropriate way will be discussed in Section 8.5.

8.3.3 PDF and other modelling uncertainties

A major QCD modelling uncertainty for W mass measurements in hadron colliders is the uncertainty in the fixed-order QCD calculation. The fixed-order QCD uncertainty is assessed with the help of global PDF sets. The central value of the default PDF set used in the generation of signal MC samples is CT10 [19], a NLO accuracy PDF set, even though the most up-to-date PDF uncertainties are available at NNLO accuracy and are considered for the W mass measurement. To keep the PDF baseline matching the variations, a reconstructed level PDF reweighting is applied to the unrolled m_T spectrum to replace the baseline PDF by an NNLO prediction: CT18ANNLO [20], a global PDF set including the ATLAS 7 TeV data of W/Z rapidity distributions [21] in the fit. The PDF baseline reweighting procedure goes as the following:

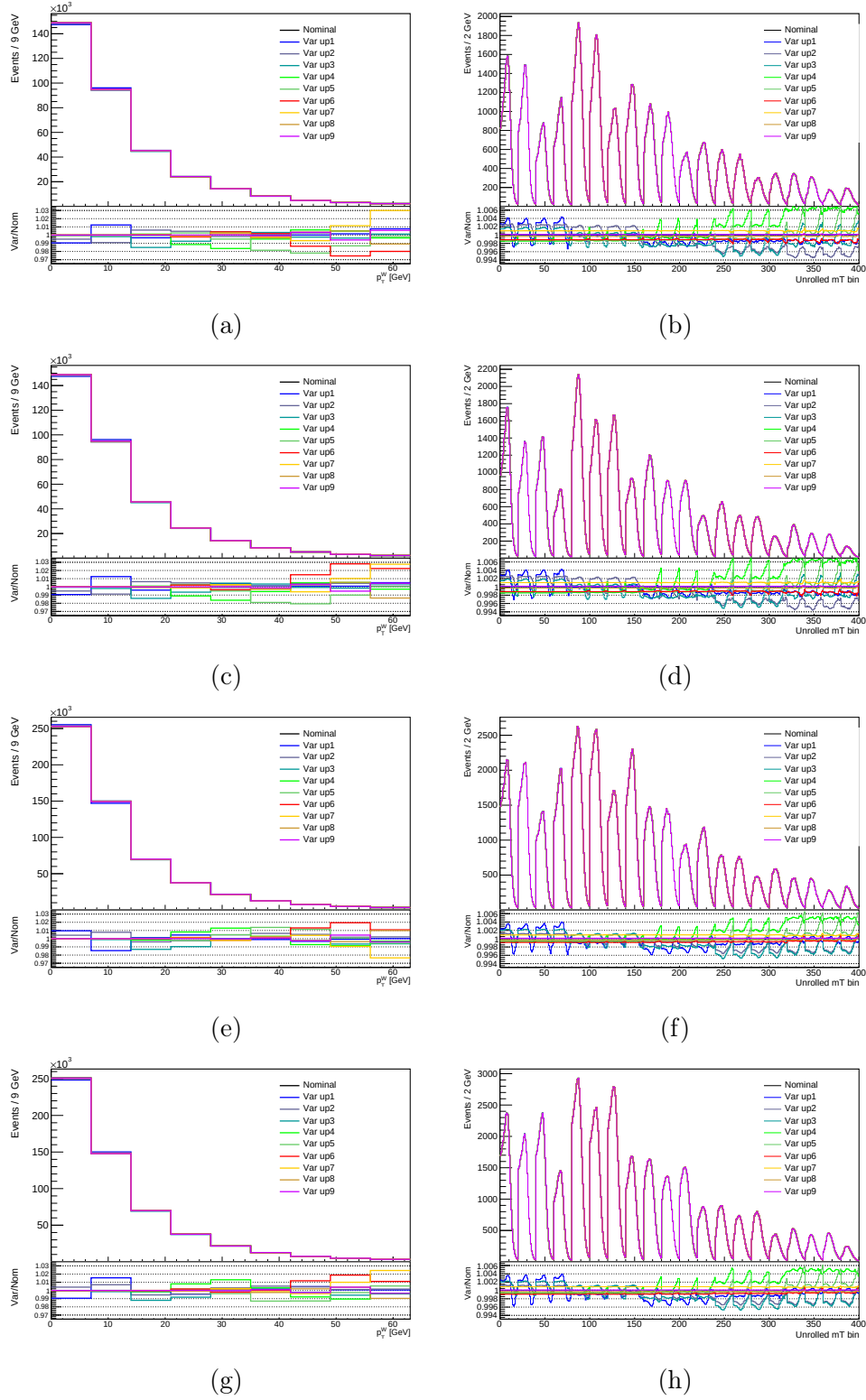


Figure 8.13 – Propagation of p_T^W uncertainty from the direct measurement at 5 TeV to the W mass fit. Orthogonal error vectors (left): at unfolded level. (right): on the reconstructed level m_T distributions. (a)(b) $W^- \rightarrow e\nu$. (c)(d) $W^- \rightarrow \mu\nu$. (e)(f) $W^+ \rightarrow e\nu$. (g)(h) $W^+ \rightarrow \mu\nu$.

