

A Search for $Z\gamma$ Resonances and TileCal Performance Studies in the ATLAS Experiment

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Declaration

I, Ntsoko Phuti Rapheeha (Student number: 833381), declare that the thesis titled: **A Search for $Z\gamma$ Resonances and TileCal Performance Studies in the ATLAS Experiment** is my original work. The thesis is being submitted to the University of the Witwatersrand, Johannesburg, in fulfillment of the requirements for the degree of Doctor of Philosophy. I make the following declarations:

1. I affirm that the submitted work is entirely my own, with proper referencing for external ideas.
2. I am aware that the University of the Witwatersrand may take disciplinary action against me if it is believed that the submitted work is not my original, unaided work or if I have neglected to acknowledge the sources of my ideas in my writing.
3. Several results presented in this thesis have been previously discussed in peer-reviewed journals. The list of publications on page [iv](#) provides details of the contributions I have made during my doctoral studies.

Johannesburg, South Africa



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Abstract

The observation of the Higgs boson by the ATLAS and CMS Collaborations at the LHC in 2012 marked a pivotal moment, completing the Standard Model, SM, of particle physics. While the SM has propelled our understanding forward with remarkable success, its brilliance is shadowed by its inability to account for several fundamental aspects of nature, such as the existence of dark matter, the mysterious origins of neutrino masses, and the glaring absence of gravity within its framework.

Motivated by the discrepancies observed in multi-lepton final states at the LHC, this thesis conducts a search for high-mass spin-0 and spin-2 resonances decaying into the $Z\gamma$ final state utilising the Run 2 dataset collected by the ATLAS detector at $\sqrt{s} = 13$ TeV, with an integrated luminosity of 140 fb^{-1} . A mild excess of 2.1σ is observed at 250 GeV, a range of interest predicted by the model that accounts for the multi-lepton anomalies. Despite the localised excesses, no significant deviation relative to the SM background is observed. Upper limits, established at a 95% confidence level, are set on the production cross-section times the decay branching ratio into $Z\gamma$ for high-mass resonances in the mass range of 220 GeV to 3400 GeV. These upper limits improve on Run 1 results by a factor of 1.9 to 4 in the mass range between 250 GeV and 2400 GeV.

This thesis also delves into performance studies of the ATLAS Tile calorimeter, TileCal. The TileCal is a hadronic sampling calorimeter consisting of scintillating tiles that are sandwiched between steel slabs. The calibration and uniformity of the response of gap and crack scintillators in TileCal are validated by studying their response to isolated muons. The response of these scintillators to isolated muons, produced by $W \rightarrow \mu\nu_\mu$ decays in pp collisions during Run 2 data-taking, is compared to Monte Carlo simulations. Additionally, preliminary luminosity measurements from TileCal D cells are used to analyse drifts in the primary ATLAS luminosity measurements provided by the LUCID detector during the 2022 data taking at $\sqrt{s} = 13.6$ TeV.

Publications

Journal Articles

- Andreas Crivellin, Yaquan Fang, Oliver Fischer, Abhaya Kumar, Elias Malwa, Bruce Mellado, Ntsoko Rapheeha, Xifeng Ruan and Qiyu Sha. “Accumulating evidence for the associated production of a new Higgs boson at the LHC”. American Physical Society, vol. 108, no. 11, p. 1150131 (2023). DOI: [10.1103/PhysRevD.108.115031](https://doi.org/10.1103/PhysRevD.108.115031).
- Aad, Georges and *et al.*. “Search for the $Z\gamma$ decay mode of new high-mass resonances in pp collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector”. Phys. Lett. B, vol. 848, no. 11, p. 138394 (2024). DOI: [10.1016/j.physletb.2023.138394](https://doi.org/10.1016/j.physletb.2023.138394)

Conference Proceedings

- N P Rapheeha, B Mellado. “Performance of the C10 cells of the Tile Calorimeter of the ATLAS detector during Run 2 data taking period”, in The Proceedings of SAIP2023, the 67th Annual Conference of the South African Institute of Physics, edited by Prof. Aletta Prinsloo, (UJ), pp. 224 - 229. ISBN: 978-0-7961-3774-6.
- N P Rapheeha, B Mellado. “Time stability of the response of gap/crack scintillators of the Tile Calorimeter of the ATLAS detector to isolated muons”, in The Proceedings of SAIP2022, the 66th Annual Conference of the South African Institute of Physics, edited by Prof. Aletta Prinsloo, (UJ), pp. 182 - 189. ISBN: 978-0-6397-4426-1.
- N P Rapheeha, X Ruan, G Mokgatitswane, S Dahbi, “Background decomposition in $Z\gamma$ events used in the search for high-mass resonances”, in The Proceedings of SAIP2021, the 65th Annual Conference of the South

African Institute of Physics, edited by Prof. Aletta Prinsloo, (UJ), pp. 108 - 113. ISBN: 978-0-620-97693-0

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Chapter 1

Introduction

Particle physics is a scientific discipline focused on understanding properties of fundamental particles and their interactions. Two types of fundamental particles exist in nature: fermions and bosons. Fermions serve as the elemental building blocks of both matter and antimatter, characterised by their half-integer quantum spin and obedience to Fermi-Dirac statistics. On the other hand, bosons are particles that facilitate interactions between fermions, distinguished by their integer quantum spin and obedience to Bose-Einstein statistics. In subsequent sections, elementary particles and their interactions will be discussed.

1.1 Fundamental Particles

1.1.1 Fermions

There are twelve fundamental fermions, divided into six quarks of different flavours, three electrically charged leptons (electron, muon, tau) and their corresponding three neutral neutrinos (electron, muon and tau neutrinos), as illustrated in Figure 1.1. Leptons can exist independently, while quarks only exist in composite states bound by gluons. Colour confinement prevents the isolation of individual quarks.

Particles like protons and neutrons are colour-neutral bound states of quarks collectively known as hadrons [1]. Quarks come in three colour charges, and they have six flavours: up, down, strange, charm, bottom and top. Three quarks hadronise to form baryons with a spin that is a factor of half, and two quarks combine to form mesons with quantum spin of zero or one. Fermions are subjected to the Pauli exclusion principle; no quarks with the same quantum numbers can combine to form a bound state.

The twelve fermions are arranged into three generations with similar properties apart from the fermion's mass. The masses increase as you go from generation I to generation III and as you go from the bottom of each generation to the top, as illustrated in Figure 1.1. Hadrons formed by generation I fermions are generally long-lived when compared to hadrons that are composed of other generations of fermions. The top quark is the heaviest fermion, boasting a mass of approximately 173.0 ± 0.4 GeV and a mean lifetime of roughly 5×10^{-25} s. It does not live long enough to hadronise and its existence is only inferred from its decay products observed in detectors [2].

1.1.2 Bosons

In the current Standard Model, SM, of particle physics there are four spin-1 bosons and one spin-0 boson. The spin-1 bosons include gluons, photons, and the W and Z gauge bosons. They have a vector field and serve as the exchanged particles in the interactions of fermions. Photons are massless particles coupled to the electric charge and they are exchanged in electromagnetic interactions. Gluons come in 8 flavours coupled to the colour charge and they serve as the exchange particles in strong nuclear interactions. The charged W , $W^\pm$, bosons and the neutral Z boson are intermediate vector bosons exchanged in weak interactions. The W s and Z bosons are the only force carrying bosons with mass, and the explanation of this phenomenon required the $SU(2)$ symmetry to break, the electroweak symmetry breaking will be discussed at length in Section 1.2.4. The Higgs boson, a spin-0 particle, is responsible for the electroweak symmetry breaking. It was observed in 2012 at the Large

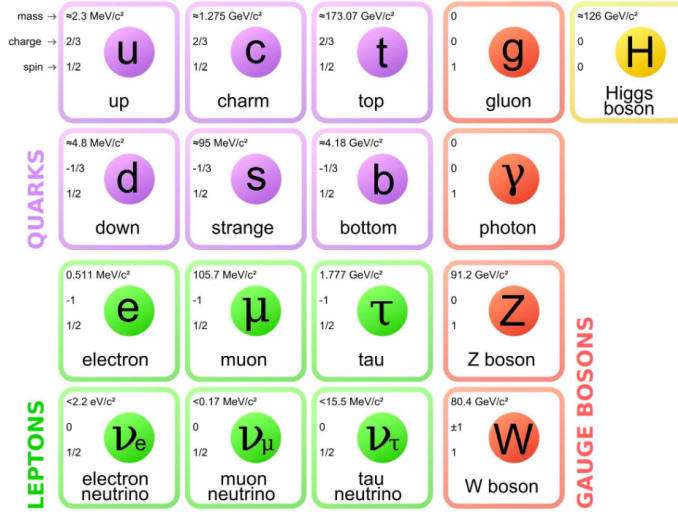


Figure 1.1 All known fundamental particles that make up the SM of particle physics, with the first three columns representing the three generations of leptons, the fourth column representing the force-carrying bosons, and the last particle being the Higgs boson [3].

Hadron Collider, LHC, by the CMS [4] and ATLAS [5] Collaborations and it is the last fundamental particle to be added to the SM.

1.2 Interactions of Fundamental Particles

The interactions governing all matter around us are primarily dictated by four fundamental forces: gravitational, electromagnetic, weak, and strong nuclear interactions. The strong and weak interactions exhibit high strength at the subatomic scale and they govern nuclear interactions. The gravitational and electromagnetic interactions are relatively long range forces when compared to the range at which the strong and weak interactions are effective.

As it can already be seen, the current SM of particle physics has its own limitations. It does not include a boson that mediates gravitational interactions. There are many theories that predict the existence of an additional boson known as the graviton, a hypothetical particle associated with the gravita-

tional field. If this hypothetical particle is to be observed, it is expected that it would possess a spin-2 quantum number, stemming from its association with a tensor gravitational field.

1.2.1 Strong Interactions

The Gell-Mann and Zweig model of quarks [6, 7] successfully classified hadrons but it falls short in explaining the dynamics of quark interactions. Quantum Chromodynamics, QCD, a quantum field theory based on the non-Abelian $SU(3)_C$ gauge group, successfully describes the dynamics of the strong interaction by confining quarks to form bound states. To describe the dynamics of quarks, let us consider a Lagrangian:

$$\mathcal{L} = \sum_f \bar{q}_f (i\gamma^\mu \partial_\mu - m_f) q_f, \quad (1.1)$$

where q_f denotes a quark field of flavour f . Replacing the partial derivative with a covariant derivative ensures that the quark field is invariant under local $SU(3)$ transformation:

$$D^\mu q_f = \left(\partial^\mu - ig_s \frac{\lambda^a}{2} G_a^\mu \right) q_f, \quad (1.2)$$

where $\frac{1}{2}\lambda^a$ ($a = 1, 2, 3, \dots, 8$) represents the Gell-Mann matrices denoting the 8 generators of the $SU(3)$ group, G_a^μ is a vector field for the 8 gauge bosons, and g_s is the coupling strength of the strong interaction [8, 9] given by:

$$g_s = \sqrt{4\pi\alpha_s}. \quad (1.3)$$

The propagation of the quark field is facilitated by introducing an $SU(3)$ gauge-invariant field tensor $G^{\mu\nu}$:

$$G_a^{\mu\nu}(x) = \partial^\mu G_a^\nu - \partial^\nu G_a^\mu + g_s f^{abc} G_b^\mu G_c^\nu, \quad (1.4)$$

where f^{abc} represent the QCD structure constants. The requirement for an invariant Lagrangian leads to the conclusion that gauge bosons must carry a colour charge, resulting in a self-coupling field. The local gauge invariant QCD Lagrangian is given by:

$$\mathcal{L}_{\text{QCD}} = -\frac{1}{4}G_a^{\mu\nu}G_{\mu\nu}^a + \sum_f \bar{q}_f(i\gamma^\mu\partial_\mu - m_f)q_f, \quad (1.5)$$

and, as consequence of gauge invariance, gluons are massless.

At low energies, the force between two coloured particles is constant, as a result, the energy required to separate quarks continues to increase until it generates a quark and anti-quark pair, turning the original hadron into a pair of hadrons. This phenomenon is called colour confinement. Accelerated quarks radiate gluons, which in turn interact with themselves and other quarks, making theoretical studies of colour confinement very complex. At high energies, the strength of quark-gluon interactions decreases, resulting in the formation of quark-gluon plasma. This property of QCD is known as the asymptotic freedom and it enables the treatment of partons as free particles.

1.2.2 Quantum Electrodynamics

James Clerk Maxwell unified electricity and magnetism into a single theory of electromagnetism which explained the nature of light. To describe the interactions of photons and matter, Dirac's equations of relativistic quantum mechanics were unified with Maxwell equations of electromagnetism into a unified Abelian gauge theory with a symmetry group $U(1)$ called Quantum Electrodynamics, QED [10]. In QED, electrically charged particles interact by exchanging photons. The coupling strength of photons to the electric charge current is given by:

$$\alpha = \frac{ke^2}{\hbar c} = \frac{1}{137}, \quad (1.6)$$

where k is the Coulomb's constant and e is the elementary electric charge. This constant is commonly referred to as the fine structure constant. The

commonly known law of conservation of electric charge arises as a consequence of gauge invariance of QED.

1.2.3 Weak Interactions

Fermions interact weakly through the exchange of W and Z bosons. There are two types of weak interactions, charged-current interactions involving electrically charged $W^\pm$ bosons that are responsible for the emission and absorption of neutrinos. The other type of weak interaction, the neutral-current interactions, involves the electrically neutral Z boson and it is responsible for momentum transfer in elastic scattering of neutrinos.

At low energies, the strength of weak interactions is quantified by the Fermi coupling constant, which scales with the square of the energy E involved in the weak process, expressed as $G_F E^2$.

1.2.4 Electroweak Theory

The electroweak theory serves as a unified model of electromagnetism and weak interactions. The QED Lagrangian maintains the $U(1)$ local gauge symmetry, forbidding vector bosons from acquiring masses. To accommodate the observed massive $W^\pm$ and Z bosons, the $SU(2)$ local gauge theory is introduced. Under $SU(2)$ local gauge transformations, a Lagrangian describing a free Dirac fermion given by:

$$\mathcal{L} = \bar{\Psi}(x)(i\not{\partial} - m)\Psi(x), \quad (1.7)$$

is made to be invariant by introducing a covariant derivative, D_μ , given by:

$$D_\mu = \partial_\mu + ig_w \vec{T} \cdot \vec{W}, \quad (1.8)$$

where $\vec{T}$ represents the three generators of the $SU(2)$ group. Each generator introduces three intermediate bosons: W_μ^1 , W_μ^2 and W_μ^3 , each with zero mass as required by the gauge theory. The charged $W^\pm$ bosons are defined as linear

combinations of the W_μ^1 and W_μ^2 intermediate bosons, while the W_μ^3 boson describes the neutral weak current in the Glashow-Weinberg-Salam, GWS, theory of electroweak interactions [11–13].

The GWS theory is based on the $SU(2)_L \otimes U(1)_Y$ gauge symmetry group, where the electromagnetic and weak interactions are united by a hypercharge, Y , defined as:

$$Y = 2(Q - T^3), \quad (1.9)$$

here Q denotes the electric charge, and T represents the third component of the weak isospin of the particle. The subscript L in $SU(2)_L$ refers to left-handed fermionic fields. From nature, it is observed that only left-handed fermions and right-handed anti-fermions interact through the weak interaction [14].

The $U(1)_Y$ symmetry group has an intermediate vector boson B_μ , which cannot possibly be originating from the electromagnetic field, since photons interact with both fermion chiralities equally. The two neutral states B_μ and W_μ^3 are given by:

$$\begin{pmatrix} W_\mu^3 \\ B_\mu \end{pmatrix} = \begin{pmatrix} \cos \theta_W & \sin \theta_W \\ -\sin \theta_W & \cos \theta_W \end{pmatrix} \begin{pmatrix} Z_\mu \\ A_\mu \end{pmatrix}, \quad (1.10)$$

where A_μ represents the photon, Z_μ represents the Z boson, and θ_W is the weak mixing angle [15]. Gauge symmetry does not allow explicit writing out of the mass terms, the Lagrangian mass terms are introduced in the GWS theory of electroweak interactions via the Brout-Englert-Higgs, BEH, mechanism through spontaneous gauge symmetry breakdown explained in Section 1.2.4.

At high energies, the coupling strength of weak interactions, given by α_W , is analogous to the fine-structure constant in the electromagnetic interactions, and it is given by:

$$\alpha_W = \sqrt{2}G_F \frac{m_W^2}{\pi} = \frac{\alpha}{\sin^2 \theta_W} \cong \frac{1}{30}, \quad (1.11)$$

where m_W represents the mass of the W boson, and θ_W denotes the weak mixing angle. From Equation 1.11 it is seen that weak interactions act at short ranges not due to their weak coupling but rather because they exchange massive bosons [16].

Spontaneous Symmetry Breaking and the BEH Mechanism

To generate the mass terms in the GWS theory, consider a Lagrangian describing a complex scalar field ϕ :

$$\mathcal{L} = \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - V(\phi) \text{ with } V(\phi) = \mu^2 \phi \phi^* + \lambda (\phi \phi^*)^2. \quad (1.12)$$

At the ground state, or vacuum state, the potential of the Lagrangian is equal to zero. For $\mu^2 > 0$, the minimum of the potential is at zero, and this is not an interesting solution, while for $\mu^2 < 0$, the minimum of the potential occurs at:

$$|\phi_{\min}| = \sqrt{\frac{-\mu^2}{2\lambda}} = \frac{v}{\sqrt{2}}. \quad (1.13)$$

This minimum has an infinite degenerate states lying on a circle of radius $|\phi_{\min}| = \frac{v}{\sqrt{2}} e^{i\alpha}$, where α is an arbitrary phase and the radius of the circle is given by $\frac{v}{\sqrt{2}}$. This is illustrated in Figure 1.2. Choosing one of these degenerate states, by fixing a value of α , leads to the spontaneous breaking of symmetry [17].

Expanding the complex scalar field ϕ around the vacuum state by a small perturbation, $\xi(x) + i\chi(x)$, results in:

$$\phi(x) \equiv \frac{1}{\sqrt{2}} [v + \xi(x) + i\chi(x)], \quad (1.14)$$

where ξ and χ are real fields, and zero at the vacuum state. The potential then takes the form:

$$V(\phi) = V(\phi_0) - \mu^2 \xi^2 + \lambda v \xi (\xi^2 + \chi^2) + \frac{\lambda}{4} (\xi^2 + \chi^2)^2. \quad (1.15)$$

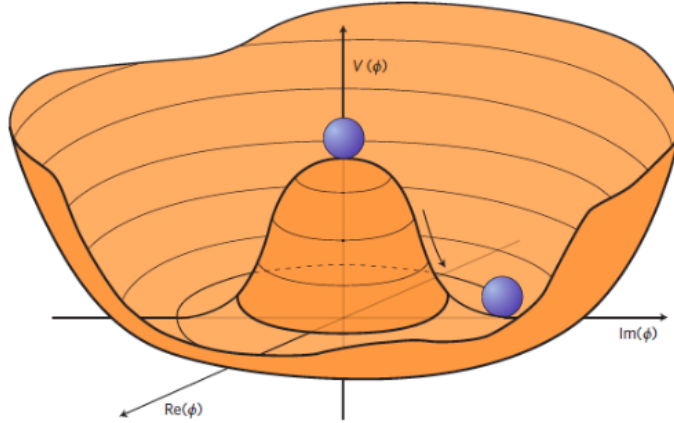


Figure 1.2 The Sombbrero, “Mexican hat”, potential that leads to spontaneous symmetry breaking. The vacuum state of the potential is described by any point chosen on the circle at the bottom brim of the hat. Movement around a brim correspond to a Nambu-Goldstone boson, while movement in the radial direction yields massive bosons [18].

The $\xi(x)$ field acquires a mass term, $\mu^2\xi^2$ describes a boson with mass $m_\xi = -2\mu^2 = \sqrt{2v^2\lambda}$, while $\chi(x)$ remains massless. The $\chi(x)$ field represents a perturbation around the plateau of the potential, therefore requiring no energy to transition between the degenerate states. The spontaneous symmetry breaking introduced to account for the massive gauge bosons introduces a massless scalar boson called Nambu-Goldstone boson, and this type of a boson is not observed in nature [19].

Spontaneous symmetry breaking breaks the global $U(1)$ symmetry, but the local gauge transformations cannot be spontaneously broken [20]. The global symmetry of the Lagrangian in Equation 1.12 can be extended to a local $U(1)$ invariance:

$$\phi(x) = e^{i\alpha(x)}\phi(x), \quad (1.16)$$

by introducing a vector field A_μ transforming according to:

$$A_\mu = A_\mu(x) + \frac{1}{e}\partial_\mu\alpha(x). \quad (1.17)$$

The local gauge invariant Lagrangian is given by:

$$\mathcal{L} = D_\mu \phi^* D^\mu \phi - \frac{1}{4} F_{\mu\nu} F^{\mu\nu} - V(\phi), \quad (1.18)$$

where $D_\mu \phi = \partial_\mu \phi + ie A_\mu \phi$ and $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ denote the covariant derivative and the field strength of the gauge field, respectively. This Lagrangian remains invariant under local $SU(2)_L \otimes U(1)_Y$ transformations. Considering a $SU(2)$ -compatible complex scalar field with neutral, $\phi^0(x)$, and electrically charged, $\phi^\pm(x)$, fields given by:

$$\phi(x) = \begin{pmatrix} \phi^+(x) \\ \phi^0(x) \end{pmatrix}, \quad (1.19)$$

and where photons solely couple to the $\phi^\pm(x)$ field. The potential, like before, has the vacuum expectation values lying around a circle of radius $\frac{v}{\sqrt{2}}$. The scalar field can be expanded around this minimum by a scalar field, $h(x)$, thereby breaking the $SU(2)_L \otimes U(1)_Y$ symmetry. As a consequence of charge conservation, only the $\phi^0(x)$ field can acquire a vacuum expectation value and become:

$$\phi(x) = e^{i\frac{\sigma_i}{2}\alpha^i(x)} \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + h(x) \end{pmatrix}, \quad (1.20)$$

where $\alpha^i(x)$ and $h(x)$ denote the 4 real fields. The $\alpha^i(x)$ fields represent the would-be massless Goldstone bosons associated with spontaneous symmetry breaking.

The local $SU(2)_L$ gauge invariance of the Lagrangian permits us to choose a gauge in which $\phi(x)$ is always real. This eliminates the Goldstone bosons, leaving us with only one real field $h(x)$, commonly called the Brout-Englert-Higgs field [21, 22]. The kinetic term in the Lagrangian in Equation 1.18 takes the form:

$$D_\mu \phi^* D^\mu \phi \rightarrow \frac{1}{2} \partial_\mu h \partial^\mu h + (v + h)^2 \left\{ \frac{g^2}{4} W_\mu^* W^\mu + \frac{g^2}{8 \cos^2 \theta_W} Z_\mu Z^\mu \right\}. \quad (1.21)$$

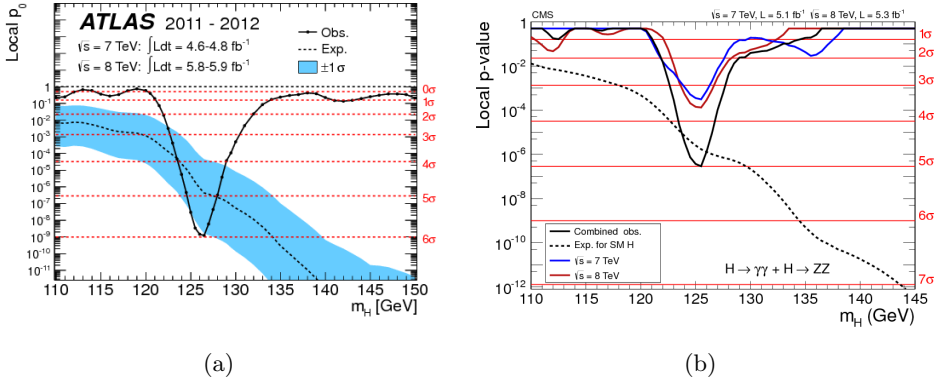


Figure 1.3 The observed local p -values for the $H \rightarrow \gamma\gamma$ and $H \rightarrow ZZ \rightarrow 4\ell$ decay channels reported by the ATLAS and CMS Collaborations. The solid black line shows the observed local p_0 values using the SM Higgs boson hypothesis, centred around a mass of approximately 125 GeV [4, 5].

The vacuum expectation of $\phi^0(x)$ causes the $W^\pm$ and Z bosons to gain mass, while photons remain massless. The mass of the boson associated with BEH field is given by:

$$m_H = \sqrt{-2\mu^2}. \quad (1.22)$$

The theory of electroweak interactions is finalised by incorporating spontaneous symmetry breaking and introducing the BEH mechanism. This completion reveals the existence of the Higgs boson, associated with the BEH field. Observing the Higgs boson was a key objective of the ATLAS and CMS experiments at the LHC, achieved in 2012 by the ATLAS and CMS Collaborations, representing a major milestone in particle physics. Figure 1.3 illustrates the observed local p -values for the $H \rightarrow \gamma\gamma$ and $H \rightarrow ZZ \rightarrow 4\ell$ decay channels reported by the ATLAS and CMS Collaborations.

1.2.5 The Standard Model Formalism

In the 1970s, numerous scientists around the world worked on a mathematical model that could comprehensively explain all fundamental particles and their

interactions. The current SM of particle physics serves as this mathematical framework, explaining all that is presently known about the strong and electroweak interactions. It is founded on the non-Abelian symmetry group $SU(3)_C \otimes SU(2)_W \otimes SU(1)_Y$, where the subscripts denote colour, weak isospin and hypercharge, respectively [23]. The gauge symmetries of these groups give rise to the fundamental interactions observed in nature.

The dynamics of a gauge field ϕ are governed by the action S , expressed as:

$$S = \int d^4x \mathcal{L}[\phi_i(x), \delta_\mu \phi_i(x)], \quad (1.23)$$

where d^4x represents integration over the four space-time coordinates, and $\mathcal{L}$ is the Lagrangian. The SM Lagrangian, $\mathcal{L}_{\text{SM}}$, takes the form:

$$\mathcal{L}_{\text{SM}} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} + i\bar{\Psi}\not{D}\Psi + h.c. + \Psi_i y_{ij} \Psi_j h + hc + |D_\mu h|^2 - V(h), \quad (1.24)$$

where $F_{\mu\nu}$ represents the field strength tensor, Ψ is the Dirac fermionic field, $\not{D}$ is the gauge covariant derivative acting on fermions, $h.c.$ denotes the Hermitian conjugate of the preceding terms, y_{ij} is the Yukawa matrix, and h is the Brout-Englert-Higgs field. The complete mathematical formulation of the SM Lagrangian is long, Equation 1.24 is an ultra-short version of the full SM Lagrangian but, it still encapsulates all the necessary information needed to explain fundamental particles and their interactions.

To understand how the SM Lagrangian explains all the fundamental particles and their interactions we need to understand what each term of the Lagrangian in Equation 1.24 represents. The term $-\frac{1}{4}F_{\mu\nu}F^{\mu\nu}$ is a scalar product of the field strength tensor, describing the interactions of all gauge bosons except the Higgs boson. It contains the necessary formulation for these particles to exist and how they interact with each other.

The terms $i\bar{\Psi}\not{D}\Psi + h.c.$ describe the interactions of gauge bosons, except the Higgs boson, with fermions, describing the electromagnetic, weak, and strong interactions. The fields $\bar{\Psi}$ and Ψ describe leptons and quarks, respectively. The term $\Psi_i y_{ij} \Psi_j h + hc$ accounts for the interactions between left-handed Ψ_i

Table 1.1 A summary of multi-lepton anomalies at the LHC [34].

Final state	Characteristics	SM backgrounds	Statistical significance
$\ell^+\ell^- + (b\text{-jets})$ [29, 35, 36]	$m_{\ell\ell} < 100$ GeV	$t\bar{t}, Wt$	$> 5\sigma$
$\ell^+\ell^- + (\text{no jets})$ [27, 37]	$m_{\ell\ell} < 100$ GeV	W^+W^-	$\approx 3\sigma$
$\ell^\pm\ell^\mp + (b\text{-jets})$ [32, 38, 39]	Moderate H_T	$t\bar{t}W^\pm, t\bar{t}t\bar{t}$	$> 3\sigma$
$\ell^\pm\ell^\mp, (\text{no } b\text{-jets})$ [31, 40, 41]	In association with h	$W^\pm h(125), WWW$	$\gtrsim 4\sigma$
$Z(\rightarrow \ell\ell)\ell, (\text{no } b\text{-jets})$ [29, 42]	$p_T^Z < 100$ GeV	$ZW^\pm$	$> 3\sigma$

and right-handed Ψ_j fermionic fields and the h field, thereby imparting mass to fermions. The Yukawa matrix represents the coupling strengths and by fixing the coupling constants, the SM correctly describes masses of fermions.

The term $|D_\mu h|^2$ explains how the gauge bosons couple to the Brout-Englert-Higgs field, describing how the Z and W bosons acquire mass. The potential $V(h)$ of the Brout-Englert-Higgs field leads to spontaneous symmetry breaking. Finally, the Higgs boson itself is coupled to the Brout-Englert-Higgs field, acquiring its mass through this spontaneous symmetry breaking mechanism.

1.3 The Multi-Lepton Anomalies at the LHC

The SM is by no means a complete theory of nature, as it was briefly discussed in Section 1.2. It does not account for dark matter, dark energy, non-vanishing neutrino masses and gravitational interactions. There is also a plethora of experimental results from the LHC, especially with two or more leptons in the final state where large disagreements with SM predictions have been observed [24–32]. These multi-lepton anomalies are summarised in Table 1.1. A more comprehensive list of anomalies in particle physics can be found in Reviews [33, 34].

The Madala hypothesis [24, 25] was initially conceived as a simplified model designed to account for the discrepancies observed between the SM predictions and experimental results published by the ATLAS and CMS Collaborations in Run 1, the first operational period of the LHC from 2010 to 2012. This

hypothesis extends the SM Lagrangian by introducing a scalar boson, H , heavier than the SM Higgs boson, h . It was proposed that this additional boson may participate in the process of electroweak symmetry breaking. The interactions within this hypothesis are described by the Lagrangian:

$$\mathcal{L} = \mathcal{L}_{\text{SM}} + \mathcal{L}_{\text{BSM}}, \quad (1.25)$$

where $\mathcal{L}_{\text{SM}}$ encompasses SM interactions, and $\mathcal{L}_{\text{BSM}}$ describes beyond the Standard Model, BSM, behavior given by:

$$\mathcal{L}_{\text{BSM}} = \mathcal{L}_H + \mathcal{L}_Y + \mathcal{L}_T + \mathcal{L}_Q, \quad (1.26)$$

the vertices $\mathcal{L}_H$, $\mathcal{L}_Y$, $\mathcal{L}_T$ and $\mathcal{L}_Q$ represent the production of H , Yukawa, trilinear, and quartic interactions, respectively. Definitions of the vertices are as follows:

$$\mathcal{L}_H = -\frac{1}{4}\beta_g\kappa_{hgg}^{SM}G_{\mu\nu}G^{\mu\nu}H + \beta_V\kappa_{hVV}^{SM}V_\mu V^\mu H, \quad (1.27)$$

$$\mathcal{L}_Y = -\frac{1}{\sqrt{2}}[y_{\bar{t}tH}\bar{t}tH + y_{\bar{b}bH}\bar{b}bH], \quad (1.28)$$

$$\mathcal{L}_T = -\frac{1}{2}v[\lambda_{Hhh}Hhh + \lambda_{h\chi\chi}h\chi\chi + \lambda_{H\chi\chi}H\chi\chi], \quad (1.29)$$

$$\mathcal{L}_Q = -\frac{1}{2}\lambda_{Hh\chi\chi}Hh\chi\chi - \frac{1}{4}[\lambda_{HHhh}HHhh + \lambda_{hh\chi\chi}hh\chi\chi + \lambda_{HH\chi\chi}HH\chi\chi], \quad (1.30)$$

where H is the proposed heavy scalar, χ is a dark matter candidate, and v denotes the vacuum expectation value, vev , a constant equal to 246 GeV.

The factor $\kappa_{hgg}^{SM} = \frac{\alpha_s}{(3\pi v)}$ couples SM gluons to the newly introduced H , and the dimensionless $\beta_g = \frac{y_{\bar{t}tH}}{t\bar{t}h}$ serve as a scaling factor relative to the SM Yukawa top coupling for H . The product $\beta_g\kappa_{hgg}^{SM}$ controls the production of H via the gluon-gluon fusion, ggF , process. The interactions of H with electroweak vector bosons are given by $\kappa_{hVV}^{SM} \simeq \frac{m_V^2}{v}$, and the branching ratios of $H \rightarrow VV$ are controlled by tuning the parameter β_V , defined in a manner similar to β_g .

The Yukawa couplings of H to heavy quarks are assumed to resemble those in the SM and are described in Equation 1.28.

The top quark couplings would explain the observed elevated production of the Higgs boson in the $h \rightarrow t\bar{t}$ channel, while the electroweak couplings would explain the observed multi-lepton production evident in both ATLAS and CMS public results. The trilinear, Equation 1.29, and quartic, Equation 1.30, couplings allow the $pp \rightarrow H \rightarrow hh, h\chi\chi$ decays. The idea was that the $H \rightarrow h\chi\chi$ decay would explain the observed distortion in the Higgs p_T spectrum, and the $H \rightarrow hh$ decay would account for excess events observed in the di-Higgs search results.

To set appropriate values of the couplings in the model, it was necessary to consider constraints from various experimental signatures predicted by the model. These constraints were derived from published experimental results by both ATLAS and CMS Collaborations during Run 1. Four classes of constraints were considered, which included measurements related to: the differential Higgs boson p_T spectrum [43–46], searches for Higgs boson pair production [47–50], top-associated Higgs boson production [51–54], and $H \rightarrow VV$ decays [55–57]. A simultaneous fit was performed by minimizing the global χ^2 value while varying the effective gluon coupling for different heavy scalar mass, m_H , hypotheses. The m_H values were varied from 260 GeV in increments of 5 GeV to 320 GeV.

Figure 1.4 displays the minimised χ^2 values for different m_H hypothesis. It is observed that the lower m_H values tend to minimise the χ^2 better with $m_H = 270$ GeV yielding the most minimised χ^2 , and the parameter β_g being best fit at a value of 1.5 ± 0.6 . Figure 1.5 shows fits to ATLAS and CMS Run 1 fiducial differential distributions of the Higgs boson p_T using the Madala hypothesis with $m_H = 270$ GeV. The combined minimised χ^2 is found to be smallest at $m_H = 272$ GeV with an upper error of 12 GeV and a lower error of 9 GeV as shown in Figure 1.4.

The Madala hypothesis extends the two Higgs doublet Model, 2HDM, by introducing two additional scalars, H and S , generating a total of five new

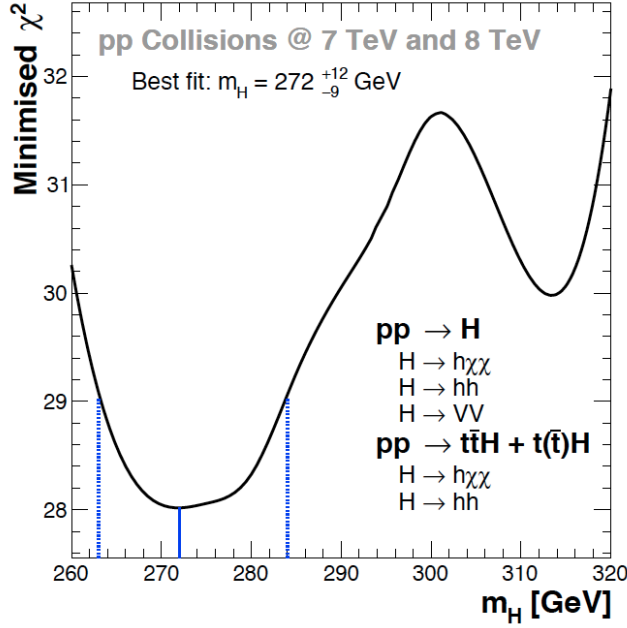


Figure 1.4 A scan of minimised χ^2 values for different m_H hypothesis. The solid blue line represents the m_H value that minimises χ^2 the most, with the two dotted blue lines indicating a 1σ deviation from central solid line [25].

scalar bosons [58, 59]. The decay process of H to dark matter candidates, $H \rightarrow h\chi\chi$, is mediated by an on-shell participation of a scalar singlet, S , which was chosen to be SM Higgs like. The mediator S facilitates the $H \rightarrow hS, SS$ decay channels, with subsequent decays $S \rightarrow \chi\chi$. The scalar S is allowed to have a mass within the range $m_h \lesssim m_s \lesssim m_H - m_h$ and $m_s > 2m_\chi$, offering a narrow yet versatile mass range for exploring various phenomenological aspects predicted by the model.

