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Probing the tWb vertex structure in t -channel
single-top-quark production and decay in pp
collisions at 13 TeV with the ATLAS detector

DOCTORAL DISSERTATION

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Certifica:

Que la presente memoria, **Probing the tWb vertex structure in t -channel single-top-quark production and decay in pp collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector**, ha sido realizada bajo su dirección en el Instituto de Física Corpuscular (IFIC), centro mixto de la Universitat de València y del CSIC, por **Mariam Chitishvili**, y constituye su Tesis para optar al título de Doctora por la Universitat de València una vez cursados los estudios en el Doctorado en Física.

Y para que así conste, en cumplimiento de la legislación vigente, presenta en el Departament de Física Atòmica, Molecular i Nuclear de la Universitat de València la referida Tesis Doctoral, y firma el presente certificado.

Valencia, a 17 de julio de 2025



Carlos Escobar Ibáñez

Declaration

This dissertation is the result of my own work, except where explicit reference is made to the work of others, and has not been submitted for another qualification to this or any other university.

Mariam Chitishvili

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Preface

The desire to understand the world around us is one of the most fundamental aspects of being human. For centuries, people have asked questions about the nature of reality—what matter is made of, how the universe works at its most basic level, and what principles lie behind the structure of everything we observe. In modern physics, these questions lead us to study the smallest building blocks of matter and the forces that govern them.

This thesis is a result of that ongoing search for understanding.

Among all known particles, the top quark holds a special place. It is the heaviest elementary particle discovered so far and has unique properties that make it a powerful tool for testing the Standard Model (SM) of particle physics. Investigating the top quark and its interactions, especially in single production processes, can provide important clues about whether the SM is complete or if new physics may be hiding just beyond its limits.

While this research focuses on a specific topic—the structure of the tWb vertex in t -channel single top-quark production—it is part of a much broader effort to connect theoretical predictions with experimental data. It reflects the way particle physics combines precise measurements with complex models to test the foundations of our understanding of nature.

This journey has required not only technical knowledge but also persistence, critical thinking, and a willingness to question assumptions. Throughout the process, I have gained a deeper appreciation for both the complexity of the subject and the collaborative nature of scientific discovery. I hope this work contributes, in its own way, to the collective progress of this field and encourages further exploration of the questions that still remain.

Chapter 1

Theoretical framework

Understanding nature at the smallest scales requires reliable theoretical framework capable of interpreting experimental observations and predicting new phenomena. Among these frameworks, the SM of Particle Physics occupies a central role due to its remarkable precision and successful explanation of various experimental results.

In this chapter, the basics of the SM are described by first presenting the fundamental particles in Section 1.1. Section 1.2 covers the electromagnetic, electroweak (EW), and strong interactions. Lastly, Sections 1.3 and 1.4 discuss the limitations of the SM, emphasising the necessity for theories beyond it in order to describe and understand nature.

1.1 The Standard Model of Particle Physics

The SM of Particle Physics is one of the most successful theories in modern science, providing an accurate description of the fundamental particles and their interactions. Its development began early in the 20th century when physicists like Paul Dirac, Richard Feynman, Julian Schwinger, and Shinichiro Tomonaga established Quantum Electrodynamics (QED) in the 1930s and 1940s [1]. The QED theory combined quantum mechanics and special relativity, significantly improving our understanding of electromagnetic interactions. In the following decades, scientists sought ways trying to unify the known fundamental forces. A major step was taken in the 1960s by Sheldon Glashow, Abdus Salam, and Steven Weinberg, who introduced the EW theory [2–4]. This theory described the electromagnetic and weak forces, and its predictions were later confirmed by the discovery of the W and Z bosons at CERN in 1983 by the UA1 and UA2 experiments [5].

At the same time, theorists were exploring the strong nuclear force, which led to the Quantum Chromodynamics (QCD) in the 1970s [6]. The QCD theory describes quarks and gluons as the fundamental particles responsible for strong interactions. The final key piece of the SM was found in 2012 at the Large Hadron Collider (LHC) [7, 8], when two of its large experiments confirmed independently the existence of the Higgs boson, a particle predicted nearly 50 years earlier [9, 10]. This discovery explained how fundamental particles acquire mass and completed the structure of the SM.

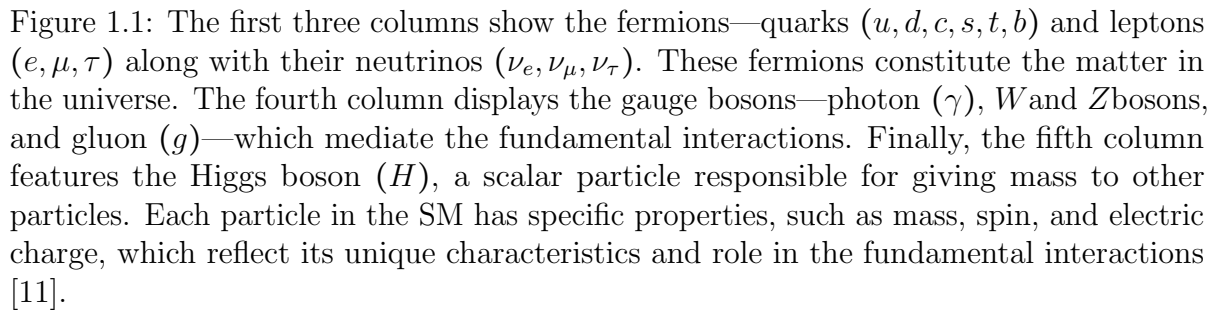
1.1.1 Overview of the Standard Model of Particle Physics

At its core, the SM is based on the concept of gauge symmetry, which describes the idea that the laws of physics should remain invariant even if certain transformations are applied to the particles being studied. In the SM, this symmetry is described by the gauge group

$$SU(3) \otimes SU(2) \otimes U(1), \quad (1.1)$$

which combines the symmetries of the strong force ($SU(3)$), weak force ($SU(2)$), and electromagnetic force ($U(1)$). The SM is a relativistic quantum field theory, where each elementary particle emerges as a localised excitation of a field extending throughout spacetime. These particles are classified into fermions and bosons as shown in Figure 1.1. Fermions, which include quarks and leptons, form the fundamental building blocks of matter. Bosons fall into two main categories: gauge bosons and the Higgs boson. The gauge bosons—photon (γ), W and Z bosons, and gluons—are vector particles that mediate the electromagnetic, weak, and strong interactions, respectively. By contrast, the Higgs boson is a scalar particle that imparts mass to other particles through the Higgs mechanism (discussed in Section 1.1.3) and does not mediate any fundamental force.

Despite its success, the SM leaves several major questions unresolved. For example, the matter–antimatter asymmetry observed in our universe cannot be fully explained by the current understanding of Charge conjugation and Parity (CP) violation — a phenomenon where the laws of physics change slightly when particles are replaced by their antiparticles and their spatial coordinates are inverted — within the SM. A more detailed discussion follows in Section 1.1.5.1. Another open problem is the so-called hierarchy problem, which highlights the large disparity between the EW and higher energy scales such as the Planck scale. Furthermore, the SM does not provide a complete unification of the electromagnetic, weak, and strong forces, and not including gravity leaving room for more comprehensive theories that can address these lingering issues.



The gauge principle serves as the foundation for theoretical models of three fundamental interactions in particle physics: strong, weak, and electromagnetic interactions involving quarks, leptons, and the related force-carrying gauge bosons. Gauge invariance is a fundamental notion in physics that describes a physical invariance under a local transformation of its gauge fields. Depending on the nature of the symmetry group, gauge symmetry can be categorised as Abelian or non-Abelian. The symmetry group in Abelian gauge theories is commutative, which means that the order of symmetry transformations has no effect on the output. Non-Abelian gauge theories, on the other hand, feature a non-commutative symmetry group, indicating that the order of symmetry transformations is important.

The free Dirac Lagrangian exhibits a global $U(1)$ symmetry. This symmetry can be

expressed through the transformation:

$$\psi \rightarrow e^{i\alpha}\psi, \quad (1.2)$$

where ψ is the Dirac field and α is a constant phase. The Lagrangian remains invariant under this global $U(1)$ transformation.

In order to extend the $U(1)$ symmetry from global to local, the phase factor α is promoted to a function of spacetime, $\alpha = \alpha(x^\mu)$, where $x^\mu, \mu \in [0, 1, 2, 3]$, are the spacetime coordinates. This step aims to preserve the invariance of the theory under a local $U(1)$ transformation, a procedure commonly known as *gauging the symmetry*.

However, applying this local transformation to the free Dirac Lagrangian,

$$\mathcal{L} = \bar{\psi}(i\gamma^\mu\partial_\mu - m)\psi \rightarrow \bar{\psi}e^{-i\alpha(x)}(i\gamma^\mu\partial_\mu - m)e^{i\alpha(x)}\psi, \quad (1.3)$$

reveals a problem. Here, ψ denotes the Dirac spinor field representing a fermion, and $\bar{\psi} = \psi^\dagger\gamma^0$ is its Dirac adjoint. The matrices γ^μ are the Dirac gamma matrices satisfying the Clifford algebra, ∂_μ is the partial derivative with respect to the spacetime coordinate x^μ , and m is the mass of the fermion. When this expression is expanded, it produces an additional term that depends on the derivative of the transformation parameter:

$$\mathcal{L} \rightarrow \bar{\psi}[i\gamma^\mu\partial_\mu - m - \gamma^\mu(\partial_\mu\alpha(x))]\psi. \quad (1.4)$$

This extra term, $\bar{\psi}\gamma^\mu\psi\partial_\mu\alpha(x)$, spoils the invariance of the Lagrangian under the local transformation. In other words, the free Dirac theory is not invariant under local $U(1)$ transformations because the derivative ∂_μ acts on both the field ψ and the now spacetime-dependent phase $\alpha(x)$.

To restore local gauge invariance, new field $A_\mu(x)$ is introduced, referred to as the *gauge field*. It is designed to transform in such a way that it cancels the unwanted term:

$$A_\mu(x) \rightarrow A_\mu(x) - \frac{1}{q}\partial_\mu\alpha(x), \quad (1.5)$$

where q is a constant that will later be identified as the electric charge of particle. Using this gauge field, the *covariant derivative* is defined:

$$D_\mu \equiv \partial_\mu + iqA_\mu. \quad (1.6)$$

The term ‘‘covariant’’ here implies that this derivative transforms in the same way as the field ψ under a local $U(1)$ transformation, thereby preserving the overall invariance of

the theory.

Mathematically, saying that a particle “carries charge” means its dynamics involve this covariant derivative with a nonzero q . If a Lagrangian of the field involves only the ordinary derivative ∂_μ (i.e., $q = 0$), then the field is neutral with respect to the $U(1)$ symmetry and does not interact with the gauge field A_μ . This idea will be further explored in Section 1.1.4, where the charges and interactions of quarks are discussed in the context of gauge theory.

$$\begin{aligned}\mathcal{L} = & \bar{\psi} e^{-i\alpha(x)} \left\{ i\gamma^\mu \partial_\mu - m - q\gamma^\mu \left[A_\mu - q\partial_\mu \alpha(x) \right] \right\} e^{i\alpha(x)} \psi = \\ & = \bar{\psi} (i\gamma^\mu D_\mu - m) \psi.\end{aligned}\tag{1.7}$$

It is important to note that the Lagrangian remains invariant not only under the local $U(1)$ transformation but also under the original global $U(1)$ symmetry. This demonstrates the successful restoration of the $U(1)$ symmetry through the addition of the field A_μ .

Now, without introducing the dynamics of the field, the obtained Lagrangian in the Euler–Lagrange equations

$$\frac{\partial \mathcal{L}}{\partial A^\mu} = -q\bar{\psi}\gamma^\mu\psi = 0,\tag{1.8}$$

would give

$$j^\mu = 0.\tag{1.9}$$

Here, $j^\mu = q\bar{\psi}\gamma^\mu\psi$ is the fermionic current associated with the local $U(1)$ gauge symmetry, representing the flow of electric charge carried by the field ψ .

This result leads to the notion that a physical field without dynamics serves as a fixed, unchanging background, lacking any active role in the system. In this regard, such a field remains constant throughout spacetime and does not contribute to the evolution or interactions within the system.

Therefore, the proper Lagrangian would look like this:

$$\mathcal{L} = \bar{\psi} (i\gamma^\mu D_\mu - m) \psi - \frac{1}{4} F^{\mu\nu} F_{\mu\nu} - j^\mu A_\mu\tag{1.10}$$

Here, $F^{\mu\nu} = \frac{1}{q} [D^\mu, D^\nu] = \partial^\mu A^\nu - \partial^\nu A^\mu$ is the field-strength tensor for the gauge field A^μ . This gauge field is coupled to the fermion field via the covariant derivative and the current term.

Starting from a non-interacting particle and introducing the $U(1)$ symmetry ensures that the original particle remains and also reveals the presence of electromagnetism within the theory. The inclusion of the A_μ field, which manifests as the photon upon quantisation, arises naturally from the $U(1)$ symmetry.

Extending the concept of gauge symmetry beyond $U(1)$ allows the analysis of particle behaviour in the EW and strong interaction contexts. Generalising the gauge theory framework to include $SU(2)$ for the EW force and $SU(3)$ for the strong force provides a unified description of gauge bosons and their interactions with fermions.

1.1.3 Spontaneous symmetry breaking in gauge theories

The gauging of symmetries in the realm of gauge theories unveils profound insights into the fundamental properties of particles and their interactions. One intriguing aspect within this framework is the phenomenon of the Spontaneous Symmetry Breaking (SSB) [12, 13], where underlying symmetries become hidden in the observable physical universe governed by gauge fields.

The Goldstone theorem [14, 15] is a significant consequence of the SSB. It states that with the spontaneous breakdown of a global continuous symmetry, massless bosons known as "Goldstone bosons" emerge. The existence of these massless particles initially posed a challenge as they had not been experimentally observed. However, within the framework of gauge theories, a deeper understanding and reconciliation between the presence of massless particles and experimental findings is achieved.

This is where the Higgs mechanism comes into play. It was independently proposed by Peter Higgs [12, 16], François Englert and Robert Brout [17], and Gerald Guralnik, Carl Hagen, and Tom Kibble [18], the Higgs mechanism provides a profound insight into the dynamics of the SSB in gauge theories. By promoting the internal symmetry to a local one, this mechanism "consumes" the Goldstone bosons and endows the gauge bosons with mass, preserving the gauge invariance of the theory. It is through this mechanism that the massless gauge bosons predicted by Yang–Mills theories [19] are transformed into observable particles with mass.

Furthermore, in certain gauge theories such as the unified model of EW interactions, the SSB responsible for the masses of gauge vector bosons arises due to the vacuum expectation value (VEV) of scalar fields, known as Higgs fields (ϕ). These scalar fields introduce a dynamical mechanism where the symmetries are spontaneously broken, resulting in the acquisition of masses by the gauge bosons. Visualisation of the Higgs potential and SSB is shown on the Figure 1.2.

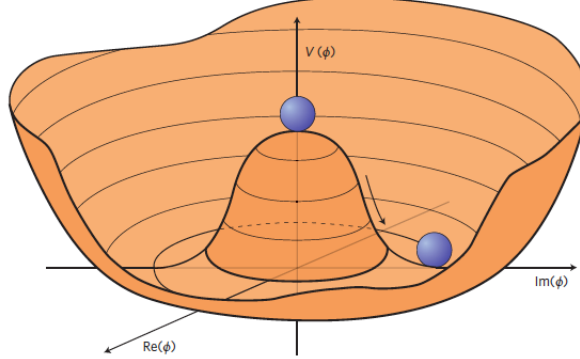


Figure 1.2: Visualisation of SSB for the scalar field (ϕ). Fluctuations around the VEV (global minimum of the potential $V(\phi)$) correspond to the physical Higgs boson, which endows certain particles with mass. The shape of the potential ensures that the true vacuum lies away from the origin [20].

The interplay between the SSB, gauge invariance, and the emergence of mass reveals the intricate relationship between fundamental symmetries and the physical world.

1.1.4 Quarks and leptons

Quarks and leptons are spin- $\frac{1}{2}$ fields in the SM, organised according to the gauge group expressed in Equation 1.1. Quarks carry colour charge (c) and transform as triplets under $SU(3)_C$, while leptons remain colour singlets [21]. The subscript C denotes that this symmetry group governs the strong interaction, with "colour" referring to the charge associated with the QCD. Each of the three generations of fermions consists of two quarks and two leptons:

- **Quarks:** Up-type quarks (up (u), charm (c), top (t)) have electric charge $+\frac{2}{3}e$, where e is the elementary charge, while down-type quarks (down (d), strange (s), bottom (b)) carry charge $-\frac{1}{3}e$. All quarks possess colour charge, which enables them to interact via the strong force and form hadrons, such as protons and neutrons [22].
- **Leptons:** Each generation includes one charged lepton (electron (e), muon (μ), tau (τ)) with electric charge $-1e$, and one neutral lepton — the corresponding neutrino: electron neutrino (ν_e), muon neutrino (ν_μ), and tau neutrino (ν_τ), all of which carry no electric charge. Leptons do not interact via the strong force [21].

Beyond their colour properties (in the case of quarks), these fermions are charged under the EW sector

$$SU(2)_L \times U(1)_Y,$$

where $SU(2)_L$ describes the weak isospin symmetry (L stands for left) and $U(1)_Y$ corresponds to the hypercharge (Y) [3, 4]. Weak isospin is an internal quantum number, denoted by T (with its third component T_3), that organizes left-handed fermions into doublets under $SU(2)_L$. Right-handed components transform as singlets (i.e., they do not carry weak isospin). Hypercharge (or the Y) is the Abelian charge associated with $U(1)_Y$.

The electric charge q is given by $q = T_3 + \frac{Y}{2}$, so that as discussed above up-type quarks (u, c, t) have $+\frac{2}{3}e$, down-type quarks (d, s, b) have $-\frac{1}{3}e$, charged leptons have $-1e$, and neutrinos have $0e$ [2].

All quarks and leptons can be grouped into three generations as already mentioned and they can be seen on Figure 1.1:

$$(u, d), \quad (c, s), \quad (t, b) \quad \text{for quarks;}$$

$$(\nu_e, e), \quad (\nu_\mu, \mu), \quad (\nu_\tau, \tau) \quad \text{for leptons.}$$

They share identical gauge interactions but differ in mass. Higher-mass quarks (e.g. t and b) and leptons (e.g. τ) undergo decays into lighter generations, which explains the stability of the first generation [21].

In the SM, neutrinos are massless, however, the observation of neutrino oscillations [23, 24] demonstrates non-zero neutrino masses, indicating the necessity for physics beyond the SM (BSM). Mechanisms such as the seesaw mechanism have been proposed to explain how neutrinos acquire small but finite masses [25, 26].

On the one hand, quarks exhibit flavour mixing, described by the Cabibbo–Kobayashi–Maskawa (CKM) matrix [27, 28], which parameterises transitions between quark flavours in weak decays (e.g. $d \rightarrow u$) subject to unitarity constraints. On the other hand, although leptons were previously thought not to mix, the discovery of neutrino oscillations led to the introduction of the Pontecorvo–Maki–Nakagawa–Sakata (PMNS) matrix, which encodes flavour transitions in the lepton sector [29, 30].

Fundamental puzzles remain, including the exact origin of the mass hierarchy among quarks and leptons, as well as the precise mechanism for generating neutrino masses.

1.1.5 CKM matrix

A key feature emerging from this generational structure of quarks is the mixing of quark flavours under weak interactions, encoded in the CKM matrix. This mixing arises from the Yukawa sector of the SM, which must be extended to include interactions among

all three quark generations. To accommodate the observed generational structure and flavour transitions, the Yukawa Lagrangian must be generalised, leading to the form:

$$\mathcal{L} = -Y_{ij}^d \bar{Q}_{Li} \phi d_{Rj} - Y_{ij}^u \bar{Q}_{Li} \epsilon \phi^* u_{Rj},$$

Where Y_{ij}^d and Y_{ij}^u are the 3x3 complex matrices representing the Yukawa couplings for down-type and up-type quarks, respectively. These matrices determine the mass and the interactions between different generations of quarks. The symbols Q_{Li} and u_{Rj}, d_{Rj} denote the left-handed quark doublets and right-handed up-type and down-type quark singlets, respectively. $\bar{Q}_{Li}$ represents the conjugate. The ϕ is the Higgs field as already defined in Section 1.1.3, and ϵ is the 2x2 antisymmetric tensor involved in the interaction with the up-type quarks.

A key consequence of this formulation is that the mass eigenstates (i.e. physical states) of the quarks do not align with the gauge eigenstates (i.e. states defined by the gauge symmetry of the SM). This misalignment allows for the possibility of quarks changing generations through interactions with the $W^\pm$ bosons, described by the CKM matrix [31].

The CKM matrix, V , can be expressed as:

$$V = \begin{pmatrix} V_{ud} & V_{us} & V_{ub} \\ V_{cd} & V_{cs} & V_{cb} \\ V_{td} & V_{ts} & V_{tb} \end{pmatrix} = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta_{13}} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta_{13}} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta_{13}} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta_{13}} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\delta_{13}} & c_{23}c_{13} \end{pmatrix}, \quad (1.11)$$

where V is a unitary matrix that describes the mixing between different quark generations. The matrix elements V_{ij} , with i, j representing the quark generations $[(u, c, t), (d, s, b)]$, are determined by diagonalising the Yukawa coupling matrices Y^u and Y^d . The unitarity of the CKM matrix, which requires that its columns and rows form orthonormal vectors, imposes six constraints on its elements. However, these constraints do not fully determine the values of the matrix, making their experimental measurement necessary.

Here in the expression $c_{ij} = \cos \theta_{ij}$ and $s_{ij} = \sin \theta_{ij}$, with i, j representing the quark generations $[(u, c, t), (d, s, b)]$. The phase δ_{13} is the sole complex phase in the SM, making it the only source of CP violation. Indeed, this led to the hypothesis of a third generation of quarks (t, b) [32], to explain CP violating phenomena.

1.1.5.1 CP invariance

The CP violation [33] is a fundamental phenomenon in particle physics that refers to the violation of the combined symmetry of charge conjugation and parity. Charge conjugation

is the transformation of a particle into its antiparticle, while parity involves flipping the spatial coordinates (essentially mirroring the system). When these two symmetries are applied together, certain interactions are expected to produce identical outcomes for particles and their antiparticles if CP symmetry holds. However, experiments have shown that this symmetry is not always preserved, particularly in weak force interactions, indicating CP violation.

The CKM matrix contains physically significant complex phase, that cannot be eliminated through any redefinition of the phases of quark fields, setting it apart as a fundamental aspect of the matrix. The presence of this complex phase is what underlies CP-violating phenomena within the quark sector. An illustrative example of this is in the decay rates of certain mesons, which are composite particles made of a quark and an antiquark. When these mesons decay through processes involving quarks, the rate at which these decays occur can differ significantly from the decay rates involving their corresponding antiquarks. This disparity in decay rates is a direct manifestation of CP violation, highlighting the profound impact of the complex structure of the CKM matrix.

1.2 Interactions

Particle interactions are fundamental processes that govern the behaviour of elementary particles. These interactions can be classified into different types based on the fundamental forces involved. In the scope of SM the electromagnetic, weak, and strong interactions are included in the following section, since gravity is not described by the SM.

1.2.1 Electromagnetic interactions

This section focuses on the phenomenological and experimental aspects of electromagnetism without repeating the gauge theory details covered in Section 1.1.2.

Electromagnetic interactions govern the forces between electrically charged particles through the exchange of photons. In its modern formulation, QED unifies quantum mechanics with special relativity to provide a highly accurate description of these interactions [34–36]. Charged fermions (e.g. electrons, muons, quarks) and photons are treated as quantum fields, allowing precise calculations of scattering amplitudes and cross-sections in perturbation theory [37].

A distinctive feature of electromagnetic interactions is their long-range nature, which arises from the masslessness of the photon. The electromagnetic potential between static

charges follows an inverse-square law $V(r) \propto \frac{1}{r}$, making it possible for charges to influence each other over macroscopic distances [38]. This property underlies common phenomena, such as the formation of atoms and molecules, the behaviour of matter in electric and magnetic fields, and the emission/absorption of electromagnetic radiation.

The QED theory stands as one of the most precisely tested areas of the SM. Observables like the anomalous magnetic moment of the electron and the Lamb shift ¹ in atomic energy levels have been measured and calculated to extraordinary precision, revealing an excellent agreement between theory and experiment [39]. These high-precision tests serve not only to validate the foundations of QED but also to constrain potential new physics BSM.

Electromagnetic processes play a critical role in collider experiments, as they dominate event topologies involving charged leptons and quarks. Photon exchanges mediate final-state radiation and contribute to background processes in searches for new particles. Accurate theoretical predictions for electromagnetic interactions are essential for distinguishing possible signals of physics BSM from the known QED background.

1.2.2 Electroweak interaction

The weak interaction [3] is one of the fundamental forces in nature, alongside electromagnetic and strong interactions. The EW theory, which combines the electromagnetic and weak forces into a unified theoretical framework, describes the EW interaction. It is formulated using quantum field theory and is based on the idea of local gauge symmetry. The EW theory is described by the Lagrangian density, $\mathcal{L}_{EW}$ which includes terms for the kinetic energy of the gauge bosons, the coupling of the gauge bosons to fermions, $\mathcal{L}_{\text{gauge}}$, and the Higgs mechanism $\mathcal{L}_{\text{Higgs}}$.

$$\mathcal{L}_{EW} = \mathcal{L}_{\text{gauge}} + \mathcal{L}_{\text{Higgs}}.$$

The construction of both terms enforces the $SU(2) \otimes SU(1)$ gauge symmetry and requires renormalizability. However, once the Higgs field acquires a VEV, SSB occurs, yielding non-zero masses for the vector bosons and fermions.

The gauge sector: The Lagrangian $\mathcal{L}_{\text{gauge}}$ incorporates gauge bosons and fermions,

¹The Lamb shift refers to a small energy difference between two hydrogen energy levels that were expected to be the same, but are not due to subtle effects from the interaction between the electron and the electromagnetic field. Measuring this difference was one of the first clear successes of QED.

based on the general formalism of gauge theories.

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{4}W_{\mu\nu}^a W^{a\mu\nu} - \frac{1}{4}B_{\mu\nu}B^{\mu\nu} + \bar{L}i\gamma^\mu D_\mu L_i + \bar{R}i\gamma^\mu D_\mu R_i, \quad (1.12)$$

where $W_{\mu\nu}^a$ represents the field strength tensor for the $SU(2)$ gauge bosons, $B_{\mu\nu}$ represents the field strength tensor for the $U(1)$ gauge boson, $\bar{L}_i$ and $\bar{R}_i$ represent the left-handed and right-handed lepton doublets or quark singlets, respectively, and D_μ is the covariant derivative. The terms $\bar{L}i\gamma^\mu D_\mu L_i$ and $\bar{R}i\gamma^\mu D_\mu R_i$ represent the interactions between the fermions and the gauge bosons, respecting the chirality of the left-handed and right-handed fermion fields. Here

$$B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu, \quad W_{\mu\nu}^a = \partial_\mu W_\nu^a - \partial_\nu W_\mu^a + g\epsilon^{abc}W_\mu^b W_\nu^c, \quad (1.13)$$

$$D_\mu = \partial_\mu - ig\frac{\sigma^a}{2}W_\mu^a - ig'\frac{Y}{2}B_\mu. \quad (1.14)$$

In these expressions, g and g' represent the coupling constant for the $SU(2)$ and $U(1)$ gauge groups, respectively, σ^a are the Pauli matrices, W_μ^a and B_μ are the gauge fields corresponding to the $SU(2)$ and $U(1)$ groups respectively, and Y is the hypercharge of the fermion field as already defined in Section 1.1.4. The gauge boson field strengths $B_{\mu\nu}$ and $W_{\mu\nu}^a$, as well as the covariant derivative D_μ , describe the interactions between the gauge bosons and fermions, allowing for the exchange of gauge particles and ensuring gauge symmetry. These terms are essential for understanding the behaviour of the gauge fields in the context of the EW theory.

The Higgs sector: The gauge principle and the need for renormalizability to be true both specify the Higgs Lagrangian as

$$\mathcal{L}_{\text{Higgs}} = (D_\mu\phi)^\dagger(D^\mu\phi) - V(\phi) = (D_\mu\phi)^\dagger(D^\mu\phi) - \mu^2\phi^\dagger\phi - \lambda(\phi^\dagger\phi)^2. \quad (1.15)$$

Here, as defined above ϕ represents the Higgs field, which is a complex scalar field that forms an EW doublet. The covariant derivative D_μ includes the gauge interactions with the W and Z bosons:

$$D_\mu = \partial_\mu + ig\frac{\sigma^a}{2}W_\mu^a + ig'\frac{Y}{2}B_\mu. \quad (1.16)$$

The first term in the Lagrangian is the kinetic term, which describes the derivative interactions of the Higgs field. The second term represents the potential energy of the Higgs field, with μ^2 being the mass parameter and λ determining the strength of the self-interactions.

The gauge boson masses arise through the interaction between the Higgs field and the gauge fields in the Lagrangian of the gauge sector. The term $(D_\mu \phi_L)^\dagger (D^\mu \phi_L)$ includes interactions between the left-handed Higgs field and the $SU(2)$ gauge fields W_μ^a , while $(D_\mu \phi_R)^\dagger (D^\mu \phi_R)$ involves interactions between the right-handed Higgs field and the $U(1)$ gauge field B_μ . These interactions result in the gauge bosons acquiring masses proportional to the magnitude of the Higgs VEV.

On the other hand, fermion masses arise from Yukawa couplings to the Higgs field. Each fermion couples to the Higgs boson with a distinct coupling strength (the Yukawa coupling), which grows larger for heavier fermions. These couplings generate mass terms proportional to the Higgs VEV, such that heavier fermions acquire correspondingly greater masses.

The masses of particles can be expressed in terms of the Higgs VEV (v) and the coupling constants (g and g') as follows:

For the W and Z gauge bosons:

$$m_W = \frac{1}{2}gv, \quad m_Z = \frac{1}{2}\sqrt{g^2 + g'^2}v. \quad (1.17)$$

For fermions:

$$m_f = \frac{y_f}{\sqrt{2}}v, \quad (1.18)$$

where m_W and m_Z are the masses of the $W^\pm$ and Z bosons, respectively, m_f is the mass of the fermion f , and y_f is the Yukawa coupling constant associated with the fermion.

The EW theory makes accurate predictions for many observables, including the masses of the W and Z bosons and their decay branching fractions [40]. Furthermore, the mixing of $SU(2)_L$ and $U(1)_Y$ gauge fields gives rise to the electromagnetic interaction once the Higgs field acquires a non-zero VEV, linking the gauge couplings g and g' to the electric charge e via the well-known relation

$$e = \frac{g g'}{\sqrt{g^2 + g'^2}}.$$

Precision measurements of these parameters, along with higher-order corrections, form a key part of EW precision tests, which constrain physics BSM [37, 41].

1.2.3 Strong interaction

The strong interaction is described by QCD, a non-Abelian gauge theory based on the symmetry group $SU(3)$ [22]. Unlike the electromagnetic force, which couples to electric charge, the strong force couples to a colour charge. Quarks carry colour charge and interact by exchanging gluons, which are the gauge bosons of QCD. Each gluon not only couples to quarks but also carries colour charge itself, leading to self-interactions among gluons.

The QCD Lagrangian can be expressed as

$$\mathcal{L}_{\text{QCD}} = -\frac{1}{4}G_{\mu\nu}^a G_a^{\mu\nu} + \sum_q \bar{\psi}_q (i\gamma^\mu D_\mu - m_q)\psi_q, \quad (1.19)$$

where: $G_{\mu\nu}^a = \partial_\mu G_\nu^a - \partial_\nu G_\mu^a + g_s f^{abc} G_\mu^b G_\nu^c$ is the field-strength tensor for the eight gluon fields G_μ^a , with indices a, b, c running from 1 to 8. ψ_q denotes the quark field of flavour q (e.g. u, d, s, c, b, t). $D_\mu = \partial_\mu - i g_s T^a G_\mu^a$ is the covariant derivative in the fundamental representation of $SU(3)$, where g_s is the strong coupling constant, and T^a are the $SU(3)$ generators.

The first term describes the kinetic energy and self-interactions of the gluons, while the second term accounts for the kinetic and mass terms of the quarks.

The strong interaction confines quarks into colour-singlet bound states - hadrons. Protons and neutrons, for example, are baryons composed of three quarks, whereas mesons are quark-antiquark pairs. These hadrons, in turn, form the atomic nucleus. At sufficiently high energies, quarks behave nearly freely inside hadrons (due to asymptotic freedom), while at low energies, the strength of the interaction increases, confining quarks permanently within hadrons [42, 43].

Quark-quark interactions proceed via gluon exchange, reflecting the fact that gluons are massless spin-1 bosons carrying colour charge themselves. Consequently, gluons can couple to one another, making QCD substantially more complex than electrodynamics. This non-Abelian nature leads to phenomena such as asymptotic freedom and confinement. In asymptotic freedom, the effective strong coupling decreases with increasing momentum transfer, rendering quarks weakly interacting at high energies.

A key feature of QCD is asymptotic freedom [42, 43], whereby the theory predicts that quarks behave as nearly free particles at extremely short distances. Conversely, at large distances or low energies, the strong coupling grows, causing colour-charged objects to be confined into hadrons. This dual behaviour underlies the observed bound states of quarks and explains why free quarks are not detected experimentally.

QCD successfully accounts for a wide variety of strong-interaction phenomena, from the internal structure of nucleons to the production of hadronic jets at colliders. The intricate interplay of quark and gluon fields remains an active field of study, with ongoing efforts to understand non-perturbative aspects such as confinement and the quark–gluon plasma formed in high-energy heavy-ion collisions.

1.2.4 The Lagrangian of the Standard Model

The full Lagrangian of the SM, which describes the electromagnetic, weak and strong interactions, can be written as:

$$\mathcal{L}_{\text{SM}} = \mathcal{L}_{\text{EW}} + \mathcal{L}_{\text{Higgs}} + \mathcal{L}_{\text{QCD}} + \mathcal{L}_{\text{Yukawa}}, \quad (1.20)$$

where each term corresponds to a specific sector of the SM discussed in previous sections. The $\mathcal{L}_{\text{Yukawa}}$ term encodes the interactions between the Higgs field and fermions, generating the fermion masses once the Higgs field develops a VEV. It is particularly interesting because it explains the observed mass hierarchy and mixing among fermions. The structure of the Yukawa couplings remains one of the most intriguing puzzles, suggesting that an underlying mechanism or symmetry might be at work.

When combined, these terms form the complete Lagrangian of the SM, providing a comprehensive description of the fundamental particles and theory interactions within the framework of quantum field theory.

Beyond its elegant mathematical structure, the SM is celebrated for its extraordinary success in describing a wide range of natural phenomena. It has accurately predicted the existence of fundamental particles—such as the W and Z bosons, the top quark, and the Higgs boson—and has provided precise descriptions of high-energy processes observed at many experiments at colliders like LEP, the Tevatron, and the LHC. Its predictive power spans many orders of magnitude in energy and distance scales.

1.3 Limitations of the Standard Model

Although the SM provides remarkably accurate predictions for a wide range of particle physics phenomena, several open questions hint at the necessity for physics beyond its boundaries. This section highlights a few limitations relevant both on a general level and in contexts where effective field theory (EFT) analyses can provide important clues. The latter is precisely the case for this thesis, where EFT serves as a framework to explore

potential deviations from SM predictions.

1. Hierarchy Problem. A central issue tied to the large mass of the top quark is the hierarchy problem. The Higgs-boson mass $\sim 125 \text{ GeV}$ appears unnaturally light relative to the Planck scale ($\sim 10^{19} \text{ GeV}$). Loop corrections, particularly those involving the top quark, give large contributions to the Higgs-boson mass, requiring significant fine-tuning to match the observed value [44, 45]. This discrepancy suggests new physics that stabilises the EW scale, such as supersymmetry or other BSM scenarios.

2. Matter–Antimatter Asymmetry. The SM incorporates CP violation, but the amount is insufficient to account for the observed baryon asymmetry of the universe [46, 47]. Since top-quark loops contribute strongly to processes involving CP violation (e.g. in Higgs or EW baryogenesis models), precise measurements of top-quark couplings offer a window into potential new CP-violating effects that could help explain the dominance of matter over antimatter.

3. Neutrino Masses. Additionally, neutrino mass remains as an important puzzle. As mentioned in Section 1.1.4, experiments have established that neutrinos have small but non-zero masses [23, 24], requiring new fields or interactions BSM. Although this challenge does not directly hinge on the top-quark sector, insights from studying large Yukawa couplings in top-quark related processes may help guide extended models that also account for neutrino masses.

4. Ultraviolet (UV) Completions and Effective field theories. The SM does not unify the EW and strong forces with gravity in a single framework, leaving open the question of a fully unified theory (e.g. grand unified theories or string theory). In the interim, EFT methods offer a systematically improvable way to parametrise possible deviations from SM predictions. Such EFT analyses are particularly powerful in top-quark processes, where new physics effects could manifest as small deviations in high-precision measurements.

5. Gravity One of the most notable limitations of the SM is its inability to incorporate gravity. Gravity, described by the general relativity and introduced by Albert Einstein [48] is not accounted in the framework of the SM. To fully understand the behaviour of particles and their interactions at all energy scales, quantum description of gravity is needed to unify gravity with the other fundamental forces.

7. Dark matter and dark energy

Dark matter and dark energy, inferred from gravitational phenomena on galactic and cosmological scales, remain unexplained within the SM. Observations of galaxy rotation curves, gravitational lensing, and the universe’s accelerating expansion collectively indicate

the presence of these components. Understanding their underlying physics represents a major challenge that cannot be resolved solely by the known particle content and interactions of the SM [49].

Each of these limitations motivates ongoing theoretical and experimental research. In many scenarios, the large mass of the top quark and its strong coupling to the Higgs boson provide a unique testing ground for BSM physics. The development of EFT frameworks and high-precision top-quark measurements at colliders are consequently key pillars of the search for new physics [50].

1.4 Effective Field Theories

EFTs [51] provide powerful frameworks for understanding and exploring physical phenomena at energies beyond the direct experimental reach. They enable controlled, model-independent strategies for incorporating potential new-physics effects, represented by higher-dimensional operators that augment the SM Lagrangian. By systematically including these operators, one can address open questions—such as the hierarchy problem, matter–antimatter asymmetry, or the origin of neutrino masses—without committing to a specific UV completion. This approach is particularly valuable in top-quark studies, where higher-order corrections can be sensitive to undiscovered particles or interactions. Through precise measurements and global fits of these EFT operators, it becomes possible to constrain or reveal deviations from the SM in a manner that remains agnostic to the details of any particular BSM model.

1.4.1 Applications of Effective Field Theory in particle physics

One fundamental concept in EFT is the effective Lagrangian $\mathcal{L}_{\text{eff}}$. It describes the relevant degrees of freedom and interactions at energies below a certain high-energy scale Λ , often referred to as the cut-off scale, beyond which new physics is presumed to appear. A common way to write it is:

$$\mathcal{L}_{\text{eff}} = \mathcal{L}_0 + \sum_i \frac{C_i}{\Lambda^{d_i-4}} \mathcal{O}_i, \quad (1.21)$$

where $\mathcal{L}_0$ is the renormalizable Lagrangian of the known low-energy theory (e.g. the SM described by Equation 1.20), $\mathcal{O}_i$ are higher-dimensional operators with mass dimension $d_i > 4$, C_i are called Wilson coefficients, which encode the strength of each operator and Λ is the high-energy scale at which new degrees of freedom are integrated

out.

The factor Λ^{d_i-4} in the denominator ensures that operators of higher dimension are suppressed at energies $E \ll \Lambda$. As E approaches or exceeds Λ , the EFT description breaks down, and a more complete theory involving the heavy (integrated-out) particles becomes necessary.

1.4.1.1 Power counting and operator construction

Power counting provides a systematic way of organising and classifying the contributions of different terms in the effective Lagrangian according to their relevance and impact on physical observables. The power counting scheme typically assigns a power of momentum, denoted as p , to each field and derivative in an operator. By counting the powers of p for each term, one can determine its relative importance. For example, a term with a higher power of p is suppressed compared to a term with a lower power.

The power counting dimensions are assigned to various quantities in the Lagrangian. For example: each field χ is assigned a power counting dimension $[\chi]$, each derivative ∂_μ is assigned a power counting dimension $[\partial_\mu] = 1$, and each coupling constant g is assigned a power counting dimension $[g] = 0$.

In scalar field theory, the kinetic term power counting dimension would be:

$$[\mathcal{L}_{\text{kinetic}}] = [\partial_\mu \chi] + [\partial^\mu \chi] = 1 + 1 = 2. \quad (1.22)$$

Operator construction involves formulating effective operators that capture the relevant physics and interactions at a particular energy scale. These operators are constructed based on symmetries, power counting rules, and the desired precision of the EFT. Each operator is assigned a power counting dimension based on the power counting dimensions of the involved fields, derivatives, and coupling constants.

By considering the power counting dimensions, one can identify the dominant contributions at a given energy scale and construct effective operators that accurately describe the physics of the system.

1.4.1.2 Matching and Wilson coefficients

Consider a specific EFT scenario that focuses on studying the interactions between a heavy particle, such as the top quark, and lighter particles within the SM. The EFT Lagrangian contains an operator $\mathcal{O}_i$ that describes these interactions, and its coefficient

is denoted as C_i .

To establish a connection with the underlying theory, such as the full SM, a matching procedure is performed. This procedure involves comparing the EFT predictions with the results obtained from the full theory. In this case, the amplitudes ($\mathcal{A}$) calculated using the EFT Lagrangian are compared with those obtained from the full SM Lagrangian.

The matching equation takes the form:

$$\mathcal{A}_{\text{EFT}} = \mathcal{A}_{\text{SM}}, \quad (1.23)$$

where $\mathcal{A}_{\text{EFT}}$ represents the amplitude calculated using the EFT Lagrangian with the operator $\mathcal{O}_i$ and the Wilson coefficient C_i , and $\mathcal{A}_{\text{SM}}$ represents the amplitude calculated using the full SM Lagrangian.

By equating these amplitudes and performing the necessary calculations, a relationship between the Wilson coefficient C_i [52] and the corresponding couplings or parameters in the full SM is derived. This matching equation permits the determination of the value of C_i based on the established physics described by the SM.

1.4.2 Limitations and challenges of the Effective Field Theory

While EFT is a powerful framework for studying physics beyond the reach of direct experimental observation, it also has certain limitations and challenges to be considered.

1. Energy range: EFT is effective within a certain energy range, beyond which new degrees of freedom or interactions may become important. The validity of EFT relies on the assumption that the energy scales of interest are much lower than the masses of the new particles or the energy scale associated with new interactions. If this assumption is violated, EFT may fail to provide an accurate description of the physical system.

2. Accuracy and precision: EFT provides a systematic and model-independent approach to describe physical phenomena, but it is inherently an approximation. The accuracy and precision of EFT predictions depend on the truncation of the expansion and the inclusion of relevant terms. The challenge lies in determining the appropriate truncation scheme and controlling the uncertainty introduced by neglecting higher-order terms.

3. UV sensitivity: EFT is sensitive to UV physics, which refers to the high-energy behaviour of the underlying theory. In some cases, UV divergences may arise when calculating loop diagrams or higher-order corrections. These divergences indicate the limitations of the EFT and the necessity of including additional degrees of freedom or a

more fundamental theory at high energies.

4. Renormalizability: While EFT can be renormalizable, ensuring that divergences can be consistently absorbed into redefinitions of parameters, not all EFTs possess this property. Renormalizability is crucial for maintaining predictive power and preserving the physical interpretation of the theory. Non-renormalizable EFTs may require the introduction of new physics or a different theoretical framework.

5. Field content and symmetry breaking: EFT assumes a specific field content and symmetry breaking pattern, which may not be the complete description of the underlying theory. The presence of new particles, interactions, or symmetry breaking mechanisms beyond the scope of the chosen EFT can limit its applicability and accuracy.

6. Model dependence: Although EFT aims to be model-independent, the choice of operators, their coefficients, and the truncation of the expansion introduce a certain level of model dependence. Different EFT frameworks or truncation schemes may lead to different results and interpretations, highlighting the need for careful model comparison and assessment.

This chapter has briefly introduced the core concepts of the SM, detailing its particle content, interactions, and the key theoretical underpinnings, as well as highlighting some of its known limitations and the motivation for EFT approaches. Although the SM provides a remarkably successful description of fundamental particles and their interactions, the open questions underscore the necessity to explore physics beyond its current scope. In the following chapter, attention will shift to the experimental techniques employed in this thesis, building upon the theoretical foundations established thus far.

Chapter 2

Top-quark physics: from discovery to multidimensional observables

In the previous chapter, the discussion focused on how the SM establishes a comprehensive yet incomplete framework for describing fundamental particles and their interactions. Building on that foundation, the following chapter presents the fascinating story of the heaviest quark, the top quark, narrative that forms a key pillar in the broader research conducted in this dissertation on the understanding of single top-quark events and their implications for EFTs.

This chapter introduces top-quark physics, beginning with Section 2.1, which covers its history and discovery. Section 2.1.1 examines the top-quark mass, while Section 2.1.2 discusses production and decay. Section 2.2 explores top-quark couplings, and Section 2.3 focuses on angular measurements probing the tWb vertex. Finally, Section 2.4, presents a path toward a complete angular resolution of the tWb vertex.

2.1 History and discovery

The theoretical predictions leading to the discovery of the top-quark originated from the SM framework as already mentioned in Section 1.1.5, which explicitly required a complete third generation of quarks to preserve consistency with the observed CP violation and the quark mixing formalism. This formalism was initially introduced by Nicola Cabibbo in 1963 [27] and later by Makoto Kobayashi and Toshihide Maskawa in 1973 [28]. In their works, it was shown that CP violation could only be accommodated within a renormalizable theory of weak interactions if at least three quark generations existed, thereby implying the existence of a top quark as the up-type partner of the

yet-undiscovered bottom quark. The subsequent discovery of the bottom quark at Fermi National Accelerator Laboratory (Fermilab, Illinois, US) in 1977 [53] strongly supported this theoretical requirement and prompted intensive experimental searches for its counterpart.

In the 1980s, several experiments, including those at Deutsches Elektronen-Synchrotron (DESY, Hamburg, Germany) and CERN, aimed to discover the top quark. However, their failure to observe it indicated that its mass was higher than the energy levels attainable by the available colliders at these laboratories. The discovery of the W and Z bosons at CERN in 1983 [5, 54, 55] provided a further impulse to search for the top quark, as its existence was crucial for the internal consistency of the SM.

A significant breakthrough was achieved with the commissioning of the Tevatron at Fermilab, a proton–antiproton ($p\bar{p}$) collider at $\sqrt{s} = 1.8\text{TeV}$ that became the first facility capable of reaching the energy threshold required for top quark production. The CDF (Collider Detector at Fermilab) and D0 experiments, utilising this collider, announced evidence for the top quark in 1994 [56]. However, it was not until 1995 that both collaborations confirmed its discovery [57]. The top-quark was found to have a mass of approximately 173 GeV [58], establishing it as the heaviest known elementary particle.

The substantial mass of the top quark—comparable to that of a gold atom—highlights its significant coupling to the Higgs field, which is central to EW symmetry breaking and the fundamental structure of the SM [59]. Its large mass implies a Yukawa coupling to the Higgs boson close to unity, $y_t \approx 1$, suggesting a direct and dominant role in the dynamics of EW symmetry breaking [59]. This coupling strength not only underscores the significance of the top quark in the SM but also makes it a sensitive testing ground for potential deviations predicted by BSM theories.

Another distinctive feature of the top-quark is its extremely short lifetime, estimated to be of $\mathcal{O}(5 \times 10^{-25}\text{s})$ [60]. This timescale is significantly shorter than the hadronisation time scale of $\mathcal{O}(10^{-24}\text{s})$ governed by QCD. As a result, the top quark decays via the weak interaction before it can form bound hadronic states, primarily through the channel $t \rightarrow Wb$ with a branching ratio exceeding 99.8% [61]. This unique property allows for direct access to the kinematic and spin information of the top quark through its decay products, without the complications introduced by non-perturbative QCD effects.

Since its discovery, the top quark has become a central focus in precision measurements and searches for new physics. The LHC continues to provide high-statistics datasets that enable detailed studies of its production mechanisms, decay properties, and interactions, offering ongoing insight into the structure and possible extensions of the SM.

2.1.1 Top-quark mass

The SM is defined by a set of 25 fundamental parameters that are not predicted by the theory itself, serve as inputs for calculating observable quantities within the framework [61] and therefore must be determined experimentally. These parameters include the three gauge couplings that govern the strengths of the electromagnetic, weak, and strong interactions (i.e. g, g', g_s); the 12 masses of quarks and fermions, three mixing angles and the CP-violation phase of the CKM matrix three mixing angles and the Dirac CP-violation phase of the PMNS¹ matrix; and the two properties of the Higgs sector, such as the VEV and the Higgs boson mass. The minimal SM contains 19 parameters, since neutrino oscillations are not considered.

Among these parameters, the mass of the top quark occupies a particularly significant role. Owing to its large mass, the top quark significantly impacts the stability of the EW vacuum, the shape of the Higgs potential, and the outcomes of precision EW analyses. Its large Yukawa coupling to the Higgs field makes it especially sensitive to new physics and central to theoretical extrapolations of the SM to high-energy scales. The precise determination of the top-quark mass is thus essential for internal consistency checks of the SM and for constraining scenarios beyond it. The most precise determinations of the top-quark mass using LHC Run 2 data were performed independently by the ATLAS and CMS collaborations, the two general-purpose detectors at the LHC at CERN (described in detail in Chapter 3). Using proton–proton (pp) collisions at $\sqrt{s} = 13$ TeV, the ATLAS measurement, considering the all-hadronic decay channel using boosted top quarks, yields a top-quark mass of $m_t = 172.95 \pm 0.53$ GeV [62]. This result together with other results are shown in the Figure 2.1, while CMS for lepton plus jets final state reports $m_t = 171.77 \pm 0.37$ GeV [63]. These results represent the most recent top-quark mass measurements to date within each experiment, highlighting the continued progress in experimental techniques and systematic uncertainty control.

The challenge, however, lies in interpreting these measurements. The mass parameter or the Monte Carlo (MC) mass m_t^{MC} , extracted through comparisons between simulated and observed kinematic distributions, is often identified with the pole mass, m_t^{pole} , a widely used renormalisation scheme in perturbative QCD. However, this identification has been the subject of ongoing debate due to theoretical ambiguities in defining the top-quark mass in quantum field theory and the limitations inherent in current MC event generators [64]. Central to this debate is the so-called "interpretation problem", which concerns the precise relation between m_t^{MC} and field-theoretically well-defined mass schemes such as the pole mass. This issue is significant, as it involves the influence of

¹The PMNS matrix is the unitary matrix relating neutrino flavour eigenstates to mass eigenstates.

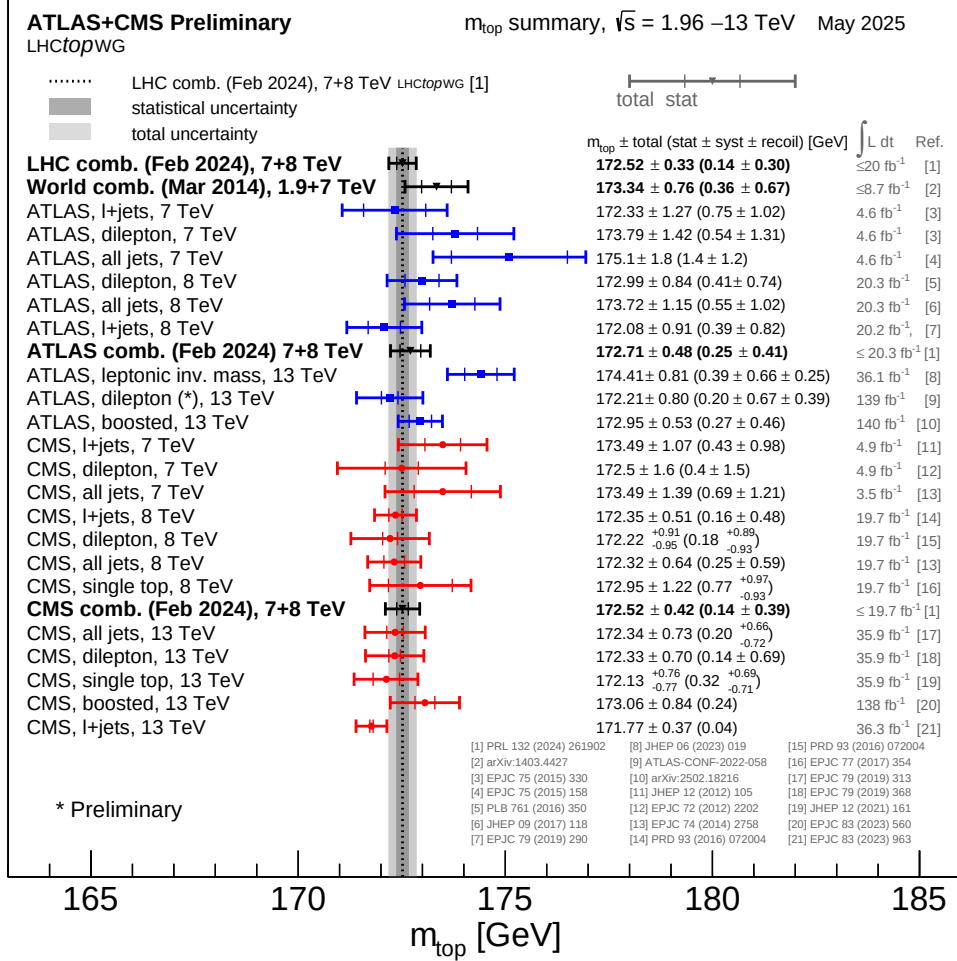


Figure 2.1: This summary presents the direct top-quark mass measurements from ATLAS and CMS, obtained through analyses of top-quark decay channels. The results are shown alongside the LHC and Tevatron+LHC combined top-quark mass determinations for comparison. Each measurement separately lists the statistical uncertainty—which includes contributions from the jet energy scale factor and the b -jet scale factor, where applicable—and the combined systematic uncertainty from all other sources. Additionally, the two most recent ATLAS results include an extra uncertainty component to account for ambiguities in modelling the recoil from parton shower emissions following top-quark decay.

QCD and EW quantum fluctuations at low-energy scales, which can substantially affect the extracted top-quark mass if not properly accounted for.

Addressing the interpretation problem associated with the top-quark mass parameter extracted from MC simulations requires substantial theoretical advancement. Fixed-order perturbative calculations, though systematically improvable, are insufficient to fully describe the complexity of top-quark production and decay, particularly the non-perturbative dynamics of colour reconnection and hadronisation. A reliable interpretation

is expected to emerge only with the development of next-to-leading logarithmic (NLL) accurate parton shower (PS) algorithms and matrix-element (ME) matched MC event generators². In order to robustly relate the reconstructed mass parameter m_t^{MC} to a field-theoretically well-defined quantity, theoretical tools must account for both perturbative corrections and realistic modelling of nonperturbative effects, such as colour neutralisation [64–66].

Among the available mass schemes, the pole mass m_t^{pole} has traditionally been used due to its gauge invariance and renormalisation scale independence at all orders in perturbation theory. However, it suffers from an intrinsic non-perturbative ambiguity of order $\Lambda_{\text{QCD}} \sim 200$ MeV, stemming from the asymptotic nature of the perturbative series [67, 68]. This limitation becomes particularly relevant as current experimental determinations of the top-quark mass approach sub-GeV precision [64]. A prominent example of this sensitivity is illustrated in Figure 2.2, which shows the phase structure of the SM vacuum in the m_t^{pole} vs. m_h^{pole} plane. The plot is based on renormalization group evolution of the Higgs quartic coupling λ , whose sign at high energy scales determines the vacuum structure of the SM. Depending on the values of the top-quark and Higgs-boson masses, the SM vacuum can be absolutely stable (green), metastable (yellow), or unstable (red). In the metastable region, quantum tunnelling to a deeper vacuum is possible. The contour ellipses represent the experimentally allowed regions at 1, 2, and 3σ , based on current measurements of m_t and m_h . The SM lies near the boundary between meta-stability and instability, implying that precise knowledge of the top-quark mass plays a crucial role in understanding the long-term fate of the EW vacuum. The dashed lines within the coloured bands indicate theoretical uncertainties from renormalisation group matching conditions, the strong coupling constant, and threshold corrections. The calculations are performed under the assumption that the top-quark mass corresponds to the pole mass, taken as $m_t^{\text{pole}} = 173.1 \pm 0.6$ GeV [69].

Consequently, alternative short-distance mass definitions—such as the $\overline{\text{MS}}$ (modified minimal subtraction) mass and the MSR (mass-subtracted renormalon³-free) mass—are increasingly adopted in precision analyses due to their improved theoretical stability and reduced sensitivity to long-distance QCD effects. The $\overline{\text{MS}}$ mass is a running mass scheme defined within dimensional regularisation, commonly used in high-order perturbative

²Matrix-element-matched MC generators combine fixed-order perturbative ME calculations with PS simulations. This matching ensures an accurate description of both hard, wide-angle radiation (captured by the ME) and soft/collinear emissions (modelled by the PS), avoiding double-counting and improving the reliability of event-level predictions.

³A type of divergence that appears in QCD calculations when summing many terms in a perturbative series, causing an unavoidable uncertainty of about Λ_{QCD} that limits the precision of quantities like the pole mass.

calculations because of its scale dependence and renormalisation-group properties [70]. The MSR mass, on the other hand, is a scale-dependent scheme that interpolates between the pole and $\overline{\text{MS}}$ masses, explicitly constructed to avoid the infrared renormalon ambiguity, making it particularly suitable for precision observables involving low-scale dynamics [71, 72].

A precise and theoretically well-defined determination of the top-quark mass is essential for consistency checks of the SM and for constraining the EW vacuum structure. Since the extraction of the top-quark mass depends critically on the accurate modelling and on the measuring of its production and decay dynamics, the next section focuses on the theoretical and experimental treatment of these processes.