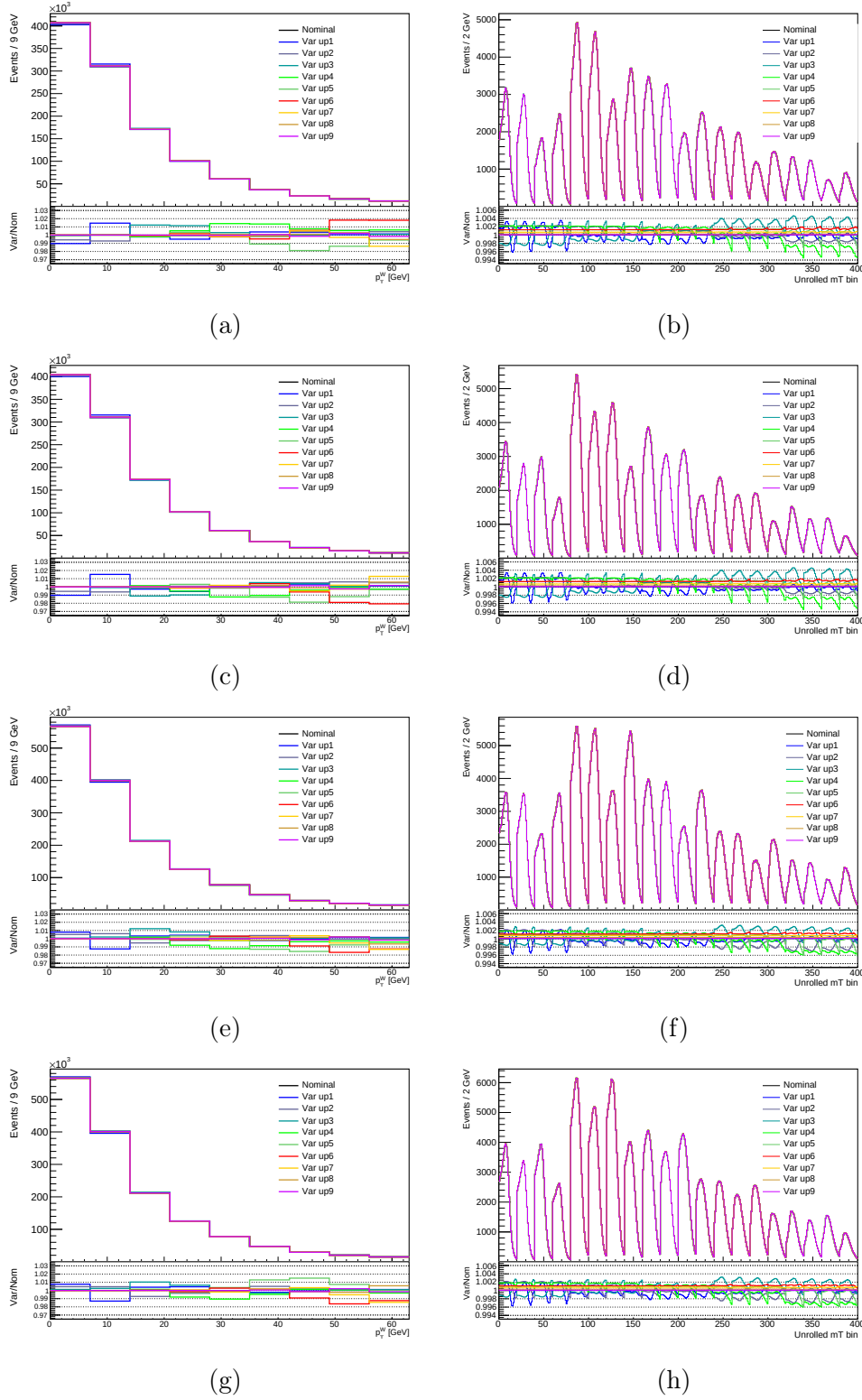


Figure 8.14 – Propagation of p_T^W uncertainty from the direct measurement at 13 TeV to the W mass fit. Orthogonal error vectors (left): at unfolded level. (right): on the reconstructed level m_T distributions. (a)(b) $W^- \rightarrow e\nu$. (c)(d) $W^- \rightarrow \mu\nu$. (e)(f) $W^+ \rightarrow e\nu$. (g)(h) $W^+ \rightarrow \mu\nu$.

- 4629 — Generate the nominal signal MC distribution using CT10 and using CT18ANNLO
- 4630 with the help of an internal Powheg+Herwig7 [22] PDF reweighting.
- 4631 — Calculate the reconstructed level ratio of CT18ANNLO/CT10 (p_T^W, y) 2D reweight-
- 4632 ing for the nominal signal MC distribution.
- 4633 — Apply the CT18ANNLO/CT10 derived in the step above to all the entire signal
- 4634 part of the probability model, including: nominal MC, the W mass templates,
- 4635 the systematic variations related to the signal apart from the PDF variations
- 4636 of CT18ANNLO.
- 4637 The effect of this PDF baseline reweighting is shown in Figures 8.15 and 8.16.

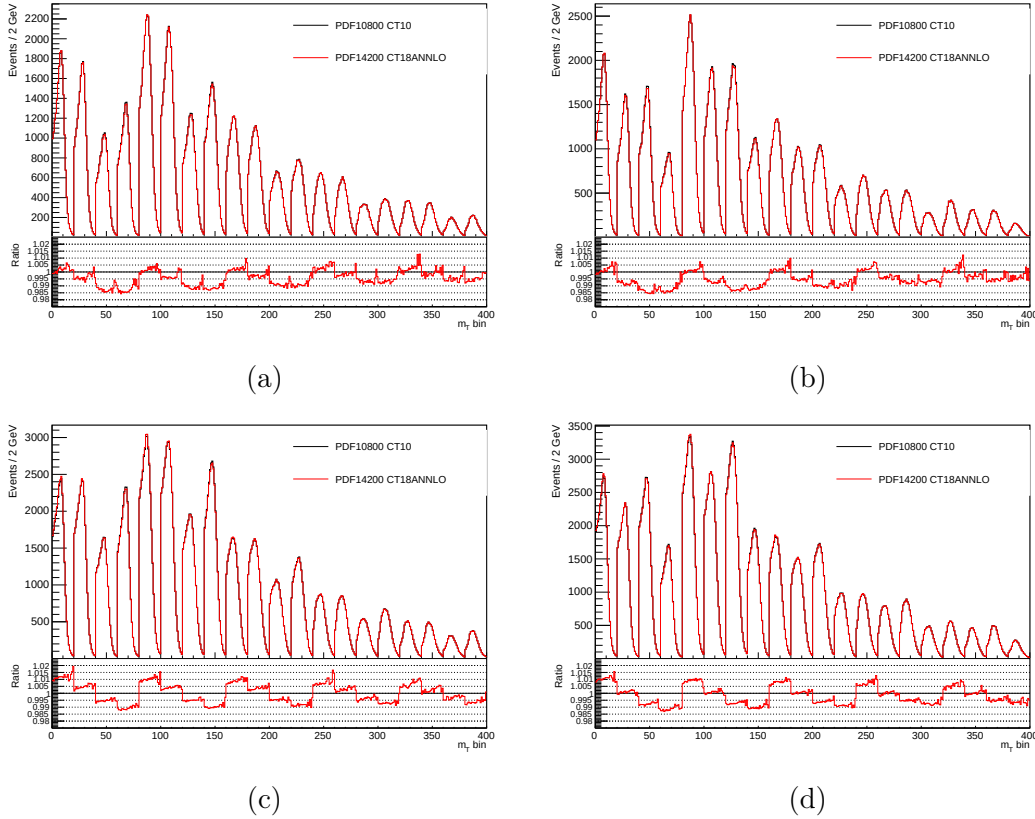


Figure 8.15 – The ratio of reconstructed level CT10ANNLO/CT10 reweighting on the unrolled m_T spectra at 5 TeV. (a) $W^- \rightarrow e\nu$ channel. (b) $W^- \rightarrow \mu\nu$ channel. (c) $W^+ \rightarrow e\nu$ channel. (d) $W^+ \rightarrow \mu\nu$ channel.

4638 For the moment, the uncertainty in the fixed-order QCD calculation is assessed
 4639 using the 29 pairs of eigenvector sets of CT18ANNLO global PDF. As a more recent
 4640 global PDF set, CT18ANNLO provides better precision than the CT10NNLO [23],
 4641 CT14NNLO [24] and MMHT2014 [25] considered in the previous ATLAS m_W mea-
 4642 surement at 7 TeV center-of-mass energy. The PDF uncertainty in the low pile-up
 4643 m_W measurement is expected to be lower than the 9 MeV reported in the 7 TeV
 4644 publication, thanks to the updated result of the global PDF analysis and the profile
 4645 likelihood technique used in the statistical analysis.