To generate the effective quartic coupling, $\lambda_{hH\chi\chi}$, thought to be responsible for the distortion observed in the Higgs boson p_T distribution, the trilinear coupling, λ_{hHS} , was generated from the $hHSS$ interaction by requiring S to acquire vev , this is illustrated in Figure 1.6. The scalar S was permitted to mix with all SM particles and χ , acting as a gateway for SM particles to

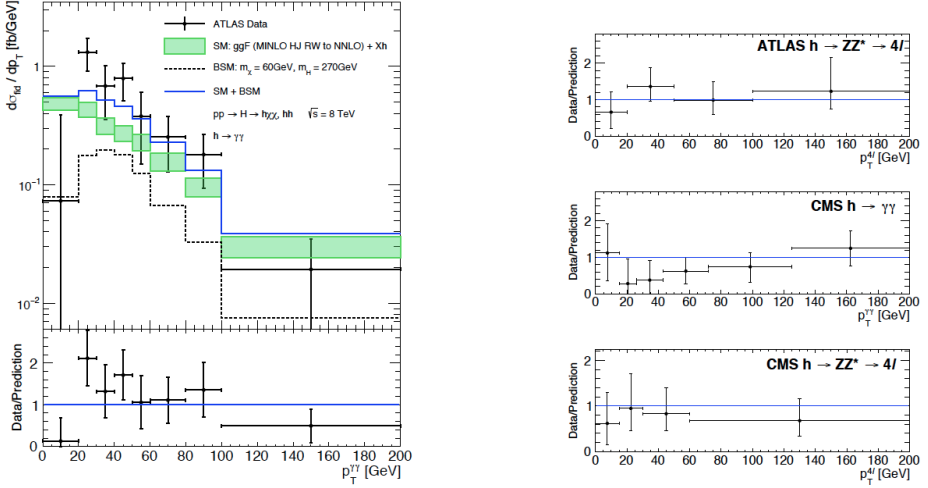


Figure 1.5 Fits to fiducial differential p_T distributions of the Higgs boson. The blue line on the left plot represent the SM+BSM model prediction in the $pp \rightarrow h \rightarrow \gamma\gamma$ channel [43]. The three ratio plots right illustrate the ratio of data and SM+BSM predictions, the top right plot is the ATLAS $pp \rightarrow h \rightarrow 4\ell$ channel [44], the middle one is the CMS $pp \rightarrow h \rightarrow \gamma\gamma$ channel [45] and the bottom one is the CMS $pp \rightarrow h \rightarrow 4\ell$ channel [46].

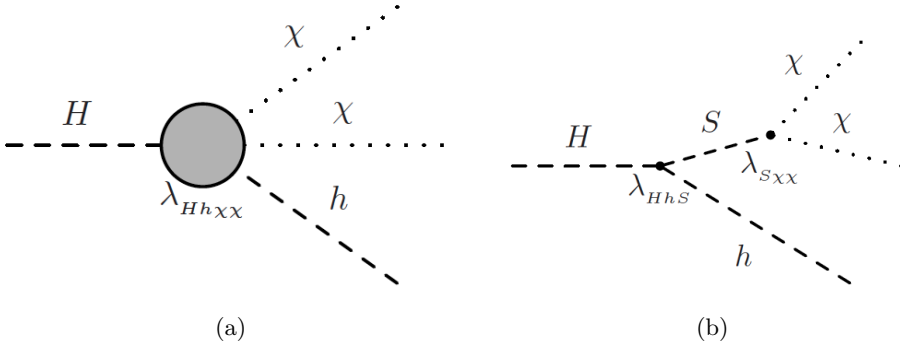


Figure 1.6 Feynman diagrams of the Madala hypothesis; on the left is the effective quartic $\lambda_{Hh\chi\chi}$ vertex and on the right the hypothesis is explained by a trilinear λ_{HhS} vertex [25].

interact with dark matter. The coupling strengths of S with SM particles closely resembled those of the SM Higgs boson, scaled down by the invisible

Table 1.2 Possible decay modes of scalars generated by the 2HDM+S framework [27].

Scalar N0.	Scalars	Decay modes
D.1	h	$b\bar{b}, \tau^+\tau^-, s\bar{s}, c\bar{c}, gg, \gamma\gamma, Z\gamma, W^+W^-, ZZ$
D.2	H	D.1, hh, SS, Sh
D.3	A	D.1, $t\bar{t}, Zh, ZH, ZS, W^\pm H^\mp$
D.4	$H^\pm$	$W^\pm h, W^\pm H^\mp, W^\pm S$
D.5	S	D.1, $\chi\chi$

decay $S \rightarrow \chi\chi$. The $S \rightarrow \chi\chi$ decay is not necessarily the only decay channel of S , and the mixing of S and h was assumed to very small to avoid violating experimental bounds. A study detailed in Reference [60] ruled out a 100% branching ratio of S into dark matter candidates.

The decay modes of particles resulting from the mixing of SM particles and the 2HDM+S scalars are summarised in Table 1.2. The mixing would allow a direct observation of S through the $S \rightarrow VV$ modes, where V denotes the electroweak gauge bosons, W, Z and γ . Direct production of S is anticipated to be small, as it is produced mainly through $pp \rightarrow SS, Sh$ modes. The background and signal modelling in $S \rightarrow VV$ searches would be similar to high-mass Higgs boson searches.

The Madala hypothesis was then used to explore multi-lepton signatures arising from the two new hypothetical scalar bosons, H and S . A study detailed in Reference [27] aimed to explain a number of anomalous features in published Run 1 experimental, which persisted into early Run 2 data. These features included the elevated production of opposite sign, OS, and same sign, SS, dileptons and three leptons events. In an effort to understand these excesses, the $pp \rightarrow H \rightarrow Sh, SS^*, SS'$ decay modes were considered. Here the $m_h = 125$ GeV SM Higgs boson primarily decays to b -quarks, while the Higgs-like scalar S predominantly decays to W and Z bosons for $m_s \approx 150$ GeV. These electroweak gauge boson decays then give rise to multi-lepton final states.

The predictions of the Madala hypothesis were compared with public ATLAS and CMS results to establish constraints on the fit parameter β_g^2 , where $\beta_g^2 = 0$ indicates the absence of a BSM signal. Assuming a heavy scalar mass of $m_H = 270$ GeV and varying the mass of the scalar S , the model's best fit was achieved at $m_s = 150 \pm 5$ GeV. The constrained values for β_g^2 were $\beta_g^2 = 1.22 \pm 0.38$ for OS di-lepton invariant mass spectra and $\beta_g^2 = 1.69 \pm 0.54$ for SS multi-leptons produced in association with b -tagged jets. These fit values closely align with $\beta_g^2 = 1.38 \pm 0.32$, derived from measurements of various Higgs boson signal strengths. Exploring a corner of phase-space with OS di-leptons accompanied by missing transverse energy, E_T^{miss} , and associated with hadronic jets could open a window to BSM physics.

Following the findings presented in Reference [27], a subsequent study detailed in Reference [29] aimed to understand in greater depth the origins of the multi-lepton anomalies observed in LHC data and their compatibility with new physics at the electroweak scale. The study assumed a simplified $H \rightarrow Sh$ model like in the study discussed in Reference [27] and utilised the model to analyse results from both Run 1 and Run 2 published by ATLAS and CMS Collaborations, focusing on multi-lepton final states. The study considered publications with following final states: SS $\ell\ell$ and $\ell\ell\ell + b$ -jets [61], OS $e\mu + b$ -jets [62], SS $e\mu, \mu\mu$ and $\ell\ell\ell + b$ -jets [63], OS $e\mu$ [64], $\ell\ell\ell + E_T^{\text{miss}}(WZ)$ [65], SS $\ell\ell$ and $\ell\ell\ell + b$ -jets [66], OS $e\mu + b$ -jets [67], and $\ell\ell\ell + E_T^{\text{miss}}(WZ)$ [68]. The presence of b -tagged jets resulting from top quark decay, along with the $h \rightarrow b\bar{b}$ decay, significantly increased the likelihood of reconstructing three b -tagged jets in the final state.

A statistical fit was performed for each of the considered publications with multi-lepton final states using a SM+BSM theoretical prediction. As per the Madala hypothesis, the masses of scalars H and S were set at $m_H = 270$ GeV and $m_S = 150$ GeV, respectively, leaving β_g^2 as a single degree of freedom of the BSM model requiring constraint. The fit parameter β_g^2 was constrained using a profile likelihood ratio. The best fit value for β_g^2 corresponds to a value that minimises the negative log-likelihood function, $-2\log \lambda(\beta_g^2)$. The local

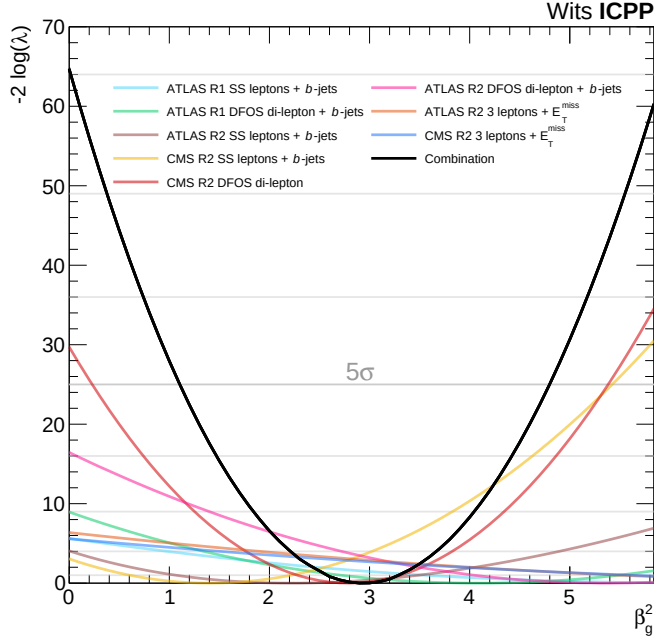


Figure 1.7 The profile likelihood ratios of individual and combined fit results. The significance is calculated where the profile likelihood crosses the y -axis ($\beta_g^2 = 0$, SM only hypothesis) [29].

Poisson significance, Z_0 , for each fit result was computed as:

$$Z_0 = \sqrt{-2 \log \lambda(0)}. \quad (1.31)$$

The combined profile likelihood ratio was computed by taking the product of the profile likelihood ratios obtained for individual measurements. The combination yielded $\beta_g^2 = 2.92 \pm 0.35$, corresponding to a combined Poisson significance of 8.04σ . Figure 1.7 shows the individual likelihood ratios for each measurement, along with the combined likelihood of the measurements. The significance was calculated by taking the square root of the y -intercept of the combined profile likelihood ratio.

The study outlined in Reference [29] demonstrated significant discrepancies between experimental results published by both the ATLAS and CMS Col-

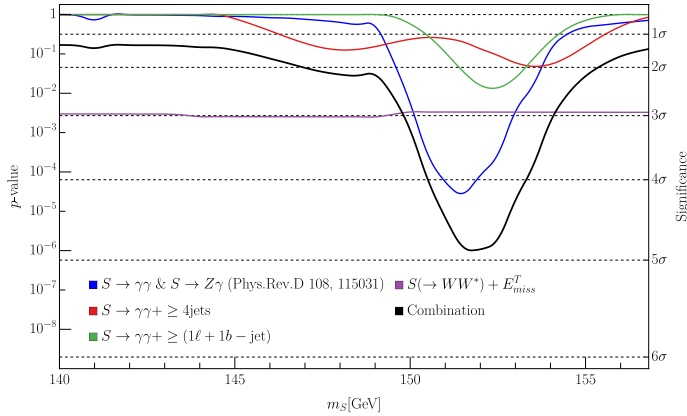


Figure 1.8 The observed global p -value as a function of m_S . The solid black line represents a combination of the individual high-mass channels [70].

laborations and the SM predictions. These discrepancies are particularly pronounced in a corner of phase space dominated by various SM processes, ruling out the possibility that the multi-lepton anomalies can be explained by mis-modelling of a particular process, such as $t\bar{t}$ production alone. Table 1.1 shows the multi-lepton final states with statistical significance reaching over 3σ , providing compelling evidence of the multi-lepton anomalies.

To explain the multi-lepton anomalies within the 2HDM+S framework, it was postulated that a scalar H should exist with a mass constrained to approximately 270 GeV. Additionally, a scalar S with a mass around 150 GeV was predicted in Reference [27].

Recent reviews of ATLAS and CMS published results have hinted at the potential existence of new scalar bosons with masses around 95 GeV and 150 GeV [69, 70]. The studies in reported excess events at a mass around 150 GeV, consistent with a scalar boson predicted in Reference [27]. The constrained mass S falls into the side-band region in Higgs searches. As a result, published ATLAS and CMS results were used to search for the existence of S in the side-bands. These searches for the signal of S utilised published ATLAS and CMS results in analysis with $\gamma\gamma$, $Z\gamma$, and $b\bar{b}$ final states [71–77].

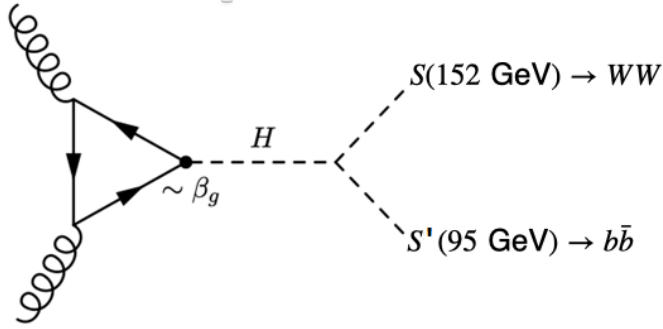


Figure 1.9 The scalar S pair production via the decays of H with $m_H = 270 \text{ GeV}$.

The study reported a global significance of 4.7σ , Figure 1.8, affirming that the observed discrepancies in the published experimental data could potentially be explained by the existence of this scalar S . The results are consistent with the $H \rightarrow SS^*$ hypothesis with $m_H = 270 \text{ GeV}$. Figure 1.9 shows the production of SS^* via the decay of H .

The associate production of S with bottom quarks can be explored by searching for asymmetric events with $\tau^+\tau^-b\bar{b}$ final states produced via the $H \rightarrow SS'$ decay. The $S' \rightarrow b\bar{b}$ could potentially explain the production of a new scalar at 95 GeV , which has been hinted by results published by LEP [78] and the mild excess seen by CMS in Run 2 results [79]. This could also explain the presence of b -jets in Table 1.1. References [35, 69, 70, 80] provide an in-depth analyses of studies highlighting anomalies involving multiple leptons at the electroweak scale.

1.4 Layout of the Thesis

The first part of the thesis is dedicated to performance and calibration studies using the Tile calorimeter, a crucial component of the ATLAS detector. Accurate measurement and calibration of the TileCal are essential for reliable luminosity measurements and overall detector performance.

The second part of the thesis, motivated by the multi-lepton anomalies, aims to explore these anomalies by searching for new high-mass resonances predicted by theoretical models, such as the Madala hypothesis.

This thesis is organised as follows:

Chapter 1, this chapter, introduces the SM of particle physics and discusses the Madala hypothesis, which interprets discrepancies between SM predictions and experimental observations at the LHC as hints of BSM physics. The model further predicted and accounted for the existence of multi-lepton anomalies in the published LHC data.

Chapter 2 provides an overview of particle accelerators, focusing on the LHC accelerator complex. Additionally, it describes the ATLAS detector and its sub-detectors.

Chapter 3 presents a detailed overview of the Tile Calorimeter, discussing its cell structure layout, read-out system, and calibration systems.

Chapter 4 discusses the role of the Tile calorimeter in luminosity measurements and presents the long-term stability uncertainty in luminosity measurements provided by the LUCID detector with respect to the Tile calorimeter.

Chapter 5 presents a study that investigates the response of gap and crack scintillators of the Tile Calorimeter to isolated muons originating from proton-proton collisions. The response of the scintillators is used to study the inter-calibration of the scintillators across all the Tile calorimeter modules.

Chapter 6 provides an overview of the types of collisions occurring at the interaction point of particle accelerators. It also introduces the coordinate system used and describes how physics objects are reconstructed by the detector.

Chapter 7 focuses on the search for a heavy scalar resonance decaying to a Z boson and a photon in the mass range between 220 and 3400 GeV. This search is motivated by the anomalies in multi-lepton final states. This search utilises proton-proton collision events at $\sqrt{s} = 13$ TeV recorded by the ATLAS detector from 2015 until 2018, with an integrated luminosity of 140.0 fb^{-1} . The study aims to identify any significant deviations from the SM predictions.

Chapter 2

Experimental Apparatus

Particle physics models are tested by means of particle physics experiments, which can either be accelerator based or non-accelerator based. The non-accelerator based experiments are mostly underground facilities that seek to study neutrino physics through collisions of neutrinos with the nuclei of the detecting medium. One of the most notable particle physics discoveries achieved using a non-accelerator experiment was the observation of neutrino oscillations by the Japan based Super-Kamiokande Collaboration in 1998 [81]. This groundbreaking observation was later confirmed by the Sudbury Neutrino Observatory in 2001 [82].

The existence of dark matter has been inferred from gravitational studies, which have observed various luminous objects moving faster than expected if they were solely influenced by the gravitational attraction of visible objects [83]. Besides its gravitational interactions, there has not been any direct observation of dark matter, to date. Various candidate particles have been proposed for dark matter, including primordial black holes [84], axions [85], sterile neutrinos [86], and weakly interacting massive particles [1]. There is a multitude of non-accelerator based experiments aimed to make a direct observation of dark matter. Deep underground detectors are generally designed to detect dark matter by observing its recoil with the nuclei of the detecting medium, just like in neutrino detectors. The focus of this chapter is on accelerator based experiments.

2.1 Particle Colliders: An Overview

Particle colliders utilise electromagnetic fields to accelerate charged particle beams to high energies. The accelerated particle beam can be directed to collide either against a fixed target or another accelerated beam. In the fixed target setup, the accelerated beam is aimed at a stationary solid or liquid target. Collisions involving high energy particle beams produce new particles, with the new particles masses dependent on the collision's centre-of-mass energy [87]. In addition to centre-of-mass energy, luminosity, introduced later, is another important parameter in collider experiments.

For the discovery of new heavy particles or for studying the subatomic structure of particles, working within a frame where the centre-of-mass energy is maximised is preferable. This frame, known as the centre-of-mass frame, is defined as a point that moves with the system's total momentum, where the centre-of-mass of the system is at rest, and the total momentum is zero. Considering a collision between two particles, p_A and p_B , with momentum $\vec{p}_A$ and $\vec{p}_B$, respectively, and energies E_A and E_B , respectively, with an angle θ between p_A and p_B . The centre-of-mass energy of such system is given by:

$$E_{\text{cm}}^2 = m_A + m_B + 2E_A E_B (1 - \beta_A \beta_B \cos \theta), \quad (2.1)$$

where $\beta = \frac{v}{c}$, and m_A and m_B represent the masses of p_A and p_B , respectively. For $v \approx c$ and $E \gg m$, Equation 2.1 simplifies to:

$$E_{\text{cm}} = \sqrt{2E_A E_B} = \sqrt{s}, \quad (2.2)$$

where s represents the total energy in the centre-of-mass system. Equation 2.2 show that the centre-of-mass energy of a fixed target experiment increases proportionally to the square root of the energy of the incoming accelerated beam.

In beam-beam colliding experiments, where two identical beams with particles, $p_A = p_B = p$, colliding head-on, $\theta = \pi$, the centre-of-mass energy, $\sqrt{s}$, for such

system is given by:

$$E_{\text{cm}} = 2E = \sqrt{s}. \quad (2.3)$$

The centre-of-mass of two identical beams colliding increases linearly, effectively doubling the energy of the incoming accelerated beam. Colliding two beams of charged particles enables researchers to achieve exceptionally high centre-of-mass energies, thereby facilitating the exploration of new physics phenomena.

The probability of a specific physics interaction occurring in a collision is called the production cross-section, σ , and it is typically quite small [88]. The cross-section serves as a fundamental measure of the scattering process. To evaluate an accelerator's performance in detecting a particular interaction, it is necessary to introduce the concept of luminosity.

In the detector, the observed rate of a specific process is directly proportional to its production cross-section. Mathematically, this relationship is expressed as:

$$R = \mathcal{L}\sigma, \quad (2.4)$$

where R is the number of interactions per unit time, σ is the event's cross-section which has units of area, barns, and $\mathcal{L}$ represents the instantaneous luminosity. Considering two accelerated bunches of particles with Gaussian horizontal width, Σ_x , and vertical width Σ_y , containing n_1 and n_2 particles, respectively, crossing each other with frequency f , $\mathcal{L}$ can be expressed as:

$$\mathcal{L} = \frac{fn_1n_2}{4\pi\Sigma_x\Sigma_y}. \quad (2.5)$$

From Equation 2.5, it is evident that high instantaneous luminosities can be attained by increasing the revolution frequency of the beams, increasing the number of particles in a bunch, and/or by tightly focusing the beams at the collision point.

2.2 The Large Hadron Collider

The LHC is a large circular particle collider managed by CERN, the European Organization for Nuclear Research, located on the boarder between Switzerland and France. As of the time of writing, it stands as the largest particle physics experiment in the world. The LHC operates in modes capable of accelerating either proton or heavy-ion beams, such as lead. In proton-proton collision mode, it accelerates two beams of protons in opposite directions to speeds approaching the speed of light, resulting in an energy of 6.5 TeV per beam and a collision energy of $\sqrt{s} = 13$ TeV. The lead-lead collision mode has a collision energy of $\sqrt{s} = 5.02$ TeV. This chapter focuses on the proton operational mode of the LHC during the Run 2, spanning from 2015 to 2018. In Run 3, the proton beams are accelerated to a collision energy of $\sqrt{s} = 13.6$ TeV.

The process of accelerating the proton beams begins with the ionisation of hydrogen atoms into protons using strong electric fields. These protons are then arranged into bunches and injected into Linac 2, a linear accelerator, where they are accelerated to 50 MeV. Afterward, the accelerated bunches enter a series of smaller circular accelerators before ultimately being injected into the LHC. The Proton Synchrotron Booster receives the accelerated bunches from Linac 2 and further boosts their energy until they reach 1.4 GeV. Following this, the bunches are fed to the Proton Synchrotron, where they are accelerated to 25 GeV before being injected into the Super Proton Synchrotron. Here, the beam is further accelerated to 450 GeV before it is finally fed into the LHC. Figure 2.1 provides an illustration of the CERN accelerator complex.

Accelerated particles emit electromagnetic radiation known as synchrotron radiation [87]. As a result, power losses due to synchrotron radiation impose a limit on the maximum energy to which particle beams can be accelerated. To find the amount of power lost through synchrotron radiation, consider a particle with charge q moving along a circular path of radius ρ . The instantaneous

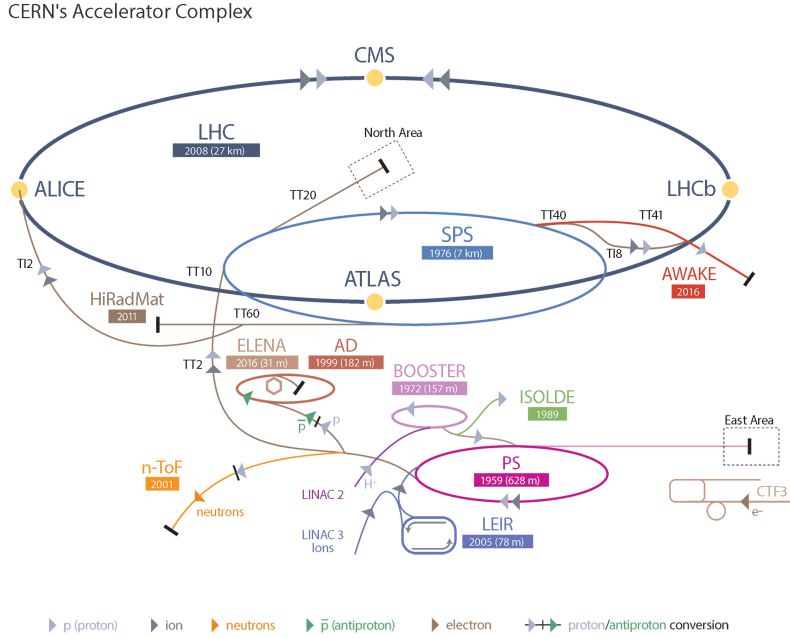


Figure 2.1 CERN accelerator complex [89].

radiated power [90], P , is given by:

$$P = \frac{2q^2c}{3\rho^2} \left(\frac{\beta}{\sqrt{1-\beta^2}} \right)^4. \quad (2.6)$$

At relativistic speeds, $\beta \approx 1$ and $\gamma = \frac{E}{mc^2} = \frac{1}{\sqrt{1-\beta^2}}$, the instantaneous radiated power becomes:

$$P = \frac{2q^2c}{3\rho^2} \gamma^4 = \frac{2q^2c}{3\rho^2} \left(\frac{E}{mc^2} \right)^4. \quad (2.7)$$

The LHC, with a circumference of approximately 27 km, is built within the same tunnel that previously housed the Large Electron-Positron, LEP, Collider [91]. One of reasons to repurpose the LEP Collider into the LHC was driven, in part, by the significantly lower energy loss through synchrotron radiation in accelerated protons compared to accelerated electrons and positrons.

To illustrate this point, consider proton and electron beams accelerated to the same energy in the same accelerator with equal radii. The ratio of the power lost due to synchrotron radiation is:

$$\frac{P_{\text{loss}}(e^-)}{P_{\text{loss}}(p)} = \left(\frac{m_p}{m_{e^-}} \right)^4 \approx 1.129 \times 10^{13}. \quad (2.8)$$

The LHC receives a beam of protons consisting of 2808 bunches, each spaced 7.5 m apart. These proton bunches are then accelerated to 6.5 TeV and focused to $16 \times 16 \mu\text{m}$ using quadrupole magnets at the interaction points. The LHC features four main interaction points where experiments are conducted: Point 1 - ATLAS, Point 2 - ALICE, Point 5 - CMS, and Point 8 - LHCb.

The ATLAS and CMS detectors serve as general-purpose detectors, with broad experimental goals encompassing searches for BSM physics, precise measurements of SM parameters, and notably, the Higgs boson observation which was achieved in 2012. The LHCb detector specialises in studying the charge-parity violation and the decays of B mesons. ALICE's mission focuses on studying the quark-gluon plasma, a state of matter present in the early stages of the Big Bang. This chapter will primarily concentrate on the ATLAS detector, but before delving into detail about the detector, it is worthwhile to discuss the coordinate system used.

2.2.1 Detector Coordinate System

The interaction point, IP, where two proton beams collide, serves as the origin of the detector's coordinate system. The IP is located at the centre of the detector, and the cylindrical symmetry of the detector is utilised to establish the orientation of the coordinate axes. The direction of the incident proton beam aligns with the z -axis, and the azimuthal angle around the beam axis is denoted by ϕ , where $\tan \phi = \frac{p_x}{p_y}$.

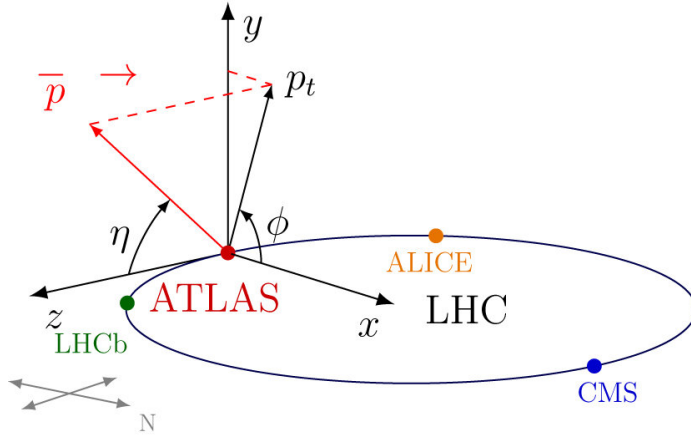


Figure 2.2 ATLAS detector coordinate system [92].

The rapidity, y , of a particle is given by the equation:

$$y = \frac{1}{2} \ln \frac{E + p_z}{E - p_z}, \quad (2.9)$$

where E and p_z represent the particle's energy and momentum along the z -axis, respectively. It serves as a measure of the velocity difference between the lab frame and a frame in which the particle moves only perpendicular to the beam. For particles traveling in the x - y plane with negligible p_z , the rapidity is zero, while for particles with $p_z c \approx E$, rapidity tends to infinity.

Accurately measuring p_z is quite challenging. As an alternative, pseudorapidity, denoted by η , is computed from more readily measurable quantities. It is defined as:

$$y \approx -\ln \left[\tan \left(\frac{\theta}{2} \right) \right] \equiv \eta, \quad (2.10)$$

where θ denotes the polar angle between the particle's path and the z -axis.

In relativistic cases where the p_T of the accelerated particle significantly exceeds its mass, pseudorapidity is approximately equal to rapidity, $\eta \approx y$. The coordinates (η, ϕ) specify the angle at which a particle was emitted from the IP. The quantity ΔR represents the angular separation or distance between two physics objects within the detector, and it is calculated as $\Delta R = \sqrt{\Delta\eta^2 + \Delta\phi^2}$.

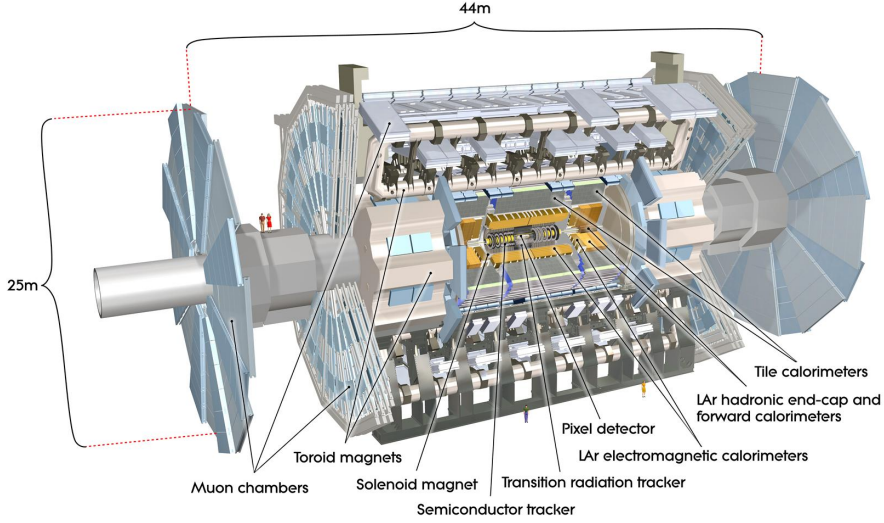


Figure 2.3 A layout of the ATLAS detector [93].

2.3 The ATLAS Detector

The ATLAS detector, with a weight of approximately 7000 tonnes, stretches 44 metres in length and has barrel diameter of 25 metres. It is split into the barrel section, where detector components surround the beam axis cylindrically, and the end-caps, positioned perpendicular to the beam axis. It features a cylindrical and forward-backward symmetry around the IP. A schematic layout of ATLAS detector and its sub-detectors is illustrated in Figure 2.3.

It is a general purpose detector that comprises of tracking detectors, which use strong magnetic fields to measure the momentum of charged particles, and calorimeters, responsible for measuring particle energies. The tracking detectors consist of the inner detector and the muon spectrometer. The calorimeters are divided into two types, the electromagnetic calorimeter, dedicated to measuring energies of photons and electrons, and the hadronic calorimeter, which measures the energies of particles that are not absorbed by the

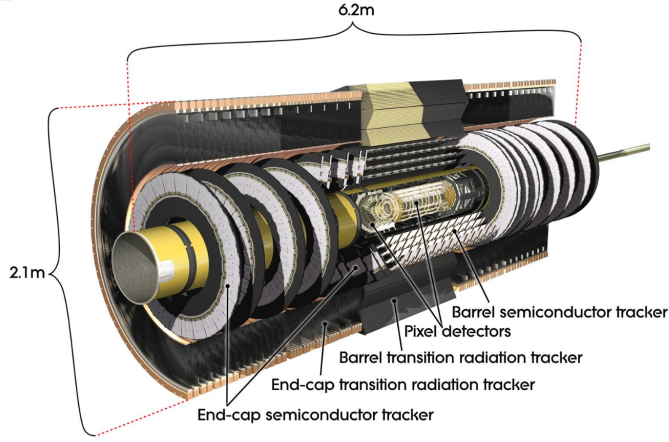


Figure 2.4 A cross-sectional view of the inner detector showing the arrangement of its sub-detectors in the barrel and end-cap regions [97].

electromagnetic calorimeter, primarily hadrons. The detector has an overall pseudorapidity coverage range of $|\eta| < 4.9$ [94].

2.3.1 The Inner Detector

The inner detector, ID, is the sub-detector closest to the IP and it operates within a two Tesla magnetic field parallel to the beam axis, provided by a superconducting solenoid. Electrically charged particles traversing the ID have their trajectories bent by the magnetic field, enabling the calculation of p_T the particles from the radius of curvature of their tracks. For precise measurements of particle tracks, the ID employs the following three complementary sub-detectors, each positioned successively further from the IP: the pixel detector, Semiconductor Tracker, and the Transition Radiation Tracker.

Figure 2.4 illustrates the layout of the ID sub-detectors in both the barrel section and the end-caps of the ATLAS detector. The pixel detector and the Semiconductor Tracker system provide exceptional momentum and vertex reconstruction capabilities [95]. The ID covers a range of $|\eta| < 2.5$ and reconstructs charged particle tracks with $p_T > 500$ MeV [96].

The high-granularity silicon pixel detector is made of three silicon pixel layers with 1456 identical modules in the barrel section [95]. The first pixel layer closest to the IP is known as the *b*-layer because it plays a very vital role in the physics of the bottom quark. ATLAS uses an Insertable B-Layer, IBL, which maintains robust tracking and vertex reconstruction capabilities even in high luminosities, hardware life-time and radiation constraints [98]. The end-cap pixel detector system consists of three discs, with a total of 288 modules positioned perpendicular to the beam axis on both sides of the IP.

The Semiconductor Tracker encompasses the pixel detector in both the barrel region and the end-caps. It is made out of two-sided layers of silicon strips with 4088 modules. It provides position measurements of charged particles with an accuracy of $17\ \mu\text{m}$ per layer of silicon strip.

The Transition Radiation Tracker system comprises of drift tubes filled with a gas mixture and equipped with a gold-plated tungsten wire at the centre, enclosed within straw tubes with an outer diameter of 4 mm. The gas mixture aids the absorption of X-rays, increases the drift velocity of electrons, and facilitates photon-quenching [99]. The barrel region contains 50 000 straws, while the end-caps house a total of 250 000 straws. Transition radiation is leveraged to differentiate electron candidates from pion candidates.

The momentum resolution of the ID is described by the equation:

$$\frac{\sigma_{p_T}}{p_T} = ap_T \oplus b = 0.36\ \text{TeV}^{-1} p_T \oplus \frac{0.013}{\sqrt{\sin \theta}}, \quad (2.11)$$

where a and b are parameters that characterise the resolution, determined through calibration and performance studies based on experimental data and simulations. The symbol $\oplus$ indicates a sum in quadrature. The ID tracks are capable of being extrapolated to both primary vertices, where particles originate, and secondary vertices, where particles decay after traversing a considerable distance.

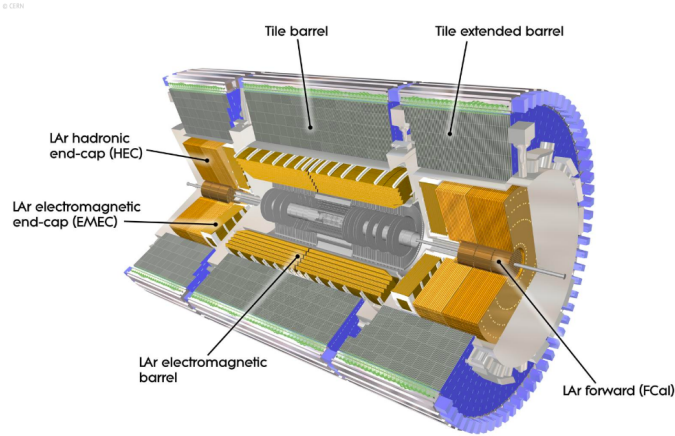


Figure 2.5 A schematic layout of the calorimeters in the ATLAS detector [100].

2.3.2 The Calorimeter Systems

The calorimeter system of the ATLAS detector measures the energy of incoming particles and determines their direction of origin. It comprises two primary types of calorimeters: electromagnetic and hadronic calorimeters. The electromagnetic calorimeter is designed to measure the energy of photons and electrons, while the hadronic calorimeters are responsible for measuring the energy of hadrons, which are essential for reconstructing jets. Additionally, the energy deposit information in the calorimeter systems is crucial for the reconstruction of E_T^{miss} , a very important signature of various new physics phenomena.

The Electromagnetic Calorimeter

The electromagnetic calorimeter, ECal, is positioned immediately after the ID. However, before particles can reach the ECal, they must traverse through a cryostat system housing the two Tesla solenoid magnet employed by the ID [101]. A presampler placed in front of the ECal serves an important role of correcting for energy losses of particles as they pass through the cryostat

system [102]. By compensating for these losses, the presampler ensures that energy measurements in the ECal are accurate and reliable.