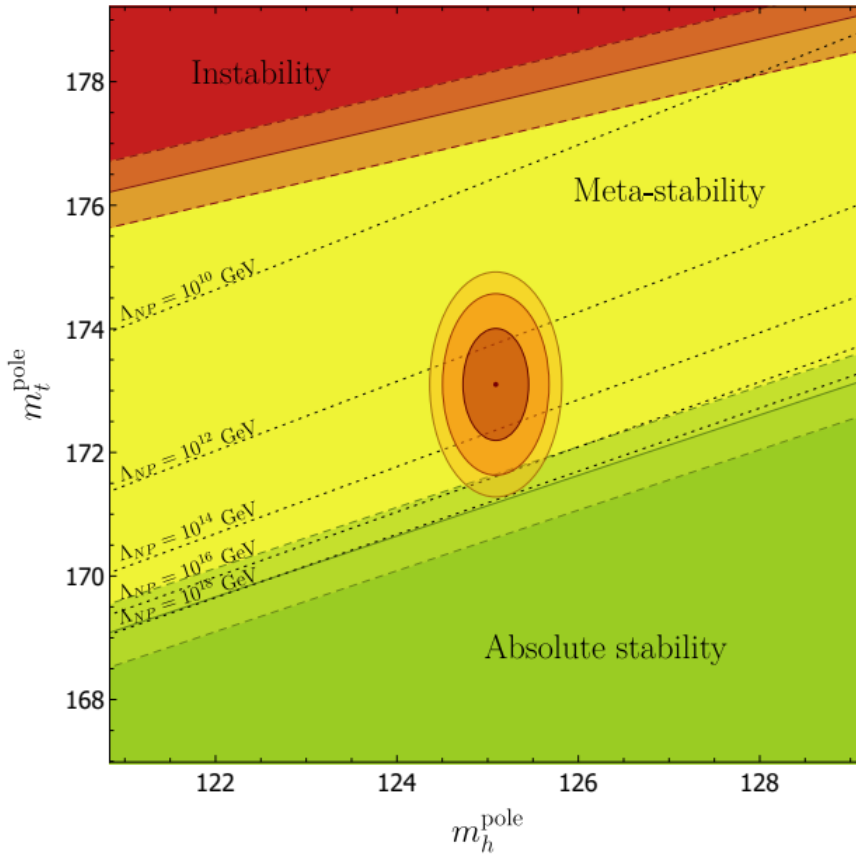


Figure 2.2: Phase diagram showing regions of absolute stability, metastability, and instability of the SM vacuum in the top-quark mass vs. Higgs boson mass plane. The ellipses correspond to the allowed regions at 1, 2, and 3 σ from experimental constraints. Shaded bands along the phase boundaries—centred on solid lines and bounded by dashed lines—reflect uncertainties associated with input parameters and theoretical calculations. The calculations assume a top-quark pole mass of $m_t^{\text{pole}} = 173.1 \pm 0.6$ GeV [69].

2.1.2 Production and decay of the top quark

The production of top quarks is predominantly as two mechanisms: top-quark pair ($t\bar{t}$) production via strong interactions shown, at leading order (LO), in Figure 2.3 and single top-quark production via EW interactions as shown in, at LO, Figure 2.4.

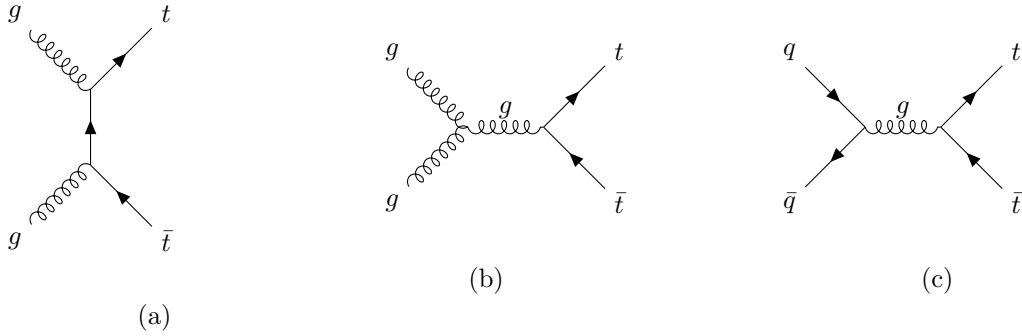


Figure 2.3: Representative Feynman diagrams at LO for the production of top-quark and anti top-quark pairs ($t\bar{t}$). The diagrams (a), (b) illustrate the production via gluon–gluon (gg) fusion, while the diagram (c) is indicative of quark–antiquark ($q\bar{q}$) annihilation.

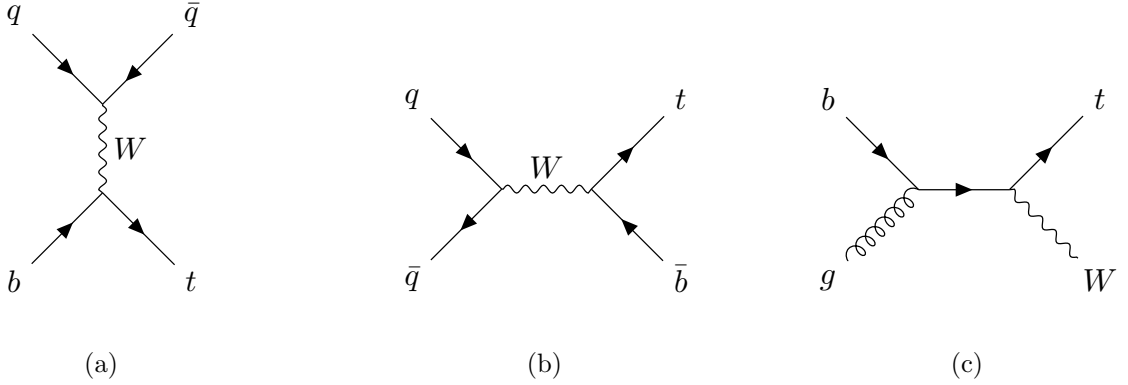


Figure 2.4: Representative Feynman diagrams depicting the LO processes for the production of a single top quark in hadron colliders categorised into three principal prediction mechanisms. These include (a) the t -channel production, (b) the s -channel production, and (c) the associated production of a top quark with a W boson.

The $t\bar{t}$ production is the primary mechanism for top-quark production in high-energy hadron colliders such as the Tevatron and the LHC. As mentioned in the introduction of this chapter, at the Tevatron, collisions occur between protons and antiprotons ($p\bar{p}$), while at the LHC, the collisions are between protons (pp). This production mechanism is primarily mediated by the strong interaction, described by QCD, and proceeds through two main partonic sub-processes:

Gluon fusion ($gg \rightarrow t\bar{t}$): This is the dominant process at the LHC due to the high gluon luminosity in pp collisions, contributing 90% of the total cross-section. It involves the fusion of two gluons, each emitted from an incoming proton, to create a $t\bar{t}$ pair.

Quark-antiquark annihilation ($q\bar{q} \rightarrow t\bar{t}$): This process, which accounted for approximately 85% of the total cross-section at the Tevatron, proceeds through the annihilation of a quark (antiquark) from one proton (antiproton) with an antiquark (quark) from the opposing antiproton (proton) beam, resulting in the production of a $t\bar{t}$ pair.

The most precise experimental measurement of the top-quark pair production cross-section at a centre-of-mass energy of 13 TeV, based on the full Run 2 dataset, yields $\sigma_{t\bar{t}} = 829 \pm 1$ (stat.) ± 13 (syst.) ± 8 (lumi.) ± 2 (beam) pb [73]. Moreover, given that the cross-section depends on the centre-of-mass energy, Figure 2.5 shows the most relevant ATLAS measurements for different energies.

The top-quark predominantly decays into a W boson and a bottom quark ($t \rightarrow Wb$) with a branching ratio (BR)⁴ close to 100%. In the SM, the partial width for the decay can be expressed in terms of the V_{tb} term of the CKM matrix.

The top-quark decay can be further categorised according to W -boson decay modes. For $t\bar{t}$ there are three final states:

1. **Dilepton channel (BR $\approx 10.5\%$):** Both W bosons decay into leptons and their corresponding neutrinos ($W \rightarrow \ell\nu$). This leads to a final state with two leptons, two neutrinos, and two b -quarks. This channel is characterised by a clean signal but has a relatively small BR because of the low leptonic BR of the W boson [61].
2. **Lepton plus jets channel (BR $\approx 43.8\%$):** One W boson decays leptonically ($W \rightarrow \ell\nu$), and the other decays hadronically ($W \rightarrow q\bar{q}'$). The final state consists of one lepton, one neutrino, two b -quarks, and two light-flavour quarks (i.e. u -, d -, c - or s -jets; also named light-jets). This channel offers a good balance between BR and background discrimination.
3. **All-jets channel (BR $\approx 45.7\%$):** Both W bosons decay hadronically ($W \rightarrow q\bar{q}'$). The final state includes two b -quarks and four light jets, leading to a fully hadronic final state. While this channel has the highest BR, it also suffers from a large QCD multijet background.

⁴The BR of a decay channel (e.g. $X \rightarrow Y$) is defined as the fraction of the total decays of a particle that proceed via that specific channel. It is calculated as the partial decay width (Γ) of a given channel divided by the total decay width (Γ_{total}) of the particle: $\text{BR}(X \rightarrow Y) = \Gamma(X \rightarrow Y)/\Gamma_{\text{total}}$. where Γ denotes the decay width.

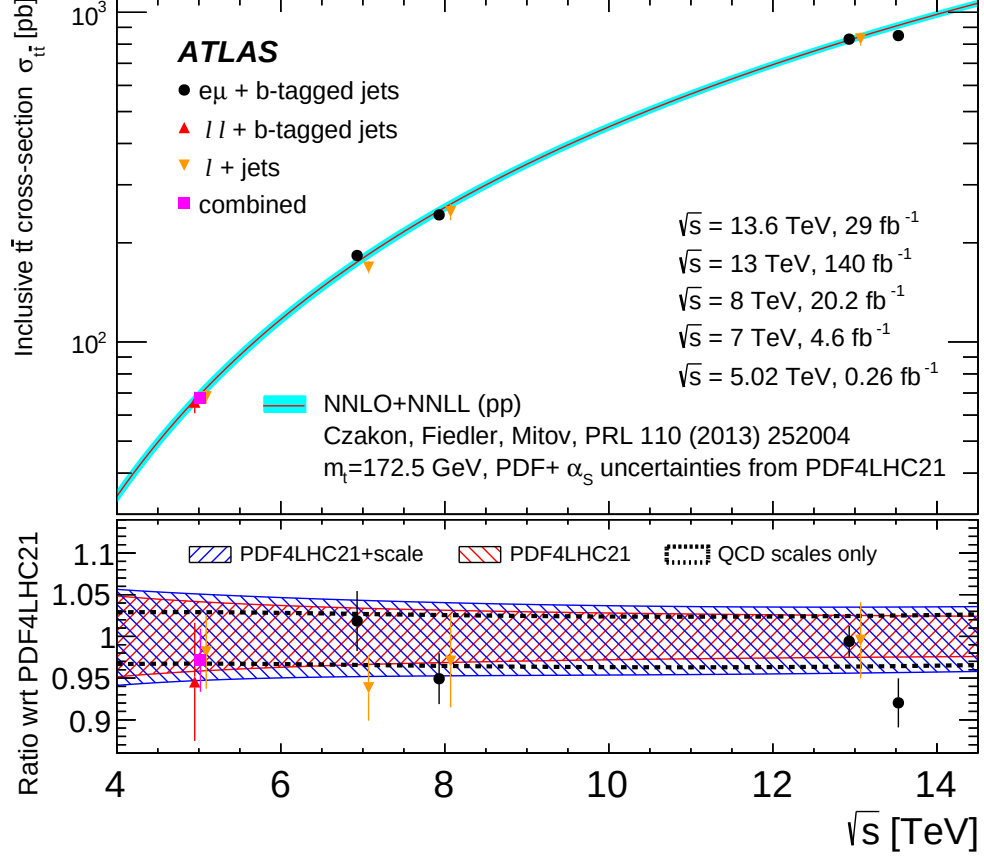


Figure 2.5: Summary of ATLAS measurements of the top-pair production cross-section as a function of the centre-of-mass energy compared to the NNLO QCD calculation complemented with NNLL resummation. The theory band represents uncertainties due to renormalisation and factorisation scales, parton density functions and the strong coupling. The measurements and the theory calculation are quoted at $m_t = 172.5 \text{ GeV}$. Measurements made at the same centre-of-mass energy are slightly offset for clarity [74].

Single top-quark production occurs mainly through three channels as already shown in Figure 2.4, each providing unique insights into the EW interaction. These are the t -channel, the associated production (known as tW -channel), and the s -channel.

t -channel: The most common single top-quark production mode and the focus of the analysis presented in this thesis. In this process, a space-like W boson is exchanged between a light-flavour quark and a bottom quark, resulting in the production of a single top-quark plus another light-flavour quark q' , also known as spectator quark. Theoretical predictions at next-to-leading order (NLO) in QCD give a cross-section of $\sigma_{t\text{-ch.}} = 215^{+2.1}_{-1.3} \text{ pb}$ at $\sqrt{s} = 13 \text{ TeV}$ [75–77]. The cross-section is proportional to $|V_{tb}|^2$ and the square of the EW coupling, highlighting sensitivity of this channel to the EW

interaction and quark mixing.

tW : In this mode, a top quark is produced in association with a W boson. This channel is of particular interest because it involves both EW and QCD interactions, and it interferes with $t\bar{t}$ production at NLO, requiring careful treatment in both theoretical calculations and experimental analyses. The predicted cross-section at $\sqrt{s} = 13$ TeV is $\sigma_{tW} = 71.7 \pm 1.8 \pm 3.4$ pb [75–77].

s -channel: A less frequent process at the LHC, in which a time-like W boson is produced via $q\bar{q}$ annihilation and subsequently decays into a top and a bottom quark. The predicted NLO cross-section at $\sqrt{s} = 13$ TeV is $\sigma_{s\text{-ch.}} = 10.3 \pm 0.29 \pm 0.27$ pb [75–77]. Due to the small signal and large backgrounds, this process has not yet been observed at the LHC with sufficient statistical significance. Like the other channels, it is also sensitive to $|V_{tb}|^2$ and the EW coupling. Figure 2.6 shown the measurements of all three single top-quark production channels as a function of the centre-of-mass energy performed by the ATLAS and CMS experiments. It also shows the combination between independent results for each channel. For all channels and all energies, theoretical predictions are included.

In single top-quark production, the top quark decays almost exclusively via the EW interaction through the tWb vertex, yielding a W boson and a b -quark. This vertex, mediated by the charged current weak interaction and proportional to the CKM matrix element V_{tb} , plays a central role in top-quark dynamics. Its structure and possible deviations from the SM are directly probed in this thesis through the study of single top-quark events and will be discussed in more detail in the following sections.

2.2 Top-quark couplings

At the LHC, the study of top-quark production in association with other particles provides a unique opportunity to probe the EW sector of the SM and to search for effects of physics BSM. NLO corrections are essential for achieving precise theoretical predictions that can be directly compared to experimental observations. These corrections include virtual loop diagrams and real boson emissions, which refine the LO results and improve the accuracy of predictions for cross-sections and kinematic distributions.

The NLO EW couplings of the top-quark to the neutral and charged gauge bosons— Z , W , and the photon γ —as well as to the Higgs boson H and the gluon g , are given below:

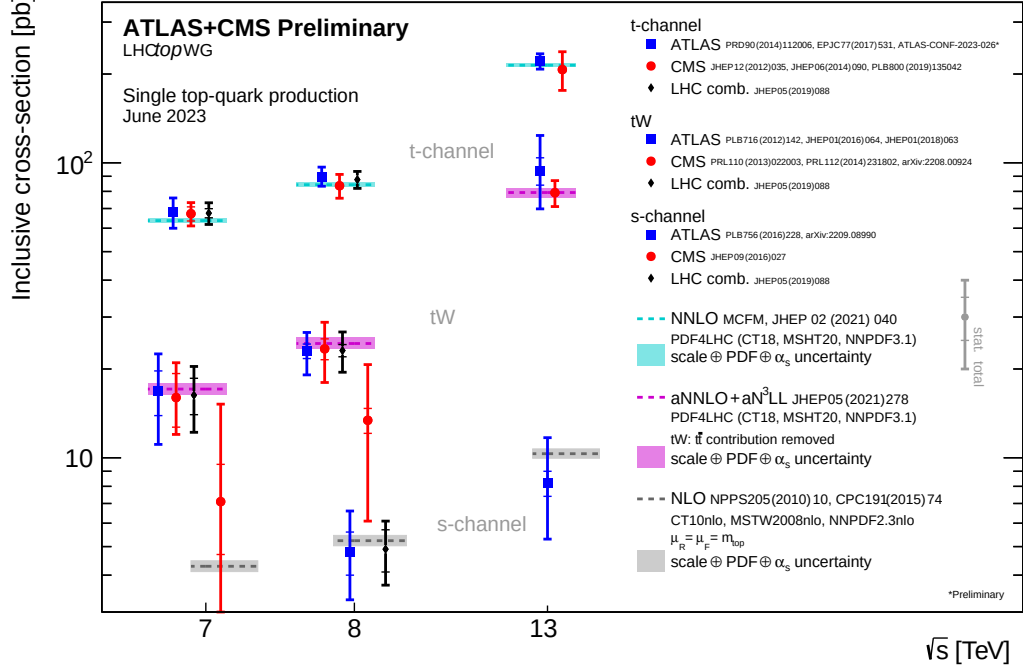


Figure 2.6: The measurements of single top-quark production cross-sections as observed in different channels by the ATLAS and CMS experiments, along with their combined data. These measurements are analysed across various centre-of-mass energy levels. The empirical data are then compared to theoretical predictions derived from several models: NLO QCD enhanced with NNLL resummation, and for the t -channel specifically, NNLO QCD. This comparison offers insights into the consistency and accuracy of theoretical models against experimental observations [74].

- **Top-quark- Z :**

$$\mathcal{L}_{tZ} = \frac{g}{4 \cos \theta_W} \bar{t}_L \gamma^\mu \left(1 - \frac{8}{3} \sin^2 \theta_W - \gamma^5 \right) t_L Z_\mu + \text{h.c.}, \quad (2.1)$$

where θ_W is the weak mixing angle, γ^5 projects onto left-handed chirality, t_L and $\bar{t}_L$ are the left-handed top-quark field and its adjoint, and the term "h.c." denotes the Hermitian conjugate.

- **Top-quark- W :**

$$\mathcal{L}_{tW} = \frac{g}{\sqrt{2}} V_{tb} \bar{b}_L \gamma^\mu t_L W_\mu^- + \text{h.c.} \quad (2.2)$$

where V_{tb} is the CKM matrix element describing the $t \rightarrow b$ transition, $\bar{b}_L$ and t_L are the left-handed bottom and top-quark spinors.

- **Top-quark- γ :**

$$\mathcal{L}_{t\gamma} = e Q_t \bar{t} \gamma^\mu t A_\mu + \text{h.c.}, \quad (2.3)$$

where A_μ is the photon field, and $\bar{t}\gamma^\mu t$ describes the vector current coupling of the top quark to the photon.

- **Top-quark- H :**

$$\mathcal{L}_{tH} = -\frac{y_t}{\sqrt{2}}\bar{t}tH + \text{h.c.}, \quad (2.4)$$

where y_t is the top-quark Yukawa coupling (related to the top-quark mass via $y_t = \sqrt{2}m_t/v$), $\bar{t}t$ represents the scalar current. This coupling not only governs the top-quark mass generation and $t\bar{t}H$ production, but also contributes—via interference with EW diagrams—to the rare tH production process.

- **Top-quark- g :**

$$\mathcal{L}_{tg} = g_s\bar{t}T^a\gamma^\mu tG_\mu^a + \text{h.c.}, \quad (2.5)$$

where g_s is the strong coupling constant, G_μ^a the gluon field, T^a are the SU(3) generators in the fundamental representation, and G_μ^a is the gluon field with colour index $a \in \{1, \dots, 8\}$. This interaction governs the dominant top-quark production mechanism at the LHC through gluon fusion and $q\bar{q}$ annihilation.

Processes such as single top-quark production in association with a Z boson (tZq), a photon ($t\gamma q$), or a Higgs boson (tH), provide direct sensitivity to the EW structure of the top-quark couplings. These channels, through their dependence on both gauge and Yukawa interactions, offer complementary information to QCD-driven production modes and serve as precision probes of the SM at NLO accuracy.

While the SM top-quark interactions are fixed by gauge and Yukawa structure, new physics effects can manifest as deviations from these couplings or the appearance of new ones, such as flavour-changing neutral currents (FCNC) interactions. These can be parametrised using an effective Lagrangian approach:

$$\mathcal{L}_{\text{FCNC}} = \sum_{q=u,c} \left(g_{tqZ}\bar{t}\gamma^\mu (v_{tq} - a_{tq}\gamma^5)qZ_\mu + g_{tq\gamma}\bar{t}\gamma^\mu qA_\mu + g_{tqH}\bar{t}qH \right), \quad (2.6)$$

where g_{tqZ} , $g_{tq\gamma}$, and g_{tqH} represent the effective coupling constants for FCNC interactions.

2.3 Lagrangian formulations and spin dynamics in the tWb vertex

The t -channel single top-quark production — the process studied in this thesis — provides access to the tWb vertex both in the production and decay of the top quark, as illustrated

in Figure 2.7. This feature makes the t -channel particularly suitable for testing SM and BSM theories in the EW sector, although complementary information can also be obtained from the s -channel and tW associated production.

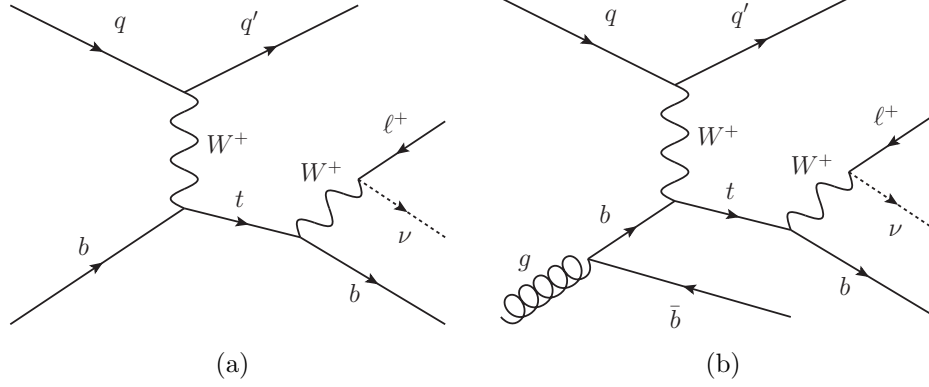


Figure 2.7: Representative LO Feynman diagrams for t -channel single top-quark production and decay in both factorisation schemes. Here q represents a u or $\bar{d}$ quark, and q' represents a d or $\bar{u}$ quark, respectively. The initial b -quark arises from (a) a sea b -quark in the five flavour scheme (5FS) ($2 \rightarrow 2$ process), or (b) a gluon splitting into a $b\bar{b}$ pair in the four flavour scheme (4FS) ($2 \rightarrow 3$ process).

As explained previously in Section 2.1, the large mass of the top quark and its decay before hadronisation make angular measurements particularly valuable, as they offer a cleaner signal for precision studies.

2.3.1 SM and Effective Lagrangian

The tWb vertex can be described in two complementary ways: via the SM CKM picture or through an EFT expansion. In the SM, the structure of the tWb vertex in single top-quark t -channel production and decay is given by the Lagrangian:

$$\mathcal{L}_{\text{SM}} = -\frac{g}{\sqrt{2}} \bar{b} \gamma^\mu \frac{1}{2} (1 - \gamma^5) t W_\mu^- + h.c., \quad (2.7)$$

where g is the weak coupling constant, and $(1 - \gamma^5)$ is the left-handed projection operator.

To investigate possible deviations from the SM, the interaction can be extended using the EFT approach, as introduced in Section 1.4.1. In this framework, new physics effects at an energy scale $\Lambda \gg v$ are parametrised by an expansion in higher-dimensional operators. The effective Lagrangian then takes the form:

$$\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{SM}}^{(4)} + \frac{1}{\Lambda} \sum_i C_i^{(5)} \mathcal{O}_i^{(5)} + \frac{1}{\Lambda^2} \sum_i C_i^{(6)} \mathcal{O}_i^{(6)} + \mathcal{O}\left(\frac{1}{\Lambda^3}\right), \quad (2.8)$$

where $\mathcal{O}_i^{(d)}$ are gauge-invariant operators of mass dimension $d > 4$, and $C_i^{(d)}$ are the corresponding Wilson coefficients. The scale Λ already defined in Section 1.4.1, represents the energy at which new heavy degrees of freedom are integrated out. The suppression by powers of $1/\Lambda$ ensures that the effects of higher-dimensional operators remain small at energies well below Λ , enabling a systematic and model-independent description of potential new physics contributions to the tWb vertex. It is important to note that the numerical bounds on the Wilson coefficients are directly tied to the chosen value of Λ , which is often fixed to 1 TeV by convention. For instance, a bound of $C_i^{(6)}/\Lambda^2 < 1 \text{ TeV}^{-2}$ with $\Lambda = 1 \text{ TeV}$ implies $C_i^{(6)} < 1$, but the same bound would correspond to $C_i^{(6)} < 100$ if $\Lambda = 0.1 \text{ TeV}$, or $C_i^{(6)} < 0.01$ for $\Lambda = 10 \text{ TeV}$. Thus, interpreting these limits requires care, as they reflect assumptions about the scale of new physics that is being integrated out.

At dimension six, there is a significant amount of operators that can be constructed from the SM degrees of freedom [78]. However, many of the possible operators are redundant and can be removed from the basis by using the equations of motion of the fields or Fierz identities. A complete, non-redundant basis exists [79] which consists of 59 independent operators, not counting flavour indices and hermitian conjugates.

To study the interactions of top quarks only a subset of these operators, involving at least one top quark, is needed. Ref. [80] gives a detailed overview of the SM Effective Field Theory (SMEFT) [50] in the top-quark sector, listing all dimension-six operators that couple to at least one top-quark. However, not all of these operators contribute to the processes studied in this analysis. In particular, there are six contributing dimension-six operators, as detailed in Table 2.1. Four of these contribute directly to the tWb vertex which governs both the single-top-quark production and decay. The general form of these operators is given by [79]

$$\begin{aligned}\mathcal{O}_{\varphi q}^{3(ij)} &= \left(\varphi^\dagger i \overleftrightarrow{D}_\mu^I \varphi \right) (\bar{q}_i \gamma^\mu \tau^I q_j), \\ \mathcal{O}_{\varphi ud}^{(ij)} &= (\tilde{\varphi}^\dagger i D_\mu \varphi) (\bar{u}_i \gamma^\mu d_j), \\ \mathcal{O}_{uW}^{(ij)} &= (\bar{q}_i \sigma^{\mu\nu} \tau^I u_j) \tilde{\varphi} W_{\mu\nu}^I, \\ \mathcal{O}_{dW}^{(ij)} &= (\bar{q}_i \sigma^{\mu\nu} \tau^I d_j) \varphi W_{\mu\nu}^I.\end{aligned}\tag{2.9}$$

The operators that contribute specifically to top-quark physics are given by setting $i = 3, j = 3$ to select the operators that involve third generation quarks. Then, writing Equation 2.9 in the physical basis, each of these operators contribute to the tWb Lagrangian as by

$$\begin{aligned}
\mathcal{O}_{\varphi q}^3 &\supset -\frac{gv^2}{\sqrt{2}}\bar{b}_L\gamma^\mu t_L W_\mu^-, \\
\mathcal{O}_{\varphi tb} &\supset \left[\frac{gv^2}{2\sqrt{2}}\bar{b}_R\gamma^\mu t_R W_\mu^- \right]^\dagger, \\
\mathcal{O}_{tW} &\supset 2v\bar{b}_L i\sigma^{\mu\nu} q_\nu t_R W_\mu^-, \\
\mathcal{O}_{bW} &\supset \left[2v\bar{b}_R i\sigma^{\mu\nu} q_\nu t_L W_\mu^- \right]^\dagger,
\end{aligned} \tag{2.10}$$

where the symbol $\supset$ indicates that the following term is a part of the operator on the left-hand side, and the symbol $\dagger$ denotes the Hermitian conjugate of the entire expression inside the brackets.

After EW symmetry breaking, the dimension-six operators give rise to anomalous couplings, leading to a modified tWb interaction Lagrangian of the form:

$$\mathcal{L}_{\text{eff}} = -\frac{g}{\sqrt{2}}\bar{b}\gamma^\mu(V_L P_L + V_R P_R)tW_\mu^- - \frac{g}{\sqrt{2}m_W}\bar{b}\sigma^{\mu\nu}(g_L P_L + g_R P_R)t\partial_\mu W_\nu^- + \text{h.c.} \tag{2.11}$$

where anomalous couplings are expressed in terms of Wilson coefficients as follows:

$$\begin{aligned}
V_L &= V_{tb} + C_{\phi q}^{(3)}\frac{v^2}{\Lambda^2}, \quad V_R = C_{\phi tb}^*\frac{v^2}{2\Lambda^2}, \\
g_L &= C_{bW}^*\frac{\sqrt{2}v^2}{\Lambda^2}, \quad g_R = C_{tW}\frac{\sqrt{2}v^2}{\Lambda^2}.
\end{aligned} \tag{2.12}$$

Where V_L and V_R represent left- and right-handed vector couplings, while g_L and g_R are left- and right-handed tensor couplings. The $P_{L,R} = \frac{1}{2}(1 \mp \gamma^5)$ symbols are the projection operators. In the SM, $V_L \approx V_{tb} \sim 1$, a CKM matrix element, while V_R , g_L , and g_R vanish. These couplings are typically expressed in units of v^2/Λ^2 , where $v = 246.22$ GeV and $\Lambda = 1$ TeV. Deviations from the SM values of these couplings could signal the presence of new physics.

The remaining two operators contributing to the top-quark physics studied in this analysis are two four-quark operators. These operators are given by, assuming minimal flavour violation (MFV),

$$\begin{aligned}
\mathcal{O}_{qq}^{1(ijkl)} &= (\bar{q}_i\gamma_\mu q_j)(\bar{q}_k\gamma^\mu q_l), \\
\mathcal{O}_{qq}^{3(ijkl)} &= (\bar{q}_i\gamma_\mu\tau^a q_j)(\bar{q}_k\gamma^\mu\tau^a q_l).
\end{aligned} \tag{2.13}$$

In the MFV framework, all sources of flavour and CP violation—both within the SM and from possible new physics—are assumed to originate solely from the Yukawa couplings. The flavour structure of any new physics must therefore follow the same hierarchical suppression patterns as in the SM, governed by the CKM matrix. While it is common in some simplified setups to retain only the third-generation Yukawa couplings, such as $\Gamma_u = \text{diag}(0, 0, y_t)$ and $\Gamma_d = \text{diag}(0, 0, y_b)$, the CKM matrix remains non-trivial, and flavour-changing processes are not eliminated but rather constrained to follow the SM structure. The four-quark operators can be rewritten as

$$\begin{aligned} c_{Qq}^{3,1} &= c_{qq}^{3(ii33)} + \frac{1}{6} \left(c_{qq}^{1(i33i)} - c_{qq}^{3(i33i)} \right), \\ c_{Qq}^{3,8} &= c_{qq}^{1(i33i)} - c_{qq}^{3(i33i)}, \end{aligned} \quad (2.14)$$

where the first operator is an SU(2) triplet and colour singlet, while the second is an SU(2) triplet and colour octet. Among the two, only the colour singlet operator interferes with the SM amplitude at leading order due to its matching colour structure. The colour octet contributes only at quadratic level, and thus its effects are more suppressed in observables dominated by interference. While both operators contribute to top-quark pair production and appear in distinct terms in the SMEFT expansion, distinguishing them from data is challenging due to the lack of interference from the colour octet and the limited sensitivity to purely quadratic terms. As a result, this analysis restricts the study to the colour singlet operator $c_{Qq}^{3,1}$, under the assumption that its contribution dominates.

2.3.1.1 Wilson coefficients

The Wilson coefficients [81] emerge from the process of integrating out heavy particles or high-energy degrees of freedom that are not accessible at the energy scales under investigation. This integration yields an effective Lagrangian that captures low-energy phenomena through operators embodying the symmetries and dynamics of the underlying high-energy theory. As introduced in Section 1.4.1 and detailed in Section 1.4.1.1, the values of these coefficients are determined through a matching procedure, wherein the amplitudes calculated in the EFT framework are equated to those of the full underlying theory, such as the SM. This ensures that the low-energy behaviour of the EFT faithfully reproduces the effects of the full theory at LO.

In the context of $t\bar{t}$ production, these coefficients, such as C_{tG} , play a critical role by modifying the gluon-top-quark vertex. Such modifications can influence the production cross-sections, thereby providing potential evidence for new physics within the gluon

sector. Similarly, processes like single-top-quark production, tW , and tZ associated production are sensitive to coefficients like C_{tW} , C_{bW} , and C_{tZ} (see Table 2.1). These coefficients alter the weak interaction vertices, potentially impacting production rates, angular distributions, and the polarisation states of the final particles.

Parameter	tt	single t	tW	tZ	t decay	$t\bar{t}Z$	$t\bar{t}W$
$C_{Qq}^{1,8}$	Λ^{-2}	—	—	—	—	Λ^{-2}	Λ^{-2}
$C_{Qq}^{3,8}$	Λ^{-2}	$\Lambda^{-4} [\Lambda^{-2}]$	—	$\Lambda^{-4} [\Lambda^{-2}]$	$\Lambda^{-4} [\Lambda^{-2}]$	Λ^{-2}	Λ^{-2}
C_{tu}^8, C_{td}^8	Λ^{-2}	—	—	—	—	Λ^{-2}	—
$C_{Qq}^{1,1}$	$\Lambda^{-4} [\Lambda^{-2}]$	—	—	—	—	$\Lambda^{-4} [\Lambda^{-2}]$	$\Lambda^{-4} [\Lambda^{-2}]$
$C_{Qq}^{3,1}$	$\Lambda^{-4} [\Lambda^{-2}]$	Λ^{-2}	—	Λ^{-2}	Λ^{-2}	$\Lambda^{-4} [\Lambda^{-2}]$	$\Lambda^{-4} [\Lambda^{-2}]$
C_{tu}^1, C_{td}^1	$\Lambda^{-4} [\Lambda^{-2}]$	—	—	—	—	$\Lambda^{-4} [\Lambda^{-2}]$	—
C_{Qu}^8, C_{Qd}^8	Λ^{-2}	—	—	—	—	Λ^{-2}	—
C_{tq}^8	Λ^{-2}	—	—	—	—	Λ^{-2}	Λ^{-2}
C_{Qu}^1, C_{Qd}^1	$\Lambda^{-4} [\Lambda^{-2}]$	—	—	—	—	$\Lambda^{-4} [\Lambda^{-2}]$	—
C_{tq}^1	$\Lambda^{-4} [\Lambda^{-2}]$	—	—	—	—	$\Lambda^{-4} [\Lambda^{-2}]$	$\Lambda^{-4} [\Lambda^{-2}]$
C_{Qq}^-	—	—	—	Λ^{-2}	—	Λ^{-2}	—
$C_{\phi Q}^3$	—	Λ^{-2}	Λ^{-2}	Λ^{-2}	Λ^{-2}	—	—
$C_{\phi tb}$	—	—	—	Λ^{-2}	—	Λ^{-2}	—
C_{tZ}	—	—	—	Λ^{-2}	—	Λ^{-2}	—
C_{tW}	—	Λ^{-2}	Λ^{-2}	Λ^{-2}	Λ^{-2}	—	—
C_{bW}	—	Λ^{-4}	Λ^{-4}	Λ^{-4}	Λ^{-4}	—	—
C_{tG}	Λ^{-2}	$[\Lambda^{-2}]$	Λ^{-2}	—	$[\Lambda^{-2}]$	Λ^{-2}	Λ^{-2}

Table 2.1: Sensitivity of various $t\bar{t}$ and single top-quark production processes to dimension-six effective operators in the SMEFT framework. The entries indicate the leading contribution of each operator to the process: linear in the Wilson coefficient C_i (order Λ^{-2}) or quadratic (order Λ^{-4}). Here, Λ denotes the scale of new physics, and C_i are the corresponding Wilson coefficients. Interference terms appear at $\mathcal{O}(\Lambda^{-2})$, while purely quadratic contributions arise at $\mathcal{O}(\Lambda^{-4})$. The entries in square brackets "[]" indicate suppressed contributions that only appear at loop level or through subleading effects, whereas entries without brackets correspond to LO contributions. "The single-t" label includes both t -channel and s -channel production modes [82].

2.3.2 Top-quark polarisation and helicity

Single top-quarks, produced in high-energy particle collisions, exhibit polarisation due to the parity-violating nature of the weak interaction. The polarisation of the top quark ($\vec{P}$) in these processes is predicted as follows: In t -channel production, the polarisation $\vec{P}$ of the single top quark is predominantly aligned with the direction of the exchanged virtual W boson [83], in good agreement with NNLO QCD SM predictions. In contrast, for s -channel production, $\vec{P}$ is expected to be zero due to the symmetric nature of the production mechanism. The polarisation in the associated tW production exhibits a complex structure, influenced by the association with a W boson.

The degree of polarisation is quantified by:

$$\vec{P} = \frac{N_+ - N_-}{N_+ + N_-}, \quad (2.15)$$

where N_+ and N_- represent the number of top quarks with spins aligned or anti-aligned with a chosen quantisation axis, respectively.

The polarisation manifests in the angular distribution of the top-quark decay products, particularly the charged lepton (ℓ) or down-type quark from the W boson decay. The differential decay rate dependent on the polarisation is expressed as:

$$\frac{1}{\Gamma} \frac{d\Gamma}{d\cos\theta} \propto (1 + \alpha P \cos\theta), \quad (2.16)$$

where θ is the angle between the top-quark spin and the momentum of the decay product in the top-quark rest frame, and α is the spin analysing power.

The ATLAS Collaboration measured the top-quark polarisation in single top-quark t -channel production at $\sqrt{s} = 13$ TeV using the full Run 2 dataset. The three components of the polarisation vector for the selected top-quark event sample are found to be $P'_x = 0.01 \pm 0.18$, $P'_y = -0.029 \pm 0.027$, and $P'_z = 0.91 \pm 0.10$. For the top-antiquark event sample, the corresponding components are $P'_x = -0.02 \pm 0.20$, $P'_y = -0.007 \pm 0.051$, and $P'_z = -0.79 \pm 0.16$.

The analysis further constrained the complex Wilson coefficient of the dimension-six operator $\mathcal{O}_{tW}$ in the framework of effective field theory. The derived 95% CL bounds are: $C_{tW} \in [-0.9, 1.4]$, $C_{itW} \in [-0.8, 0.2]$ (see Figure 2.8) [84].

The SM predicts the spin analysing power α of top-quark decay products with high precision, as summarised in Table 2.2. At LO, the charged lepton is a perfect spin analyser with $\alpha_\ell = 1$ [85], while the down-type quark exhibits slightly lower sensitivity. Other decay

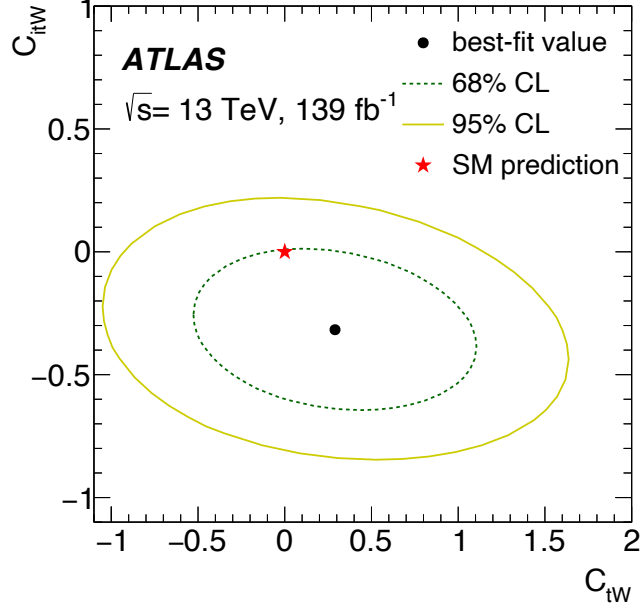


Figure 2.8: The observed best-fit value (indicated by the dot) for the Wilson coefficients C_{tW} and C_{itW} is shown, along with the uncertainty contours at 68% CL (dashed) and 95% CL (solid). These confidence regions are obtained from a simultaneous fit to both parameters. The red star marks the SM prediction [84].

products such as the neutrino, b -quark, and W boson show moderate anti-correlation with the top-quark spin.

Helicity fractions are defined as partial decay rate of each state over the total rate. In the decay of a top quark to a W boson and a bottom quark, the W boson can be emitted with different helicities: left-handed, right-handed, or longitudinal as shown in the Figure 2.10. The SM predicts the helicity fractions of the W boson due to the vector-axial (V-A) structure of the weak interaction, with the W boson primarily being longitudinally polarised or left-handed when the b quark is treated as massless.

The normalised differential decay rate as a function of the cosine of the angle θ^* between the momentum of the charged lepton in the W boson rest frame and the W boson momentum in the top-quark rest frame is given by:

$$\frac{1}{\Gamma} \frac{d\Gamma}{d\cos\theta^*} = \frac{3}{8}(1 - \cos\theta^*)^2 F_L + \frac{3}{4}\sin^2\theta^* F_0 + \frac{3}{8}(1 + \cos\theta^*)^2 F_R, \quad (2.17)$$

where F_L , F_0 , and F_R are the helicity fractions for left-handed, longitudinal, and right-handed W bosons, respectively.

The measurements of the W -boson helicity fractions in top-quark decays serve as powerful probes of the SM and potential new physics effects. For a top-quark mass of

Decay product i	α^{LO}	α^{NLO}
Charged lepton ℓ^+	+1.00	+0.998
Down-type quark (from W)	+1.00	+0.930
Neutrino ν	-0.330	-0.314
b -quark	-0.390	-0.393
W boson	-0.390	-0.393

Table 2.2: Analysing power α of top-quark decay products at LO and NLO in QCD. The values indicate the degree to which each decay product is correlated with the top-quark spin. The charged lepton provides the maximal spin sensitivity, while the b -quark and W boson exhibit moderate anti-correlation [86].

$m_t = 172.8 \pm 1.3$ GeV, NNLO QCD predictions yield $F_0 = 0.687 \pm 0.005$, $F_L = 0.311 \pm 0.005$, and $F_R = 0.0017 \pm 0.0001$ [87]. These calculations assume a b -quark mass of $m_b = 4.8$ GeV and a W -boson mass of $m_W = 80.401 \pm 0.043$ GeV, derived via NNLO computations.

The most precise Run 2 measurements by the ATLAS collaboration using $t\bar{t}$ events in the dilepton channel at $\sqrt{s} = 13$ TeV with 139 fb^{-1} of data yield [88]:

$$\begin{aligned}
F_0 &= 0.684 \pm 0.005 \text{ (stat.)} \pm 0.014 \text{ (syst.)}, \\
F_L &= 0.318 \pm 0.003 \text{ (stat.)} \pm 0.008 \text{ (syst.)}, \\
F_R &= -0.002 \pm 0.002 \text{ (stat.)} \pm 0.014 \text{ (syst.)},
\end{aligned}$$

in excellent agreement with SM expectations.

These measurements constrain the structure of the tWb vertex, which can be modified by physics beyond the SM. Anomalous couplings such as V_R , g_L , and g_R parametrise deviations from the purely left-handed SM structure, and are constrained by the measured helicity fractions. A global fit from ATLAS and CMS shown in the Figures 2.9 at $\sqrt{s} = 8$ TeV places 95% CLs of $V_R \in [-0.11, 0.16]$, $g_L \in [-0.08, 0.05]$, and $g_R \in [-0.04, 0.02]$ [89].

Alternatively, in the framework of the SMEFT, the helicity fractions are sensitive to several dimension-six operators—such as O_{tW} , $O_{\phi q}^{(3)}$, and O_{tB} —which induce deviations from the SM via suppressed contributions of order Λ^{-2} . Precision measurements of the helicity fractions thus play a crucial role in constraining the Wilson coefficients associated with these operators.

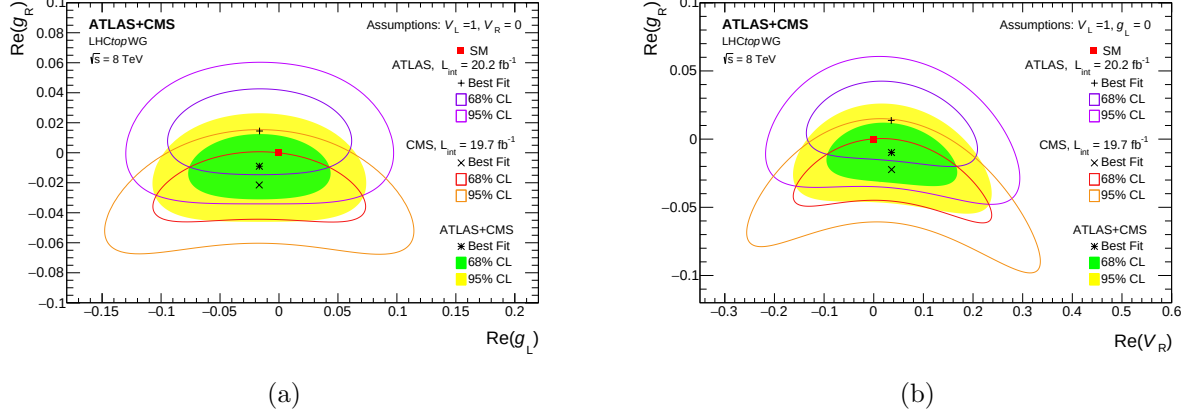


Figure 2.9: Allowed regions for anomalous tWb couplings in the planes of (a) left- and right-handed tensor couplings, and (b) right-handed vector and tensor couplings. The contours represent the 68% and 95% confidence level limits derived from measurements of the W boson helicity fractions by ATLAS, CMS, and their combination. The CMS constraints are based on the pre-combined input measurements. All anomalous couplings are assumed to be real [89].

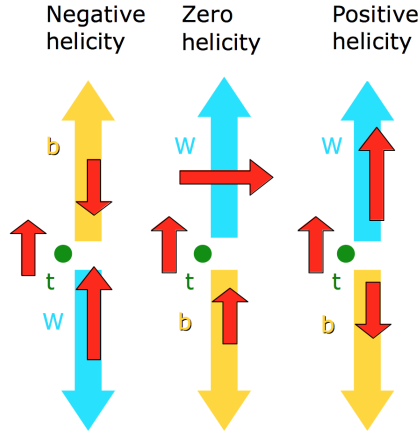


Figure 2.10: Visualisation of the helicity states of the W boson in top-quark decays in the top-quark rest frame. The red arrows indicate the spin direction (polarisation) of the top-quark, while the yellow and blue arrows represent the momentum directions of the b -quark and the W boson, respectively. The three panels correspond to the allowed helicity states of the W boson: left-handed (negative helicity), longitudinal (zero helicity), and right-handed (positive helicity). In the SM, the left-handed and longitudinal states dominate, while the right-handed configuration is highly suppressed.

2.4 Towards complete angular resolution of the tWb vertex

This chapter presents a sequence of progressively detailed angular analyses of the top-quark decay process, culminating in the fully differential analysis based on four angular variables, introduced in the Section 2.4.1, which forms the foundation of the main work presented in this thesis. While this formalism offers the most complete description of the kinematic and spin-dependent structure of the decay, its development was preceded by a series of studies employing single-variable, double-differential, and triple-differential approaches described in Sections 2.4.2, 2.4.3 and 2.4.4, respectively. These earlier analyses played an essential role in building the theoretical understanding and methodological framework required for the full multidimensional treatment. They are presented here to contextualise the development of the techniques and to highlight how the precision analysis strategy adopted in this work has emerged through successive refinements of angular observables in top-quark decay.

2.4.1 Fully differential top-quark decay formalism

The angular structure of the decay process $t \rightarrow Wb \rightarrow \ell \nu b$ can be fully characterised using four angles [90]. The first pair, (θ, ϕ) , denotes the spherical coordinates of the W boson momentum in the rest frame of the top quark, defined with respect to a chosen coordinate system (x, y, z) . A specific frame is adopted by selecting the proton beam direction that shares the same sign in p_z as the spectator quark. Following the prescription in Ref. [91], a right-handed coordinate system is constructed such that the $\hat{z}$ -axis is aligned with the spectator quark momentum $\vec{p}_s$ in the top-quark rest frame, the $\hat{y}$ -axis is defined as the cross product $\hat{z} \times \hat{p}_q$, where $\vec{p}_q$ is the direction of the incoming proton in the same frame, and the $\hat{x}$ -axis completes the orthogonal right-handed basis.

The second pair of angles, (θ^*, ϕ^*) , describes the direction of the charged lepton in the W -boson rest frame. These angles are defined in a new system (x', y', z') , obtained via a standard boost from the top-quark rest frame such that $\hat{z}'$ is aligned with the W -boson direction, and $\hat{y}'$ lies within the plane of the original xy -axes. This setup is illustrated in Figure 2.11.

Using the Jacob–Wick helicity formalism [92], the decay amplitude for the full process can be factorised in the narrow width approximation as:

$$A_{MF_2F_3F_4} = \sum_{F_1} a_{F_1F_2} b_{F_3F_4} D_{MF}^{1/2*}(\phi, \theta, 0) D_{F_1F}^{1*}(\phi^*, \theta^*, 0), \quad (2.18)$$

where F_1 , F_2 , F_3 , and F_4 are the helicities of the W boson, b -quark, charged lepton, and neutrino, respectively; subindex M denotes the spin component of the top quark; and $F = F_1 - F_2$; and $F = F_3 - F_4$.

The angular structure is governed by the Wigner D -matrices [93]:

$$D_{m'm}^j(\alpha, \beta, \gamma) = \langle jm' | e^{-i\alpha J_z} e^{-i\beta J_y} e^{-i\gamma J_z} | jm \rangle, \quad (2.19)$$

and the coefficients $a_{F_1 F_2}$ and $b_{F_3 F_4}$ encode the vertex dynamics.

For top-quark decay, angular momentum conservation limits the number of non-zero coefficients to four: $a_{1\frac{1}{2}}$, $a_{0\frac{1}{2}}$, $a_{0-\frac{1}{2}}$, and $a_{-1-\frac{1}{2}}$. In the SM and under the assumption of massless leptons, only one reduced coefficient $b_{F_3 F_4}$ contributes, and its value can be factored out from the angular dependence.

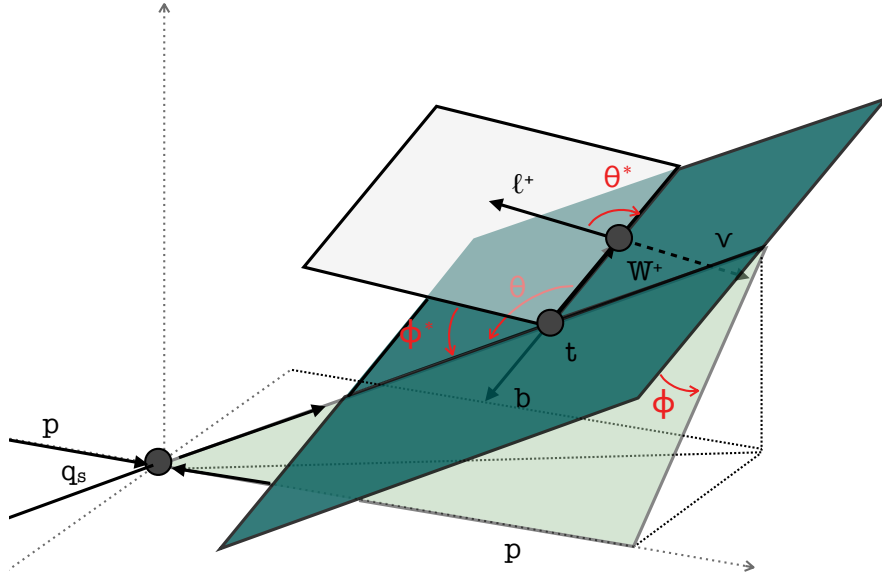


Figure 2.11: The angles θ and ϕ describe the momentum of the W boson in the top-quark rest frame, while θ^* and ϕ^* describe the momentum of the charged lepton in the W -boson rest frame.

The normalised quadruple differential decay rate is expanded using an orthonormal basis of functions $M_{m'm}^{j_1 j_2}$, defined as:

$$M_{m'm}^{j_1 j_2}(\phi, \theta, \phi^*, \theta^*) = \frac{1}{4\pi} \sqrt{(2j_1 + 1)(2j_2 + 1)} D_{m'm}^{j_1}(\phi, \theta, 0) D_{m'0}^{j_2}(\phi^*, \theta^*, 0). \quad (2.20)$$

The angular distribution can thus be written as:

$$\frac{1}{\Gamma} \frac{d\Gamma}{d\Omega d\Omega^*} = \sum_{j_1, j_2, m', m} c_{m'm}^{j_1 j_2} M_{m'm}^{j_1 j_2}(\phi, \theta, \phi^*, \theta^*), \quad (2.21)$$

where each M -function is constructed from a product of Wigner D -functions and forms an orthonormal basis for angular distributions. The coefficients $c_{m'm}^{j_1 j_2}$ are directly observable, and their structure reflects the interplay of polarisation, helicity fractions, and interference terms. Only a finite number of coefficients contribute due to the spins of the involved particles: $j_1 \in 0, 1$ and $j_2 \in 0, 1, 2$. Thus, the non-zero coefficients in the expansion are

$$\begin{aligned} c_{00}^{00} &= \frac{1}{4\pi}, \\ c_{00}^{10} &= \frac{1}{4\sqrt{3}\pi} P_z \left[|a_{1\frac{1}{2}}|^2 - |a_{0\frac{1}{2}}|^2 + |a_{0-\frac{1}{2}}|^2 - |a_{-1-\frac{1}{2}}|^2 \right] / \mathcal{N}, \\ c_{10}^{10} &= -(c_{-10}^{10})^* = -\frac{1}{4\sqrt{6}\pi} (P_x + iP_y) \left[|a_{1\frac{1}{2}}|^2 - |a_{0\frac{1}{2}}|^2 + |a_{0-\frac{1}{2}}|^2 - |a_{-1-\frac{1}{2}}|^2 \right] / \mathcal{N}, \\ c_{00}^{01} &= \lambda \frac{\sqrt{3}}{8\pi} \left[|a_{1\frac{1}{2}}|^2 - |a_{-1-\frac{1}{2}}|^2 \right] / \mathcal{N}, \\ c_{00}^{11} &= \lambda \frac{1}{8\pi} P_z \left[|a_{1\frac{1}{2}}|^2 + |a_{-1-\frac{1}{2}}|^2 \right] / \mathcal{N}, \\ c_{10}^{11} &= -(c_{-10}^{11})^* = -\lambda \frac{1}{8\sqrt{2}\pi} (P_x + iP_y) \left[|a_{1\frac{1}{2}}|^2 + |a_{-1-\frac{1}{2}}|^2 \right] / \mathcal{N}, \\ c_{01}^{11} &= (c_{0-1}^{11})^* = \lambda \frac{1}{4\sqrt{2}\pi} P_z \left[a_{0\frac{1}{2}} a_{1\frac{1}{2}}^* + a_{-1-\frac{1}{2}} a_{0-\frac{1}{2}}^* \right] / \mathcal{N}, \\ c_{11}^{11} &= -(c_{-1-1}^{11})^* = -\lambda \frac{1}{8\pi} (P_x + iP_y) \left[a_{0\frac{1}{2}} a_{1\frac{1}{2}}^* + a_{-1-\frac{1}{2}} a_{0-\frac{1}{2}}^* \right] / \mathcal{N}, \\ c_{1-1}^{11} &= -(c_{-11}^{11})^* = -\lambda \frac{1}{8\pi} (P_x + iP_y) \left[a_{1\frac{1}{2}} a_{0\frac{1}{2}}^* + a_{0-\frac{1}{2}} a_{-1-\frac{1}{2}}^* \right] / \mathcal{N}, \\ c_{00}^{02} &= \frac{1}{8\sqrt{5}\pi} \left[|a_{1\frac{1}{2}}|^2 - 2|a_{0\frac{1}{2}}|^2 - 2|a_{0-\frac{1}{2}}|^2 + |a_{-1-\frac{1}{2}}|^2 \right] / \mathcal{N}, \\ c_{00}^{12} &= \frac{1}{8\sqrt{15}\pi} P_z \left[|a_{1\frac{1}{2}}|^2 + 2|a_{0\frac{1}{2}}|^2 - 2|a_{0-\frac{1}{2}}|^2 - |a_{-1-\frac{1}{2}}|^2 \right] / \mathcal{N}, \\ c_{10}^{12} &= -(c_{-10}^{12})^* = -\frac{1}{8\sqrt{30}\pi} (P_x + iP_y) \left[|a_{1\frac{1}{2}}|^2 + 2|a_{0\frac{1}{2}}|^2 - 2|a_{0-\frac{1}{2}}|^2 - |a_{-1-\frac{1}{2}}|^2 \right] / \mathcal{N}, \\ c_{01}^{12} &= (c_{0-1}^{12})^* = \frac{1}{4\sqrt{10}\pi} P_z \left[a_{0\frac{1}{2}} a_{1\frac{1}{2}}^* - a_{-1-\frac{1}{2}} a_{0-\frac{1}{2}}^* \right] / \mathcal{N}, \\ c_{11}^{12} &= -(c_{-1-1}^{12})^* = -\frac{1}{8\sqrt{5}\pi} (P_x + iP_y) \left[a_{0\frac{1}{2}} a_{1\frac{1}{2}}^* - a_{-1-\frac{1}{2}} a_{0-\frac{1}{2}}^* \right] / \mathcal{N}, \\ c_{1-1}^{12} &= -(c_{-11}^{12})^* = -\frac{1}{8\sqrt{5}\pi} (P_x + iP_y) \left[a_{1\frac{1}{2}} a_{0\frac{1}{2}}^* - a_{0-\frac{1}{2}} a_{-1-\frac{1}{2}}^* \right] / \mathcal{N}. \end{aligned} \quad (2.22)$$

where $\mathcal{N}$ is corresponding to the number of events. Each coefficient has a clear physical interpretation. For instance, c_{00}^{10} probes the longitudinal top-quark polarisation, P_z , while c_{00}^{01} is sensitive to the W -boson helicity asymmetry ($F_+ - F_-$). The c_{01}^{11} and c_{01}^{12} coefficients are sensitive to the interference between longitudinal and transverse helicity amplitudes and allow for measurement of relative phases, especially those arising from complex EFT couplings.