4646 Another source of QCD modelling uncertainty is the spin-correlation uncertainty.
 4647 For the nominal signal process, the spin-correlation coefficients in the signal MC is
 4648 reweighted from the Powheg prediction to the DYTURBO prediction using CT10 PDF

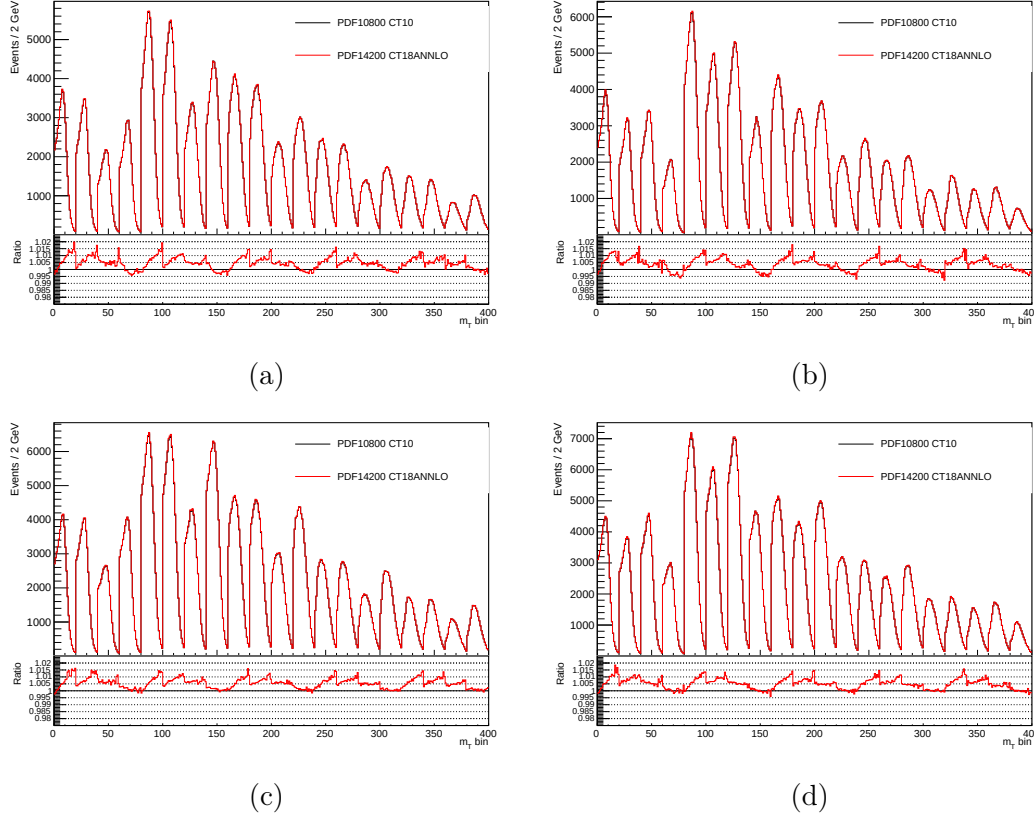


Figure 8.16 – The ratio of reconstructed level CT10ANNLO/CT10 reweighting on the unrolled m_T spectra at 13 TeV. (a) $W^- \rightarrow e\nu$ channel. (b) $W^- \rightarrow \mu\nu$ channel. (c) $W^+ \rightarrow e\nu$ channel. (d) $W^+ \rightarrow \mu\nu$ channel.

at NNLO accuracy in a 2D mapping of spin-correlation coefficients in bins of W-boson p_T^W and rapidity. Different from the estimation of spin-correlation uncertainty in the 7 TeV paper, where the uncertainties in the prediction of W-boson angular coefficients are propagated from the experimental uncertainty of the Z-boson production measurement [26] while taking into account the non-closure between the measurement and the prediction of A_2 as a function of p_T^Z , for the low pile-up analysis a variation of spin-correlation coefficients will be produced by reweighting the Powheg prediction to DYTURBO prediction at NLO accuracy. However, whether or not the NLO vs NNLO difference of spin-correlation coefficients can be used as the systematic variation for the analysis needs validation in the future.

As for baseline electroweak corrections, the QED ISR is included in Pythia8 parton [27] shower, while the QED FSR is taken into account by interfacing the Powheg+Pythia8 [28]–[31] prediction to PHOTOS++ [32]. The electroweak uncertainty will be evaluated by Winhac [33] at reconstructed level. The uncertainty in QED FSR correction is obtained by comparing the default photon emission matrix elements of PHOTOS++ with the calculation of Yennie–Frautschi–Suura formalism available in Winhac. The other sources of electroweak uncertainty include the missing pure weak correction, the interference between QED ISR and QED FSR, as well as the effect of lepton pair production in the final state. As for an improvement in the low pile-up m_W analysis, the electroweak correction uncertainties will be folded to

the reconstructed level to give a more conservative and realistic representation of the uncertainty with respect to the previous 7 TeV measurement.

8.4 Summary of preliminary results

The report of preliminary results of the low pile-up W mass measurement focuses on the uncertainty of the measurement. The uncertainty decomposition for the profile likelihood fit is presented following the prescription of citation (the uncertainty decomposition paper), and will be explained in details in the next chapter. In brief, the bin-by-bin uncorrelated statistical uncertainty can be estimated by fluctuating the (pseudo) data to make bootstrap toys and by repeating the fit to these toys using the full probability model. The spread of the toys gives the contribution from the statistical uncertainty. A source of systematic uncertainty that affects the kinematic distribution in a bin-by-bin correlated manner can be estimated by pulling the corresponding global observable from the central value to ± 1 standard deviation, then repeating the fit using the full probability model. The offset in POI with respect to the nominal fit to the Azimov data corresponds to the contribution from this systematic source.

The current status of the systematic uncertainties considered in the measurement is summarized in Table 8.3.

Experimental systematics		Theoretical systematics	
Source	Status	Source	Status
Recoil calibration	Done	PDF uncertainties	CT18NNLO
Lepton SF	Done	Spin-correlation coefficients	On-going
Lepton calibration	Done	EW corrections	On-going
Luminosity	Done	Background cross-sections	Done
Multijet systematics	Done	p_T^W modelling	Done

Table 8.3 – The status of the systematic uncertainties in the W mass measurement.

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8.4.1 Separate fit in individual channels

With the current scheme of categorization, each W-boson decay channel comprises 400 bins of (m_T, u_T, η_ℓ) distribution. The NP reduction is carried out at the level of all channels and categories. Data scaling has been implemented. 400 bootstrap toys of pseudo data are produced to estimate the contribution of bin-by-bin uncorrelated uncertainty (mainly statistical uncertainty). The details of the uncertainty breakdown in each channel are presented in Tables 8.4, 8.5 and 8.6.