The ECal uses liquid Argon, LAr, as its active medium. This choice of material offers several advantages, such as excellent energy resolution and radiation hardness. In the barrel region, the ECal covers the range $|\eta| < 1.4$, while the two end-cap calorimeters cover the range $1.4 < |\eta| < 3.2$. The barrel and the two end-cap sections of the ECAL are called EMA and EMEC, respectively.

The Hadronic Calorimeter

The barrel hadronic calorimeter of the ATLAS detector is called the Tile Calorimeter, TileCal. It is a sampling calorimeter that uses iron as an absorbing medium and plastic scintillating tiles as the active medium [103, 104]. It is split into the central barrel and extended barrels, positioned on either side of the central barrel. The central barrel covers pseudorapidity range $|\eta| < 1$, while the extended barrels cover the $0.8 < |\eta| < 1.7$ range. The TileCal will be discussed in detail in Chapter 3.

The hadronic end-caps, HEC, utilise a LAr sampling calorimeter with copper as the absorbing material. LAr is chosen for its intrinsic radiation-hard properties, making it suitable for the demanding environment of the HEC. The HEC is made up of two co-axial wheels. The inner wheel is made from 25 mm thick copper and the outer wheel from 50 mm thick copper. They cover a pseudorapidity range of $1.5 < |\eta| < 3.2$ [101], overlapping with the region covered by the TileCal.

The Forward Calorimeter

The forward calorimeter, FCal, is positioned approximately 4.7 m away from the IP on each side, enabling the detection of physics objects in the high η region. The FCal serves multiple purposes; forward jets tagging, studying E_T^{miss} and providing background shielding for the muon spectrometer [105]. Covering the region of $3.1 < |\eta| < 4.9$, the FCal is particularly instrumental in

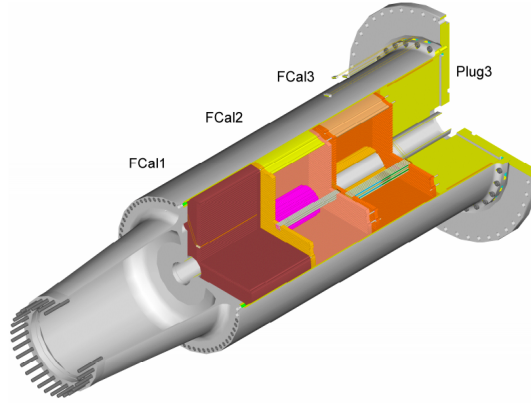


Figure 2.6 Layout of the forward calorimeter [106].

reconstructing high- p_T neutrinos or other physics objects that weakly interact with the detector material. These objects may escape without detection, but their presence can be inferred from events exhibiting a significant p_T imbalance in the forward region, making the FCal vital for reconstructing potential new physics signatures.

Due to the high particle flux in the forward region, largely resulting from elastic collisions, the materials used in the FCal must be radiation hard. The FCal is made from three cylindrical modules behind each other: FCal-1, FCal-2 and FCal-3 as seen in Figure 2.6. The FCal-1, functioning as an electromagnetic calorimeter, is the closest module to the IP and it uses copper as the absorbing material and LAr as the active material. The FCal-2 and FCal-3 are dedicated to measuring energy of hadrons and they use LAr as the active material and tungsten as the absorber material. Behind the FCal-3, there is a copper plug serving as a radiation shield for the muon spectrometer

The energy resolution of the calorimeter systems is expressed as:

$$\frac{\sigma_E}{E} = \frac{a}{\sqrt{E}} \oplus \frac{b}{E} \oplus c, \quad (2.12)$$

Table 2.1 Energy resolution of the ATLAS calorimeter sub-detectors.

Sub-detector	Resolution
ECal	$\frac{\sigma_{E_{\text{ECal}}}}{E} = \frac{10\%}{\sqrt{E}} \oplus 0.7\%$ (both barrel and end-caps)
TileCal	$\frac{\sigma_{E_{\text{TileCal}}}}{E} = \frac{80\%}{\sqrt{E}} \oplus 15\%$
HEC	$\frac{\sigma_{E_{\text{HEC}}}}{E} = \frac{50\%}{\sqrt{E}} \oplus 3\%$
FCal	$\frac{\sigma_{E_{\text{FCal}}}}{E} = \frac{100\%}{\sqrt{E}} \oplus 10\%$

where a represents the stochastic term, E denotes the deposited energy in GeV, b corrects for non-uniformity in the calorimeter response, and c is a constant noise term [107, 108]. Table 2.1 shows resolution of calorimeter subsystems of the ATLAS detector. The non-uniformity term b is not explicitly included in the table but is part of the general resolution formula provided.

2.3.3 The Muon Spectrometer

The Muon Spectrometer, MS, is designed to detect muons, which, being heavier particles compared to electrons, interact minimally through electromagnetic interactions and do not participate in the strong interaction. Consequently, they typically traverse the ID and calorimeters with minimal energy loss. Positioned outside the ID and calorimeter systems, the MS incorporates three massive superconducting toroidal magnets, each made up of eight coils. It is composed of 4000 muon chambers which work together with the magnets to determine the charge and direction of passing muons [109].

Three concentric cylinders are used for the barrel chambers, covering the $|\eta| < 1$ region, with the exception of $|\eta| = 0$, where there is gap to allow passage of detector electronics cables. The end-cap region has four disk chambers covering the $1 < |\eta| < 2.7$ region. The Monitored Drift Tubes, MDTs, are used in the majority of the chambers, except in the end-cap region, $2 < |\eta| < 2.7$, where high particle fluxes necessitate the use of radiation-hard Cathode Strip Chambers.

The momentum resolution of the muon spectrometer is expressed as:

$$\frac{\sigma(p)}{p} = \frac{P_0}{p_T} \oplus P_1 \oplus P_2 \cdot p_T, \quad (2.13)$$

where the parameter P_0 represents a coefficient related to the energy loss in the material before the muon spectrometer, P_1 denotes multiple scattering and P_2 is the intrinsic resolution term. These parameters are determined through a global fit method described in detail in Reference [110].

2.3.4 Luminosity Detectors

Luminosity measurements are crucial for understanding the performance of the LHC and precise luminosity measurements are crucial in precisely determining the $t\bar{t}$, W , and Z boson production cross-sections. The primary bunch-by-bunch luminosity measurements during LHC Run 2 were provided by the Luminosity Cherenkov Integrating Detector, LUCID [111]. It was complemented by measurements from the ATLAS beam conditions monitor, BCM, diamond detectors, and offline measurements of the track multiplicity of reconstructed charged particles in randomly selected bunch-crossings.

The LUCID detector is made from two modules placed approximately 17 meters from the IP, in each forward arm of the ATLAS detector. Each module is made of radiation hard 4 quartz fibre bundles and 16 photomultiplier tubes, PMTs [112]. It uses the PMTs to measure the instantaneous luminosity of inelastic events by collecting the Cherenkov radiation produced when charged particles pass through the modules [113]. Further details regarding luminosity detectors and the algorithms utilised by the ATLAS detector are discussed in Chapter 4.

2.3.5 The Trigger System

There are limits on data transmission, processing and storage capabilities, so not every event per bunch crossing is saved. To decide which events to store,

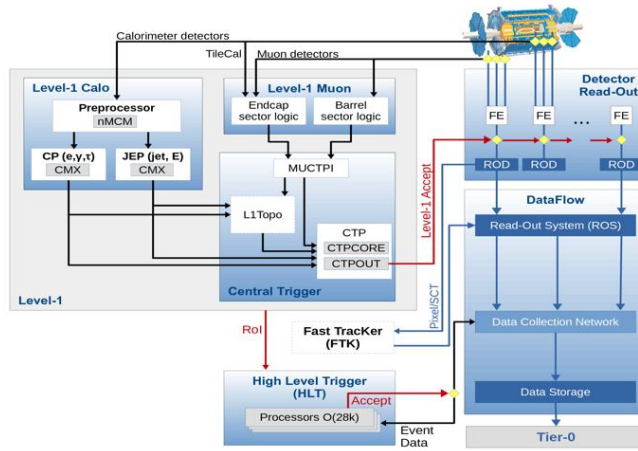


Figure 2.7 A schematic diagram of the trigger system used in ATLAS during Run 2 data-taking [115].

ATLAS employs a trigger system which rapidly selects interesting physics events to be recorded for offline analyses. The trigger system during Run 2 consisted of a hardware-based Level-1, L1, trigger and a software-based high level trigger, HLT [114].

The L1 trigger utilises signals from detector electronics to identify the Regions-of-interest, RoIs, within the detector, using information from the calorimeter cells and the muon system. Its primary function is to decrease the event rate from the LHC’s approximate rate of 40 MHz to an average recording rate of 100 kHz, making the decision to keep or discard an event within $2.5 \mu\text{s}$. The HLT processes information from the RoIs generated by L1 trigger. It uses a series of algorithms to further decrease the event rate from the L1 trigger to approximately 1 kHz on average, accomplishing this within a processing time of around 200 ms [116].

Chapter 3

The ATLAS Tile Calorimeter

The ATLAS TileCal, briefly introduced in Section 2.3.2, is a hadronic sampling calorimeter composed of tiles of plastic scintillators sandwiched between low-carbon steel plates acting as absorbers. It covers the barrel region of the ATLAS detector up to $|\eta| < 1.7$ and is segmented into three parts along the beam axis. The TileCal consists of one fixed long barrel, 5.8 metres long, covering the central area of the collision point $|\eta| < 1.0$, and two extended barrels, each measuring 2.6 metres in length, covering the forward and backward regions, $0.8 < |\eta| < 1.7$. Symmetrically positioned with respect to the IP, the positive η side is referred to as the A-side, while the negative η side is called the C-side.

3.1 Layout of TileCal Cells

The long barrel of the TileCal is segmented radially into three layers: A, BC, and D, approximately 1.4, 3.9, and 1.8 nuclear interaction lengths thick, respectively, at $\eta = 0$. The extended barrels are segmented into layers A, B and D, with thicknesses of approximately 1.6, 2.6 and 3.3 nuclear interaction lengths, respectively [117]. The full azimuthal coverage around the beam axis is achieved using 64 wedge-shaped modules, depicted schematically in Figure 3.1, each covering $\Delta\phi = 0.1$ radians. The modules are oriented radially and perpendicular to the beam axis. The TileCal has total of 256 of such

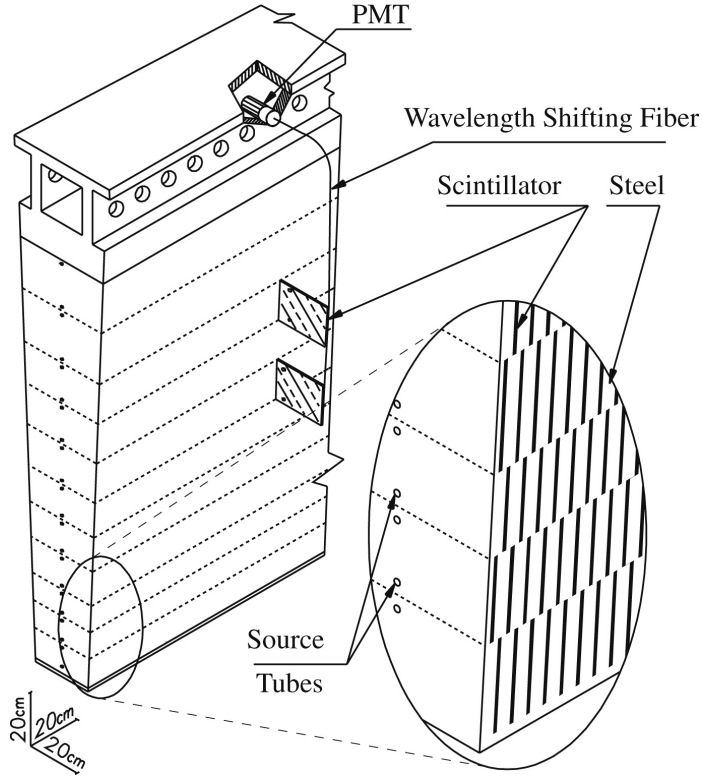


Figure 3.1 A schematic representation of a wedge-shaped TileCal module [117], showing the mechanical assembly and optical readouts. The holes accommodate radioactive source tubes traversing the modules perpendicular to the iron and plastic scintillators for calibration purposes.

modules, evenly distributed between the A-side and C-side, with 128 module on each. Each module has a segmentation of $\Delta\eta = 0.1$ in the first and second radial layers, and $\Delta\eta = 0.2$ in the third radial layer.

The ϕ , η and radial segmentation define the three dimensional stacks that define TileCal cells within a module. The cells located within radial layers A and B(C) are named according to their respective η locations. For example, cells BC4 are located in the BC layer at $\eta = 0.4$, while cells A13 are found in the A layer at $\eta = 1.3$. The D layer has $\Delta\eta = 0.2$ segmentation, with cells indexed sequentially from the centre of the cell. For instance, cells D0 and

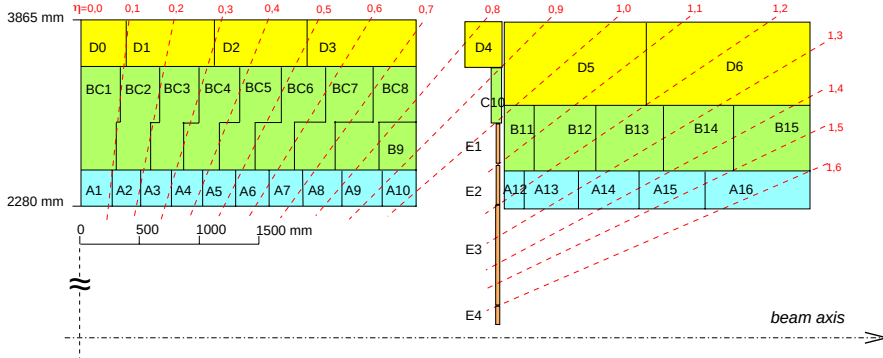


Figure 3.2 Schematic layout of TileCal cells in the ϕ - η plane on the A-side of the detector [118]. Solid lines show cell boundaries formed by grouping optical fibres from the tiles for readout by PMTs. The gap and crack scintillators are shown in orange.

D5 will occupy the intervals $-0.1 < \eta < 0.1$ and $0.9 < \eta < 1.1$, respectively. Figure 3.2 shows the schematic layout of η and radial segmentation of the TileCal cell assembly on the A-side.

The transition region, $0.8 < |\eta| < 1.6$, is covered by the Intermediate Tile Calorimeter, ITC. The ITC is a plug detector designed to fill the gap between the long and extended barrels with the aim of correcting for energy losses experienced by particles when pass the passive material that fills this region. The region $0.8 < |\eta| < 1.0$ of the ITC, known as the plug, consist of 311 mm thick steel-scintillator stacks, D4, in the region $0.8 < |\eta| < 0.9$ with similar design to the regular TileCal cells. The region $0.9 < |\eta| < 1.0$ features 96 mm stacks, C10, in the z -direction. The reduced thickness of the C10 and D4 cells or their special geometry is to accommodate services and read-out electronics for other ATLAS detector systems [119, 120].

The gap and crack region of the ITC, located in the region $1.0 < |\eta| < 1.6$, consists of cells made of single scintillating tiles without any steel absorbers, due to space limitations. These cells, often just simply referred to as scintillators, are trapezoidal in shape. Within the region $1.0 < |\eta| < 1.2$, the cells are known as gap scintillators, made of E1 and E2 cells, while those within the

region $1.2 < |\eta| < 1.6$ are called crack scintillators, consisting of E3 and E4 cells.

The space constraints in the transition region limit the maximum thickness of the gap and crack scintillators. Gap scintillators are limited to a thickness of 12 mm, while crack scintillators have a thickness of 4 to 6 mm. The gap scintillators primarily provide hadronic shower sampling, while the crack scintillators play an important role in measuring energy escaping the ECal. The Technical Design Report [120] provides a detailed description of the ATLAS TileCal and its construction, optical instrumentation and installation into the ATLAS detector are detailed in References [119, 121].

3.2 TileCal Read-out System

The TileCal read-out system plays a crucial role in the detection high-energy particles within the calorimeter. When high-energy particles interact with the steel absorber material in TileCal cells, they generate showers of lower energy particles, which, in turn, produce ultraviolet light upon absorption by the plastic scintillators. The emitted light from the scintillation process is captured by wavelength shifting, WLS, fibres coupled to the radial edges of plastic scintillators. These fibres are bundled together to form the optical read-out of a cell, and they operate independently of each other.

The scintillation light is transmitted by the WLS fibres to two PMTs, one PMT per radial side of a cell, to maximise the light yield and provide redundancy in the light readout. The PMTs play the important role of converting the scintillation light into electrical signals. In special cells made from a single scintillator without steel absorbers, such as the gap and crack scintillators, the WLS fibres are only coupled to a single PMT.

The PMTs and front-end electronics for each module are housed within drawers located at the outer radius of the calorimeter. The read-out architecture divides the TileCal detector into four partitions. The two extended barrels form part of the EBA and EBC partitions, while the long barrel is divided

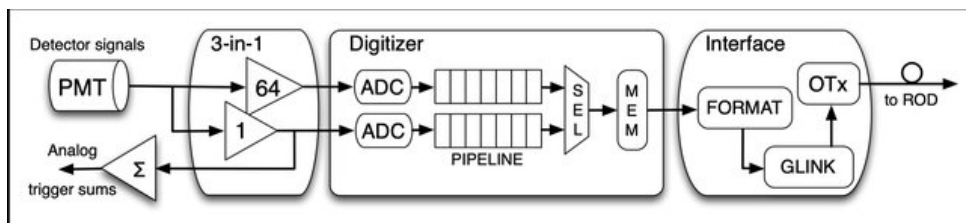


Figure 3.3 Schematic of TileCal front-end read-out electronics in Run 2 [123].

into two partitions, namely the LBA and LBC, depending on whether the partition falls on the A or C side of the detector.

Figure 3.3 illustrates the TileCal read-out architecture used during Run 2. The read-out electronics have been refurbished for Run 3, in order for them to cope with the increase in average luminosity, 5 to 7 times larger than the nominal Run 2 values [122]. In this chapter the Run 2 read-out architecture will be briefly discussed.

During the digitisation of the scintillation light into electrical signals, the PMTs send a signal to a specialised component known as the 3-in-1 card, named for its multifunctionality. The 3-in-1 card is made of a signal shaper, charge injection system, CIS, and gain splitting branches. The PMTs, along with the 3-in-1 card, light mixer, and high voltage divider, are grouped together in a device called a PMT-block [124]. The CIS is used to calibrate the front-end electronics and it will be discussed in more detail in Section 3.3.

The analogue PMT signals are shaped and amplified in two gain settings, with a nominal gain ratio of 64:1 [122]. An automatic gain selection mechanism assigns a gain setting to the PMT signals. The signals from either the low-gain or high-gain branch are simultaneously digitised by a 10-bit ADC every 25 ns [125]. The automatic gain selection mechanism uses the high-gain signal by default, but if the signal is saturated, it switches to the low-gain signal. The selected gain information is then sent to the back-end electronics for event processing.

The analogue electrical signal produced in TileCal cells has an amplitude corresponding to the amount of energy deposited in the cell, the signal shaper ensures all signals have the same width. The Optimal Filtering method is used to reconstruct the sampled signals [126, 127]. This method extracts the signal amplitude, a time offset τ with respect to the 40 MHz LHC clock, and the pedestal of the signal pulse in units of ADC counts for each sampled pulse shape. The signal reconstruction procedure is performed twice: in real time by the Read-Out Drivers for use in the triggers and offline after the data has been recorded for use in data analyses.

3.3 TileCal Calibration Systems

Three calibration systems are used to monitor the signal reconstruction process in TileCal cells. The calibration systems are used to maintain a time-independent electromagnetic energy scale, and account for changes in the hardware and electronics due to radiation damage, ageing and faults.

The Caesium, Cs, calibration system sets the global electromagnetic scale to monitor the optical and electrical responses of each PMT in the ATLAS environment [128]. A hydraulic system moves a capsule containing a ^{137}Cs radioactive γ -source through stainless steel tubes traversing every TileCal module, the holes for the tubes are illustrated in Figure 3.2. Corrections derived from Minimum Bias Trigger Scintillators are used to calibrate cells that do not have the ^{137}Cs system [129]. The 0.662 MeV photons from the ^{137}Cs radioactive source generate scintillation light in the cells. The system's design and configuration allows for precise examination and monitoring of the cell's optical path, from the scintillating tiles to the PMTs with better than 1% precision. The observed PMT signal indicates the scintillator and optical fibre quality of the cell.

The process of moving the ^{137}Cs source through the entire TileCal typically takes around six hours. This calibration procedure is not performed frequently

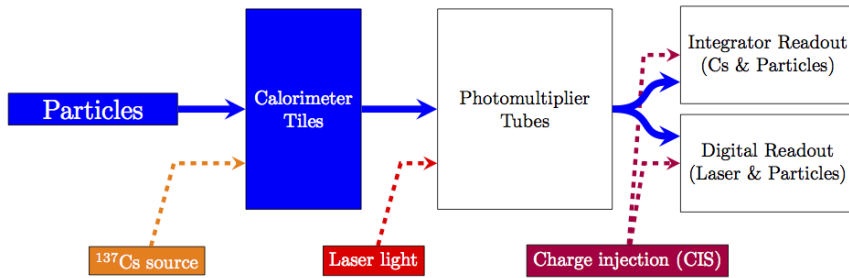


Figure 3.4 A schematic of the TileCal signal reconstruction path [130]. The path of the physics signal is shown by the thick solid lines and calibration paths of the different systems are indicated with dashed lines.

as it requires the ATLAS detector not to participate in taking any collision data for periods of up to six hours.

The laser calibration system monitors and corrects for PMT response variations in between Cs scans and also monitors timing during data-taking campaigns [118, 131]. The laser calibration constants are updated regularly to address the fast drift of PMT responses caused by changing instantaneous luminosity. It is not possible to perform the Cs calibration in gap and crack scintillators, so the first laser runs before data-taking of each year serves as the reference run.

The CIS monitors the front-end electronics by injecting a signal of known charge into the ADCs. It determines the calibration constants by analysing the relationship between the resulting electronic response, the ADC-count spectrum, and the injected signal pulse. Figure 3.4 shows a flow diagram that summarises the components of the TileCal read-out system tested by the different calibration systems.

A set of calibration constants f are used to convert the reconstructed channel amplitude in units of ADC counts to units energy in units of GeV. The conversion is expressed as:

$$E[\text{GeV}] = f_{\text{pC} \rightarrow \text{GeV}} \cdot f_{\text{Cs}} \cdot f_{\text{Las}} \cdot f_{\text{ADC} \rightarrow \text{pC}} \cdot A[\text{ADC}], \quad (3.1)$$

where $f_{\text{pC} \rightarrow \text{GeV}}$ is a calibration constant determined from dedicated electron and muon test beam analyses. The constants f_{Cs} , f_{Las} and $f_{\text{ADC} \rightarrow \text{GeV}}$ are obtained from the Cs, laser and CIS calibration systems, respectively.

Chapter 4

Long-Term Stability of Luminosity Measurements

ATLAS luminosity measurements rely on multiple luminosity detectors and algorithms, which have complementary capabilities and different systematic uncertainties. These luminosity algorithms provide measurements of a quantity that is proportional to the instantaneous luminosity, using a procedure specific to that particular detector. The luminosity detectors and algorithms can be categorised into those that provide bunch-by-bunch and bunch-integrated luminosity measurements.

The detectors and algorithms providing bunch-by-bunch luminosity measurements were introduced in Section 2.3.4. The ATLAS calorimeter systems provide bunch-integrated luminosity measurements. This is done by relating LAr gap currents in EMEC and FCal, and PMT currents of the TileCal to the instantaneous luminosity.

In this chapter, we briefly discuss the calibration transfer procedure, where luminosity measurements performed during low-luminosity runs are calibrated to high luminosities in physics data-taking regime using the van der Meer, vdM, technique. Several factors can affect the long-term stability of a given luminosity detector. The primary aim of this chapter is to utilise TileCal luminosity measurements to determine the uncertainty associated with long-term drifts in luminosity measurements provided by the LUCID detector during the 2022 data-taking campaign.

4.1 Luminosity Measurements at the LHC

The instantaneous luminosity resulting from a pair of colliding bunches, $\mathcal{L}_b$, can be expressed as:

$$\mathcal{L}_b = \frac{R}{\sigma}, \quad (4.1)$$

where R is the rate of events produced per unit of time for a process with cross-section, σ . In the context of inelastic pp collisions at the LHC with a cross-section σ_{inel} , Equation 4.1 can be expressed as:

$$\mathcal{L}_b = f_{\text{LHC}} \frac{\mu}{\sigma_{\text{inel}}}, \quad (4.2)$$

where f_{LHC} is the LHC bunch revolution frequency, and μ is the pileup parameter giving the number of inelastic interactions per bunch crossing.

The total instantaneous luminosity within a bunch crossing is given by:

$$\mathcal{L}_{\text{inst}} = \sum_{b=1}^{n_b} \mathcal{L}_b = n_b \langle \mathcal{L}_b \rangle = f_{\text{LHC}} \frac{n_b \langle \mu \rangle}{\sigma_{\text{inel}}}, \quad (4.3)$$

where n_b denotes the number of colliding bunch pairs, $\langle \mathcal{L}_b \rangle$ is the average luminosity per bunch pair, and $\langle \mu \rangle$ is the average pileup parameter for all colliding bunch pairs.

Any detector sensitive to inelastic pp interactions can be used as a luminometer. The ATLAS experiment employs multiple luminometers to measure the instantaneous luminosity by observing the visible interaction rate per bunch crossing, μ_{vis} . For a given luminometer $\mu_{\text{vis}} = \epsilon \mu$, where ϵ is the efficiency of the detector and algorithm under consideration and visible cross-section for the same detector and algorithm is defined as $\sigma_{\text{vis}} = \epsilon \sigma_{\text{inel}}$. The measured luminosity, using Equation 4.2, can be expressed as:

$$\mathcal{L}_b = f_{\text{LHC}} \left(\frac{\mu_{\text{vis}}}{\sigma_{\text{vis}}} \right), \quad (4.4)$$

where σ_{vis} is a calibration constant that represents the absolute luminosity calibration for a given luminometer. The process of determining σ_{vis} is known as calibration transfer.

4.1.1 The van der Meer Formalism

LUCID, the primary luminometer used in ATLAS, is calibrated in vdM runs where the absolute luminosity can be inferred from direct measurements of the beam parameters [132–135]. The luminosity per bunch crossing, expressed in terms of colliding-beam parameters, is given by:

$$\mathcal{L}_b = f_{\text{LHC}} n_1 n_2 \int \hat{\rho}_1(x, y) \hat{\rho}_2(x, y) dx dy, \quad (4.5)$$

where the colliding beams are assumed to collide with zero crossing angle, n_1 and $\hat{\rho}_1(x, y)$, and n_2 and $\hat{\rho}_2(x, y)$ are the number of protons and the normalised particle density in the transverse and horizontal plane of beam-1 and beam-2, respectively. The particle densities can be factorised into independent horizontal and vertical component distributions. The horizontal component of the distribution is given by:

$$\hat{\rho}_1(x, y) = \rho_x(x) \rho_y(y), \quad (4.6)$$

and the vertical component is defined in a similar manner. Equation 4.5 can be written as:

$$\mathcal{L}_b = f_{\text{LHC}} n_1 n_2 \int \rho_{x1}(x) \rho_{x2}(x) dx \int \rho_{y1}(y) \rho_{y2}(y) dy. \quad (4.7)$$

Utilising a technique suggested by van der Meer [132], one can compute the overlap integral in the x -direction as follows:

$$\int \rho_{x1}(x) \rho_{x2}(x) dx = \frac{R_x(0)}{\int R_x(\Delta x) d\Delta x}, \quad (4.8)$$

where $R(\Delta x)$ represents an uncalibrated instantaneous luminosity observed during scans with a horizontal separation between the beams of Δx , and $R(0)$ denotes the instantaneous luminosity when the two colliding beams have zero horizontal separation. The overlap integral in the y -direction is defined in a manner analogous to the definition in the x -direction. The convolved beam size in the horizontal plane is expressed as:

$$\Sigma_x = \frac{1}{2\pi} \frac{\int R_x(\Delta x) d\Delta x}{R_x(0)}, \quad (4.9)$$

while in the vertical plane, it is given by:

$$\Sigma_y = \frac{1}{2\pi} \frac{\int R_y(\Delta y) d\Delta y}{R_y(0)}, \quad (4.10)$$

as detailed in Reference [136]. The absolute luminosity using beam parameters in Equation 4.1.1 can then be expressed as:

$$\mathcal{L}_b = f_{\text{LHC}} \frac{n_1 n_2}{2\pi \Sigma_x \Sigma_y}. \quad (4.11)$$

The calibration constant, σ_{vis} , of a given luminosity algorithm can be determined by combining Equation 4.4 and 4.11 to obtain:

$$\sigma_{\text{vis}} = \mu_{\text{vis}}^{\text{max}} \frac{n_1 n_2}{2\pi \Sigma_x \Sigma_y}, \quad (4.12)$$

where $\mu_{\text{vis}}^{\text{max}}$ represents the visible interaction rate per bunch crossing measured at zero beam separation.

Figure 4.1 shows a typical horizontal scan from a vdM scan performed in November 2022 using the default LUCID luminosity algorithm Bi2HitOR. At each scan step, μ_{vis} , is corrected for afterglow effects, instrumental noise, and beam gas background, described in detail in Reference [138]. A Gaussian function is fitted to the distribution and the peak of the function provides the measurement of $\mu_{\text{vis}}^{\text{max}}$, and Σ_x is determined as per Equation 4.9.

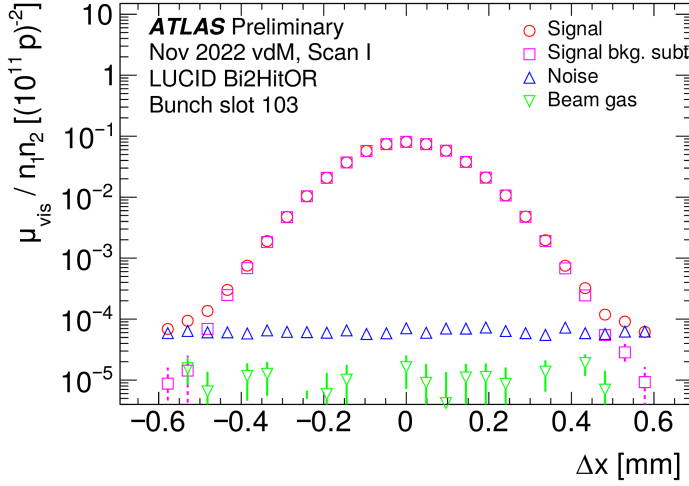


Figure 4.1 Distribution of $\frac{\mu_{\text{vis}}}{n_1 n_2}$ during horizontal vdM scan. The measured rate is shown by the red circles, the background subtracted rate by the pink squares, the estimated background from noise and afterglow by the blue upward-pointing triangles and that from beam–gas interactions by the green downward-pointing triangles. The error bars show statistical uncertainties, which are in some cases smaller than the symbol sizes. [137].

The LUCID calibration constant, σ_{vis} , is determined under low- μ data-taking conditions, involving a limited number of isolated bunches with peak instantaneous luminosities that are at least three orders of magnitude smaller than those in typical physics data-taking regime. There is a strong non-linear dependence of σ_{vis} on run conditions such as pileup [138]. Luminosity values obtained from track-counting algorithm are robust against pileup and are used to extrapolate σ_{vis} from the low pileup vdM conditions to the high pileup physics data-taking conditions. The calibration transfer procedure is validated using calorimeter luminosity measurements.

The luminosity detectors measure the rate μ_{vis} over a finite time interval called a luminosity block, LB, during which data-taking conditions are assumed to remain constant. The LB durations, measured in seconds, are defined by the ATLAS central trigger processor [139, 140]. Similar to Run 2, in Run 3, the LB duration is typically around one minute during normal detector operation.

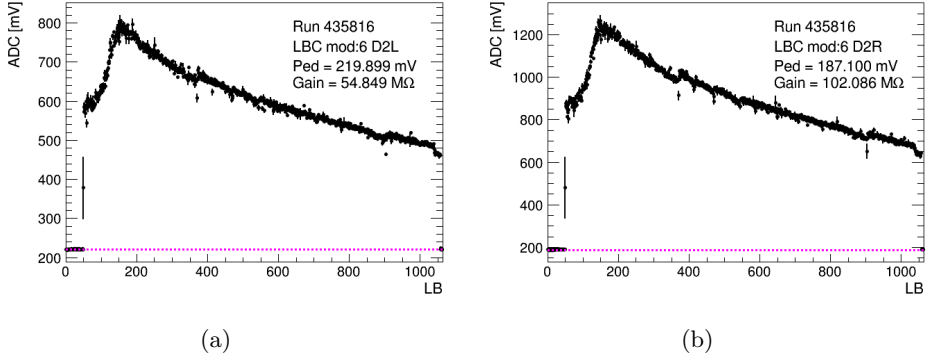


Figure 4.2 Distributions of ADC counts in module 6 of the LBC segment of the TileCal during Run 435816. The magenta line represent the pedestal.

The ATLAS central trigger processor sometimes shortens the LB duration when there are problems affecting the detector's data acquisition.

4.2 TileCal as a Luminometer

Originally, the TileCal was not designed for use as a luminometer. However, the anode current, I_{PMT} , drawn by each PMT, is proportional to the number of particles passing through a specific TileCal cell and is thus linked to the bunch-integrated luminosity.

The anode current is computed from the raw response of the PMTs which is measured in ADC counts, as discussed in Section 3.2. The average anode current of a PMT is computed by the relation:

$$\langle I_{\text{PMT}} \rangle [\text{nA}] = \frac{\langle \text{ADC} \rangle [\text{mV}] - p [\text{mV}]}{\text{Int. gain} [\text{M}\Omega]}, \quad (4.13)$$

where $\langle \text{ADC} \rangle$ represents the average ADC counts in a LB, p is the pedestal and Int. gain is the integrator gain constant calculated using the CIS calibration.

Figure 4.2 illustrates the distributions of PMT responses measured in ADC counts in module 6 of the LBC segment during Run 435816. The optical

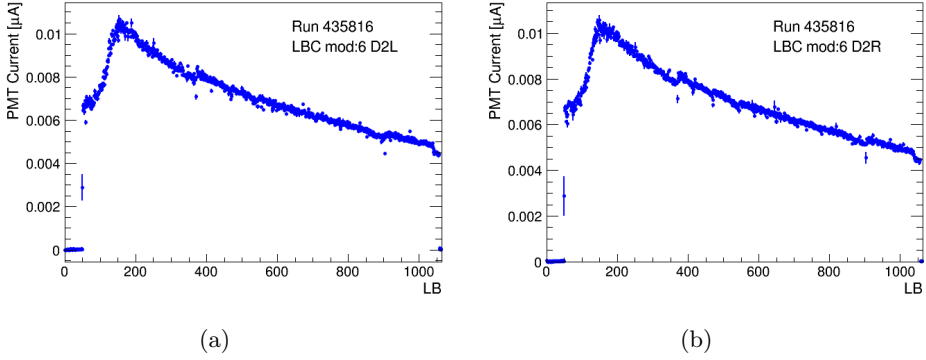


Figure 4.3 Distributions of raw PMT currents in module 6 of the LBC segment of the TileCal during run 435816.

response of the cells is read out by two PMTs, as discussed in Section 3.2. The pedestal, indicating the baseline ADC measurement primarily due to electronic noise, is calculated as the mean ADC counts before the beams are brought into collisions. Figure 4.2(a) and 4.2(b) show the response in the left and right PMTs, respectively.

The anode currents drawn by left and right PMTs of cell D2 located in module 6 are shown in Figure 4.3(a) and 4.3(b), respectively. The anode currents, I_{PMT} , are converted to luminosity through calibration constants. The relationship between I_{PMT} and luminosity is expressed as:

$$\mathcal{L}_{\text{PMT}} = \alpha \langle I_{\text{PMT}} \rangle, \quad (4.14)$$

where α is a calibration constant, and $\langle I_{\text{PMT}} \rangle$ is average anode current in a LB.

Figure 4.4 displays a linear dependence between the drawn PMT currents and luminosity across all three radial segments of TileCal. The PMT currents vary with η location of a cell and the radial sampling layer location. Cells located at $\eta \approx 1.25$ experience higher particle fluxes and thus higher currents. At a fixed η value, currents are highest in cells located in the innermost sampling

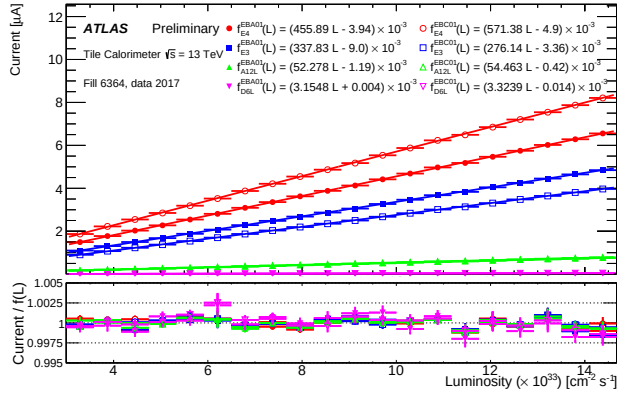


Figure 4.4 The dependence of the PMT anode current on track-counting Luminosity for a few channels of the Tile Calorimeter during fill 6364 in 2017 data-taking period [141].

layer and lowest at the outermost sampling layer due to progressive absorption of hadronic showers.