2.4.2 Charge asymmetry and CP violation in single top-quark events

Angular observables in single-top-quark t -channel events provide a sensitive probe for potential CP-violating effects in the tWb vertex. One such observable is the forward-backward asymmetry, A_{FB}^N , defined using the angle θ_N between the lepton from W -boson decay and the top-quark spin direction, measured in the top-quark rest frame:

$$A_{\text{FB}}^N = \frac{N(\cos \theta_N > z) - N(\cos \theta_N < z)}{N(\cos \theta_N > z) + N(\cos \theta_N < z)}. \quad (2.23)$$

where normal direction N is defined as $\vec{N} = \vec{s}_t \times \vec{q}$, $\vec{s}_t$ being spin direction (direction of the spectator quark in the top quark rest frame) and $\vec{q}$ the W boson momentum in the top quark rest frame (helicity basis). And z is an arbitrary but fixed point in the angular distribution. In this analysis three out of four angles in the Equation 2.21 is integrated, keeping dependence only on parameter θ . This asymmetry is particularly sensitive to the imaginary part of the anomalous coupling g_R , with the LO V_R prediction yielding $A_{\text{FB}}^N = 0.64 P \text{Im}(g_R)$ assuming $V_L = 1$ and $V_R = g_L = 0$. At one-loop level, EW corrections in the SM induce a non-zero contribution to an effective right-handed coupling, yielding $g_R^{\text{eff}} = (-7.17 - 1.23i) \times 10^{-3}$. This is not a bare Lagrangian parameter, but rather an effective coupling arising from EW loop effects. The imaginary component, responsible for T-odd observables, contributes approximately 17% of the total magnitude. Note that this contribution is reversed under charge conjugation ($g_R^{\text{eff}} \rightarrow g_R^*$ for anti-top quarks), and does not correspond to CP violation [94].

The ATLAS experiment performed this measurement using 4.7 fb^{-1} of $\sqrt{s} = 7 \text{ TeV}$ pp collision data from 2011 [94]. Events were selected in the single-lepton channel, and angular distributions were corrected using a matrix-based unfolding⁵ procedure.

The measured asymmetry was found to be $A_{\text{FB}}^N = 0.31 \pm 0.065 \text{ (stat.)} \pm 0.03 \text{ (syst.)}$,

⁵Unfolding refers to the correction of reconstructed data for detector resolution, efficiency, and acceptance effects, in order to infer the underlying particle-level distribution.

consistent with the SM prediction within uncertainties. The extracted constraint on the imaginary part of g_R , shown in Figure 2.12, limits sizeable CP violating contributions beyond the V_R coupling.

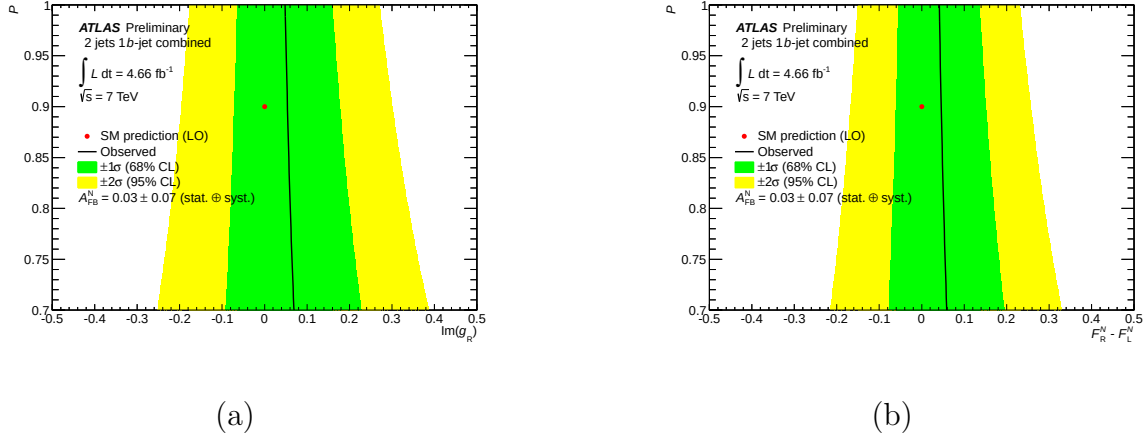


Figure 2.12: (a) Constraints on polarisation as a function $\text{Im}(g_R)$ from the A_{FB}^N measurement. (b) Allowed parameter space in polarisation as function $F_R^N - F_L^N$, both with 68% and 95% CLs [94].

2.4.3 Double-Differential Analysis

The two-angle study of performed by the ATLAS collaboration [95] exploits the joint distribution of the polar and azimuthal decay angles, θ^* and ϕ^* , of the W boson in the single-top-quark t -channel events. Expressed in spherical harmonics $Y_\ell^m(\theta^*, \phi^*)$ ⁶, the normalised rate in this analysis forms the basis for later developments as shown in Equation 2.18, where the angular structure is generalised using Wigner D -functions to incorporate the full set of decay angles. The Wigner D -functions extend the spherical harmonics formalism to higher-dimensional spin-angular correlations, enabling a more complete treatment of interference effects and spin dynamics:

$$\rho(\theta^*, \phi^*; \hat{z}, P) = \sum_{\ell \leq 2} \sum_{m=-\ell}^{\ell} a_{\ell,m}(\hat{z}, P) Y_\ell^m(\theta^*, \phi^*), \quad (2.24)$$

with the parameter vector $\hat{z} = (F_T, \frac{F_R}{F_T}, \frac{F_0^{(R)}}{F_0}, \delta_+, \delta_-)$, where $F_T = F_L + F_R$ is the total transverse helicity fraction of W -boson, F_R/F_T is the fraction of transverse W -bosons that are right-handed, $F_0^{(R)}/F_0$ is the fraction of longitudinal W -bosons produced in association with a right-handed b -quark. and δ_+ , δ_- are the relative phases between the

⁶Here P denotes the magnitude of the top-quark polarisation. The phases $\delta_\pm$ are defined exactly as in Ref. [95].

longitudinal and transverse helicity amplitudes. The coefficients $a_{\ell,m}$ used in analysis correspond to a reduced subset of the full set of angular coefficients $c_{m'm}^{j_1 j_2}$ in Equations 2.22 and 2.21. In particular, they arise from integrating over the angles (θ, ϕ) associated with the W boson direction in the top-quark rest frame, leaving a residual dependence only on (θ^*, ϕ^*) . Mathematically, this integration projects out the spherical harmonics basis, and the resulting coefficients correspond to j_2 and m in the full expansion, with $j_1 = 0$.

Explicitly, the non-zero amplitudes are:

$$\begin{aligned} a_{0,0} &= \frac{1}{\sqrt{4\pi}}, & a_{2,0} &= \frac{1}{\sqrt{20\pi}} (3F_T - 1), \\ a_{1,0} &= \frac{\sqrt{3}}{\sqrt{4\pi}} F_T \left(\frac{F_R}{F_T} + \frac{1}{2} \right), & a_{1,1} &= -a_{1,-1}^* (\propto \text{Re}, \text{Im}[g_R/V_L]), \\ a_{2,1} &= -a_{2,-1}^* (\propto \text{Re}, \text{Im}[g_R/V_L]), & a_{2,2} &= a_{2,-2} = 0. \end{aligned}$$

Using 4.6 fb^{-1} of pp collisions at $\sqrt{s} = 7 \text{ TeV}$ [95], the ATLAS collaboration performed a likelihood fit that folds the parton-level⁷ distribution with detector efficiency, resolution and background densities expressed in the same harmonic basis. Setting $V_R = g_L = 0$, the fit returns

$$\begin{aligned} F_T &= 0.37 \pm 0.05(\text{stat}) \pm 0.05(\text{syst}), \\ \delta_- &= -0.014\pi \pm 0.023\pi(\text{stat}) \pm 0.028\pi(\text{syst}), \\ \text{Re}[g_R/V_L] &= -0.13 \pm 0.07(\text{stat}) \pm 0.10(\text{syst}), \\ \text{Im}[g_R/V_L] &= 0.03 \pm 0.06(\text{stat}) \pm 0.07(\text{syst}). \end{aligned} \tag{2.25}$$

At the 95% CL, the anomalous couplings are shown in Figure 2.13, with ratios constrained to:

$$-0.36 < \text{Re} \left[\frac{g_R}{V_L} \right] < 0.10, \quad -0.17 < \text{Im} \left[\frac{g_R}{V_L} \right] < 0.23.$$

The likelihood projections for F_T and δ_- are shown in Figure 2.14, as these parameters play a crucial role in characterising the interference between different helicity states.

2.4.4 Triple-differential analysis

This analysis investigates the triple differential decay rate of single top-quark events in the t -channel, focusing on constraints on anomalous couplings and the structure of the tWb

⁷Parton-level refers to the state of quarks and gluons as generated by the matrix element calculation, before any hadronisation or detector effects are applied.

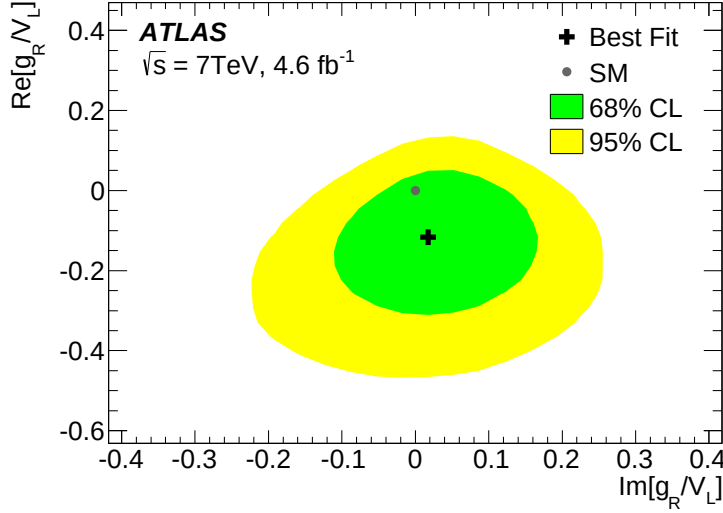


Figure 2.13: Two-dimensional likelihood projection of $\text{Re}[g_R/V_L]$ versus $\text{Im}[g_R/V_L]$, including systematic uncertainties. The black cross denotes the observed best-fit value, and the grey point marks the SM prediction. The 68% and 95% CL regions are shown in green and yellow, respectively [95].

vertex. This study also uses Run 1 data recorded by ATLAS detector at a centre-of-mass energy of $\sqrt{s} = 8$ TeV, corresponding to an integrated luminosity of 20.2 fb^{-1} [96].

The differential decay rate is analysed as a function of three angular variables $(\theta, \theta^*, \phi^*)$ that capture both the polarisation of the top quark and the helicity states of its decay products. The decay distribution is expanded using a basis of orthonormal M -functions, similarly to the full-differential analysis formalism described in Equation 2.20, enabling the extraction of angular coefficients via the orthogonal series density estimation (OSDE) [97]. The angular coefficients used in the triple-differential analysis correspond to the subset of $c_{m'm}^{j_1 j_2}$ in Equations 2.22 and 2.21 with $m' = 0$. This results from integrating over the azimuthal angle ϕ of the W boson, which eliminates terms with $m' \neq 0$ due to the orthogonality of the Wigner D -functions. The resulting angular distribution depends on $(\theta, \theta^*, \phi^*)$ and retains sensitivity to both polarisation and interference effects while simplifying the parameter space.

To correct for detector effects, the analysis applies a Fourier-dual space deconvolution method [98], using a migration matrix derived from MC simulations. The true angular coefficients are extracted through a regularised χ^2 minimisation, allowing for the reconstruction of physical parameters such as helicity fractions and anomalous couplings.

The key results of this analysis are:

- Right-handed b -quark helicity fraction: $\frac{F_0^R}{F_0} + < 0.041$ (68% CL), consistent with the

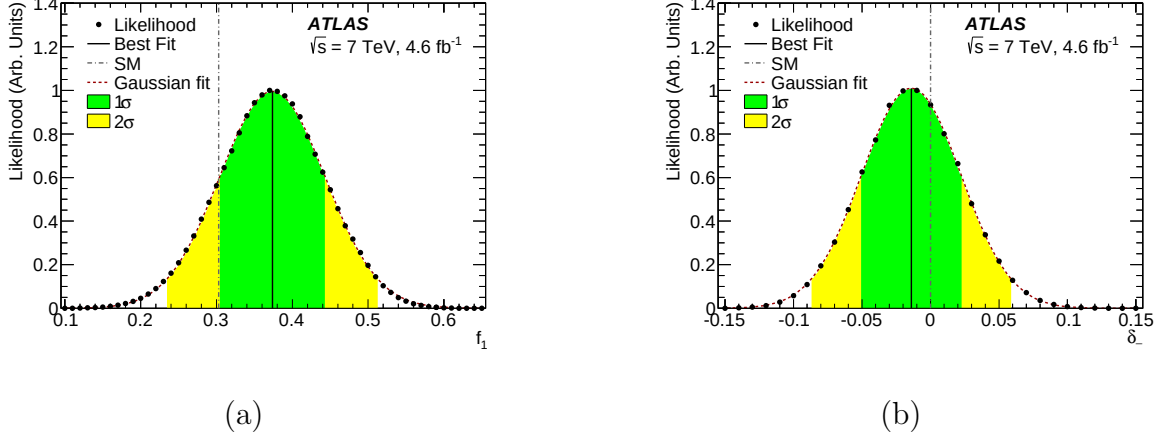


Figure 2.14: Projections of the likelihood function constructed from the signal region probability density and data events into (a) $F_T \equiv f_1$, (b) δ_- , with systematic uncertainties included. The black points represent the highest evaluated likelihood in each bin of the projected variable. Gaussian fits to the one-dimensional projections are shown as the red curve. Regions in green and yellow correspond to the 68% and 95% confidence level regions, respectively. A black line indicates the observed value, while the grey line or point indicates the SM expectation [95].

SM prediction $\frac{F_0^R}{F_0} = 6 \times 10^{-5}$.

- Right-handed W boson helicity fraction: $\frac{F_T^R}{F_T} < 0.053$ (68% CL), consistent with the SM expectation $\frac{F_T^R}{F_T} = 0.001$.
- W bosons produced with right-handed b -quarks, yielding 95% CL limits on the complex coupling ratio: $\text{Re}[g_R/V_L] \in [-0.12, 0.17]$, $\text{Im}[g_R/V_L] \in [-0.07, 0.06]$.
- The polarisation of single top quarks in the t -channel is measured to satisfy $P > 0.72$ at 95% CL.

The individual likelihood profiles for the helicity fractions $\frac{F_0^{(R)}}{F_0}$ and $\frac{F_R}{F_T}$ are shown in Figure 2.15, illustrating the consistency of the observed values with the SM expectations within the 68% and 95% CL regions.

The observed constraints on the anomalous coupling ratios $\text{Re}[g_R/V_L]$ and $\text{Im}[g_R/V_L]$ are summarised in Figure 2.16, which shows the joint likelihood contours including both statistical and systematic uncertainties. These results provide stringent bounds on anomalous couplings in top quark interactions and demonstrate no significant deviation from SM expectations.

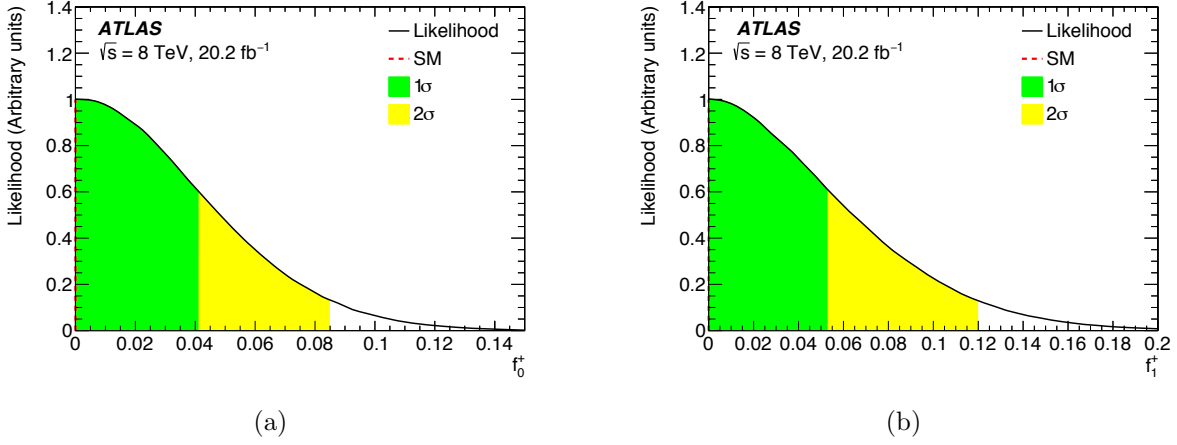


Figure 2.15: Likelihood profiles for the parameters $\frac{F_0^{(R)}}{F_0} \equiv f_0^+$ and $\frac{F_R}{F_T} \equiv f_1^+$. Panels (a) and (b) show the one-dimensional likelihoods for f_0^+ and f_1^+ , respectively, where the black line represents the profiled likelihood and the red dashed line indicates the SM expectation. The 68% and 95% CL regions are shown in green and yellow, respectively [96].

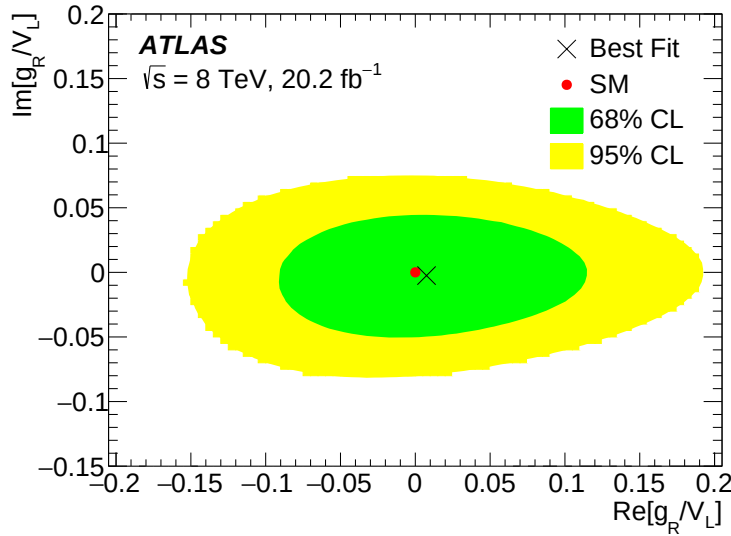


Figure 2.16: Joint likelihood of $\text{Re}[g_R/V_L]$ vs. $\text{Im}[g_R/V_L]$ with statistical and systematic uncertainties. The red marker indicates the SM point [96].

This result provides a direct handle on the chirality structure of the tWb vertex and sets meaningful constraints on the EFT operators that induce right-handed currents. In particular, the limits on g_R/V_L map onto combinations of dimension-six Wilson coefficients, including C_{tW} , which are also probed in the quadruple-differential analysis presented later in this thesis Section 5.8. The three-angle measurement thus serves as a complementary approach, targeting a reduced angular basis but still sensitive to key EFT degrees of

freedom.

This chapter presented a detailed study of top quark spin observables in the t -channel single top production using angular distributions of the charged lepton. The analysis extracted polarisation components and helicity fractions, and used these to constrain anomalous couplings in the framework of an effective field theory. All measured observables are consistent with SM predictions, and no significant deviations were observed. The resulting limits on EFT Wilson coefficients provide robust constraints on potential new physics in the top quark sector.

Chapter 3

Manifesting theories into empirical realities

The quest to translate abstract theoretical constructs into measurable phenomena lies at the heart of modern particle physics. CERN serves as a global nerve centre for high-energy physics research and technological innovation. Established in 1954 by twelve European countries, CERN has grown into one of the most prominent scientific institutions worldwide, fostering international collaboration in fundamental physics. It has been the site of numerous landmark achievements, including the invention of the World Wide Web (WWW) [99], and the discovery of the W and Z bosons [5], which confirmed the EW unification. More recently, the observation of the Higgs boson in 2012 by the ATLAS and CMS collaborations [9, 10] marked an important confirmation of the SM as already mentioned in Chapter 1. Beyond these discoveries, CERN continues to advance technologies in particle detection, data analysis, and accelerator systems, while serving as a training ground for physicists and engineers across the globe.

This thesis is conducted within the framework of the ATLAS experiment, one of the flagship projects hosted by CERN. The experiment operates at the LHC — the most powerful particle accelerator constructed to date [7, 8]. The LHC accelerates and collides protons as well as heavy ions at unprecedented energies to explore the building blocks of matter and the fundamental forces governing their interactions. This accelerator complex has become synonymous with cutting-edge research in particle physics, driving both theoretical advances and experimental verifications with unprecedented precision.

This chapter provides an overview of the experimental foundations underpinning the measurements and analyses presented in this thesis. Section 3.1 begins with a description of the accelerator system, including the concept of luminosity and the major experiments

operating at the LHC. Section 3.2 introduces the ATLAS detector, with a particular focus on its tracker detector, which plays a central role in charged particle tracking. Section 3.3 covers ATLAS detector performance, the procedures for alignment of the mentioned tracker and the challenges associated with momentum reconstruction. In Section 3.4, refined strategy for evaluating and correcting momentum biases are discussed as part of the work developed within this thesis. These components together establish the experimental framework required for the high-precision measurements discussed in later chapters.

3.1 The accelerator system

Regardless of the final energy or particle type, all beams are accelerated in stages through a series of interconnected machines that make up the long-established accelerator complex at CERN. The structure of this system is shown in Figure 3.1.

The CERN accelerator complex is a multi-stage system designed to accelerate particles through a sequence of machines, each optimised for a specific energy range and beam manipulation purpose. The chain begins with hydrogen ion extraction and culminates in proton collisions at the LHC, reaching a centre-of-mass energy of 13.6 TeV. The layout and interconnectivity of this system are also illustrated in Figure 3.1.

The acceleration sequence begins at Linac4, which replaced the older Linac2 in 2020. Linac4 accelerates negative hydrogen ions (H^-) to an energy of 160 MeV over a length of 86 m using a combination of Radio Frequency Quadrupole, drift-tube linac, and Cell Coupled Drift Tube Linac sections operating at 352.2 MHz [100]. This energy is optimised for charge-exchange injection into the next stage, the Proton Synchrotron Booster (PSB), improving injection efficiency and minimising emittance growth.

The PSB consists of four vertically stacked synchrotron rings, each with a circumference of 157 m. It receives the 160 MeV H^- beam from Linac4, where the ions are stripped of their electrons to form protons through a charge-exchange injection process. The PSB accelerates these protons to 2 GeV, delivering them to the Proton Synchrotron (PS). The upgrade from the previous 1.4 GeV top energy to 2 GeV, enabled by Linac4, has significantly improved longitudinal brightness and reduced space-charge effects, both of which are critical for high-intensity beam operation [101].

The PS, a 628 m circumference synchrotron, has been operational since 1959 but has undergone numerous upgrades. It accepts 2 GeV protons from the PSB and accelerates them to 25 GeV. The PS also performs sophisticated bunch manipulations including bunch splitting and spacing adjustment required by downstream accelerators [102]. In

March 2025, the PS successfully completed standalone recommissioning, marking its readiness for Run 3 operation [103].

The Super Proton Synchrotron (SPS), with a 6.9 km circumference, serves as the final injector before the LHC. It receives 25 GeV protons from the PS and accelerates them to 450 GeV using Radio Frequency cavities operating at 200 MHz. The SPS is also equipped with fast extraction systems to transfer beams into the LHC and beam dump lines. The machine has 744 dipole and 216 quadrupole magnets [104], and is critical for both pp and heavy-ion beam preparation.

The LHC is the final stage in the CERN accelerator chain. It has a circumference of 26.7 km and is located between 45 m and 170 m underground. Beams are injected into the LHC at an energy of 450 GeV and are accelerated to 6.8 TeV per beam. The LHC employs a complex system of superconducting magnets operating at 1.9 K, including 1232 dipole magnets used to bend the beam, and over 400 quadrupole magnets used for beam focusing. The dipoles, each 14.3 m in length, generate magnetic fields up to 8.33 T. Together with additional corrector magnets and collimators, these elements form the magnetic lattice that guides and shapes the beam through the accelerator ring. The LHC supports up to 2808 bunches per beam with a 25 ns spacing, enabling peak luminosities on the order of $2 \times 10^{34} \text{ cm}^{-2}\text{s}^{-1}$ during Run 3 operations. Physics data-taking resumed in March 2025 following a successful recommissioning campaign [105].

The LHC project was officially approved in 1994, with construction commencing shortly thereafter. It was built in the pre-existing 26.7 km tunnel formerly occupied by the Large Electron–Positron Collider (LEP), situated 45–170 m underground beneath the Franco–Swiss border near Geneva. Reusing this infrastructure minimised costs and accelerated project realisation. The LHC was designed as a twin-beam synchrotron, employing a two-in-one magnet configuration in which two counter-circulating beams share a common cryostat. This twin-aperture dipole design allows the beams to be steered using a single cold mass, optimising spatial and cryogenic efficiency [7].

Particle bunches in the LHC cross at a frequency of 40 MHz — every 25 ns — at the interaction points, setting the fundamental timing scale for data acquisition systems. Since the beginning of its operational phase in 2009, the LHC has undergone several data-taking periods, known as "runs", interleaved with extended technical shutdowns for maintenance and upgrades.

Run 1, spanning from 2010 to 2013, involved pp collisions at centre-of-mass energies of $\sqrt{s} = 7$ and 8 TeV. It marked the first full-scale physics exploitation of the machine and culminated in the discovery of the Higgs boson. Run 2 followed from 2015 to 2018, during which the LHC delivered collisions at $\sqrt{s} = 13$ TeV. The dataset collected during

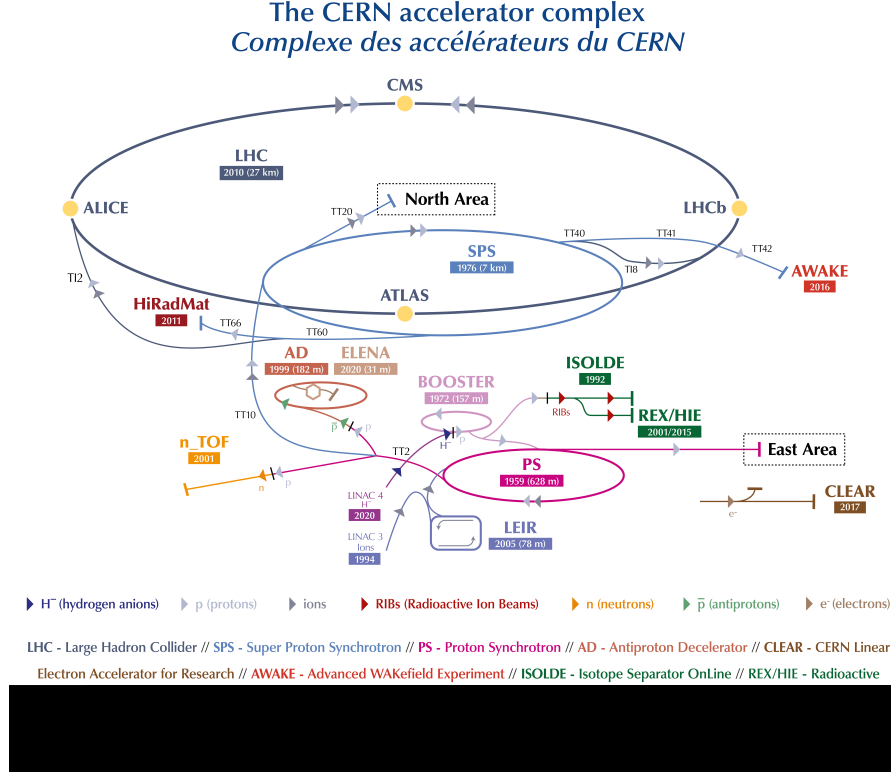


Figure 3.1: The CERN accelerator complex [106].

this period forms a central foundation for the measurements and analyses presented in this thesis. The current phase, Run 3, began in July 2022 and is expected to conclude in 2026. This phase is operating at a record collision energy of $\sqrt{s} = 13.6$ TeV, with significantly improved beam parameters and detector capabilities [107].

Following each major run, the LHC enters a period of extended downtime, known as a "Long Shutdown". These intervals serve several purposes: they enable upgrades to accelerator and detector systems, address wear and tear, and provide the physics community with valuable time to process and analyse the massive datasets accumulated during the prior operational cycle.

The Long Shutdown scheduled after Run 3, expected to begin in 2026, is of particular importance, as it will initiate the transition to the High-Luminosity LHC (HL-LHC) [108]. The HL-LHC represents a major upgrade to the machine, aiming to increase the instantaneous luminosity by a factor of 5–7 and the integrated luminosity by an order of magnitude compared to the original design. In collider physics, luminosity — defined as the number of interactions per unit area per unit time — is a critical parameter, as it directly impacts the statistical sensitivity of measurements and the discovery potential for rare processes. The HL-LHC is scheduled to begin operations around 2030 and is

expected to deliver more than 3000 fb^{-1} of data by the end of its lifetime [109]. Table 3.1 provides a summary of key beam and machine parameters across different LHC runs and the planned HL-LHC configuration.

Parameter	Design	Run 1	Run 2	Run 3	HL-LHC
Beam energy [TeV]	7	3.5 – 4	6.5	6.8	7
Centre-of-mass energy ($\sqrt{s}$) [TeV]	14	7 – 8	13	13.6	14
Bunch spacing [ns]	25	50	25	25	25
Bunch intensity [10^{11} ppb]	1.15	1.6	1.2	up to 1.8	2.2
Number of bunches (n_b)	2808	1400	2500	2800	2800
Transverse emittance (ϵ) [μm]	3.5	2.2	2.2	2.5	2.5
Amplitude function at interaction point (β^*) [cm]	55	80	30→25	30→25	down to 15
Crossing angle [μrad]	285	–	300→260	300→260	TBD
Peak luminosity [$10^{34} \text{ cm}^{-2}\text{s}^{-1}$]	2.1	0.8	2.0	2.0	5.0
Peak pile-up	25	45	60	55	150
Nominal magnetic field (B) [T]	8.73	4.16 – 7.76	7.73	8.73	8.73 / 12
Injection energy [GeV]			450		
Circumference length [km]			26.7		
Radius [km]			4.24		
Number of dipole magnets			1232		
Length of dipole magnets [m]			14.3		
Number of quadrupole magnets			395		
Total mass [tons]			27.5		

Table 3.1: Overview of key accelerator parameters for the LHC, including the original design values, the configurations used during Run 1, Run 2, and Run 3, and the projected parameters for Run 4 as part of the HL-LHC upgrade [7, 110, 111].

3.1.1 Luminosity

In the complex realm of particle interactions, the processes of interest are often extremely rare, necessitating the statistical analysis of a large number of collisions to extract meaningful signals. Every 25 ns, proton bunches containing approximately 10^{11} protons are brought into collision (see Table 3.1). The event rate $\frac{dN}{dt}$ for a given process is governed by the process cross-section σ and the instantaneous luminosity L of the collider:

$$\frac{dN}{dt} = \sigma \cdot L. \quad (3.1)$$

The instantaneous luminosity L can be expressed in terms of the beam parameters as:

$$L = \frac{N_{b1} N_{b2} n_b f_{\text{rev}}}{4\pi \sigma_x \sigma_y}, \quad (3.2)$$

where N_{b1} and N_{b2} are the number of protons per bunch in each beam, n_b is the number of bunches per beam (identical for synchronisation), f_{rev} is the revolution frequency of the accelerator, and σ_x and σ_y represent the transverse root mean square (RMS) sizes of the colliding bunches in the horizontal and vertical planes, respectively.

To determine the production cross-section of a process from event counts, one must measure the integrated luminosity:

$$L_{\text{int}} = \int L(t) dt. \quad (3.3)$$

The integrated luminosity is commonly reported in units of inverse femtobarns, where $1 \text{ fb}^{-1} = 10^{39} \text{ cm}^{-2}$, although other units may be used depending on the precision and dataset scale. A related and practically important quantity is the mean number of interactions per bunch crossing (name as pile-up), denoted by μ . It represents the average number of inelastic pp collisions occurring each time two proton bunches cross at an interaction point. This value is calculated per bunch from the instantaneous per-bunch luminosity using

$$\mu = \frac{L_{\text{bunch}} \cdot \sigma_{\text{inel}}}{f_{\text{rev}}},$$

where L_{bunch} is the instantaneous luminosity per bunch and σ_{inel} is the inelastic pp cross-section (taken to be approximately 80 mb at 13 TeV).

Over the full course of Run 2, the LHC delivered approximately 160 fb^{-1} of integrated luminosity to ATLAS and CMS. Of this, ATLAS successfully recorded about 139 fb^{-1} for offline analysis [112]. The analysis presented in this thesis is based on the Run-2 dataset collected at a centre-of-mass energy of 13 TeV. Figure 3.2 (a) displays the cumulative distributions of the integrated luminosity registered by the ATLAS detector during LHC stable beams as a function of time for the Run-1 (2011-2012) and Run-2 (2015-2018) data-taking periods.

3.1.2 Pile-Up

The luminosity in pp collisions can be increased in three principal ways, as described by Equation 3.2: by increasing the number of protons per bunch, reducing the transverse width of the colliding beams, or increasing the number of bunches per beam. Each of

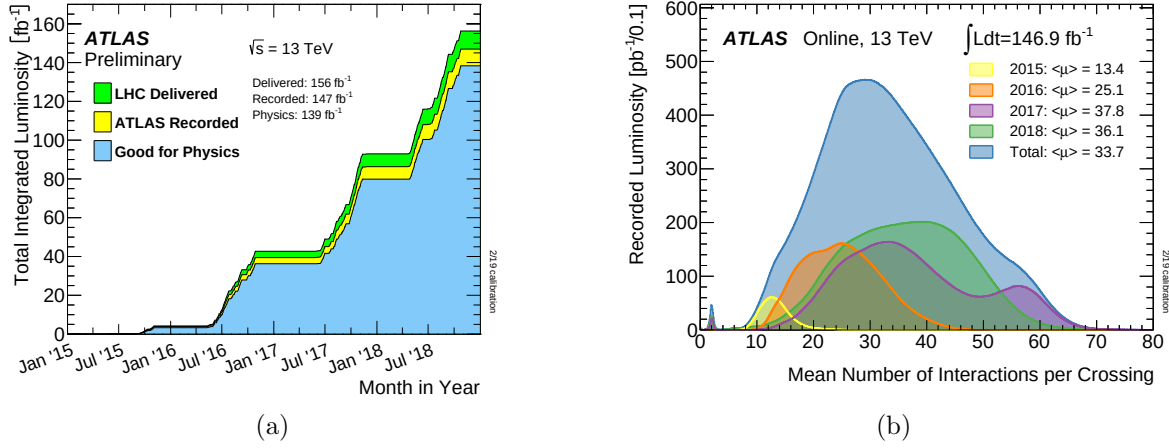


Figure 3.2: (a) Integrated luminosity delivered to (green) and recorded (yellow) by the ATLAS detector during Run2. The integrated luminosity of the collisions usable for physics analyses is shown in blue. (b) The pile-up (μ) distribution measured by the ATLAS detector during Run2 [112].

these methods results in a higher density of collisions per unit time and thus contributes to an increase in the number of particle interactions within a given bunch crossing — a phenomenon known as pile-up.

Pile-up can be categorised as either in-time, referring to multiple interactions occurring simultaneously within the same bunch crossing, or out-of-time, where the detector integrates activity from neighbouring bunch crossings due to its finite temporal resolution. Increasing the number of protons per bunch or reducing the beam width primarily enhances in-time pile-up, while decreasing the bunch spacing (i.e. increasing the number of bunches) contributes more significantly to out-of-time pile-up effects [113].

Pile-up is typically composed of secondary, low transverse momentum (p_T) interactions that are not directly relevant to the process under study. These additional interactions complicate the reconstruction of physics objects — particularly jets and missing transverse energy — and can degrade the resolution and performance of many analyses. As a result, controlling and mitigating pile-up is a key challenge in high-luminosity hadron collider environments.

The ATLAS detector employs a range of techniques to mitigate pile-up effects, including jet area-based subtraction [114], constituent-level pile-up corrections [115], and calorimeter signal timing to identify out-of-time contributions [116]. These strategies have proven effective in maintaining detector performance under increasing luminosity conditions.

Figure 3.2 (b) shows the luminosity-weighted distribution of the mean number of

interactions per bunch crossing, denoted by μ , for each year of data-taking during Run 2. The steady rise in μ — reaching average values above 36 by the end of 2018 [117] — reflects the increasing instantaneous luminosity delivered by the LHC and presents a progressively more challenging pile-up environment. These high pile-up conditions are directly relevant to the dataset analysed in this thesis, which is based on the full Run 2 pp collision data at $\sqrt{s} = 13$ TeV.

3.1.3 The Experiments at the Large Hadron Collider

The LHC is the most powerful particle accelerator constructed to date, located beneath the Franco–Swiss border in the tunnel previously occupied by the LEP collider [7] as mentioned before. It is designed to accelerate and collide proton and heavy-ion beams, including Pb–Pb collisions at energies suitable for studying strongly interacting matter.

The LHC at CERN hosts several experiments. The four primary detectors—ATLAS, CMS, ALICE, and LHCb—serve as the main instruments for data collection and analysis. Additionally, several smaller experiments contribute to specific research areas.

The **A Toroidal LHC ApparatuS (ATLAS)** is the largest general-purpose detector at the LHC, designed to investigate a wide range of physics phenomena, including the search for the Higgs boson, studies of the SM, and searches for new physics BSM. The ATLAS detector measures 46 m in length and 25 m in diameter, with a total weight of approximately 7000 tonnes. It is situated about 100 m underground at Point 1 on the LHC ring. The detector comprises multiple subsystems arranged in concentric layers: an inner tracking detector, electromagnetic and hadronic calorimeters, a muon spectrometer, and a magnet system consisting of a central solenoid and large toroidal magnets. This configuration allows for precise measurements of particle trajectories, energies, and momenta across a broad pseudorapidity¹ range [118].

The **Compact Muon Solenoid (CMS)** is the other general-purpose detector, complementary to ATLAS in its design and research objectives. CMS is optimised for high-precision measurements of muons, electrons, photons, and hadrons. The detector is 21.6 m long, 15 m in diameter, and weighs about 14000 tonnes, making it the heaviest of the LHC experiments. The CMS detector features a 3.8 T superconducting solenoid magnet, which enables precise momentum measurements of charged particles. Its layered structure includes an inner silicon tracker, electromagnetic and hadronic calorimeters, and extensive muon detection systems [119].

The **Large Ion Collider Experiment (ALICE)** is specifically designed to study

¹Pseudorapidity is defined as $\eta = -\ln \tan(\theta/2)$, where θ is the polar angle relative to the beam axis.

the quark–gluon plasma, a state of matter believed to have existed shortly after the Big Bang. ALICE focuses on heavy-ion collisions, particularly lead–lead interactions, to investigate the properties of quark–gluon plasma. The detector measures 26 m in length, 16 m in height and width, and weighs approximately 10000 tonnes. The ALICE complex setup includes a variety of subsystems for tracking, particle identification, and calorimetry, optimised for handling the high particle multiplicities produced in heavy-ion collisions [120].

The **Large Hadron Collider beauty (LHCb)** experiment specialises in studying the decays of b -hadrons to understand matter–antimatter asymmetries and to search for indirect evidence of new physics. Unlike the central detectors ATLAS and CMS, LHCb is a forward spectrometer covering the pseudorapidity range of approximately $2 < \eta < 5$. The detector is about 20 m long and 10 m high, with a design optimised for precise vertex reconstruction and particle identification. Key components include a silicon vertex locator, tracking stations, ring-imaging Cherenkov detectors, calorimeters, and muon chambers [121].

In addition to the four main detectors, the LHC hosts several smaller experiments, each focused on specific physics objectives:

- **Total, Elastic and Diffractive Cross-section Measurement (TOTEM):** Measures the total pp cross-section, elastic scattering, and diffractive processes using forward detectors positioned close to the beamline [122].
- **Large Hadron Collider forward (LHCf):** Aims to study neutral particles produced in the very forward region of pp collisions to improve the modelling of ultra-high-energy cosmic ray showers [123].
- **Monopole and Exotics Detector at the LHC (MoEDAL):** Searches for magnetic monopoles and other highly ionising particles using passive nuclear track detectors and trapping volumes [124].
- **ForwArd Search ExpeRiment (FASER):** Located 480 m downstream of the ATLAS interaction point, it is designed to detect long-lived, light, weakly interacting particles produced in the far-forward direction [125].
- **Scattering and Neutrino Detector at the LHC (SND@LHC):** A novel detector dedicated to measuring neutrinos produced at the LHC, located near the beam collision axis and optimised for forward physics studies [126].

Each of these experiments contributes uniquely to the physics programme of the LHC, enhancing our understanding of fundamental particles and forces.

3.2 The ATLAS detector

The ATLAS detector [118] is a general-purpose particle detector designed to explore a broad range of physics phenomena, from precision measurements within the SM to searches for new physics BSM. This ambitious scientific endeavour is enabled by the ATLAS collaboration, comprising over 5000 members from more than 180 institutions across 40 countries, including approximately 3000 scientific authors. The ATLAS collaboration represents one of the largest and most complex collaborative scientific efforts ever undertaken.

At its core lies the **Inner Detector (ID)**, which reconstructs the trajectories of charged particles with high precision in a 2T magnetic field. Surrounding the ID are the **calorimeter systems**, subdivided into an **electromagnetic calorimeter**, optimised for the measurement of photons and electrons, and a **hadronic calorimeter**, designed to measure the energy of strongly interacting particles such as hadrons and jets. The outermost subsystem is the **Muon Spectrometer (MS)**, which detects and measures the momentum of muons — the only charged particles capable of traversing the full detector without being absorbed.

A simplified schematic of the ATLAS detector layout is presented in Figure 3.3, illustrating the overall structure and placement of these key subdetectors.

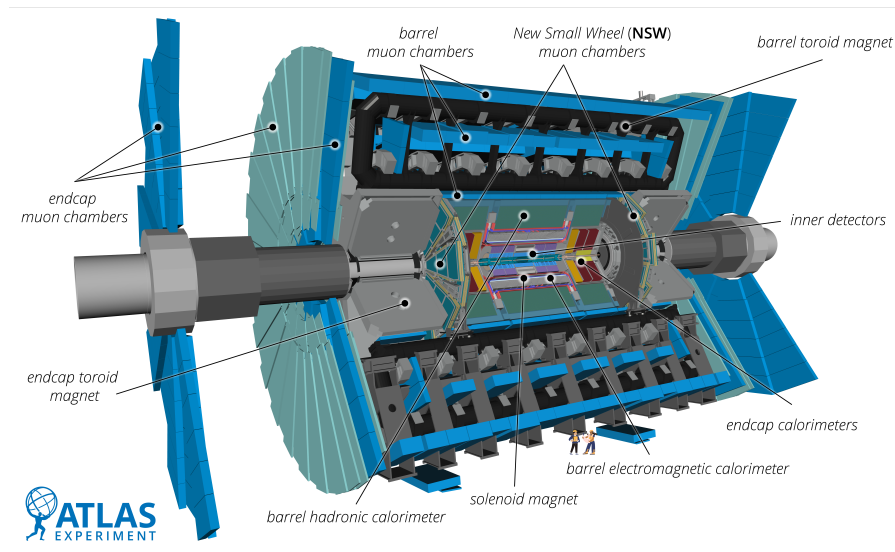


Figure 3.3: Illustrative diagram of the ATLAS detector, with annotations indicating its key components [127].

3.2.1 Coordinate System of the ATLAS Detector

The ATLAS detector is described using a right-handed Cartesian coordinate system (x, y, z) , with the origin located at the nominal interaction point at the center of the detector. The z -axis is defined to lie along the beam direction, the x -axis points radially inward toward the center of the LHC ring, and the y -axis points vertically upward.

In practice, it is often more convenient to use a cylindrical coordinate system (r, ϕ, θ) . Here, r denotes the transverse radial distance from the beam axis, ϕ is the azimuthal angle measured in the x - y plane, and θ is the polar angle measured from the beam axis (the z -axis). The transverse plane is defined as the plane perpendicular to the beam direction, and is central to many analyses.

To describe particle kinematics in a way that is invariant under Lorentz boosts along the beam axis, the concept of rapidity y is used. It is defined as:

$$y = \frac{1}{2} \ln \left(\frac{E + p_z}{E - p_z} \right), \quad (3.4)$$

where E is the energy of the particle and p_z is its momentum component along the beam axis.

At LHC energies, most particles are ultra-relativistic, meaning their rest mass can be neglected. In this limit, rapidity is commonly approximated by the pseudorapidity η , which depends only on the polar angle θ :

$$\eta = -\ln \tan(\theta/2). \quad (3.5)$$

This definition implies that $\eta = 0$ corresponds to transverse emission ($\theta = \pi/2$), while $|\eta| \rightarrow \infty$ corresponds to emission along the beamline ($\theta \rightarrow 0$ or π). The ATLAS detector provides coverage for particles within the range $|\eta| < 4.9$, enabling nearly full solid-angle detection.

The angular separation between two particles is typically quantified by the variable ΔR , defined as:

$$\Delta R = \sqrt{(\Delta\eta)^2 + (\Delta\phi)^2}. \quad (3.6)$$

Another important kinematic quantity is the p_T , which represents the component of a particle's momentum in the transverse plane. It is defined as:

$$p_T = \sqrt{p_x^2 + p_y^2}, \quad (3.7)$$

and is particularly useful in hadron collider environments where the initial net momentum in the transverse plane is approximately zero. The transverse momentum is typically reconstructed using energy deposits in the calorimeters and plays a crucial role in the identification and characterisation of final-state particles.

The full momentum of a particle in the detector is often described using the triplet (p_T, η, ϕ) . A related quantity of interest is the missing transverse momentum, which is defined as the negative vectorial sum of all reconstructed transverse momenta in the event (see Section 5.7 for more details). This variable is essential in analyses involving undetected particles, such as neutrinos or potential candidates for dark matter.

3.2.2 ATLAS Inner Detector

The ID is a high-precision tracking spectrometer located directly around the beam pipe and constitutes the closest detector system to the interaction point. It is embedded in a 2 T axial magnetic field generated by the central solenoid, which induces curvature in the trajectories of charged particles, allowing for accurate measurement of their momenta and impact parameters. This capability is essential for the reconstruction of both primary and secondary vertices.

The ID provides full tracking coverage in the pseudorapidity range up to $|\eta| < 2.5$, with excellent spatial resolution and high track-reconstruction efficiency. It measures approximately 7 m in length and 2.3 m in diameter. The system is composed of four nested subdetectors: the Insertable B-Layer (IBL), the Pixel Detector, the Semiconductor Tracker (SCT), and the Transition Radiation Tracker (TRT). These subdetectors are arranged in a barrel-and-endcap geometry to ensure uniform coverage over the full interaction region.

A schematic illustration of the ID layout in the barrel region is shown in Figure 3.4, while Table 3.2 summarises the key features and design parameters of each subdetector.

A brief description of each component is provided below:

- **IBL:** The IBL is the innermost subdetector of the ID, positioned closest to the interaction point. It was installed during Long Shutdown 1 and has been operational since the start of Run 2. Its primary function is to enhance vertex resolution and improve track reconstruction performance by providing an additional two-dimensional space point per track near the beamline. The IBL consists of 280 Pixel modules mounted on 14 azimuthal staves, which are tilted to overlap by 1.82° to ensure full coverage in the ϕ direction. It extends the tracking acceptance up to

$$|\eta| < 3.$$

- **Pixel Detector:** The Pixel Detector surrounds the IBL and forms the innermost component of the original ID design. It consists of 1774 identical modules arranged into three barrel layers concentric with the beam pipe, and two endcap sections, each comprising three disks oriented perpendicular to the z -axis. The detector provides full azimuthal coverage and extends the tracking acceptance up to $|\eta| < 2.5$. It delivers three high-precision, two-dimensional space points per track, playing a critical role in vertex reconstruction and pattern recognition.
- **SCT:** The SCT shares a similar geometry with the Pixel Detector and is composed of 4088 modules arranged in four concentric barrel layers and two endcap assemblies, each containing nine disks. Each module consists of two single-sided silicon microstrip sensors mounted back-to-back with a stereo angle of 40 mrad, enabling two-dimensional position measurements in the r - ϕ and z directions. The SCT provides full azimuthal coverage and extends up to $|\eta| < 2.5$. On average, it contributes four precise two-dimensional space points per charged-particle track, significantly enhancing track resolution and redundancy.
- **TRT:** The TRT is the outermost subdetector of the ID and is composed of 350,848 gas-filled straw tubes, each 4 mm in diameter. These tubes are arranged parallel to the beam axis in the barrel region and radially in the two endcaps. In the barrel, the spaces between the straws are filled with polymer fibres, while foils are used in the endcaps. As particles traverse the interfaces between different materials, they can emit transition radiation, which is absorbed by the detector gas, resulting in an enhanced readout signal. This feature allows the TRT to provide electron-pion discrimination in addition to its tracking capability. It delivers approximately 30 one-dimensional measurements per track in the r - ϕ plane and covers the pseudorapidity range up to $|\eta| < 2$, although it does not provide measurements along the z -axis.

3.2.3 Calorimeters

Surrounding the ID, the ATLAS calorimeter system is designed to absorb and measure the energy of most particles produced in pp collisions. These calorimeters play a crucial role in identifying and reconstructing the energy of electrons, photons, and hadrons. Among SM particles, only neutrinos and high p_T muons escape full containment in the calorimeters. Neutrinos are inferred indirectly through the imbalance in transverse

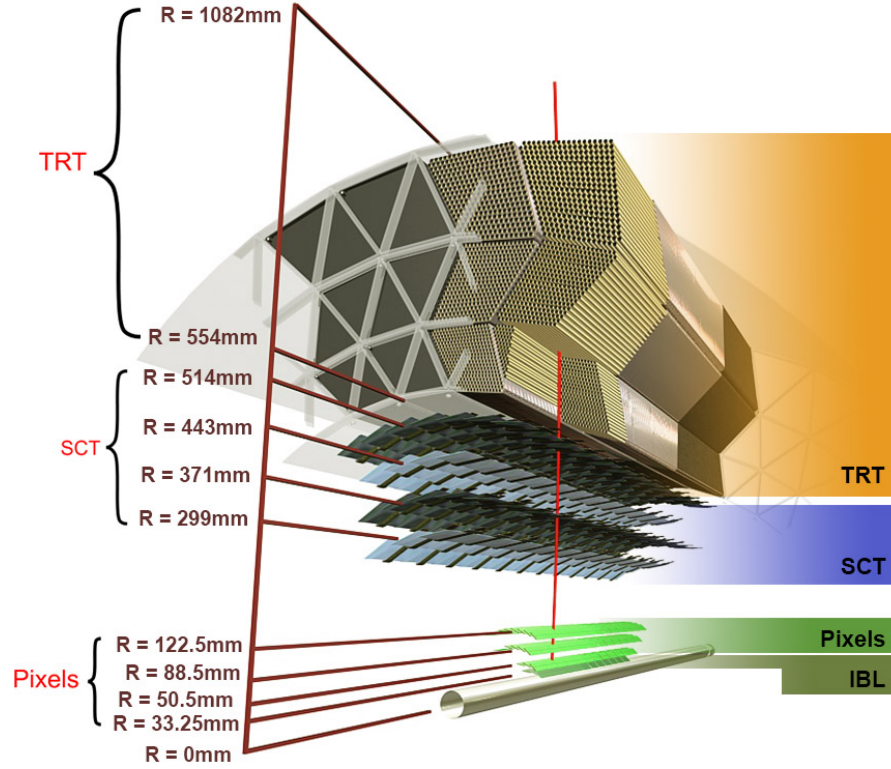


Figure 3.4: The ID barrel region is depicted in a schematic, showcasing the central beam-pipe, followed by the IBL, three Pixel layers, then the SCT four cylindrical layers, and finally the TRT 72 straw layers. [128]

energy—reconstructed as missing transverse momentum—while high- p_T muons are subsequently tracked in the MS.

The calorimeter system employs a sampling design, consisting of alternating layers of passive absorber material and active detector medium. The absorber initiates particle showers by halting incident particles, while the active material detects the resulting secondary particles and converts their passage into measurable signals.

The calorimeters are divided into three main subsystems: the Electromagnetic Calorimeter (ECAL), which measures the energy of electrons and photons; the Hadronic Calorimeter (HCAL), which absorbs and reconstructs the energy of strongly interacting particles such as charged pions and protons; and the Forward Calorimeter (FCal), optimised for the detection of high-energy particles produced at small angles relative to the beam axis. A schematic overview of the calorimeter subcomponents is shown in Figure 3.5.

ECAL: The ECAL covers the pseudorapidity range up to $|\eta| < 3.2$ and is designed to achieve a high energy resolution $\sigma(E)/E = 10\%/\sqrt{E} \oplus 170 \text{ MeV}/\sqrt{E} \oplus 0.7\%$ [130].

Subdetector	Element size [μm]	Intrinsic resolution [μm]
IBL	50×250	10×60
Pixel	50×400	10×115
SCT	80	17×580
TRT	4000	130

Table 3.2: Summary of the main characteristics of the ID subdetectors [113, 129]. In addition to its element size, the intrinsic resolution along r - ϕ and z is reported, except for the TRT which is insensitive in the z direction. The element sizes in the SCT refer to the spacing between readout strips (barrel modules), while in the TRT they refer to the diameter of its tubes.

It is segmented into a central barrel region ($|\eta| < 1.475$) and two endcap regions ($1.375 < |\eta| < 3.2$), with the barrel itself divided into two halves separated at $z = 0$. The endcaps are composed of two coaxial wheels that extend the coverage into the forward regions.

The calorimeter employs lead as the absorber material, arranged in an accordion geometry to guarantee full azimuthal coverage without gaps in ϕ . The active medium is liquid argon (LAr), forming what is referred to as the LAr calorimeter. To maintain the LAr at cryogenic temperatures, the entire ECAL system is enclosed within a cryostat.

HCAL: The HCAL covers the region up to $|\eta| < 3.2$, with a design energy resolution of approximately $\sigma(E)/E = 52.9\%/\sqrt{E} \oplus 5.7\%$ [130]. Like the ECAL, it is segmented into a barrel ($|\eta| < 1.7$) and two endcap regions ($1.5 < |\eta| < 3.2$).

The central barrel region employs steel as the absorber and plastic scintillator tiles as the active material, forming what is known as the Tile Calorimeter (TileCal). Its structure consists of a central barrel flanked by two extended barrels to cover the full azimuthal and pseudorapidity acceptance. The endcaps use liquid argon as the active medium and copper as the absorber, arranged into two wheels per side. These share the same cryostat infrastructure as the ECAL endcaps.

The HCAL is designed to be sufficiently thick to contain hadronic showers and suppress punch-through of high-energy particles into the MS. It plays a vital role in ensuring accurate reconstruction of hadronic energy and contributes significantly to the resolution of missing transverse energy.

FCALs: The FCALs are integrated into the liquid argon endcap cryostats and extend calorimetric coverage into the high-pseudorapidity region, spanning $3.1 < |\eta| < 4.9$ [130]. Each FCAL unit consists of three longitudinal modules: the first uses copper as the

absorber material for the measurement of electromagnetic showers from electrons and photons, while the remaining two use tungsten absorbers to measure the energy of hadronic particles.

The FCAL provides essential coverage in the forward region, improving the overall hermeticity of the detector. It plays a key role in reconstructing missing transverse momentum and enables efficient identification of forward jets, both of which are important in precision measurements and searches for phenomena BSM.

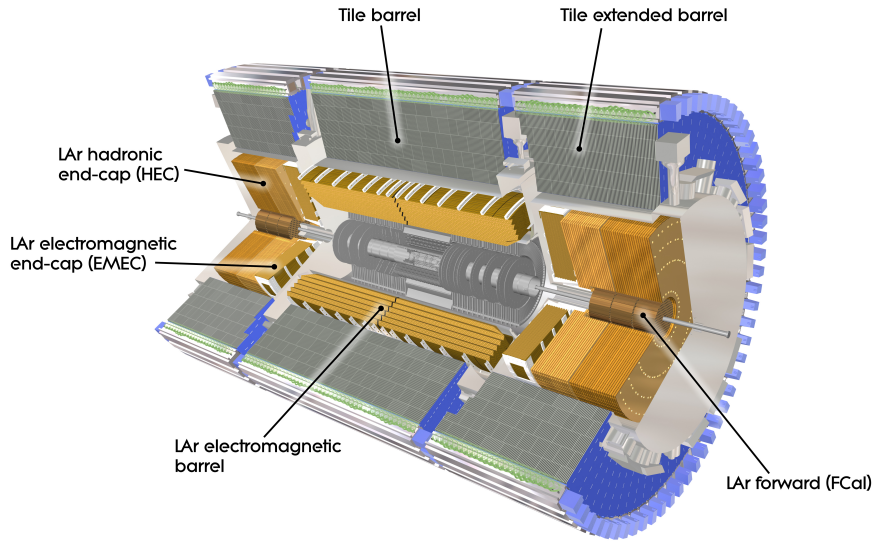


Figure 3.5: Schematic illustration of the ATLAS calorimeter system, showing the subcomponents of the electromagnetic, hadronic, and forward calorimeters [130].

3.2.4 The Muon Spectrometer

The MS is a crucial subsystem of the ATLAS detector, designed for precise reconstruction and identification of muons and for the trigger (see Section 3.2.6). It provides full azimuthal coverage and spans the pseudorapidity range up to $|\eta| < 2.7$, using three cylindrical barrel layers and six endcap disks. As muons deposit only minimal energy in the calorimeters, the MS enables their detection by measuring the curvature of their trajectories in a toroidal magnetic field, similar in concept to the solenoidal field used in the ID.