Even though the numbers of events are comparable between the electron and the muon channels in the data, the statistical uncertainties in the muon channels are much higher due to the large muon calibration systematic variations and the strong correlations between the POI and the NPs for muon calibration systematics.

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Channel	Total	Stat.	Syst.
5 TeV $W^- \rightarrow e\nu$	44.52	37.53 ± 1.33	25.93
5 TeV $W^+ \rightarrow e\nu$	43.21	31.53 ± 1.11	28.23
5 TeV $W^- \rightarrow \mu\nu$	59.57	51.14 ± 1.81	32.02
5 TeV $W^+ \rightarrow \mu\nu$	61.85	49.05 ± 1.73	36.76
13 TeV $W^- \rightarrow e\nu$	39.47	28.10 ± 0.99	25.67
13 TeV $W^+ \rightarrow e\nu$	38.08	28.66 ± 1.01	26.25
13 TeV $W^- \rightarrow \mu\nu$	59.55	45.20 ± 1.60	35.95
13 TeV $W^+ \rightarrow \mu\nu$	61.26	48.56 ± 1.72	38.33

Table 8.4 – The total uncertainty, statistical uncertainty and total systematic uncertainty of the fit in separate channels. Unit in [MeV]. The error on the statistical uncertainty is caused by the limited number of bootstrap toys.

Channel	Lepton SF	Lepton Calib.	Recoil	Lumi.	MJ
5 TeV $W^- \rightarrow e\nu$	6.50	20.42	6.59	1.51	1.93
5 TeV $W^+ \rightarrow e\nu$	5.11	21.76	7.04	0.86	0.82
5 TeV $W^- \rightarrow \mu\nu$	10.78	26.45	5.28	1.71	1.72
5 TeV $W^+ \rightarrow \mu\nu$	11.17	29.57	6.87	0.38	0.68
13 TeV $W^- \rightarrow e\nu$	6.59	17.47	6.57	1.25	4.22
13 TeV $W^+ \rightarrow e\nu$	5.39	18.18	6.41	0.77	2.32
13 TeV $W^- \rightarrow \mu\nu$	8.36	31.01	4.97	1.71	3.34
13 TeV $W^+ \rightarrow \mu\nu$	8.59	32.73	5.07	1.05	2.00

Table 8.5 – The decomposition of the experimental uncertainty for the fit in separate channels. Unit in [MeV].

8.4.2 Joint fit of all the channels

The joint fit of all the eight channels involves a minimization of 3200 bins in total, with more than 500 NPs after NP reduction. 2000 bootstrap toys of pseudo data are produced to estimate the contribution of the bin-by-bin uncorrelated uncertainty. The decomposition of the uncertainty is shown in Tables 8.7, 8.8 and 8.9.

The fixed-order perturbative QCD uncertainty estimated using the CT18NNLO global PDF set is around 5 MeV. With the measurement of p_T^W distributions, the p_T^W uncertainty is expected to be well under control in the W mass fit using the m_T spectrum, and is around 2 MeV, much lower than that in the previous ATLAS 7 TeV measurement where large uncertainty is brought by the p_T^Z to p_T^W extrapolation. The leading systematic uncertainty in the low pile-up measurement is the lepton calibration systematics, including the electron energy scale uncertainty, the muon energy scale uncertainty and the muon sagitta bias uncertainty. More investigations are needed to improve the lepton calibration for the W mass measurement.

8.5 Prospects

So far, preliminary fit for the low pile-up W mass measurement has been carried out based on the simulation and data-driven multijet background. The total uncertainty

Channel	MC background	PDF	p_T^W
5 TeV $W^- \rightarrow e\nu$	3.41	12.16	2.13
5 TeV $W^+ \rightarrow e\nu$	2.56	15.43	1.41
5 TeV $W^- \rightarrow \mu\nu$	3.99	12.44	2.28
5 TeV $W^+ \rightarrow \mu\nu$	3.97	16.80	2.50
13 TeV $W^- \rightarrow e\nu$	3.29	15.11	3.00
13 TeV $W^+ \rightarrow e\nu$	2.50	16.40	2.68
13 TeV $W^- \rightarrow \mu\nu$	4.01	14.20	2.17
13 TeV $W^+ \rightarrow \mu\nu$	2.70	16.69	2.75

Table 8.6 – The decomposition of the modelling uncertainty for the fit in separate channels. Unit in [MeV].

Total	Stat.	Syst.
18.02	14.78 ± 0.23	10.40

Table 8.7 – The total uncertainty, statistical uncertainty and total systematic uncertainty of the joint fit. Unit in [MeV]. The error on the statistical uncertainty is caused by the limited number of bootstrap toys.

of the fit is about 20 MeV when the W mass is extracted from the unrolled m_T distribution. Further optimization of the categorization in u_T and η_ℓ may help to bring down the total uncertainty by a certain amount. Also, some studies related to statistics are on-going, which are expected to facilitate the combination of statistically correlated observables in the profile likelihood fit.

One possible solution to improve the electron calibration uncertainty, which currently contributes as high as 7 MeV to the total uncertainty, would be to extrapolate from the full Run 2 dataset to help reduce the uncertainty in electron energy scale and resolution. The sources of high pile-up systematics, such as the temperature effect and difference in the topo-cluster noise threshold, however, will need to be handled carefully during the extrapolation.

The application of the data-driven multijet estimation to the m_W measurement is still preliminary. More efforts are anticipated on the validation and improvement of the method, including the optimization of the function form for the extrapolation of the yield, the smoothing of the multijet shape in the Signal Region and so on.

On the physics modelling side, Powheg EW [34][35] may become an alternative tool for the estimation of electroweak uncertainty. As for the QCD modelling, possible improvements in the PDF uncertainties can be brought by the measurements of W and Z differential cross sections. The effect of missing higher-order QCD corrections has to be accounted for in the future. The spread of the central value of the fit using different global PDF set needs to be studied as well, which involves the development of a flexible tool for the PDF reweighting. The propagation of p_T^W uncertainty from the direct measurement is sufficient for preliminary studies. The only caveat is that in order to ensure a good description of the recoil in the data, the function of p_T^W reweighting has to be re-derived each time when the reconstructed level PDF baseline reweighting is applied, which is not yet implemented at this stage. In the next phase, the result of p_T^W measurement can be used for a parton shower tuning, so that the correlation

Elec. SF	Elec. Calib.	Muon SF	Muon Calib.	Recoil	Lumi.	MJ
1.70	7.22	2.75	3.03	2.31	0.76	0.85

Table 8.8 – The decomposition of the experimental uncertainty for the joint fit. Unit in [MeV].

MC background	PDF	p_T^W
1.44	4.92	1.84

Table 8.9 – The decomposition of the modelling uncertainty for the joint fit. Unit in [MeV].

between p_T^W and W rapidity will be appropriately modelled in the simulation.