In principle all TileCal cells can be used in luminosity measurements. The cell position in the detector influences its sensitivity and exposure to radiation. In Figure 4.4, the magenta lines shows that the outermost radial layer, D-layer, has insufficient sensitivity to be calibrated during the low luminosity vdM scans. The D-cells exhibit the least response variation over time due to radiation damage, making them preferred for long-term luminosity measurements. The D5 and D6 cells are used in primary luminosity measurements, while D1 and D2 cells are utilised for systematic comparisons.

Similar to EMEC and FCal currents, TileCal currents are cross-calibrated to track-counting luminosity measurements. Figure 4.5 illustrates the ratio of the PMT currents of the TileCal over the track-counting luminosity. The ratios in the left and right PMTs are shown in Figure 4.5(a) and 4.5(b), respectively. The anchoring constant for a PMT in each module is calculated as:

$$\alpha_{\text{module, PMT}} = \frac{\mathcal{L}_{\text{Track-counting}}}{\langle I_{\text{PMT}} \rangle_{\text{module}}}. \quad (4.15)$$

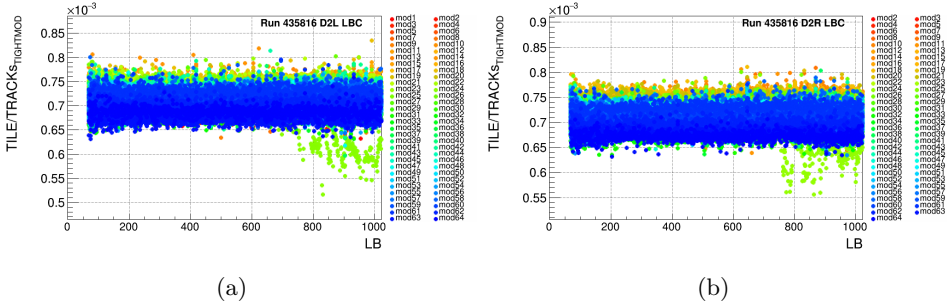


Figure 4.5 Distributions of raw PMT currents over track-counting luminosity in Run 435816. Each plot displays the ratio of PMT current to track counting luminosity for different TileCal modules, module numbers are specified in the legend.

Figure 4.5 shows the inverse of the anchoring constant used in cross-calibrating TileCal PMT currents to track-counting luminosity.

The derived anchoring constants for each PMT are used to compute luminosities for any PMT current using Equation 4.14. The physics Run 435816 is used to derive the TileCal anchoring constants in the 2022 data-taking campaign. Run 435816 was chosen as the anchoring run due to its wide luminosity range and smooth changes in instantaneous luminosity, making it ideal for deriving the calibration constants. The TileCal cell luminosity within a module is determined by calculating the average of luminosity measurements from the left and right PMTs. The luminosity of each cell is calculated as the average luminosity across all modules. Figure 4.6 shows the comparison between the track-counting luminosity and the TileCal luminosity in the D2 cell. The figure shows a notable agreement between the track-counting and TileCal luminosities, indicating the success of the cross-calibration procedure in accurately converting anode PMT currents to TileCal luminosities.

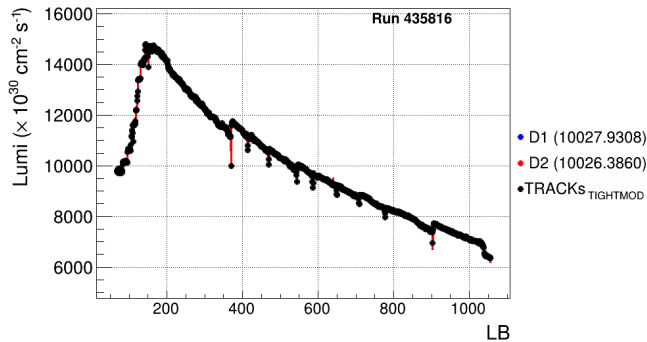


Figure 4.6 Comparison between TileCal D2 cell and track-counting luminosities.

4.3 Long-Term Stability of Luminosity Measurements

The calibration of LUCID during vdM scans is performed at low values of μ , $\mu \approx 0.5$, in contrast to higher μ values, typically $\mu \approx 50$, encountered during physics data-taking runs. In physics data-taking runs, the number of bunch trains and the crossing angle also increase significantly. The natural luminosity decay in long LHC physics fills, where both the luminosity and the bunch average pileup, $\langle\mu\rangle$, decrease by a factor of two or three, is utilised to study the relative response of different luminosity algorithms as a function of $\langle\mu\rangle$.

Figure 4.7 shows the ratio of bunch-averaged track-counting luminosity to LUCID luminosity as a function of $\langle\mu\rangle$. The default LUCID luminosity algorithm, Bi2HitOR, is used to record LUCID luminosities.

Figure 4.7 shows that the track-counting/LUCID ratios are typically $\mathcal{O}(10\%)$ smaller than unity at the start of the physics fill. The ratio exhibits a strong linear dependence on $\langle\mu_{\text{LUCIDBi2HitOR}}\rangle$, simplifying the correction of the LUCID response as a function of $\langle\mu\rangle$ in high-luminosity LHC fills each year of data-taking. The LUCID corrections are derived in the form:

$$\langle\mu_{\text{corr}}\rangle = p_0\langle\mu_{\text{uncorr}}\rangle + p_1(\langle\mu_{\text{uncorr}}\rangle)^2, \quad (4.16)$$

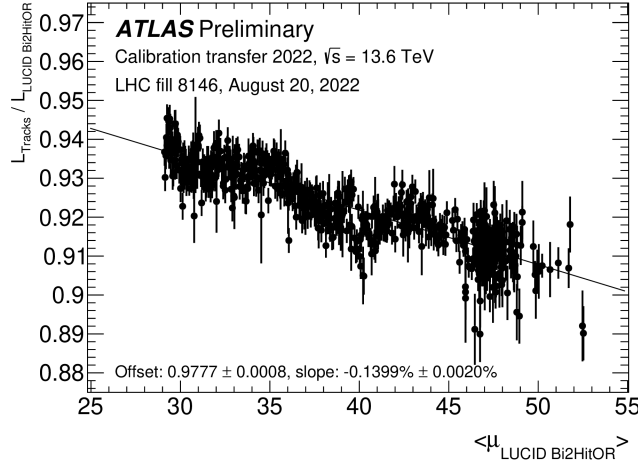


Figure 4.7 Ratio of the bunch-averaged luminosity L_{Tracks} , reported by the track-counting luminosity algorithm, to $L_{\text{LUCID Bi2HitOR}}$, obtained using the LUCID Bi2HitOR algorithm before any pileup corrections, as a function of the bunch-averaged pileup parameter, $\langle \mu_{\text{LUCID Bi2HitOR}} \rangle$, during LHC fill 8146 [137].

where $\langle \mu_{\text{uncorr}} \rangle$ and $\langle \mu_{\text{corr}} \rangle$ are the uncorrected and corrected LUCID $\langle \mu \rangle$, $\langle \mu_{\text{LUCID Bi2HitOR}} \rangle$, respectively. The parameters p_0 and p_1 are obtained from linear a fit to track-counting/LUCID luminosity ratios in Figure 4.7.

The correction of LUCID using Equation 4.3 ensures agreement between track-counting and LUCID luminosities during high-luminosity reference fills. A single high-luminosity reference fill is generally sufficient to correct LUCID for an entire year of data-taking, assuming no time dependence. The ^{207}Bi calibration system employed by LUCID approximately validates this assumption. The gain monitoring and relative calibration of LUCID PMTs using a ^{207}Bi source is described in detail in Reference [111].

Drifts in LUCID response not fully corrected by the ^{207}Bi calibration system or drifts in track-counting luminosity measurements used for LUCID corrections can introduce time-dependent bias in LUCID luminosity measurements. These potential biases are studied and quantified by comparing run-integrated

luminosity measurements from LUCID with independent luminosity measurements from EMEC, FCal, and TileCal luminosity algorithms.

The run-integrated luminosities are compared for physics runs with at least 100 LB in the good run list, GRL. The GRL contains a list of runs and LBs considered good for physics analysis, defined based on the data quality criteria such as detector performance, stability, and conditions during data taking. Since computation of TileCal PMT currents uses the pedestal, defined as the mean number ADCs before beams are brought into collisions, runs that started after pp collisions commenced in the detector are excluded.

Fractional differences between run-integrated luminosities from TileCal D-cells and those from other luminometers are depicted in Figures 4.8, 4.9, 4.10 and 4.11. These figures show the ratio of luminosity measurements from TileCal D-cells, excluding D0 cell, to those from other luminometers as a function of run number. The figures provide insight into trends and variations in TileCal luminosities over time relative to luminosity measurements from other algorithms. In instances where there is an outlier from the trend and the ratios in closely align, it suggests that the luminosity measurements from TileCal are contributing to the outlier in that specific run.

The variation in TileCal luminosity measurements compared to luminosity measurements from other algorithms typically remains within 2%, except for early runs with a large deviation relative to FCal luminosity measurements. Studies in Reference [138] have identified a reproducible non-linearity in FCal luminosity measurements that strongly depends on the instantaneous luminosity. The downward slope observed in Figures 4.8, 4.9, 4.10 and 4.11 indicates clear evidence of PMT ageing in TileCal cells. The ageing is caused by radiation exposure and its severity is dependent on the η position of a cell. It can be seen that D1 cells exhibit the least damage, while the most pronounced damage is observed in D5 cells.

The long-term stability of LUCID luminosity measurements is assessed by comparing the run-integrated luminosity measurements from calorimeter algorithms with those from LUCID. The long-term stability uncertainty of lumi-

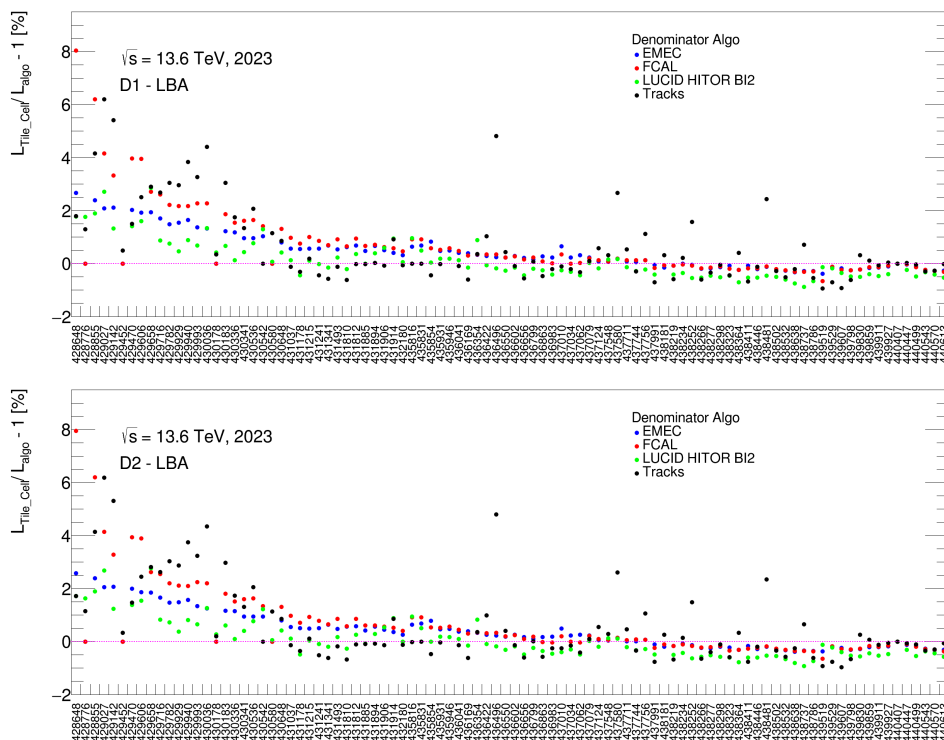


Figure 4.8 Run-integrated ratios of luminosity measured by EMEC and FCAL to that measured by TileCal D1 and D2 cells on the A-side of the detector.

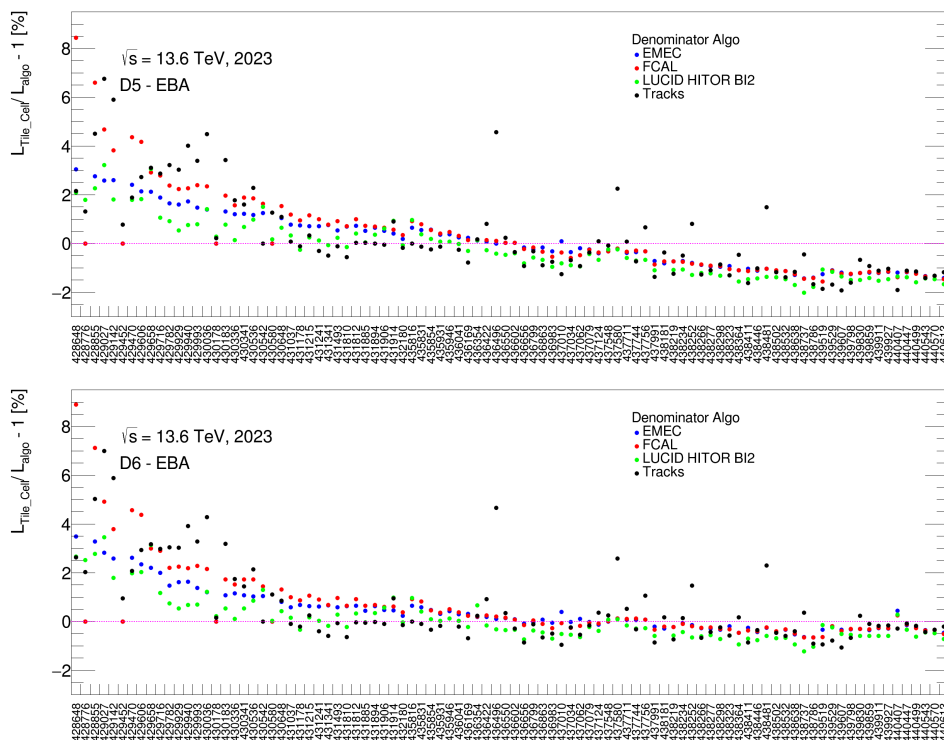


Figure 4.9 Run-integrated ratios of luminosity measured by EMEC and FCAL to that measured by TileCal D1 and D2 cells on the A-side of the detector.

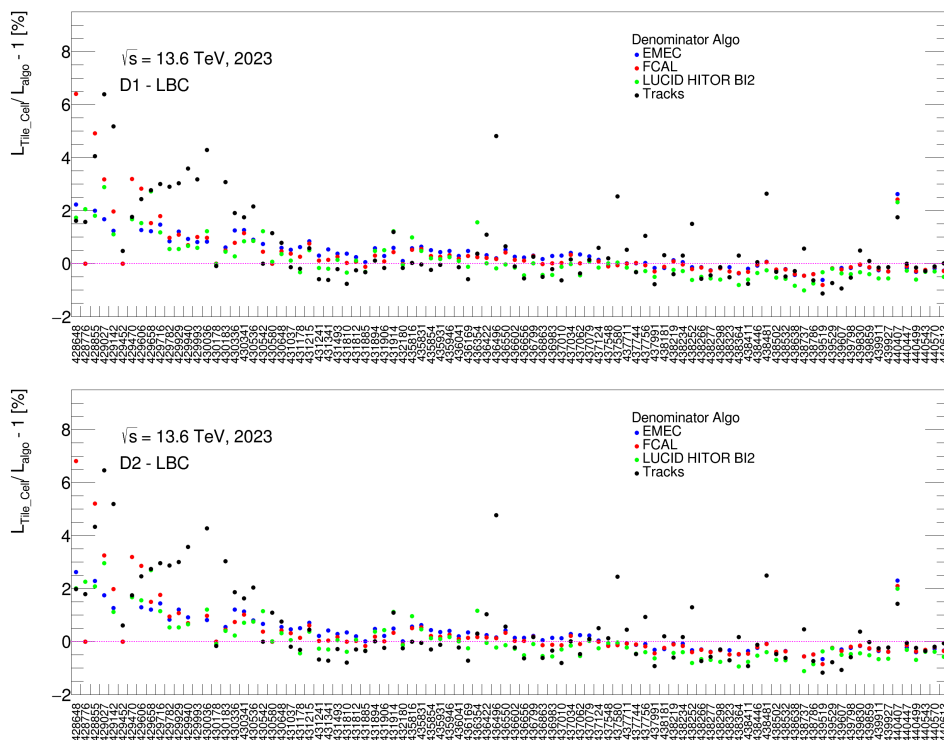


Figure 4.10 Run-integrated ratios of luminosity measured by EMEC and FCal to that measured by TileCal D1 and D2 cells on the C-side of the detector.

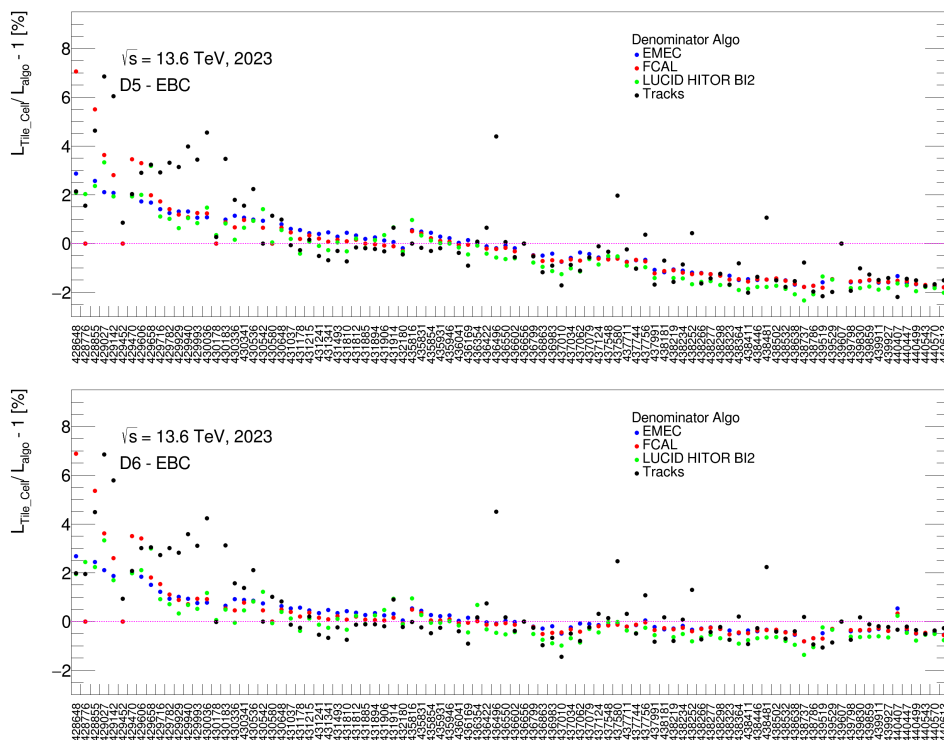


Figure 4.11 Run-integrated ratios of luminosity measured by EMEC and FCal to that measured by TileCal D5 and D6 cells on the C-side of the detector.

luminosity measurements in a data-taking year is defined as the largest measured mean difference between LUCID and calorimeter-based measurements.

Figures 4.8, 4.9, 4.10 and 4.11 show the run-to-run spread between luminosity measurements, but it shows all the runs with equal weight. The run-to-run comparison of TileCal and LUCID luminosity measurements during the 2022 data-taking is summarised in Figure 4.12. Figure 4.12 gives an alternative presentation of fractional differences between TileCal and LUCID measurements in Figures 4.8, 4.9, 4.10 and 4.11, with the exception that the fractional luminosity differences are weighted by the integrated luminosity of each run. This representation gives a higher weight to longer runs, runs with large integrated luminosity.

The comparison between the TileCal luminosity and LUCID measurements are made separately for the A and C side of the detector. The mean and RMS values in the legend represent fractional differences in integrated luminosity over the entire data-taking year and the spread in the fractional differences, respectively. The mean deviation between luminosity measurements from TileCal cells and those from LUCID range between 0.065% and 0.805%. To establish the preliminary long-term stability uncertainty, we compare the mean deviation in luminosity measurements from TileCal D6 cells with the mean deviation in FCal and EMEC luminosity measurements. The decision to use D6 cells for comparison was based historical usage, as they were utilised in Run 2 measurements [138]. FCal measured the largest deviation of 0.41%, EMEC has a deviation of 0.27% and TileCal D6 cell have a deviation of 0.233%. These preliminary TileCal results are without application of corrections from the laser calibration to correct for ageing of the PMTs.

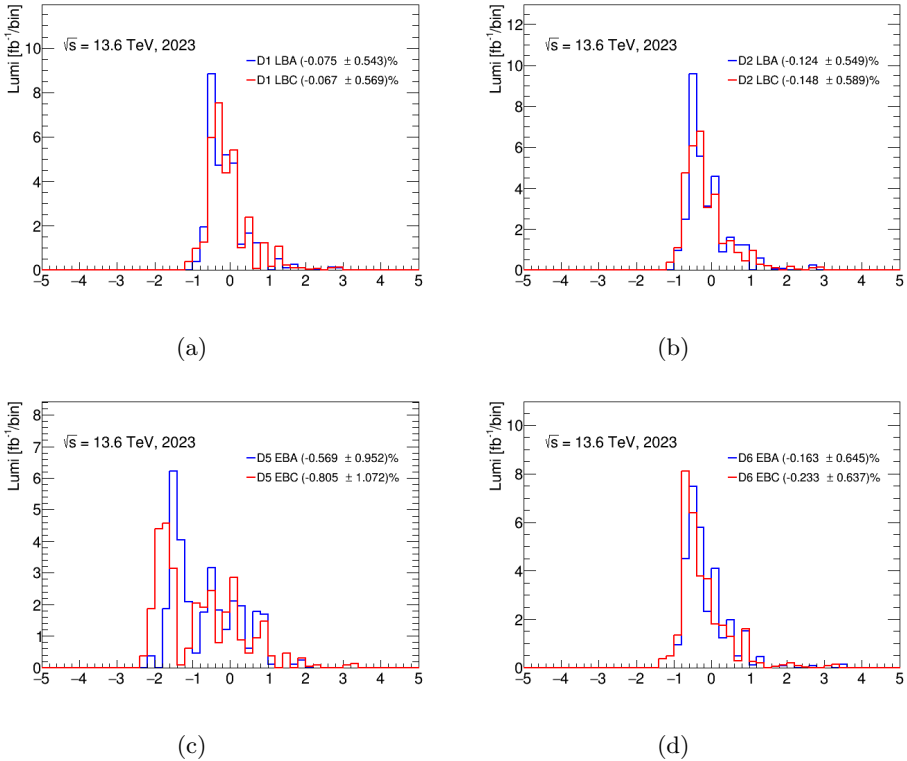


Figure 4.12 Fractional differences in run-integrated luminosity between TileCal D1(a), D2(b), D5(c) and D6(d) luminosity measurements and the baseline LUCID luminosity measurements, weighted by the integrated luminosity of each run. The mean and RMS values for each cell are indicated in the legend.

Chapter 5

Performance of Gap and Crack Tile Scintillators

As mentioned previously, calorimeters play a pivotal role in general-purpose hadron collider detectors. In the context of the LHC’s physics program, the TileCal plays a critical role in the reconstruction of isolated hadrons, jets, τ -leptons, and E_T^{miss} arising from non-interacting particles. It also provides essential inputs for electron isolation studies and muon triggering.

This chapter is dedicated to the comprehensive assessment of TileCal’s performance during the Run 2 ATLAS data-taking period. Performance studies allow us to verify if the signal reconstruction and calibration methods, discussed in detail in Section 3.3, were applied correctly across all modules. In this chapter, single isolated particles are used to assess the performance of the TileCal by studying the response of gap and crack scintillators.

5.1 Introduction

This study evaluates the performance of the TileCal by analysing the response of gap and crack scintillators to isolated muons. The location and the role of the gap and crack scintillators within TileCal were discussed extensively in Chapter 3. Figure 3.2 shows the location of gap and crack scintillators, highlighted in orange. The muons used as probe particles in this study are

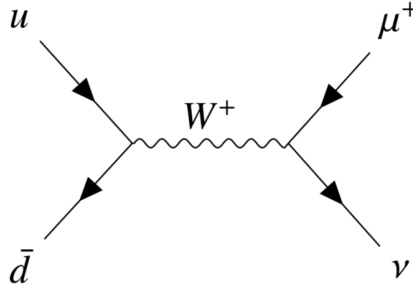


Figure 5.1 Feynman diagram illustrating the production of a W^+ boson in pp collisions, which then decays leptonically into a muon and a muon neutrino. The diagram for W^- is identical but with the charges flipped.

required to originate from the decay of a W boson into a muon and a muon neutrino, $W \rightarrow \mu\nu_\mu$. The Feynman diagram in Figure 5.1 illustrates the $W^+ \rightarrow \mu^+\nu_\mu$ process.

The response of a TileCal cell as muons traverse through it is quantified by the amount of energy, ΔE , deposited per unit length, Δx , in that particular cell. The energy loss of muons as they traverse through matter follows a well-understood process [142], where ionisation is the dominant mechanism for muons with energies below 100 GeV. The probability distribution of the response typically follows the Landau function for thin scintillators.

Muons originating from the $W \rightarrow \mu\nu_\mu$ process are selected due to their distinct experimental characteristics within the detector, allowing for efficient background suppression. In the $W \rightarrow \mu\nu_\mu$ channel, the electroweak background is predominantly composed of $W \rightarrow \tau\nu_\tau$ and $Z/\gamma^* \rightarrow \mu^+\mu^-$ processes.

Background events from the $Z/\gamma^* \rightarrow \mu^+\mu^-$ process are suppressed by requiring the selected events to have exactly one muon candidate. The decay of the τ -lepton into a muon and a τ neutrino is a complex process that involves multiple particles, leading to a displaced τ decay vertex and, on average, a muon with low p_T [143, 144]. The p_T of the selected muon candidates in addition to the position of the primary vertex is used to suppress these background events. In events where heavy quarks undergo semi-leptonic decays, there is a possi-

Table 5.1 Data-taking periods of the LHC in Run 2 and their corresponding integrated luminosity.

Period	$\int \mathcal{L} dt [\text{fb}^{-1}]$
2015	3.2
2016	33.0
2017	44.3
2018	58.5

bility of hadrons to be misidentified as leptons. These multi-jet background processes substantially contribute to the total background. The presence of $E_{\text{T}}^{\text{miss}}$ in the events is used to suppress these multi-jet backgrounds.

5.2 Datasets

Experimental pp collision data collected during the LHC Run 2 data-taking at $\sqrt{s} = 13$ TeV is used. Table 5.1 summarises the data-taking years of Run 2 alongside their corresponding integrated luminosities.

The conditions of the TileCal are subject to change over the data-taking years due to various factors such as module malfunctions, radiation damage, and updates to calibration constants. The response of the gap and crack scintillators of the TileCal is analysed in three distinct time periods: 2015-2016, 2017, and 2018 data-taking periods. Data from the 2015 data-taking period is combined with data collected in the 2016 data-taking period due to their low integrated luminosities.

Monte Carlo, MC, simulated samples of $W \rightarrow \mu\nu_\mu$ events are used as a reference for studying the response from experimental data. Three sets of MC event samples were generated to match pileup conditions of the three data-taking periods: 2015-2016, 2017, and 2018. The event generator SHERPA [145] was used to generate the W boson production events, which are then interfaced with PYTHIA8 [146] for parton showering. The ATLAS geometry, track-

Table 5.2 Event selection based on the W decays.

	Variable	Requirement
1	Number of Muons	$N_{\text{muons}} = 1$
2	Transverse invariant mass	$40 < M_T < 140 \text{ GeV}$
3	Missing transverse energy	$30 < E_T^{\text{miss}} < 120 \text{ GeV}$
4	Track isolation	$\sum p_T _{\Delta R=0.4} < 1 \text{ GeV}$
5	Calorimeter isolation	$E_{\text{LAR}} _{\Delta R=0.4} < 1.5 \text{ GeV}$
6	Momentum of the muon	$p^\mu \leq 80 \text{ GeV}$
7	Transverse momentum of the muon	$p_T^\mu > 28 \text{ GeV}$
8	Average interactions per bunch crossing	$10 \leq \langle \mu \rangle \leq 40$

ing tools and physics models describing particle interactions with the detector material are provided by the GEANT4 toolkit [147, 148].

5.3 Event Selection

Table 5.2 summarises the selected events used in this analysis. The selections were chosen to optimise the number of events with a single muon originating from the $W \rightarrow \mu\nu_\mu$ process while minimising the number of events from background processes.

The primary single muon trigger was used to select the candidate muons. The ATLAS trigger system uses a trigger menu that defines a trigger selection criteria that a candidate physics object has to satisfy. Due to changing conditions at the LHC and continuous improvements to the trigger algorithms, the trigger menu is subject to frequent changes. For events recorded in 2015, the HLT_mu20_loose_L1MU15 trigger was used [149]. The software-based HLT was configured to require the muon candidates to have $p_T > 20 \text{ GeV}$ and satisfy the Loose isolation requirement. The HLT trigger is seeded by a hardware-based L1 trigger, which selects muon candidates with p_T above 15 GeV.

The instantaneous luminosity increased significantly in events recorded between 2016 and 2018 as compared to those recorded in 2015. The muon trigger HLT_mu26_ivarmedmedium was used during this period, selecting muon candidates with $p_T > 26$ GeV that pass the Medium isolation requirement [140]. The requirement for the muons to have $p_{T^\mu} > 28$ GeV is set to ensure that the selected muons are well within the trigger efficiency plateau. This higher p_{T^μ} threshold ensures robust muon selection across different data-taking periods and trigger configurations, particularly given the variations in trigger thresholds over time. Detailed descriptions and performance evaluations of the ATLAS muon triggers in Run 2 can be found in Reference [150].

In this analysis, selected events are required to contain exactly one muon that passes either the Medium or Tight identification requirement [151]. This criterion ensures that the number of muon candidates is restricted to minimise contamination from electroweak background events.

A cut is applied on the reconstructed transverse mass, m_T , of the system defined as:

$$m_T = \sqrt{2p_T^\mu E_T^{\text{miss}}(1 - \cos[\Delta\phi(\mathbf{p}_T^\mu, \mathbf{E}_T^{\text{miss}})]), \quad (5.1)$$

where $\mathbf{p}_T^\mu$ is a vector representing the candidate muon's p_T and $\mathbf{E}_T^{\text{miss}}$ is a vector representing the E_T^{miss} . The scalar variables are $p_T^\mu \equiv |\mathbf{p}_T^\mu|$ and $E_T^{\text{miss}} \equiv |\mathbf{E}_T^{\text{miss}}|$. The selected events were required to have $40 < m_T < 140$ GeV to exclude QCD multi-jet processes and to ensure that the energy loss of muons in the cells is primarily due to ionisation.

To further reduce backgrounds with low E_T^{miss} from jet production at low E_T^{miss} and $t\bar{t}$ process which can contain neutrinos of higher momentum, events with $30 < E_T^{\text{miss}} < 120$ GeV were selected. Cuts on the track, $\sum p_T|_{\Delta R=0.4}$, and calorimeter, $E_{\text{LAr}}|_{\Delta R=0.4}$, based isolation variables were applied to ensure that the selected muon candidate is well isolated from the hadronic activity around it.

To reduce the dependency of the cell's response of the muon's momentum, the selected muon candidate were required to have $p^\mu \leq 80$ GeV. This selection

also minimises the effects of rare energy loss processes such as bremsstrahlung or the energetic delta, δ , rays.

Event selection requirements from Table 5.2 were imposed on both the data and MC samples. The events from the simulated samples were normalised to the $W \rightarrow \mu\nu_\mu$ production cross-section multiplied by the integrated luminosity of the respective data-taking period. The simulated events were also reweighted to match the $\langle\mu\rangle$ distribution in data to remove any discrepancies that are correlated with pileup.

To demonstrate that the selected events mainly consist of $W \rightarrow \mu\nu$ events and that the simulation accurately represents the ATLAS geometry, a comparison is made between events in data and MC. Figures 5.2 and 5.3 show the comparison between data and MC distributions of some of the kinematic variables used in the analysis. The distributions show a notable agreement between data and MC, the little discrepancies that exist are due to background from multi-jet production.

In addition to the event selection criteria outlined in Table 5.2, the energy deposited by the passing muon candidate in the gap and crack cells was required to be above 60 MeV in order to suppress electronic noise signals. The selected muon was required to traverse a minimum path length of 11 mm in gap cells and 5 mm in crack cells. The track path length is defined as a straight line between two opposite surfaces of the scintillator. The required absolute azimuthal angle between the muon's track and cell's centre was $|\Delta\phi(\mu, \text{cell})| < 0.046$. The measured deposited energy in a cell depends on the position of local energy deposit, thus, the requirement on $\Delta\phi(\mu, \text{cell})$ ensured that the tracks of the selected muon candidate pass close to the centre of the cell. Table 5.3 summarises the cell level selections.

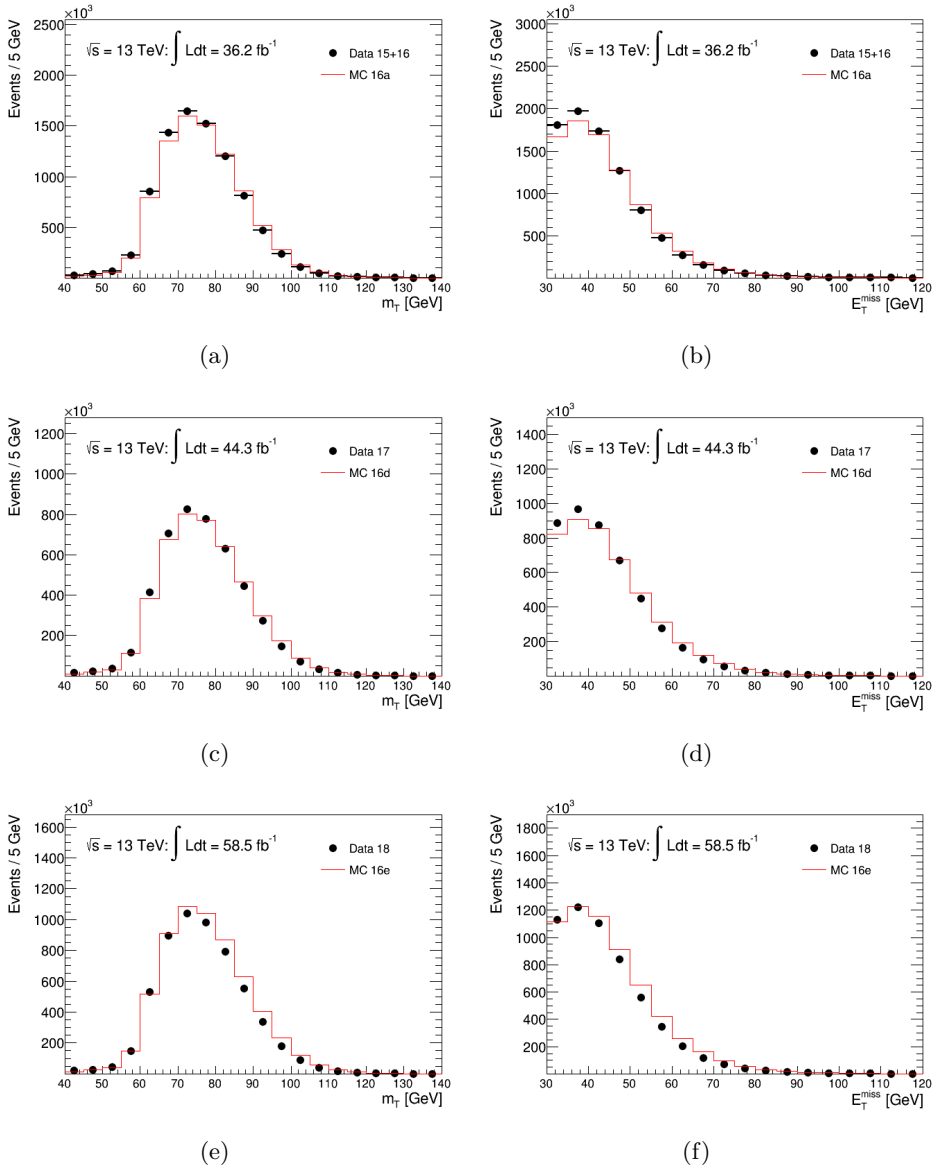


Figure 5.2 Distributions of the transverse invariant mass, m_T , on the left reconstructed muon candidates and the reconstructed E_T^{miss} on the right. The plots compare data taken during the 2015-2016, 2017, and 2018 periods. The top, middle, and bottom rows represent the respective data-taking periods.