The spectrometer employs multiple detector technologies. Monitored Drift Tubes (MDTs) are used throughout most of the system to provide high-resolution position measurements in the bending plane. In the forward region ($2 < |\eta| < 2.7$), Cathode Strip

Chambers (CSCs) are used due to their finer granularity and ability to handle higher particle flux. Resistive Plate Chambers (RPCs) in the barrel and Thin Gap Chambers (TGCs) in the endcaps serve as both trigger detectors and as providers of complementary spatial measurements in the coordinate orthogonal to bending. An upgrade during Long Shutdown 1 introduced the New Small Wheels, incorporating Micromegas and small-strip TGCs (sTGCs) to improve tracking and trigger performance in the forward region. A schematic view of the MS can be found in Figure 3.6.

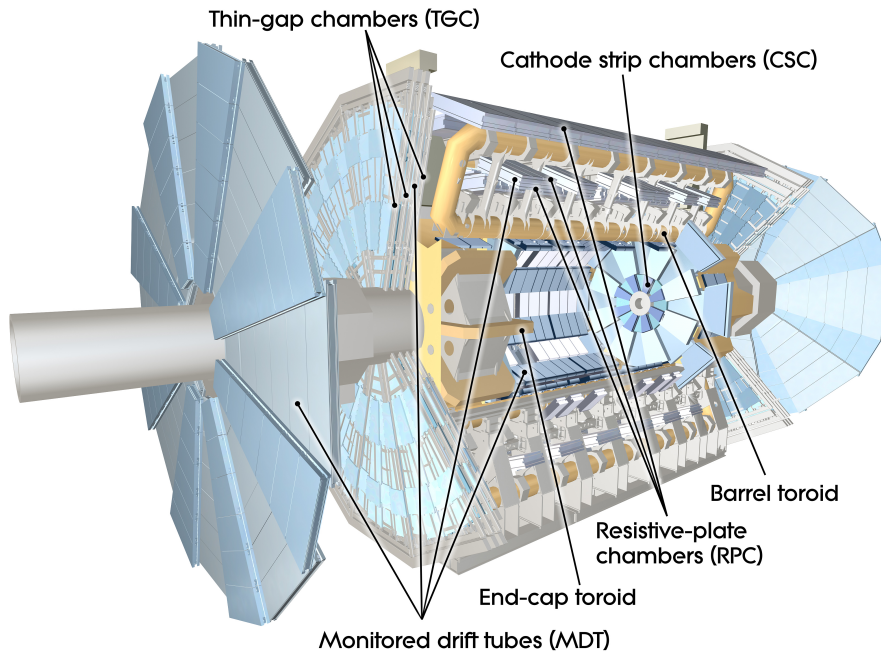


Figure 3.6: Cut-away view of the MS, including the different types of detectors. The toroid magnets are also displayed [131].

3.2.5 Magnet System

The ATLAS magnet system is a hybrid configuration that enables independent magnetic field generation for the ID and the MS. Within the ID, a superconducting solenoid magnet produces a uniform axial magnetic field of 2 T, which enables precise momentum measurements of charged particles in the central tracking volume.

The MS relies on a large air-core toroidal system composed of a barrel toroid and two endcap toroids. Each toroid consists of eight superconducting coils arranged symmetrically to generate a strong and non-uniform magnetic field optimised for muon momentum reconstruction. The magnetic field reaches peak values of approximately 3.9 T in the barrel and 4.1 T in the endcap regions. This configuration provides bending power over

a wide pseudorapidity range while minimising material in front of the muon detectors, which is essential for maintaining track resolution. Figure 3.7 shows a schematic view of the magnet system.

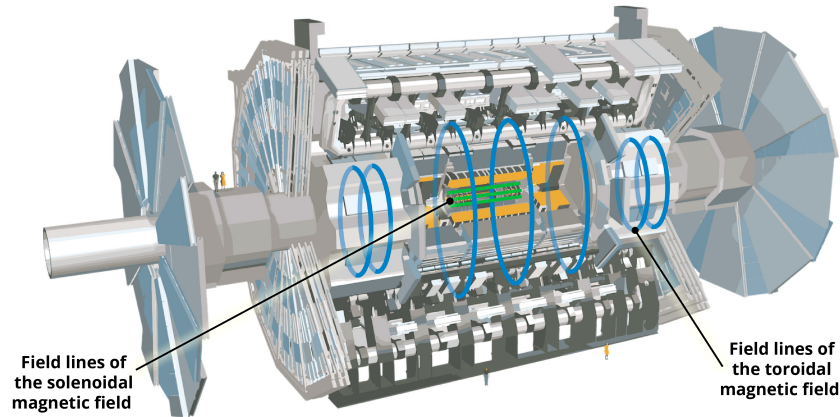


Figure 3.7: Schematic view of the ATLAS magnet system [132]. The field lines of the solenoidal magnet are shown in green, while the field lines of the toroidal magnets are shown in blue.

3.2.6 Trigger and Data Acquisition

The ATLAS trigger and data acquisition (DAQ) system [133–135] plays a central role in selecting and recording events of interest from the vast amount of data produced by the detector. During Run 2 of the LHC, proton bunches crossed every 25 ns, corresponding to a bunch-crossing rate of 40 MHz and resulting in approximately 25 interactions per crossing. This generated a raw data rate exceeding 60 TB/s, which far exceeds the capacity for storage and offline processing.

To reduce this data volume to manageable levels, ATLAS employs a two-tiered online trigger system composed of the Level-1 (L1) trigger and the High-Level Trigger (HLT). The L1 trigger is implemented in custom hardware and uses reduced granularity information from all calorimeters and dedicated muon trigger chambers to make a fast decision within $2.5 \mu\text{s}$. If an event is accepted, L1 generates Regions of Interest (RoIs), which include the coordinates and transverse momenta of candidate objects. Accepted events are forwarded to the Readout System (ROS) through optical links connected to the front-end electronics of the detector via Readout Drivers (RODs). The L1 system can sustain an output rate of up to 100 kHz.

The HLT is a software-based system that further refines the L1 selection using full

detector information available in the ROS. It performs detailed reconstruction of event topologies and applies algorithms to identify final-state particles. The HLT reduces the event rate to approximately 1.5 kHz, corresponding to a data throughput of about 1.5 GB/s. These selected events are recorded to disk and subsequently transferred to the Worldwide LHC Computing Grid (WLCG), where they are processed and made available for physics analyses within the ATLAS collaboration. An overview of the trigger and DAQ system is shown in Figure 3.8.

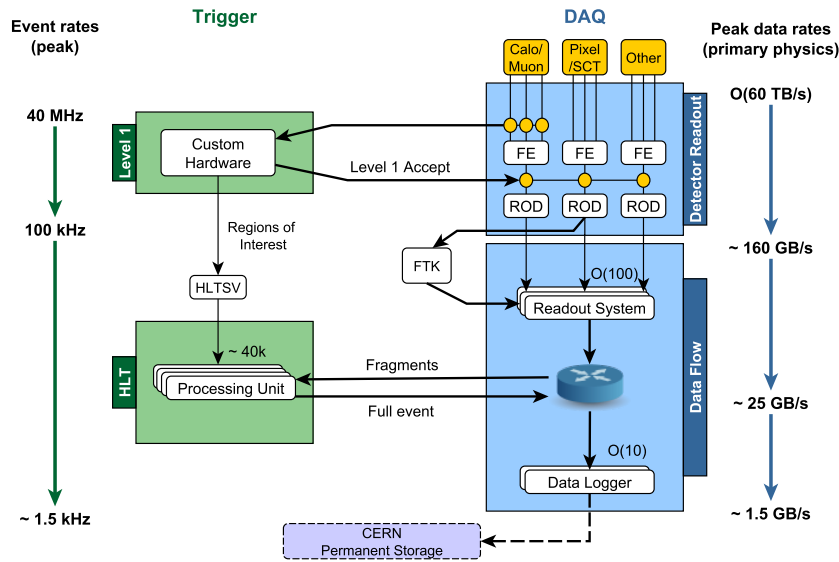


Figure 3.8: Functional diagram of the ATLAS Trigger and DAQ system during Run 2, indicating the expected peak rates and bandwidths through each component. Events accepted by the L1 hardware trigger (top left) are transferred to the HLT (bottom left) via the farm supervisor node (HLTSV), which also assembles RoIs. Simultaneously, event data from the detector front-end electronics are transmitted to the ROS through optical links originating from the RODs in response to the L1 accept signal. These data are buffered within the ROS and accessed by algorithms executed in the HLT. Events accepted by the HLT are then sent to permanent storage through the Data Logger [136].

3.2.7 Computing System

The LHC produces an enormous volume of data annually, estimated between 50 and 70 PB. This volume exceeds the processing and storage capabilities of CERN alone. To meet this demand, a global distributed computing infrastructure known as Grid computing was adopted [137]. This model combines resources from institutions around the world, integrating computing, storage, and networking capacity. It led to the development of

the WLCG [138], which enables reliable, large-scale data processing and analysis for all LHC experiments.

The WLCG is organised into a hierarchical structure consisting of four tiers. The Tier-0 level, located at CERN, is responsible for the archival of raw data, prompt reconstruction, and data distribution to Tier-1 sites. The Tier-1 network includes thirteen major computing centres globally. These facilities store and reprocess data, and distribute processed samples to Tier-2 centres. The Tier-2 layer comprises approximately 160 research institutes and universities, which support user-level data analysis and MC event simulation. Finally, the Tier-3 level consists of local computing resources dedicated to individual scientists or small research teams for specific analysis tasks. The Instituto de Física Corpuscular (IFIC) in Valencia holds a Tier-2 computing centre as well as a Tier-3

This globally distributed computing model enables the high-throughput processing and efficient management of LHC data, ensuring that researchers worldwide have timely access to the resources needed for large-scale analysis.

3.3 Performance of the ATLAS Detector

Precise reconstruction of the trajectories of charged particles produced in pp collisions at the LHC is fundamental to the physics programme of the ATLAS experiment. This reconstruction is critical for a wide range of analyses, including SM measurements and searches for physics BSM.

The ATLAS ID is the central tracking system responsible for reconstructing the paths of charged particles. As these particles traverse the detector volume, they leave discrete energy deposits—referred to as hits—in the sensitive components of the ID. These hits are used to reconstruct trajectories and extract key track parameters, such as momentum, direction, and impact parameter. The precision of these measurements depends on several factors, including the intrinsic resolution of the detector sensors, the accuracy of the magnetic field description, and, importantly, the knowledge of the positions and orientations of the detector elements.

The alignment of the ID aims to determine the true geometry of the active tracking elements with high precision and to monitor any changes over time. Inaccurate knowledge of the detector geometry or uncorrected mechanical shifts—referred to as misalignments—can significantly degrade the resolution of reconstructed track parameters. Furthermore, even when residuals are minimised, certain correlated geometric deformations, known as weak modes, can introduce systematic biases in track

reconstruction. These biases may arise from real structural deformations or limitations in the alignment procedure itself.

Understanding and mitigating such weak modes is a central aspect of ensuring optimal tracking performance in ATLAS [128].

3.3.1 ID Alignment

Having explored the ID subsystems, it is crucial to address the concept of alignment [128]. The ATLAS detector employs a track-based alignment procedure that minimises the residuals between the reconstructed particle trajectories and the corresponding measured hit positions in the tracking detectors.

The ID is comprised of a multitude of subsystems, each capable of being aligned independently. With nearly 750000 degrees of freedom to consider, the ID alignment process is structured hierarchically, spanning from global mechanical assembly structures down to local sensor alignments, and encompasses different timescales of detector movements. The alignment routine utilises a track-based algorithm that minimises a χ^2 function of the track-to-hit residuals. However, this method has limited sensitivity to certain global detector deformations that may alter the reconstructed track parameters without significantly affecting the overall χ^2 of the fitted tracks. These deformations are referred to as weak modes as mentioned already. Mitigating these weak modes is an essential aspect of the alignment procedure.

The accurate functioning of these subsystems is not only dependent on their individual performance but also on their relative spatial orientations. In other words, the position of each sensor must be precisely known in the global reference frame of the detector for the data to be interpreted correctly.

Alignment in the experiment is of paramount importance. Even a small misalignment can introduce substantial systematic uncertainties into track measurements. For instance, if the drift tubes in the ID are misaligned by merely a millimeter, it could yield an uncertainty in the momentum measurement of $\Delta p_T/p_T \sim \mathcal{O}(10^{-4})$, which, in turn, impacts event reconstruction and the subsequent estimation of physical observables.

In the context of the ATLAS experiment, ID alignment aims to minimise these kinds of uncertainties. The goal is to ensure that the positions of all sensitive elements are known to an accuracy that is much smaller than their intrinsic resolutions. Therefore, alignment serves as a foundational pre-processing step that calibrates the ATLAS detector for high-precision measurements.

The fundamental principles underlying this alignment methodology are outlined in the following paragraphs.

Local Coordinate System: The positions of the subdetectors that constitute the ATLAS ID are defined relative to the global coordinate system (x, y, z) , as described in Section 3.2.1. In contrast, the positions of hits and track parameters within individual subdetector modules are expressed in a local coordinate system (x', y', z') . This local system is centred at the geometric centre of each module, with the x' -axis aligned along the most sensitive measurement direction of the module. A schematic illustration of the global and local coordinate systems is provided in Figure 3.9.

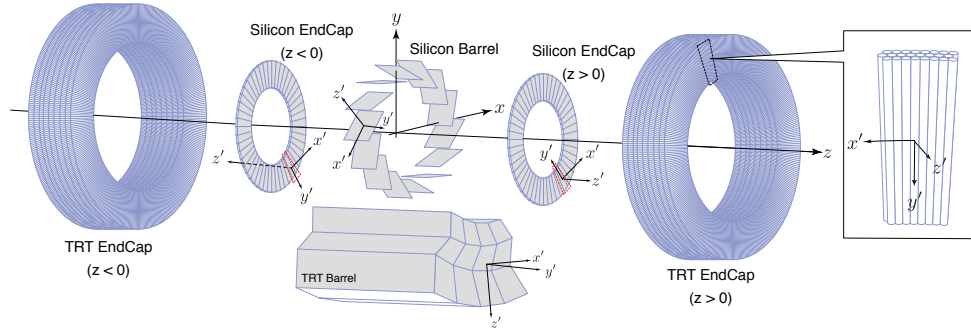


Figure 3.9: Schematic illustration of the ATLAS global coordinate system (x, y, z) , along with the local coordinate system (x', y', z') used for individual subdetector modules [128].

Track Parameters: In the presence of a magnetic field, charged particles follow a helical trajectory, which can be described by a set of five track parameters defined with respect to the global coordinate system. These parameters are typically evaluated at the perigee, the point of closest approach of the track to the global z -axis.

The five parameters are:

- ϕ_0 : the azimuthal angle of the momentum vector at the perigee, measured in the transverse plane with respect to the x -axis, in the range $[-\pi, \pi)$,
- θ : the polar angle with respect to the z -axis, in the range $[0, \pi]$,
- d_0 : the transverse impact parameter, which is the distance of closest approach to the beamline in the x - y plane,
- z_0 : the longitudinal impact parameter, corresponding to the z -position of the perigee,
- q/p : the electric charge of the particle divided by its momentum, which determines the curvature of the helical trajectory.

Additionally, the polar angle θ is often replaced by the $\eta = -\ln \tan(\theta/2)$, which provides a boost-invariant measure along the beam axis.

Alignment Parameters: The position and orientation of a rigid detector element can be fully described by six degrees of freedom, referred to as the alignment parameters:

$$\vec{a} = (T_x, T_y, T_z, R_x, R_y, R_z), \quad (3.8)$$

where T_x, T_y, T_z denote translations along the local coordinate axes, and R_x, R_y, R_z represent rotations about these axes. Each module of the ID can be aligned by adjusting these six parameters relative to its nominal geometry.

Residuals: The track-to-hit residual, denoted by r , is defined as the distance between the measured position of a hit in a detector element and the corresponding intersection point of the reconstructed track. It can be expressed as:

$$r = m - e(\vec{\tau}, \vec{a}), \quad (3.9)$$

where m is the position of the measured hit, and $e(\vec{\tau}, \vec{a})$ is the extrapolated intersection point of the reconstructed track. The latter depends on the track parameters τ and the alignment parameters $\vec{a}$. Residuals are evaluated within the local coordinate system of each subdetector module.

A schematic illustration of the track-to-hit residuals produced by a charged particle traversing a planar, rigid detector is shown in Figure 3.10.

Track-Based Alignment Fit: The alignment of the ID is performed by minimising a global χ^2 function constructed from the track-to-hit residuals. This track-based approach uses a global fitting algorithm [139] to determine corrections to the positions of detector modules and mechanical structures.

The χ^2 function is defined as:

$$\chi^2 = \sum_e \sum_t \sum_h \left(\frac{r_{t,h}(\vec{\tau}, \vec{a})}{\sigma_h} \right)^2 + C(\vec{\tau}) + C'(\vec{a}), \quad (3.10)$$

where the sums run over all events e , reconstructed tracks t , and associated hits h . The quantity $r_{t,h}$ represents the residual between the reconstructed track and the hit position, which depends on both the track parameters $\vec{\tau}$ and the alignment parameters $\vec{a}$. The term σ_h denotes the uncertainty associated with each hit. The additional terms $C(\vec{\tau})$ and $C'(\vec{a})$ introduce constraints on the track and alignment parameters, respectively. These constraints encode prior knowledge and help stabilise the fit in regions where the algorithm may be insensitive to certain misalignments or global deformations. Such limitations will be discussed further in the following subsection.

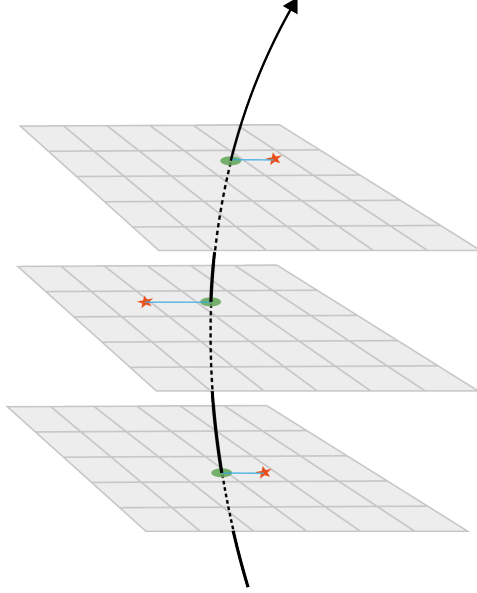


Figure 3.10: Schematic illustration of a charged particle traversing successive detector planes. The hit positions recorded by each detector layer are marked with red stars. The reconstructed particle trajectory is shown as a black line. Green ellipses represent the points at which the fitted track intersects each detector surface, and blue lines indicate the residuals, defined as the distances between the measured hits and the corresponding track intersections [128].

3.3.2 Weak Modes

Unlike standard misalignments that modify the track-to-hit residuals and thus impact the overall χ^2 of the track-based fit, weak modes are correlated detector deformations that leave the χ^2 function largely unchanged. Despite this, they can introduce systematic biases in both the reconstructed track parameters and the alignment corrections. Figure 3.11 provides a classification of the most relevant weak modes in the ID, along with their possible impact on physics observables.

Momentum-related biases induced by coherent detector deformations can generally be grouped into two broad categories:

- **Sagitta distortions**, which are deformations in the bending plane that affect the curvature of reconstructed tracks differently for positively and negatively charged particles (see Figure 3.12 (a)).
- **Length scale distortions**, which occur along the particle trajectory and bias the curvature in the same way for both charge signs (see Figure 3.12 (b)).

As described in Section 3.3.1, external constraints introduced via the terms $C(\vec{\tau})$ and $C'(\vec{a})$ can be applied to reduce the influence of weak modes. Residual momentum biases,

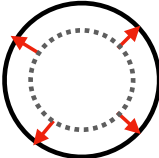
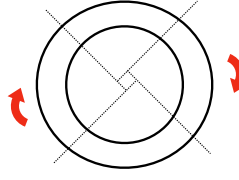
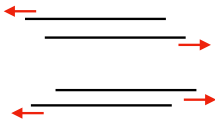
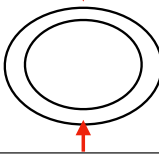
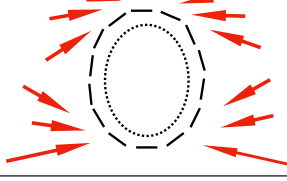
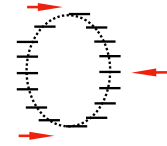
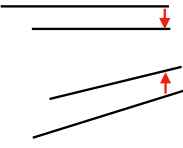
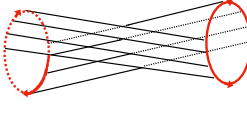
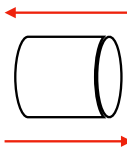
	ΔR	$\Delta \phi$	ΔZ
R	Radial Expansion (distance scale) 	Curl (charge asymmetry) 	Telescope (CM boost) 
ϕ	Elliptical (vertex mass) 	Clamshell (vertex displacement) 	Skew (Z momentum) 
Z	Bowing (total momentum) 	Twist (vertexing) 	Z expansion (distance scale) 

Figure 3.11: Summary table of the possible systematic deformations, the so-called weak modes [128].

after alignment corrections, are typically small enough that their effects can be accounted for through direct correction of the track parameters.

Understanding the behaviour of weak modes requires knowledge of how a charged particle propagates in a magnetic field, where its trajectory takes the form of a helix. The transverse momentum of a particle traversing a cylindrical detector of radius R is related to the sagitta s , which quantifies the deviation of the track from a straight trajectory, as:

$$p_T = 0.3 q B \rho = 0.3 q B \sqrt{8sR + 2s^2}, \quad (3.11)$$

where q is the electric charge, B is the magnetic field strength, and ρ is the radius of curvature. For high- p_T particles, the sagitta is small relative to R , allowing the expression to be approximated as:

$$p_T \approx \frac{0.3 q B R^2}{8s}. \quad (3.12)$$

From this relation, it is clear that the reconstructed transverse momentum can be affected by biases in the magnetic field, the sagitta, or the detector geometry itself. Misalignments between the tracking volume and the magnetic field—referred to as magnetic field biases—were investigated in Run 1 [140]. Detailed studies of sagitta distortions were performed in Run 2, as discussed in Ref. [141].

Radial distortions represent another class of weak modes and are characterised by a linear expansion or contraction of detector layers with respect to radius (see Figure 3.12 (b)). These can be constrained using well-known physics processes. For instance, the invariant mass of muon pairs from the decay of the Z boson ($Z \rightarrow \mu^+\mu^-$) provides a high-statistics and well-calibrated reference. Other resonances, such as the J/ψ and Υ , are also used to further constrain such distortions.

A detailed treatment of momentum biases and correction strategies is presented in the following subsection.

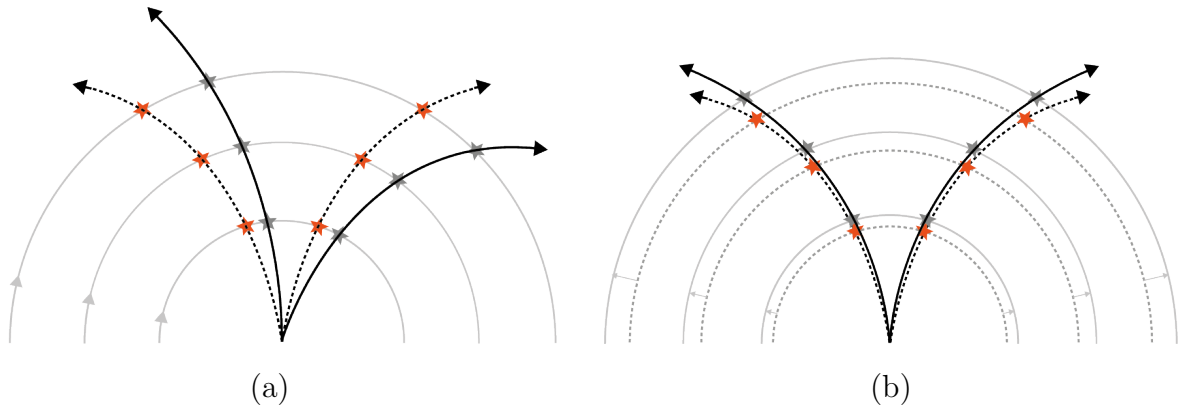


Figure 3.12: Schematic illustration of two representative weak modes: a detector curl (a), which induces charge-asymmetric effects, and a radial expansion (b), which affects the reconstruction symmetrically for both charges. Grey lines indicate the nominal detector geometry, while arrows represent the impact of each deformation on momentum reconstruction. Red stars denote the true hit positions, and grey stars indicate the reconstructed positions. The black curves show the trajectories of oppositely charged particles [128].

Momentum Biases: Momentum biases in track reconstruction can arise due to coherent detector deformations that do not significantly affect the overall χ^2 of the track-based fit. These biases fall into two main categories: sagitta distortions, which are charge-antisymmetric and affect track curvature differently for particles of opposite charge, and length scale biases, which are charge-symmetric and alter the curvature uniformly.

To constrain such effects, several complementary observables are employed. The E/p ratio, defined as the calorimeter energy E divided by the track momentum p , provides sensitivity to sagitta-like distortions. This observable is constructed for electrons and positrons by comparing the energy measured in the ECAL to the momentum reconstructed by the tracking system. Additionally, low-mass resonances such as the J/ψ offer constraints at lower p_T .

Decays of the Z boson into muons ($Z \rightarrow \mu^+\mu^-$) serve as a key reference for monitoring and correcting sagitta biases. These events offer multiple advantages: a large and clean sample, minimal energy loss as muons traverse the detector, and extended path lengths that reach the MS. As a result, $Z \rightarrow \mu^+\mu^-$ decays are highly effective for probing track curvature across a wide kinematic range and for identifying residual misalignments that lead to momentum biases.

3.3.3 Sagitta bias measurement using $Z \rightarrow \mu^+\mu^-$ decays

The trajectory curvature of a charged particle in a magnetic field allows for the determination of its momentum. Within a uniform magnetic field, the particle traces a circular path with a radius r . The Lorentz force $\vec{F}$ influencing it is given by:

$$\vec{F} = q\vec{v} \times \vec{B}. \quad (3.13)$$

Here, $\vec{v}$ denotes the particle velocity, q represents the charge, and $\vec{B}$ the encompassing magnetic field. For trajectories perpendicular to this magnetic field, the formula becomes:

$$|\vec{F}| = q|\vec{v}|\sin(\theta)|\vec{B}|, \quad \frac{mv^2}{r} = qvB. \quad (3.14)$$

In this case, θ is the angle between the magnetic field and the particle momentum vector, m is the particle mass and B is the magnitude of $\vec{B}$. Recognising that $mv = p_T$, the transverse momentum can be described as:

$$p_T = qBr. \quad (3.15)$$

Through basic geometric principles, r can be expressed in terms of the measurement radius R of the detector and the sagitta s , which is the track deflection from a straight path:

$$r = \frac{R^2}{8s} + \frac{s}{2}. \quad (3.16)$$

For scenarios where $R \gg s$, the equation simplifies to:

$$p_T \approx qB \frac{R^2}{8s}. \quad (3.17)$$

Assuming that actual and estimated tracks possess distinct sagittas, define ϵ_s as the sagitta discrepancy, so that $s = s_0 + \epsilon_s$. The inverse of the reconstructed p_T is expressed as:

$$\frac{1}{p_T} = \frac{8s}{qBR^2} = \frac{8}{qBR^2}(s_0 + \epsilon_s). \quad (3.18)$$

Thus, the equation for the reconstructed p_T becomes

$$p_T = p_{T0}(1 + qp_{T0}\delta_s)^{-1}, \quad (3.19)$$

where $p_{T0} = \frac{qBR^2}{8s_0}$, and $\delta_s = \frac{8\epsilon_s}{BR^2}$ represents the sagitta deviation.

For the longitudinal calculations, from Equation 3.19, assuming p_{z0} and $|p_0|$ are positive, the expression becomes:

$$p_z = p_T \sinh(\eta) = p_{T0}(1 + qp_{T0}\delta_s)^{-1} \sinh\left(\operatorname{arctanh}\left(\frac{p_{z0}}{|p_0|}\right)\right). \quad (3.20)$$

This simplifies to:

$$p_z = p_{z0}(1 + qp_{T0}\delta_s)^{-1}. \quad (3.21)$$

For the three-momentum:

$$p_T^2 + p_z^2 = p^2, \quad (3.22)$$

$$p_{T0}^2(1 + qp_{T0}\delta_s)^{-2} + p_{z0}^2(1 + qp_{T0}\delta_s)^{-2} = p_0^2(1 + qp_{T0}\delta_s)^{-2}, \quad (3.23)$$

$$p = p_0(1 + qp_{T0}\delta_s)^{-1}, \quad (3.24)$$

where p and p_0 are reconstructed and true momentum values, respectively. The sagitta correction δ_s is proportional to the transverse momentum p_{T0} .

During Run 2, the ATLAS ID employed the baseline mass method for measuring sagitta bias. This method involves iterative steps to minimise the effective difference between reconstructed and invariant dimuon mass.

The decay $Z \rightarrow \mu^+\mu^-$ is selected using the following event criteria. The invariant mass of the dimuon pair is constrained to lie within the range $70 \text{ GeV} < m_{\mu\mu} < 110 \text{ GeV}$, ensuring selection of events near the Z -boson peak. Each muon is required to have a $p_T > 20 \text{ GeV}$ and lie within $|\eta| < 2.5$. The opening angle between the muons must satisfy $\gamma(\mu^+\mu^-) > 40.5^\circ$. Additionally, the impact parameter differences are required to fulfil $\Delta d_0(\mu^+\mu^-) < 0.1 \text{ mm}$ and $\Delta z_0(\mu^+\mu^-) < 0.6 \text{ mm}$, to ensure the muons originate from a common interaction point.

3.3.4 Mass method for measuring sagitta bias

The invariant mass squared of a pair of relativistic muons is determined by their momenta and the angle γ between them [142] as:

$$m^2 = 2p^+p^-(1 - \cos \gamma), \quad (3.25)$$

where p^+ and p^- denote the magnitudes of the momenta of the positively and negatively charged muons, respectively.

Using Equation 3.25, the invariant mass of a reconstructed dimuon system is computed from the muon momenta and their relative angle. This provides the foundation for the mass-based sagitta bias measurement. By introducing sagitta distortions into the reconstructed transverse momentum of each muon, the corrected invariant mass becomes:

$$m_{\mu\mu 0}^2 = 2|p_0^+||p_0^-|(1 - \cos \gamma) = 2|p^+||p^-|(1 - \cos \gamma)(1 + p_T^+\delta_s^+)^{-1}(1 - p_T^-\delta_s^-)^{-1}, \quad (3.26)$$

where $\delta_s^\pm \equiv \delta_s^\pm(\eta^\pm, \phi^\pm)$ represent the sagitta bias corrections for positively and negatively charged muons. The measured invariant mass is given by $m_{\mu\mu}^2 = 2|p^+||p^-|(1 - \cos \gamma)$, allowing simplification:

$$m_{\mu\mu 0}^2 = m_{\mu\mu}^2(1 - p_T^+\delta_s^+ + p_T^-\delta_s^-). \quad (3.27)$$

Rearranging yields the following relation:

$$\frac{m_{\mu\mu}^2 - m_{\mu\mu 0}^2}{m_{\mu\mu}^2} = (p_T^- \delta_s^- - p_T^+ \delta_s^+). \quad (3.28)$$

This equation involves two unknown bias terms, so a unique solution cannot be extracted directly. To proceed, it is assumed that both muons contribute equally to the observed mass discrepancy; further discussion related to muon momentum behaviour can be found in Appendix A.

$$\delta_s^- = \frac{1}{2} \frac{m_{\mu\mu}^2 - m_{\mu\mu 0}^2}{m_{\mu\mu}^2} \frac{1}{p_T^-}, \quad \delta_s^+ = -\frac{1}{2} \frac{m_{\mu\mu}^2 - m_{\mu\mu 0}^2}{m_{\mu\mu}^2} \frac{1}{p_T^+}. \quad (3.29)$$

In the case of $Z \rightarrow \mu^+ \mu^-$ decays, the true invariant mass is taken to be the known Z -boson mass, m_Z , giving:

$$m_{\mu\mu}^2 - m_Z^2 \approx m_Z^2 (p_T^+ \delta_s(\eta^+, \phi^+) - p_T^- \delta_s(\eta^-, \phi^-)). \quad (3.30)$$

The sagitta bias is estimated for each dimuon event and iteratively refined using local averages over detector regions in (η, ϕ) , where distortions are spatially correlated. The updated value at iteration i is computed from:

$$\delta_{s,i}(\eta, \phi) = -q \frac{m_{\mu\mu}^2 - m_Z^2}{2m_Z^2} \frac{(1 + qp_T \langle \delta_{s,i-1}(\eta, \phi) \rangle)}{p_T} + \langle \delta_{s,i-1}(\eta, \phi) \rangle. \quad (3.31)$$

Here, q is the muon charge, $\delta_{s,i}(\eta, \phi)$ is the sagitta correction at iteration i , and $\langle \delta_{s,i-1}(\eta, \phi) \rangle$ is the average correction from the previous iteration. The procedure converges after a few cycles when changes between steps become negligible.

Figure 3.13 shows the average sagitta bias obtained in simulation. As expected from MC samples without misalignment, the result is close to zero.

Although the mass method is effective in extracting spatial sagitta patterns, it is inherently sensitive only to relative differences between oppositely charged muons. As Equation 3.28 shows, the right-hand side depends on the difference in contributions from positive and negative transverse momenta. When both muon p_T spectra are similar, global biases tend to cancel and remain undetected.

This limitation is illustrated in Figure 3.14, where a uniform artificial bias of 0.4 TeV^{-1} is injected across all (η, ϕ) bins. The method fails to recover the injected value. In contrast, when a shape-dependent bias is introduced, the method captures its spatial

variation but not the absolute scale, as shown in Figure 3.15. A modification of the selection to use only the leading muon in each event improves sensitivity to the global shift, as demonstrated in Figure 3.16.

Due to this insensitivity to absolute scale, an alternative strategy — the VarMin method — is developed within the context of this thesis. This method is designed to improve sensitivity to global sagitta biases that are not well constrained by the standard mass method.

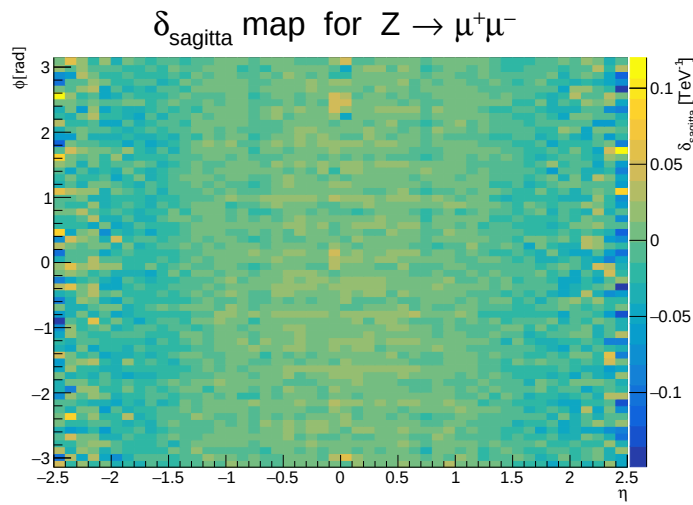


Figure 3.13: Delta sagitta map that represents sagitta corrections measured for MC simulated events versus (η, ϕ) bins. The figure shows result of one iteration using the mass method.

3.4 VarMin method for improving the measurement of the sagitta bias

Seeking improvements over the traditional mass method, a novel approach known as the **VarMin** method is proposed [143, 144]. This innovative technique is predicated on the minimisation of the squared invariant mass of the dimuon system, resulting in enhancements in both computational efficiency and detection sensitivity. While it operates on the same physical principles as the previous method, its distinction lies in the discretisation of the output aligned with the geometric configuration of the detector (η, ϕ) , allowing it to detect global scale biases that were previously elusive.

At the core of the VarMin method is the optimisation of the variance of the corrected dimuon mass, isolating the sagitta biases that minimise this mass, as delineated in

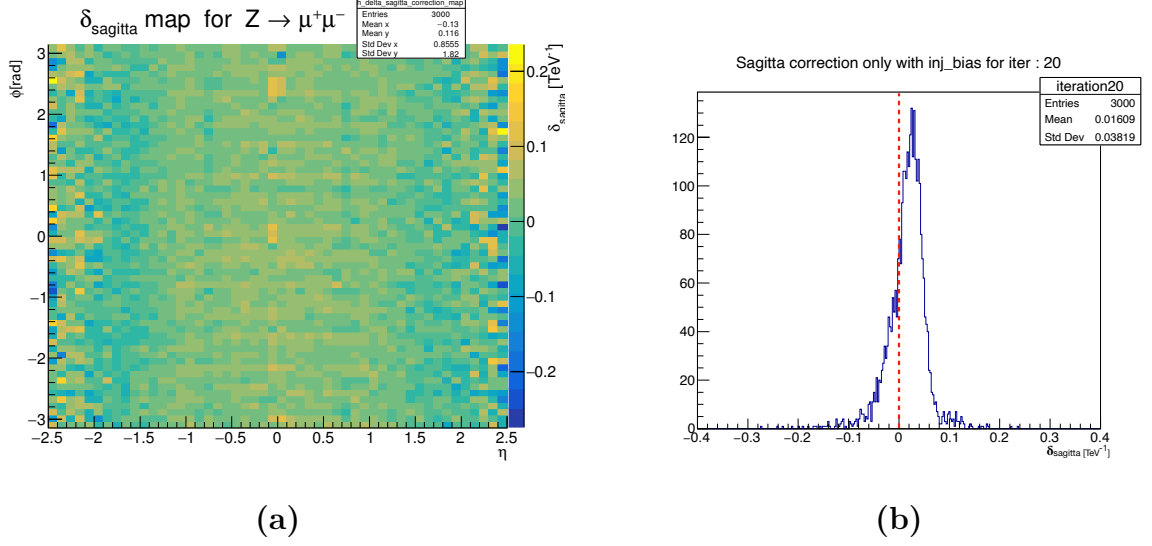


Figure 3.14: Test results of the mass method under a uniform global bias. Figure (a) shows the sagitta bias map after injecting an artificial bias of constant strength 0.4 TeV^{-1} uniformly across all (η, ϕ) regions. Figure (b) presents the projection of the two-dimensional map in Figure (a) onto the ϕ -axis, obtained by averaging over all η bins.

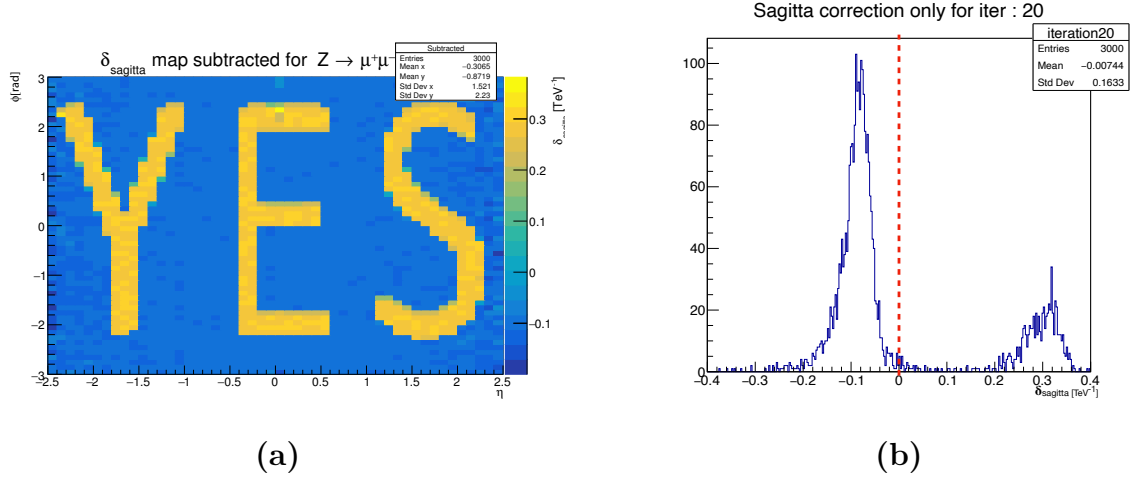


Figure 3.15: Figure (a) shows the sagitta bias as a function of (η, ϕ) , where an artificial bias with a "Yes"-shaped pattern and a magnitude of 0.4 TeV^{-1} has been injected. The result corresponds to 20 iterations of the mass method. In Figure (b), the smaller peak near 3.5 corresponds to bins where the artificial bias was applied, while the larger peak arises from bins without injected bias. Notably, this larger peak does not centre at zero, indicating that the mass method also introduces a global shift in the background, even in the absence of true bias in those regions.

Equation 3.33. The method relies on a set of correction parameters structured into a vector form:

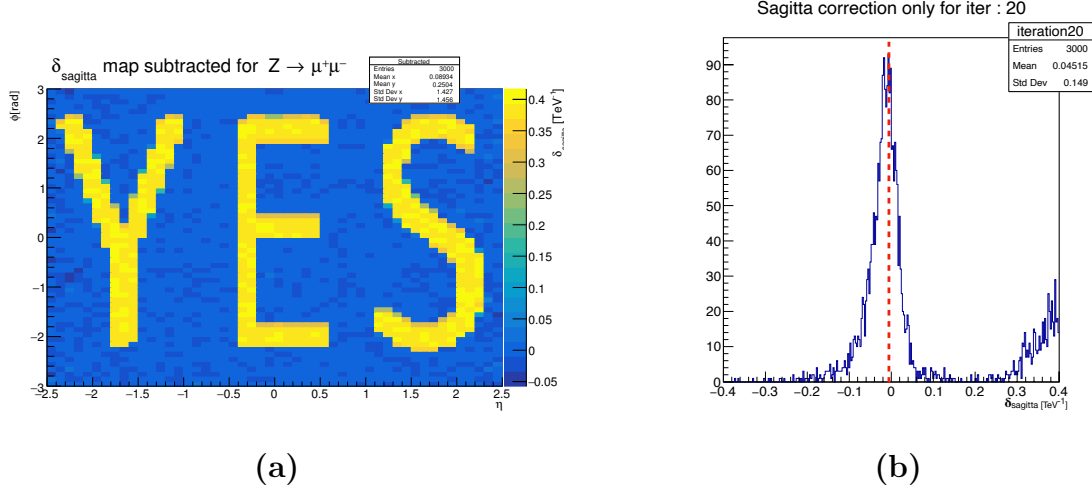


Figure 3.16: Test results of the mass method with single-muon events. Figure (a) sagitta bias map for the mass method using leading muon only. Injected artificial bias has a "Yes" shape and is of strength 0.4 TeV^{-1} . The map is for 20 iterations, subtracted from a pure MC delta sagitta map. Figure (b) represents the Z projection of the sagitta bias map. The larger peak corresponds to the background which is distributed around zero. The smaller peak corresponds to bins in "Yes" shape.

$$\vec{\delta} = \begin{pmatrix} \delta_1 \\ \vdots \\ \delta_D \end{pmatrix}. \quad (3.32)$$

Here, D denotes the total number of distinct (η, ϕ) bins, and $\vec{\delta}$ represents the sagitta bias corrections defined over these detector regions. This formulation enables rewriting Equation 3.28 in a compact matrix form, suitable for variance minimisation.

$$(m_{\mu\mu,i}^{\text{corr}})^2 = \Delta m_{\mu\mu,i}^2 - \vec{e}_i \cdot \vec{\delta}. \quad (3.33)$$

Here, $(m_{\mu\mu,i}^{\text{corr}})^2$ denotes the fully corrected dimuon invariant mass squared for event i , and $\Delta m_{\mu\mu,i}^2 \equiv m_{\mu\mu,i}^2 - m_Z^2$ is the difference between the reconstructed dimuon mass squared and the nominal Z -boson mass squared, with m_Z taken from the Particle Data Group [61]. The vector $\vec{e}_i$ is a sparse D -dimensional vector that encodes the contributions of the muons in event i to the sagitta bias correction. The number of dimensions D corresponds to the total number of (η, ϕ) detector bins. Each vector $\vec{e}_i$ contains one or two non-zero components, depending on whether both muons fall into the same or different detector regions. The structure of the e_i is written as:

$$\vec{e}_i \equiv \begin{pmatrix} 0 \\ \cdot \\ \cdot \\ \cdot \\ \cdot \\ m_{\mu\mu i}^2 p_i^+ \\ 0 \\ \cdot \\ \cdot \\ \cdot \end{pmatrix} - \begin{pmatrix} 0 \\ \cdot \\ \cdot \\ \cdot \\ \cdot \\ 0 \\ m_{\mu\mu i}^2 p_i^- \\ \cdot \\ \cdot \\ \cdot \end{pmatrix}. \quad (3.34)$$

The method introduces the variance of the corrected dimuon mass as follows:

$$\begin{aligned} \text{Var} \left[\left(m_{\mu\mu, i}^{\text{corr}} \right)^2 \right] &= \text{Var} \left[\Delta m_{\mu\mu, i}^2 - \vec{e}_i \cdot \vec{\delta} \right] \\ &= \text{Var} \left[\Delta m_{\mu\mu, i}^2 \right] + \text{Var} \left[\vec{e}_i \cdot \vec{\delta} \right] - 2 \text{Cov} \left[\Delta m_{\mu\mu, i}^2, \vec{e}_i \cdot \vec{\delta} \right] \\ &= \text{Var} \left[\Delta m_{\mu\mu, i}^2 \right] + \text{Cov} \left[\vec{e}_i \cdot \vec{\delta}, \vec{e}_i \cdot \vec{\delta} \right] - 2 \text{Cov} \left[\Delta m_{\mu\mu, i}^2, \vec{e}_i \right] \cdot \vec{\delta} \\ &= \text{Var} \left[\Delta m_{\mu\mu, i}^2 \right] + \sum_{s, t} \text{Cov} \left[(\vec{e}_i)_s, (\vec{e}_i)_t \right] \delta_s \delta_t - 2 \sum_s \text{Cov} \left[\Delta m_{\mu\mu, i}^2, (\vec{e}_i)_s \right] \delta_s. \end{aligned} \quad (3.35)$$

The variance of the corrected dimuon mass is then expressed as a function of the sagitta bias corrections δ_s , where s and t index the (η, ϕ) detector bins. The optimal values of these corrections are obtained by minimising this variance, which is achieved by setting the partial derivative of the variance with respect to each individual correction parameter δ_k to zero.

$$\begin{aligned} \frac{\partial}{\partial \delta_k} \text{Var} \left[\left(m_{\mu\mu, i}^{\text{corr}} \right)^2 \right] &= 0, \\ \sum_t \text{Cov} \left[(\vec{e}_i)_k, (\vec{e}_i)_t \right] \delta_t + \sum_s \text{Cov} \left[(\vec{e}_i)_s, (\vec{e}_i)_k \right] \delta_s - 2 \text{Cov} \left[\Delta m_{\mu\mu, i}^2, (\vec{e}_i)_k \right] &= 0. \end{aligned} \quad (3.36)$$

In order for Equation 3.36 to stand, the following equation has to be satisfied

$$\sum_t \text{Cov} \left[(\vec{e}_i)_k, (\vec{e}_i)_t \right] \delta_t = \text{Cov} \left[\Delta m_{\mu\mu, i}^2, (\vec{e}_i)_k \right]. \quad (3.37)$$

This minimisation condition leads to a system of linear equations, where the matrix M and the vector $\vec{k}$ are constructed from event-by-event contributions. Solving the linear system

$$M\vec{\delta} = \vec{k}, \quad (3.38)$$

yields the set of sagitta bias corrections $\vec{\delta}$ that minimise the variance of the corrected dimuon mass. Each component of $\vec{\delta}$ corresponds to a correction in a specific (η, ϕ) detector bin, ensuring a spatially resolved treatment of the bias.

Here, M is a matrix of dimension $D \times D$ and $\vec{k}$ vector of dimension D , with components

$$\begin{aligned} k_t &\equiv \text{Cov} [\Delta m_{\mu\mu,i}^2, (\vec{e}_i)_t] \equiv \text{Cov} [m_{\mu\mu,i}^2 - m_Z^2, (\vec{e}_i)_t] \\ &= \left(\frac{1}{N} \sum_i (m_{\mu\mu,i}^2 - m_Z^2)(\vec{e}_i)_t \right) - \left(\frac{1}{N} \sum_i (m_{\mu\mu,i}^2 - m_Z^2) \right) \left(\frac{1}{N} \sum_i (\vec{e}_i)_t \right). \end{aligned} \quad (3.39)$$

$$\begin{aligned} M_{s,t} &\equiv \text{Cov} [(\vec{e}_i)_s, (\vec{e}_i)_t] \\ &= \left(\frac{1}{N} \sum_i (\vec{e}_i)_s (\vec{e}_i)_t \right) - \left(\frac{1}{N} \sum_i (\vec{e}_i)_s \right) \left(\frac{1}{N} \sum_i (\vec{e}_i)_t \right). \end{aligned} \quad (3.40)$$

The VarMin method begins by populating four preliminary objects within the event loop and then solves a system of linear equations to output the $\vec{\delta}$ vector, representing sagitta corrections for each detector bin. This vector is plotted in a two-dimensional distribution to adjust momentum measurements iteratively, typically requiring three cycles for optimal alignment. The VarMin method, currently employed in place of the mass method, provides an improved approach for measuring sagitta bias.

To validate the new VarMin method, a known bias is injected into the MC simulated samples and the method ability to detect both the shape and magnitude of this bias is assessed. The results, illustrated in Figure 3.17, confirm that the VarMin method accurately identifies the injected bias with high precision, unaffected by background noise, since the MC simulated samples are expected to be free of bias. It is important to note that while MC simulations aim for zero bias, inherent fluctuations often introduce a small "noise bias", which is accounted for and removed from the test cases to ensure clarity in the results. This subtraction procedure is discussed in more detail in Appendix B, where results for one, two, three, and ten-iterations are presented.

The mass method required up to ten iterations to converge, whereas the new VarMin

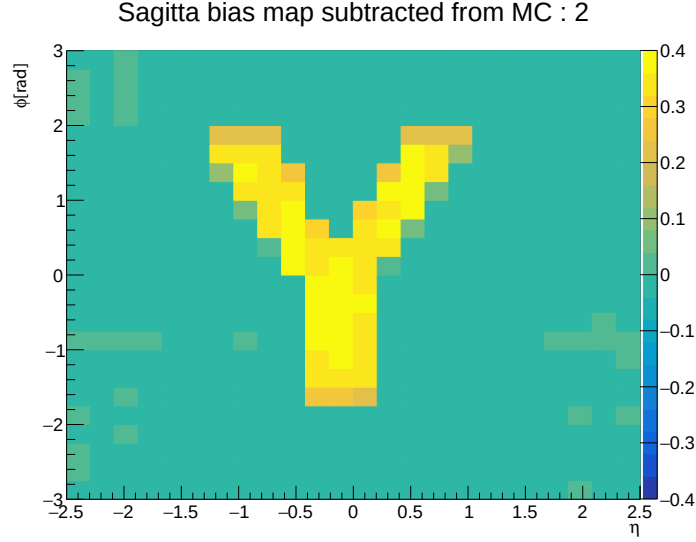


Figure 3.17: Sagitta bias map for pure MC simulation subtracted from the δ_s with “Y” shaped bias injected. In one iteration the “Y” shape converges to 0.4 while the background converges to zero.

method achieves convergence much more rapidly. The results displayed in Figure 3.17 are based on just a single iteration, highlighting the efficiency of the method. For a direct comparison of convergence rates, Figure 3.18 presents side-by-side convergence plots for both the mass method and the VarMin method, clearly demonstrating that the VarMin method reaches convergence significantly quicker.

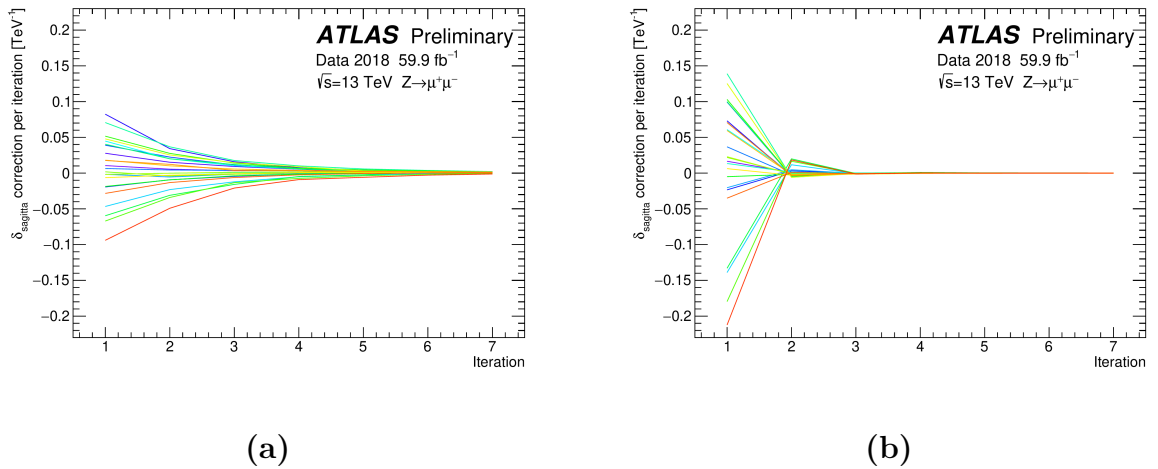


Figure 3.18: Convergence of δ_s corrections per iteration for the mass (a) and (b) VarMin methods [142].

The mass method calculated statistical uncertainty by assessing the deviation in the mean of the delta sagitta distributions. However, the VarMin method, by assigning

a singular delta sagitta correction for each (η, ϕ) bin, does not inherently provide an estimate of the random uncertainty. To address this, the bootstrapping method is employed. Bootstrapping is employed to generate multiple pseudo-experiments by resampling the available data with replacement. Each pseudo-experiment mimics an independent measurement, allowing the statistical fluctuations of the sagitta bias values to be studied. The distribution of results across these pseudo-experiments provides an estimate of the statistical uncertainty associated with the measured sagitta bias.

Figure 3.19 illustrates the propagation of uncertainties across the detector geometry. This methodological approach involves creating n pseudo-experiment sets, with the variance in results across these sets offering a measure of statistical uncertainty.

In the VarMin method, a Poisson distribution is applied to weight the events, generating varied pseudo-datasets as defined by:

$$P(n) = \frac{e^{-\mu} \mu^n}{n!}, \quad (3.41)$$

where $P(n)$ denotes the probability of observing n events when the mean expected number of occurrences is μ .

To estimate the statistical uncertainty, the standard deviation (σ) of the bootstrap event samples is used. A set of 100 bootstrap event samples yields an uncertainty range of approximately ± 0.005 .

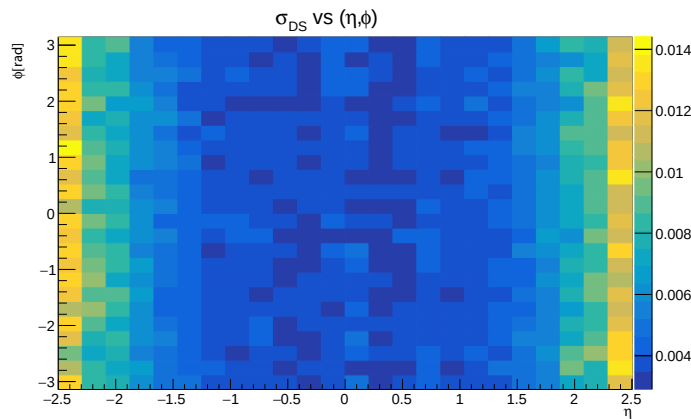


Figure 3.19: Statistical uncertainty propagation for the VarMin method across detector geometry.

Figure 3.20 presents the statistical uncertainty estimation for the mass method, comparing the two figures indicates that the statistical uncertainty from both the mass and VarMin methods are comparable. This suggests that the VarMin approach is as robust as the mass method with respect to statistical uncertainty.

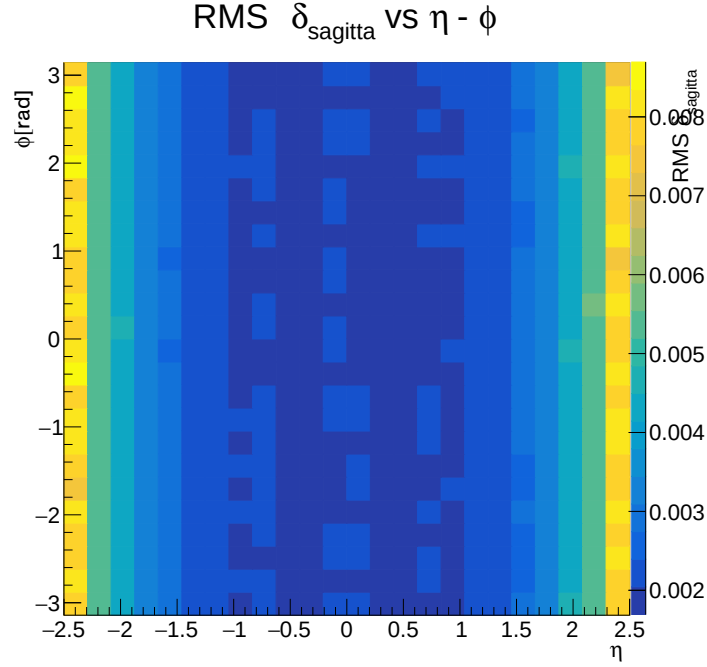


Figure 3.20: Statistical uncertainty map for mass method.

The VarMin method establishes a robust and unbiased framework for the measurement of sagitta biases in the ATLAS ID. By minimising the variance of the corrected dimuon mass, it achieves sensitivity to both localised and global distortions that could otherwise compromise momentum measurements. Its validation on MC simulated samples and subsequent application to real ATLAS collision data confirm its stability and consistency over time, as illustrated in Figure 3.21.