The W mass measurement using low pile-up data requires all the above aspects and methods to be properly understood and scrutinized. The preliminary study covers most of the above topics to some extent, revealing the great potential of low pile-up data for the precision determination of W mass. It has been shown that the low pile-up m_W measurement will achieve a comparable precision as to the existing 7 TeV results [7], with most systematic sources better under control. In particular, the improved u_T resolution in the low pile-up conditions enhances the statistical sensitivity of the m_T distribution to m_W and facilitates a direct p_T^W measurement instead of extrapolating p_T^W from p_T^Z , which has been found to effectively constrain the corresponding modelling uncertainty in the m_W measurement. Furthermore, the use of profile likelihood fit in the statistical analysis constrains all systematic sources using the power of data statistics.

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4904 Uncertainty components in profile 4905 likelihood fit

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9.1 Uncertainty decomposition for profile likelihood fit

9.1.1 Measurement and uncertainty decomposition

In order to facilitate physics interpretation of the result, the measurements of physics quantities are often reported not only the quoting total uncertainty, but also including the breakdown into components with reasonable physics meanings, such as statistical uncertainty and one or multiple sources of systematic uncertainty.

In this study, the statistical uncertainty is defined by its frequentist interpretation, that is the standard deviation of a statistical estimator when the exact same experiment is repeated with the same systematics on independent data samples of the identical expected size. However, depending on the design and the technical details of the statistical analysis, the decomposition of the total uncertainty can be carried out in different ways based on the choice of test statistic used for the measurement.

In the simplest case, the statistical analysis compares the measurement of a quantity or distribution to a model that is parameterized by physical constants to be determined. The auxiliary parameters of systematic effects are fixed to their best estimates in the model. The measured central values of the physics constants are derived from the maximum of the likelihood, while the curvature of the likelihood contour near the maximum reflects the lower bound of statistical uncertainty of the measurement according to the Cramér–Rao bound. The systematic uncertainties relevant to the measurement are assessed by repeating the procedure using varied models for which auxiliary parameters are changed within their uncertainty. Each of these variations represents a source of uncertainty. And the resulting uncertainties are often uncorrelated by construction, such that the total uncertainty is the quadratic sum of statistical uncertainty and all the components of systematic uncertainties. An example of the measurement using this analysis strategy is the ATLAS 7 TeV m_W measurement [1].

An improved method of statistical analysis involves parameterizing the model with both the physical constants and the sources of systematic uncertainty [2][3]. The best values of the physical constants and the uncertainty parameters are determined simultaneously. The maximum of the likelihood in this model therefore corresponds to the global optimum for all fit parameters, while the curvature of the likelihood contour represents the total uncertainty, or in other words, the combined fluctuations of the data and all sources of uncertainties. The current practice for breaking down the statistical and systematic uncertainties in these fits is referred to as impact method [4]–

[6], following which the contribution of a given source of systematics comes from quadratic difference between the total uncertainty of the fit including or excluding this source in the model. The statistical uncertainty in the impact method is, on the other hand, determined by the uncertainty of the fit excluding the presence of all systematics in the model. Although the impacts quantify the inflation in the total uncertainty when a new source of systematic is added to a measurement, by no means they are supposed to be interpreted as the contribution of the systematic sources to the total uncertainty of a complete measurement. The supporting reasons are that the quadratic sum of statistical uncertainty and all the impacts do not reproduce the total uncertainty, and that the impacts do not obey the usual uncertainty decomposition formulas even if the uncertainties are genuinely Gaussian, as it has been pointed out in [7]. The solution to a coherent uncertainty decomposition will be provided in Section 9.1.2.

In the sections below, the measurements and the predictions in a model will be denoted as $\vec{m}$ and $\vec{t}$, labelled by indices i . The predictions are functions of the parameters of interest (POIs) θ and nuisance parameters (NPs) for systematics $\vec{a}$ labelled r . The sources of uncertainty are denoted as $\vec{a}$, corresponding to $\vec{a}$.

9.1.2 Uncertainty decomposition from shifted observables

Once the form of the probability model is given for a measurement, post-fit central value of a POI is solely dependent on the global observable and the (pseudo) data in the input, such that:

$$\hat{\theta} = \hat{\theta}(\vec{m}, \vec{a}) \quad (9.1)$$

In the most common cases, for instance, the extraction of a physical constant from a reconstructed level kinematic distribution, the measurements and global observables are uncorrelated at source level. If so, regardless of the distribution of the uncertainty distributions in the probability model, the post-fit uncertainty of the POI can be expressed via a linear error propagation:

$$\Delta\theta = \sqrt{\sum_i \left(\frac{\partial \hat{\theta}}{\partial m_i} \Delta m_i \right)^2 + \sum_r \left(\frac{\partial \hat{\theta}}{\partial a_r} \Delta a_r \right)^2} \quad (9.2)$$

where Δm_i and Δa_r are the standard deviations of m_i and a_r , respectively. The first sum under the square root originates from the fluctuations of the data, while the second term sums up the contributions from all the systematic sources. Following the spirit of this formula, the contribution of a given source of uncertainty can be estimated by shifting the corresponding measurement (for data statistics) or global observable (for systematics) by one standard deviation in the expression of the likelihood, then repeating the fit with everything else remaining the same in the probability model. The difference between the post-fit value of $\hat{\theta}$ before and after shifting the measurement or global observable is the corresponding uncertainty. This prescription allows a coherent uncertainty decomposition, since the quadratic sum of all uncertainty sources adds up to the total uncertainty, and the contribution of each uncertainty source is uniquely determined.

Instead of shifting measurements or global observables by one standard deviation, an equivalent procedure of uncertainty decomposition based on Equation 9.2 would be to fluctuate all the bins m_i together according to the statistical uncertainty in

each bin, or to fluctuate the a_r 's in interest according to their distributions given by the auxiliary measurements. This procedure generates bootstrap toys of the data or those of the global observables. The spread of the bootstrap toys of m_i gives the data statistical uncertainty. Similarly, the spread of the bootstrap toys of a_r corresponds to the systematic uncertainty arising from the sources of which the global observables are fluctuated. The bootstrap error ($Error_{BS}$) in the involved uncertainty components (σ_{BS}) due to limited number of bootstrap toys (N_{BS}) is estimated by

$$Error_{BS} \approx \frac{\sigma_{BS}}{\sqrt{2N_{BS}}} \quad (9.3)$$

In particular, the bootstrap equivalence of the shifted observable approach can be more convenient when the number of measurements or global observables are way more than the number of toys needed to achieve the desired numerical precision.

9.2 Illustrative example: W mass fit on pseudo data

In order to validate the method of uncertainty decomposition discussed in Section 9.1.2, a numerical exercise featuring m_W measurement is provided in this section. The statistical analysis of the measurement is a shape analysis that uses profile likelihood fit, where the fit to the distribution is parameterized by a POI and several NPs accounting for the effect of systematics. While the effect of varying the hypothetical value of m_W is parameterized by the POI in the profile likelihood fit, three representative sources of systematic uncertainty of a m_W measurement encountered in hadron colliders [8]–[11] are parameterized by the corresponding NPs: the lepton momentum scale uncertainty, the hadronic recoil resolution uncertainty and the p_T^W modelling uncertainty. The mass of W-boson can be extracted from either p_T^ℓ or m_T distributions. However, the sensitivity to certain types of systematics is very different depending on the choice of the kinematic spectrum.