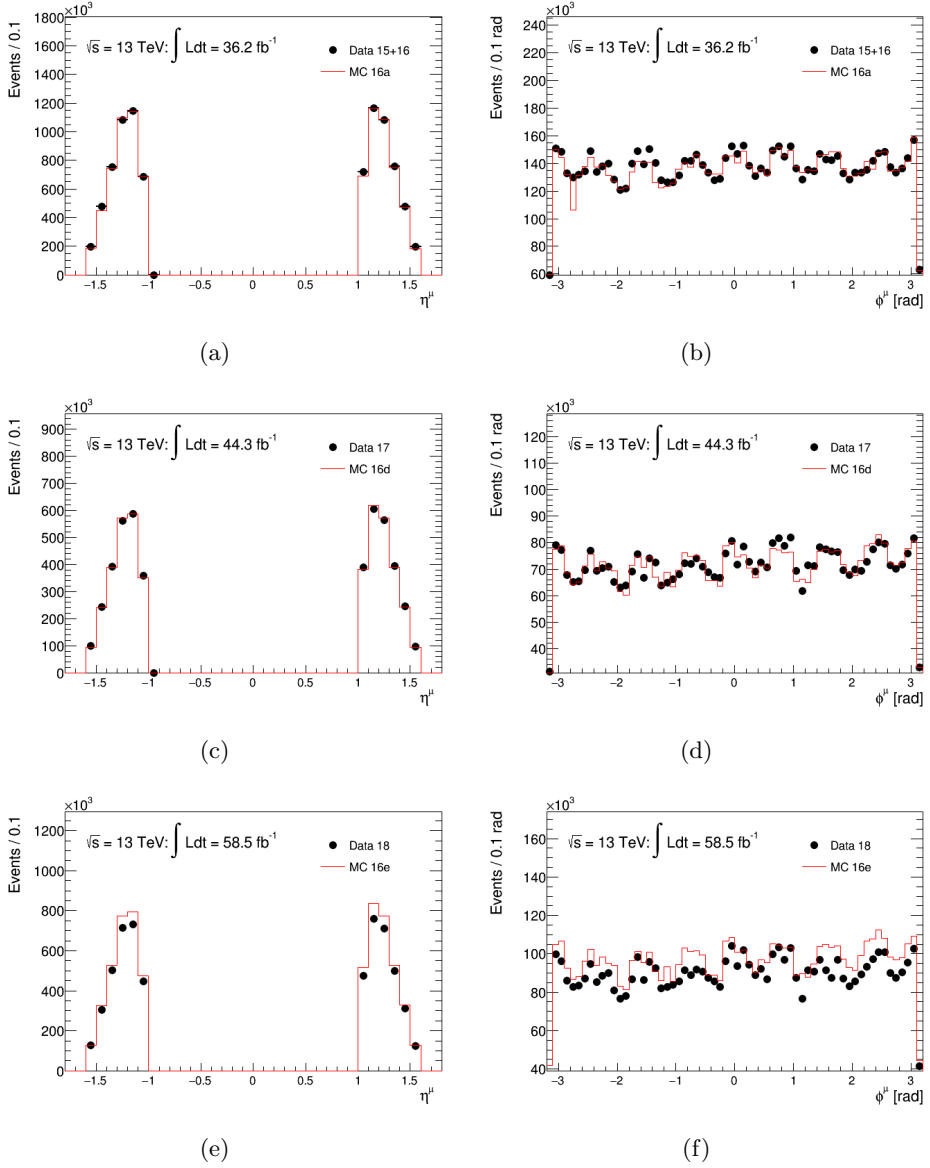


Figure 5.3 Distributions of the reconstructed muons in the η and the azimuthal, ϕ , directions. The plots compare data taken during the 2015–2016, 2017, and 2018 periods. The top, middle, and bottom rows represent the respective data-taking periods.

Table 5.3 Cell level selections.

	Variable	Requirement
1	Muon path length	E1, E2: $dx > 11$ mm; E3, E4: $dx > 5$ mm
2	Cell energy	$\Delta E > 60$ MeV
3	Track impact point	$ \Delta\phi(\mu, \text{cell}) < 0.046$

5.4 Calorimeter Response

The ratio $\Delta E/\Delta x$ serves as a measure of a cell's response to passing muons, acting as a proxy for dE/dx , particularly in thin layers of scintillating material. The dE/dx is a well understood phenomena, with its average described by the Bethe Block formula [152] and its probability distribution follows a Landau function [142].

Ideally, the response of a cell, $\Delta E/\Delta x$, can be modeled as a convolution of a Landau function and a Gaussian distribution, where the former accounts for the actual energy loss and the latter for noise. Previous attempts to fit the $\Delta E/\Delta x$ distribution with such model have failed to pass a χ^2 quality of fit test. The failure is attributed to the noise dependence of factors like electronic, pileup noise and cell size, which cannot be modelled by a single Gaussian distribution.

To measure cell response, the truncated mean of the $\Delta E/\Delta x$ distribution was used. This mean was computed by excluding 1% of entries with the highest $\Delta E/\Delta x$ values, effectively suppressing the impact of rare energy loss processes such as bremsstrahlung or energetic δ rays. The truncated mean exhibits a slight non-linear scaling with the path length Δx [153]. This non-linearity and other residual non-uniformities, such as the differences in momentum and incident angle spectra, are to a large extent reproduced by the MC simulations. To compensate for these effects, the response of a given cell was obtained by taking the ratio of the observed truncated mean $\Delta E/\Delta x$ with the one from

the MC simulation:

$$R \equiv \frac{\left\langle \frac{\Delta E}{\Delta x} \right\rangle_F^{\text{Data}}}{\left\langle \frac{\Delta E}{\Delta x} \right\rangle_F^{\text{MC}}}, \quad (5.2)$$

where F is the fraction of events truncated, set to be 1%. The truncated mean ratio R quantifies the consistency between experimental data and MC simulation. The statistical uncertainty of R is calculated by propagating the uncertainties of $\Delta E/\Delta x$ distributions from data and MC simulations:

$$\sigma_R = \left| \frac{r^{\text{data}}}{r^{\text{MC}}} \right| \times \sqrt{\left(\frac{\sigma_r^{\text{data}}}{r^{\text{data}}} \right)^2 + \left(\frac{\sigma_r^{\text{MC}}}{r^{\text{MC}}} \right)^2}, \quad (5.3)$$

where r is a shorthand notation for the truncated mean, $\langle \Delta E/\Delta x \rangle_F$. The uncertainty of the truncated mean, σ_r , is taken as the standard error of mean defined as:

$$\sigma_r = \frac{\text{std. dev.}}{\sqrt{N_{\text{eff}}}}, \quad (5.4)$$

where std. dev. is the standard deviation of the truncated $\langle \Delta E/\Delta x \rangle_F$ distribution, and N_{eff} is the effective number of entries of the distribution.

The double ratio, R , given in Equation 5.2, can be used to study the calorimeter response with respect to various detector geometrical properties. In this chapter, we examine the response of gap and crack scintillators as a function of ϕ during the 2015+2016, 2017 and 2018 data-taking periods.

The ^{137}Cs calibration system is used to achieve intercalibration of cell responses, equalising the cell responses to be within 0.2% to 0.3% for standard TileCal cells and 2.4% for TileCal cells composed solely of scintillating material [128, 154]. Non-uniformities in the calibration and the response may arise due to optical properties of the scintillators that degrade over time due to radiation exposure.

Due to the cylindrical symmetry of the TileCal and the even distribution of particles produced in pp collisions, it is expected that R computed using collision muons crossing any cell, c , to be equal across all modules. The R

uniformity of each cell was evaluated over the entire azimuthal direction on both the A and C sides of the detector. This evaluation involved maximising a Gaussian likelihood function over the modules of each cell, where the likelihood function accounts for both observed R values and their statistical uncertainties. The maximised Gaussian likelihood function is given by:

$$\mathcal{L}_c = \prod_{m=1}^{64} \frac{1}{\sqrt{2\pi}\sqrt{\sigma_{c,m}^2 + s_c^2}} \exp \left[-\frac{1}{2} \frac{(R_{c,m} - \mu_c)^2}{\sigma_{c,m}^2 + s_c^2} \right], \quad (5.5)$$

where the product runs over the 64 modules in detector slices of $\Delta\phi \approx 0.1$. The coordinate centre of each module is given by $\phi(m) = -\frac{\pi}{64} + \frac{2\pi}{64} \times m$. The $R_{c,m}$ and $\sigma_{c,m}$ are the observed R and its statistical uncertainty for a given cell in module m .

The unbinned log likelihood $-2 \log \mathcal{L}$ is maximised to find $\hat{\mu}_C$ and $\hat{s}_C$, which are the average response, $\langle R \rangle$, of data-MC agreement and systematic uncertainty attributed to the non-uniformity across the modules, respectively. The $\hat{s}_C$ parameter becomes non-zero by a combination of various factors such as light collection spread in different modules and miscalibration of certain modules for a given cell type.

Figures 5.4 and 5.5 show the distribution of R as a function of the module number in gap, E1 and E2 cells, and crack, E3 and E4 cells, scintillators on the A-side of the detector for different data-taking periods. Distributions in the gap and crack scintillators on the C-side of the detector are shown in Appendix A. The red line represents the fitted average response across all modules, while the yellow band represents the systematic uncertainty attributed to the non-uniformity across the modules.

Tables 5.4 and 5.4 summarise the obtained fit parameters for the gap and crack scintillators on both the A-side and C-side of the detector. The response deviations and limits on non-uniformity over ϕ are also reported, demonstrating a good uniformity over ϕ and symmetry between the A and C sides of the detector.

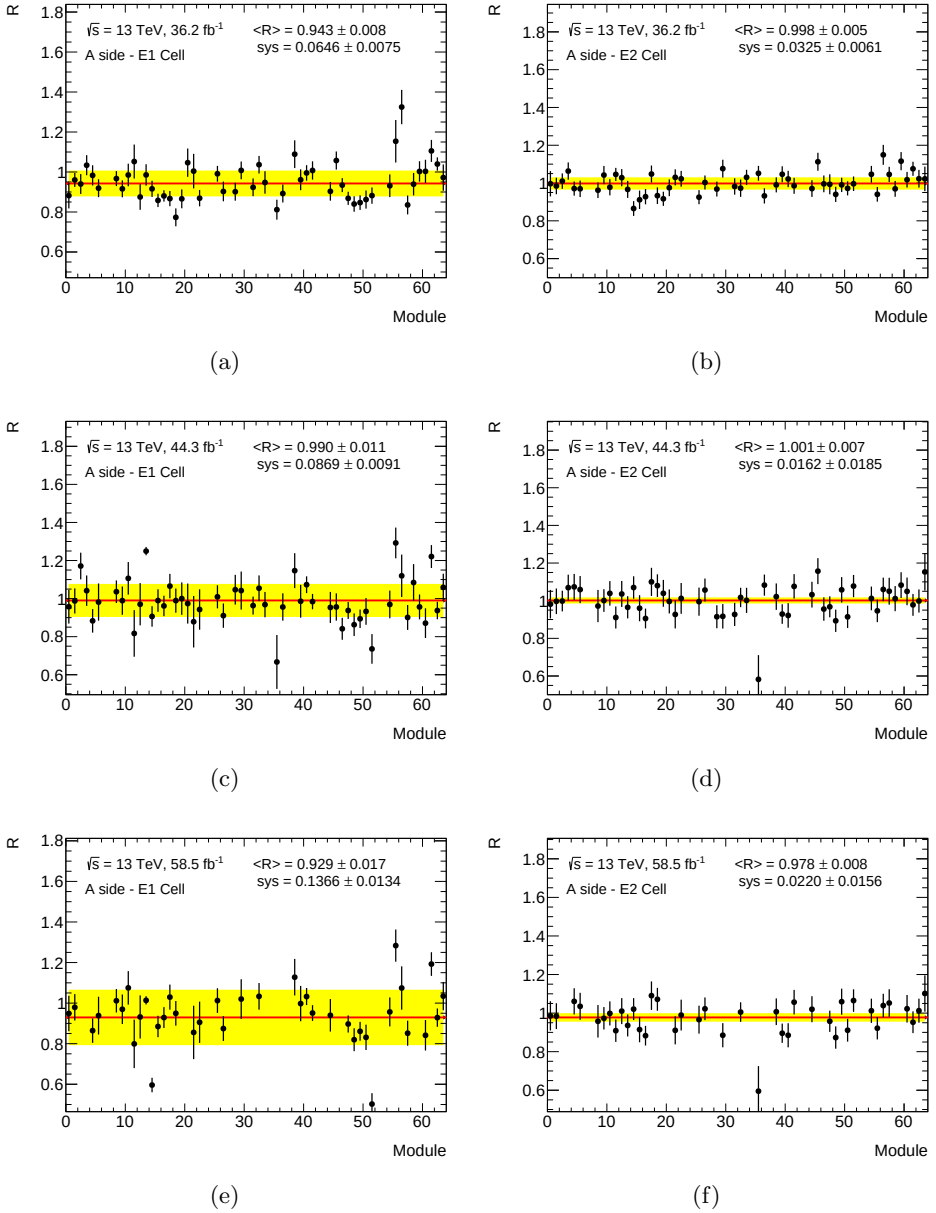


Figure 5.4 Distributions of module number vs the average response in E1 cells on the C side. The plots show results for 2015+2016, 2017, and 2018 data-taking periods. The top, middle, and bottom rows represent the respective years.

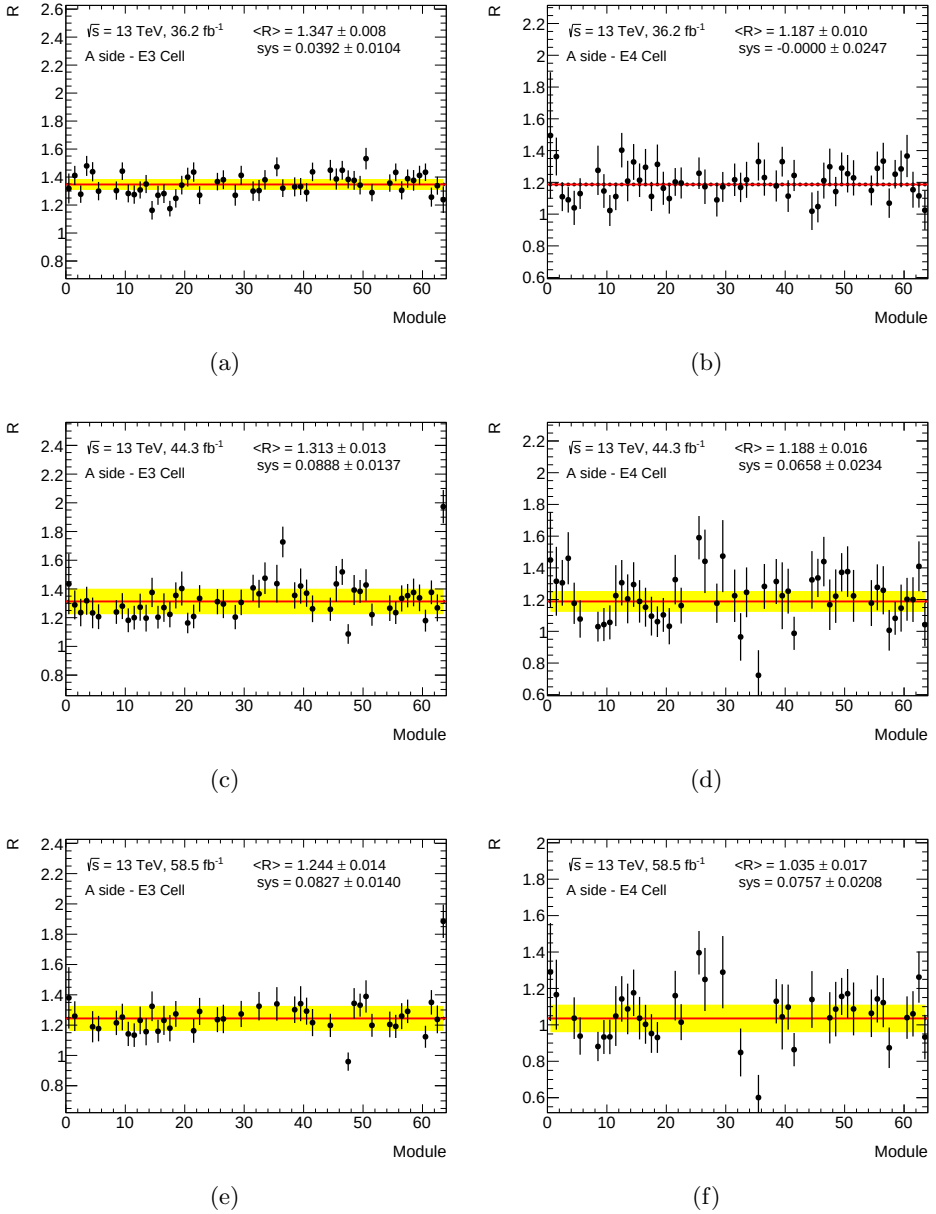


Figure 5.5 Distributions of module number vs the average response in E1 cells on the C side. The plots show results for 2015+2016, 2017, and 2018 data-taking periods. The top, middle, and bottom rows represent the respective years.

Table 5.4 Fit parameters μ and s obtained for E1, E2, E3 and E4 cells on the A-side of the detector during 2015+2016, 2017 and 2018 data-taking periods.

	2015+2015		2017		2018	
	$\hat{\mu}$	$\hat{s}$	$\hat{\mu}$	$\hat{s}$	$\hat{\mu}$	$\hat{s}$
E1	0.943 ± 0.008	0.065 ± 0.008	0.990 ± 0.011	0.087 ± 0.009	0.929 ± 0.017	0.137 ± 0.013
E2	0.982 ± 0.010	0.000 ± 0.022	1.001 ± 0.007	0.016 ± 0.019	0.978 ± 0.008	0.022 ± 0.0156
E3	1.347 ± 0.008	0.0391 ± 0.0104	1.313 ± 0.013	0.0889 ± 0.0137	1.244 ± 0.014	0.0827 ± 0.014
E4	1.187 ± 0.010	0.00 ± 0.0247	1.188 ± 0.016	0.0658 ± 0.0234	1.035 ± 0.017	0.0757 ± 0.0208

Table 5.5 Fit parameters μ and s obtained for E1, E2, E3 and E4 cells on the C-side of the detector during 2015+2016, 2017 and 2018 data-taking periods.

	2015+2015		2017		2018	
	$\hat{\mu}$	$\hat{s}$	$\hat{\mu}$	$\hat{s}$	$\hat{\mu}$	$\hat{s}$
E1	0.982 ± 0.013	0.123 ± 0.011	1.005 ± 0.015	0.127 ± 0.012	0.941 ± 0.021	0.171 ± 0.016
E2	0.951 ± 0.006	0.042 ± 0.006	0.983 ± 0.008	0.052 ± 0.008	0.959 ± 0.009	0.061 ± 0.009
E3	1.313 ± 0.008	0.053 ± 0.009	1.251 ± 0.009	0.020 ± 0.025	1.189 ± 0.010	0.030 ± 0.019
E4	1.207 ± 0.030	0.068 ± 0.036	1.121 ± 0.016	0.084 ± 0.018	0.982 ± 0.016	0.086 ± 0.018

The E1 cells show a significant drop in the response during the 2018 data-taking campaign. This drop is observed on both sides of the detector. The observed drop is attributed to degradation of the optical transmission of the scintillator tiles due to prolonged exposure to intense radiation. The maximum deviation of 2% in the response in E2 cells is observed between 2017 and 2018 data-taking periods. The deviation is not as significant as the one observed in the E1 cells. E3 cells show a drop in the response over time. The upper limit on the average non-uniformity in ϕ , quantified by the spread in response amongst modules of a given cell type, is found to be 0.171 in the E1 cells on the C-side of the detector during 2018 data-taking period.

Each module m of the considered calorimeter cell type c was tested for azimuthal intercalibration by requiring a deviation $|\Delta_{m,c}|$ to be less than three standard deviations from the fitted parameter $\hat{\mu}_C$ obtained from the likelihood

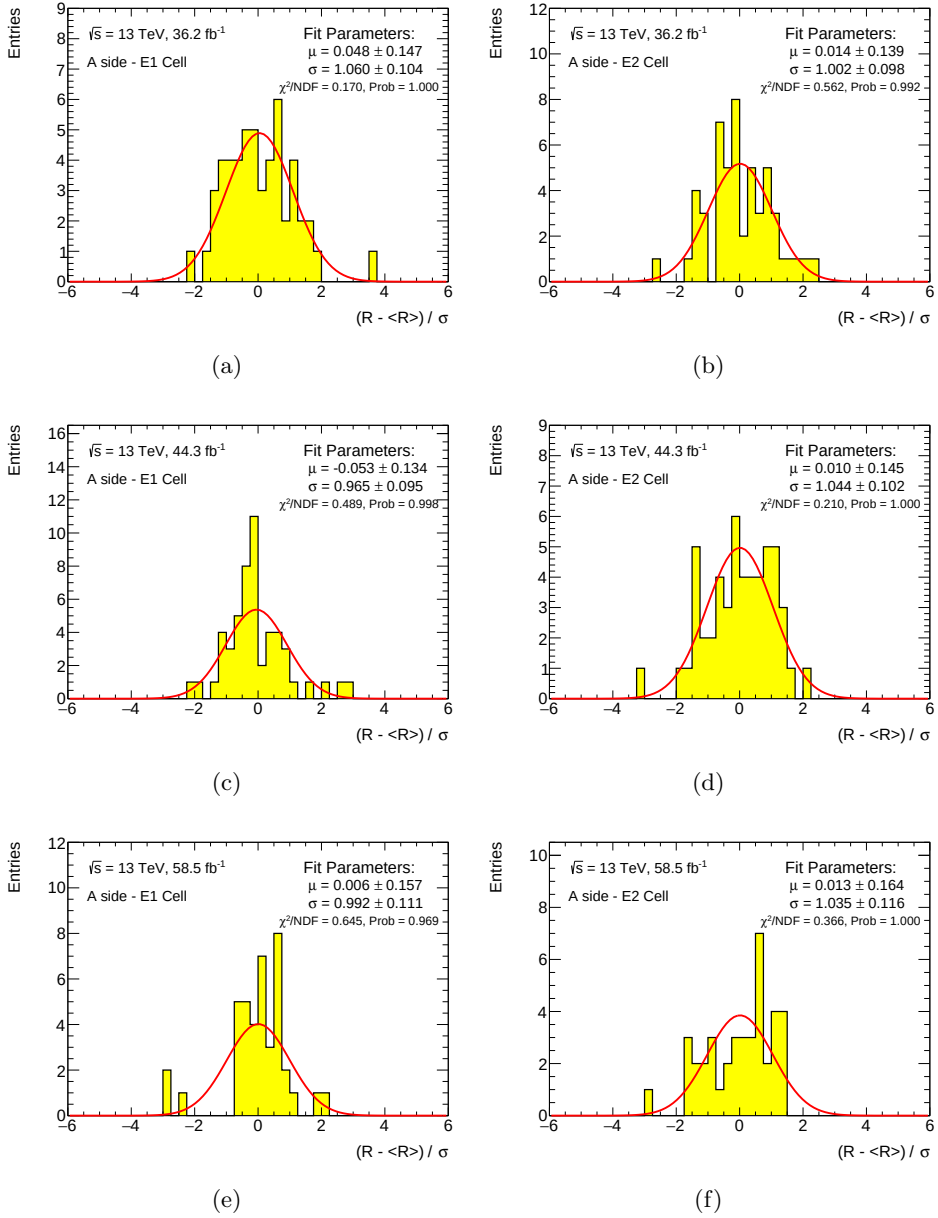


Figure 5.6 Pull distributions obtained using the fitted parameters $\hat{\mu}$ and $\hat{\sigma}$ for gap scintillators. The plots show results for 2015+2016 (top), 2017 (middle) and 2018 (bottom) data-taking periods.

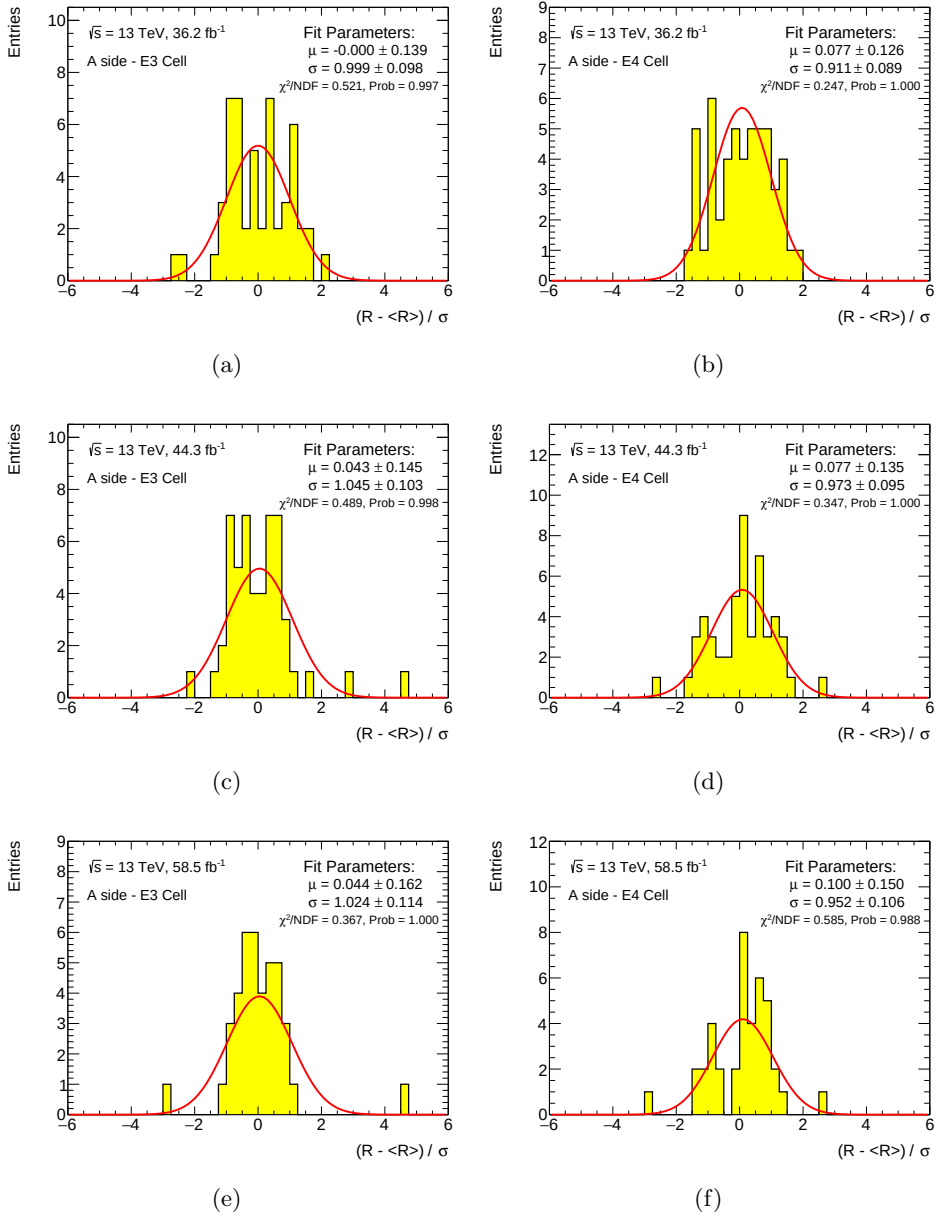


Figure 5.7 Pull distributions obtained using the fitted parameters $\hat{\mu}$ and $\hat{\sigma}$ for crack scintillators. The plots show results for 2015+2016 (top), 2017 (middle) and 2018 (bottom) data-taking periods.

fit. The deviation is defined as:

$$\Delta_{m,c} = \frac{R_{c,m} - \langle R \rangle}{\sqrt{\sigma_{c,m}^2 + s_c^2}}, \quad (5.6)$$

where $\langle R \rangle$ is obtained through likelihood fits. By construction, it is expected that the pull will be Gaussian distributed with a mean of zero and a unit width.

Figures 5.6 and 5.7 show the pull distributions used to check for azimuthal intercalibration in gap and crack scintillators, on the A-side of the detector. Distribution on the C-side of the detector are shown in Appendix A. The total number of modules in gap and crack scintillators observed to deviate more than three standard deviations from the fitted parameter $\hat{\mu}_C$ was found to be 4 in 2015 and 2016, 3 in 2017 and only 2 in 2018 data-taking period.

Chapter 6

Reconstruction of Physics Objects

When proton-proton collisions occur at the LHC, the constituent particles of protons, quarks and gluons, interact, leading to two main types of collisions: elastic and inelastic. Elastic collisions leave the incident proton energies unchanged and produce no new particles, while inelastic collisions involve changes in momentum of the incident protons, leading to the creation of new particles. This phenomenon is fundamentally explained by the principles of QCD and electroweak interactions [92, 155]. In the context of particle physics, inelastic collisions are of particular interest due to the creation of new particles.

6.1 Inelastic Collisions

Inelastic collisions can be categorised into two main classes: diffractive and non-diffractive collisions. According to the Regge theory of strong interactions, the diffractive dissociation is thought to be a result of colour neutral interactions mediated by the hypothetical Pomeron [156].

The dominant diffractive dissociation modes include the single and double diffractive processes. In single diffractive processes, one of the incident protons comes out unscathed after the collision while the other proton dissociates and then hadronise to form a system of particles. The double diffractive processes occur when both the incident protons dissociate and hadronise, forming products that generally exit the system at relatively small angles. The diffractive

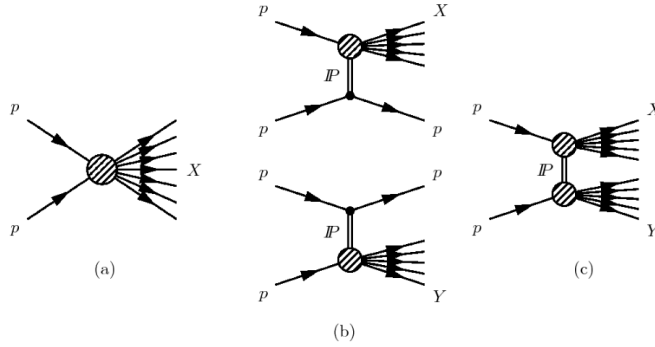


Figure 6.1 Schematic diagram of (a) non-diffractive, (b) a single diffractive process where one proton hadronises after an exchange of the hypothetical Pomeron and (c) double diffractive process. $X(Y)$ represents a hadronised system of particles [158].

processes in detectors are characterised by the absence of hadronic activity in certain η regions of the final state phase space [157]. All diffractive processes involving minimal momentum transfer, typically a few hundred MeV^2 , along with elastic scattering, are classified as soft processes.

The second class of inelastic scattering processes is the non-diffractive dissociation where partons from the two colliding protons interact directly, leading to creation of new particles. Unlike diffractive interactions, non-diffractive processes result in created particles being distributed evenly throughout the detector volume, leaving no η gaps. The non-diffractive dissociations can be further classified as either soft processes, involving minimal momentum transfer, or hard processes, characterised by momentum transfers on the order of one GeV^2 [157].

The total cross-section of pp collisions is given as:

$$\sigma_{\text{tot}} = \sigma_{\text{elastic}} + \sigma_{\text{inelastic}} = \sigma_{\text{elastic}} + \sigma_{\text{diffractive}} + \sigma_{\text{non-diffractive}}, \quad (6.1)$$

where σ_{elastic} represents the cross-section of elastic collisions, and the cross-section of all inelastic collisions is represented by $\sigma_{\text{inelastic}}$.

6.2 Reconstruction of Hard Scatter Objects

In experimental high-energy particle physics, the main focus is often directed towards studying products generated by hard scattering process. This process is characterised by substantial momentum transfers between the colliding protons, resulting in the production of new particles. These newly formed particles are subsequently reconstructed from their energy deposits in the calorimeter systems and the tracks they leave behind in tracking detectors.

The reconstructed hard scatter objects need to be isolated from background activity originating from initial and final state radiation, hadrons produced by beam-beam remnants, and multiple parton interactions [159]. The presence of pileup events also affects the reconstruction of hard scatter objects. Subsequent sections will delve into reconstruction of various physics objects used in analyses.

6.3 Reconstruction of Photons and Electrons

When photons and electrons interact with the material of the electromagnetic calorimeter, they create electromagnetic showers that deposit significant energy in the calorimeter cells. Electrons and photons have similar energy deposits in the ECal, as such, their reconstruction is performed in parallel.

6.3.1 Energy Cluster Reconstruction

The process of reconstructing electron and photon candidates begins by using clustering algorithms designed to group calorimeter cells in both the barrel region and end-caps of the detector. The energy deposits in the ECal cells are reconstructed using a sliding-window algorithm, aggregating energy deposits within a fixed-rectangular window [160]. The resultant topological seed clusters are required to have $p_T > 2.5$ GeV and dimensions of $\Delta\eta \times \Delta\phi = 0.075 \times 0.123$, with their position adjusted to contain local max-

imum of the reconstructed transverse energy, E_T [161]. A duplicate removal algorithm is then applied to eliminate overlapping energy clusters.

The reconstructed topological seed clusters are then matched with tracks in the ID to determine whether they belong to an electron or a photon candidate [162]. Two types photon candidates can be reconstructed, converted and unconverted photons. Converted photons are prompt photons that were converted to electron-positron pairs, while unconverted photons remain unchanged. The classification of a candidate as either an electron or an unconverted photon depends on whether it has matching tracks in the ID.

6.3.2 Track Reconstruction

The trajectory of a particle with electric charge q passing through a magnetic field can be described by five helix parameters: $d_0, z_0, \phi, \cot \theta$ and q/p_T , illustrated in Figure 6.2. The parameter $\cot \theta$ gives the inverse of the slope of the track and q/p_T provides information about the particle's charge relative to its momentum in the transverse plane. Track fitting algorithms perform fits that best estimate these track parameters. Details of various tracking algorithms used in ATLAS are discussed in detail in Reference [162].

The reconstruction of tracks is a two step process that begins by finding tracks from hits in the tracking detector using a pattern recognition algorithm and then performing a χ^2 fit to select good tracks. During the track reconstruction process, tracks with more points are given priority over those with fewer points. Tracks matched to seed clusters are further refined using the Gaussian Sum Filter [164], which improves the electron track parameters estimations, especially in cases involving significant energy loss due to Bremsstrahlung radiation. Finally the tracks are matched with calorimeter seed clusters that passed the shower shape requirements to complete the reconstruction process [165].

It is possible for one seed cluster to be matched to more than one track, in such cases a matching algorithm is employed. The algorithm considers the distance

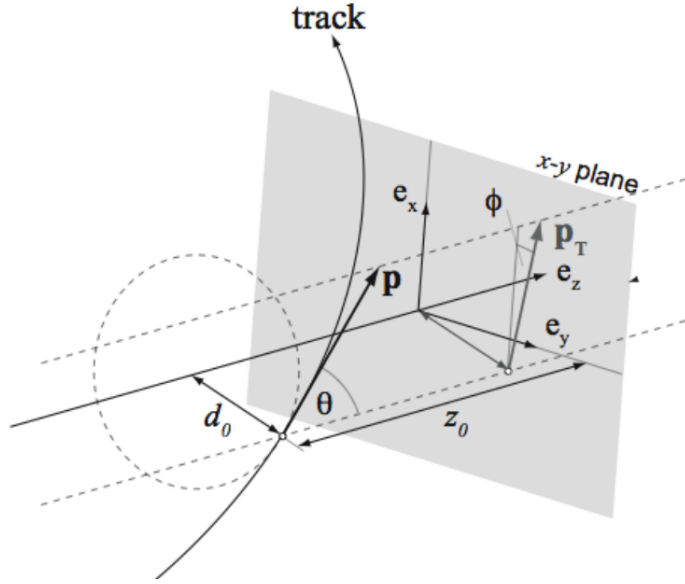


Figure 6.2 Track parameters that describe the path of a charged particle in a magnetic field. The parameter d_0 is the transverse impact parameter, z_0 , represents the longitudinal impact parameter, ϕ_0 denotes the azimuthal angle in the x - y plane, θ gives the polar angle and $\mathbf{p}$ is the particle's momentum [163].

between the track and the barycentre of the seed cluster and the number of hits in the silicon detector. An object with a seed cluster matched to a single track, with a minimum of four hits in the pixel detector, is classified as an electron candidate [166]. When two reconstructed tracks form a vertex consistent with that of a decaying massless particle, the seed clusters are classified as converted photon candidates. Converted photon candidates are reconstructed from seed clusters matched with tracks that have no hits in the IBL.

The final seed clusters are then rebuilt using new cluster sizes determined by the particle hypothesis, whether it is an electron or photon candidate, and the calorimeter region. The newly built clusters are then calibrated to account for energy leakages behind the ECal and energy losses in the upstream material [167]. In the barrel region, the final seed clusters are sized 3×5 for unconverted photons, and 3×7 for both electrons and converted photons.

The final seed for both electron and photon candidates are sized 5×5 in the end-caps.