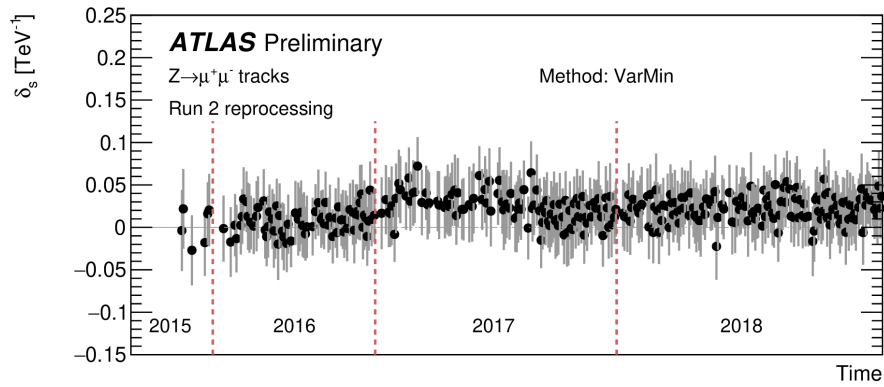


Figure 3.21: Annual measurements of sagitta bias in data collected by the ATLAS detector, demonstrating the efficiency of the alignment corrections over time [142].

In the broader context of precision measurements at the LHC, the development and deployment of the VarMin method reflect the continuous refinement of experimental techniques required to meet the increasing demands of modern particle physics. Accurate alignment and momentum calibration are not merely technical necessities, but fundamental pillars supporting the search for new physics phenomena.

Chapter 4

Recording data and simulating events with the ATLAS detector

In particle physics, pp collisions are studied to deepen the understanding of fundamental particles and their interactions. At the LHC, these collisions are provided at unprecedented energies, and the ATLAS detector records the resulting events.

To enable a statistical interpretation of the collision data, theoretical predictions are obtained using MC event generators. The SM serves as the baseline theoretical framework for describing particle interactions.

For real collision data, the detector response is evaluated using reconstruction algorithms that identify and reconstruct the final-state objects. For simulated events, MC-generated particles undergo a full simulation chain that replicates the physics of the collisions, the interaction with the detector material, the effects of magnetic fields, and the electronic response of the detector. Subsequently, the simulated events are processed with the same reconstruction algorithms as those applied to the real data, ensuring a consistent comparison between data and simulation.

This chapter introduces the foundational aspects of simulating events in the ATLAS detector. It begins with a discussion of the phenomenology of pp collisions, including the parton model, hard scattering, PS and hadronisation, and the underlying event presented in Section 4.1. The data acquisition process by ATLAS is briefly reviewed in Section 4.2, followed by Section 4.3 with a detailed explanation of the simulation chain, from event generation to detector simulation. The MC tools used in the analysis are introduced and contextualised in terms of the analysis requirements in Section 4.3.2. Then, Section 4.4 outlines the data and simulated samples used in the quadruple-differential analysis, including those generated with EFT parameter variations explained in Section 4.4.1.

Finally, it presents the morphing procedure used to interpolate event predictions across the EFT parameter space based on the M-function decomposition shown in Section 4.5.

4.1 Phenomenology of pp Collisions

The analysis presented in this thesis is based on data collected from pp collisions recorded by the ATLAS detector. Understanding the phenomenology of such collisions requires recognising that protons are composite objects, consisting of quarks and gluons rather than elementary particles. These constituents are collectively referred to as partons, and include the valence quarks—two up quarks and one down quark—as well as a dynamic sea of quark–antiquark pairs generated via gluon splitting. Gluons are responsible for binding quarks together and can themselves split into additional quark–antiquark pairs. The description of the proton in terms of these point-like partons forms the basis of the parton model [145].

At the energies achieved by the LHC, the interactions between protons are best understood as interactions between their constituent partons. Depending on the energy scale, collisions may be elastic or inelastic. Elastic collisions dominate at lower energies, while deep inelastic scattering becomes prominent at higher energies, enabling detailed studies of proton structure.

The subsections that follow provide an overview of the essential components of pp collision phenomenology, including the parton model, hard scattering processes, PS evolution, hadronisation, and the underlying event.

4.1.1 The Parton model

The parton model, introduced by Richard Feynman in 1969 [145], offers a simplified yet powerful framework to describe the internal structure of hadrons in terms of partons. Each parton carries a fraction of the total momentum of the hadron, and during the high-energy interaction, partons behave as quasi-free particles.

The momentum distribution of partons inside the proton is characterised by parton distribution functions (PDFs). These functions depend on the fraction of the proton momentum carried by the parton, denoted x , and the energy scale Q^2 of the interaction. If protons consisted solely of three non-interacting quarks, the PDFs would reduce to sharp peaks at $x = 1/3$. However, due to the complexity of QCD, continuous gluon emissions and sea quark contributions result in more diffuse distributions. PDFs are extracted from experimental data through global fits and are evolved using perturbative QCD techniques.

Well-known PDF sets are provided by the Martin–Stirling–Thorne–Watt (MSTW) [146], Neural Network PDF (NNPDF) [147], and Coordinated Theoretical-Experimental Project on QCD (CTEQ) [148] collaborations. An example of a PDF set at different energy scales is shown in Figure 4.1.

4.1.2 Hard Scattering

Hard scattering refers to pp interactions with large momentum transfer, resulting in the production of high-energy final-state particles. In the theoretical description, the factorisation theorem¹ allows the separation of the hadronic cross-section into a convolution of partonic cross-sections and PDFs. The cross-section is expressed:

$$\sigma(p_1, p_2) = \sum_{i,j} \int dx_1 dx_2 f_i(x_1, \mu_F^2) f_j(x_2, \mu_F^2) \hat{\sigma}_{ij}, \quad (4.1)$$

where x_1 and x_2 are the fractions of momenta of the considered parton within p_1 and p_2 , respectively; $f_i(x_1, \mu_F^2)$ and $f_j(x_2, \mu_F^2)$ represent the associated PDFs; μ_F is the factorisation scale determining the scale at which perturbative calculations are valid; and $\hat{\sigma}_{ij}$ is the partonic cross-section between the two considered partons. The indices i and j run over all partons in p_1 and p_2 , respectively.

The partonic cross-section $\hat{\sigma}_{ij}$ can be computed within perturbative QCD as an expansion in the strong coupling constant α_s as:

$$\hat{\sigma}_{ij} = \hat{\sigma}_{ij,\text{LO}} + \alpha_s(\mu_R) \hat{\sigma}_{ij,\text{NLO}} + \alpha_s^2(\mu_R) \hat{\sigma}_{ij,\text{NNLO}} + \dots, \quad (4.2)$$

where the renormalisation scale μ_R cures the ultraviolet divergences in QCD due to large momentum in the loops. Since calculating the entire series is unfeasible, it has to be truncated at a particular order, fixing μ_R to a particular value. This and the μ_F choice introduce uncertainties which are often estimated by varying these parameters.

Simulations of hard-scattering processes are performed using matrix-element generators such as MADGRAPH [149] and POWHEG [150–152], which provide accurate modelling of the primary partonic interactions at various orders in perturbation theory.

¹The factorisation theorem states that, at sufficiently high energies, the cross-section of a hadronic process can be expressed as the convolution of universal, process-independent PDFs and a process-dependent partonic cross-section calculable using perturbative QCD.

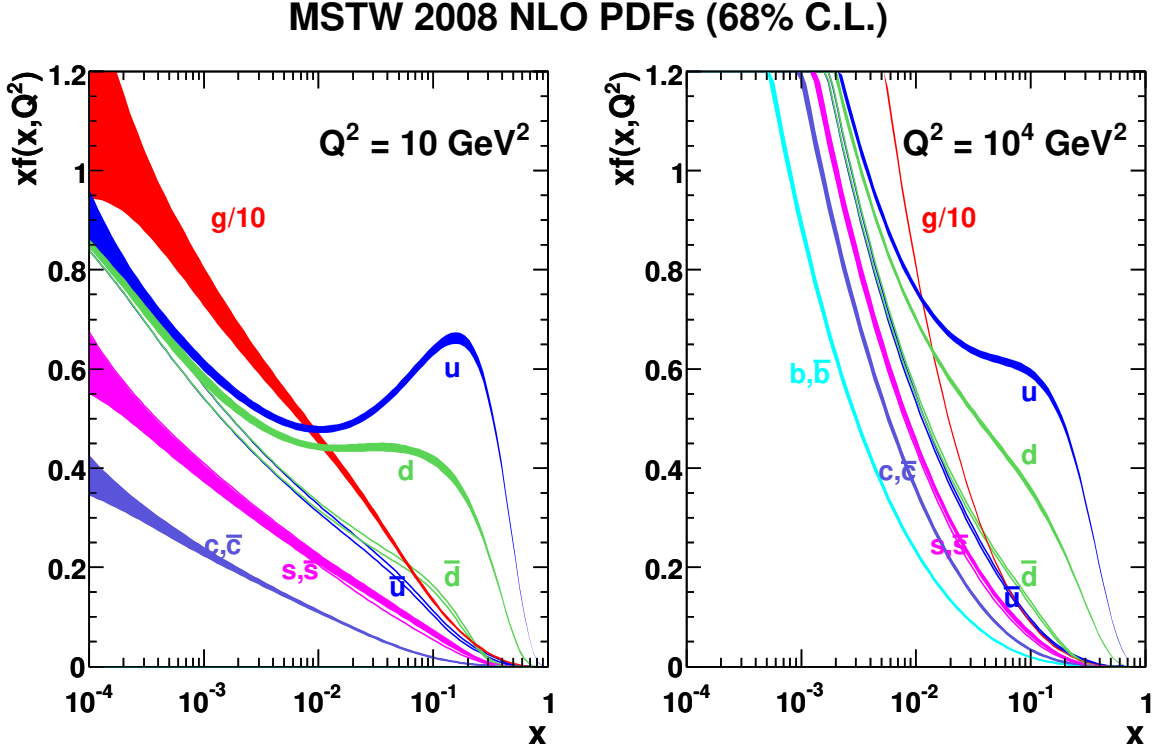


Figure 4.1: The PDFs are evaluated using the MSTW 2008 NLO set [146] at two different energy scales: $Q^2 = 10 \text{ GeV}^2$ (left) and $Q^2 = 10^4 \text{ GeV}^2$ (right) [146].

4.1.3 Parton Shower and Hadronisation

Following the hard interaction, outgoing partons undergo additional radiation through the emission of gluons and quark–antiquark pairs, a phenomenon known as the PS. This process redistributes the energy of the initial partons into a larger number of lower-energy partons.

As the PS evolves and the typical energy scale drops, perturbative calculations become unreliable. At this stage, partons hadronise into colour-neutral hadrons. MC event generators such as PYTHIA [153], HERWIG [154], and SHERPA [155] simulate the PS, hadronisation, and subsequent decays of unstable particles. Hadronisation models, such as the Lund string model [156] used in PYTHIA and the cluster model [157] used in HERWIG and SHERPA, provide phenomenological descriptions of this non-perturbative process. Further decay chains are simulated using dedicated tools such as EVTGEN [158].

4.1.4 Underlying Event

In addition to the products of the primary hard scattering, a typical pp collision includes other sources of hadronic activity collectively referred to as the underlying event. Contributions arise from multiple parton interactions (MPI), beam–beam remnants, and initial- and final-state radiation (ISR/FSR) [159].

Accurate modelling of the underlying event is necessary for precise experimental measurements and for comparisons with theoretical predictions. The complexity of the underlying event, dominated by soft QCD processes, requires non-perturbative modelling within event generators. PYTHIA, HERWIG, and SHERPA each provide different approaches and parameter tunings for the simulation of the underlying event. Figure 4.2 illustrates the components of a typical underlying event in a pp collision, highlighting the interplay between hard-scattering products and additional soft activity.

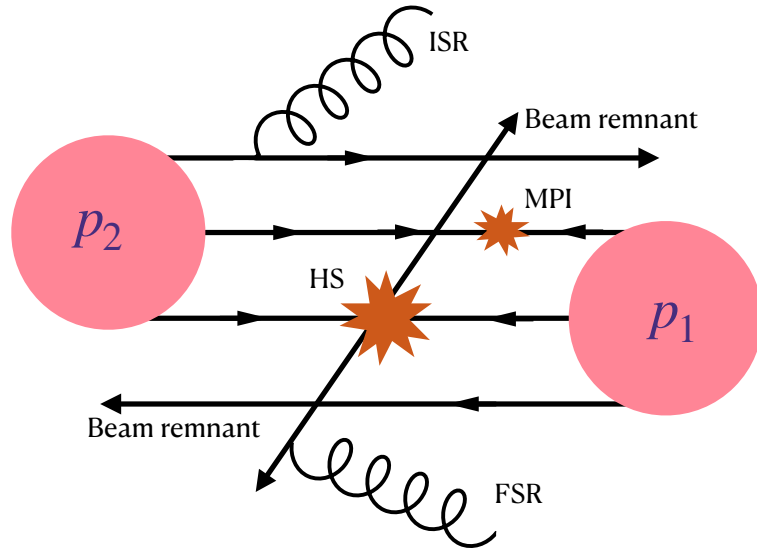


Figure 4.2: Schematic representation of a generic pp collision, each p_1 and p_2 provide a parton participating in a hard scattering (HS), resulting in two outgoing final-state partons. The underlying event consists of ISR, FSR, beam–beam remnants, and additional MPI.

4.2 Data Collection

The collection of data with the ATLAS detector begins with pp collisions delivered by the LHC, producing a wide variety of final-state particles in the detector. These interactions occur at a very high rate, and only a small fraction can be retained for further analysis.

Once selected for storage, the data must be transformed from raw detector signals into meaningful physics observables. This process follows a well-structured model that ensures consistency, reproducibility, and compatibility across all stages of the analysis chain.

The ATLAS data model [113] provides a layered framework for recording, reconstructing, and interpreting events. It defines how data evolve from the detector readout format to high-level physics objects. A key feature of this model is the parallel treatment of real and simulated data: both pass through analogous stages of reconstruction and formatting, which enables direct comparison between measured events and theoretical predictions. Figure 4.3 illustrates the alignment of these processing paths, ensuring that systematic uncertainties can be consistently evaluated.

The final dataset used in this thesis corresponds to events collected during LHC Run 2. Only a subset of all recorded events is considered suitable for physics analysis. This selection is based on detailed monitoring of detector conditions and reconstruction performance. ATLAS employs dedicated data quality algorithms to assess whether each component of the detector functioned as expected during data-taking. Periods that meet the strictest quality criteria are compiled into Good Run Lists (GRLs) [160], which represent the foundation of all published analyses.

For the Run 2 data, approximately 95.6% of the delivered luminosity satisfies the “All Good” selection criteria, which require nominal performance from all major sub-detectors and the correct reconstruction of key physics objects such as electrons, muons, jets, and missing transverse momentum. The dataset used in this analysis includes only events from these GRLs, ensuring that the physics results are based on data with fully validated integrity and detector performance.

4.3 Simulating events

MC simulations are essential for modelling both signal and background processes in ATLAS analyses. Events are generated according to SM cross-sections using various MC tools that rely on random sampling, requiring large statistics to capture the probabilistic nature of particle collisions. The resulting collections, known as MC event samples, replicate the experimental conditions with high fidelity.

The ATLAS simulation chain consists of three key stages [161]:

1. **Event generation:** A matrix-element generator simulates the hard-scattering process and immediate decays. This includes PS, hadronisation, and pile-up. At this stage, the detector geometry is not yet considered.

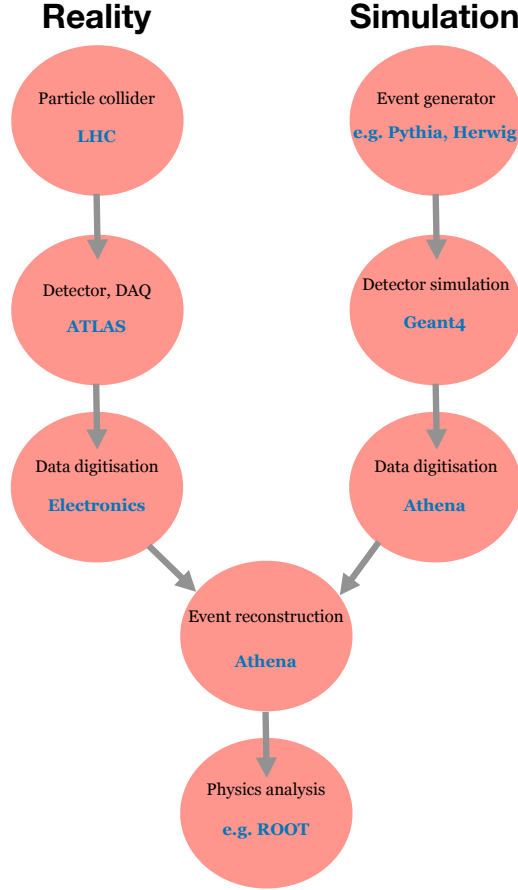


Figure 4.3: Comparison of the paths followed by data recorded by the ATLAS detector and the simulated samples. At each step, the format of the simulated data is the same as the recorded data.

2. **Detector simulation:** The resulting particles are propagated through a detector model using GEANT4 [161], simulating all relevant interactions with detector materials such as energy loss, scattering, and secondary decays.
3. **Digitisation:** The simulated energy deposits are converted into signals mimicking those produced in the actual detector electronics.

The final output reproduces the format of real events recorded by the ATLAS DAQ system (see Section 3.2.6), enabling identical reconstruction workflows. Figure 4.3 illustrates the alignment between real and simulated data.

4.3.1 Steps for simulating event data

The simulation of event data is essential for accurately interpreting measurements performed with the ATLAS detector. The procedure accounts not only for the primary hard-scattering process but also for additional interactions such as multiple pp collisions per bunch crossing (pile-up) (see Section 3.1.2) and residual activity from adjacent crossings. A detailed depiction of the simulation workflow is shown in Figures 4.3 and 4.4.

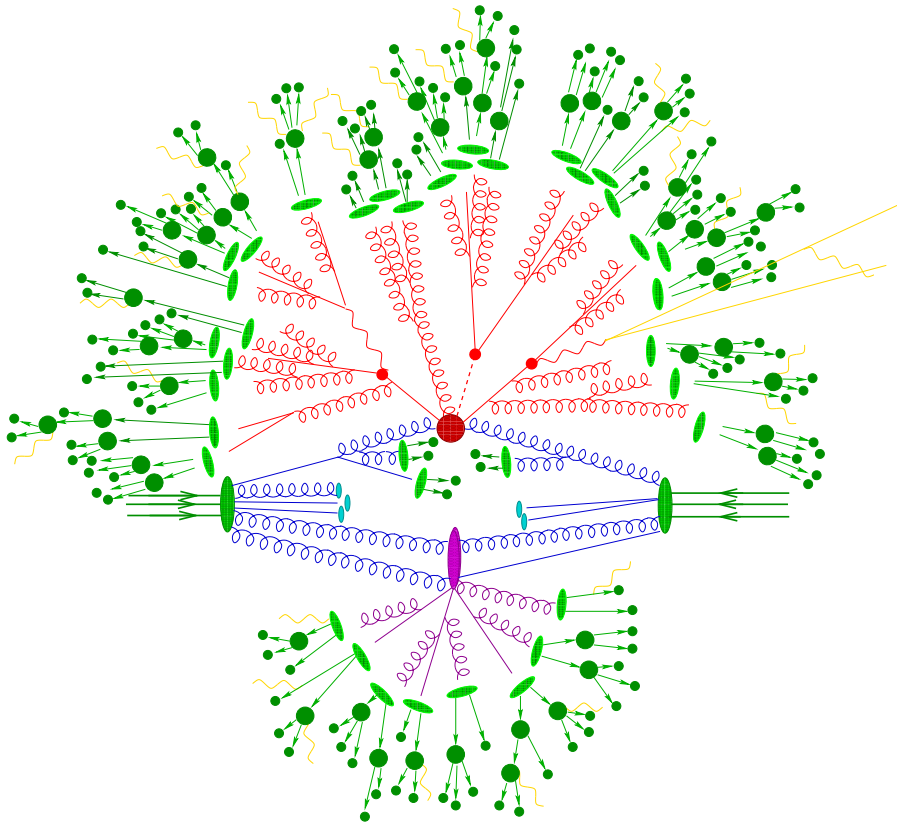


Figure 4.4: Event structure of a typical LHC collision. The central red region denotes the hard-scattering interaction, producing multiple outgoing partons (small red nodes), followed by QCD radiation generated through PS. A secondary partonic interaction, shown in purple, occurs before the onset of hadron formation, which is represented in light green. Subsequent decays of unstable hadrons are depicted in dark green. Electromagnetic radiation, shown in yellow, may occur at various stages throughout the event evolution [162].

To ensure comprehensive coverage of the phase space, simulated samples are produced independently of the natural cross-sections of the processes involved. Rather than generating events proportional to their expected occurrence rates, a sufficiently large number of events for each process are generated, followed by appropriate cross-section re-weighting. This approach provides a robust description of both dominant and rare

processes but introduces events with negative weights, a consequence of higher-order corrections in MC simulations.

4.3.1.1 Hard-Scattering Process

The simulation chain begins with the modelling of the hard-scattering event. MEs are computed at a fixed order in perturbative QCD, typically at LO or NLO, depending on the required accuracy. The complexity of the calculation increases with the order considered, impacting computational cost. From the MEs, cross-sections are derived and events are generated at parton level, without reference yet to the detector geometry. Following the hard interaction, PSs and hadronisation models are applied to incorporate radiative corrections and non-perturbative effects.

4.3.1.2 Parton Shower and Hadronisation

After the hard-scattering stage, PS are simulated to account for soft and collinear QCD and QED radiation. All colour-charged and electrically charged partons can emit additional partons or photons, respectively, creating complex final states with wide energy and momentum spectra. PS algorithms typically employ different ordering schemes, such as transverse momentum (p_T) ordering in PYTHIA [153], or angular and dipole ordering in HERWIG [154].

The PS evolution is terminated at a low-energy cut-off, around the hadronisation scale (~ 1 GeV), where perturbative methods lose validity. Hadronisation models, which translate coloured partons into colour-neutral hadrons, are based on QCD-inspired phenomenological descriptions. Two primary models are employed: the Lund string model [156], implemented in PYTHIA, and the cluster model [157], used by HERWIG and optionally by SHERPA.

4.3.1.3 Underlying Event

Beyond the radiation associated with the hard scattering, additional hadronic activity referred to as the underlying event (UE) must be modelled. UE encompasses contributions from MPI, remnants of the colliding protons, and soft initial- and final-state radiation. These effects are simulated with event generators such as PYTHIA, HERWIG, and SHERPA, with their parameters tuned to experimental data [159].

4.3.1.4 Hadron decay

Unstable hadrons produced during hadronisation undergo decays before being detected. Particles with lifetimes corresponding to decay lengths shorter than 10 mm are considered unstable. Modelling hadronic decays is essential, as a significant fraction of final-state particles originate from such decays. Decay patterns and branching ratios are implemented based on experimental data.

4.3.1.5 Pile-Up modelling

Pile-up interactions, arising from multiple pp collisions within the same or neighbouring bunch crossings, are simulated by overlaying additional minimum-bias events onto the hard-scattering event. In the ATLAS experiment, PYTHIA8 [153] is used to generate these minimum-bias events. The distribution of pile-up events is re-weighted to match the pile-up conditions observed during data-taking.

4.3.1.6 Simulation of the ATLAS Detector and Digitisation

Following the generation of the complete particle content, the interaction of these particles with the ATLAS detector is simulated using the GEANT4 toolkit [163]. This step models the passage of particles through the detailed geometry of the detector, accounting for processes such as ionisation, scattering, energy loss, photon conversions, and interactions with both sensitive and non-sensitive materials. Magnetic field effects are also included.

The digitisation step simulates the electronic signals produced by particle interactions within the readout systems of the detector. Simulated hits are transformed into electronic signals (digits), including the effects of detector noise. Digitisation is performed within the Athena framework and ensures that the output format mirrors that of real data collected by the ATLAS DAQ system. Trigger simulation is incorporated before digitisation to reproduce the conditions under which events are selected.

Events that have passed through the full simulation and reconstruction chain are referred to as reconstruction-level events. These events are used in the subsequent physics analyses, including tasks such as lepton origin identification and background estimation.

4.3.1.7 Full and Fast Detector Simulations

Given the computational demands of full GEANT4-based simulation, ATLAS also employs a faster alternative for some analyses. Full Simulation (FS) provides a highly detailed

response at the cost of significant CPU resources, particularly in the calorimeters. In contrast, the Atlfast-II (AFII) fast simulation [164, 165] uses parameterised responses in the calorimeters while preserving detailed GEANT4 simulation in the inner detector and muon spectrometer.

The AFII approach enables faster generation of large MC samples while maintaining sufficient accuracy for many analyses. Both FS and AFII samples are used in the present study depending on the precision requirements of specific processes.

4.3.2 Monte Carlo Generators

The event simulation chain in high-energy physics relies on a range of MC generators that implement the various stages of event modelling, including ME calculations, PS, hadronisation, UE simulation, and hadronic decays. The most relevant generators used in this analysis are summarised below.

- **MadGraph5_aMC@NLO** [149]: Short for “Matrix element Automatic Generator”, a flexible NLO generator used to compute MEs for a wide range of scattering processes. It supports multiple initial- and final-state parton multiplicities and both the four- and five-flavour schemes (4FS and 5FS). The generator automatically generates Feynman diagrams, performs phase-space integration, and produces weighted events which can be interfaced with external PS generators such as PYTHIA or HERWIG. In this analysis, MadGraph5_aMC@NLO is used primarily for the simulation of complex final states in backgrounds such as tZq , tWZ , tHq , tHW , and four-top production.
- **POWHEG-BOX** [150–152]: Short for “Positive Weight Hardest Emission Generator”, a next-to-leading order generator designed to consistently match matrix-element calculations with PS while ensuring positive event weights. It simulates the hardest radiation using NLO precision before interfacing with a parton-shower generator for softer emissions. While the POWHEG formalism avoids negative weights in most processes, some, including single top-quark t -channel production [166], may still introduce them. In this analysis, POWHEG-BOX is used to generate the t -channel signal as well as major backgrounds such as $t\bar{t}$, $t\bar{t}H$, tW , and s -channel single top processes.
- **Herwig 7** [154, 167]: Its name comes from “Hadron Emission Reactions With Interfering Gluons”. A general-purpose generator that provides both matrix-element and parton-shower capabilities at LO and NLO. Its PSs are ordered either angularly or using a dipole formalism, and it employs the cluster model for hadronisation. The

underlying event is simulated using an eikonal multiple-interaction model. In this analysis, Herwig 7 is used primarily to evaluate systematic uncertainties associated with PS and hadronisation modelling through alternative samples.

- **Pythia 8** [153]: A widely used LO event generator that provides PS, hadronisation, and UE modelling. The PS is ordered in transverse momentum, and the Lund string model [156] is used for hadronisation. Pile-up and minimum-bias events are also simulated with this tool. In this analysis, PYTHIA 8 is employed for most PS and pile-up simulations, as well as for UE modelling in signal and background samples.
- **Sherpa** [155, 168]: A multi-purpose event generator capable of simulating full events, from matrix-element generation to PS and hadronisation. Sherpa supports multi-leg matrix-element calculations matched to PS using the ME+PS@NLO scheme, and includes its own hadronisation and UE models. In this analysis, Sherpa is used for background samples involving high final-state multiplicity, such as V + jets, diboson, and triboson processes.
- **MadSpin** [169, 170]: A tool designed to decay heavy resonances in simulated events while preserving spin correlations between decay products. In this analysis, MADSPIN is used to decay the top quark and W boson in the t -channel signal sample, maintaining angular correlations in the final state relevant for differential observables.
- **EvtGen** [158]: A decay package specialised in modelling the decays of heavy-flavour hadrons, particularly B - and D -mesons. It includes detailed decay models for various hadronic and leptonic final states. In this analysis, EVTGEN is used in conjunction with PYTHIA 8 to simulate the decays of b - and c -hadrons in signal and background processes involving heavy flavour.

4.4 Data and simulated event samples in the Quadruple differential analysis

This section presents the datasets and simulated samples used in the quadruple differential analysis, including both real data collected by the ATLAS detector and MC simulations used to model signal and background processes.

4.4.1 Data

Data were collected from pp collisions from 2015 to 2018, with an energy of $\sqrt{s} = 13$ TeV, and were recorded by the ATLAS detector. The selection process involved a unified data stream with un-prescaled single-lepton triggers, as described in Refs. [171–173]. Single-electron events detected in the muon data stream were intentionally excluded to prevent duplication of data entries.

The analysed datasets correspond to an integrated luminosity of 140 fb^{-1} , with a reported uncertainty of 0.83% for the combined data from 2015–2018 [174]. These measurements relied primarily on the LUCID-2 detector [175], with additional verification from the ATLAS IDs and calorimeters. Annual luminosity measurements varied slightly, with uncertainties ranging from 0.89% to 1.13%, showing partial correlations.

Detailed information on both the specific and total integrated luminosities, their uncertainties, and other relevant data are presented in Table 4.1.

Year	Periods	Run numbers	Number of events	Integrated luminosity [pb^{-1}]	Avg. pile-up (μ)
2015	D–H, J	276262–284484	220.58M	$3244.54 \pm 1.13\%$	~ 13.4
2016	A–G, I, K, L	297730–311481	1057.84M	$33402.2 \pm 0.89\%$	~ 25.1
2017	B–F, H, I, K	325713–340453	1340.80M	$44630.6 \pm 1.13\%$	~ 37.8
2018	B–D, F, I, K, L, M, O, Q	348885–364292	1716.77M	$58791.6 \pm 1.10\%$	~ 39.2
2015–2018	All	276262–364292	4335.99M	$140068.94 \pm 0.83\%$	~ 34

Table 4.1: Integrated luminosity, average pile-up, and number of events per year for Run 2 data taking. Run numbers and relative uncertainties are also shown.

In this analysis, pile-up conditions reflect the full Run 2 dataset, with $\langle\mu\rangle$ ranging from approximately 13 to 40. Simulated events are reweighted to match the pile-up distribution observed in data, ensuring consistent modelling of the detector response and reconstructed objects across all data-taking periods.

4.4.2 Simulated event samples

Samples of events used in this analysis were generated using different MC event generators interfaced with various PS and hadronisation algorithms. After generation, the simulation of the detector and trigger response was performed using the **Athena** software framework[176], employing either the GEANT4 toolkit for full simulation [163] or the **Atlfast2** fast-simulation framework [164, 165]. The FS samples are used by default, while **Atlfast2** also known as (AFII) samples are employed for certain processes or to evaluate systematic uncertainties when FS is not available.

Pile-up effects were modelled by overlaying simulated events with inelastic pp collisions generated using PYTHIA 8.186, interfaced with the NNPDF2.3 LO PDF set[177] and the A3 tune[178]. The MC events were re-weighted to reproduce the distribution of the average number of interactions per bunch crossing observed in data, with a global rescaling factor of 1.03 ± 0.04 applied to match the visible inelastic cross-section [179].

Simulated samples for the t -channel single top-quark signal and a variety of background processes were used to model detector efficiencies, resolution effects, and systematic uncertainties. A detailed summary of the simulated samples is provided in Table 4.2.

Process	Generator	Order (scheme)	PDF set	Parton shower	PDF set (tune)
Signal					
t-channel	POWHEG Box v2	NLO (4FS)	NNPDF3.0NLO nf4	PYTHIA8 [8.230]	NNPDF2.3LO (A14 tune)
Background					
$t\bar{t}$	POWHEG Box v2	NLO (5FS)	NNPDF3.0NLO	PYTHIA8 [8.230]	NNPDF2.3LO (A14 tune)
V+jets	SHERPA 2.2.11	NLO+LO	NNPDF3.0NNLO	-	-
Diboson	SHERPA 2.2.1-2	NLO+LO	NNPDF3.0NNLO	-	-
Triboson	SHERPA 2.2.2	NLO+LO	NNPDF3.0NNLO	-	-
$t\bar{t}Z$	MadGraph5_aMC@NLO 2.3.3	NLO	NNPDF3.0NLO	PYTHIA8 [8.230]	NNPDF2.3LO (A14 tune)
$t\bar{t}W$	POWHEG Box v2	NLO	NNPDF3.0NLO	PYTHIA8 [8.230]	NNPDF2.3LO (A14 tune)
$t\bar{t}H$	POWHEG Box v2	NLO (5FS)	NNPDF3.0NLO	PYTHIA8 [8.230]	NNPDF2.3LO (A14 tune)
tW	POWHEG Box v2	NLO (5FS, DR)	NNPDF3.0NLO	PYTHIA8 [8.230]	NNPDF2.3LO (A14 tune)
s-channel	POWHEG Box v2	NLO	NNPDF3.0NLO	PYTHIA8 [8.230]	NNPDF2.3LO (A14 tune)
tZq	MadGraph5_aMC@NLO 2.3.3	NLO	NNPDF3.0NLO	PYTHIA8 [8.230]	NNPDF2.3LO (A14 tune)
tWZ	MadGraph5_aMC@NLO 2.3.3	NLO	NNPDF3.0NLO	PYTHIA8 [8.230]	NNPDF2.3LO (A14 tune)
tHq	MadGraph5_aMC@NLO 2.6.2	NLO (4FS)	NNPDF3.0NLO nf4	PYTHIA8 [8.230]	NNPDF2.3LO (A14 tune)
tWH	MadGraph5_aMC@NLO 2.8.1	NLO (5FS, DR)	NNPDF3.0NLO	PYTHIA8 [8.245p3]	NNPDF2.3LO (A14 tune)
ttt	MadGraph5_aMC@NLO 2.2.2	NLO	NNPDF3.1NLO	PYTHIA8 [8.186]	NNPDF2.3LO (A14 tune)
$tt\bar{t}$	MadGraph5_aMC@NLO 2.3.3	NLO	NNPDF3.1NLO	PYTHIA8 [8.230]	NNPDF2.3LO (A14 tune)
ggH	POWHEG Box v1	NLO (CT10)	CT10	PYTHIA8 [8.210]	CTEQ6L1 (AZNLO tune)
qqH	POWHEG Box v1	NLO	CT10	PYTHIA8 [8.186]	-
WH	PYTHIA8 [8.186]	LO	NNPDF2.3LO	-	-
ZH	PYTHIA8 [8.186]	LO	NNPDF2.3LO	-	-

Table 4.2: Summary of the baseline simulated signal and background event samples.

4.4.2.1 Signal process

The baseline simulated t -channel single top-quark events were produced using the POWHEG Box v2 generator, which provides MEs at NLO in the strong coupling constant α_s in the 4FS ² with the NNPDF3.0NLO nf4 PDF set [180–185]. The functional form of the μ_R) and μ_F was set to $\sqrt{m_b^2 + p_{T,b}^2}$ following the recommendation of Ref. [181], where

²The event generation is done in the 4FS, where the bottom quark is treated as a massive quark

m_b and $p_{T,b}$ are the mass and p_T of the b -quark from the initial gluon splitting, so-called "second b -quark." This choice of scales ensures that the kinematics of the bottom quark originating from gluon splitting are properly accounted for. The top quark and the W boson were decayed at LO using MadSpin [186, 187] to preserve all spin correlations, which is crucial for accurately simulating the angular distributions and polarisations of decay products. For this sample, it was required that the Breit-Wigner (BW) distribution, generated in MadSpin, of the invariant mass of the decay products was distributed up to 50 times the particle width around the particle mass (i.e. it was set at a cutoff of 50 GeV). This configuration ensures an accurate simulation of the resonance behaviour, accounting for the finite lifetime of the particle and capturing the correct resonance behaviour in simulations.

4.4.2.2 Background processes

Two major backgrounds and others

- **Top-quark pair ($t\bar{t}$):** A major background in t -channel single top-quark analyses, particularly in dileptonic final states where one lepton may be undetected, mimicking the signal signature. Semileptonic $t\bar{t}$ events also contribute when fewer than four jets are reconstructed.

The baseline $t\bar{t}$ events were generated using the POWHEG BOX v2 generator, with MEs at NLO in α_s in the SFS and the NNPDF3.0NLO PDF set [185]. The renormalisation and factorisation scales were dynamically set to $\sqrt{m_{\text{top}}^2 + p_T^2}$, and the h_{damp} parameter was set to $1.5 m_t$ to control high- p_T radiation [188]. Events were interfaced with PYTHIA 8.230 for PS and hadronisation using the A14 tune and the NNPDF2.3LO PDF set [189]. Top quarks were decayed internally by POWHEG, and bottom and charm hadrons were decayed using EVTGEN 1.6.0 [190].

The event sample was normalised to the NNLO+NNLL cross-section prediction of $\sigma_{t\bar{t}} = 832 \pm 51$ pb at $\sqrt{s} = 13$ TeV for $m_t = 172.5$ GeV, calculated with TOP++ 2.0 [191, 192]. PDF and α_s uncertainties were estimated using the PDF4LHC15 prescription [193].

Modelling uncertainties were evaluated using alternative simulation samples and re-weighting techniques. The uncertainty from PS and hadronisation modelling was assessed by comparing the nominal POWHEG +PYTHIA sample to one interfaced with HERWIG 7, using the H7UE tune and MMHT2014LO PDF set [194, 195]. Initial- and final-state radiation (ISR/FSR) uncertainties were evaluated by varying the renormalisation and factorisation scales by factors of 0.5 and 2.0, and by using A14 Var3c ISR tune variations [196, 197].

The impact of the NLO matching scheme was assessed by changing the `POWHEG:pThard` parameter from 0 to 1, altering the definition of the shower veto region [198]. Uncertainties in b -quark recoil modelling were evaluated using alternative simulations where the gluon recoil scheme was modified from the default `recoilToColor` to `recoilToTop`, affecting the b -jet fragmentation and hadron kinematics.

All alternative samples assumed a top-quark mass of $m_t = 172.5\text{GeV}$ and a 100% branching ratio to $t \rightarrow Wb$, and were processed through either full GEANT4 or `Atlfast2` detector simulation.

- **W +jets production:** The W +jets production is the second major background for the single top-quark t -channel process. QCD W +jets production is simulated with the Sherpa (v2.2.11) [199] ME+PS generator. Additional hard parton emissions [200] are matched to a PS based on Catani-Seymour dipoles [201]. The simulation uses the NNPDF3.0 NNLO PDFs [185], along with a set of tuned PS parameters specifically developed by the Sherpa authors for this version.

The ME+PS matching [202] is employed for different jet multiplicities, which are then merged into an inclusive sample using an improved CKKW matching procedure [203, 204], extended to NLO accuracy using the MEPS@NLO prescription [205]. ME+PS matching combines the accuracy of ME calculations for hard scatterings with the detailed description of PS simulations for subsequent parton emissions, ensuring no double counting and a smooth transition between the two.

These particular simulations are NLO accurate for up to two additional partons and LO accurate for up to four additional partons. The virtual QCD correction for MEs at NLO accuracy are provided by the OpenLoops library [206, 207].

- **Other backgrounds:**

Single top-quark tW and s -channel productions: The single top-quark tW associated production and the s -channel process are subdominant backgrounds in the t -channel single top-quark analysis, but are included due to their similar final-state topologies. Both processes are simulated using the POWHEG BOX (v2) generator, which provides MEs at NLO accuracy in the strong coupling constant α_s in the 5FS, using the NNPDF3.0 NLO PDF set [185]. The renormalisation and factorisation scales, μ_R and μ_F , are set to the default functional form $\sqrt{m^2 + p_T^2}$, where m and p_T refer to the mass and transverse momentum of the produced particle.

For tW production, the diagram removal (DR) scheme [208] is employed to handle interference with $t\bar{t}$ production [188], removing the overlapping doubly-resonant diagrams from the amplitude. This choice avoids interference at the cost of gauge

invariance but remains consistent within the precision goals of this analysis. While the t - and s -channel processes do not interfere even at NLO in QCD, interference effects between tW and $t\bar{t}$ can be significant and are thus carefully treated.

Events are interfaced with PYTHIA 8.230 [189] for PS, hadronisation, and underlying event modelling using the A14 tune and the NNPDF2.3 LO PDF set. Bottom and charm hadron decays are simulated using EVTGEN (v1.6.0). Top-quark decays are handled at LO using MADSPIN [169, 170] to preserve spin correlations in the final state.

The predicted inclusive cross-section for tW production is $\sigma(tW) = 79.3^{+2.9}_{-2.8}$ pb at $\sqrt{s} = 13$ TeV, computed at NLO QCD including next-to-next-to-leading logarithmic (NNLL) soft-gluon resummation [209]. For the s -channel, the cross-section is calculated at NLO using HATHOR (v2.1) [210, 211], yielding $\sigma(s\text{-channel}) = 10.32^{+0.40}_{-0.36}$ pb for inclusive top and anti-top production, with values of 6.35 pb and 3.97 pb for top and anti-top separately.

The quoted theoretical uncertainties include variations of μ_R and μ_F and PDF+ α_s uncertainties derived using the MSTW2008 [146], CT10 [212], and NNPDF2.3 [213] PDF sets. These are added in quadrature to the scale uncertainty following the PDF4LHC recommendations [214].

Additional simulation samples and internal weights are used to evaluate systematic uncertainties related to initial- and final-state radiation (ISR/FSR), PDF variations, and PS models. ISR is evaluated by varying μ_R and μ_F in the ME and PS, using the Var3c up/down variations of the A14 tune. FSR uncertainties are assessed using weights that vary the renormalisation scale in the shower by factors of 0.5 and 2.0. All modelling uncertainties are evaluated using either full GEANT4 or **At1fast2** simulation of the ATLAS detector.

Multijet, Z +jets, and diboson (WW , ZZ , WZ) processes: Multijet events, arising from QCD processes, contribute mainly through non-prompt leptons and are estimated using a data-driven matrix method described in Section 4.4.3. The Z +jets production is simulated similarly to W +jets using the SHERPA (v2.2.11) ME+PS generator [199], employing Catani–Seymour dipole-based PS and matched MEs. Diboson processes (WW , ZZ , WZ) are also simulated using SHERPA (v2.2), with NLO accuracy for up to one additional parton and LO for up to three partons, incorporating OpenLoops [206] for one-loop MEs.

$t\bar{t}$ and single top-quark productions associated with vector or Higgs bosons ($t\bar{t}Z$, $t\bar{t}W$, $t\bar{t}H$, tZq , $tHq+tWH$, and tWZ): The $t\bar{t}Z$ and $t\bar{t}W$ processes are simulated using MADGRAPH5_aMC@NLO (v2.3.3) with matrix elements calculated at NLO in QCD, using the NNPDF3.0 NLO PDF set. The events

are interfaced with PYTHIA 8 (v8.210) for PS and hadronisation, employing the A14 tune and NNPDF2.3 LO PDF set. The cross-sections are normalised to the predictions obtained at NLO in QCD and electroweak corrections as reported in Ref. [215], with ttZ further corrected for off-shell contributions down to 5 GeV [216]. The ttH production is modelled using the POWHEG BOX generator at NLO with the NNPDF3.0 NLO PDF set, and interfaced with PYTHIA 8 (v8.230) using the A14 tune. Cross-sections are taken from NLO+NLO EW predictions [215].

The tZq process is generated with MADGRAPH5_aMC@NLO (v2.3.3) using the NNPDF3.0 NLO PDF set, interfaced with PYTHIA 8 (v8.212). The 4FS is used, where all quark masses are set to zero except for the top and bottom quarks. Events include at least one lepton and off-resonant contributions are excluded.

The associated production of a single top quark with a Higgs boson (tHq and tHW) is simulated with MADGRAPH5_aMC@NLO (v2.6.2 for tHq and v2.8.1 for tHW) at NLO, using the NNPDF3.0 NLO PDF set. Events are interfaced with PYTHIA 8 using the A14 tune. Both 4FS and 5FS schemes are employed. For tHW , the DR1 scheme is used to suppress interference with ttH production.

The tWZ process is simulated with MADGRAPH5_aMC@NLO (v2.3.3), using the NNPDF3.0 NLO PDF set and interfaced with PYTHIA 8 (v8.212) for showering and hadronisation. The DR1 scheme is used to mitigate interference with ttZ production, and events include an inclusive top decay and a Z boson decaying to a lepton pair.

4.4.3 Multijet background estimation

The multijet (or “fake”) background, arising from jets misidentified as leptons or leptons from heavy-flavour decays, is estimated using a data-driven matrix method. Two categories of lepton identification are needed and therefore defined: *LOOSE* and *TIGHT*, where *TIGHT* is a subset of *LOOSE*. Each category includes both real and fake leptons, and the total number of events is modelled as:

$$\begin{aligned} N_{\text{LOOSE}} &= N_{\text{LOOSE}}^{\text{real}} + N_{\text{LOOSE}}^{\text{fake}}, \\ N_{\text{TIGHT}} &= N_{\text{TIGHT}}^{\text{real}} + N_{\text{TIGHT}}^{\text{fake}}. \end{aligned} \tag{4.3}$$

Real and fake efficiencies, defined as the probability of a *LOOSE* lepton passing the *TIGHT* selection, are measured in dedicated control regions:

$$\epsilon_{\text{real}}(\vec{k}) = \frac{N_{\text{TIGHT}}^{\text{real}}(\vec{k})}{N_{\text{LOOSE}}^{\text{real}}(\vec{k})}, \quad \epsilon_{\text{fake}}(\vec{k}) = \frac{N_{\text{TIGHT}}^{\text{fake}}(\vec{k})}{N_{\text{LOOSE}}^{\text{fake}}(\vec{k})}, \tag{4.4}$$

where $\vec{k}$ denotes kinematic variables such as p_T and $|\eta|$.

Using these efficiencies, the number of fake leptons in the TIGHT sample is estimated via:

$$N_{\text{TIGHT}}^{\text{fake}} = \int \frac{\epsilon_{\text{fake}}(\vec{k})}{\epsilon_{\text{real}}(\vec{k}) - \epsilon_{\text{fake}}(\vec{k})} [\epsilon_{\text{real}}(\vec{k}) N_{\text{LOOSE}}(\vec{k}) - N_{\text{TIGHT}}(\vec{k})] d\vec{k}. \quad (4.5)$$

This method allows for an event-by-event weight to be assigned, depending on the lepton selection, enabling accurate modelling of non-prompt background contributions directly from data. This is the method used in the analysis presented in this thesis to estimate the multijet background.

4.4.4 Simulated samples with EFT parameter variations

As discussed in Chapter 2, the top-quark sector provides a particularly sensitive probe for new physics effects parametrised via EFT operators. In this framework, deviations from the SM can be described by Wilson coefficients associated with higher-dimensional operators, which can modify both production and decay dynamics. To evaluate how these coefficients influence observables—particularly the four angular variables used in this analysis to describe the kinematics of top-quark decays—it is necessary to generate MC samples for a range of EFT scenarios. However, due to the large number of possible combinations of Wilson coefficients—each potentially varied over multiple values and possibly contributing simultaneously—the parameter space becomes high-dimensional, making it computationally impractical to generate dedicated MC samples for every scenario.

A more efficient approach is to use the re-weighting method. This method allows us to adjust the weights of events in an existing MC simulation to reflect the impact of new EFT values. By reweighting events, the need to generate a large number of dedicated MC samples is avoided, enabling the use of existing simulated datasets.

In this analysis, MadGraph is employed to perform the reweighting procedure. MadGraph helps generate the necessary MEs for our specific processes. For example, in the signal process $qb \rightarrow q't \rightarrow q'Wb \rightarrow q'(l^+\nu_l)b$ the interaction amplitudes at LO for the involving sub-processes would be:

$$\mathcal{M}_1 = \bar{u}(k_1) \left(\frac{ig}{\sqrt{2}} \gamma^\mu (1 - \gamma^5) \right) \frac{-i(g_{\mu\nu} - \frac{q_\mu q_\nu}{M_W^2})}{q^2 - M_W^2} \left(\frac{ig}{\sqrt{2}} \gamma^\nu (1 - \gamma^5) \right) u(p_1). \quad (4.6)$$

$$\mathcal{M}_2 = \bar{u}(k_4) \left(\frac{ig}{\sqrt{2}} \gamma^\mu (1 - \gamma^5) \right) u(k_2) \epsilon_\mu^*(k_3). \quad (4.7)$$

$$\mathcal{M}_3 = \bar{v}(k_5) \left(\frac{ig}{\sqrt{2}} \gamma^\mu (1 - \gamma^5) \right) u(k_6) \epsilon_\mu(k_3). \quad (4.8)$$

To implement the re-weighting, the MEs are calculated for both the original MC scenario and the modified EFT scenario with varied Wilson coefficients. The re-weighting formula for the events is then given by the ratio of these MEs:

$$w_{\text{new}} = \frac{|\mathcal{M}_{\text{new}}|^2}{|\mathcal{M}_{\text{old}}|^2} w_{\text{old}}$$

Here, $|\mathcal{M}_{\text{new}}|^2$ and $|\mathcal{M}_{\text{old}}|^2$ are the squared MEs for the new EFT scenario and the original MC scenario, respectively, and w_{old} is the original event weight. The MEs are summed over all helicities, meaning they consider contributions from all possible spin configurations of the particles involved. This comprehensive summation ensures that the calculated weights accurately reflect the physical processes under the new EFT values.

For the t -channel signal sample, a grid of 582 re-weighted points in the six-dimensional parameter space $(C_{tW}, C_{itW}, C_{bW}, C_{ibW}, C_{\phi tb}, C_{qQ})$ is generated, as summarised in Table 2.1. The grid of Wilson coefficient values is subject to a constraint to avoid being too far from the SM compared to current limits:

$$0.5 < \sum_k \kappa_k^2 < 16.5,$$

where κ_k are the indices of the corresponding Wilson coefficients C_k in the grid loop. Here, κ_k indices for C_k coefficients are $\{\pm 3, \pm 2, 0\}$, except for C_{qQ} , which uses $\{\pm 2, 0\}$. The coefficients are given by:

$$C_k = \beta_k \kappa_k$$

Bases β_k are set to 1.0 for C_{tW}, C_{bW}, C_{ibW} ; 0.5 for C_{itW}, C_{qQ} ; and 5.0 for $C_{\phi tb}$.

The W -boson helicity fraction F_R is zero when the relationship $\frac{m_W}{2} C_{\phi tb} = \sqrt{2} m_t C_{bW}$ holds. Given the stringent constraints on F_R , the values of C_{bW} are selected to be either on or parallel to this line.

$$C_{bW} = \frac{m_W}{2\sqrt{2}m_t} C_{\phi tb} + \beta_{bW} \kappa_{bW} = 0.16 C_{\phi tb} + \beta_{bW} \kappa_{bW} \quad (4.9)$$

Since C_{tW} , C_{itW} , C_{bW} , C_{ibW} , and $C_{\phi tb}$ contribute to both top-quark production and decay, a full parameterisation of the morphing re-weighting points falls within the same $\bar{\kappa}$ radius of 16.5, resulting in a final set of 743 re-weighting points evaluated in the $t\bar{t}$ EFT sample.

For the $t\bar{t}$ events, most of the EFT operators considered affect the top-quark production and decay. Since $t\bar{t}$ events are major background events of this t -channel analysis, a designated re-weighted sample for $t\bar{t}$ events is produced. It uses the same parameter settings as described above, with the C_{qQ} coefficient replaced by the C_{tG} coefficient, as C_{qQ} only affects single-top-quark production. The re-weighting points are not changed significantly since C_{tG} also only enters in the production side of $t\bar{t}$ events. Thus, its κ value only takes 3 values: $\pm 1, 0$. Its base is assigned to be 0.5 instead of 1.5 for C_{qQ} .

Wilson Coefficient	Sampled Values
C_{tW}	0.0, ± 2.0 , ± 3.0
C_{itW}	0.0, ± 1.0 , ± 1.5
$C_{\phi tb}$	0.0, ± 10.0 , ± 15.0
C_{bW}	$0.16C_{\phi tb} + (0.0, \pm 2.0, \pm 3.0)$
C_{ibW}	0.0, ± 2.0 , ± 3.0
C_{gqQ}	0.0, ± 3.0
C_{tG}	0.0, ± 0.5

Table 4.3: Sampled values for each Wilson coefficient used in this analysis.

Following the development of a robust re-weighting framework capable of efficiently probing the impact of various EFT operators on the relevant observables, the analysis proceeds with the implementation of a sophisticated morphing procedure to predict experimental outcomes under different theoretical scenarios.

Top Quark Width Correction

The top quark width, Γ_{top} , encodes the total decay rate of the top quark and plays a critical role in the normalisation and shape of matrix elements involving resonant top-quark propagators. As the only quark that decays before hadronisation, its width is not only theoretically well-defined but also directly relevant to precision measurements and EFT interpretations. Modifications to the EW top-quark couplings—particularly through higher-dimensional operators—can induce deviations in Γ_{top} from its SM value.

In EFT-based event generation, however, tools like **MadGraph** compute matrix elements

with a fixed input width and do not automatically adjust Γ_{top} when Wilson coefficients are varied. This leads to inconsistencies, especially in processes such as single top quark or $t\bar{t}$ production, where the propagator sensitivity to the width affects both kinematic distributions and total rates. To address this, a correction factor is introduced at the reweighting stage by rescaling event weights according to the ratio $\Gamma_{\text{top}}^{\text{SM}}/\Gamma_{\text{top}}(\vec{C})$. The correction enters linearly in processes with a single-top-quark propagator (e.g. t -channel production), and quadratically in $t\bar{t}$ production:

$$w_{\text{new}}^{t\text{-chan}} = w_{\text{old}} \cdot \frac{\Gamma_{\text{top}}^{\text{SM}}}{\Gamma_{\text{top}}(\vec{C})}, \quad w_{\text{new}}^{t\bar{t}} = w_{\text{old}} \cdot \left(\frac{\Gamma_{\text{top}}^{\text{SM}}}{\Gamma_{\text{top}}(\vec{C})} \right)^2.$$

The width is evaluated using `MadWidth`³ for each EFT parameter point. Among the contributing operators, C_{tW} has the dominant impact on Γ_{top} , and the dependence can be approximated as:

$$\frac{\Gamma_{\text{top}}}{\Gamma_{\text{top}}^{\text{SM}}} \approx 1 + 0.1665 C_{tW} + 0.0113(C_{tW}^2 + C_{itW}^2).$$

This correction ensures the internal consistency of the simulation and accurate modelling of width-sensitive observables under EFT hypotheses.

4.5 Morphing Procedure

In order to accurately predict the behaviour of observables under different EFT hypotheses, the analysis presented in this thesis employs a morphing procedure. The morphing procedure enables the modelling of how observables change with respect to variations in the Wilson coefficients.

The primary motivation for morphing arises from the need to handle the vast parameter space introduced by EFTs. When studying processes such as single top quark production, the cross-section and other observables depend on several Wilson coefficients. Simulating events for every possible combination of these coefficients is computationally infeasible. Instead, morphing allows the use of a finite number of simulated points (re-weighting points) to construct a continuous and smooth model that describes the behaviour of observables across the entire parameter space.

The M-function coefficients $c_{m'm}^{j_1 j_2}$, defined in Equation 2.21, are proportional to the cross-section σ of the process. Cross section, representing total production rate is

³`MadWidth` is a module of `MadGraph5_aMC@NLO` used to compute tree-level particle decay widths, including contributions from new physics operators.

expressed with all the configurations of all the states, over SM and beyond $\sigma \propto |M|^2$. For the t-channel single-top events, the tWb vertex appears twice: once in the production and once in the decay of the top quark.

$$|M|^2 \sim \left| M_{\text{SM}} + \sum_i C_i M_i + \sum_{i,j} C_i C_j M_{ij} + \dots \right|^2. \quad (4.10)$$

Therefore the polynomial dependence in the cross-section can involve up to factor of four of the Wilson coefficients, such as C_{tW} , C_{itW} , C_{bW} , C_{ibW} , and C_{ptb} . For coefficients like C_{qQ} and C_{tG} that only affect the production of single-top-quark or $t\bar{t}$ events, the dependence is at most second-order. The total number of independent terms in a fourth-order polynomial with six parameters is given by:

$$\binom{n+k}{k} = \binom{6+4}{4} = 210,$$

where $n = 6$ (number of parameters) and $k = 4$ (order of the polynomial). Out of these, 203 terms are independent after accounting for higher-order dependencies.

To determine the coefficients of this polynomial accurately, one needs to sample the parameter space efficiently. This is achieved using the re-weighting module of the MADGRAPH generator.

From the grid of re-weighting points, a set of orthonormal polynomials $P_i(\vec{C})$ are constructed ensuring the independence of each polynomial making the morphing function accurate:

$$\sum_k P_i(\vec{C}_k) P_j(\vec{C}_k) = \delta_{ij}. \quad (4.11)$$

Here, $\vec{C}_k$ represents the values of the Wilson coefficients at the k -th re-weighting point, δ_{ij} is the Kronecker delta.

The design matrix M is constructed by evaluating each orthonormal polynomial at the re-weighting points:

$$M_{ik} = P_i(\vec{C}_k). \quad (4.12)$$

To morph any observable O , its value at each re-weighted point O_k is computed to form a vector $\vec{O}$. The morphing function for the observable O at LO is given by:

$$f_O^{\text{LO}}(\vec{C}) = \sum_i \eta_i P_i(\vec{C}) = \vec{\eta}^T \vec{P}(\vec{C}). \quad (4.13)$$

The coefficients $\vec{\eta}$ of the morphing function for each of the basis polynomials are obtained by solving the equation $\vec{\eta} = M\vec{O}$. The covariance matrix associated with $\vec{\eta}$ is $\mathbb{C}_\eta = M\mathbb{C}M^T$, where $\mathbb{C}$ is the covariance matrix of the observable O for each re-weighting point.

The morphing functions are applied to Monte Carlo samples corresponding to single top-quark t -channel production, as well as semileptonic and dileptonic decays of top-quark pair ($t\bar{t}$) events. In the analysis, the morphed observables are the event counts N in the signal region and the 23 M-function coefficients. The top and anti-top channels are treated independently, with the morphing functions produced separately for each.

To compute the overall correction factor from LO to NLO, the difference between the MadGraph sample at the SM point and the NLO PowhegBox+Pythia sample is determined as follows:

$$f_O^{\text{NLO}}(\vec{C}) = f_O^{\text{LO}}(\vec{C}) + (f_O^{\text{NLO}}(\vec{C}_{\text{SM}}) - f_O^{\text{LO}}(\vec{C}_{\text{SM}})). \quad (4.14)$$

In this context, for the single top-quark t -channel process, PowhegBox+Pythia samples are used for both top and anti-top-quark production. For the $t\bar{t}$ process, a corresponding PowhegBox+Pythia sample is employed to model the top-quark pair production.