Since the measurement of m_W is a typical shape analysis, the conclusions drawn from this example can be in principle generalized to all kinds of shape analysis.

9.2.1 MC samples

The signal process being simulated is the charge-current Drell-Yan process [12] $pp \rightarrow W^- \rightarrow \mu^- \nu$ at $\sqrt{s}=13$ TeV. The samples are generated using Madgraph+Pythia8 [13][14], of which the details of the generation is listed in Table 9.1. No background process is considered in this exercise.

The m_W templates are produced with Breit-Wigner reweighting [15] of the nominal signal MC. The systematic variation of p_T^W modelling is implemented by first using a linear reweighting as a function of the p_T^W : $w(p_T^W) = 0.96 + 8 \times 10^{-4} \times p_T^W [\text{GeV}]$ before the event selection, then taking the shape effect on the underlying p_T^W distribution. The reweighting is fixed at 1.04 for $p_T^W > 100$ GeV. At reconstructed level, the p_T of the bare muon is smeared by 2 % following a Gaussian distribution. A source of systematic uncertainty in the muon momentum scale is considered. This variation changes the muon momentum scale by 0.5 per-mille. For simplicity, the muon reconstruction and identification is considered to be 100 % efficient. The hadronic recoil $\vec{u}_T$ is taken to be the opposite of $\vec{p}_T^W$ smeared by a constant 6 GeV in both directions of the transverse plane. The variation in the recoil resolution uncertainty is made by changing the

Event Generator	
$pp \rightarrow W^- \rightarrow \mu^- \nu_\mu$ at $\sqrt{s}=13$ [TeV]	
Integrated luminosity	76.42 [pb^{-1}]
Number of events	10,000,000
Matrix elements	Madgraph at LO
Input m_W	80.419 [GeV]
Input Γ_W	2.0476 [GeV]
Parton shower & QED FSR	Pythia8
MC sample weight	0.05

Table 9.1 – Madgraph+Pythia8 event generation for MC samples. Events with an off-shell boson are excluded in the event generation at parton level, leading to a total cross-section of 6543 pb.

recoil resolutions by 5 per-mille in both directions. The recoil resolution uncertainty is taken as the second source of experimental systematics.

The bare muon is considered to be the muon at reconstructed level. The neutrino missing energy $\vec{E}_T^{\text{miss}}$ is calculated vectorially from the reconstructed hadronic recoil and the reconstructed muon. The m_T spectrum is calculated from the reconstructed level muon p_T and $\vec{E}_T^{\text{miss}}$. As listed in Table 9.2, the event selection is chosen to be as close as possible to the design of a realistic W mass measurement at hadron colliders. The cut-flow efficiency for these signal events is around 29 %.

The reconstructed level muon p_T and m_T distributions are shown in Figure 9.1. All W mass templates and systematic variations are normalized to the integral of the nominal signal MC in the fit range. The distortion on the spectra brought by the lepton energy scale variation resembles the effect of changing the hypothetical value of m_W . The muon p_T distribution is sensitive to the p_T^W modelling uncertainty but almost unaffected by the recoil resolution variation. In contrast, the m_T distribution is insensitive to p_T^W variation by construction, but more sensitive to the recoil resolution uncertainty, since the hadronic recoil enters the computation of m_T .

Event selection	
η_ℓ selection	$[-2.5, 2.5]$
p_T^ℓ selection	>25 [GeV]
E_T^{miss} selection	>25 [GeV]
m_T selection	>50 [GeV]
u_T selection	<25 [GeV]

Table 9.2 – Event selection for the Madgraph+Pythia8 samples.

9.2.2 Validation of uncertainty decomposition

The profile likelihood fit is performed using HistFactory [16] and RooFit [17], the same framework for the statistical analysis used in the low pile-up m_W measurement. The output of the fit returns both the post-fit central values and the post-fit uncertainties of all fit parameters. Each uncertainty component of the profile likelihood fit is obtained by first repeating the fit to the bootstrap toys of the pseudo data for

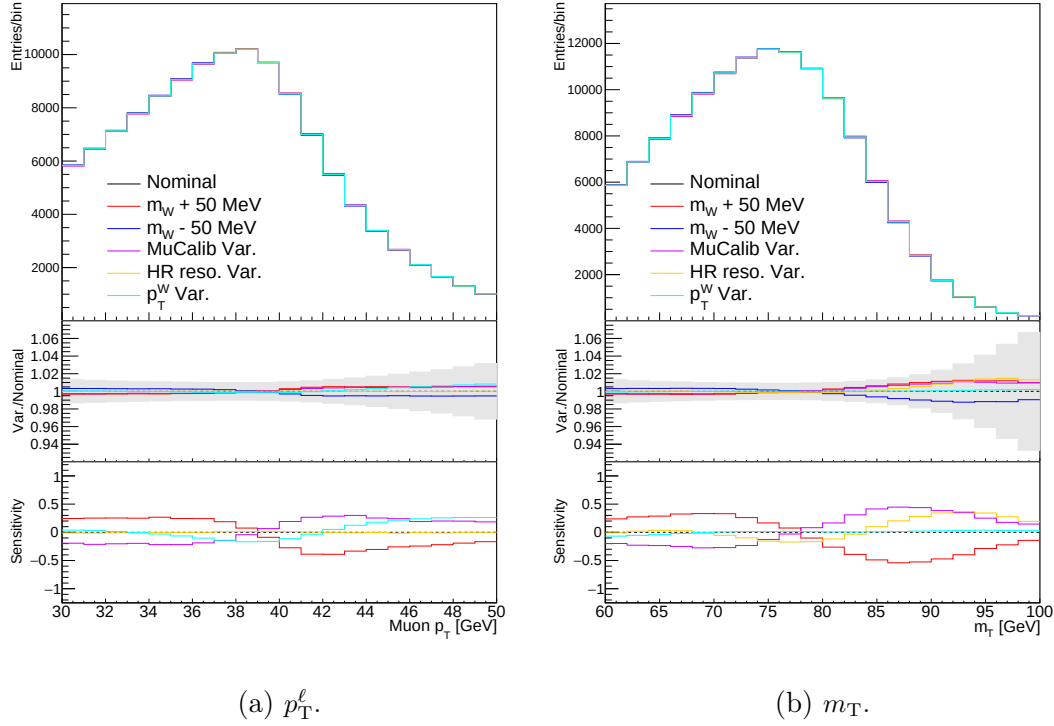


Figure 9.1 – Reconstructed muon p_T and m_T distributions of the Madgraph + Pythia8 samples. The integrated luminosity is 76.42 pb^{-1} . (top): Kinematic spectra. (middle): The variation to nominal ratio with statistical uncertainty indicated by the error band. (bottom): Sensitivity (the difference between variation and nominal, divided by statistical uncertainty) of W mass templates (unit: $[50 \text{ GeV}]^{-1}$) and systematic variations.

statistical uncertainty, or those of the global observables for systematics, then calculating the spread of the offsets in the POI. Changing the relative luminosity of the samples corresponds to a change of data statistics, for which the effect is emulated by repeating the fit with an overall luminosity scale factor multiplied to all reconstructed distributions. By using the luminosity scale factor, it doesn't require to generate more samples to match an increased relative luminosity. The setup of these validation fits are listed in Table 9.3.