6.3.3 Photon Identification

Seed clusters of prompt-photon candidates need to be isolated from seed clusters of background photons originating from decays of neutral hadrons or energy deposits of jets in the ECal. Algorithms used in photon identification use calorimetric variables to separate background photons from prompt-photon candidates. ATLAS employs an algorithm that provides two levels of photon identification, the Loose and Tight working points [165]. These working points are defined in such a way that one is a subset of the other. The Loose working point is obtained by using information about shower shapes and energy deposits in the hadronic calorimeters to separate prompt-photon candidates from background photons. The Tight working point tightens selection requirements on the Loose selection variables and also adds information from the segmented strip layer. Converted photon candidates generally have broad shower shapes in the calorimeter cells, so the optimisation of converted and unconverted photons is performed separately.

6.3.4 Electron Identification

The electron identification algorithm, similar to the photon identification algorithm, employs calorimetric variables to distinguish prompt electron candidates from potential background electrons. The algorithm uses hit counts in the IBL to distinguish electrons from converted photon candidates. The algorithm provides three working points for electron identification: Loose, Medium and Tight working points [168]. The Loose working point uses information from the first two layers of the ECal and TileCal to distinguish prompt-electrons from background electrons. The Medium working points adds tighter selection cuts to the Loose working point selections by using information from the TRT and the third layer of the ECal. The Tight working

point incorporates additional information used in matching tracks with seed clusters and further tightens the selection criteria from the Medium working point. For electrons with $E_T > 125$ GeV, electrons have similar signatures to background energy deposits, the Tight and Medium working points are similar [166].

6.3.5 Isolation of Electrons and Photons

Isolation requirements are crucial in distinguishing prompt objects from background energy deposits. A traditional method used in ATLAS to isolate photons and electrons utilises a cone with a radius, $\Delta R = \sqrt{\Delta\eta^2 + \Delta\phi^2}$, drawn around the direction of the photon/electron candidates [169]. Isolation variables quantify the amount of activity around the candidate object, allowing for the discrimination of prompt objects from other background energy deposits.

The calorimeter isolation energy, E_T^{iso} , requirement is obtained by summing of E_T of the topological seed clusters within a cone of size ΔR to compute the raw isolation energy variable $E_{T,\text{raw}}^{\text{iso}}$ [170]. The deposited energy of electron/photon candidates is quantified by $E_{T,\text{core}}$, illustrated as a yellow rectangle in Figure 6.3. The energy leaking into the cone is quantified by $E_{T,\text{leakage}}$, and the energy contribution from pileup events is given by $E_{T,\text{pileup}}$ [166]. The calorimeter based isolation variable $E_{T,\text{core}}^{\text{iso}}$ is then given by:

$$E_{T,\text{cone}}^{\text{isol}} = E_{T,\text{raw}}^{\text{isol}} - E_{T,\text{core}} - E_{T,\text{leakage}} - E_{T,\text{pileup}}. \quad (6.2)$$

The track based isolation, p_T^{iso} , requirement is obtained by summing the p_T of all tracks, excluding the p_T of tracks matched to candidate electrons or converted photons. The considered tracks must to have $E_T > 1$ GeV and be placed $|z_0 \sin \theta| < 3$ mm from the reconstructed primary vertex. The requirement on $|z_0 \sin \theta|$ aims to select tracks originating from a primary vertex deemed to be relevant for the process. To isolate electron candidates from other decay products that are very close to the direction of the electron candidate, the ΔR used is defined to vary in size as a function of the electron's

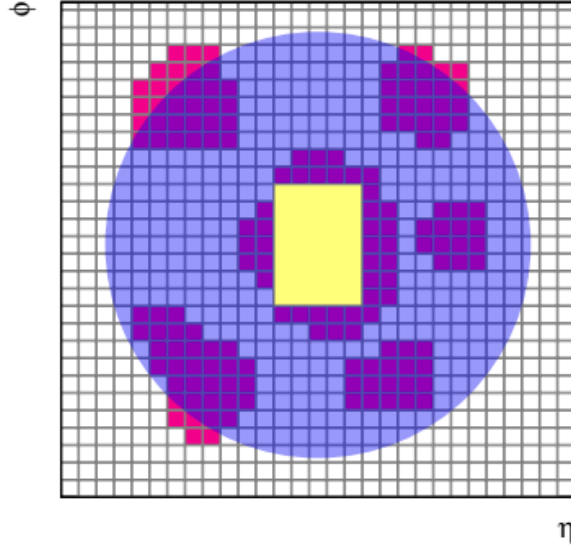


Figure 6.3 A schematic illustrating the calorimeter based isolation requirement with the grid representing the calorimeter cells in the η - ϕ direction. The red energy blobs represent topological clusters, with the one in the middle being the candidate object. The circle around the topological clusters represents a cone of radius ΔR [166].

p_T . The p_T dependent cone size is given by:

$$\Delta R = \min \left(\frac{10 \text{ GeV}}{p_T [\text{GeV}]}, R_{\max} \right), \quad (6.3)$$

where $R_{\max}$ is the maximum cone size, typically $R_{\max} = 0.2$. The value 10 GeV in the equation represents a threshold that ensures the isolation cone size decreases for higher momentum electrons, providing better isolation for high- p_T electrons while maintaining sufficient separation from other particles for lower- p_T electrons.

The choice of isolation requirements depends on the physics analysis since different physics processes have different background sources. The Loose and Tight isolation requirements are used in studies documented in this thesis. The Loose isolation requirement takes into consideration E_T^{iso} and p_T^{iso} of the

photon candidate computed within a cone of $\Delta R = 0.2$. It is defined as:

$$E_T^{\text{iso}}|_{\Delta R < 0.2} < 0.065 \cdot E_T \text{ GeV and } p_T^{\text{iso}}|_{\Delta R < 0.2} < 0.05 \cdot E_T. \quad (6.4)$$

The Tight isolation requirement is determined based on E_T^{iso} computed within a cone of $\Delta R = 0.4$ and p_T^{iso} computed within a cone of $\Delta R = 0.2$, and it is defined as:

$$E_T^{\text{iso}}|_{\Delta R < 0.4} < 0.022 \cdot E_T + 2.45 \text{ GeV and } p_T^{\text{iso}}|_{\Delta R < 0.2} < 0.05 \cdot E_T. \quad (6.5)$$

The electron isolation working points are defined in two ways: those with fixed cuts on the E_T^{iso} and p_T^{iso} isolation variables or those targeting to obtain a fixed value of efficiency by varying cuts on the E_T^{iso} and p_T^{iso} isolation variables. The tag-and-probe method detailed in Reference [166] is used to measure the electron isolation efficiencies. The isolation working points designed to obtain a targeted isolation efficiency used in ATLAS include the LooseTrackOnly, Loose and Tight working points, each designed to obtain 99%, 98% and 95% isolation efficiencies, respectively.

The fixed cut isolation working points include the FixCutLoose and FixCutTight working points. The FixCutLoose working point is defined as:

$$E_T^{\text{iso}}|_{\Delta R < 0.2} < 0.20 \cdot E_T \text{ GeV and } p_T^{\text{iso}}|_{R_{\text{max}}=0.2} < 0.05 \cdot E_T, \quad (6.6)$$

and the FixCutTight working point is defined as:

$$E_T^{\text{iso}}|_{\Delta R < 0.2} < 0.06 \cdot E_T \text{ GeV and } p_T^{\text{iso}}|_{R_{\text{max}}=0.2} < 0.05 \cdot E_T. \quad (6.7)$$

The isolation working points are defined using a cone of size $\Delta R = 0.2$ for E_T^{iso} and $R_{\text{max}} = 0.2$ for p_T^{iso} .

6.4 Muon Reconstruction

Muon candidates are reconstructed by combining tracks from the ID and MS, along with energy deposits in the calorimeters indicative of minimum ionising particles. The reconstruction process for muon candidates in the MS and ID is carried out independently. In the ID, muon candidates are reconstructed in a manner similar to electron candidates. In the MS, the reconstruction process begins by identifying hit patterns within the muon chambers to form segments. Candidate muon tracks are then constructed by globally fitting hits from different layers of the MS [171]. Finally, muon candidates are reconstructed using information provided by the ID, MS and calorimeter systems.

6.4.1 Muon Identification

Four types of muons are reconstructed for use in physics analyses, depending on which part of the detector was utilised in their reconstruction.

- Combined muons are reconstructed by performing a global χ^2 fit on reconstructed muon tracks from both the ID and MS. These muons serve as the default type used in most analyses.
- Stand-alone muons are reconstructed solely from tracks found in the MS, primarily in the region $2.5 < |\eta| < 2.7$, where there is no coverage by the ID. These muons may also originate from high- p_T objects decaying in-flight [172].
- Segment-tagged muons are reconstructed using tracks from the ID along with a matching segment in the MDTs or cathode strip chambers. These type of muons are reconstructed in regions of poor detector coverage. They are typically low- p_T muons that do not traverse the entire detector.
- Calorimeter-tagged muons are reconstructed utilising ID tracks and signatures of minimum ionising particle energy deposits in the calorimeters.

There are cases where different reconstructed muons share a track. To resolve the overlap, priority is given to combined muons, followed by segment-tagged

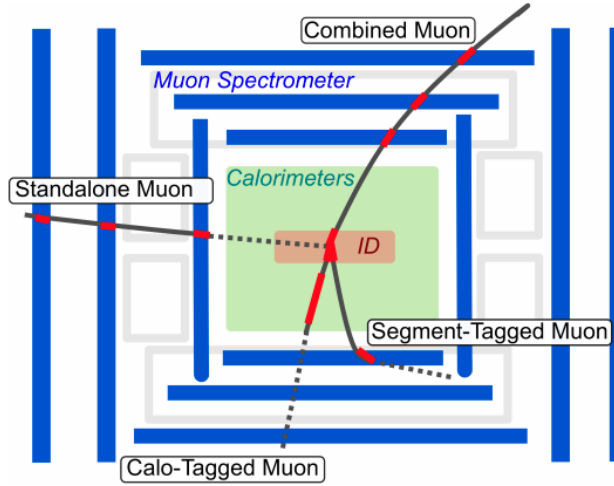


Figure 6.4 A schematic diagram illustrating the various detector components utilised in muon reconstruction in ATLAS. Different types of muons are reconstructed depending on which sub-detectors are employed for reconstruction [173].

muons, and finally by calorimeter-tagged muons. This thesis utilises combined and stand-alone muons.

The reconstructed muon candidates need to be distinguished from background muons produced by in-flight decays of kaons and pions. Tracks of kaons and pions have a characteristic “kink” in their reconstructed topology, leading to a poor global χ^2 fit quality [171], aiding in distinguishing prompt and background muons. Four identification working points are used to separate prompt-muons from background muons:

- The Loose identification criterion is designed to maximise the reconstruction efficiency of muons, it restricts calorimeter-tagged and stand-alone muons to $|\eta| < 1.0$, where the reconstruction of combined muons is not most efficient. For regions $2.5 < |\eta| < 2.7$ stand-alone muons are used.
- The Medium identification criterion serves as the default identification working point for muon identification in ATLAS. It uses combined and

stand-alone muons to minimise systematic uncertainties. The combined muons are required to have a minimum of 3 hits in at least 2 layers of the MDTs or the cathode strip chambers, while stand-alone muons are required to have at least 3 hits across all layers of the MDTs or the cathode strip chambers.

- The Tight identification criterion is developed by using only the combined muons with additional cuts on the normalised χ^2 fit, aiming to reduce fake muons.
- The High- p_T identification criterion is developed to improve the resolution of the reconstructed muons with $p_T > 100$ GeV. The combined muons used in this criterion are required to have a minimum of 3 hits in the MDTs or cathode strip chambers, with regions of the MS having less active material being vetoed.

The calorimeter-based isolation variable, $E_T^{\text{topocone20}}$, calculates the sum of topological clusters within a cone of radius $\Delta R = 0.2$, while excluding clusters within $\Delta R < 0.1$ to account for potential energy deposits from the candidate muon.

6.4.2 Muon Isolation

Muon isolation is performed by quantifying the activity of the detector around the muon candidates. The muon isolation working points are defined using calorimeter-based and track-based isolation variables. The calorimeter-based isolation variable, $E_T^{\text{topocone20}}$, calculates the sum of topological clusters within a cone of radius $\Delta R = 0.2$, excluding clusters within $\Delta R < 0.1$ to account for potential energy deposits from the candidate muon. The track-based isolation variable, $p_T^{\text{varconetr30}}$, sums tracks with $p_T > 1$ GeV, excluding the muon candidate's track in a cone of size ΔR . It is performed similar to the isolation of electrons where the cone size ΔR is p_T dependent to improve the identification of muons from boosted particle decays. The p_T dependent radius of the cone

used in track-based isolation is given by:

$$\Delta R = \min\left(\frac{10 \text{ GeV}}{p_T^\mu}, 0.3\right), \quad (6.8)$$

with p_T^μ being the p_T of the muon and 0.3 being the maximum ΔR .

The muon isolation working points, similar to electrons', are designed by requiring fixed cuts on the calorimeter-based and track-based isolation variables or by targeting to obtain a fixed value of efficiency. The LooseTrackOnly isolation working point uses $p_T^{\text{varconetr30}}/p_T^\mu$ as discriminating variable to obtain a 99% efficiency. The other isolation working points that aim to obtain a fixed efficiency include the Loose and Tight working points which use $E_T^{\text{topocone20}}/p_T^\mu$ and $p_T^{\text{varconetr30}}/p_T^\mu$ as discriminating variables to obtain efficiency of 99% and 96%, respectively.

The fixed cut isolation working points include the FixCutLoose, FixCutTightTrackOnly and FixedCutTight points. The FixCutLoose working point is defined as:

$$E_T^{\text{topocone20}} < 0.30 \text{ and } p_T^{\text{varconetr30}}/p_T^\mu < 0.15, \quad (6.9)$$

the FixCutTightTrackOnly working point is defined as:

$$p_T^{\text{varconetr30}}/p_T^\mu < 0.06, \quad (6.10)$$

and the FixedCutTight working point is defined as:

$$E_T^{\text{topocone20}} < 0.20 \text{ and } p_T^{\text{varconetr30}}/p_T^\mu < 0.05. \quad (6.11)$$

The isolation working points are defined using a cone of size $\Delta R = 0.2$ for $E_T^{\text{topocone20}}$ and $R_{\text{max}} = 0.3$ for $p_T^{\text{varconetr30}}/p_T^\mu$.

6.5 Jets Reconstruction

Hard scattering processes generate quarks and gluons, which will eventually hadronise, forming collimated bunches of hadrons flying through the detector in the same direction. These bunches of collimated streams of hadrons, called jets, serve as experimental signatures of quarks and gluons.

6.5.1 Jet Reconstruction Algorithms

Various methods are employed to reconstruct jets within particle detectors. This section discusses two commonly used approaches: the cone algorithm and the clustering algorithm.

The process of reconstructing jets using cone algorithms begins with the formation of topological clusters from calorimeter cells. These clusters group together cells with substantial energy deposits to suppress noise from electronics and soft interactions. A detailed description on the formation of these topological clusters can be found in Reference [170]. The cone algorithms prioritise cells based on their energy-to-noise ratio, defined as $E_{\text{cell}}/\sigma_{\text{cell}}^{\text{noise}}$. These algorithms identify stable cones centred on seed clusters, overcoming challenges related to infrared and collinear safety through a split-merge procedure [174].

Clustering algorithms operate through a sequential recombination of pairs of clusters to construct a jet. This process involves evaluating the distances, d_{ij} , between objects i and j , as well as d_{iB} , the distance between object i and the beam B .

Clustering algorithms, on the other hand, use a sequential recombination of pairs of clusters to construct a jet. This process involves evaluating the distance, d_{ij} , between objects i and j , as well as d_{iB} , the distance between object i and the beam B . The algorithms work by comparing d_{ij} with d_{iB} , if d_{ij} is less than d_{iB} objects i and j are combined to form a pseudo-jet. This pseudo-jet is then compared to another object j , if the distance exceeds d_{iB} , the pseudo-jet is classified as a jet. This process continues until no objects are left [175]. The

expression for d_{ij} is defined as:

$$d_{ij} = \min(p_{T,i}^{2p}, p_{T,j}^{2p}) \frac{\Delta_{ij}^2}{R^2}; \quad d_{iB} = p_{T,i}^{2p}, \quad (6.12)$$

where $\Delta_{ij}^2 = (y_i - y_j)^2 + (\phi_i - \phi_j)^2$ gives the spacial separation between objects i and j , $p_{T,i}$ is the p_T , y_i is the rapidity, and ϕ_i is the azimuth of object i . The radius parameter R and p give a relation between the energy and geometrical scaling. If $p = 1$:

$$d_{ij} = \min(p_{T,i}^2, p_{T,j}^2) \frac{\Delta_{ij}^2}{R^2}; \quad d_{iB} = p_{T,i}^2, \quad (6.13)$$

the algorithm is called the k_t algorithm, its distance measures the approximate inversion of QCD branching process. if $p = 0$:

$$d_{ij} = \min(p_{T,i}^0, p_{T,j}^0) \frac{\Delta_{ij}^2}{R^2} = \frac{\Delta_{ij}^2}{R^2}; \quad d_{iB} = 1, \quad (6.14)$$

the algorithm that uses this parameter is known as the Cambridge/Aachen algorithm, it only uses the spatial separation, Δ_{ij}^2 , of the two objects when building jets, it does not take their p_T into consideration. It clusters the closest radiation first. If $p = -1$:

$$d_{ij} = \min\left(\frac{1}{p_{T,i}^2}, \frac{1}{p_{T,j}^2}\right) \frac{\Delta_{ij}^2}{R^2}; \quad d_{iB} = \frac{1}{p_{T,i}^2}, \quad (6.15)$$

the resulting algorithm is called the anti- k_t , it is the default algorithm used in ATLAS, with $R = 0.4$ or 0.6 . High p_T particles result in small spatial separation between the two objects used to build a jet, resulting in soft particles clustering around the hard particle before they can cluster amongst themselves. This produces a perfect canonical jet, i.e. the anti- k_t algorithm is a perfect cone algorithm.

High p_T jets produced from the hard scattering process must be distinguished from background jets. Sources of background jets include calorimeter noise and non-collision events such as cosmic ray muons overlapping in-time with

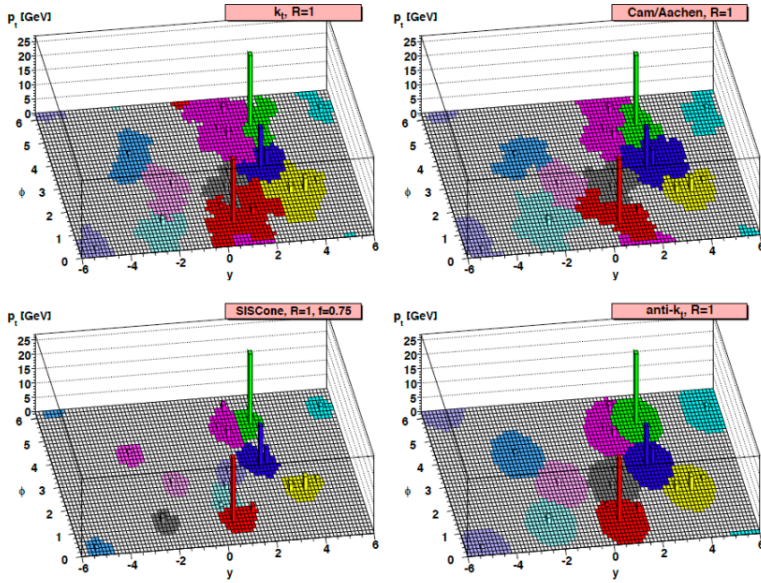


Figure 6.5 Illustration of active catchment areas parton-level events generated with Herwig with random soft events of jets reconstructed with k_t , Cambridge/Aachen, SiSCone and anti- k_t algorithm [176].

collision events and protons that form a halo around the beam core [177]. ATLAS uses the Loose and Tight jet selection requirements to effectively reject jets originating from background processes, while maintaining a high efficiency for jets produced in hard scattering collisions [178]. The Tight jet selection requirement is defined to be used in physics analyses that are sensitive to non-collision background jets.

The efficiency of jet selection requirements is evaluated using a tag-and-probe method detailed in Reference [179]. The Loose jet selection requirement ensures that over 99.5% of jets originating from hard scattering collisions with $p_T > 20$ GeV are retained, while jets with $p_T > 100$ GeV are selected with an efficiency exceeding 99.9%. The Tight jet selection requirement ensures the selection of over 95% of jets with $p_T > 20$ GeV originating from the hard scattering process, the selection efficiency further increases to above 99.5% for jets with $p_T > 100$ GeV.

6.6 Missing Transverse Energy Reconstruction

Missing transverse energy, denoted as E_T^{miss} , mainly originates from particles that do not interact with the detector material. The E_T^{miss} signatures can also arise from factors such as limited coverage of the detector material or dead regions within the detector. ATLAS uses energy deposits in calorimeters and muon tracks to reconstruct E_T^{miss} . It is reconstructed as the net imbalance of momentum in the transverse plane, computed by taking the negative vector sum of the transverse momentum of all detected objects:

$$\mathbf{E}_{x(y)}^{\text{miss}} = - \left(\sum_{i \in \{\text{hard objects}\}} \mathbf{p}_{x(y)}^i + \sum_{i \in \{\text{soft events}\}} \mathbf{p}_{x(y)}^i \right). \quad (6.16)$$

The objects used in the reconstruction of $\mathbf{E}_T^{\text{miss}}$ include both calibrated hard and soft events. The reconstruction sequence begins with electrons, followed by photons, hadronically decaying τ leptons, and then jets. Lower priority reconstructed objects are fully rejected if they share calorimeter signals with a higher priority objects [180].

The magnitude of $\mathbf{E}_T^{\text{miss}}$ and ϕ^{miss} , are calculated as follows:

$$E_T^{\text{miss}} = |\mathbf{E}_{x(y)}^{\text{miss}}| = \sqrt{(E_x^{\text{miss}})^2 + (E_y^{\text{miss}})^2}, \quad (6.17)$$

$$\phi^{\text{miss}} = \tan^{-1} \theta \left(\frac{E_y^{\text{miss}}}{E_x^{\text{miss}}} \right), \quad (6.18)$$

where ϕ^{miss} gives the direction of E_T^{miss} in the transverse plane.

6.6.1 Object selection for E_T^{miss} reconstruction

The following event selection requirements are applied to reconstructed particles and jets used in the calculation of E_T^{miss} :

- Candidate electrons are required to meet the Medium identification requirement, the Loose isolation criteria, have $p_T > 10$ GeV, and be found

within $|\eta| < 1.37$ or $1.52 < |\eta| < 2.47$ to exclude events in the transition region [181].

- Photon candidates are required to meet the Tight identification requirement, have $p_T > 25$ GeV, and be found in $|\eta| < 1.37$ or $1.52 < |\eta| < 2.47$.
- Hadronically decaying τ leptons appear in the detector as narrow jets with low track multiplicity, the jets used in the reconstruction must be found within $|\eta| < 1.37$ or $1.52 < |\eta| < 2.47$ with $p_T > 25$ GeV. The jet candidates must pass the Medium quality selection [182].
- In the $|\eta| < 2.5$ region, reconstructed combined muon candidates are used in the E_T^{miss} calculation. Stand-alone muon candidates are used in the $2.5 < |\eta| < 2.7$ region. These muon candidates must meet the Medium identification quality requirement, have $p_T > 10$ GeV and the combined muon candidates are further required to be associated with a hard scatter vertex.
- The jets considered in the E_T^{miss} calculation are reconstructed by the default anti- k_t algorithm with $R = 0.4$. These jets are reconstructed within $|\eta| < 4.5$ and they are required to have $p_T > 20$ GeV. Electrons and photons emitted in the transition region are reconstructed as jets. The Jet Vertex Tagger is employed to suppress pileup jets with $p_T < 60$ GeV, found in $|\eta| < 2.4$ and not associated with the hard scatter primary vertex [183].
- The soft term consists of calorimeter signals or ID tracks associated with the hard scatter vertex but not linked to any hard objects. The calorimeter-based soft term is constructed from signals from topological clusters in the calorimeter. The track-based soft term employs ID tracks not associated with the reconstructed hard objects, making it insensitive to pileup. The ID tracks used in the track-based soft term reconstruction are required to satisfy the following criteria: $p_T > 400$ MeV, have a transverse impact parameter $|d_0|/\sigma d_0 < 2$, longitudinal impact parameter $|z_0 \sin \theta| < 3$ mm, $\Delta\phi(\text{track}, e/\gamma) > 0.05$ or $\Delta\eta(\text{track}, e/\gamma) > 0.2$ and $\Delta R(\text{track}, \tau) > 0.2$.

Chapter 7

High-mass Resonance Search in the $Z\gamma$ Final State at $\sqrt{s} = 13$ TeV With the ATLAS Detector

The high-mass $Z\gamma$ resonance search is part of a wide series of searches aiming to find evidence of new BSM physics by exploring new particles decaying into boson pairs, as predicted by various extensions of the Standard Model [184–186]. This chapter focuses on the search for a narrow high-mass resonance with either spin-0 or spin-2, in the mass range between 220 GeV and 3400 GeV, using pp collision data collected by the ATLAS detector during the LHC Run 2 data taking.

The high-mass resonance, X , is searched through its decay into a final state comprising of a Z boson and a photon. The Z boson further undergoes a leptonic decay into a pair of oppositely charged electrons or muons. The Feynman diagram in Figure 7.1 illustrates the $X \rightarrow Z(\ell\ell)\gamma$ decay process. The decays of the Z boson into a pair of muons and electrons are selected because they can be efficiently triggered, offer good invariant mass resolution, and have a clean experimental signature in the detector. This cleanliness allows them to be distinguished from background events produced in pp collisions more easily as compared to hadronic, neutrino, and τ lepton decays.

The dominant background events are expected to stem from SM $Z + \gamma$ and $Z + \text{jets}$ events. In the SM, the production of $Z + \gamma$ background events pri-

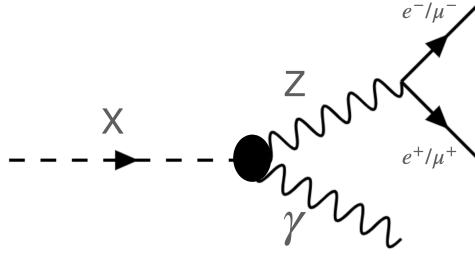


Figure 7.1 Feynman diagram of the $X \rightarrow Z\gamma$ process. The Z boson further decays to electron or muon pairs.

marily occurs due to final-state radiation, FSR, from leptons or initial-state radiation, ISR, from quarks. In such events, the prompt photon candidate is characterised by a narrow energy cluster in the ECal, typically well isolated from hadronic activity. The remaining background is contributed by the production of Z bosons in association with jets, $Z + \text{jet}$, where one jet is mistakenly identified as a photon. The misidentified photon candidate typically originates from the decay of a neutral meson, typically a π^0 , which carries a significant portion of the initial parton's energy. These misidentified photon candidates usually have large energy deposits in the ECal with non-negligible leakage into the TileCal.

Searches for X decaying into the $Z\gamma$ final state have previously been conducted by both the ATLAS and CMS Collaborations. The ATLAS Collaboration performed the search for X decaying into $Z\gamma$ final state, where the Z boson decays leptonically, using data collected from 2015 to 2016, corresponding to an integrated luminosity of 36.1 fb^{-1} [187]. No significant excess above the expected SM background was observed, and upper limits on the production cross-section times the branching ratio, $\sigma(pp \rightarrow X) \cdot \mathcal{B}(X \rightarrow Z\gamma)$, were set. For spin-0 narrow resonances, these limits ranged from 88 fb to 2.8 fb at 95% confidence level, CL, in a mass range between 250 GeV and 2400 GeV. Additionally, the ATLAS Collaboration conducted a search for X decaying into the $Z\gamma$ final state, with the Z boson decaying hadronically [188]. Once again, no significant excess above the expected SM background was observed,

and upper limits on $\sigma(pp \rightarrow X) \cdot \mathcal{B}(X \rightarrow Z\gamma)$ at 95% varied from 10 fb to 0.05 fb in the mass range from 1.0 TeV to 6.8 TeV.

The study performed by the CMS Collaboration used data collected from 2015 to 2016, with an integrated luminosity of 36.1 fb^{-1} [189]. The study combined results from channels where the Z boson decayed leptonically and hadronically. For narrow spin-0 resonances with masses between 0.3 TeV and 4.0 TeV, the upper limits on $\sigma(pp \rightarrow X) \cdot \mathcal{B}(X \rightarrow Z\gamma)$ at 95% ranged from 50 fb to 0.3 fb.

This chapter presents a comprehensive investigation into the search for X decaying into the $Z\gamma$ final state, focusing on cases where the Z boson undergoes leptonic decays. The methodology employed in this study closely mirrors that of previous searches conducted by the ATLAS Collaboration. However, key distinctions include the increase in integrated luminosity, refinements in reconstruction techniques, and an extension in the search mass range.

7.1 Datasets

The dataset used for analysis was obtained from pp collisions at $\sqrt{s} = 13$ TeV, recorded by the ATLAS detector at the LHC from 2015 to 2018. The pp collision events were triggered using the single-lepton, dilepton and single high- p_T photon triggers.

The single-lepton trigger had a minimum threshold of 26 GeV. The threshold of dilepton triggers varied depending on whether the events contained muons and electrons. The dielectron trigger required each electron to have a minimum p_T of at least 17 GeV, whereas the dimuon trigger had asymmetric p_T thresholds of 22 GeV and 8 GeV for the muon pairs. Events that satisfied at least one of these trigger requirements were retained for analysis. Following the application of data quality requirements, the dataset resulted in an integrated luminosity of 140 fb^{-1} [138].

Simulated event samples were employed to optimise the search strategy, assess the event selection efficiency, and study the background processes. Both the data and MC samples used in this analysis correspond to 25 ns bunch

crossing intervals. All MC samples were generated with the MC16 settings and subsequently reweighted to match the distribution of $\langle\mu\rangle$ observed in the experimental data.

7.1.1 Simulated Signal Samples

The simulated signal samples were generated and processed with a detailed software that simulates the ATLAS detector [190] based on GEANT4 [147].

Spin-0 signal samples produced via the gluon-gluon fusion, ggF , were simulated with masses of 200, 300, 500, 1000, 1500, 2000, 2500, 3000, and 3500 GeV. These samples were generated with an intrinsic narrow width of 4 MeV. The events were generated using the POWHEG framework [191–194].

Spin-2 signal samples were produced via the gluon-gluon, ggX , and the $q\bar{q}X$, quark–antiquark, annihilation processes with masses of 200, 250, 300, 500, 750, 1000, 1500, 2000, 2500, 3000, and 3500 GeV. These samples were also generated with an intrinsic narrow width of 4 MeV. For both processes, events were generated at LO in QCD using the Higgs Characterisation Model [195] implemented using MadGraph5_aMC@NLO [196].

For both spin-0 and spin-2 signal processes, the generated events were interfaced with PYTHIA8 [146] for parton showering, hadronisation and multiparton interactions, with the AZNLO set of tuned parameters. The PDF sets employed were the CTEQ6L1 [197] for spin-0 processes and NNPDF2.3 [198] for spin-2 processes.

7.1.2 Simulated Background Samples

The background samples for the non-resonant $Z + \gamma$ process were separately generated for the electron and muon final states. These samples were produced using SHERPA 2.2.2 [145] event generator, using Comix [199] and OpenLoops [200] generators to calculate the matrix elements. The matrix elements were integrated with the SHERPA parton shower [201] using the MEPS@LO

Table 7.1 Selection criteria for physics objects and events used in the analysis.

	Muon	Electron	$e(\gamma)$	Photon
p_T	> 10 GeV	> 10 GeV	> 50 GeV	> 15 GeV
$ \eta $	< 2.7	< 2.47 Exclude $[1.37, 1.52]$	< 2.47 Exclude $[1.37, 1.52]$	< 2.37 Exclude $[1.37, 1.52]$
$ d_0 /\sigma_{d_0}$	< 3	< 5		
$ z_0 \sin \theta $	< 0.5 mm	< 0.5 mm		
Identification	Medium	Mixed	MVA	Tight
Isolation	Track-based Tight	Track-based Tight		Loose
$\Delta R(\text{track}, \gamma)$			< 0.1	
$e^\pm e^\mp$ or $\mu^\pm \mu^\mp$		≥ 2		
$e(\gamma)$			$\Delta R(e, \gamma) < 1$ $ p_T^e - p_T^\gamma /p_T^{e(\gamma)} > 5\%$	
Categorisation		$m_{\ell\ell}$ closest to $m_Z = 91.2$ GeV		
		$ m_{\ell\ell}^{\text{FSRcorrected}} - m_Z < 15$ GeV		
Event selections		Trigger matching and overlap removal		
		$p_T^\gamma/m_{Z\gamma} > 0.2$; $200 < m_{Z\gamma} < 3500$ GeV		

prescription [202], to generate real emissions up to three partons at LO in QCD.

7.2 Object and Event Selection

The physics objects utilised in this analysis, namely photons, jets, electrons, and muons, were defined in Chapter 6. The selected events were required to have at least one primary vertex candidate reconstructed from tracks with $p_T > 500$ MeV. The candidate vertex with the highest sum of the square transverse momenta from its associated tracks was selected as the primary vertex.

The $X \rightarrow Z\gamma$ events were selected by associating two oppositely charged electrons or muons to reconstruct the Z boson and photon candidates linked to the primary vertex. Table 7.1 summarises the object and event selection criteria used in the analysis.

7.2.1 Photons

The selected photon candidates were required to have $p_T > 15$ GeV and $|\eta| < 2.47$, excluding events in the transition region, $1.32 < |\eta| < 1.52$. The Loose isolation requirement is applied to photon candidates at the object selection level, to allow for background composition studies described in Section 7.3.

7.2.2 Leptons

Muon candidates were required to have $p_T > 10$ GeV and be found within $|\eta| < 2.7$. Electron candidates were required to satisfy $p_T > 10$ GeV and $|\eta| < 2.47$, excluding candidates found in the transition region.

Muon candidates were required to meet the Medium identification criterion. In scenarios where the $ee\gamma$ system is highly boosted, distinguishing the two electrons becomes challenging because their energy deposits in the calorimeter are very close to each other. This challenge often leads to the misreconstruction of the sub-leading electron candidate as a photon. To mitigate this issue, a multivariate analysis based identification, MVA ID, criterion was developed using calorimeter shower shape variables and track-based variables. Electron candidates were required to satisfy the Mixed ID identification criterion, which combined the MVA ID and Loose ID criterion [203].

Events containing sub-leading electrons misreconstructed as photon candidates, $e(\gamma)$, were required to have $p_T > 50$ GeV, with η requirements similar to those for electrons. The $e(\gamma)$ candidates were required to have at least one track within $\Delta R < 0.1$ from the photon candidate.

A track-based Tight isolation requirement [151, 204] was applied to ensure the selected leptons were isolated from additional detector activity. The candidate leptons were required to be associated with the primary vertex by imposing a constrain of $|z_0 \sin \theta| < 0.5$ mm. To suppress secondary leptons originating from heavy flavour hadron decay chains, the significance of the transverse impact parameter, $|d_0/\sigma d_0|$, where σd_0 denotes the uncertainty in

the d_0 measurement, was required to be less than 5 and 3 for electron and muon candidates, respectively.

The candidate Z boson was reconstructed from two same-flavor oppositely-charged lepton candidates, either $e^\pm e^\mp$ or $\mu^\pm \mu^\mp$. In the case where the events contain $e(\gamma)$ candidate, additional object selection criteria were applied to suppress photon backgrounds and increase the selection efficiency of real electrons. The angular distance between the selected electron and $e(\gamma)$, $\Delta R(e, e(\gamma))$ was required to be less than one and the relative difference between the p_T of the electron candidate and $e(\gamma)$, $|p_T^e - p_T^{e(\gamma)}|/p_T^{e(\gamma)}$ was required to be greater than 5%. In scenarios where multiple Z boson candidates were reconstructed within an event, the system with a dilepton invariant mass, $m_{\ell\ell}$, nearest to the pole Z boson mass of 91.2 GeV was selected, after applying FSR corrections [205,206].

7.2.3 Overlap Removal

In instances where two physics objects shared a detector response signal, an overlap removal procedure was implemented. The ATLAS detector in the Run 2 had an excellent muon reconstruction resolution, as such it was not necessary to have muon overlap requirements on a pair of muons [207].

The reconstructed lepton candidates must exhibit good angular and track isolation. If two electron candidates shared track parameters, the candidate with the lower p_T was discarded. In cases where two electron candidates were situated within $|\Delta\eta| < 0.075$ and $|\Delta\phi| < 0.125$ of each other, the p_T of the leading electron candidate was used to optimise the overlap removal procedure. In events where the leading electron had $p_T < 500$ GeV, electron clusters closer than $|\Delta\eta| < 0.05$ and $|\Delta\phi| < 0.05$ were discarded. In cases where the leading electron had $p_T > 500$ GeV, electron clusters closer than $|\Delta\eta| < 0.05$ and $|\Delta\phi| < 0.05$ were discarded. If the track associated with an electron candidate was within $\Delta R < 0.02$ of a track associated with a muon candidate that passed all muon selection requirements, the muon candidate was given precedence over the electron candidate.