Another coefficient that is introduced in the fit is coefficient $C_{\phi q}$, and is not morphed. Instead, the dependency of $C_{\phi q}$ can be derived from the t -channel scale factor as described in Equation 4.15. $C_{\phi q}$ represents interaction between the Higgs field ϕ and the quark doublet q .

$$f^{t\text{-channel}} = |V_L|^2 = \left(1 + \frac{v^2}{\Lambda^2} C_{\phi q}\right)^2. \quad (4.15)$$

To conclude the chapter on the morphing procedure and its implementation in this analysis, a set of representative plots is presented to demonstrate how Wilson coefficients affect key observables used in the fit. Each figure includes results separately for both top-quark and top-antiquark channels, reflecting the charge-specific structure of the EFT contributions and the analysis strategy.

Figure 4.5 shows the total morphed event yields in the signal region for both top-quark (upper) and top-antiquark (lower) channels. The variation with respect to one Wilson coefficient at a time is displayed, while the remaining coefficients are set to zero. Each subplot stacks the contributions from all subprocesses, with the solid black line indicating

the total prediction and the triangle marking the SM point. This figure illustrates how different operators affect the overall event normalization.

Two-dimensional interpolations in EFT parameter space are shown in Figure 4.6. The left and right panels display results for pairs of coefficients (C_{ptb}, C_{bw}) and (C_{ptb}, C_{ibw}) , respectively, for both top-quark (upper) and top-antiquark (lower) channels. These plots reveal non-trivial correlations and interference effects between operators, highlighting the importance of multidimensional modelling.

Figure 4.7 displays the morphed real part of the angular M-function coefficient $c_{1,0}^{1,0}$, separately for top-quark and top-antiquark events. As before, each panel shows the dependence on one Wilson coefficient, with others fixed at their SM values. Below each panel, the ratio to the SM prediction is presented to emphasize the size of the deviation. These results illustrate how EFT modifications impact angular distributions, which are crucial in this analysis.

The morphing technique successfully captures the dependence of both total yields and shape-sensitive observables on Wilson coefficients. This approach enables efficient interpolation across the parameter space and underpins the statistical interpretation presented in subsequent chapters.

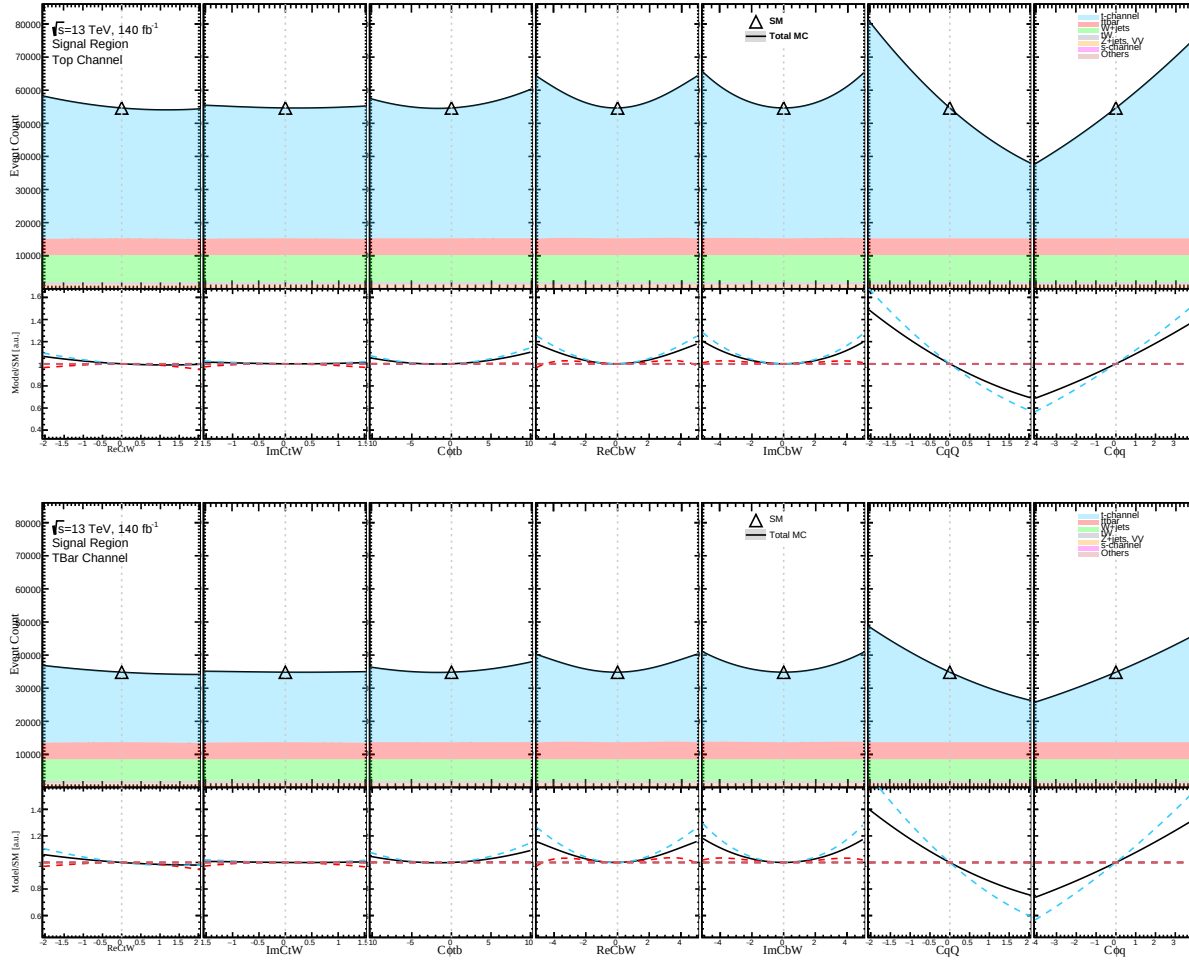


Figure 4.5: Morphed event counts for the top-quark channel (upper) and anti-top-quark (lower) channel. Each subplot corresponds to the morphed event counts as a function of one of the Wilson coefficients, with all others fixed at zero. All subprocesses are stacked on top of each other, with the solid black line representing the sum of each subprocess, and the triangle point corresponding to the SM. An MC dataset corresponding to the full LHC Run 2 data is used. The scale factors for the three major processes are fixed at $\vec{S} = \{1, 1, 1\}$. For the $C_{\phi q}$ reweighted points are not generated, instead plot shows how Equation 4.15 influenced the event counts.

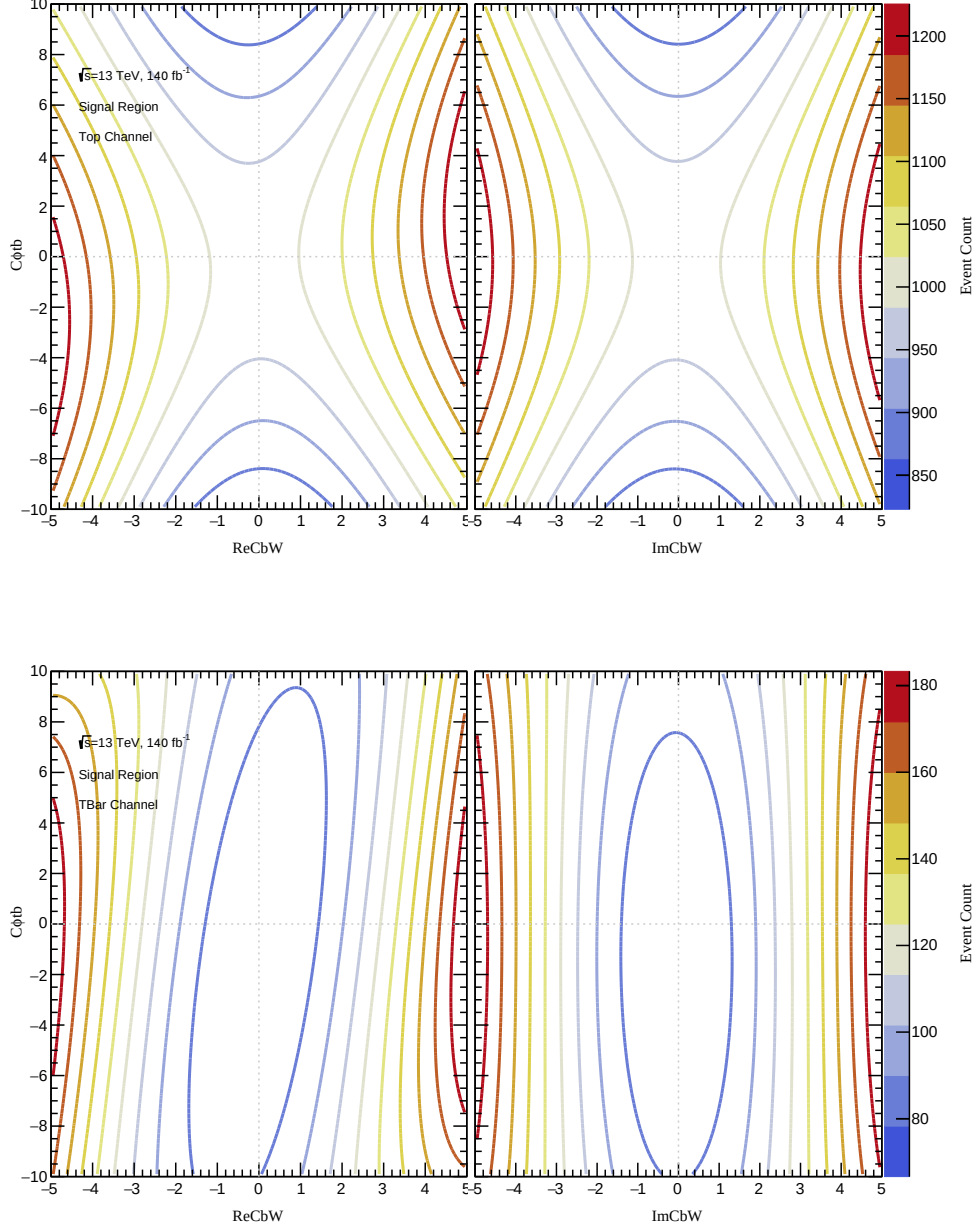


Figure 4.6: Morphed event counts for top-quark channel (up) and top-antiquark (bottom) channel $N^{\pm, \text{SR}}(C_{ptb}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{ptb}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $1 \vec{S} = \{1, 1, 1\}$.

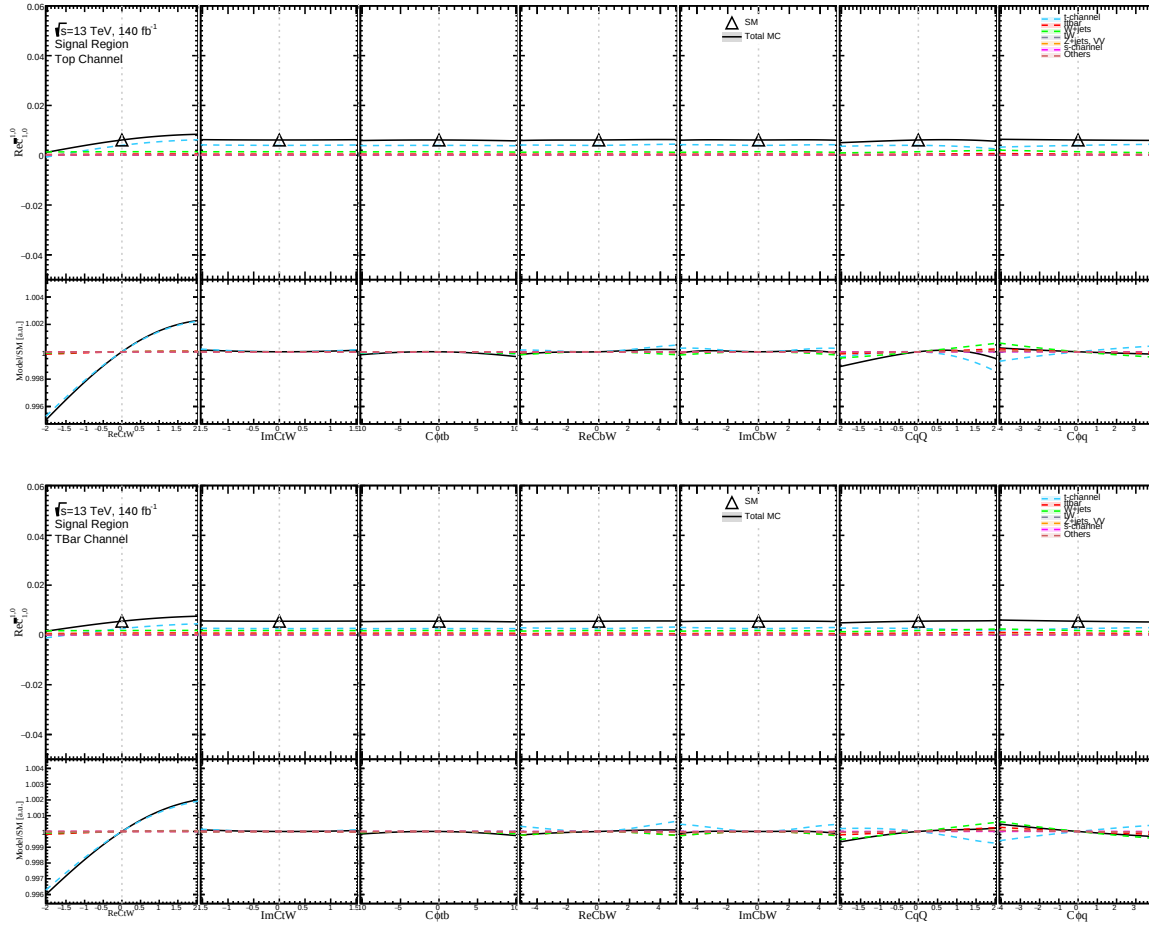


Figure 4.7: Morphed $\Re\langle c_{1,0}^{1,0} \rangle$ for top-quark channel (up) and anti-top-quark (bottom) channel $\langle \text{Re } c_{1,0}^{1,0} \rangle^{\pm, \text{SR}}(C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses overlap each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point. The scale factors of the three major processes are fixed at 1 $\tilde{S} = \{1, 1, 1\}$.

Chapter 5

Object definition, identification and reconstruction

In the ATLAS experiment, object definition, identification and reconstruction are an essential processes for converting raw detector data into useful information. This starts with tracking charged particles using the ID. The ID measures the paths of particles accurately by reconstructing their tracks from hits in the detector layers, using algorithms that take into consideration magnetic field effects, interactions with detector material, and multiple scattering as discussed in Chapter 3. Vertex reconstruction is another crucial part, where the primary and secondary vertices are identified. Primary vertices are where particle collisions occur, and secondary vertices are related to decays of unstable particles. These techniques allow precise determination of vertices, which is important for identifying collision events and decay chains. By measuring the paths of charged particles using the inner detector, and combining this information with the energy deposits in the electromagnetic and hadronic calorimeters, as well as the momentum measurements from the MS, the ATLAS detector can reconstruct and identify physics objects such as electrons, muons, jets, and missing transverse energy. The calorimeters measure the energy of electromagnetic and hadronic showers, while the MS provides precise momentum measurements of muons that penetrate the calorimeters. The $E_{\text{T}}^{\text{miss}}$ is inferred using transverse momentum conservation, serving as an indicator of undetected particles such as neutrinos [113, 217].

This chapter describes the reconstruction, identification, and selection criteria applied to physics objects used in the analysis. The chapter begins with a discussion of tracking and vertexing in the Section 5.1, which underpins the reconstruction of charged particles. The trigger strategy employed to select relevant events is summarised in Section 5.2. The reconstruction and selection of primary analysis objects—electrons (Section 5.3),

muons (Section 5.4), and τ -leptons (Section 5.5)—are described in detail, including their isolation and impact parameter requirements.

The reconstruction of jets is outlined in Section 5.6, with further details on the identification of b -jets via multivariate tagging algorithms provided in Section 5.6.1. The calculation of the missing transverse momentum is discussed in Section 5.7, including the treatment of soft contributions and pile-up mitigation. Finally, Section 5.8 presents the procedure used for resolving overlaps between reconstructed objects to avoid ambiguities and double-counting.

5.1 Tracking and vertexing

As explained in Chapter 3, Section 3.2.2, charged-particle trajectories are reconstructed in the ID using hits from the silicon-based Pixel and SCT detectors, along with the TRT. Vertex reconstruction is performed by clustering tracks based on spatial proximity and applying a χ^2 -minimisation procedure. The primary vertex is defined as the vertex with the highest scalar sum of squared transverse momenta of associated tracks with $p_T > 400$ MeV.

5.1.1 Track reconstruction

Track reconstruction in this analysis follows the standard ATLAS procedure. Space-points are first formed from hits in the Pixel and SCT sub-detectors and are used to seed track candidates. For tracks within $|\eta| < 2.0$, the TRT provides additional one-dimensional measurements that enhance the momentum resolution and assist in electron identification via transition radiation signatures.

Tracks are subject to quality requirements, including thresholds on the number of silicon hits, impact parameter significance, and fit quality, to suppress fake tracks and ensure good resolution. Only tracks passing these criteria are considered in the reconstruction and selection of physics objects. In this analysis, track information is used for lepton isolation, jet reconstruction, impact parameter-based b -jet identification, and pile-up rejection through vertex association. Track-based observables also contribute to the definition of inputs used in the fit.

5.2 Triggers

As introduced in Chapter 3, Section 3.2.6, the ATLAS trigger system reduces the 40 MHz bunch crossing rate to an event storage rate of about 1 kHz using a two-level selection: a hardware-based L1 trigger and a software-based HLT. In this analysis, only un-prescaled single-lepton triggers are used, selected to ensure high and stable efficiency under varying pile-up conditions across Run 2 data taking.

The list of trigger paths used is summarised in Table 5.1, and the triggers are combined using a logical **OR**. To ensure accurate object-level selection, the reconstructed lepton is required to be geometrically matched to the corresponding trigger object in the event.

Year	Single-electron trigger	Single-muon trigger
2015	HLT_e24_lhmedium_L1EM20VH	HLT_mu20_iloose_L1MU15
	HLT_e60_lhmedium	HLT_mu50
	HLT_e120_lhloose	
2016–2018	HLT_e26_lhtight_nod0_ivarloose	HLT_mu26_ivarmedium
	HLT_e60_lhmedium_nod0	HLT_mu50
	HLT_e140_lhloose_nod0	

Table 5.1: Un-prescaled single-lepton trigger selections of the ATLAS trigger menu per lepton flavour per year. These single-lepton triggers are combined using a logical **OR**.

Electrons are triggered by requiring a calorimeter cluster matched to an ID track and satisfying identification criteria based on a likelihood discriminant. These criteria are categorised into working points (WP) known as **loose**, **medium**, and **tight**, which reflect increasingly stringent requirements on shower shapes, track–cluster matching, and isolation variables. In 2015, L1 required $E_T > 20$ GeV, followed by HLT selections with $E_T > 24$ GeV and **medium** identification. From 2016 onward, the threshold was raised to $E_T > 26$ GeV with **tight** identification and isolation. To maintain high efficiency at large p_T , complementary triggers were used across all years, requiring either **medium** electrons with $E_T > 60$ GeV or **loose** (non-isolated) electrons with $E_T > 120$ GeV in 2015 and $E_T > 140$ GeV in later years.

Muons are triggered by matching tracks in the ID and muon spectrometer. In 2015, muons required loose isolation and $p_T > 20$ GeV; in 2016–2018, the threshold increased to $p_T > 26$ GeV with tighter isolation. To retain efficiency at high p_T , an additional trigger

without isolation was used, requiring $p_T > 50$ GeV.

Trigger efficiencies in simulation are corrected using data-to-MC scale factors, which are valid down to 1 GeV above the respective trigger thresholds. To ensure operation in the efficiency plateau region, leptons are required to have $p_T > 25$ GeV (electrons) or $p_T > 21$ GeV (muons) in 2015, and $p_T > 27$ GeV from 2016 onward. In this analysis, however, a more stringent requirement of $p_T > 30$ GeV is applied to all leptons, as detailed in the following sections.

5.3 Electrons

Electrons in this analysis are reconstructed using standard ATLAS algorithms and selected based on identification, isolation, and impact parameter requirements, as summarised in Table 5.2 (left). These criteria are designed to enhance prompt lepton purity while suppressing backgrounds from non-prompt sources such as jets, photon conversions, and heavy-flavour decays.

Two identification WPs are used: `looseAndBLayerLH`¹ for pre-selected electrons, and `tightLH`². These likelihood-based WPs offer superior background rejection compared to cut-based selections.

Pre-selected electrons need to pass $p_T > 10$ GeV and $|\eta| < 2.47$. Furthermore, there is a requirement on the electron isolation, corresponding to `PLVTight`³ isolation WP, which is defined by a multivariate analysis (MVA) likelihood discriminant combining shower shape and track information to distinguish prompt electrons from electron candidates from hadronic jets, γ -conversions, and heavy-flavour hadron decays.

To ensure accurate association with the primary vertex, the associated track must satisfy $|z_0 \sin(\theta)| < 0.5$ mm and $|d_0|/\sigma(d_0) < 5$, using the definitions of longitudinal and transverse impact parameters introduced in Chapter 3, Section 3.2.3. Electron candidates with calorimeter clusters in the transition region between the barrel and endcap ($1.37 < |\eta^{\text{clust}}| < 1.52$) are excluded.

Additional selection criteria are applied in certain analysis regions to improve object

¹`looseAndBLayerLH` is a likelihood-based electron identification working point with loose selection criteria, including a requirement for a hit in the innermost (IBL or B-layer) pixel detector.

²`tightLH` is a likelihood-based working point with stringent requirements on shower shape, track-cluster matching, and track quality variables, optimised for prompt electron selection.

³`PLVTight` is an isolation working point based on a multivariate discriminant combining calorimeter and tracking information, tuned to reject non-prompt leptons from jets, photon conversions, and heavy-flavour decays.

purity. These include a requirement on the Electron Charge ID Selector (ECIDS)⁴ to suppress charge mis-identification, and selections on `DFCommonAddAmbiguity` and `ambiguityType`, as defined in the Electron Ambiguity Tool⁵ to veto e/γ ambiguity candidates consistent with photon conversions.

Scale factors for reconstruction, identification, and isolation are applied to simulated electrons to correct for efficiency differences between data and MC simulated events.

5.4 Muons

Muons in this analysis are reconstructed using standard ATLAS algorithms and selected based on identification, isolation, and track quality requirements. Pre-selected muons, used for overlap removal and initial event filtering, are required to have $p_T > 7$ GeV and $|\eta| < 2.5$, and must satisfy the `medium` identification WP recommendation. This working point imposes requirements on the number of hits in the ID and MS sub-detectors, as well as on the significance of the charge-over-momentum ratio q/p [218, 219]. Muons that fail the reconstruction quality criteria—typically due to limited momentum resolution—result in the corresponding event being removed, following the established recommendations.

To suppress contributions from non-prompt or misidentified muons, an additional isolation requirement is applied beyond the pre-selection stage. The `PLVTight` isolation working point is used, which combines calorimetric and tracking variables in a multivariate discriminant.

The associated track must satisfy the impact parameter requirements $|z_0 \sin(\theta)| < 0.5$ mm and $|d_0|/\sigma(d_0) < 3.0$, in accordance with the track parameter definitions and performance considerations discussed in Chapter 3, Section 3.3.1. These constraints enhance vertex association and reduce contamination from pile-up and displaced background sources.

Data-to-simulation scale factors for muon identification and isolation are applied to simulated events to correct for efficiency differences observed between data and MC simulated events [218]. A summary of the muon selection criteria is provided in Table 5.2 (right).

⁴ECIDS is a likelihood-based classifier used to suppress electrons with mis-identified charge, particularly relevant at high p_T .

⁵The Electron Ambiguity Tool removes ambiguous candidates likely originating from photon conversions by checking for multiple tracks near the electron cluster.

	Pre-selected electron	Pre-selected muon
Identification	looseAndBLayerLH	medium
Acceptance	$p_T > 10 \text{ GeV}$, $ \eta^{\text{clust}} < 2.47$ except $1.37 < \eta^{\text{clust}} < 1.52$	$p_T > 10 \text{ GeV}$, $ \eta < 2.5$
Impact parameter	$ d_0/\sigma(d_0) < 5.0$ $ z_0 \sin(\theta) < 0.5$	$ d_0/\sigma(d_0) < 3.0$ $ z_0 \sin(\theta) < 0.5$
	Electron	Muon
Identification	tightLH	medium
Isolation	PLVTight	PLVTight
Extra selections	ECIDS, e/γ ambiguity-cuts	

Table 5.2: Summary of the electron and muon object definitions. The selection requirements for actual electrons/muons are applied in addition to the pre-selected objects used for the overlap removal

5.5 Taus

Hadronically decaying tau leptons (τ_{had}) are reconstructed and identified according to the recommendations given from the corresponding ATLAS combined performance group. Candidates are required to have $p_T > 20 \text{ GeV}$, $|\eta^{\text{clust}}| < 2.5$, excluding the transition region $1.37 < |\eta^{\text{clust}}| < 1.52$, and must be associated with either one or three tracks. To suppress backgrounds from jets and electrons, τ_{had} candidates are required to pass the **medium** (or **loose**) jet identification criteria based on a recurrent neural network (RNN), as well as a loose electron veto based on a boosted decision tree (BDT). No dedicated muon veto is applied. The MVA-based tau energy scale is used for energy calibration, and corresponding scale factors are applied to simulated events to correct for differences in reconstruction and identification efficiency between data and MC simulation.

Leptonic decays of tau leptons (i.e., into electrons or muons) are implicitly included in the analysis through the respective electron and muon object selections described earlier. No dedicated treatment of leptonic tau decays is applied beyond those selections.

In this analysis, events containing reconstructed τ_{had} candidates are vetoed. The full definition of τ_{had} object criteria used for pre-selection and vetoing is summarised in Table 5.3.

	τ_{had}
Acceptance	$p_{\text{T}} > 20 \text{ GeV}$, $ \eta^{\text{clust}} < 2.5$ except $1.37 < \eta^{\text{clust}} < 1.52$
Number of tracks	1 or 3
Identification	RNN Medium (Loose)
Electron veto	electron BDT Loose (Loose)

Table 5.3: Summary of the τ_{had} object definitions with the loose criteria in parentheses. The selection requirements for actual τ_{had} are applied in addition to the pre-selected objects used for overlap removal.

5.6 Jets

Jets are collimated sprays of particles resulting from the hadronisation of quarks and gluons produced in high-energy pp collisions. Although quarks and gluons cannot be observed directly due to colour confinement, their presence is inferred from the energy deposits left by their hadronic decay products in the calorimeters. A schematic representation of a jet is shown in Figure 5.1.

In this analysis, jets are reconstructed using the anti- k_t clustering algorithm [221] with a radius parameter of $R = 0.4$, applied to particle-flow objects [222]. The reconstruction is performed using the **FastJet** package [223], resulting in the **AntiKt4EMPFlowJets** collection recommended for Run 2 ATLAS analyses.

To suppress contributions from pile-up and ensure jets originate from the hard-scatter interaction, a Jet Vertex Tagger (JVT) [224, 225] is applied, following the recommendations of the Jet/ $E_{\text{T}}^{\text{miss}}$ ATLAS combined performance group. Jets are retained if they have $p_{\text{T}} > 20 \text{ GeV}$ and $|\eta| < 4.5$. For jets with $p_{\text{T}} < 60 \text{ GeV}$ and $|\eta| < 2.4$, a requirement of $\text{JVT} > 0.5$ is imposed, corresponding to the **Tight** JVT working point. This selection efficiently rejects pile-up jets while maintaining high efficiency for jets from the primary vertex.

For jets in the forward region ($2.5 < |\eta| < 4.5$) and $p_{\text{T}} < 120 \text{ GeV}$, a forward JVT (**fJVT**) working point is applied, requiring **fJVT** < 0.4 in addition to a timing requirement, to further suppress pile-up contributions in regions without tracking coverage.

Jet calibration is performed using the standard procedure described in Ref. [226]. This corrects the reconstructed jet energy to match, on average, the true energy of the

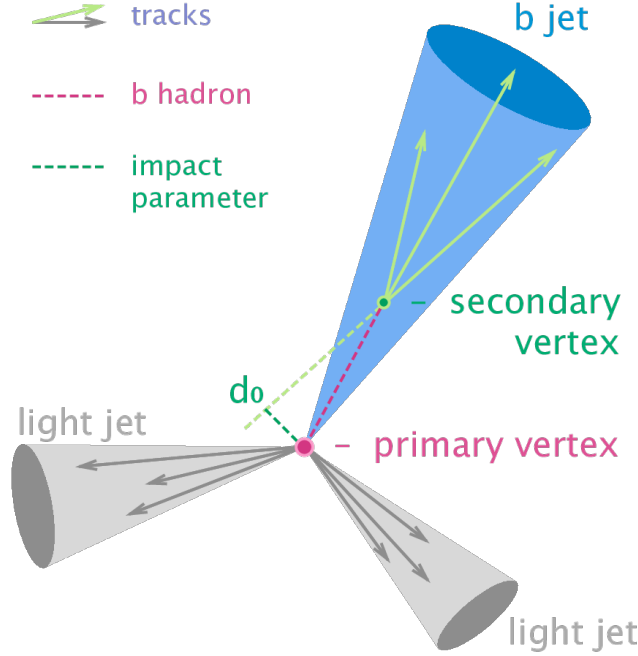


Figure 5.1: Schematic illustration of a *b*-jet showing the displacement of the secondary vertex from the primary vertex due to the long-lived *b*-hadron, contrasted with light jets. Adapted from the Ref [220].

jet at particle level in simulated data. An additional step, known as in-situ calibration, uses data-driven techniques to refine the jet energy scale by comparing the jet response to well-measured reference objects.

In this analysis, an additional requirement of $p_T > 30$ GeV is applied to all jets used in event selection. The full jet selection is summarised in Table 5.4.

5.6.1 *b*-tagged jets

The identification of jets originating from the hadronisation of *b*-quarks (named *b*-jets) is referred to as *b*-tagging. The purpose of *b*-tagging is to distinguish *b*-jets from jets initiated by *c*-quarks, known as forward jets or by gluons and light-flavour quarks (named light-flavour jets or simply light-jets). The flavour identification is achieved using dedicated algorithms that exploit the long lifetime and decay characteristics of *b*-hadrons.

In this analysis, the MVA *b*-tagging algorithm DL1r is used. It combines information from the impact parameters of displaced tracks and the topological properties of secondary and tertiary decay vertices reconstructed within the jet [227–229]. A jet is considered *b*-tagged if the value of the DL1r discriminant exceeds a WP. Four WPs are defined for the DL1r tagger, corresponding to *b*-jet selection efficiencies of 85%, 77%, 70%, and 60%

in simulated $t\bar{t}$ events. In this analysis, the 60% WP is used for nominal event selection. This choice provides a tighter selection, reducing background contamination at the cost of lower b -jet acceptance. Additionally, a pseudo-continuous b -tagging approach is employed in the analysis by dividing the continuous DL1r discriminant into multiple intervals. This allows for a more fine-grained classification of jets based on their b -likelihood and supports the use of MVA techniques. The corresponding scale factors are derived using the eigenvector method, with 45 decorrelated variations provided for b -jets.

To ensure reliable modelling of b -tagging in simulation, the performance of the DL1r tagger is measured in collision data. The b -jet tagging efficiency [230] and the forward jet mis-tagging rate [231] are evaluated using $t\bar{t}$ events, while the light-jet mis-tagging rate calibration relies on events containing Z boson, using a method similar to that presented in Ref. [232]. Correction factors derived from these calibration analyses are applied to simulated events to account for residual differences between data and MC simulation.

Pre-selected jet	
Collection	AntiKt4EMPFlowJets
Acceptance	$p_T > 20 \text{ GeV}, \eta < 4.5$
Jet Vertex Tagger	> 0.5 if $ \eta < 2.4$ and $p_T < 60 \text{ GeV}$
	< 0.4 if $2.5 < \eta < 4.5$ and $p_T < 120 \text{ GeV}$
b -tagging jet	
Acceptance	$p_T > 20 \text{ GeV}, \eta < 2.5$
b -tagging	DL1r algorithm
	Available WPs: 85% , 77% , 70% and 60%

Table 5.4: A summary of the jet selection criteria and b -tagging requirements, including the use of the 60% working point for b -jet identification in this analysis.

5.7 Missing transverse momentum

The missing transverse momentum, denoted by the vector $\vec{p}_T^{\text{miss}}$, represents the imbalance in the total transverse momentum of visible particles in a pp collision event. It is defined as the negative vector sum of the transverse momenta of all reconstructed and calibrated physics objects, including electrons, muons, photons, and jets [233]. The magnitude of this vector, E_T^{miss} , is commonly used as an observable in searches for undetectable final-state particles.

In addition to the contributions from these so-called hard objects, a soft term is included in the calculation of $\vec{p}_T^{\text{miss}}$ to account for low-momentum particles that are not reconstructed as part of any identified final-state object. This soft term is constructed from charged-particle tracks associated with the primary vertex, providing robustness against pile-up contamination and improving the resolution of the E_T^{miss} measurement [234, 235].

Accurate reconstruction of E_T^{miss} is essential in analyses sensitive to non-interacting particles, such as neutrinos. In such cases, significant E_T^{miss} may be the only indication of their presence. The performance of the E_T^{miss} reconstruction is influenced by SFs, including the energy scale and resolution of the individual input objects, detector coverage and efficiency, and the quality of the track-based soft term. As such, the reliability of E_T^{miss} as a discriminating observable depends on careful calibration and systematic treatment of its associated uncertainties.

5.8 Overlap removal

To avoid double-counting of physics objects, overlaps between objects are resolved after their identification. Ambiguities between overlapping reconstructed objects are resolved using a harmonised procedure, in which pre-selected loose leptons and jets are considered in a fixed sequence of operations designed to prevent double-counting.

Any electron found with a track overlapping with any other electron is removed. Similarly, any calorimeter muon sharing a track with an electron is removed. Any electron sharing a track with a muon is also removed. Any jet found within a distance ΔR of 0.2 from an electron is removed. Any electron subsequently found within ΔR of 0.4 from a jet is removed. Any jet with fewer than three tracks associated with it, found within ΔR of 0.2 from a muon, is removed. Additionally, any jet with fewer than three tracks, having a muon ID track ghost-associated with it, is removed. Any muon subsequently found within ΔR of 0.4 from a jet is also removed.

Any τ_{had} found within a ΔR of 0.2 from an electron is removed. Furthermore, any τ_{had} found within a ΔR of 0.2 from any type of muon with p_T greater than 2 GeV is removed. If the τ_{had} p_T is greater than 50 GeV, it is only removed if it overlaps with a combined-type muon.

The definitions and selections described in this chapter form the foundation of the event reconstruction strategy used in this analysis. By applying consistent object definitions, quality criteria, and overlap removal procedures, the analysis ensures reliable

identification of final-state particles and accurate modelling of detector-level observables. These reconstructed objects serve as inputs to the event selection and signal extraction procedures discussed in the following chapters.

Chapter 6

From angular distributions to EFT: A fully differential analysis of single- top-quark decays

This chapter presents a detailed analysis of single top-quark events in the t -channel, utilising the full kinematic reconstruction of the decay chain through a four-angle differential approach explained in Section 2.4. The method exploits the full angular structure of the top-quark decay to maximise sensitivity to deviations from the SM arising from dimension-six operators in the SMEFT framework. This analysis represents the central research contribution of the present dissertation.

As already discussed in previous chapters, the analysis is based on pp collision data collected at a centre-of-mass energy of $\sqrt{s} = 13$ TeV, corresponding to an integrated luminosity of 140 fb^{-1} . The t -channel single-top-quark production mode offers a unique laboratory for probing the EW structure of the top-quark sector, particularly through the tWb vertex. Although some operators induce only subleading contributions in this channel, the angular observables employed here provide enhanced sensitivity by accessing interference terms and subtle distortions in the kinematic distributions.

This chapter presents the fully differential angular analysis of single-top-quark decays developed for extracting constraints on EFT coefficients. The chapter starts with the introduction in Section 6.1, common event reconstruction and selection in Section 6.2 followed by analysis regions definition in Section 6.3. In Section 6.4, event yields, kinematic and angular distributions are presented. Section 6.5 contains detailed information about the systematic uncertainties that analysis is sensitive to. In Section 6.7, the analysis method and statistical framework are presented. Finally, in Section 6.8 the fit results of the analysis are presented, followed by the discussion and conclusions.

6.1 Introduction

The analysis of the quadruple differential cross-section for single-top-quark events in the t -channel is a major step forward in understanding the angular distributions of top-quark decays. Earlier ATLAS analyses, discussed in Chapter 2 set the stage for this work. Single-angle studies, described in Section 2.4.2, examine the distributions of individual angles, demonstrating the sensitivity of angular observables to SM predictions. However, these studies could not fully capture how the angles interact. Two-angle analyses presented in Section 2.4.3 improved on this by looking at correlations between pairs of angles, and three-angle studies in Section 2.4.4 went further by exploring more complex relationships. Despite these advancements, previous studies did not fully utilise the complete angular information available. This thesis introduces a more comprehensive framework, based on Ref.[90] that captures the full phase space of the decay, leading to a more thorough and precise analysis.

The use of Fourier-dual space, described in Section 2.4.4, allows for a systematic description of the angular distributions. As explained in Section 2.4.1, expanding the decay dynamics into a set of independent components enables a systematic study of potential effects from new physics interactions. This approach relies on representing angular distributions in terms of M -function coefficients, labelled as $c_{m'm}^{j_1 j_2}$ (see Equation 2.21). At LO and parton level, these distributions can be described by 23 non-constant coefficients $c_{m'm}^{j_1 j_2}$, with $j_1 < 1$ and $j_2 < 2$. These coefficients, directly influenced by the Wilson coefficients, establish a connection to the properties of the tWb vertex.

In this analysis, the `dim6top`¹ model is used to generate MC simulations essential for SM and EFT physics analysis. The morphing procedure discussed in Section 4.5 serves as the ultimate power-up for analysis sensitivity, expanding the phase space of Wilson coefficients and making sure no important signal contribution gets missed.

6.2 Event reconstruction and selection

The reconstruction of final-state particles, such as the top quark and the W boson, is a critical step in this analysis. This process utilises the measured kinematic properties of detected objects, including leptons, jets, and E_T^{miss} , to infer the properties of the original

¹The `dim6top` model is a Universal FeynRules Output (UFO) model implementation of the SMEFT with a focus on dimension-six operators affecting top-quark interactions. It includes modifications to the tWb vertex, four-fermion interactions, and anomalous top-quark couplings, parametrised by Wilson coefficients.

particles produced in the event. Advanced algorithms are employed to reconstruct these particles, taking into account detector resolution, acceptance effects, and systematic uncertainties. These methods ensure a precise identification of particle candidates, which forms the basis for further analysis.

This section focuses on the event selection criteria and the categorisation of events based on the reconstructed particle properties.

6.2.1 Reconstruction of the top quark

Two main methods are utilised for reconstructing the kinematics of the top quark: the likelihood-based (LH) method and the analytical method. Each approach offers distinct advantages, and their comparison provides a deeper understanding of the reconstruction process.

In the analytical method, the W -boson and neutrino kinematics are reconstructed using kinematic equations under the assumption of an on-shell W boson. The key constraint is provided by momentum conservation:

$$(p_W)^2 = (p_l + p_\nu)^2 = m_W^2, \quad (6.1)$$

where p_w, p_l, p_ν represent the four-momentum of the W boson, the lepton and the neutrino, respectively, and m_W is the mass of the W boson. To account for the missing transverse energy, E_T^{miss} , the transverse momentum of the neutrino is assumed to match it:

$$\vec{p}_\nu = (E_T^{\text{miss}} \cos \phi_T^{\text{miss}}, E_T^{\text{miss}} \sin \phi_T^{\text{miss}}, p_z(\nu)), \quad (6.2)$$

where $\vec{p}_\nu$ is the three-momentum of the neutrino, ϕ_T^{miss} is the azimuthal angle of the missing transverse energy. The neutrino energy, E_ν , is then calculated using the energy-momentum relation:

$$E_\nu = \sqrt{(E_T^{\text{miss}})^2 + (p_z(\nu))^2}. \quad (6.3)$$

The longitudinal component $p_z(\nu)$ is determined by imposing the W -boson mass constraint:

$$m_W^2 = m_l^2 + 2E_l E_\nu - 2\vec{p}_l \cdot \vec{p}_\nu, \quad (6.4)$$

where $\vec{p}_l$ is the three-momentum of the charged lepton. This results in a quadratic equation for $p_z(\nu)$:

$$A(p_z(\nu))^2 + Bp_z(\nu) + C = 0, \quad (6.5)$$

with coefficients:

$$A = (E_l)^2 - (p_z(l))^2, \quad (6.6)$$

$$B = 2p_z(l) \left[m_l^2 - m_W^2 - 2E_T^{\text{miss}} (p_x(l) \cos \phi_T^{\text{miss}} + p_y(l) \sin \phi_T^{\text{miss}}) \right], \quad (6.7)$$

$$C = (E_l)^2 (m_W^2 - m_l^2) - \left[m_W^2 - 2E_T^{\text{miss}} (p_x(l) \cos \phi_T^{\text{miss}} + p_y(l) \sin \phi_T^{\text{miss}}) \right]^2. \quad (6.8)$$

The solutions for $p_z(\nu)$ are given by:

$$p_z(\nu) = \frac{-B \pm \sqrt{B^2 - 4AC}}{2A}. \quad (6.9)$$

When real solutions exist, the one with the smaller magnitude is chosen. If no real solution exists ($\Delta < 0$), the E_T^{miss} , is recalibrated iteratively until a valid numerical solution is obtained.

Once the neutrino momentum is reconstructed, the four-momentum of the W boson is calculated as:

$$p_W = p_l + p_\nu. \quad (6.10)$$

Subsequently, the top-quark momentum is reconstructed by combining the W -boson four-momentum with the b -jet momentum:

$$p_t = p_W + p_b. \quad (6.11)$$

On the other hand, the LH method improves upon the analytical approach by incorporating detector resolution effects and constraints into a global likelihood framework. Thanks to the KLfitter² package [236] a likelihood function that fits the kinematic

²The KLfitter (short for Kinematic Likelihood Fitter) is a toolkit used in high-energy physics analyses, especially within the ATLAS experiment, to reconstruct the kinematics of particle collisions in a way that maximises the likelihood of the observed final state originating from a specific theoretical process, commonly used in top-quark production.

variables is constructed in order to measured detector-level observables. The likelihood function is expressed as:

$$\ln \mathcal{L} = \ln B(m_W|M_W, \Gamma_W) + \ln B(m_t|M_t, \Gamma_t) + \sum_i \ln W(E_i^{\text{meas}}|E_i), \quad (6.12)$$

where:

- $B(m_W|M_W, \Gamma_W)$: Breit-Wigner distribution for the W -boson mass,
- $B(m_t|M_t, \Gamma_t)$: Breit-Wigner distribution for the top-quark mass,
- $W(E_i^{\text{meas}}|E_i)$: Transfer functions³ mapping true-particle energies to detector-level measurements.

The likelihood function accounts for experimental uncertainties through the transfer functions, which are modelled as double-Gaussian distributions. By maximising the likelihood, the LH method ensures the most probable kinematic reconstruction.

While the analytical method offers a straightforward approach by solving kinematic equations directly (used in previous analyses [84, 95, 96] as described in Section 2.3), it is limited by its reliance on idealised assumptions about E_T^{miss} and detector resolution. In contrast, the LH method provides a more robust reconstruction by incorporating detector effects and experimental constraints, reducing biases in reconstructed observables. Although computationally more demanding, the LH method is preferred for precision studies due to its superior accuracy and reliability, and therefore the method used in this analysis.

6.2.2 Signal selection

The signal candidates considered in this analysis are single-top-quarks produced in the t -channel, decaying into one lepton and one b -jet. The decay chain can be expressed as:

³The transfer functions are typically modelled as:

$$W(E_i^{\text{meas}}|E_i) = \frac{1}{\sqrt{2\pi}(\sigma_1 + p_3\sigma_2)} \left[e^{-\frac{(\Delta E - \mu_1)^2}{2\sigma_1^2}} + p_3 e^{-\frac{(\Delta E - \mu_2)^2}{2\sigma_2^2}} \right], \quad (6.13)$$

where:

$$\Delta E = \frac{E_i - E_i^{\text{meas}}}{E_i}, \quad (6.14)$$

where σ_1 and σ_2 are the widths of the Gaussian components, μ_1 and μ_2 are the means, and p_3 is the relative weight of the second Gaussian. These functions account for the detector resolution effects and are calibrated using MC simulations. By integrating these functions into the likelihood, the LH method accurately accounts for experimental distortions, enabling precise reconstruction of the kinematics of the top quark.

$$t \rightarrow Wb, \quad W \rightarrow \ell\nu, \quad \ell \in e, \mu, \tau.$$

For events where the W -boson decays into a τ -lepton, the analysis considers them if they decay leptonically, i.e. into an electron or a muon, which occur approximately 30% of the time. These secondary decays can be written as:

$$\tau \rightarrow e\nu_e\nu_\tau, \quad \tau \rightarrow \mu\nu_\mu\nu_\tau.$$

These decay chains produce a final state kinematically similar to direct W -boson decay into electron and muon. However, due to the presence of two additional neutrinos from the τ -lepton decay, reconstructing the W -boson kinematics becomes more challenging, leading to broader distributions and reduced precision in the measurements. Events in which the τ lepton decays hadronically are vetoed, as discussed in Section 5.5.

To ensure the integrity of the data used in this analysis, events are required to satisfy the GRL criteria and general quality requirements. These measures remove mis-reconstructed events and reject non-collision backgrounds.

6.2.3 Background estimation

In the analysis of single-top-quark t -channel events, several background processes produce final states that closely resemble the signal. Among these, the W +jets and $t\bar{t}$ processes are the most significant contributors, while additional contributions arise from non-prompt leptons (also named as fakes), which are treated separately.

As described in Section 4.4, the W +jets background arises from QCD processes where a W boson is produced in association with one or more hadronic jets. In this analysis, the relevant W -boson decay is the leptonic mode, via $W \rightarrow \ell\nu$, resulting in a charged lepton (electron or muon) and a neutrino that escapes detection, which contributes to E_T^{miss} . The accompanying jets originate from the radiation of quarks or gluons during the hard scatter or from ISR.

This background poses several challenges for analysis. The final-state topology of W +jets events can closely resemble the signal, especially in cases where one of the jets is misidentified as originating from a b -quark and passes b -tagging criteria. Furthermore, the production cross-section of W +jets is significantly larger than that of the t -channel

single-top-quark process, making it one of the most dominant sources of background contamination. Adding to the complexity, W +jets events often exhibit diverse jet multiplicities, which can overlap with signal and other backgrounds, complicating efforts to separate them based on jet-based selection criteria alone. To address these challenges, the W +jets background is estimated using dedicated control regions (CRs). These regions are designed to enhance the contribution of W +jets processes while simultaneously suppressing the signal and other backgrounds. This approach allows for precise modelling and normalisation of W +jets contributions, enabling accurate background subtraction in the signal region (SR).

The production of $t\bar{t}$ events occurs through strong interactions, primarily via gluon-gluon fusion as described in Section 2.1.2, and represents a major background in this analysis. In $t\bar{t}$ events, both top quarks decay via $t \rightarrow Wb$, resulting in two W bosons and two b -jets in the final state. The decay modes of the W bosons determine the classification of $t\bar{t}$ events: single-lepton (also known as l +jets), dilepton and all hadronic as explained in Section 2.1.2. The single-lepton channel poses the greatest challenge for this analysis, as its final state closely resembles the t -channel single-top-quark signal. While the higher jet multiplicity in $t\bar{t}$ events can sometimes provide a distinguishing feature, misclassification of jets or uncertainties in b -tagging can result in $t\bar{t}$ events entering the SR, complicating their separation from the signal.

The estimation of the $t\bar{t}$ background involves several challenges. The $t\bar{t}$ production cross-section is much larger than that of the t -channel single-top-quark process, making it a dominant source of background. The complex topology of $t\bar{t}$ events, with additional jets and b -jets, requires detailed modelling to ensure accurate discrimination from the signal. Furthermore, the efficiency of b -tagging and jet reconstruction is crucial for separating $t\bar{t}$ from the signal, and any miscalibration or misidentification in these processes can introduce biases into the background estimate. Accurate modelling and CR-based normalisation is therefore essential to address these challenges effectively.

For both W +jets and $t\bar{t}$ backgrounds, CRs are constructed to enhance these processes while suppressing signal-like events. These CRs are defined by imposing particular selection criteria, leveraging differences in kinematic distributions (e.g., jet multiplicities, invariant masses, E_T^{miss}) between the signal and backgrounds. Data-driven corrections or normalisation factors derived from these CRs are applied to the SR to ensure accurate background subtraction.

In this analysis, non-prompt backgrounds refer to processes where a jet is misidentified as an electron, or where a muon originates from the semi-leptonic decay of heavy-flavour quarks. These fakes primarily arise in QCD interactions, where the detector

misclassified hadronic activity or secondary decays as prompt leptons. Unlike prompt leptons, which are directly produced in EW decays of particles such as the W or Z boson, fake leptons originate from secondary processes or misidentifications, making their modelling significantly more complex. The estimation of non-prompt backgrounds using MC simulations presents several challenges. First, the precise simulation of detector effects required to model misidentification is inherently difficult. Small inaccuracies in these models can lead to large biases in the predicted fake rates. Additionally, QCD processes, which are the primary source of fakes, have an extremely large cross-section compared to the fake rate. This necessitates prohibitively large MC simulated samples to achieve statistically significant results. To address these challenges, this analysis adopts a data-driven method based on a matrix method. This approach is well-suited for estimating fake contributions in a robust and model-independent manner. The matrix method utilises a parametrisation of lepton misidentification efficiencies, referred to as fake efficiencies, which are derived from a multijet-enriched CR (see Section 6.3) where fakes are the dominant source of leptons. These efficiencies are then used to reweight events in the data that fail lepton identification criteria. By applying this reweighting on an event-by-event basis, the matrix method effectively estimates the contribution of fake leptons while minimising dependence on MC simulations. This strategy not only mitigates biases associated with detector misclassification but also leverages real data to accurately capture the behaviour of non-prompt backgrounds.

6.3 Region definition

The goal of this analysis is to measure EFT parameters by selecting single-top-quark t -channel events while minimising contributions from other processes, particularly $t\bar{t}$ and W +jets backgrounds. To achieve this, events are categorised into one SR and two CRs ($t\bar{t}$ CR and W +jets CR). These regions are designed to allow for the simultaneous estimation of background contributions and the measurement of EFT coefficients using a data-driven method.

The event selection process begins with a preselection phase aimed at reducing contributions from misidentified jets. Candidate events are required to have one charged tight lepton (electron or muon) with transverse momentum $p_T > 30$ GeV and pseudorapidity $|\eta| < 2.5$. Events with electrons and muons from leptonically decay taus are considered while events with τ_{had} are vetoed as already explained in Section 5.5. Events with additional secondary low- p_T charged loose leptons ($p_T > 10$ GeV, $|\eta| < 2.5$) are vetoed. Additionally, two jets are required in the final state, each with $p_T > 30$ GeV

and $|\eta| < 4.5$. For jets in the transition region between $2.75 < |\eta| < 3.5$, the transverse momentum threshold is raised to 35 GeV. At least one of these jets must be tightly b -tagged, passing the 60% WP as discussed in Section 5.6.1. This preselection also requires E_T^{miss} greater than 35 GeV and a transverse mass of the lepton- E_T^{miss} system, $m_T(\ell, E_T^{\text{miss}})$, exceeding 60 GeV. An additional lepton- p_T requirement is applied, where it must satisfy $p_T(\ell) > 50 \times \left(1 - \frac{\pi - |\phi(j_1, \ell)|}{\pi - 1}\right)$, with $\phi(j_1, \ell)$ being the azimuthal angle difference between the leading jet and the lepton.

Events that fail these criteria are enriched in non-prompt lepton events and are used to validate the estimation of fake contributions. Events passing the preselection are further categorised into three regions: the SR and two CRs, $t\bar{t}$ CR and W +jets CR (also referred to as CR1 and CR2, respectively). These regions allow for the isolation of signal events and the study of dominant background processes.

In the SR, additional requirements are applied to further isolate t -channel signal events. The invariant mass of the lepton- b -jet system, m_{lb} , must be less than 153 GeV in order to exclude events in which the top quark of the t -channel signal process is off-shell. The invariant mass of the lepton- E_T^{miss} - b -jet system, $m_{lE_T^{\text{miss}}b}$, must lie within the range [120.6, 234.6] GeV, to reject background events from processes not involving top quarks. A trapezoidal kinematic requirement is imposed on the pseudorapidity of the spectator jet (η_j) and the reconstructed top-quark pseudorapidity ($\eta_{lE_T^{\text{miss}}b}$) in order to reject background events, which have leptons in the forward region in the events with central reconstructed top-quarks. This requirement is defined as:

$$\begin{aligned} & \left(\eta_j < \left(a_1 \eta_{lE_T^{\text{miss}}b} + b_1 \right) \right) \cap \left(\eta_j > \left(a_1 \eta_{lE_T^{\text{miss}}b} - b_1 \right) \right) \cap \\ & \left(\eta_j > \left(a_2 \eta_{lE_T^{\text{miss}}b} + b_2 \right) \right) \cup \left(\eta_j < \left(a_2 \eta_{lE_T^{\text{miss}}b} - b_2 \right) \right), \end{aligned} \quad (6.15)$$

with optimised parameters $a_1 = 3$, $b_1 = 10.5$, $a_2 = 0.25$, and $b_2 = 2.5$. The illustration of the trapezoidal requirement is shown in Figure 6.1.

Furthermore, the invariant mass of the reconstructed top quark and spectator jet, $m_{jlE_T^{\text{miss}}b}$, must exceed 320 GeV to also reject background events from processes not involving top quarks. The sum of the energy deposited in the calorimeter, E_{calo} , must be greater than 190 GeV since the E_{calo} distributions of the backgrounds peak at lower values than the t -channel signature. Finally, the log-likelihood⁴ value from the KLFitter algorithm, $\ln \mathcal{L}_K$, must be larger than -36, in order to reject events not fully compatible

⁴The term “log-likelihood” refers specifically to the natural logarithm, denoted as $\ln$, i.e., the logarithm to the base e , of the likelihood function. This convention is standard in statistical inference due to the simplifying properties of the natural logarithm in analytical calculations, particularly with exponential families and when maximising likelihoods.

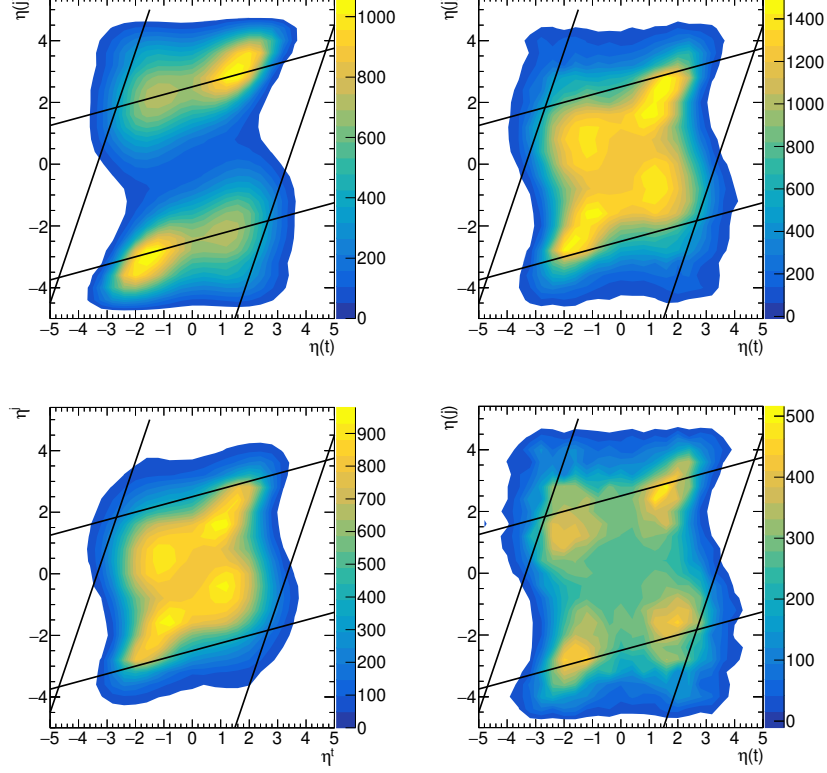


Figure 6.1: Distribution of the pseudorapidity of the top quark versus the pseudorapidity of the spectator jet for events with 1 b -jet and 1 light-jet that survived the pre-selection criteria. The upper left the plot shows the distribution for signal events and the upper right for the combination of all backgrounds. The bottom left shows the distribution for the $t\bar{t}$ background while the bottom left shows the distribution for W +jets. The black lines show the trapezoidal requirements as explained in the text.

with a physics process with a good reconstructed top-quark candidate. The distribution of $\ln \mathcal{L}_K$ is shown in Figure 6.2.

The first CR is enriched in $t\bar{t}$ events and is defined by requiring two tightly b -tagged jets that pass the 60% WP, in addition to the preselection criteria. The second CR, enriched in W +jets events, includes events that fail at least one SR requirement and contain only one tightly b -tagged jet, passing the 60% b -tagging WP.

The event selection is designed to isolate well-modelled single-top-quark t -channel events and suppress backgrounds. The presence of a forward untagged jet exploits the spectator-quark topology specific to t -channel production, while requirements on η and event topology suppress poorly modelled or non-resonant backgrounds such as W +jets and multijet events.