Probability model	Full model	Stat. only
NPs	Lepton momentum scale	—
	Recoil resolution	
	p_{T}^W systematics	
Luminosity scale factor	0.1, 0.4, 0.7, 1.0, 2.0, 5.0, 10.0	
Fit range	$30 < p_{\text{T}}^{\ell} < 50$ [GeV] $60 < m_{\text{T}} < 100$ [GeV]	

Table 9.3 – Configuration of the m_W fits. The luminosity scale factor of 1.0 corresponds to $76.42 \text{ [pb}^{-1}\text{]}$.

The responses of uncertainty components towards relative luminosity are presented in Figures 9.2 and 9.3, with error bars computed using Equation 9.3. The quadratic sum of all uncertainty components reproduces the total uncertainty returned by the

profile likelihood fit. As one would expect, more statistics reduces both total uncertainty and statistical uncertainty of the fit. However, increasing the statistical power does not help the fit effectively constrain the muon energy scale uncertainty, mainly because it is difficult to disentangle the effect of varying muon energy scale and that of changing the assumed value of m_W . The contribution of recoil resolution systematics is low for a fit based on muon p_T distribution regardless of the statistics, while the p_T^W behaves similarly for a fit based on the m_T spectrum. The uncertainty decomposition also reveals that with more data, certain sources of systematics can be effectively constrained by the profile likelihood fit, for instance the p_T^W uncertainty in the muon p_T fit in this exercise, or the recoil resolution uncertainty in the m_T fit. Tables 9.4 and 9.5 present the uncertainty decomposition using impact approach and the shifted observable approach at two different values of the luminosity scale factor. In the shifted observable approach, the quadratic sum of the statistical uncertainty and the systematic uncertainty reproduces the total uncertainty within the tolerance of numerical errors. The statistical uncertainty estimated by the impact approach decreases by the square root of the increase in luminosity scale factor ($\sqrt{10}$), while the statistical uncertainty given by the shifted observable approach decreases slower than that rate.

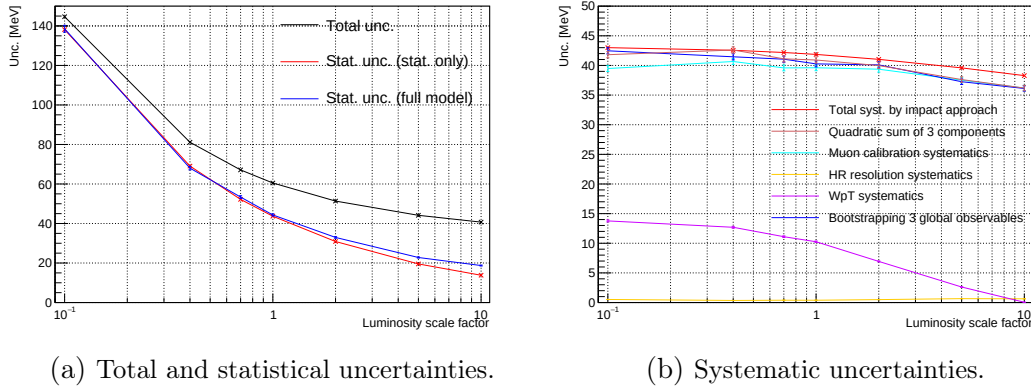


Figure 9.2 – Uncertainty decomposition for the profile likelihood fit based on muon p_T spectrum. For the total systematic uncertainty, the result of bootstrapping three global observables simultaneously is compatible with the quadratic sum of bootstrapping the three components individually.

Muon p_T fit			
Measurement	σ_{stat}	σ_{syst}	σ_{tot}
1.0 lumi. impact	43.68	41.86	60.50
1.0 lumi. shifted obs.	43.96 ± 0.55	40.99 ± 1.02	60.50
10.0 lumi. impact	13.81	38.29	40.70
10.0 lumi. shifted obs.	18.76 ± 0.23	36.19 ± 0.45	40.70

Table 9.4 – The uncertainty break-down for the muon p_T fit at two different values of the luminosity scale factor, using the impact approach and the shifted observable approach for profile likelihood fit (unit in [MeV]). The errors arise from the limited number of bootstrap toys. The baseline luminosity is $76.42 \text{ [pb}^{-1}\text{]}$.

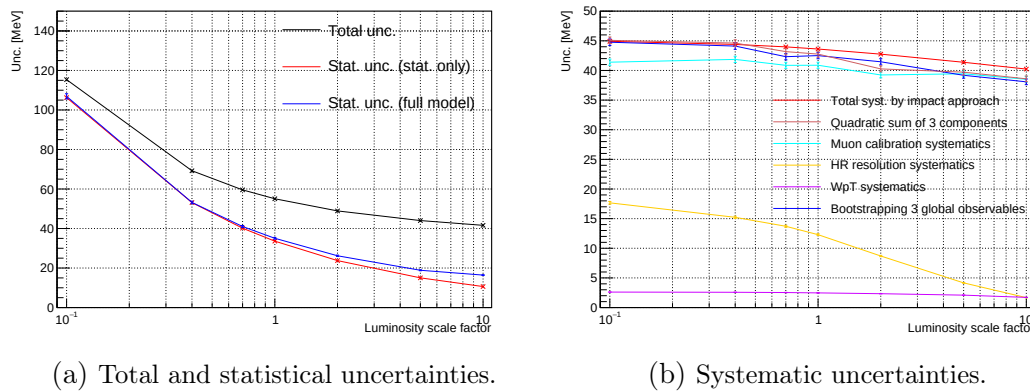


Figure 9.3 – Uncertainty decomposition for the profile likelihood fit based on m_T spectrum. For the total systematic uncertainty, the result of bootstrapping three global observables simultaneously is compatible with the quadratic sum of bootstrapping the three components individually.

m_T fit			
Measurement	σ_{stat}	σ_{syst}	σ_{tot}
1.0 lumi. impact	33.60	43.62	55.06
1.0 lumi. shifted obs.	34.41 ± 0.43	39.14 ± 0.98	55.06
10.0 lumi. impact	10.63	40.23	41.61
10.0 lumi. shifted obs.	15.73 ± 0.20	38.59 ± 0.48	41.61

Table 9.5 – The uncertainty break-down for the m_T fit at two different values of the luminosity scale factor, using the impact approach and the shifted observable approach for profile likelihood fit (unit in [MeV]). The errors arise from the limited number of bootstrap toys. The baseline luminosity is 76.42 [pb⁻¹].