Photon candidates within $\Delta R < 0.3$ of the selected $m_{\ell\ell}$ were rejected to suppress background contributions from $Z+\gamma$ events and from photons potentially mimicking electrons. Jets within $\Delta R < 0.2$ of photons, pairs of electrons, or muons were discarded. The $e(\gamma)$ candidate was also considered in the overlap removal procedure. Overlap of the $e(\gamma)$ candidate with muons was treated in a manner similar to the electron overlap removal procedure. However, no overlap removal procedure applied between $e(\gamma)$ and electron candidates due to their close proximity.

7.2.4 Final Event Selection

After the overlap removal procedure, further event selection requirements were imposed. The final selected leptons were required to pass further stricter trigger thresholds. The additional trigger threshold were defined as:

$$p_T > p_T \text{ (trigger threshold)} + 1 \text{ GeV}, \quad (7.1)$$

for both single and dilepton triggers. This requirement helps to mitigate trigger turn-on inefficiencies. The selected photon candidate used in the reconstruction of the $\ell\ell\gamma$ candidate must match the trigger object that activated the photon high- p_T trigger.

The reconstructed $m_{\ell\ell}$ used to reconstruct the $m_{Z(\ell\ell)\gamma}$ was required to fall within 15 GeV of the Z boson pole mass. Previous studies have demonstrated that this constraint significantly improves the $Z\gamma$ mass resolution [205, 206].

The invariant mass, $m_{Z\gamma}$, of the selected $X \rightarrow Z\gamma$ candidate was computed using the reconstructed lepton and photon candidates. Selected events were required to fall within the invariant mass window, $200 < m_{Z\gamma} < 3500$ GeV, also termed the signal region. This range covers the mass spectrum where a scalar boson predicted in Section 1.3 could be detected.

To reduce the amount of the non-resonant $Z+\gamma$ background within the signal region, a variable called the relative photon p_T was introduced. It is defined as the photon p_T relative to $m_{Z\gamma}$, expressed as $p_T/m_{Z\gamma}$. A scan was made

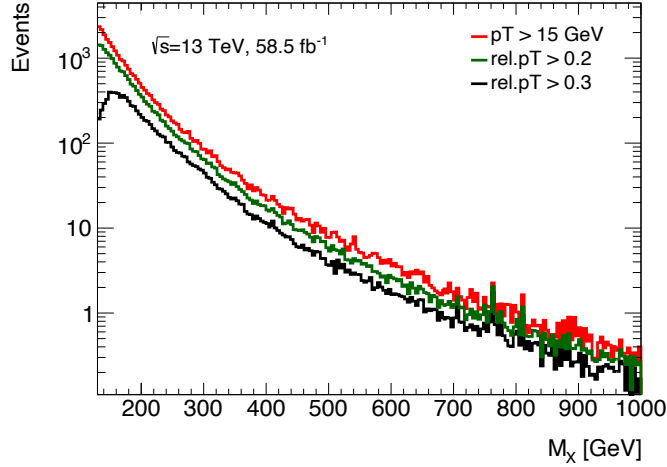


Figure 7.2 Distribution of $m_{Z\gamma}$ from simulated $Z + \gamma$ events for $p_T > 15$ GeV, $p_T/m_{Z\gamma} > 0.2$ and $p_T/m_{Z\gamma} > 0.3$ requirements.

over the possible $p_T/m_{Z\gamma}$ values to find an optimal value that maximises the increase in signal sensitivity, $S/\sqrt{S+B}$, where S and B denote the signal and background yields, respectively. Signal yields were estimated from simulated ggF samples, while the background yields were obtained from the simulated $Z + \gamma$ samples, considering background events within a 3% mass window around the mass of the simulated signal sample.

The optimal values of $p_T/m_{Z\gamma}$ maximising the signal sensitivity were found at $p_T/m_{Z\gamma} > 0.2$ and $p_T/m_{Z\gamma} > 0.3$. The effect of applying these two values alongside applying the minimal photon p_T cut required in the object selection is shown in Figure 7.2. The requirement $p_T/m_{Z\gamma} > 0.3$ exhibited a turn-on effect at approximately 200 GeV, as such the requirement $p_T/m_{Z\gamma} > 0.2$ was used for the analysis.

7.3 Background Composition Studies

In this Section, we employ data-driven methods to quantify the background composition of the selected dataset within the signal region. Specifically, we

use the 2D-sideband, 2D-SB, and template-fit methods, which have been previously used in various Run 1 measurements having prompt photons in the final state [208, 209].

The total background yields a smooth $m_{Z\gamma}$ distribution that can be effectively described by an analytic function. Parameters of the distribution are to be obtained directly from fits to the $m_{Z\gamma}$ distribution from data. The knowledge of the background composition plays an insignificant role in determining the background shape. It only becomes crucial when combining the $Z + \gamma$ and $Z + \text{jet}$ samples for performing spurious signal studies [210], which will be discussed in detail in Section 7.4.2.

7.3.1 Background Estimation With the 2D-SB Method

The 2D-SB is a counting method that relies on the definition of a two dimensional plane, as shown in Figure 7.3, based on the isolation and identification variables of the selected prompt photon candidate within the $\ell\ell\gamma$ triplet [211]. The signal region, A, characterised by an enhanced $Z\gamma$ contribution and three control regions B, C and D enriched with $Z + \text{jet}$ events are defined as follows:

- Tight and isolated, region A: Photon candidates in this region are both well isolated from hadronic activity and pass the Tight selection criteria.
- Tight but non-isolated, region B: Photon candidates here are not isolated from the hadronic activity but satisfy the Tight selection criteria.
- Non-tight, isolated region C: Photon candidates in this category are isolated from the hadronic activity, similar to region A, but fail the Tight identification criteria but satisfy some looser identification criteria.
- Non-tight, non-isolated region D: Photon candidates in this region are non-isolated, similar to those in region B, and pass the same identification requirements of region C.

These four regions are populated with events that satisfy all the object selection requirements of the analysis except for the photon identification and

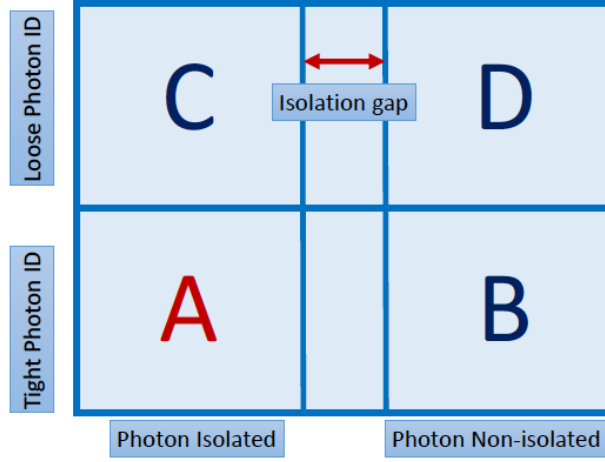


Figure 7.3 Definition of the ABCD regions based on photon identification and isolation variables.

isolation requirements. The looser photon identification requirements used to define regions C and D were derived from the nominal Tight identification criteria by relaxing the requirements on the four shower shape variables: ΔE , w_{s3} , E_{ratio} , and F_{side} , computed from a few strips of the first layer of the ECal, creating a new variable called Loose'4, more details about its definition can be found in Reference [210].

The $Z + \gamma$ yield, $N_A^{Z\gamma}$, in region A is estimated based on the observed number of events in data sample across four regions represented as N_k^{data} , where $k \in \{A, B, C, D\}$. The estimation is derived using the relation:

$$N_A^{Z\gamma} = N_A^{\text{data}} - (N_B^{\text{data}} - c_B N_A^{Z\gamma}) \frac{(N_C^{\text{data}} - c_C N_A^{Z\gamma})}{(N_D^{\text{data}} - c_D N_A^{Z\gamma})} R^{Zj}, \quad (7.2)$$

where $c_k \equiv \frac{N_k^{Z\gamma}}{N_A^{Z\gamma}}$ are signal leakage fractions defined as the ratio of the $Z + \gamma$ events in regions k to those in region A, extracted from the simulated $Z + \gamma$ samples.

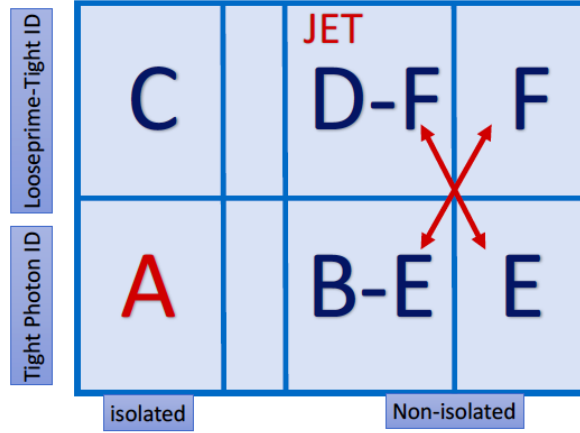


Figure 7.4 Definition of the ABCDEF regions to directly obtain R^{Zj} directly from data.

The correlation between the isolation and identification variables for the jets faking photons in $Z + \text{jet}$ events is quantified by the parameter R^{Zj} , defined as:

$$R^{Zj} \equiv \frac{N_A^{Zj} N_D^{Zj}}{N_B^{Zj} N_C^{Zj}}. \quad (7.3)$$

The value of R^{Zj} is expected to be 1 in the absence of correlations. The correlation parameter cannot be directly computed from data since it requires background events that fall into the signal region. While the parameter R^{Zj} can be obtained from the $Z + \text{jet}$ full MC simulated sample, the limited statistics in the simulated sample and a general lack of credibility of MC in these type of studies necessitate the use of data-driven methods. Therefore, alternative approaches are devised to estimate this parameter using data-driven techniques.

Data-driven Estimation of R Using $Z + \gamma$ Events

The estimation of R^{Zj} can be achieved by defining two new non-isolated regions called region E and region F, as illustrated in Figure 7.4. These additional regions are defined as follows:

- Extra tight but non-isolated region, E: In this region, the photon candidates can pass the Tight identification requirement but only in high E_T^{cone20} region, where the isolation threshold is defined as:
 $E_T^{\text{iso}} > 0.065 \times p_T + 3.5 \text{ GeV}.$
- Extra Loose'4 and non-isolated region, F: In this region, the photon candidates are allowed to pass the Tight identification but only in the high E_T^{cone20} region.

These two regions are defined so that R^{Zj} from Equation 7.3 can be estimated directly from data instead of MC simulations. Equation 7.3 is then represented as Equation 7.4:

$$R^{\text{data}} = \frac{N_{B-E}^{\text{data}} N_F^{\text{data}}}{N_{D-F}^{\text{data}} N_E^{\text{data}}}. \quad (7.4)$$

Data-driven Estimation of R Using $\gamma\gamma$ Events

To validate the accuracy of the purity calculated with $R^{\text{data}, Z\gamma}$ estimated using Equation 7.4, a new data-driven method is used. The main idea of the method is to use an $X + \text{jet}$ sample, where X is a well defined object, in order to estimate R in the same regions as those used to define R^{Zj} in Equation 7.3.

A $\gamma\gamma$ data sample composed of events containing two real photons, one real photon and a hadronic jet, and two hadronic jets is used. Each photon candidate is classified as either belonging to a category A, B, C or D, depending on whether it fails or passes the identification and isolation criteria, as shown in Figure 7.5. The two photon candidates are considered sequentially; first, the 2D-SB method is applied to extract events for which the leading p_T candidate is a true photon. Then, the method is applied to the subleading p_T

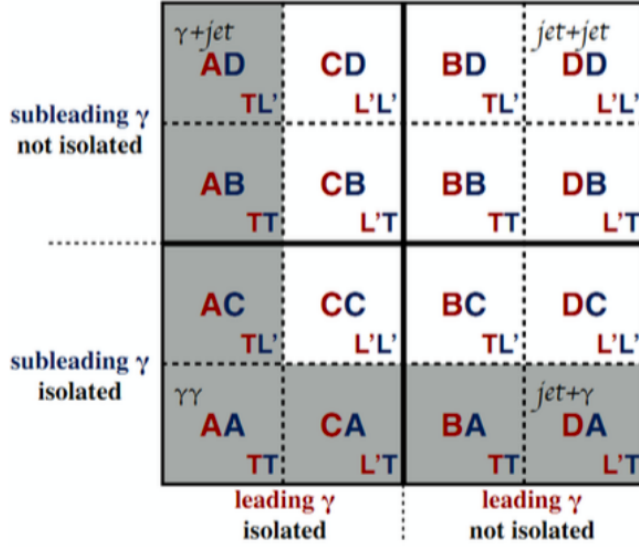


Figure 7.5 Definition of the 16 regions used in the $2 \times 2D$ -SB method of estimating purity of $\gamma\gamma$ events. The notations TT, TL, LL, and LT denote different categories based on the identification and isolation criteria of the leading and subleading photons. TT refers to both photons passing the Tight identification and isolation criteria. TL refers to the leading photon passing the Tight criteria and the subleading photon passing the Loose criteria. LL indicates both photons passing the Loose criteria. LT refers to the leading photon passing the Loose criteria and the subleading photon passing the Tight criteria.

photon candidate, knowing that the leading photon is tightly identified and well isolated, creating 16 orthogonal regions illustrated in Figure 7.5.

An $X + \text{jet}$ sample is formed by fixing either the leading or subleading photon to be well isolated and satisfy the Tight identification requirement, making it X . To obtain the desired $X + \text{jet}$ sample from the $\gamma\gamma$ data sample, pure $\gamma\gamma$ events are subtracted from the data sample. The number of $X + \text{jet}$ events in a region of interest is given by:

$$N_j^{\text{data}} = N_j^{\text{data}} - N_j^{\text{MC}} \times k - \text{factor}, \quad (7.5)$$

where j represents a region of interest, and $k\text{-factor} = N_{AA}^{\text{data}} \times \frac{\text{purity}}{N_{AA}^{\text{MC}}}$. The purity of $\gamma\gamma$ events was computed using the 2×2D-sideband method used in ATLAS studies having a diphoton system in the final state [208]. The method performs a data-driven background evaluation by extrapolating the background from the control regions defined in the sidebands of isolation and identification variables, in which the photons either pass or fail the Tight ID criteria or fail the isolation selection. The purity of $\gamma\gamma$ events is computed in a region where both the leading and subleading photons are well isolated and identified.

The data-driven correlation parameter in the case where the leading p_T photon candidate is classified as being Tight and isolated, $X = \gamma$, is computed as:

$$R^{\gamma+\text{jet}} = \frac{N_{AA}^{\gamma+\text{jet}}/N_{AB}^{\gamma+\text{jet}}}{N_{AC}^{\gamma+\text{jet}}/N_{AD}^{\gamma+\text{jet}}}. \quad (7.6)$$

In the case where the subleading p_T photon candidate is selected as X , the correlation parameter is defined as:

$$R^{\text{jet}+\gamma} = \frac{N_{AA}^{\text{jet}+\gamma}/N_{BA}^{\text{jet}+\gamma}}{N_{CA}^{\text{jet}+\gamma}/N_{DA}^{\text{jet}+\gamma}}. \quad (7.7)$$

The obtained values of $R^{\text{data}, \gamma\gamma}$ will depend on whether leading or subleading photon in the $\gamma\gamma$ data sample was selected as X when forming a $X + \text{jet}$ sample. To eliminate this potential bias, the selection of X will be done at random.

The following event selection requirements were used to select $\gamma\gamma$ events used to compute $R^{\text{data}, \gamma\gamma}$. The photon pairs were required to have a reconstructed invariant mass between 130 GeV and 3500 GeV. The p_T of the leading photon was required to be at least 40 GeV and the subleading photon was required to have a minimum p_T of 30 GeV. The photon isolation and identification requirements are the same as the ones used for selecting $Z\gamma$ events.

The measured purity of $\gamma\gamma$ events significantly impacts the computed value of $R^{\text{data}, \gamma\gamma}$. Figure 7.6 (b) illustrates the values of $R^{\text{data}, \gamma\gamma}$ using the central

Table 7.2 Summary of R values obtained from data-driven methods and Z +jet MC simulations at different photon p_T bins.

p_T^γ bin [GeV]	R (ABCDEF method)	R(Z +jet)	R ($\gamma\gamma$)
$40 \leq p_T < 90$	1.32579 ± 0.0524253	1.30308 ± 0.0390146	1.80002 ± 0.0372719
$90 \leq p_T < 140$	1.23662 ± 0.0705408	1.54748 ± 0.12818	1.98543 ± 0.0823212
$140 \leq p_T < 190$	1.116064 ± 0.10278	1.69762 ± 0.272027	1.95898 ± 0.12796
$p_T > 190$	1.59579 ± 0.165076	1.59876 ± 0.518419	2.02765 ± 0.132232

values of the $\gamma\gamma$ purity, as well as lower and upper values of purity obtained by subtracting or adding the statistical uncertainty of $\gamma\gamma$ purity from the central value. The largest difference between R values computed using the central of purity and the values computed using lower and upper values is propagated as the systematic uncertainty of $R^{\text{data}, \gamma\gamma}$.

Results of R Computed Using $Z + \gamma$ and $\gamma\gamma$ Events

To compare and validate the background estimations, a sample of $Z + \text{jet}$ events was generated. These events were simulated at NLO in QCD using the POWHEG framework [191–194], followed by showering with PYTHIA8 [146] event generators. The matrix element calculations employed the CT10 PDF set [212], while non-perturbative effects were modeled using the AZNLO parameter set alongside the CTEQ6L1 PDF set [197]. This MC serves as the basis for computing R^{Zj} , as outlined in Equation 7.3.

Table 7.2 summarises values of R obtained using the data-driven methods and the simulated $Z + \text{jet}$ sample at different bins of photon p_T . The errors associated with the R values obtained using the $\gamma\gamma$ events include both statistical and systematic uncertainties.

In the low-mass region, defined by photon p_T less than 90 GeV, the R value computed from $Z + \text{jet}$ MC simulated events closely aligns with the one computed using the data-driven ABCDEF method. As the photon p_T increases,

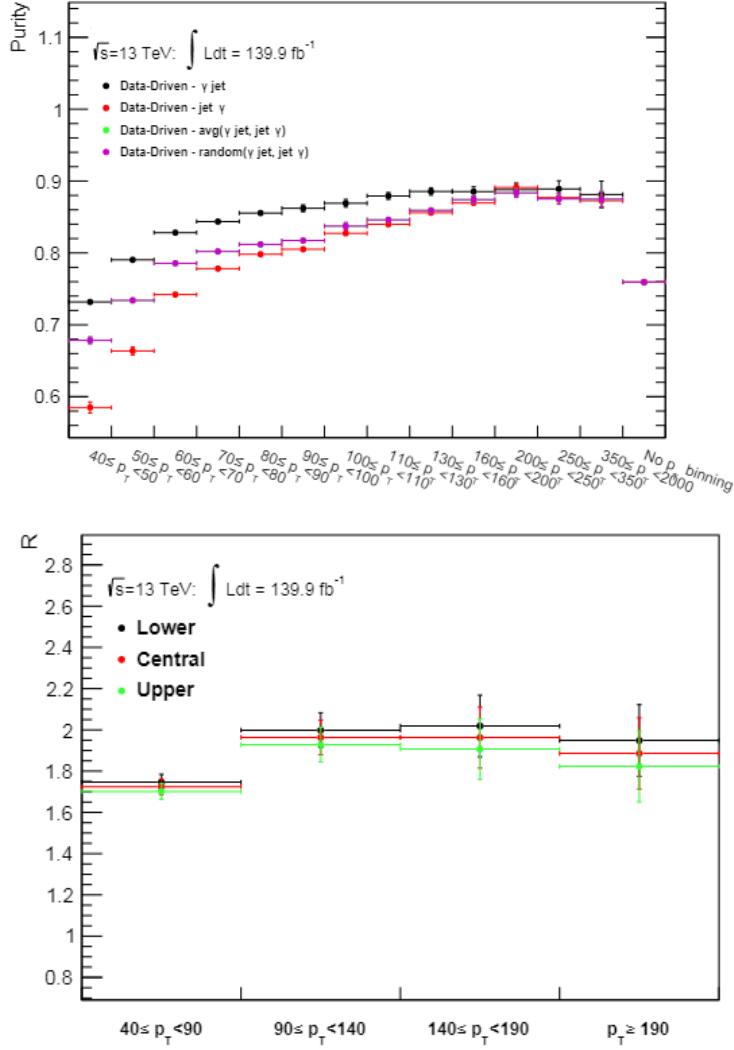


Figure 7.6 Top: diphoton purity at different photon p_T bins. Bottom: the R parameter computed using the central value of $\gamma\gamma$ purity $\pm$ the statistical error of the purity. The error bars represent the statistical uncertainty.

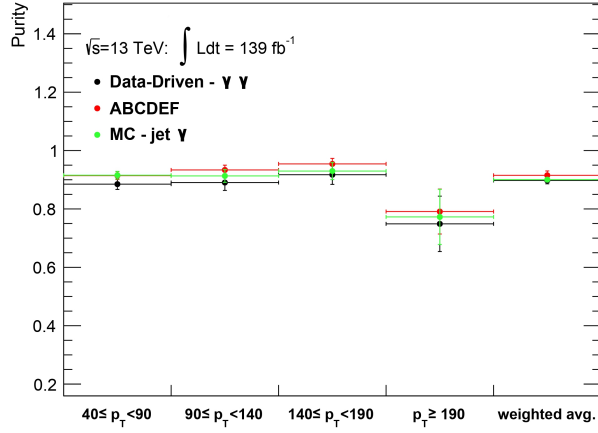


Figure 7.7 The observed $Z\gamma$ purity using R values computed using the different methods in different photon p_T bins.

going to higher masses, the R value computed with the $Z + \text{jet}$ MC gets closer to R computed with the $\gamma\gamma$ data-driven method.

The comparison of $Z + \gamma$ purity measurements in each photon p_T bin using different methods of computing R is illustrated in Figure 7.7. The analysis reveals that the choice of R value has minimal impact on the measured purity. Additionally, it is observed that as photon p_T increases, purity values obtained from the $\gamma\gamma$ data-driven method tend to converge towards those obtained from the MC-derived R values.

The three methods of estimating R resulted in the measured $Z\gamma$ that is relatively the same, except for the minor differences highlighted above.

7.3.2 Background Estimation Using the Template-fit Method

The template-fit method estimates the background event composition in the signal region by performing a binned template fit to the calorimeter isolation variable distribution of the photon candidates. This approach involves performing a simultaneous fit to both the signal and control regions, which are defined similarly to regions A and C of the 2D-SB method.

The simultaneous fit determines the yield and shape of the non-resonant $Z + \gamma$ events in both the signal and control region. Consequently, it determines the yield and shape of the $Z + \text{jet}$ distribution in those regions. The parameters of the fit are defined as follows:

- $N_{\text{SR}}^{\text{Sig}}$ represents the signal yield in the signal region.
- $N_{\text{SR}}^{\text{Bkg}}$ quantifies the amount of background in the signal region.
- $N_{\text{ibin}}^{\text{Bkg}}$ denotes the background yield in each bin of the calorimeter isolation variable, $E_{\text{T}}^{\text{cone20}}$, in the control region.

Based on the assumption that the $Z + \gamma$ shape in the signal region, denoted as $\text{Fraction}_{\text{ibin}}$, and the leakage of $Z + \gamma$ events between the signal and control regions in data are identical to those obtained from simulated $Z + \gamma$ background events, and further assuming that the background shape is consistent in both the signal and control regions, we derive the following equations:

$$N_{\text{ibinSR}} = N_{\text{SR}}^{\text{Sig}} \times \text{Fraction}_{\text{ibin}} + N_{\text{SR}}^{\text{Bkg}} \times \frac{N_{\text{ibin}}^{\text{Bkg}}}{\sum_{\text{ibin}=1}^{\text{Nbins}}}, \quad (7.8)$$

and

$$N_{\text{ibinCR}} = N_{\text{SR}}^{\text{Sig}} \times \text{Fraction}_{\text{ibin}} \times \text{Leakage}_{\text{ibin}} + N_{\text{ibin}}^{\text{Bkg}}, \quad (7.9)$$

where $N_{\text{ibin}}^{\text{SR}}$ and $N_{\text{ibin}}^{\text{CR}}$ represent the event counts in each $E_{\text{T}}^{\text{cone20}}$ bin of the signal and control regions, respectively. These equations are obtained for both the signal and control regions.

If $N_{\text{bin}} \geq 2$, a simultaneous fit can be performed to determine all parameters of the fit. Figures 7.8 and 7.9 show the outcome of applying the template-fit procedure to the distribution of the isolation variable of photon candidates in the signal region, left, and the control region, right, for ee and $\mu\mu$ channels, respectively. In these figures, signal represents $Z + \gamma$ events, and background corresponds to $Z + \text{jet}$ events. Table 7.3 summarises the purity, i.e., composition, of $Z + \gamma$ events in the signal region.

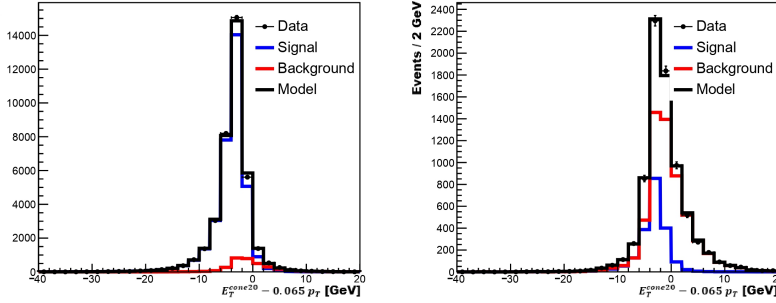


Figure 7.8 Template-fit to the distribution of the isolation variable of photon candidates in ee events.

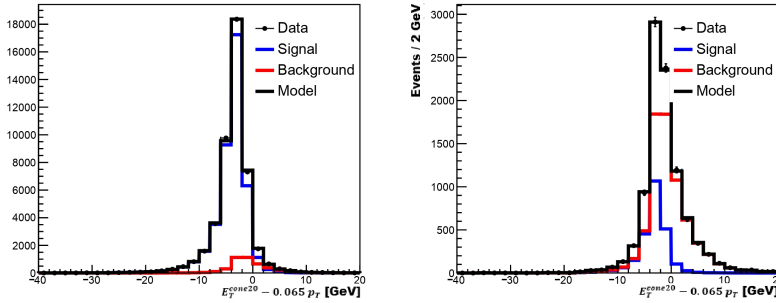


Figure 7.9 Template-fit to the distribution of the isolation variable of photon candidates in $\mu\mu$ events.

Table 7.3 Nominal purity of $Z + \gamma$ events extracted using the template-fit method for both ee and $\mu\mu$ channels.

Channel	Nominal purity	N_S	N_B	N_{Data}
ee	0.914 ± 0.004	34206.84 ± 233.66	3200.81 ± 146.96	37407.64 ± 193.34
$\mu\mu$	0.906 ± 0.004	41394.79 ± 262.08	4297.29 ± 169.94	45692.08 ± 213.74

Table 7.4 summarises the ratio of $Z + \gamma$ events to all background events, purity, obtained using different methods. The relative difference between the obtained purity values for each channel is less than 1%. Purity values obtained using

Table 7.4 The $Z + \gamma$ purity extracted using template-fit and 2D-SB methods for ee and $\mu\mu$ channels.

Channel	Template-fit	2D-SB
ee	0.914 ± 0.004	0.911 ± 0.0155
$\mu\mu$	0.906 ± 0.004	0.916 ± 0.0108

the template-fit method have smaller errors associated with them; therefore, they were used throughout the analysis.

7.4 Background and Signal Modelling

The signal and background yields were extracted via a maximum-likelihood fit to the $m_{Z\gamma}$ distribution obtained from data. The combined statistical model used to describe the $m_{Z\gamma}$ distribution is defined as follows:

$$f_{S+B} = N_{\text{sig}} \cdot f_{\text{sig}} + N_{\text{spur}} \cdot f_{\text{sig}} + N_{\text{bkg}} \cdot f_{\text{bkg}}, \quad (7.10)$$

where f_{sig} and f_{bkg} denote the normalised analytical functions representing the signal and background models, respectively. The parameters N_{sig} , N_{spur} and N_{bkg} correspond to the number of fitted signal, spurious signal and background events, respectively.

7.4.1 Signal Modelling

The signal shape is modeled using a Double-Sided Crystal Ball, DSCB, function [213,214]. The DSCB function combines a Gaussian core with exponential tails on both sides of the distribution mode. Mathematically, the DSCB function is expressed as:

$$\text{DSCB} = N \cdot \begin{cases} e^{-t^2/2}, & \text{if } -\alpha_{\text{low}} \leq t \leq \alpha_{\text{high}} \\ e^{-\frac{\alpha_{\text{low}}^2}{2}} \left[\frac{1}{R_{\text{low}}} (R_{\text{low}} - \alpha_{\text{low}} - t) \right]^{-n_{\text{low}}}, & \text{if } t < -\alpha_{\text{low}} \\ e^{-\frac{\alpha_{\text{high}}^2}{2}} \left[\frac{1}{R_{\text{high}}} (R_{\text{high}} - \alpha_{\text{high}} - t) \right]^{-n_{\text{high}}}, & \text{if } t > \alpha_{\text{high}} \end{cases} \quad (7.11)$$

where $t = (x - \mu_{\text{CB}}) / \sigma_{\text{CB}}$, $R_{\text{low}} = \frac{\alpha_{\text{low}}}{n_{\text{low}}}$, and $R_{\text{high}} = \frac{\alpha_{\text{high}}}{n_{\text{high}}}$. The normalisation parameter is given by N , μ_{CB} and σ_{CB} represent the mean and width of the Gaussian distribution, respectively. Parameters α_{low} and α_{high} denote the transition points from the Gaussian core to the exponential tails on the low and high sides, while n_{low} and n_{high} represent the exponents of the exponential tails on the low and high sides, respectively.

The parameters of the DSCB function were determined by fitting it to the simulated signal samples mentioned in Section 7.1.1. These parameters were then interpolated as a function of m_X to obtain the signal shape at mass points where signal samples were not available [203]. Figure 7.10 shows the $m_{Z\gamma}$ distribution of the spin-0 resonance at $m_X = 1000$ GeV. The parametric model obtained from the fit matches the simulated signal, as indicated by the agreement between the markers representing the simulated events and the solid, for $ee\gamma$, and dashed, for $\mu\mu\gamma$, lines representing the fitted models

7.4.2 Background Modelling

The background shape was obtained by fitting analytical functions to the $m_{Z\gamma}$ distribution of the background processes. This distribution was constructed using a combination of the simulated $Z + \gamma$ sample and $Z + \text{jet}$ sample derived from a control region in data. The data-driven $Z + \text{jet}$ sample was obtained by applying all event selection criteria to the data sample, except for the Tight identification criterion for the photon candidate, which was intentionally failed.

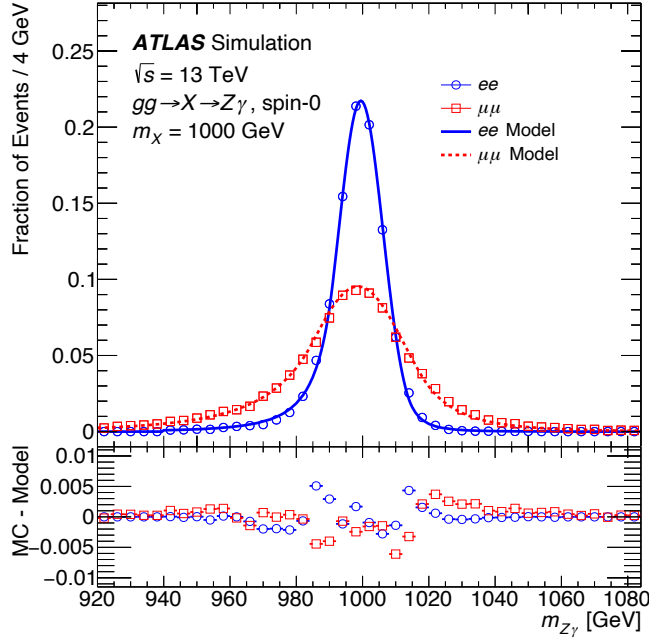


Figure 7.10 Distributions of the $m_{Z\gamma}$ for the spin-0 resonance at $m_X = 1000$ GeV, generated via the ggF process [203].

Various candidate analytic functions were considered for the background model selection. The following functions were considered to find the best-fitting model for describing the background distribution:

FK0 $f = x^{a_0}(1 - \sqrt[3]{x})^{a_1}$, where $x = \frac{m_{Z\gamma}}{13000}$

Dijet $f = x^{a_0}(1 - x)_1^a$, where $x = \frac{m_{Z\gamma}}{13000}$

DijetHighMass $f = x^{a_0}(1 - \sqrt[3]{x})^{a_1}$, assuming $a_0 < 0$, where $x = \frac{m_{Z\gamma}}{13000}$

Dijet2HighMass $f = x^{a_0}(1 - x)^{a_1}$, assuming $a_0 < 0$, where $x = \frac{m_{Z\gamma}}{13000}$

The selection of an appropriate analytic function for the background model is critical, as it should have the least number of free parameters to prevent a decrease in sensitivity. To mitigate bias in selecting the analytic function, spurious signal studies were performed to quantify any bias in the extracted signal yield. The procedure for performing spurious signal studies is described

Table 7.5 Analytic background functions satisfying the criteria of spurious signal study.

	Background model	Free parameters	S	max. $ \frac{S}{\delta S} $	χ^2/DoF	$P(\chi^2)$ [%]
$X \rightarrow Z(ee)\gamma$	DijetHighMass	2	-11.5	-37.8%	18.1	0.00
	Dijet2HighMass	2	9.65	27.5%	1.52	0.00
	FK0	2	-11.5	-37.8%	17.8	0.00
	Dijet	2	9.63	23.3%	1.51	0.00
$X \rightarrow Z(\mu\mu)\gamma$	Dijet2HighMass	2	10.0	38.3%	1.46	0.00
	Dijet	2	10.1	38.2%	1.47	0.00
	DijetHighMass	2	7.69	28.7%	5.91	0.00
	FK0	2	7.68	28.7%	5.87	0.00

in detail in Reference [215]. These studies help assess the impact of different background models on the analysis sensitivity and ensure that the chosen model does not introduce a significant bias in the results.

The candidate analytic functions were fitted to the total background distribution, and their parameters were determined as a function of $m_{Z\gamma}$. The scan points of the fit were chosen from 220 GeV to 3400 GeV, a range narrower than the signal region window, 200-3500 GeV, to guarantee the stability of the fits. The analytic function with the least number of free parameters and the smallest $S/\sigma S$ value were chosen, where S represents the maximum fitted signal yield and σS is the statistical uncertainty of S .

Table 7.5 summarises the results of spurious signal studies obtained after fitting the analytic functions to the $Z(ee)\gamma$ and $Z(\mu\mu)\gamma$ background distributions. The Dijet function was selected to model both the electron and muon channels. The spurious signal yields ranged from 10.2 to 0.003 events for the electron channel and from 10.7 to 0.007 events for the muon channel. These ranges of spurious signal yield were used in defining systematic uncertainties of the background modelling, parameterised as an exponential function of $m_{Z\gamma}$.

7.5 Statistical Procedure

To search for a localised excess of events over a smoothly falling SM background in the $m_{Z\gamma}$ distribution of the data, an unbinned profile-likelihood-ratio fit [216] was employed. The parameter of interest, $\sigma(pp \rightarrow X) \cdot \mathcal{B}(X \rightarrow Z\gamma)$, was estimated within the mass range of 220 GeV to 3400 GeV.

The expected number of signal events, N_{sig} , from Equation 7.4 can be expressed as $N_{\text{sig}} = \mathcal{L} \cdot \epsilon \cdot \sigma(pp \rightarrow X) \cdot \mathcal{B}(X \rightarrow Z\gamma)$, where $\mathcal{L}$ is the integrated luminosity of the dataset and ϵ is the parameterised signal efficiency as a function of m_X . The signal efficiency was defined as the number of events meeting all event selection requirements over the total number of events. Equation 7.4 can therefore be expressed as:

$$f_{S+B} = \mathcal{L} \cdot \epsilon \cdot \sigma(pp \rightarrow X) \cdot \mathcal{B}(X \rightarrow Z\gamma) \cdot f_{\text{sig}} + N_{\text{spur}} \cdot f_{\text{sig}} + N_{\text{bkg}} \cdot f_{\text{bkg}}. \quad (7.12)$$

The extended likelihood function $\mathcal{L}(\alpha, \theta)$ was defined as the product of a Poisson term, representing the observed data, the $f_{\text{tot}}(m_{Z\gamma}^i, \alpha, \theta)$ representing the PDF of the $m_{Z\gamma}$ distribution for each candidate event i , and constraint terms $G(\theta)$ for the nuisance parameters:

$$\mathcal{L}((\alpha, \theta) | \{m_{Z\gamma}^i\}_{i=1..n}) = \frac{e^{-N(\alpha, \theta)} N^n(\alpha, \theta)}{n!} \left(\prod_{i=1}^n f_{\text{tot}}(m_{Z\gamma}^i, \alpha, \theta) \right) \times G(\theta), \quad (7.13)$$

where $\alpha = \sigma(pp \rightarrow X) \cdot \mathcal{B}(X \rightarrow Z\gamma)$ is the parameter of interest, θ is a set of nuisance parameters, n is the observed numbers of events, and N is the expected event yield, a sum of N_{sig} , N_{bkg} and the spurious signal, $N_{\text{spur}}\theta_{\text{spur}}$, events. The parameter θ_{spur} is the nuisance parameter associated with the spurious signal.