The event counts in $t\bar{t}$ CR and W +jets CR are used as sensitive variables in a

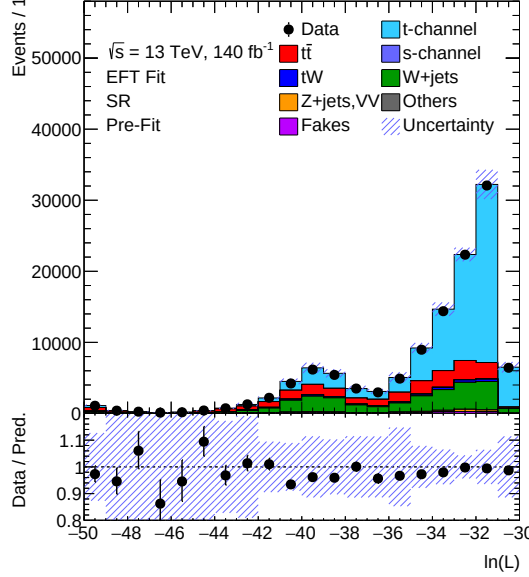


Figure 6.2: Distribution of the log-likelihood in the signal region. The prediction is compared to data, shown as the black points with statistical uncertainties. No overall normalisation scale factors are considered. The uncertainty bands correspond to the uncertainties due to the size of the simulated event samples. The lower plots show the ratio of data to prediction in each bin.

simultaneous fit to estimate the yields of $t\bar{t}$ and W +jets events in a data-driven approach, alongside measuring the EFT coefficients. Selection criteria for the preselection, SR, and CRs are summarised in Table 6.1. These stringent selection steps ensure high purity of the SR, enabling accurate measurements of EFT parameters.

6.4 Event yields, and kinematic and angular distributions

This section displays the preselection and SR, as well as the two CRs, CR1 and CR2, expected and observed event yields, and kinematic and angular distributions. Events must fulfil strict requirements to enter the preselection area, which include an isolated lepton, two jets, and one b -tagged jet. As already mentioned the SR is then specified by additional selection criteria that guarantee the presence of a top quark and a W boson. Two CRs are established in order to evaluate the reliability of the primary backgrounds, namely the $t\bar{t}$ and W +jets processes. The $t\bar{t}$ -enriched CR (CR1) applies the same criteria as the SR but requires two b -tagged jets, whereas the W +jets-enriched CR (CR2) consists of events that fail at least one of the SR selections.

Preselection region	SR	$t\bar{t}$ CR	W +jets CR
	Exactly 1 charged tight lepton (electron and muon) ($p_T > 30$ GeV and $ \eta < 2.5$)		
	Exactly zero τ_{had} ($p_T > 10$ GeV and $ \eta < 2.5$)		
	Exactly 2 jets ($p_T > 30$ GeV and $ \eta < 4.5$; $p_T > 35$ GeV within $2.7 < \eta < 3.5$)		
	At least one of the jets has to be b -tagged ($ \eta < 2.5$; 60% WP)		
	$E_T^{\text{miss}} > 35$ GeV		
	$m_T(lE_T^{\text{miss}}) > 60$ GeV		
	$p_T(\ell) > 50 \left(1 - \frac{\pi - \Delta\phi(j_1, \ell) }{\pi - 1} \right)$ GeV		
	Exactly 1 b -jet ($ \eta < 2.5$; 60%WP)	Exactly 2 b -jet ($ \eta < 2.5$; 60%WP)	Exactly 1 b -jet ($ \eta < 2.5$; 60%WP)
	$m_b < 153$ GeV $m_{lE_T^{\text{miss}b}} \in [120.6, 234.6]$ GeV trapez. requirement $m_{jlE_T^{\text{miss}}} > 320$ GeV $E_{\text{calo}} > 190$ GeV		Fail any selection of the SR
	$\ln \mathcal{L}_K > -36$		

Table 6.1: Summary of the selection criteria for defining the preselection, SR and the two CRs, CR1 and CR2.

6.4.1 Event yields

Table 6.2 presents the total expected and observed event yields for several regions, including preselection and SR, as well as $t\bar{t}$ and W +jets CRs, i.e. CR1 and CR2, respectively. This table pertains to the combined electron and muon channels, and the uncertainties reported are purely statistical. In addition, the table also includes the signal-to-background ratio (S/B) and data-to-expected-MC-events ratio (Data/Pred.) for each region. Good agreement between data and prediction is found within the statistical uncertainty in all regions.

6.4.2 Kinematic distributions in the signal region

This section of the study provides a presentation of the kinematic distributions in the SR. The distributions presented include both the electron and muon channels, and also combine the top quarks and antiquarks. All predicted distributions are scaled according to the theoretical cross-sections, while the multijet distribution is normalised based on the data-driven approach as explained in Section 4.4.3. The prediction uncertainty bands displayed in the kinematic distributions correspond to the uncertainties resulting from the limited size of the simulated event samples, which are added in quadrature.

A detailed overview of kinematic observables in the SR is presented in Figures 6.3

Process	Preselection region	SR	W +jets CR	$t\bar{t}$ CR
t -channel	262000 ± 5000	61300 ± 1700	176000 ± 4000	15800 ± 700
s -channel	13670 ± 230	210 ± 10	9280 ± 130	4110 ± 130
tW	99000 ± 4000	1390 ± 40	92300 ± 3300	4400 ± 210
$t\bar{t}$	770000 ± 40000	10200 ± 500	603000 ± 35000	144000 ± 8000
W +jets	609000 ± 31000	14600 ± 1400	565000 ± 27000	20100 ± 900
Z +jets, VV	58600 ± 5100	1510 ± 160	52800 ± 4300	3100 ± 400
others	28600 ± 800	870 ± 30	25900 ± 600	1500 ± 200
Total Exp.	1840000 ± 50000	90100 ± 2500	1520000 ± 40000	193000 ± 9000
Data	1866198	89565	1547678	199970
S/B	0.166 ± 0.006	2.129 ± 0.149	0.131 ± 0.005	0.0892 ± 0.006
Data/Pred.	1.014 ± 0.028	0.994 ± 0.028	1.018 ± 0.027	1.036 ± 0.048

Table 6.2: The event yields in the preselection and SR, as well as in the $t\bar{t}$ control and W +jets CRs. The event yield predictions are obtained using simulated event samples along with their theoretical cross-sections, except for the multijet prediction which is estimated from a data-driven likelihood fit as explained in Section 4.4.3. These event yields are computed without considering any overall normalisation scale factors. The uncertainties displayed in the results are solely statistical. The individual predictions are rounded to integers, whereas the "Total Exp." represents the sum of full-precision individual predictions, rounded to the nearest integer. The table also includes the expected S/B and Data/Pred. ratios.

to 6.8. The spectator jet kinematics—its transverse momentum, pseudorapidity, and azimuthal angle—are shown in Figure 6.3. These variables are essential in characterising the forward nature of the light jet, which is a key feature of t -channel single-top-quark production. Figure 6.4 displays the corresponding distributions for the b -tagged jet. Since the b -jet arises from the decay of the top quark, its properties help constrain the top-quark reconstruction and distinguish signal from background, particularly W +jets. The lepton kinematics are presented in Figure 6.5, as the charged lepton is a clean signature of the leptonic W -boson decay. The angular separations $\Delta\phi(j_1, \ell)$, $\Delta\phi(\ell, E_T^{\text{miss}})$, and $\Delta\phi(W, j_b)$ shown in Figure 6.6 provide sensitivity to the event topology and help differentiate between background processes with different angular correlations. Figure 6.7 includes mass-related observables such as $m_{j\ell E_T^{\text{miss}}b}$ and m_{lb} , which are useful for validating the reconstruction of the top quark and probing its decay products. Additionally, $\Delta R(\ell, j)$ serves as an angular separation variable sensitive to event geometry and reconstruction

quality. Finally, Figure 6.8 shows the scalar sum of transverse momenta (H_T), E_T^{miss} , and the m_{jb} invariant mass, which together offer a global view of event energy balance and further enhance discrimination between signal and background processes. Across all distributions, predicted templates are compared with observed data, and the ratio panels provide a quantitative assessment of agreement within statistical uncertainties. There are some disagreements between data and prediction in some bins of these distributions but these are covered by considering the systematics, included in Section 6.5.

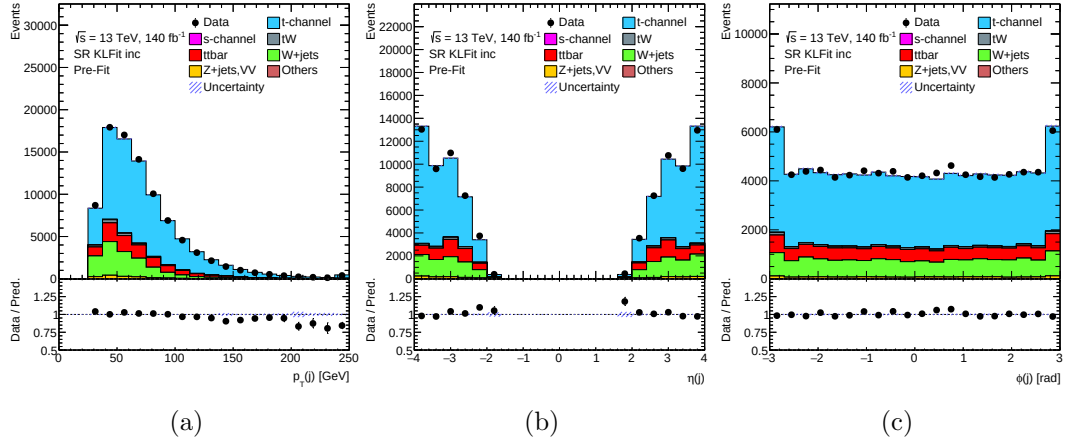


Figure 6.3: Distributions in the SR of (a) spectator jet p_T , (b) spectator jet η , and (c) b -jet ϕ . The predicted distributions are compared with the observed data, depicted as black points with statistical uncertainties. No overall normalisation scale factors are applied to the data or the predictions. The uncertainty bands reflect the statistical uncertainties of the predicted distributions due to the limited size of the simulated event samples. The first and last bins include underflow and overflow events, respectively. The lower panels of each distribution show the ratio of the observed data to the predicted distributions in each bin, offering a quantitative measure of the agreement between the prediction and expectation.

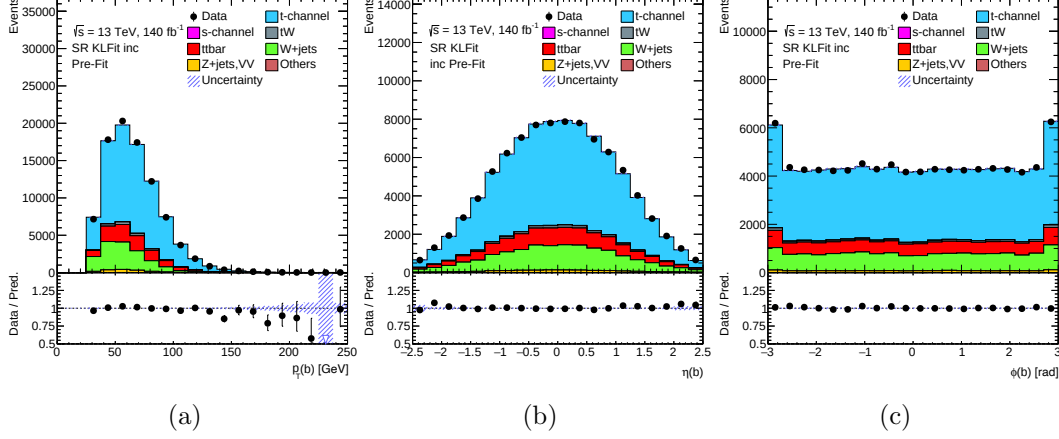


Figure 6.4: Distributions in the SR of (a) b -jet p_T , (b) b -jet η , and (c) b -jet ϕ . The predicted distributions are compared with the observed data, depicted as black points with statistical uncertainties. No overall normalisation scale factors are applied to the data or the predictions. The uncertainty bands reflect the statistical uncertainties of the predicted distributions due to the limited size of the simulated event samples. The first and last bins include underflow and overflow events, respectively. The lower panels of each distribution show the ratio of the observed data to the predicted distributions in each bin, offering a quantitative measure of the agreement between the prediction and expectation.

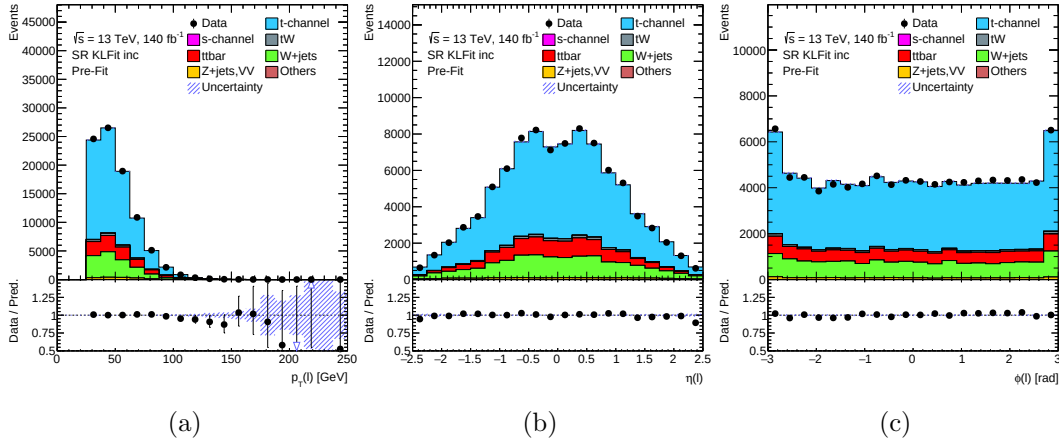


Figure 6.5: Distributions in the SR of (a) lepton p_T , (b) lepton η , and (c) lepton ϕ . The predicted distributions are compared with the observed data, depicted as black points with statistical uncertainties. No overall normalisation scale factors are applied to the data or the predictions. The uncertainty bands reflect the statistical uncertainties of the predicted distributions due to the limited size of the simulated event samples. The first and last bins include underflow and overflow events, respectively. The lower panels of each distribution show the ratio of the observed data to the predicted distributions in each bin, offering a quantitative measure of the agreement between the prediction and expectation.

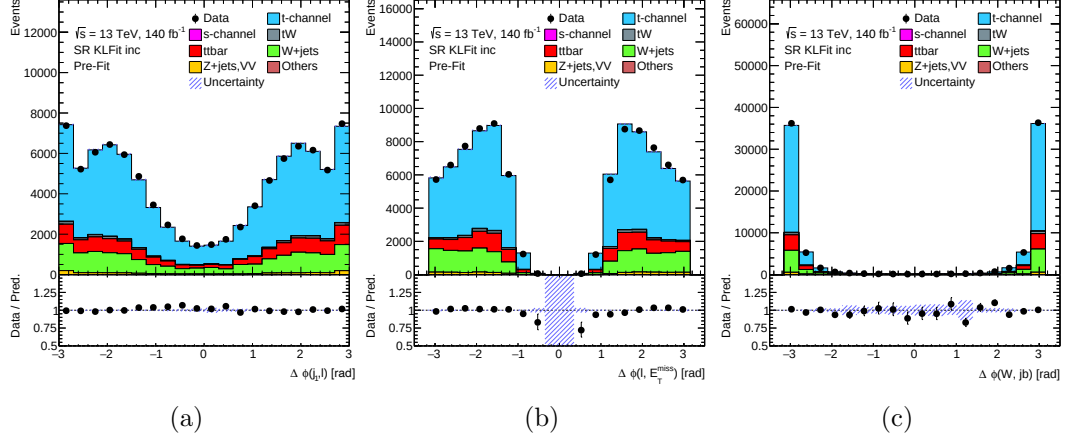


Figure 6.6: Distributions in the SR of (a) $\Delta\phi(j_1, \ell)$, (b) $\Delta\phi(\ell, E_T^{\text{miss}})$, and (c) $\Delta\phi(W, j_b)$. The predicted distributions are compared with the observed data, depicted as black points with statistical uncertainties. No overall normalisation scale factors are applied to the data or the predictions. The uncertainty bands reflect the statistical uncertainties of the predicted distributions due to the limited size of the simulated event samples. The first and last bins include underflow and overflow events, respectively. The lower panels of each distribution show the ratio of the observed data to the predicted distributions in each bin, offering a quantitative measure of the agreement between the prediction and expectation.

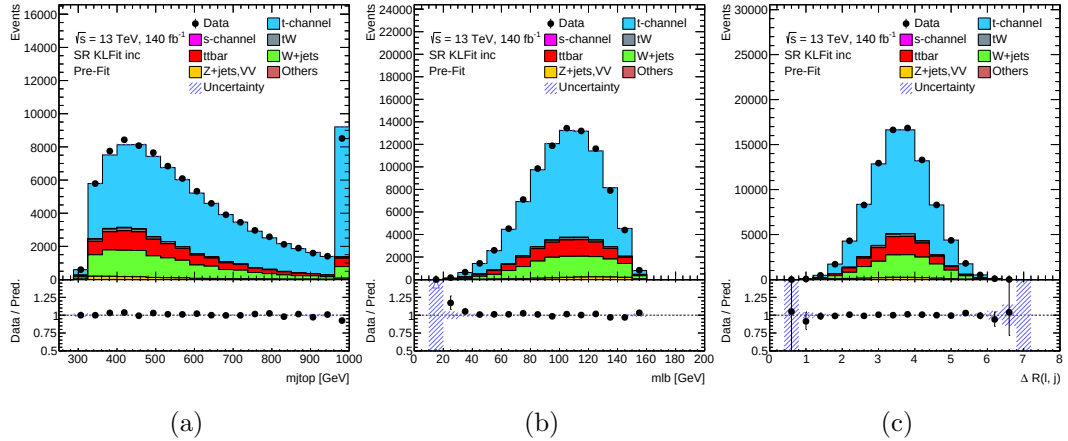


Figure 6.7: Distributions in the SR of (a) $m_{j\ell E_T^{\text{miss}} b}$, (b) m_{lb} , and (c) $\Delta R(\ell, j)$. The predicted distributions are compared with the observed data, depicted as black points with statistical uncertainties. No overall normalisation scale factors are applied to the data or the predictions. The uncertainty bands reflect the statistical uncertainties of the predicted distributions due to the limited size of the simulated event samples. The first and last bins include underflow and overflow events, respectively. The lower panels of each distribution show the ratio of the observed data to the predicted distributions in each bin, offering a quantitative measure of the agreement between the prediction and expectation.

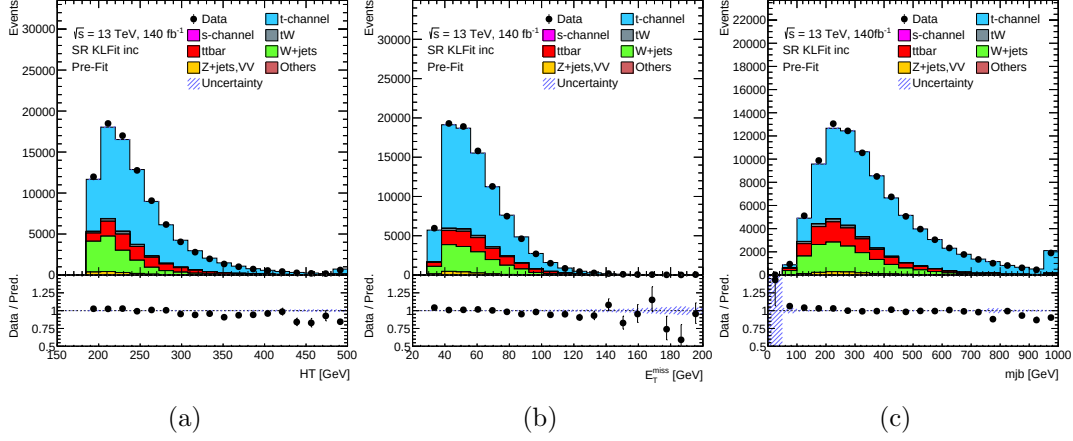


Figure 6.8: Distributions in the SR of (a) H_T , (b) E_T^{miss} , and (c) m_{jb} . The predicted distributions are compared with the observed data, depicted as black points with statistical uncertainties. No overall normalisation scale factors are applied to the data or the predictions. The uncertainty bands reflect the statistical uncertainties of the predicted distributions due to the limited size of the simulated event samples. The first and last bins include underflow and overflow events, respectively. The lower panels of each distribution show the ratio of the observed data to the predicted distributions in each bin, offering a quantitative measure of the agreement between the prediction and expectation.

6.4.3 Angular distributions in the signal and control regions

The angular distributions of $\cos\theta$, $\cos\theta^*$, ϕ , and ϕ^* , previously described in Chapter 2, Section 2.4.1, are presented in Figures 6.9, 6.10 and 6.11 for the SR and two CRs, respectively. These distributions are shown before any statistical fitting is performed (that is why the pre-fit label is shown), allowing for a direct comparison between the expected and observed distributions. The angular variables $\cos\theta$, $\cos\theta^*$, ϕ , and ϕ^* , shown in Figures 6.9–6.11, play a central role in the analysis as they form the basis of the quadruple-differential decay structure used to extract EFT-sensitive observables as discussed in Section 2.4. These angles are defined in the rest frames of the top quark and the W boson, and their distributions are sensitive to the spin correlations and production dynamics of the single-top-quark process. Figure 6.9 presents these angular distributions in the SR, while Figures 6.10 and 6.11 show them in the W +jets and $t\bar{t}$ CRs, respectively. Comparing the shapes of these distributions across regions highlights how different event selections affect the kinematic acceptance and composition of each sample. For instance, the W +jets region tends to show flatter angular distributions due to the absence of a resonant top quark, whereas the $t\bar{t}$ region exhibits distinct modulations from the decay of two correlated top quarks. In contrast, the SR displays angular shapes consistent with a single-top-quark decay topology and characteristic of a forward light-jet production.

These variations are critical for validating the modelling of the angular observables under different background-dominated conditions. Moreover, as these angles are directly used in the quadruple-differential likelihood fit, understanding their behaviour across regions supports the robustness of the EFT interpretation. For completeness, angular distributions split for top quarks and top antiquarks in the SR are also provided in Appendix G.

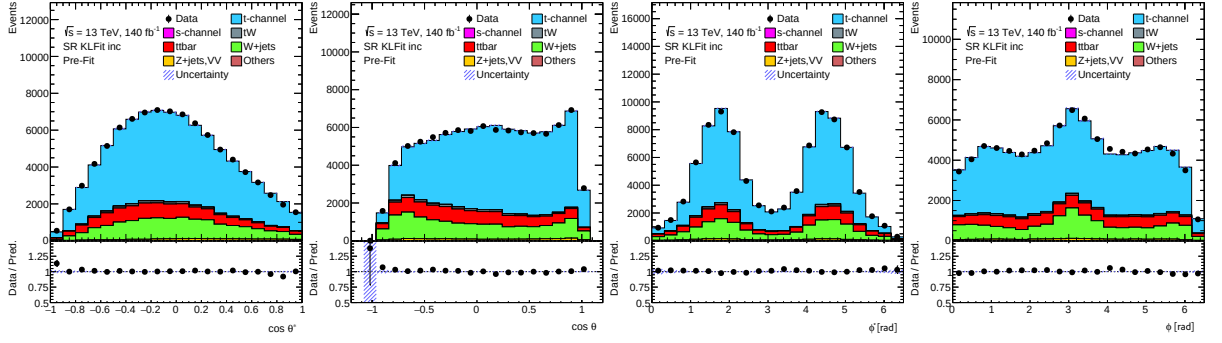


Figure 6.9: Distributions of four angles, namely $\cos \theta$, $\cos \theta^*$, ϕ , and ϕ^* in the SR inclusive channel. The predicted distributions are compared with the observed data, depicted as black points with statistical uncertainties. No overall normalisation scale factors are applied to the data or the predictions. The uncertainty bands reflect the statistical uncertainties of the predicted distributions due to the limited size of the simulated event samples. The first and last bins include underflow and overflow events, respectively. The lower panels of each distribution show the ratio of the observed data to the predicted distributions in each bin, offering a quantitative measure of the agreement between the prediction and expectation.

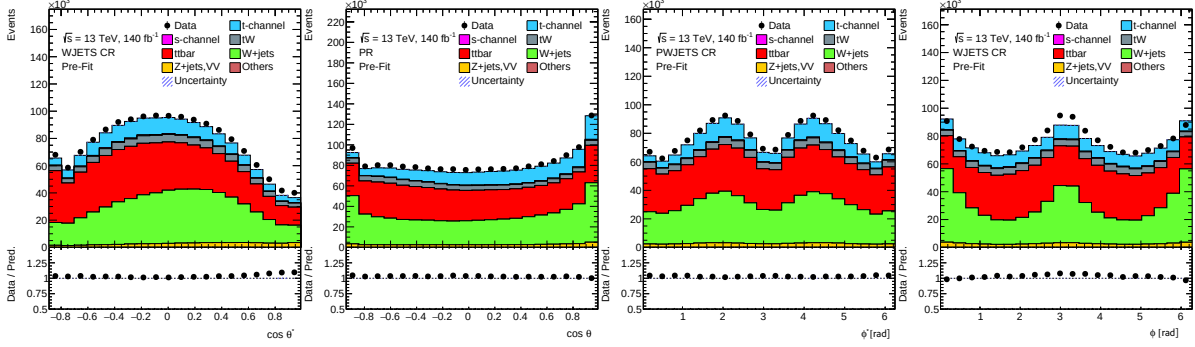


Figure 6.10: distributions of four angles, namely $\cos \theta$, $\cos \theta^*$, ϕ , and ϕ^* in the $W + \text{jets}$ CR. The predicted distributions are compared with the observed data, depicted as black points with statistical uncertainties. No overall normalisation scale factors are applied to the data or the predictions. The uncertainty bands reflect the statistical uncertainties of the predicted distributions due to the limited size of the simulated event samples. The first and last bins include underflow and overflow events, respectively. The lower panels of each distribution show the ratio of the observed data to the predicted distributions in each bin, offering a quantitative measure of the agreement between the prediction and expectation.

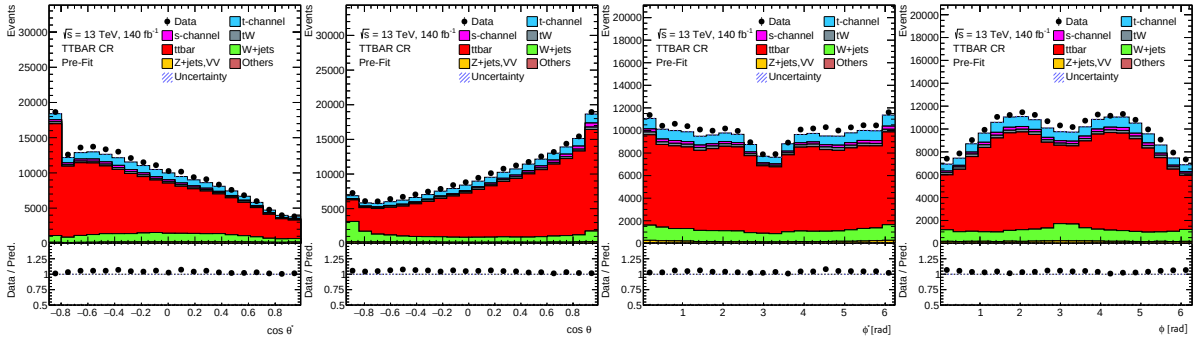


Figure 6.11: distributions of four angles, namely $\cos \theta$, $\cos \theta^*$, ϕ , and ϕ^* in the $t\bar{t}$ CR. The predicted distributions are compared with the observed data, depicted as black points with statistical uncertainties. No overall normalisation scale factors are applied to the data or the predictions. The uncertainty bands reflect the statistical uncertainties of the predicted distributions due to the limited size of the simulated event samples. The first and last bins include underflow and overflow events, respectively. The lower panels of each distribution show the ratio of the observed data to the predicted distributions in each bin, offering a quantitative measure of the agreement between the prediction and expectation.

6.5 Sources of systematic uncertainties

Sources of systematic uncertainty are a critical component of any high-precision measurement in particle physics, particularly in analyses like this one that rely on multidimensional differential distributions and detailed modelling of both signal and background processes. These uncertainties encapsulate our limited knowledge of detector performance, theoretical predictions, and MC modelling, and can significantly affect the extracted physics results if not properly accounted for. In general, systematic uncertainties impact either the overall event rate (normalisation) or the shape of observable distributions—such as kinematic or angular variables—which are central to this analysis.

In this analysis, systematic uncertainties are categorised into two main classes: experimental uncertainties, which are primarily detector-related and typically process-independent, and modelling uncertainties, which originate from theoretical assumptions or generator-specific configurations and are generally process-dependent. Each source of uncertainty is introduced into the statistical model through a dedicated nuisance parameter (NP), denoted as θ , and is treated through profiling in the likelihood fit. This approach allows the data itself to constrain the uncertainty within its statistical sensitivity.

Experimental uncertainties include contributions from the Jet Energy Scale (JES), Jet Energy Resolution (JER), pile-up, object identification efficiencies, and $E_{\text{T}}^{\text{miss}}$ among others. These are evaluated independently and applied across all relevant processes. Modelling uncertainties, on the other hand, are assessed by comparing different MC generators, varying renormalisation and factorisation scales, or altering PS and hadronisation settings. These are studied individually for major processes such as t -channel, $t\bar{t}$, W +jets, and tW .

To allow a consistent interpretation of the fitted NPs in the likelihood model, systematic variations are symmetrised. For uncertainties with both "up" and "down" variations, the symmetrised shift is calculated as (up – down)/2 and mirrored around the nominal prediction. For one-sided systematics—typically from generator comparisons—the variation is mirrored to provide a balanced treatment. This standardised approach simplifies the interpretation of NP pulls and constraints in the fit.

Modelling uncertainties derived from two-point comparisons (e.g., using alternative generators or different PS schemes) often show statistical fluctuations that do not reflect real differences in physics modelling. To address this, a bootstrap-based pruning method is explored to remove statistically insignificant variations. However, it is not used in the final results to avoid underestimating uncertainties. Instead, these systematics are

decorrelated between shape and normalisation components, and across charge channels when appropriate, to retain sensitivity and avoid artificial constraints.

In EFT-based measurements like this one, which depend on subtle distortions in multidimensional angular distributions, a careful and comprehensive treatment of systematics is essential. Any bias in the modelling of these effects can lead to incorrect constraints on new physics parameters. Therefore, the systematic treatment adopted in this analysis aims to be both conservative and robust, enabling reliable extraction of EFT coefficients from the data.

6.5.1 Sources of experimental systematic uncertainties

Experimental uncertainties are an essential aspect of particle physics analyses, arising from limitations in the detector response, calibration procedures, and object reconstruction efficiencies. In ATLAS, these uncertainties are primarily detector-related and impact how well physical objects are measured and identified. This section outlines the experimental uncertainties considered in this analysis, which stem from multiple sources and are evaluated using well-established procedures to ensure the reliability of the measurements.

- The integrated **luminosity** overall uncertainty is $\pm 0.83\%$ [237], and is applied as a single rate uncertainty to all signal and background processes that are scaled by luminosity in the fit. This uncertainty reflects the precision with which the total amount of data collected by the ATLAS detector is known (see Section 3.1.1). Although the luminosity is measured independently for each data-taking year, it is treated as a single effective uncertainty in this analysis. The uncertainty is propagated by scaling the expected yields of all simulated samples accordingly, ensuring consistency with the observed data.
- The **pile-up** systematic uncertainty accounts for mismodelling of the pile-up conditions in simulation (more about pile-up is given in Section 3.1.2). Two components are included: (1) an uncertainty in the reweighting procedure, evaluated by varying the average number of interactions per bunch crossing by $\pm 4\%$ around the nominal scaling factor of $1/1.03$ [238]; and (2) pile-up-related effects on the JES, captured by four JES NPs. These uncertainties are symmetrised and included as independent NPs in the fit.
- The **JES** systematic uncertainty reflects our limited knowledge of how accurately the energy of reconstructed jets matches the true particle-level energy. It arises from multiple sources, including detector calibration, modelling of jet fragmentation and hadronisation, and the pile-up environment. These uncertainties are derived using a combination of test-beam studies, LHC collision data, and simulation.

In this analysis, JES uncertainties are treated using the CategoryReduction scheme [239], which organises correlated components into a reduced set of 31 statistically independent NPs. These cover statistical and detector effects, modelling uncertainties, flavour-dependent responses, η inter-calibration, and extreme kinematic regions such as high jet p_T . Each JES component is symmetrised and treated as an independent variation in the fit, allowing the analysis to account for how changes in jet calibration affect both event selection and the shapes of key observables.

- The **JER** uncertainty reflects the finite precision with which the energy of jets is reconstructed in the detector. It is evaluated using dijet events, where two high- p_T jets are produced back-to-back, allowing the JER to be inferred from the width of the asymmetry distribution between the two jets as a function of p_T and η [240]. To reduce the number of uncertainty components, an eigenvector decomposition is performed, resulting in 12 principal NPs. An additional NP is included to account for residual differences between data and simulation, giving a total of 13 NPs. This approach, referred to as the “FullJER” scheme [239], is used in the fit with all variations symmetrised.
- **Jet Vertex Tagger (JVT) and Forward JVT (fJVT):** The JVT and fJVT uncertainties account for residual contamination from pile-up jets that survive pile-up suppression techniques. These algorithms are used to discriminate jets originating from the hard-scatter interaction against those from additional pp collisions. The uncertainty is evaluated using the `JetJvtEfficiency` tool [241], which provides variations in the tagging efficiency based on comparisons between data and simulation. The JVT is applied in the central region ($|\eta| < 2.5$), while fJVT targets the forward region ($2.5 < |\eta| < 4.5$). The associated uncertainties are included as NPs in the fit and are symmetrised.
- **Flavour tagging:** Flavour-tagging uncertainties arise from differences between data and simulation in the efficiency of identifying b -, c -, and light-flavour jets. These efficiencies are measured in CRs using data-driven methods: dileptonic $t\bar{t}$ events for b -jets, W -boson decays for c -jets, and dijet events for light jets. From these measurements, scale factors and their uncertainties are derived and applied to simulated events using the `BTaggingEfficiencyTool` [242]. The uncertainties are parameterised as a function of jet p_T and η , with a total of 45 NPs for b -jets, and 20 each for c - and light-flavour jets. Additional NPs are included to cover extrapolations to high- p_T regions. These are symmetrised and treated independently in the fit.
- **Electrons:** Uncertainties in electron reconstruction, identification, isolation, and

trigger efficiency are derived from $Z \rightarrow e^+e^-$ and $J/\psi \rightarrow e^+e^-$ events using the tag-and-probe method. Corrections for energy scale and resolution are obtained from comparisons between data and MC using Z -boson peak distributions and E/p measurements in $W \rightarrow e\nu$ events. A simplified correlation model is used to reduce the number of NPs, assuming full correlation in η , resulting in a total of three independent NPs [243]. These account for differences in detector response and calibration and are symmetrised before inclusion in the fit.

- **Muons:** Muon-related uncertainties include efficiency corrections for identification, isolation, and triggering, derived from tag-and-probe studies in $Z \rightarrow \mu^+\mu^-$ and $J/\psi \rightarrow \mu^+\mu^-$ events [244]. Additional uncertainties arise from the track-to-vertex association efficiency and are applied in the same manner. Momentum scale and resolution corrections are also applied, based on studies of muon tracks in the ID and MS. A separate uncertainty accounts for the charge-dependent sagitta bias observed in data. In total, seven NPs are included to cover all muon-related systematics, each symmetrised and independently varied in the fit.
- **Taus:** Since this analysis vetoes events containing hadronically decaying taus, uncertainties related to the tau energy scale are considered to account for potential mismodelling in tau identification and kinematics. These are derived from tag-and-probe studies of μ - τ_{had} pairs, using χ^2 fits to the visible mass in data, and are further cross-checked with variations in E/p and detector simulation configurations using GEANT4 [245]. Only tau energy scale uncertainties are included, as tau identification systematics have negligible effect under the veto strategy. These are treated as symmetrised NPs in the fit.
- The $E_{\text{T}}^{\text{miss}}$ includes a component from soft tracks not associated with any reconstructed object, known as the “soft term”. Uncertainties in this term arise from imperfect modelling of low- p_{T} activity and detector response in unclustered regions. They are evaluated by studying the p_{T} balance between the hard and soft terms in data and simulation [233]. Separate NPs are defined for the scale and resolution of the soft term. For each systematic variation, $E_{\text{T}}^{\text{miss}}$ is recalculated event-by-event to properly propagate the uncertainty through to selection and observable distributions.

6.5.2 Source of modelling systematic uncertainties

Sources of modelling uncertainty arise from theoretical assumptions, MC generator choices, and configuration settings used in the simulation of both signal and background processes. These uncertainties are evaluated by either comparing nominal predictions to alternative

MC samples, varying generator parameters such as QCD renormalisation (μ_R) and factorisation scales (μ_F) or PS settings, or by assigning fixed uncertainties to theoretical cross-sections. In each case, the differences between nominal and varied predictions are computed bin-by-bin and symmetrised to define systematic variations. These are then implemented as NPs affecting either the shape or normalisation of the distributions in the likelihood fit. All nominal and alternative samples used in this analysis are generated using the AFII framework to ensure consistency in detector modelling.

- Single-top-quark t -channel modelling:** The modelling of the t -channel process is tested using several variations that probe different aspects of theoretical uncertainty. The (μ_R) and (μ_F) scales are varied independently by factors of 0.5 and 2, both in the hard-scatter and in the PS stages, to evaluate ISR effects. The PS scale is further tested by varying the Var3c eigentune of the A14 tune [196]. The FSR modelling is assessed by modifying the renormalisation scale used in FSR. Additional comparisons are made between nominal POWHEG BOX +PYTHIA 8 and simulated samples generated with POWHEG:pThard=1 to probe matrix element and shower matching, and between PYTHIA 8 and HERWIG 7.1.3 to test hadronisation. The impact of the top-quark mass is evaluated by comparing samples with $m_t = 172$ GeV and $m_t = 173$ GeV. Since these variations are based on two-point comparisons, which are sensitive to statistical fluctuations, a bootstrap study is performed to assess bin-level consistency, although no pruning is ultimately applied. All variations are symmetrised and treated as uncorrelated between shape and normalisation components in the fit.
- $t\bar{t}$ modelling:** The modelling of $t\bar{t}$ production is probed using similar variations as the t -channel process. Scale uncertainties are introduced by varying (μ_R) and (μ_F) in the hard-scatter and PS. The shower scale is also modified through the A14 Var3c eigentune [196]. Matching uncertainties are assessed via POWHEG:pThard variations, and hadronisation is tested by comparing PYTHIA 8 with HERWIG 7.1.3. The h_{damp} parameter is varied from $1.5 \cdot m_t$ to $3 \cdot m_t$ to study its effect on high- p_T emissions. An additional uncertainty from the gluon recoil scheme in $t \rightarrow Wb$ decays is introduced by comparing the `recoilToColor` and `recoilToTop` settings. Finally, the top-quark p_T modelling is corrected to match NNLO+EW predictions, and the full difference between corrected and uncorrected samples is taken as an uncertainty [246]. As with t -channel, shape and normalisation components are treated separately and decorrelated in the fit. The overall theoretical uncertainty is put at 4%.
- W +jets modelling:** Uncertainties in W +jets production are assessed by varying the μ_R and μ_F scales in the ME and PS by factors of 0.5 and 2. Additionally, an

uncertainty in the jet-flavour composition is introduced by varying the fractions of b - and c -jets by $\pm 10\%$ while keeping the total cross-section constant. These variations account for potential mismodelling in Drell–Yan-like W +jets events and are symmetrised before being included as shape uncertainties in the fit.

- **tW modelling:** The tW single-top-quark process uses the same modelling variations as t -channel, including scale variations $((\mu_R), (\mu_F))$, the Var3c eigentune of the A14 [196] variation for the shower scale, and FSR uncertainties. Although this process is subdominant in the analysis, the corresponding uncertainties are included for completeness. The overall theoretical uncertainty is put at 3.5%.
- **Z +jets modelling:** A conservative 20% uncertainty is applied to the Z +jets background cross-section to account for theoretical uncertainties in its modelling from SHERPA. No dedicated alternative samples are used in this case.
- **Parton distribution functions (PDFs):** PDF uncertainties reflect the limited knowledge of parton momentum distributions inside the proton. They are evaluated using the PDF4LHC15 set, consisting of 30 eigenvector variations [214], applied via event-by-event reweighting of the nominal POWHEG+PYTHIA8 samples. The full envelope is treated as a correlated uncertainty and included in the likelihood fit.
- **Multijet background:** The uncertainty in the estimation of fake leptons is derived by comparing the prediction from a data-driven matrix method (see Section 4.4.3). The full difference is taken as the 1σ uncertainty and scaled by a factor of two to conservatively cover possible mismodelling in the fake rate. This is implemented as a normalisation uncertainty in the fit.
- **Minor backgrounds:** To account for theoretical uncertainties in the normalisation of background processes, overall cross-section variations are assigned as systematic uncertainties. For the $t\bar{t}H$ process, QCD scale and PDF uncertainties are considered at $+5.8\%/-9.2\%$ and $\pm 3.6\%$, respectively. The $t\bar{t}W$ process receives $+12.9\%/-11.5\%$ for QCD scale and $\pm 3.4\%$ for PDF uncertainties, while $t\bar{t}Z$ and tZq are both assigned $+9.6\%/-11.3\%$ from scale and $\pm 4\%$ from PDF variations. A conservative $\pm 50\%$ uncertainty is applied to the tWZ process to account for limited theoretical precision. Diboson backgrounds (WW, WZ, ZZ) are assigned a $\pm 30\%$ uncertainty, and rare processes such as $t\bar{t}t$, $t\bar{t}t\bar{t}$, other Higgs modes, tWH , and tHq are collectively assigned a conservative $\pm 50\%$ uncertainty due to their limited modelling accuracy.

6.6 Analysis method

The primary aim of this analysis is to place constraints on the Wilson coefficients, denoted as $\vec{C}$, which parametrise deviations from the SM within the EFT framework.

As explained in Chapter 2, EFT offers a systematic way to probe potential new-physics effects at energy scales beyond the direct reach of the LHC through higher-dimensional operators. In single-top-quark production, such effects manifest themselves both in the total event yields and in subtle distortions of the multi-dimensional angular distributions of the top-quark decay products. The statistical model therefore exploits information from signal, background, and CRs, while incorporating experimental and theoretical systematic uncertainties that affect both the normalisation and the shapes of the predicted distributions.

From effective operators to fit parameters. The deviations probed in this work originate from a well-defined subset of dimension-six operators that modify the tWb vertex, as introduced in Chapter 2, Section 2.3.1:

$$\mathcal{O}_{\phi Q}^{(3)}, \quad \mathcal{O}_{\phi tb}, \quad \mathcal{O}_{bW}, \quad \mathcal{O}_{tW}, \quad \mathcal{O}_{Qq}^{3,1}. \quad (6.16)$$

They introduce eight a priori independent Wilson coefficients: two real ($c_{\phi Q}$, c_{Qq}) and three complex ($c_{\phi tb}$, c_{bW} , c_{tW}). Not all of these complex phases can be accessed experimentally. At LO and in the massless b -quark limit, the top-quark decay amplitude factorises into four helicity structures, two involving a left-handed b -quark and two a right-handed one [247]. This naturally groups the coefficients into a left-handed sector, (V_L, c_{tW}) , and a right-handed sector, $(c_{\phi tb}, c_{bW})$. Observables are sensitive only to the relative phase between the right-handed coefficients; absolute phases remain unmeasurable unless one retains b -quark-mass suppressed terms of order m_b/m_W or m_b/m_t , whose numerical impact lies well below the current precision.

To remove this redundant degree of freedom, the common convention of taking $c_{\phi tb}$ to be real and positive is adopted. With this choice the parameter space probed in the likelihood fit reduces to seven real quantities,

$$\vec{C} = (C_{\phi q}, C_{tW}, C_{itW}, C_{bW}, C_{ibW}, C_{\phi tb}, C_{qQ}), \quad (6.17)$$

which enter the theoretical predictions through

$$\begin{aligned} V_L &= V_{tb} + \frac{v^2}{\Lambda^2} C_{\phi q}, & c_{tW} &= C_{tW} + i C_{itW}, \\ c_{\phi tb} &= |C_{\phi tb}| + i0, & c_{bW} &= C_{bW} + i C_{ibW}, \\ c_{Qq}^{3,1} &= C_{qQ}. \end{aligned} \quad (6.18)$$

Fixing $c_{\phi tb}$ in this way is a purely technical choice: multiplying both $c_{\phi tb}$ and c_{bW} by a

common phase leaves all observables unchanged and therefore does not bias the inference.

Rate versus shape sensitivity. Among the operators listed in Equation 6.16, $\mathcal{O}_{\phi Q}^{(3)}$ shares the same Lorentz structure as the SM tWb interaction. Its effect is confined to an overall rescaling of the cross-section and hence cannot be disentangled from V_L using shape information alone. Consequently, normalised angular observables constrain only ratios such as C_{tW}/V_L and C_{itW}/V_L , whereas absolute information on V_L (and thus on $C_{\phi q}$) is extracted from the inclusive rate. A simultaneous likelihood fit to both rate and shape observables is therefore essential to exploit the full constraining power of the data and to provide comprehensive limits on new physics in the tWb sector.

6.7 The profile likelihood fit of the analysis

A profile likelihood fitting is a fundamental tool in physics analyses, enabling precise extraction of physics parameters while systematically incorporating uncertainties. In this analysis, a maximum likelihood fit is performed to constrain Wilson coefficients, leveraging the M-function formalism (see Section 2.4.1) to capture the full angular structure of the single-top-quark t -channel process. Unlike traditional binned likelihood fit approaches, which rely on histograms and event counts in predefined bins, this analysis directly fits the normalised M-function coefficients, providing a more unbinned and information-enriched approach.

6.7.1 Computation of the M-functions

In order to evaluate the M-function coefficients $c_{m'm}^{j_1, j_2}$ discussed in Section 2.4.1, it is necessary to determine the M-functions themselves. The M -functions are built from two Wigner D -functions previously defined in Equation 2.20, which on the other hand are elements of Wigner D -Matrix. For efficiency reasons in this analysis, entire D -Matrix is calculated and cached. The workflow of the calculation is as follows.

The Wigner D -matrix describes the transformation of quantum states under three-dimensional rotations. Although in the case of this analysis two consecutive transformations of two D -Functions are considered.

$$D_{m'm}^j(\phi, \theta, 0) = \langle j, m' | e^{-i\phi J_z^{(j)}} e^{-i\theta J_y^{(j)}} | j, m \rangle. \quad (6.19)$$

A rotation about the y -axis by an angle θ , and a rotation about the z -axis by an

angle ϕ . The matrix elements of the Wigner D -matrix can be explicitly computed by exponentiating the angular momentum generators $J_y^{(j)}$ and $J_z^{(j)}$.

The first factor, $e^{-i\phi J_z^{(j)}}$, is diagonal in the basis of $J_z^{(j)}$, meaning its exponentiation is straightforward:

$$(e^{-i\phi J_z^{(j)}})_{m'm} = e^{-i\phi} \delta_{m'm}. \quad (6.20)$$

This indicates that the effect of a rotation around the z -axis is simply a phase shift proportional to the eigenvalue m . The second factor, $e^{-i\theta J_y^{(j)}}$, is not diagonal in the standard basis and must be computed via matrix exponentiation. This process can be significantly accelerated using the similarity transformation:

$$e^{-i\theta J_y^{(j)}} = A \cdot e^{-i\theta J_z^{(j)}} \cdot A^T, \quad (6.21)$$

where A is the matrix that diagonalises $J_y^{(j)}$. Since $J_z^{(j)}$ is diagonal, its exponentiation is immediate:

$$e^{-i\theta J_z^{(j)}} = \text{diag}(e^{-i\theta j}, e^{-i\theta(j-1)}, \dots, e^{-i\theta(-j)}). \quad (6.22)$$

Detailed calculation of matrices $J_y^{(j)}$ is given in the Appendix E, for specific values of j , the matrices $J_y^{(j)}$ are:

$$J_y^{(1)} = \begin{pmatrix} 0 & -i\sqrt{2}/2 & 0 \\ i\sqrt{2}/2 & 0 & -i\sqrt{2}/2 \\ 0 & i\sqrt{2}/2 & 0 \end{pmatrix}, \quad (6.23)$$

$$J_y^{(2)} = \begin{pmatrix} 0 & -i & 0 & 0 & 0 \\ i & 0 & -i\sqrt{3}/2 & 0 & 0 \\ 0 & i\sqrt{3}/2 & 0 & -i\sqrt{3}/2 & 0 \\ 0 & 0 & i\sqrt{3}/2 & 0 & -i \\ 0 & 0 & 0 & i & 0 \end{pmatrix}. \quad (6.24)$$

Thus, the Wigner D -matrix elements can be computed efficiently by applying the

transformation A and using the diagonal exponentiation. The full D -matrix is then obtained by multiplying these exponentiated components together.

6.7.2 Components of the likelihood

This analysis employs a maximum log-likelihood fitting procedure on a carefully selected subset of the full data—referred to as dataset $\mathcal{D}$ —using a model defined by adjustable parameters. In practice, the observable dataset $\mathcal{D}$ is constructed by applying both signal-enhancing selection criteria and additional requirements (such as event quality, kinematic thresholds, and detector acceptance conditions) to the full data sample starting from event selection, as discussed in Section 6.2, to region definition, described in Section 6.3. The fitting function models $\mathcal{D}$ by accounting for the overall event counts as well as the detailed angular information captured by the M -function coefficients. In the fit, the model predictions depend on key physics parameters—such as the Wilson coefficients $\vec{C}$ —and on normalisation factors encapsulated in the scale vector $\vec{S}$.

Modelling the event counts

In single-top-quark production via the t -channel at the LHC, top quarks and top antiquark are produced through different partonic subprocesses, primarily due to the asymmetric PDFs of the proton. The dominant mechanism for producing single top quarks involves valence up-type quarks, while top antiquarks production relies on down-type sea antiquarks, which are less abundant. Consequently, the cross-section for top-quark production is approximately twice that of top antiquark production [84]. Additionally, the differing initial-state kinematics lead to variations in the reconstructed observables and polarisation features of top-quark and top-antiquark events [84]. To capture these effects accurately, the analysis is performed separately for events with positively and negatively charged leptons, which reflect the charge of the parent top quark: $t \rightarrow W^+b \rightarrow \ell^+\nu b$ and $\bar{t} \rightarrow W^-\bar{b} \rightarrow \ell^-\bar{\nu}\bar{b}$. This separation enhances sensitivity to the production asymmetry and allows for independent extraction of top-quark and top-antiquark polarisation observables [84].

For each top-quark sign (with $s \in \{+, -\}$), the model predicts the number of events in:

- The SR: Where the contribution from single-top-quark production is most prominent.
- The CRs ($t\bar{t}$ CR and W +jets CR, also named CR1 and CR2, respectively): These are designed to be enriched in $t\bar{t}$ and W +jets background processes, respectively, to give the information about the size of the main background processes and to help constrain the NPs associated with systematic uncertainties.

The event counts are collectively denoted for a given lepton charge by

$$\vec{N}^{(s)} = \left(N^{(s,\text{SR})}, N^{(s,\text{CR1})}, N^{(s,\text{CR2})} \right).$$

Each event count is built up as a sum of contributions from different processes. For the SR, the prediction is written as:

$$N^{(s,\text{SR})}(\vec{C}, \vec{S}) = f^{t\text{-channel}} N_{t\text{-channel}}^{(s,\text{SR})}(\vec{C}) + f^{W+\text{jets}} N_{W+\text{jets}}^{(s,\text{SR})} + f^{t\bar{t}} N_{t\bar{t}}^{(s,\text{SR})}(\vec{C}) + \sum_{\text{bkg}} N_{\text{bkg}}^{(s,\text{SR})}. \quad (6.25)$$

Here $f^{t\text{-channel}}$, $f^{W+\text{jets}}$, and $f^{t\bar{t}}$ are scale factors that adjust the contribution from each process (i.e. t -channel, W +jets and $t\bar{t}$, respectively). The $N_{t\text{-channel}}^{(s,\text{SR})}(\vec{C})$ and $N_{t\bar{t}}^{(s,\text{SR})}(\vec{C})$ are the event counts for the single-top-quark and $t\bar{t}$ processes, respectively. They are “morphed” functions—that is, they include the dependence on the Wilson coefficients $\vec{C}$. The $N_{W+\text{jets}}^{(s,\text{SR})}$ is the contribution from the W +jets CR that does not have dependence on Wilson coefficients (since any top quark is involved). The $\sum_{\text{bkg}} N_{\text{bkg}}^{(s,\text{SR})}$ term accounts for other background contributions, under the assumption that they do not exhibit an explicit $\vec{C}$ dependence.

For the CRs, which are primarily used to calibrate background levels and systematic uncertainties, the model prediction is:

$$N^{(s,\text{CR})}(\vec{S}) = f^{t\text{-channel}} N_{t\text{-channel}}^{(s,\text{CR})} + f^{W+\text{jets}} N_{W+\text{jets}}^{(s,\text{CR})} + f^{t\bar{t}} N_{t\bar{t}}^{(s,\text{CR})} + \sum_{\text{bkg}} N_{\text{bkg}}^{(s,\text{CR})}.$$

In these regions, the processes depend on scale factor $\vec{S}$.

As shown in Chapter 4, Figure 4.5 displays how the morphed event counts for the top-quark ($s = +$) and top-antiquark ($s = -$) channels vary as functions of individual Wilson coefficients. Furthermore, Figure 4.6 shows the two-dimensional dependence on pairs of coefficients (such as $c_{\phi tb}$ and c_{bw}), revealing their expected correlation.

Modelling the M -function coefficients

The M -function coefficients are denoted by $\vec{c}^{(s)}$ (see Equation 2.21), and the model includes a set of 2×23 coefficients, with the index i running over 23 components for each top-quark (or decaying lepton) sign s .

For each coefficient $c_i^{(s)}$, the measured value is computed from Equation 2.21 on an event-by-event basis. In practice, for every event the relevant angular variables ϕ , θ , ϕ^* , and θ^* are measured and the corresponding M -function is evaluated, as explained in

Section 2.4.1. Each event contributes to the coefficient $c_i^{(s)}$ by taking its event weight ρ and multiplying it by the complex conjugate of the evaluated M-function. In other words, the event-by-event contribution is given by the term $\rho \cdot \left(M_{m'm}^{j_1 j_2}(\phi, \theta, \phi^*, \theta^*) \right)^*$. The final value of $c_i^{(s)}$ is then obtained by summing this contribution over all events in the dataset.

Since Fourier decomposition is a linear operation, the model for these coefficients, $c_i^{(s)}(\vec{C}, \vec{S})$, is then constructed as a linear combination of contributions from several processes.

$$\vec{c}^{(s, \text{SR})}(\vec{C}, \vec{S}) = f^{t\text{-channel}} \vec{c}_{t\text{-channel}}^{(s, \text{SR})}(\vec{C}) + f^{W+\text{jets}} \vec{c}_{W+\text{jets}}^{(s, \text{SR})} + f^{t\bar{t}} \vec{c}_{t\bar{t}}^{(s, \text{SR})}(\vec{C}) + \sum_{\text{bkg}} \vec{c}_{\text{bkg}}^{(s, \text{SR})}. \quad (6.26)$$

- **Signal (t -channel) contribution:** This term depends on the Wilson coefficients $\vec{C}$ because new physics would alter the angular distribution of the single-top-quark signal.
- **$t\bar{t}$ and W + jets backgrounds:** These are included with their respective scale factors and, in the case of $t\bar{t}$, with an additional dependence on $\vec{C}$.
- **Other backgrounds:** These are additional, usually less dominant, contributions. The effect of subdominant processes—including tW , s -channel, tH , and other rare production modes—on the coefficients $\vec{C}$ is neglected in this analysis.

In order to focus on the shape of the angular distributions rather than the absolute event yield, the M -function coefficients are normalised by the total event count in the SR:

$$\langle \vec{c} \rangle^{(s, \text{SR})}(\vec{C}, \vec{S}) = \frac{\vec{c}^{(s, \text{SR})}(\vec{C}, \vec{S})}{N^{(s, \text{SR})}(\vec{C}, \vec{S})}.$$

This normalisation ensures that the sensitivity to the EFT parameters is encapsulated in the variation (or “shape”) of the angular distributions rather than in their overall normalisation.

Normalised M -function coefficients are presented as functions of the Wilson coefficients in Figures 4.5 and 4.7. These distributions illustrate how the angular distributions—and hence the M -function coefficients—respond to variations in the EFT parameters.

6.7.3 Construction of the likelihood function

This analysis uses data that is partitioned into several regions: two SRs and four CRs. The CRs for $t\bar{t}$ and W +jets are also built separately for top-quark and top-antiquark production. This results in four distinct CRs—one for each background process and each

top-quark sign configuration. The observed event counts in each region are denoted by $N^{(s,j)}$ with j representing the region. For example, j can be SR for the SR, CR1 for the W +jets dominated region, or CR2 for the $t\bar{t}$ dominated region. The model predicts these counts, and the predictions are denoted as $N_j^{(s),\text{model}}(\vec{C}, \vec{S})$, where $\vec{C}$ and $\vec{S}$ are vectors of parameters that govern theoretical expectations, as defined earlier in this chapter.

In addition to the simple event counts, the analysis uses a set of observables known as M -function coefficients. These coefficients, denoted by $\vec{c}^{(s)}$, are derived from distributions—such as angular distributions—in the SR. To simplify the notation, the normalised version is used, represented as $\langle \vec{c} \rangle^{(s,\text{SR})}$; here, the normalisation ensures that the model compares shapes rather than absolute scales, thereby focusing on the detailed features of the decay process. The coefficients are predicted as $\vec{c}^{(s),\text{model}}(\vec{C}, \vec{S})$.

The goal is to construct a likelihood function that quantifies how well the model, with its parameters and uncertainties, describes the observed data. In statistical analysis, the likelihood function represents the probability of observing the data given a set of parameters. The best-fit parameters of the model are determined by maximising this function, or equivalently, by minimising its negative logarithm.