5084 The dependence of post-fit uncertainties of the NPs on the relative luminosity is
 5085 shown in Figure 9.4. The responses of NPs are consistent with those of the corre-
 5086 sponding source of systematic uncertainty in Figures 9.2 and 9.3.

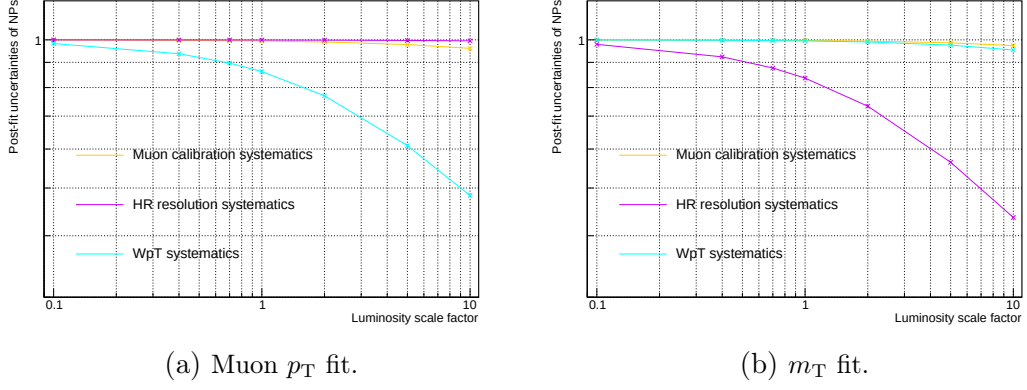


Figure 9.4 – Post-fit uncertainties of the NPs at different values of the luminosity scale factor.

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Chapter 10

Conclusion

This thesis presents the work on the production and qualification of NSW Micromegas for the ATLAS Phase-I upgrade, as well as the measurements of W-boson properties at 5.02 and 13 TeV center-of-mass energies using low pile-up data recorded by ATLAS.

The increasing luminosity of LHC during Run 3 demands improved rate capability of detectors on the inner end-caps of the ATLAS muon system. The old inner end-caps of the muon spectrometer have been replaced by the NSWs during the Phase-I upgrade. The NSWs comprise two complementary detector technologies to combine the fast trigger capability given by the sTGC, and the tracking capabilities obtained with the Micromegas. Micromegas, one of the main micro-pattern gaseous detectors, are capable of sustaining a high flux rate of 15 kHz cm^{-2} while keeping a high tracking precision of $100 \mu\text{m}$ per detecting layer, thanks to the micro-mesh technology and the separation of the conversion region and the amplification region. All the LM1 Micromegas modules are produced and qualified at Saclay. With the setup of the cosmic bench at Saclay, in particular in the presence of the three layers of 2D external trackers, the performance of the LM1 Micromegas modules are examined with a high granularity in terms of efficiency and gain. The 2D maps of efficiency and gain are produced per detecting layer for each LM1 Micromegas module. The efficiency loss due to the structure of interconnection in the module has also been estimated. An improved mixture of working gas for the NSW Micromegas was initially proposed by Saclay and tested on the cosmic bench. With the improved working gas compared to the default choice used on the cosmic bench characterization for most of the time, the NSW Micromegas benefits from attaining the same gain in a lower high voltage, therefore suppressing the sparking rate and the loss of efficiency. Comprehensive follow-up studies have been carried out at CERN in parallel to and after the production of the NSW modules, proving the feasibility of this improved working gas mixture for the NSW Micromegas in Run 3. So far, no significant aging effect has been observed with the improved gas mixture. All Micromegas modules have been delivered and mounted on the NSWs. With the NSW commissioned and operational, ATLAS will retain the performance of muon reconstruction at higher instantaneous luminosity during the Run 3 data taking.

The measurements of W-boson properties benefit from the low pile-up data taking conditions in specialized datasets, thanks to which the resolution of the hadronic recoil is largely improved with respect to the higher pile-up datasets. The thesis covers the description of the Run 2 low pile-up datasets, as well as an introduction to the

dedicated detector calibration for the low pile-up datasets. The multijet background contamination in the W event selection is estimated by a data-driven method, with improvements in the determination of the yield: A shape correction to the multijet templates in the control region, and the compensation for a bias introduced by the reconstruction of the hadronic recoil. The measurement of the p_T^W spectrum is mostly driven by the need for a precise determination of W-boson mass, which allows to perform stringent tests of the Standard Model. The low pile-up p_T^W measurement adopts the Iteration Bayesian Unfolding to infer the underlying p_T^W distribution from the reconstructed level hadronic recoil spectrum. A data-driven boson- p_T correction is implemented in the analysis to correct for the large mis-modelling of boson- p_T distribution in the default simulation. Bias uncertainties associated with the boson- p_T correction are assessed, and are taken into account in the optimization of the regularization parameters for the unfolding. The unfolded p_T^W spectra in the electron channels are combined and compared with various theoretical predictions. Several discrepancies between the measurement and the predictions indicates the needs for an improved QCD modelling in the predictions.

The low pile-up m_W measurement benefits from the calibration work of the p_T^W measurement. A profile likelihood fit has been set up for the extraction of m_W from the reconstructed level kinematic distributions, and the effects of systematic uncertainties are parameterized into the probability model of the fit. Even though the final strategy for the boson- p_T modelling has not yet been decided, the results of the low pile-up p_T^W measurement are tentatively propagated to the m_W measurement as a preliminary study, revealing an effective constraint on the modelling of p_T^W distribution. Efforts are also spent on the improvement of the statistical analysis, including the reduction of the number of nuisance parameter in the probability model using a standard procedure for experimental systematics, as well as the implementation of the data scaling technique for MC statistical uncertainty. The possibilities of combining the results from the profile likelihood fit using statistically correlated observables are briefly discussed. In addition, a coherent uncertainty decomposition for profile likelihood fits has been introduced for the first time, and applied to the preliminary study of low pile-up m_W measurement. The formulation and the validation of the uncertainty decomposition approach are summarized in a stand-alone chapter.

As an on-going analysis, the low pile-up m_W measurement requires not only further improvement in the detector calibration, but also in the theoretical modelling of W-boson production and decay, as well as the estimation of relevant uncertainties. The low pile-up m_W measurement is expected to demonstrate the power of the profile likelihood fit in terms of constraining the theoretical uncertainties, and provide new results complementary to the previous ATLAS measurement using 7 TeV dataset, furthermore improving the world average of direct m_W measurements. In the meantime, the low pile-up data has shown its great potential for precision electroweak measurements in ATLAS. In particular, the improved detector resolution with the low pile-up conditions enhances the statistical sensitivity of the measurement and helps better constrain p_T^W modelling uncertainties.