The PDF describing $m_{Z\gamma}$, $f_{\text{tot}}(m_{Z\gamma}^i, \alpha, \theta)$, is constructed from the PDFs describing the signal and background models. It is formulated as:

$$f_{\text{tot}}(m_{Z\gamma}^i, \alpha, \theta) = \frac{1}{N} \sum_c \left\{ \left[N_{\text{sig}}^{(c)}(m_X, \alpha, \theta_{\text{sig}}) + N_{\text{spur}}^{(c)}(m_X, \theta_{\text{spur}}) \right] \times \right. \\ \left. f_{\text{sig}}^{(c)}(m_{Z\gamma}^i, \theta_{\text{sig}}) + N_{\text{bkg}}^{(c)} \times f_{\text{bkg}}^{(c)}(m_{Z\gamma}^i, \theta_{\text{bkg}}) \right\}, \quad (7.14)$$

where c is the index indicating the category, the decay channel of the Z boson, and θ represents a set of nuisance parameters. The parameters θ_{bkg} represent nuisance parameters that determine the background shape, while θ_{sig} denotes nuisance parameters associated with the uncertainties in parameterising the signal efficiency and acceptance. The nuisance parameters associated with uncertainties in the event yield or the $m_{Z\gamma}$ resolution are modelled with log-normal PDFs, while the nuisance parameters associated with the $m_{Z\gamma}$ signal peak position are modelled with Gaussian PDFs. Similarly, the parameters θ_{spur} are also assigned Gaussian PDFs.

The compatibility between data and a hypothesis assuming $\alpha = \mu$ is quantified by a p -value. This observed p -value gives the probability for data generated under the $\alpha = \mu$ hypothesis to be in the same or worse agreement with that hypothesis. A large p -value would indicate that the data is consistent with the hypothesis. The calculation of the p -value involves determining a likelihood-ratio test statistic $q(\mu)$, which is expressed as:

$$q(\mu) = -2 \ln \lambda(\mu) = -2 \ln \frac{L(\mu, \hat{\hat{\theta}}(\mu))}{L(\hat{\mu}, \hat{\theta})}, \quad (7.15)$$

where μ represents the parameter of interest, allowed to vary independently, and θ is the nuisance parameters. The parameters $\hat{\mu}$ and $\hat{\theta}$ are chosen to maximise the likelihood function in a background-only fit, and $\hat{\hat{\theta}}(\mu)$ is the best-fit value of θ under the hypothesis assuming signal strength μ [216].

The p -value of the null hypothesis is computed as:

$$p(\mu) = \int_{q_{\mu, \text{obs}}}^{\text{inf}} f(q_{\mu} | m_H, \mu, \hat{\hat{\theta}}(\mu)) dq_{\mu}, \quad (7.16)$$

where $q_{\mu,\text{obs}}$ gives the value of $q(\mu)$ observed from data, and $f(q_\mu)$ denotes the PDF of q_μ under the assumption of the signal strength μ .

Assuming that the contribution of a signal process to the mean number of events is non-negative, any physically meaningful estimator for μ must also be non-negative. In a model where $\mu \geq 0$, if one finds data such that $\hat{\mu} < 0$, it implies that the best agreement between data and any physically allowed value of μ must occur at $\mu = 0$. For purposes of establishing upper limits on $\sigma(pp \rightarrow X) \times \mathcal{B}(X \rightarrow Z\gamma)$, the test statistic $\tilde{q}_\mu$ is defined as:

$$\tilde{q}_\mu = \begin{cases} -2\ln\tilde{\lambda}(\mu) & \hat{\mu} \leq \mu \\ 0 & \hat{\mu} > \mu \end{cases} = \begin{cases} -2\ln\frac{L(\mu,\hat{\theta}(\mu))}{L(0,\hat{\theta}(0))} & \hat{\mu} < 0 \\ -2\ln\frac{L(\mu,\hat{\theta}(\mu))}{L(\hat{\mu},\hat{\theta})} & 0 \leq \hat{\mu} \leq \mu \\ 0 & \hat{\mu} > \mu \end{cases} . \quad (7.17)$$

The rationale behind setting $\tilde{q}_\mu = 0$ for $\hat{\mu} > \mu$ is that when establishing upper limits, data with $\hat{\mu} > \mu$ is not regarded as less compatible with μ than the obtained data. As a result, such data is not considered in the rejection region of the test. From the definition of the test statistic, higher values of $\tilde{q}_\mu$ represent greater incompatibility between data and the hypothesised value of μ .

7.6 Results

In the analysis, both statistical and systematic uncertainties are considered to ensure accurate and reliable results. Table 7.6 summarises the estimated systematic uncertainties for the $X \rightarrow Z\gamma$ search in the ee and $\mu\mu$ channels and the mass range $m_X = 220\text{--}3400$ GeV.

Simultaneous fits were performed on the unbinned distributions of $m_{Z\gamma}$ events for both $X \rightarrow ee\gamma$ and $X \rightarrow \mu\mu\gamma$ channels. Figure 7.11 shows fits to the $m_{Z\gamma}$ distributions of data events using the signal-plus-background models. The bottom panels in the figure show the statistical significance, which is defined

Table 7.6 The main sources of systematic uncertainty for the $X \rightarrow Z\gamma$ search. The gluon–gluon fusion spin-0 signal samples produced for $m_X = 220\text{--}3400$ GeV are used to evaluate the systematic uncertainty. The uncertainty ranges span the variations among different categories and different m_X resonance masses. The uncertainty due to the spurious signal is reported as an absolute number of events. In the table, “ID” for photons and electrons refers to identification efficiency uncertainties, “ISO” refers to isolation efficiency uncertainties, “TRIG” refers to trigger efficiency uncertainties, “RECO” refers to muon reconstruction efficiency uncertainty and “TTVA” refers to muon track-to-vertex-association efficiency uncertainty [203].

Category	$\mu\mu\gamma$	$ee\gamma$
Luminosity	0.83%	
<i>Signal Efficiency</i>		
Photon ID/ISO/TRIG efficiency	1.0%–1.5%	1.0%–1.7%
Muon ISO efficiency	1.0%–1.2%	–
Muon RECO efficiency	0.22%–6%	–
Muon TTVA efficiency	0.14%–0.23%	–
Muon TRIG efficiency	0.6%–1.0%	–
Electron ID/ISO/RECO/TRIG efficiency	–	2.9%–4%
MVA/Mixed electron ID efficiency	–	1.0%–1.1%
Pile-up	< 0.016%	–
<i>Signal modelling effect on μ_{CB}</i>		
Electron and photon energy scale	0.33%–0.4%	0.15%–0.7%
Muon momentum scale/sagitta bias	< 0.023%	–
<i>Signal modelling effect on σ_{CB}</i>		
Electron and photon energy resolution	2.5%–10%	7%–60%
Muon ID resolution	0.4%–1.8%	–
Muon MS resolution	0.6%–1.9%	–
Extra smearing of muon p_{T}	2.4%	–
<i>Background modelling, in number of events</i>		
Spurious signal	0.01–10.7	0.003–10.2

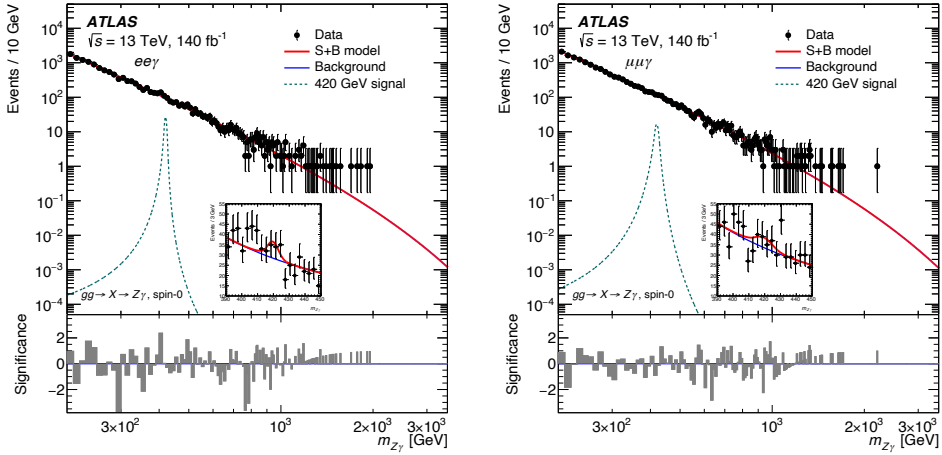


Figure 7.11 The $m_{Z\gamma}$ distributions of events meeting all the high-mass selection criteria in data for the $ee\gamma$, left, and $\mu\mu\gamma$, right, channels. The black points represent the data with the statistical uncertainty. The background component (shown as a dashed blue line) and the signal component (represented by a dashed dark cyan line) are depicted in the signal-plus-background unbinned fit (illustrated as a solid red line) to the data. These fits were performed independently in each channel [203].

as the residual of the data with relative to the background-only fit divided by the statistical uncertainty of the data.

For the spin-0 heavy resonances, the most significant excess is observed at 420 GeV, with a local significance of 2.1σ in the electron channel and 1.1σ in the muon channel. The lower mass region is zoomed in and displayed in the inset plot, where the excesses at 420 GeV can be seen. There is also an excess of about 2.1σ at around 250 GeV, as seen in Figure 7.12, the range of interest predicted by the Madala hypothesis that predicted the multi-lepton anomalies.

The computation of p -values is elaborated in Section 7.5. Figure 7.12 shows the observed local p -values as a function of m_X for the $ee\gamma$, $\mu\mu\gamma$ and the combined $ee\gamma + \mu\mu\gamma$ channels. The observed local p_0 -values can equivalently be expressed in terms of the number of standard deviations, referred to as the

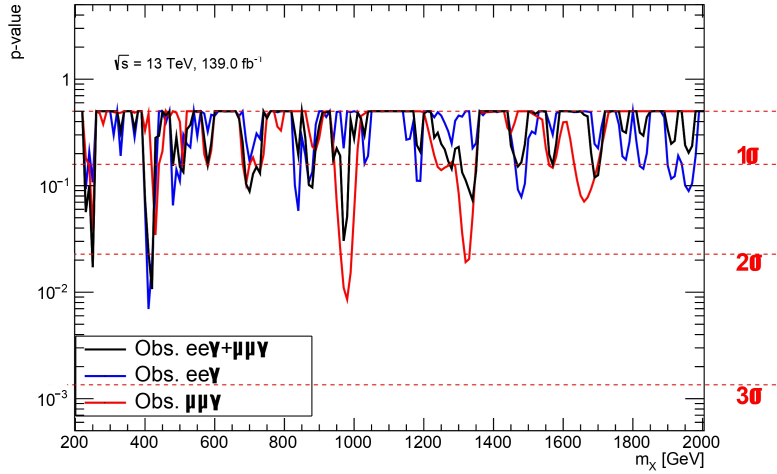


Figure 7.12 Distribution of the observed local p -values as a function of m_X .

local significance, Z . In the asymptotic regime, the distribution of $\tilde{q}_\mu$ follows a half χ^2 distribution for one degree of freedom [216], enabling the calculation of the local significance using:

$$Z = \sqrt{\tilde{q}_\mu}. \quad (7.18)$$

No significant excess with respect to the background-only hypothesis was observed. Consequently, upper limits at 95% CL were set on $\sigma(pp \rightarrow X) \cdot \mathcal{B}(X \rightarrow Z\gamma)$. The modified CL_s frequentist approach uses the PDFs of the test statistic $\tilde{q}_\mu$ to set upper limits on the cross-section [217] using a closed-form asymptotic formulae [216]. The limits at 95% CL are determined by identifying the signal cross-section for which the CL_s value is equals 0.05. From Figure 7.11, it can be seen that the highest-mass point in the data are at approximately 2.0 TeV for both the $ee(\gamma)$ and $\mu\mu\gamma$ channels. Limits beyond those mass points were set using Asimov toy-generated events [216].

The observed and expected upper limits on $\sigma(pp \rightarrow X) \cdot \mathcal{B}(X \rightarrow Z\gamma)$ as a function of m_X for a spin-0 heavy resonance signal produced via the ggF

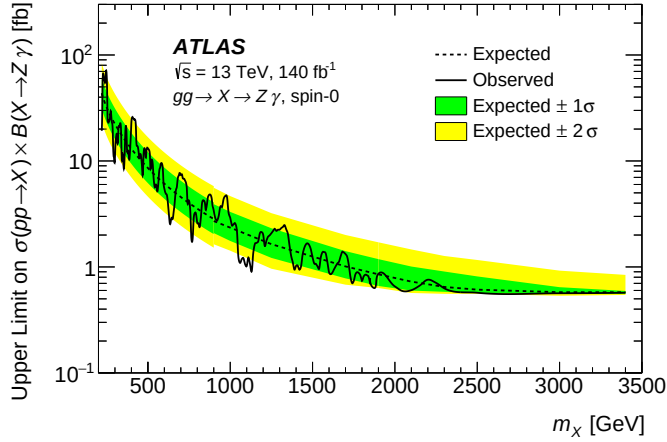


Figure 7.13 Observed, solid line, and expected, dashed line, 95% CL upper limits on $\sigma(pp \rightarrow X) \cdot \mathcal{B}(X \rightarrow Z\gamma)$ of a narrow-width spin-0 resonance X produced via ggF decaying into $Z\gamma$ as a function of m_X [203].

process is shown in Figure 7.13. Additionally, Figure 7.14 shows the observed and expected limits interpreted in terms of a spin-2 heavy resonance produced via ggX , Figure 7.14 (a), and $q\bar{q}X$, Figure 7.14 (b), processes. The dashed lines on Figures 7.13 and 7.14 represent the expected limits, while the green and yellow bands correspond to the $\pm 1\sigma$ and $\pm 2\sigma$ intervals for the expected upper limits. The observed upper limits are indicated by solid black lines.

Table 7.7 presents the observed and expected upper limits on $\sigma(pp \rightarrow X) \cdot \mathcal{B}(X \rightarrow Z\gamma)$ 95% CL. The upper limits are provided for spin-0 and spin-2 heavy resonances. The ranges of the upper limits are provided, indicating the variation from high to low mass points. The observed spin-0 upper limits on $\sigma(pp \rightarrow X) \cdot \mathcal{B}(X \rightarrow Z\gamma)$ improve the previous observed limits, reported in Reference [187], by a factor of 1.9 to 4 in the m_X range between 250 GeV and 2400 GeV

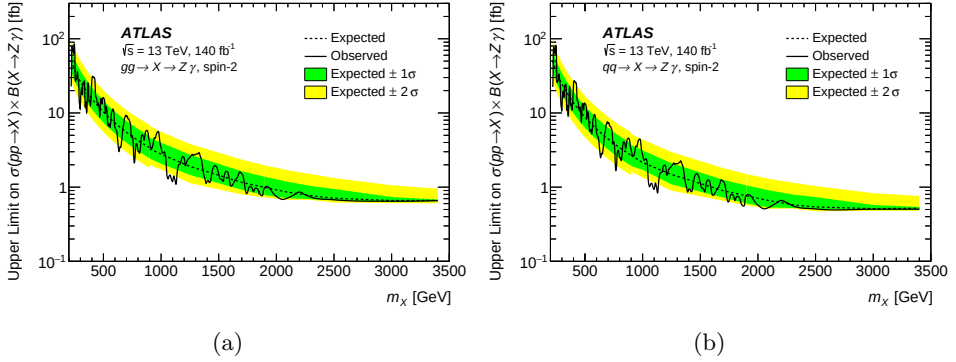


Figure 7.14 Observed, solid line, and expected, dashed line, 95% CL limits on $\sigma(pp \rightarrow X) \cdot \mathcal{B}(X \rightarrow Z\gamma)$ of a narrow-width spin-2 resonance X produced via (a) ggX (b) $q\bar{q}X$ processes and decaying into $Z\gamma$ function of m_X [203].

Table 7.7 Observed and expected upper limits on $\sigma(pp \rightarrow X) \cdot \mathcal{B}(X \rightarrow Z\gamma)$ at 95% CL.

95% CL upper limits on $\sigma(pp \rightarrow X) \cdot \mathcal{B}(X \rightarrow Z\gamma)$	Observed	Expected
ggF spin-0	65.5 fb – 0.6 fb	43.3 fb – 0.6 fb
ggX spin-2	77.4 fb – 0.6 fb	50.8 fb – 0.6 fb
$q\bar{q}X$ spin-2	76.1 fb – 0.5 fb	50.3 fb – 0.5 fb

Conclusions

The discovery of the Higgs boson in 2012 at the LHC by the ATLAS and CMS Collaborations marked a significant milestone in particle physics, providing crucial insights into the mechanism of electroweak symmetry breaking, and completing the SM of particle physics. However, subsequent studies of the properties of the Higgs boson revealed discrepancies between the observed experimental data and SM predictions, particularly in final states containing multiple leptons.

To address these inconsistencies, the Madala hypothesis, a model based on the 2HDM+S framework, was proposed. This model introduces a heavy scalar boson, H , potentially involved in the electroweak symmetry breaking. As discussed in Section 1.3, the decay processes of H , were understood through the introduction of a mediating scalar singlet S , which has globally re-scaled Higgs-like couplings. Previous studies have shown that fixing the mass of H to approximately 270 GeV and the mass of S fixed to around 150 GeV enables the Madala hypothesis to account for numerous multi-lepton anomalies observed in the published LHC results. The hypothesis also predicts signatures that can be probed in the LHC data.

The thesis is composed of two main parts: performances studies utilising the TileCal in Run 2 and Run 3, and a search for high-mass resonances in the $Z\gamma$ final state in pp collisions at $\sqrt{s} = 13$ TeV, recorded with the ATLAS detector.

Alongside the centre-of-mass energy, luminosity is one of the most important parameters that characterise the performance of a particle accelerator. Integrated luminosity gives the number of pp collisions in a data sample, and the uncertainty on the integrated luminosity is often a leading source of uncer-

tainty in precise measurement of $t\bar{t}$, W , and Z boson production cross-sections. An accurate estimate of luminosity is also often required in studies searching for new physics phenomena to determine background levels and sensitivity. The LUCID detector provides primary luminosity measurements for the ATLAS detector, other luminometers are used for calibrations and systematic comparisons. It is calibrated during vdM runs where the pileup levels are very low, and track-counting luminosities are used to transport the calibrations to the physics data-taking regime at high luminosity. Drifts in LUCID response during the 2022 data-taking campaign were quantified through the long-term stability uncertainty. Luminosity measurements from the TileCal D6 cells were compared with LUCID luminosity measurements to quantify the long-term stability uncertainty. The preliminary results found a mean deviation of 0.233%.

The performance and stability of gap and crack scintillators of the TileCal were evaluated using single isolated muons originating from $W \rightarrow \mu\nu_\mu$ events. The events were produced from pp collisions at $\sqrt{s} = 13$ TeV and recorded by the ATLAS detector during the Run 2 data-taking period spanning from 2015 to 2018. An unbinned negative log likelihood function was minimised to obtain the average response of each cell and a systematic uncertainty term during the 2015-2016, 2017 and 2018 data-taking years. The systematic uncertainty term captured the variation of the response of a given cell type across different modules. The upper limit on the average non-uniformity in ϕ was found to be 17.1% in E1 cells on the C-side of the detector during the 2018 data-taking period. The deviant modules were flagged by the calibration systems, indicating that the calibration systems were performing well. It was also found that E3 and E4 cells have the most noise and suffer greatly from radiation damage. The performance of the gap and crack scintillators was found to be in accordance with expectation.

The search for high-mass $Z\gamma$ resonances with a spin-0 or spin-2 utilised pp collision data collected by the ATLAS detector at $\sqrt{s} = 13$ TeV during the Run 2 data-taking period, corresponding to an integrated luminosity of 140 fb^{-1} . The high-mass resonance is searched, in a mass range from 220 to 3400

GeV, covering a mass range where a heavy scalar boson predicted by the hypothesis predicted the multi-lepton anomalies can potentially be observed. The study performed an unbinned profile-likelihood-ratio simultaneous fits to search for localised excesses on the smoothly falling background on the $m_{Z\gamma}$ distribution of data events in the $ee\gamma$ and $\mu\mu\gamma$ channels. A mild excess of 2.1σ was observed at 250 GeV, in the region of interest of the multi-lepton anomalies. The study set upper limits at 95% CL on $\sigma(pp \rightarrow X) \cdot \mathcal{B}(X \rightarrow Z\gamma)$ as a function of m_X in mass range 220 GeV to 3400 GeV. In the spin-0 interpretations, the observed upper limits on $\sigma(pp \rightarrow X) \cdot \mathcal{B}(X \rightarrow Z\gamma)$ ranged between 64 fb and 0.6 fb. For the spin-2 interpretation where X is produced via $q\bar{q}X$ and ggX processes, the observed upper limits varied between 75.1 fb and 0.5 fb, and between 77.4 fb and 0.6 fb, respectively. These observed upper limits on $\sigma(pp \rightarrow X) \cdot \mathcal{B}(X \rightarrow Z\gamma)$ improved the previous observed limits by a factor of 1.9 to 4 in the m_X range between 250 GeV and 2400 GeV.

Appendix A

Performance of Gap and Crack Scintillators

This appendix contains additional plots that are directly relevant to Chapter 5. Figures A.1 and A.2 show distributions of the average response of modules in gap, E1 and E2 cells, and crack, E3 and E4 cells, scintillators. The red line in the figures represents the fitted average response across all modules and the yellow band represents the systematic uncertainty attributed to the non-uniformity across the modules.

Figures A.3 and A.4 show the pull distributions used to check for azimuthal intercalibration in gap and crack scintillators on the C-side of the detector. The total number of modules in the gap and crack scintillators found to deviate more than three standard deviation from the mean of the pull distribution. The number of modules that a three standard deviation from the fitted parameter $\hat{\mu}_C$ was found to be four in 2015 and 2016 and only two in 2017 and 2018 data-taking periods.

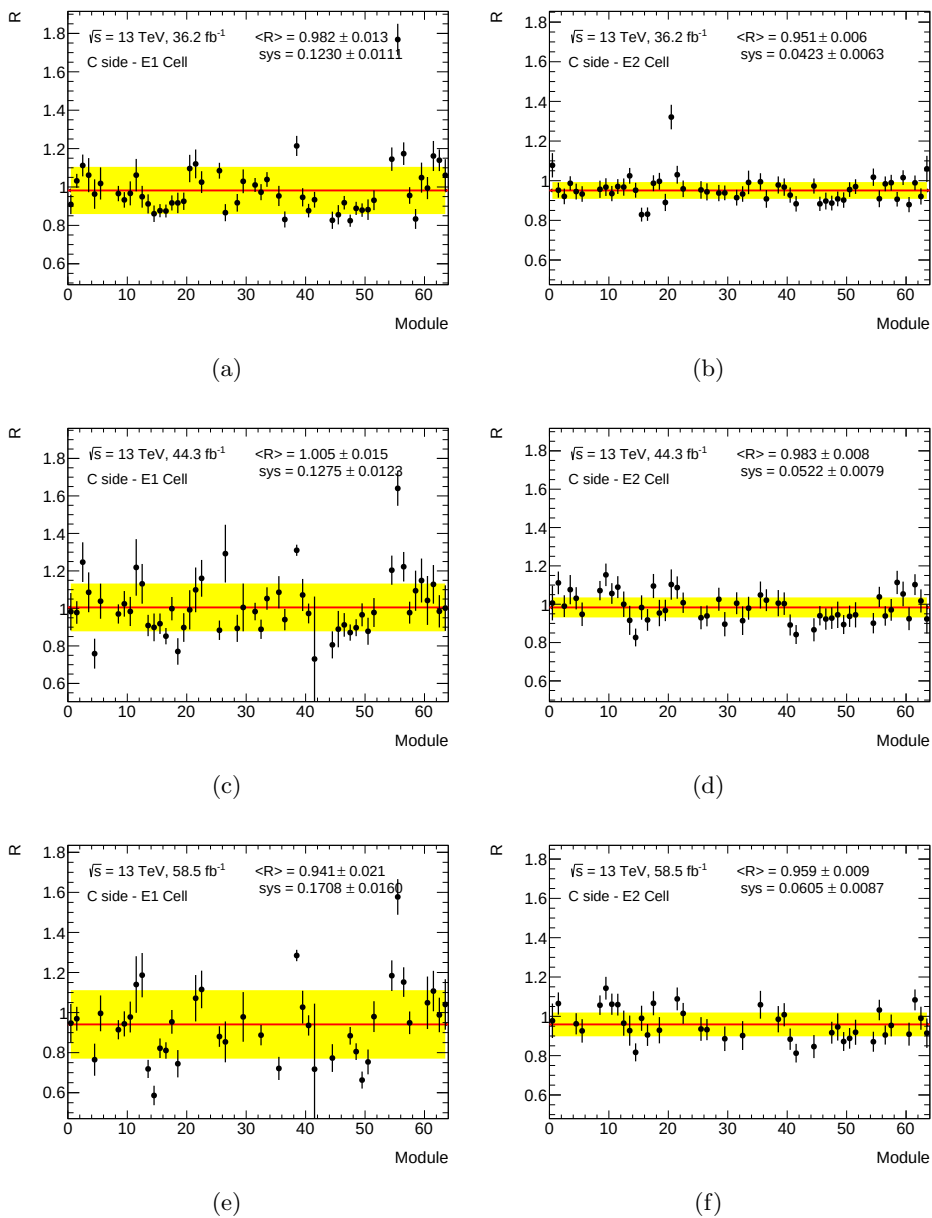


Figure A.1 Distributions of module numbers vs the average response in E1 cells on the C side. The top, middle, and bottom rows represent the respective years.

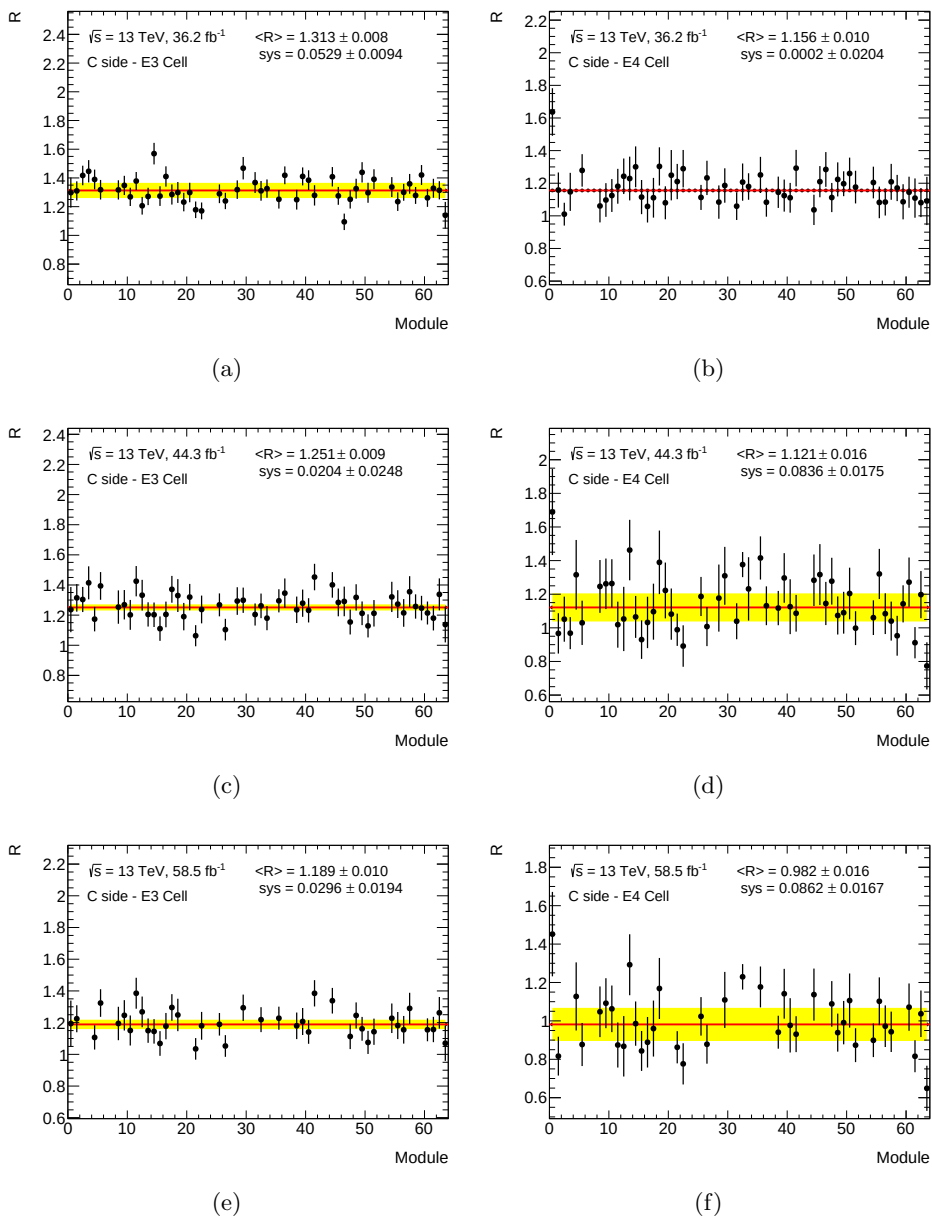


Figure A.2 Distributions of module number vs the average response in E1 cells on the C side. The top, middle, and bottom rows represent the respective data-taking years.

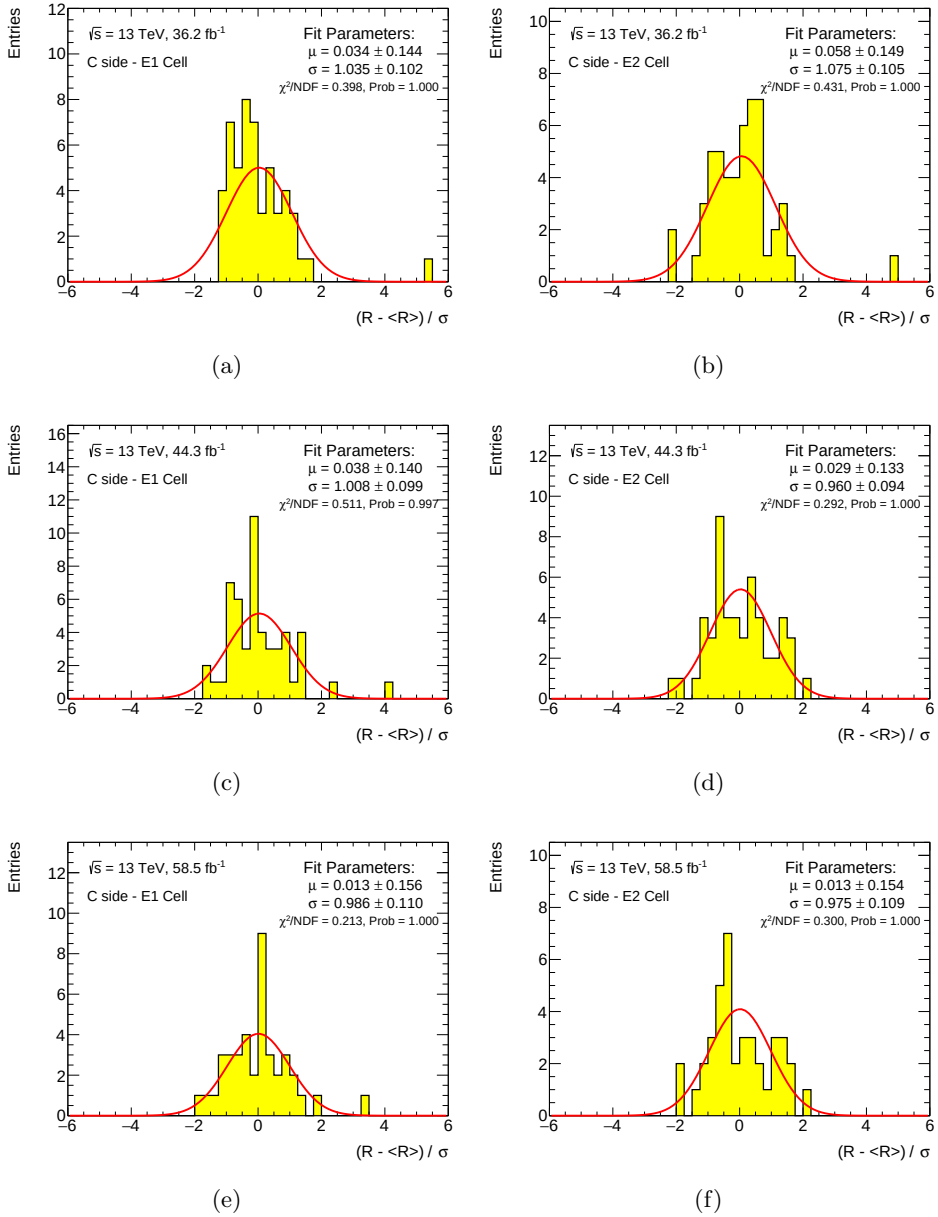


Figure A.3 Pull distributions obtained using the fitted parameters $\hat{\mu}$ and $\hat{s}$ for gap scintillators. The top, middle, and bottom rows represent the respective data-taking years.

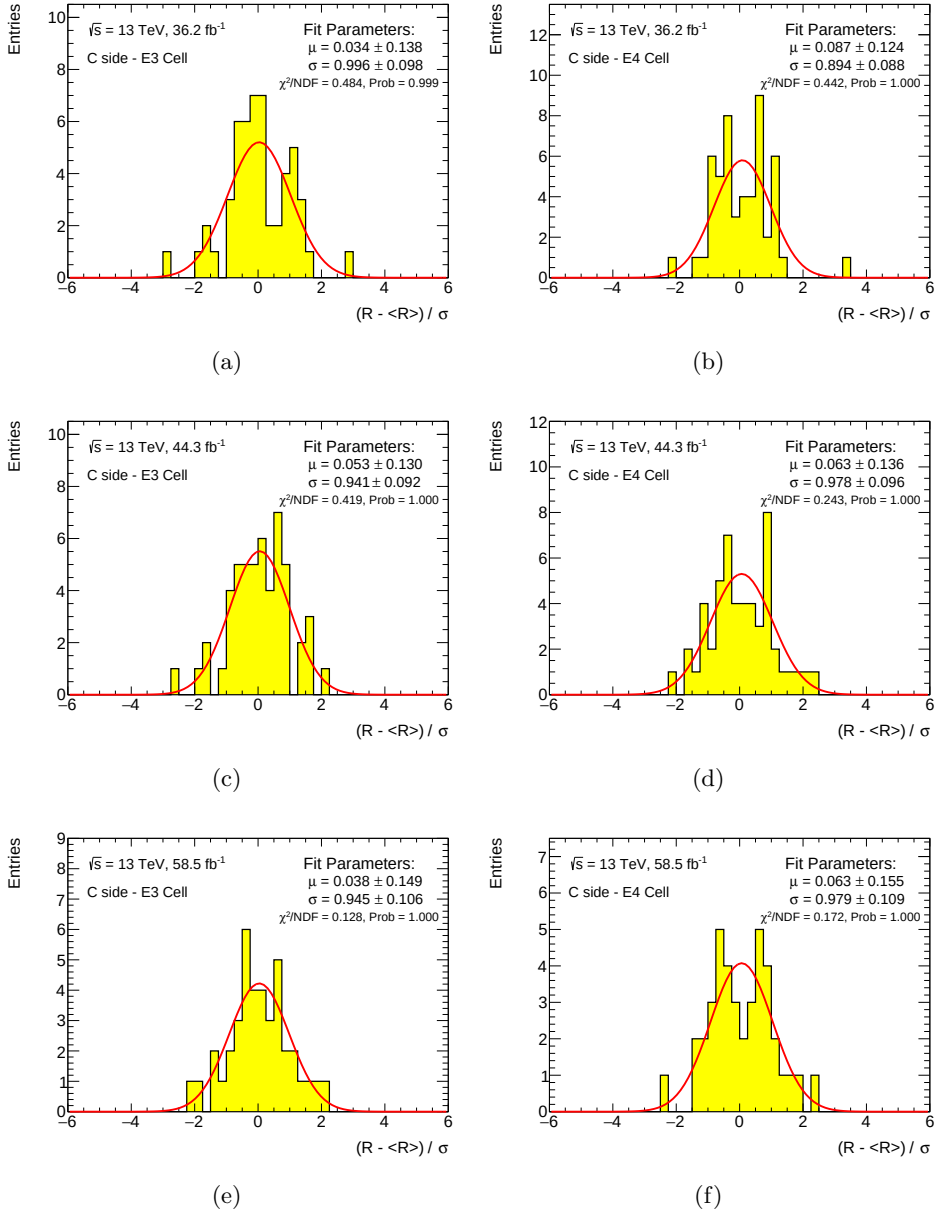


Figure A.4 Pull distributions obtained using the fitted parameters $\hat{\mu}$ and $\hat{s}$ for crack scintillators. The top, middle, and bottom rows represent the respective data-taking years.

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