The statistical part of the likelihood is considered first, neglecting systematic uncertainties initially. It comprises two main components:

1. The M -function coefficients:

It is assumed that the differences between the measured coefficients $\vec{c}^{(s)}$ and their model predictions $\vec{c}^{(s),\text{model}}(\vec{C}, \vec{S})$ follow a multivariate Gaussian distribution. This is a common assumption when the uncertainties are well understood and the number of data points is large. The comparison is quantified by the term

$$-\frac{1}{2} \left(\vec{c}^{(s)} - \vec{c}^{(s),\text{model}}(\vec{C}, \vec{S}) \right)^T \left[\mathbb{C}^{(s)} \right]^{-1} \left(\vec{c}^{(s)} - \vec{c}^{(s),\text{model}}(\vec{C}, \vec{S}) \right),$$

where $\mathbb{C}^{(s)}$ is the covariance matrix, a 23×23 matrix that encodes the uncertainties and their correlations among the coefficients. This covariance is a sum of uncertainties coming from the data ($\mathbb{C}_{\text{data}}^{(s)}$) and from the base MC model ($\mathbb{C}_{\text{model}}^{(s)}$):

$$\mathbb{C}^{(s)} = \mathbb{C}_{\text{data}}^{(s)} + \mathbb{C}_{\text{model}}^{(s)}.$$

This term quantifies the deviation of the observed coefficients from the model predictions, while taking into account the dependencies among the coefficients.

2. The event counts:

For the event counts in each region, the observed counts $N^{(s,j)}$ are compared with

the model predictions $N^{(s,j),\text{model}}(\vec{C}, \vec{S})$. Although the true statistical nature of count data is Poissonian, when the counts are large a Gaussian approximation is often used. In this approximation, the variance is taken to be approximately equal to the observed count itself, leading to a term like

$$-\frac{1}{2} \sum_{j \in \{\text{SR}, \text{CR1}, \text{CR2}\}} \frac{\left(N^{(s,j),\text{model}}(\vec{C}, \vec{S}) - N^{(s,j)}\right)^2}{N^{(s,j)}}.$$

This term penalises deviations between the observed and predicted event counts, weighted by the statistical uncertainty in each region.

Putting these two pieces together, the log-likelihood function that captures our statistical uncertainties is given by:

$$\begin{aligned} \ln \mathcal{L}_{\text{stat}}(\vec{C}, \vec{S}) = & -\frac{1}{2} \sum_{s \in \{+, -\}} \left\{ \left(\vec{c}^{(s)} - \vec{c}^{(s),\text{model}}(\vec{C}, \vec{S}) \right)^T [\mathbb{C}^{(s)}]^{-1} \right. \\ & \left. \times \left(\vec{c}^{(s)} - \vec{c}^{(s),\text{model}}(\vec{C}, \vec{S}) \right) + \sum_j \frac{\left(N^{(s,j),\text{model}}(\vec{C}, \vec{S}) - N^{(s,j)}\right)^2}{N^{(s,j)}} \right\}. \end{aligned} \quad (6.27)$$

In any real-world analysis, systematic uncertainties are unavoidable. As already discussed and detailed in Section 6.5, these uncertainties might arise from various sources such as detector calibration, modelling of the physics process, or theoretical assumptions. To handle these, NPs are introduced that allow the model predictions to adjust in response to systematic effects. As already mentioned before, this process is known as profiling.

For each source of systematic uncertainty (denoted by the subindex u), the corresponding variation is defined by its one-standard-deviation magnitude, σ_u , which quantifies the size of the uncertainty. Then, using this information, new simulations are generated where the parameter u is shifted by $\pm\sigma_u$ (often by reweighting the events from the nominal simulation). From these shifted simulations, the changes in the observables are measured—both for the M -function coefficients (denoted $\vec{c}_u^{(s)}$) and the event counts (denoted $\vec{N}_u^{(s)}$). The differences between these shifted values of the model and the nominal values (labelled as $\vec{c}_{\text{SM}}^{(s),\text{model}}$ and $\vec{N}_{\text{SM}}^{(s),\text{model}}$) are then computed:

$$\begin{aligned} \Delta \vec{c}_u^{(s)} &= \vec{c}_u^{(s)} - \vec{c}_{\text{SM}}^{(s),\text{model}}, \\ \Delta \vec{N}_u^{(s)} &= \vec{N}_u^{(s)} - \vec{N}_{\text{SM}}^{(s),\text{model}}. \end{aligned}$$

It is assumed that these differences change linearly with the size of the systematic

uncertainty. To include these effects in the model, a NP θ_u is introduced for each uncertainty. The model predictions are then modified by adding these systematic shifts scaled by θ_u :

$$\begin{aligned}\vec{c}^{(s),\text{model}}(\vec{C}, \vec{S}, \vec{\theta}) &= \vec{c}^{(s),\text{model}}(\vec{C}, \vec{S}) + \sum_u \theta_u \Delta \vec{c}_u^{(s)}, \\ \vec{N}^{(s,j),\text{model}}(\vec{C}, \vec{S}, \vec{\theta}) &= \vec{N}^{(s,j),\text{model}}(\vec{C}, \vec{S}) + \sum_u \theta_u \Delta \vec{N}_u^{(s,j)}.\end{aligned}$$

Here, $\theta_u = 1$ means the systematic is shifted by one sigma, and the values of θ_u are allowed to vary in the fit, but with a constraint that they should stay close to zero (their nominal value). This constraint is imposed by a Gaussian prior $e^{-\theta_u^2/2}$, which effectively penalises large deviations from zero unless the data strongly indicate them.

Combining the statistical likelihood with the treatment of systematic uncertainties, the full log-likelihood function becomes dependent on the NPs, denoted by the vector $\vec{\theta}$, which encapsulates all sources of systematic uncertainty considered in the model.

$$\begin{aligned}\ln \mathcal{L}(\vec{C}, \vec{S}, \vec{\theta}) &= -\frac{1}{2} \left\{ \sum_{s \in \{+, -\}} \left[(\vec{c}^{(s)} - \vec{c}^{(s),\text{model}}(\vec{C}, \vec{S}, \vec{\theta}))^T [\mathbb{C}^{(s)}]^{-1} \times \right. \right. \\ &\quad \left. \times (\vec{c}^{(s)} - \vec{c}^{(s),\text{model}}(\vec{C}, \vec{S}, \vec{\theta})) + \sum_{j \in \{\text{SR}, \text{CR1}, \text{CR2}\}} \frac{(N^{(s,j),\text{model}}(\vec{C}, \vec{S}, \vec{\theta}) - N^{(s,j)})^2}{N^{(s,j)}} \right] + \vec{\theta} \cdot \vec{\theta} \right\}.\end{aligned}\tag{6.28}$$

In this equation, the first term inside the sum accounts for the differences in the M -function coefficients, now modified by the systematic shifts, while the second term does the same for the event counts. The final term, $\vec{\theta} \cdot \vec{\theta}$, is the contribution from the Gaussian priors on the NPs.

6.7.4 Fitting procedure

This analysis employs a two-step fitting procedure, applied first to an Asimov dataset and subsequently to real data. The Asimov dataset, constructed to represent the expected statistical outcome under the nominal (null) hypothesis, serves as a deterministic benchmark free of statistical fluctuations. It is used to evaluate the expected sensitivity of the measurement, validate the stability of the fit, and optimise methodological choices—such as binning strategies and regularisation—prior to unblinding. By relying on this synthetic dataset during the development phase, potential biases introduced by

premature inspection of real data are avoided, ensuring a more robust and objective inference process.

An Asimov dataset is a synthetic dataset constructed using the expected values of all observables under a given model. In practice, this means filling each analysis bin with the exact expected number of events from signal and/or background instead of random Poisson fluctuations. For unbinned analyses, one can generate a binned dataset fine-grained enough to capture the underlying distribution. The Asimov dataset is named after science fiction author Isaac Asimov (in reference to a story about a single representative voter), reflecting the idea that this dataset represents the “average” outcome of many experiments. The Asimov data is generated for a specific set of model parameters such that if one fits it, the maximum-likelihood estimates (MLEs) of those parameters equal their true input values. In other words, performing a likelihood fit to the Asimov dataset returns $\hat{\mu} = \mu_0$ and $\hat{\theta} = \theta_0$, where μ_0 and θ_0 are the nominal input values used to construct the Asimov dataset—typically corresponding to the SM prediction for the signal strength and the central values of the NPs, respectively.

The main motivation for using Asimov data is that it allows for efficient estimation of sensitivities, significance, limits and tests statistical model reliability itself, without using real data and therefore avoiding any bias because of this.

Here some considerations:

- The likelihood fit is performed using the negative log-likelihood function, $-\ln \mathcal{L}$, to determine the most probable values of the parameters of interest (POIs) $\vec{C}$, the scale factors $\vec{S}$, and the NPs $\vec{\theta}$. The minimisation is carried out using the Minuit2⁵ library [248] with $6 + 3 + 206$ free-floating parameters, yielding a best-fit likelihood value, $-\ln \hat{\mathcal{L}}$. For most parameters, where the likelihood function remains approximately quadratic, confidence intervals can be determined using a standard parabolic approximation. However, for certain Wilson coefficients, such as $C_{\phi tb}$, C_{bw} , and C_{ibw} , the log-likelihood exhibits a quadratic dependency that renders the usual parabolic uncertainty estimation from the RooFit/RooStats library [249], a toolkit for modelling the expected distribution of events in physics analysis. While the Minos⁶ algorithm could, in principle, determine asymmetric confidence intervals

⁵Minuit2 is a numerical minimisation library developed at CERN, widely used in high-energy physics for function minimisation and parameter estimation. It is part of the ROOT framework and provides robust algorithms for likelihood and χ^2 minimisation, along with uncertainty evaluation and parameter correlations.

⁶Minos is an algorithm within the Minuit2 minimisation package that provides asymmetric confidence intervals for fitted parameters. Unlike simpler methods that assume a parabolic likelihood shape (e.g., based on the Hessian matrix), Minos accounts for non-linearities and correlations by scanning the likelihood function, yielding more accurate uncertainty estimates, especially in cases with non-Gaussian

in such cases, it fails to converge in this analysis due to the Hessian matrix⁷ not always being positive definite. To resolve this issue, an alternative approach is used to extract parameter uncertainties.

- To obtain reliable confidence intervals, a combination of Asimov data and likelihood scans is used. Typically, one-dimensional likelihood scans are performed to determine confidence intervals for $\vec{C}$ and $\vec{S}$. In this procedure, a given parameter is fixed at a scan point while the remaining $6 + 3 + 206 - 1$ parameters are allowed to float. The log-likelihood is then minimised using Minuit2, and the resulting $\ln \mathcal{L}[C_i]$ is used to extract the 1σ and 2σ confidence intervals, corresponding to points where $-2\Delta \ln \mathcal{L}(C_i) = -2\ln[\mathcal{L}(C_i)/\ln \hat{\mathcal{L}}]$ takes values of 1 and 2, respectively.
- The correlation matrix provided by Minuit2 is also extracted from the full combined fit. However, for $C_{\phi tb}$, C_{bw} , and C_{ibw} , where the log-likelihood dependency is quadratic, the results from the Minuit2 correlation matrix are no longer meaningful. This is because it is estimated using the Hessian matrix, which depends on second-order partial derivatives. To fully understand the correlation between each element of $\vec{C}$ and $\vec{S}$, a two-dimensional scan is performed, by fixing two parameters at their scanning point C_i, C_j , allowing the rest to float freely, and then implementing the Minuit2 minimisation to get $\mathcal{L}[C_i, C_j]$.
- In standard ATLAS analyses, the impact of NPs is often assessed by comparing the best-fit values of the POIs when θ_α is fixed at 0 and ± 1 (pre-fit) or shifted to its 1σ best-fit value (post-fit). However, due to the non-parabolic nature of the $-\ln \mathcal{L}$ in this analysis, this approach is not valid. The best-fit values for some POIs lie in a region where the likelihood function has a “flat-bottom” shape, making comparisons at fixed values unreliable. The impact of NPs is assessed by comparing the value of $-\ln \mathcal{L}$ before and after allowing specific categories of θ_α to vary freely in the fit. These categories correspond to different sources of systematic uncertainty.

Finally, the MC statistical uncertainties also play a role in the likelihood fit, affecting the precision of the extracted results. To quantify this impact, the change in $-\ln \mathcal{L}$ is evaluated before and after incorporating $\mathbb{C}_{\text{model}}$ into the total covariance matrix $\mathbb{C}$. This comparison indicates the extent to which MC statistical uncertainties affect the extracted parameter constraints.

behaviour or constrained parameters.

⁷In Minuit2, the correlation matrix ρ is defined as $\rho_{ij} = \frac{V_{ij}}{\sqrt{V_{ii}}\sqrt{V_{jj}}}$, where V_{ij} is the covariance matrix derived from inverting the Hessian matrix $\mathbb{V} = \mathbb{H}^{-1}$. The Hessian matrix is a square matrix of the second-order partial derivatives for the minimised function f : $H_{ij} = \frac{\partial^2 f}{\partial p_i \partial p_j}$.

6.8 Results

The likelihood function is analysed using likelihood scans, where the statistical test is defined as the profile log-likelihood ratio. Here, a shorthand notation is adopted, where C_i refers to an element in either the POI vector $\vec{C}$ or the scale-factor vector $\vec{S}$:

$$-2\Delta \ln \mathcal{L}(C_i) = -2 \ln \frac{\hat{\mathcal{L}}(C_i)}{\hat{\mathcal{L}}} = -2 \ln \frac{\mathcal{L}(C_i, \hat{C}_j[C_i], \hat{\theta}[C_i])}{\mathcal{L}(\hat{C}, \hat{\theta})}, \quad (6.29)$$

where $\hat{C}$ and $\hat{\theta}$ represent the global maximum-likelihood estimates of the POIs and NPs, respectively. The terms $\hat{C}_j[C_i]$ and $\theta[C_i]$ correspond to the values of the parameters $C_j \neq C_i$ and $\vec{\theta}$ that maximise the likelihood function while keeping C_i fixed.⁸

Having established the overall fitting framework, attention now turns to a thorough examination of the Wilson coefficients and the corresponding physical effects associated with each operator:

C_{tW} and C_{itW} (Top-quark dipole operator $\mathcal{O}_{tW}$):

As explained in Chapter 2, in the SM, the tWb vertex is purely left-handed, following a V–A structure with $V_{tb} \approx 1$ and lacking any dipole contributions. A nonzero C_{tW} therefore signals the presence of new physics that alters this standard interaction by adding a term similar to a magnetic moment coupling between the top quark and the W boson.

This additional interaction can be separated into two components: a real part, $\text{Re } C_{tW}$ (simply C_{tW}), which corresponds to CP-conserving modifications, and an imaginary part, $\text{Im } C_{tW}$ (or C_{itW}), which introduces CP-violating effects. After EW symmetry breaking, an anomalous tWb dipole operator might be expressed as

$$\frac{g^2}{m_W} \bar{b} i \sigma^{\mu\nu} q_\nu (g_L P_L + g_R P_R) t W_\mu^- + \text{h.c.},$$

where g_L and g_R are effective tensor couplings that, in the SMEFT framework, are related to C_{tW} (see Equation 2.12). Thus, while the SM expects a purely left-handed tWb vertex with no dipole term, any observed nonzero C_{tW} —whether through its CP-even real part or its CP-odd imaginary part—would indicate an alteration of the conventional interaction and could point to new physics beyond the SM.

The C_{tW} coefficient also affects the top-quark decay $t \rightarrow Wb$, which is reflected

⁸For the Asimov dataset, the expected values are set to $\hat{C} = \vec{0}$, $\hat{S} = \vec{1}$, and $\hat{\theta} = \vec{0}$.

in the W -boson helicity fractions. In the SM, the W boson from top-quark decay is predominantly longitudinal ($F_0 \sim 70\%$) or left-handed ($F_L \sim 30\%$), with a negligible right-handed fraction ($F_R \sim 0\%$) [250]. An anomalous dipole coupling can modify these fractions—for example, a nonzero $\text{Re } C_{tW}$ primarily shifts F_0 and F_L away from their SM values, and since F_R is extremely small in the SM, any measurable right-handed W -boson fraction would be a clear indication of new physics. Furthermore, the ATLAS collaboration explicitly notes that F_0 and F_L are especially sensitive to C_{tW} [250], while a CP-violating $\text{Im } C_{tW}$ could manifest via subtle angular asymmetries in the decay distributions, such as T-odd correlations between the decay products.

Recent limits from LHC, LEP/SLD and Tevatron on C_{tW} are obtained in a global SMEFT fit considering both linear and quadratic terms. The 95% probability interval is found to be $-0.19 < C_{tW} < 0.50 \text{ TeV}^{-2}$, as reported in Ref. [251].

The ATLAS collaboration also set constraints on the Wilson coefficients in the SMEFT framework using measurements of top-quark polarisation in single-top-quark t -channel production at $\sqrt{s} = 13 \text{ TeV}$, resulting in the 95% confidence intervals $-0.9 < C_{tW} < 1.4 \text{ TeV}^{-2}$ and $-0.8 < C_{itW} < 0.2 \text{ TeV}^{-2}$, as reported in Ref. [252].

The CMS collaboration constrains on C_{tW} in the SMEFT framework, using data collected at $\sqrt{s} = 13 \text{ TeV}$ at integrated luminosity of 138 fb^{-1} , $-0.31 < C_{tW} < 0.22 \text{ TeV}^{-2}$ at the 95% confidence interval [253].

In this analysis, data collected at $\sqrt{s} = 13 \text{ TeV}$ with an integrated luminosity of 140 fb^{-1} are used to constrain the Wilson coefficient C_{tW} via a simultaneous fit that includes seven SMEFT coefficients, three normalisation scale factors, and 206 NPs. The resulting 95% confidence interval is

$$-0.22 < C_{tW} < 0.01 \text{ TeV}^{-2},$$

consistent with the SM expectation of zero. This result improves upon earlier constraints by combining rate and shape information, especially angular observables sensitive to the helicity structure of top-quark decay. The likelihood scans based on Asimov and real data are shown in Figure 6.12, demonstrating both the expected and observed sensitivity.

The CP-violating component $C_{itW} = \text{Im } C_{tW}$ is similarly constrained, yielding

$$-0.07 < C_{itW} < 0.05 \text{ TeV}^{-2},$$

at 95% confidence level. The corresponding likelihood scans in Figure 6.13 confirm

compatibility with the SM and illustrate the sensitivity of angular observables to CP-violating effects. These limits also improve current public constraints.

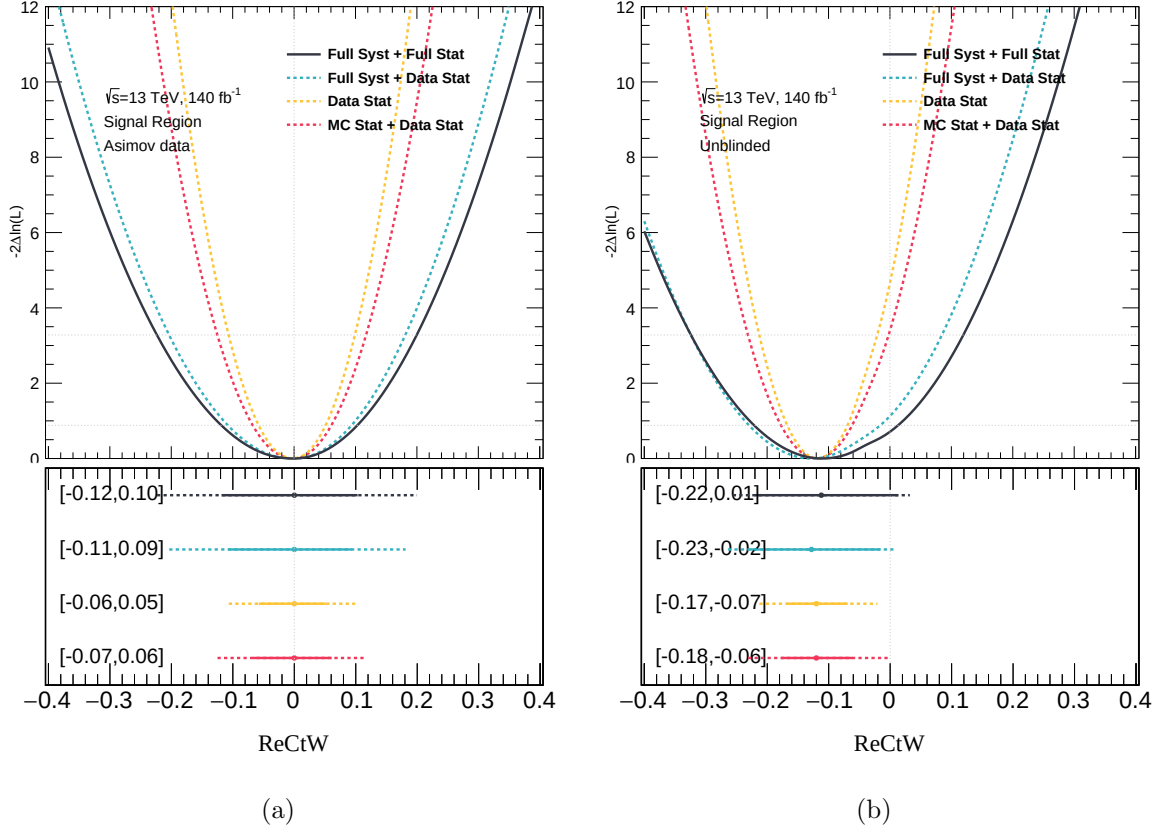


Figure 6.12: Likelihood scans of the parameter of interest C_{tw} , with all other parameters profiled, obtained from a fit to (a) Asimov data and (b) to real data. The different curves represent estimates obtained under varying combinations of statistical and systematic uncertainties. The solid line shows the results when considering all the possible uncertainties.

C_{bW} and C_{ibW} (Bottom-quark dipole operator $\mathcal{O}_{bW}$):

These are the analogues of C_{tW} for the bottom quark. The $\mathcal{O}_{bW}$ operator introduces a tensor coupling g_L involving a right-handed b -quark. A nonzero C_{bW} would allow the top quark to decay into a right-handed b -quark plus a W boson, something highly suppressed in the SM (SM decays yield almost exclusively left-handed b -quarks due to V–A structure as already described in Section 2.3.2). This would modify the decay kinematics – in particular, it would split the longitudinal W -boson fraction into two components depending on helicity of the b -quark. The fully differential decay distribution can distinguish whether the W^+ -boson from top-quark decay is accompanied by a b^+ (spin 1/2) or b^- (spin $-1/2$) quark [90]. The C_{bW} coefficient, thus alters the angular correlations between the W -boson and b -quark spins, potentially yielding a measurable difference

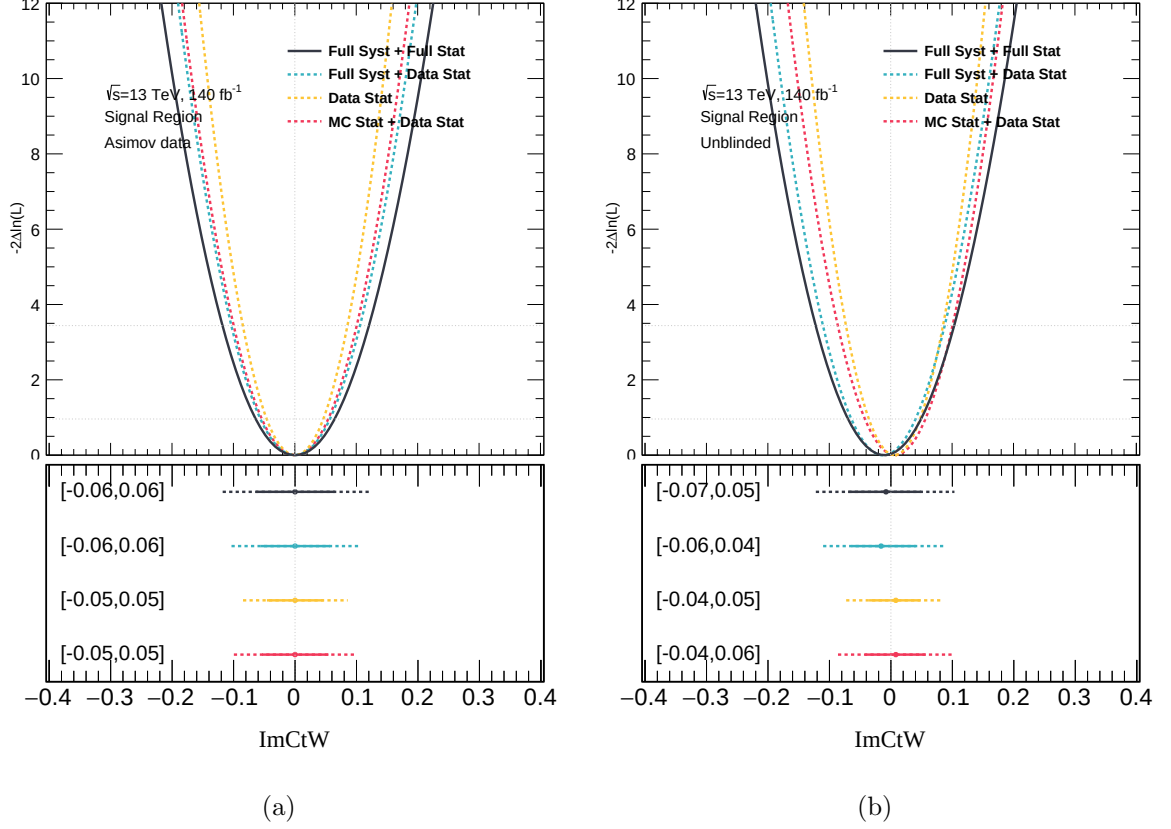


Figure 6.13: Likelihood scans of the parameter of interest C_{itw} , with all other parameters profiled, obtained from a fit to (a) Asimov data and (b) to real data. The different curves represent estimates obtained under varying combinations of statistical and systematic uncertainties. The solid line shows the results when considering all the possible uncertainties.

between the two components of the longitudinal W -boson fraction. In practice, however, contributions of O_{bW} in single-top-quark t -channel are suppressed. In the production process, an insertion of O_{bW} would require a right-handed b -quark in the initial state, which is kinematically disfavoured, or else two insertions (hence order $1/\Lambda^4$) to interfere with the dominant left-handed amplitude [90]. Similarly in decay, the effect is proportional to the small b -quark mass. Therefore, C_{bW} and C_{ibW} are expected to have a much smaller impact (and current limits permit them to be of order unity) compared to C_{tW} .

The global combined fit from LHC, LEP/SLD and Tevatron on C_{bW} at 95% probability interval is $-0.91 < C_{bW} < 0.91 \text{ TeV}^{-2}$, as reported in Ref. [251].

Constraints on C_{bW} by the CMS collaboration in the SMEFT framework, using data collected at $\sqrt{s} = 13 \text{ TeV}$, corresponding to an integrated luminosity of 138 fb^{-1} , is found to be $-0.39 < C_{bW} < 0.40 \text{ TeV}^{-2}$ at 95% confidence interval [253].

In this analysis, using data collected at $\sqrt{s} = 13$ TeV with an integrated luminosity of 140 fb^{-1} , the Wilson coefficients C_{bW} and C_{ibW} are simultaneously fitted alongside all other POIs. The fit to real data yields a 95% confidence interval of

$$0.61 < C_{bW} < 1.25 \text{ TeV}^{-2}, \quad -0.29 < C_{ibW} < 0.25 \text{ TeV}^{-2}.$$

The corresponding likelihood scans are shown in Figures 6.14 and 6.15, comparing results from Asimov data and observed data. While the constraints on C_{bW} are weaker due to its suppressed contribution in the t -channel, the shape and angular information in the analysis still provide sensitivity to this operator. The fitted interval for C_{bW} excludes the SM value of zero at 95% confidence level, indicating a preference for a non-zero contribution. The likelihood scan exhibits a pronounced double-well structure (see Section 6.8.1 for farther description), with two local minima symmetrically displaced from the SM point. This behaviour reflects a degeneracy in the way the operator modifies angular distributions and highlights the limited power of the available observables to fully resolve such effects. Nonetheless, despite the presence of this double-minimum structure, the fit converges numerically to a single best-fit value within the resolution of the algorithm. The imaginary part, C_{ibW} , which introduces CP-violating effects, is found to be consistent with zero within the current experimental precision.

$C_{\phi tb}$ (Higgs-boson–top-quark coupling):

This Wilson coefficient multiplies an operator that involves the Higgs field and two right-handed quarks (t_R and b_R). After EW symmetry breaking, it effectively generates a right-handed vector coupling V_R in the tWb vertex (i.e. a term $\bar{b}_R \gamma^\mu t_R W_\mu$). A nonzero $C_{\phi tb}$ thus directly contributes to V_R in the Lagrangian shown in Equation 2.8. Phenomenologically, this would produce a right-handed tWb current, which would increase the probability of finding a right-handed b -quark in top-quark decay without requiring a momentum flip (unlike the dipole operators). A sizable V_R would tend to reduce the asymmetry in the top-quark decay that stems from the V–A structure and could increase the fraction of transversely polarised W bosons. However, like C_{bW} , the effects of this operator are suppressed by the bottom mass in many observables [90].

As for the recent measurements, the global combined fit from LHC, LEP/SLD and Tevatron on $C_{\phi tb}$ in 95% probability interval is found to be $-7.55 < C_{\phi tb} < 7.05 \text{ TeV}^{-2}$ [251].

The CMS collaboration obtained constraints on $C_{\phi tb}$ in the SMEFT framework, using data collected at $\sqrt{s} = 13$ TeV, and corresponding to an integrated luminosity of 138 fb^{-1} . The 95% confidence interval is found to be $-1.69 < C_{\phi tb} < 1.70 \text{ TeV}^{-2}$, as reported in Ref.

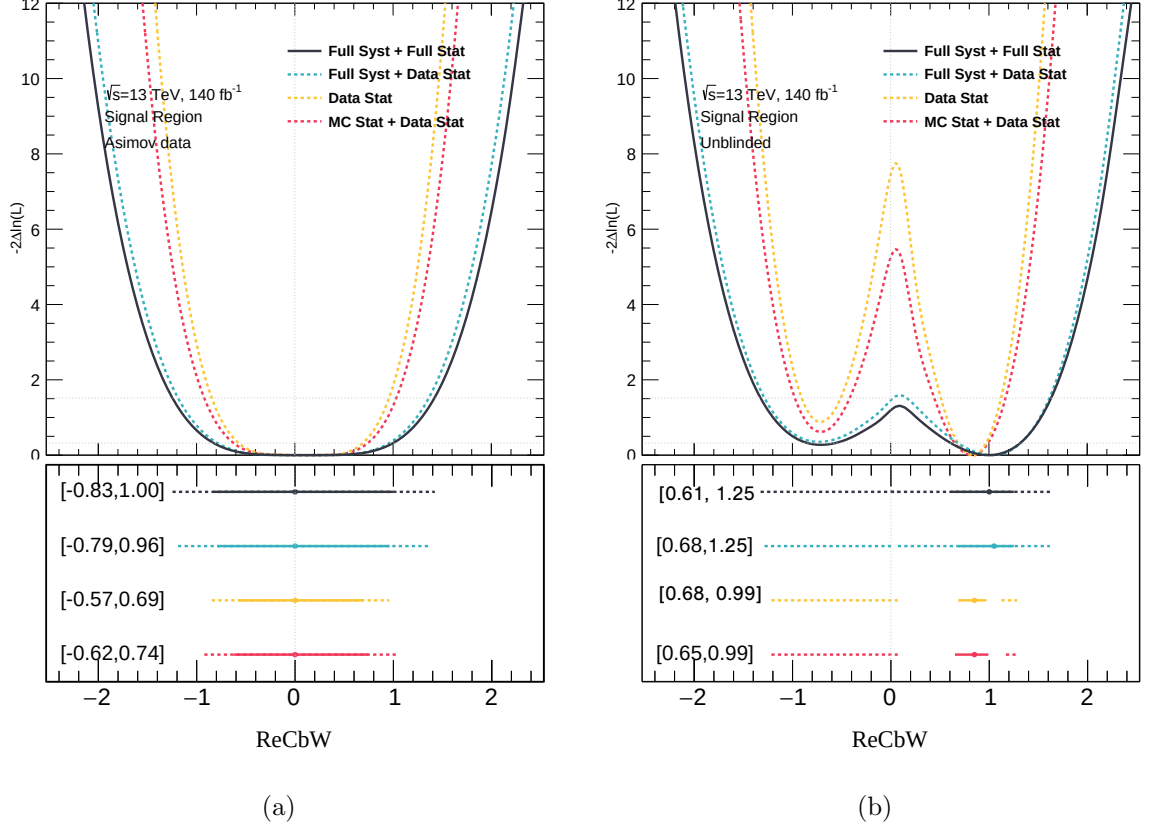


Figure 6.14: Likelihood scans of the parameter of interest C_{bw} , with all other parameters profiled, obtained from a fit to (a) Asimov data and (b) to real data. The different curves represent estimates obtained under varying combinations of statistical and systematic uncertainties. The solid line shows the results when considering all the possible uncertainties.

[253].

In this analysis, using data collected at $\sqrt{s} = 13$ TeV and an integrated luminosity of 140 fb^{-1} , $C_{\phi tb}$ is profiled alongside all other SMEFT parameters. From the fit to real data, the coefficient is constrained to the 95% confidence interval

$$0.34 < C_{\phi tb} < 5.04 \text{ TeV}^{-2}; \quad -6.06 < C_{\phi tb} < -3.00 \text{ TeV}^{-2}.$$

The corresponding likelihood scans, shown in Figure 6.16, compare the constraints obtained from the Asimov dataset (representing the SM expectation) with those extracted from real data. Two distinct intervals are observed, both excluding the SM value at the 95% confidence level. This behaviour arises from the pronounced double-well structure in the likelihood profile—further examined in Section 6.8.1. The presence of two stable and well-separated minima in the likelihood scan is attributed to a strong correlation between

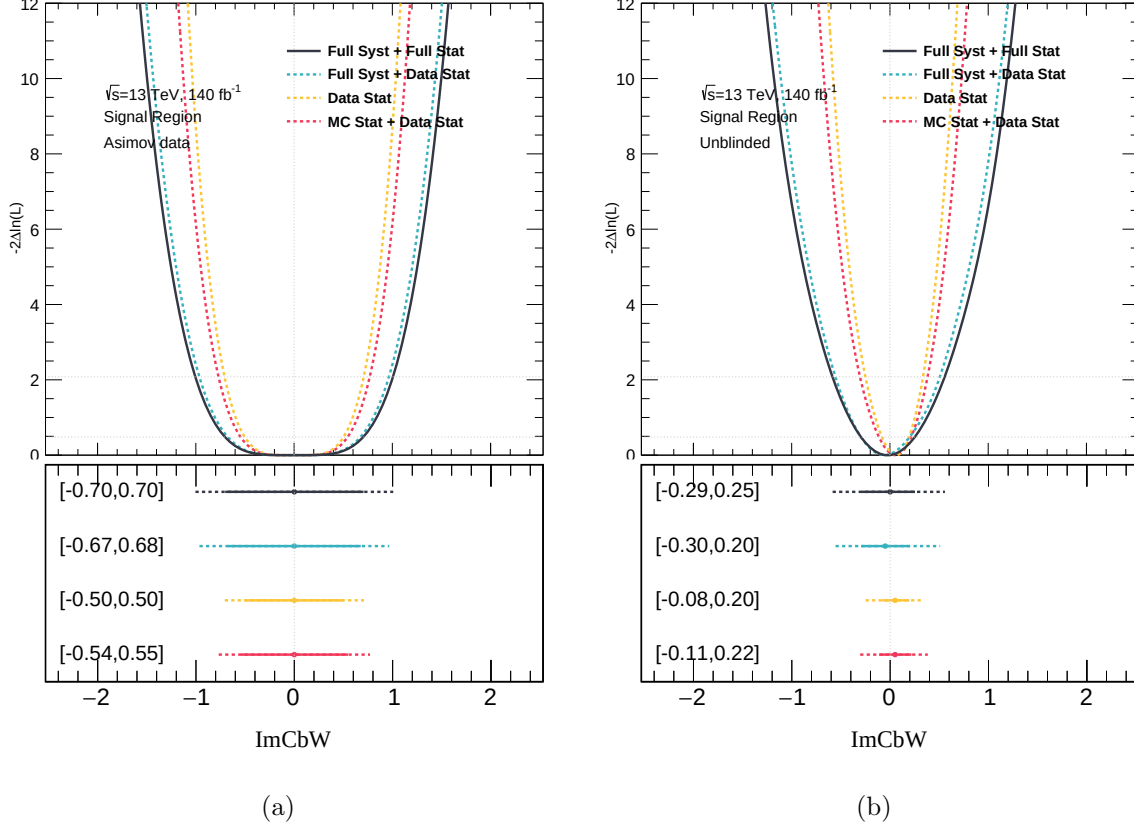


Figure 6.15: Likelihood scans of the parameter of interest C_{ibw} , with all other parameters profiled obtained from a fit to (a) Asimov data and (b) to real data. The different curves represent estimates obtained under varying combinations of statistical and systematic uncertainties. The solid line shows the results when considering all the possible uncertainties.

the C_{bW} and $C_{\phi tb}$ operators, which both exhibit similar double-minimum structures. This correlation leads to a complex likelihood topology where distinct combinations of parameter values produce comparable distortions in the angular observables. As a result, the current dataset lacks the statistical power to disentangle these effects, yielding symmetric degenerate solutions on opposite sides of the SM point.

$C_{\phi q}$ (Higgs-boson–quark doublet coupling):

This coefficient (specifically the triplet operator $O_{\phi q}^{(3)}$) modifies the left-handed tWb coupling. Essentially, $O_{\phi q}$ adds a term $\left(\phi^\dagger i \overleftrightarrow{D}_\mu^I \phi\right)(\bar{q}_i \gamma^\mu \tau^I q_j)$, which shifts the coefficient of $\bar{b}_L \gamma^\mu t_L W_\mu$ after EW symmetry breaking. A positive or negative $C_{\phi q}$ will increase or decrease the effective V_L . In the single-top-quark production cross-section, this can scale the overall rate (since the t -channel rate $\propto |V_L|^2$). However, $C_{\phi q}$ has little effect on normalised angular distributions [250]. All decay angles and spin asymmetries that

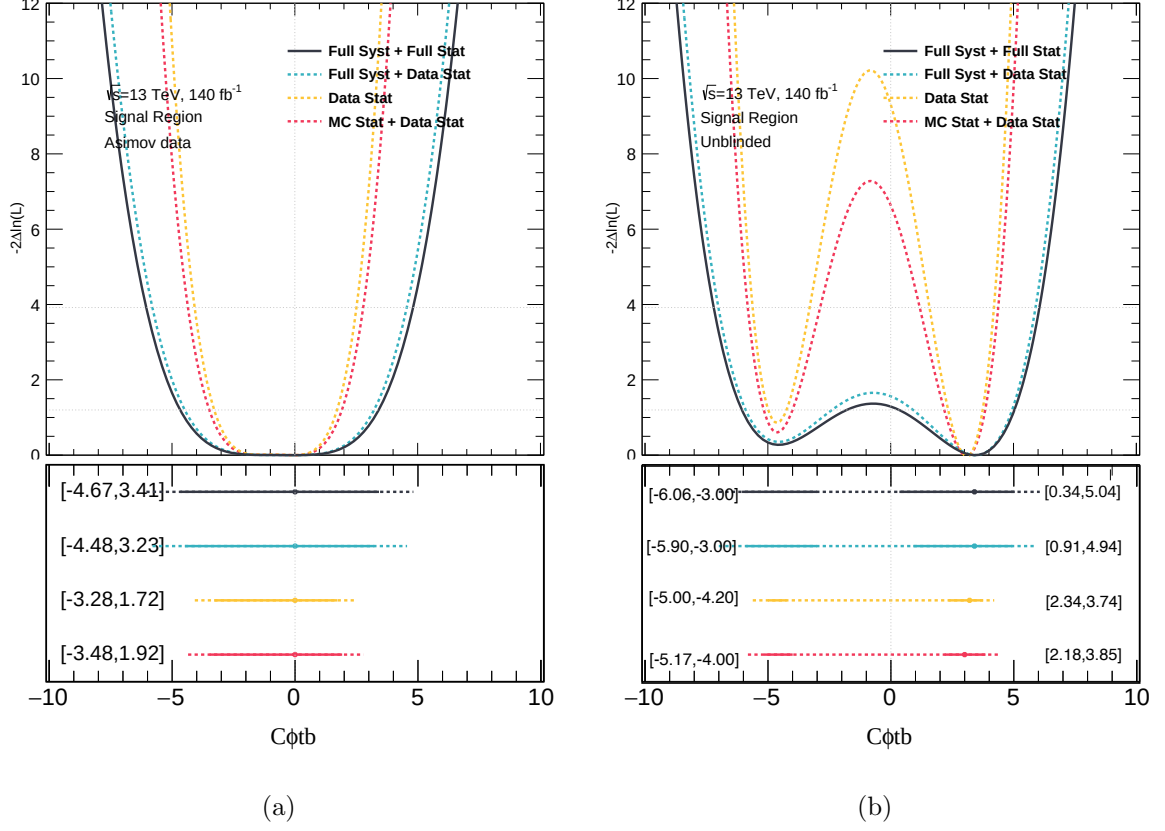


Figure 6.16: Likelihood scans of the parameter of interest $C_{\phi tb}$, with all other parameters profiled, obtained from a fit to (a) Asimov data and (b) to real data. The different curves represent estimates obtained under varying combinations of statistical and systematic uncertainties. The solid line shows the results when considering all the possible uncertainties.

depend on the ratio of amplitudes remain nearly unchanged if one simply rescales the left-handed amplitude. Thus, while $C_{\phi q}$ can affect the total and differential cross-section magnitude, it does not induce distinctive new angular patterns to a first approximation. In practice, analyses often treat $C_{\phi q}$ as a nuisance that is constrained by overall rate measurements, while angular observables are more sensitive to the dipole and right-handed couplings.

Current limits on $C_{\phi q}$ in 95% probability interval measures $-0.94 < C_{\phi q} < 0.81 \text{ TeV}^{-2}$ [251] in combined global fit from LHC, LEP/SLD and Tevatron.

The ATLAS collaboration constrains this coefficient by using measurements of the single-top-quark t -channel cross-section at 13 TeV with 95% confidence interval to be $-0.87 < C_{\phi q} < 1.42 \text{ TeV}^{-2}$ [254].

The limits for $C_{\phi q}$ from the CMS collaboration in the SMEFT framework, using data

collected at $\sqrt{s} = 13$ TeV and an integrated luminosity of 138 fb^{-1} , is $-0.09 < C_{\phi q} < 1.35 \text{ TeV}^{-2}$ for the 95% confidence interval [253].

In this analysis, using data collected at $\sqrt{s} = 13$ TeV by the ATLAS detector and corresponding to an integrated luminosity of 140 fb^{-1} , a fit to real data including simultaneous profiling of all parameters results in the constraint

$$-0.16 < C_{\phi q} < 0.95 \text{ TeV}^{-2},$$

at the 95% confidence level. This range is compatible with previous measurements and confirms that the data remain consistent with the SM expectation. Figure 6.17 presents the likelihood scan for this coefficient, comparing the Asimov dataset (SM baseline) with the actual data. The shape of the likelihood shows that the data slightly prefer a positive deviation, although the SM value of zero remains within the allowed region.

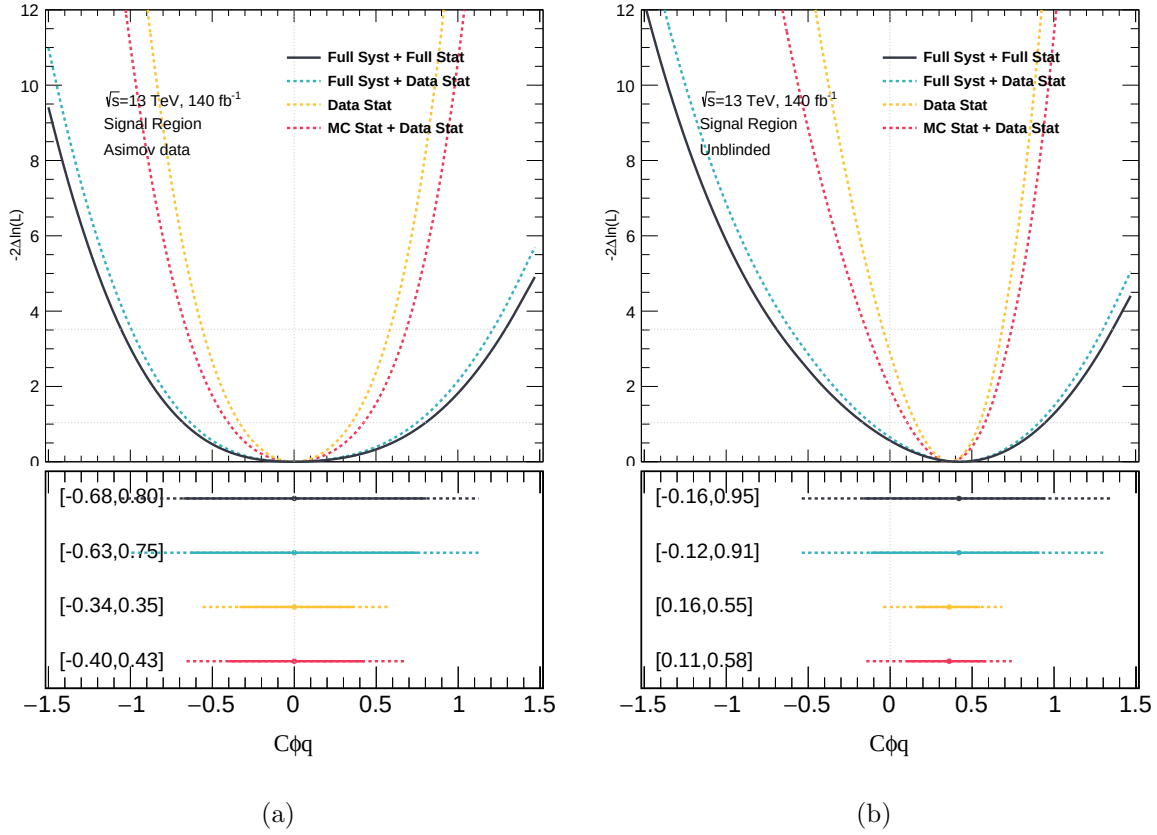


Figure 6.17: Likelihood scan of the parameter of interest $C_{\phi q}$ with all other parameters profiled obtained from a fit to Asimov data (a) and to real data (b). The different curves represent estimates obtained under varying combinations of statistical and systematic uncertainties. The solid line shows the results when considering all the possible uncertainties.

C_{qQ} (Four-fermion operator):

The C_{qQ} coefficient represents a four-quark contact interaction (in the t -channel context, an operator like $(\bar{q}q)(\bar{q}q)$ involving the third generation). Such an operator could provide an alternate production channel for single top quark: e.g. an $ub \rightarrow dt$ contact term that bypasses the W -boson exchange. This would modify the kinematics of top-quark production (potentially altering the p_T of the spectator jet or the angular distribution of the top-quark). However, some studies show that the specific O_{qq} , relevant to single-top-quark production, has a negligible effect on the angular distributions and spin polarisations at LO [250]. It would primarily manifest as an extra contribution to the cross-section.

The ATLAS experiment constrains C_{qQ} using single top-quark production cross-section measurement at 13 TeV at 95% confidence interval to be $-0.37 < C_{qQ} < 0.06 \text{ TeV}^{-2}$ [254].

The CMS Collaboration C_{qQ} measurement in the SMEFT framework using data collected in pp collisions at $\sqrt{s} = 13 \text{ TeV}$, corresponding to an integrated luminosity of 138 fb^{-1} . The 95% confidence interval is found to be $-0.04 < C_{qQ} < 0.03 \text{ TeV}^{-2}$ [253].

In this analysis, at $\sqrt{s} = 13 \text{ TeV}$ and an integrated luminosity of 140 fb^{-1} , the coefficient C_{qQ} is constrained through a global fit alongside six other EFT parameters and relevant nuisance terms. The fit to real data yields a 95% confidence interval of

$$0.22 < C_{qQ} < 0.46 \text{ TeV}^{-2},$$

indicating a clear deviation from zero within the current uncertainties. The likelihood scans, shown in Figure 6.18, illustrate both the expected and observed constraints on C_{qQ} , with the best-fit value favouring a nonzero coupling.

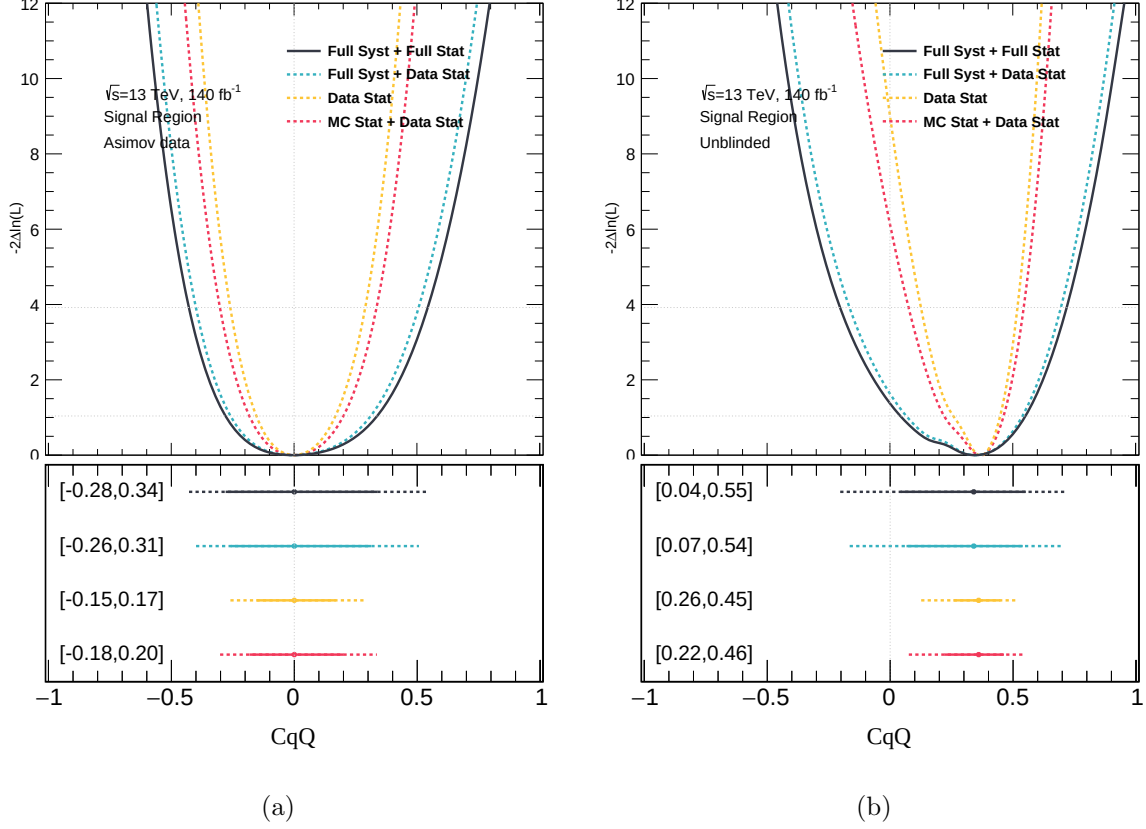


Figure 6.18: Likelihood scan of the parameter of interest C_{qq} with all other parameters profiled obtained from a fit to Asimov data (a) and to real data (b). The different curves represent estimates obtained under varying combinations of statistical and systematic uncertainties. The solid line shows the results when considering all the possible uncertainties.

6.8.1 Fit diagnostics and model validation

To ensure the validity of the parameter extraction and to understand the interplay between different components of the model, several post-fit diagnostic tools are employed. As explained in Section 6.7.4, the likelihood functions in this analysis often exhibit flat-bottomed profiles near the best-fit region. In such cases, traditional one-dimensional scans and uncertainties may underestimate or obscure correlations between parameters. Therefore, two-dimensional (2D) likelihood scans become particularly valuable, as they allow for a more comprehensive view of the joint behaviour of parameter pairs and help to visualise degeneracies or flat directions in the parameter space.

In addition to the 2D scans, the global correlation matrix is used to quantify the linear correlation structure among all parameters of interest and NPs as shown in Figure 6.19. Pull distributions provide insight into how NPs shift during the fit and how strongly

they are constrained by the data. Finally, ranking distributions are employed to assess the relative impact of individual systematic uncertainties on the signal strength or EFT coefficients. Together, these tools serve to validate the model, highlight the sensitivity structure of the analysis, and ensure that the extracted results are statistically and systematically well-understood.

2D Likelihood scans: To investigate potential correlations between different Wilson coefficients, two-dimensional likelihood scans are performed, extending the one-dimensional scans discussed earlier. These scans are constructed using the test statistic

$$-2\Delta \ln \mathcal{L}(C_i, C_j) = -2 \ln \left(\frac{\mathcal{L}(C_i, C_j, \hat{\hat{C}}_k[C_i, C_j])}{\mathcal{L}(\hat{C})} \right),$$

where $\hat{\hat{C}}_k[C_i, C_j]$ denotes the profiled values of all other parameters (nuisance and EFT) conditional on the given pair (C_i, C_j) , and $\hat{C}$ represents the global best-fit point. These scans provide a visual representation of how the data constrain pairs of parameters, capturing both their individual sensitivities and their mutual correlations.

The two-dimensional likelihood scans shown in Figure 6.20 display the correlations between selected pairs of Wilson coefficients after profiling all other parameters. In Figure 6.20(a), the coefficients C_{qQ} and $C_{\phi q}$ exhibit a clear positive correlation, indicated by the tilted elliptical shape of the confidence region. The observed correlation stems from the fact that both operators affect the left-handed quark current in the production vertex, leading to similar distortions in the event kinematics and angular distributions.

Figure 6.20(b) shows the likelihood scan in the plane of $C_{\phi tb}$ and C_{bW} , which also reveals a strong positive correlation. This arises due to overlapping effects these operators have on the top-quark decay structure, particularly through their influence on the angular observables sensitive to the tWb vertex. In both scans, the best-fit points (red stars) deviate from the SM expectation (blue triangles), though remain within the 95% confidence level. The elongated contours in both planes reflect degeneracies between the operator effects, where variations in one coefficient can be partially compensated by the other.

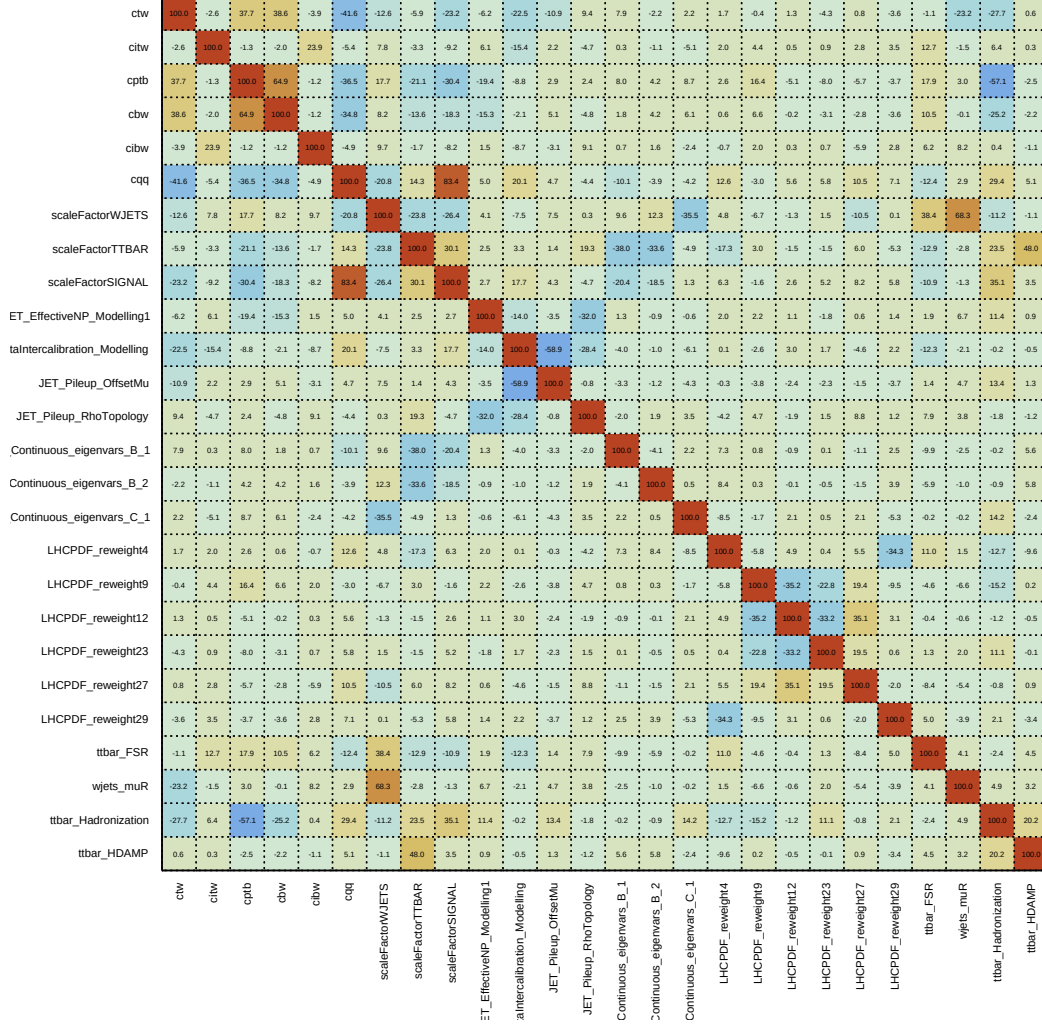


Figure 6.19: The correlation matrix displays the relationships between parameters with absolute correlations exceeding 30%. The colour gradient represents the strength and direction of the correlation: red indicates a perfect positive correlation (+100%), yellow corresponds to zero correlation (0%), and blue represents a perfect negative correlation (−100%).

Figure 6.21(a) displays the likelihood contour between C_{bW} and C_{ibW} . The near-circular shape of the contours suggests minimal correlation between the CP-conserving and CP-violating components of the b -quark dipole operator. This supports the interpretation that the imaginary and real parts are effectively constrained independently in the fit.

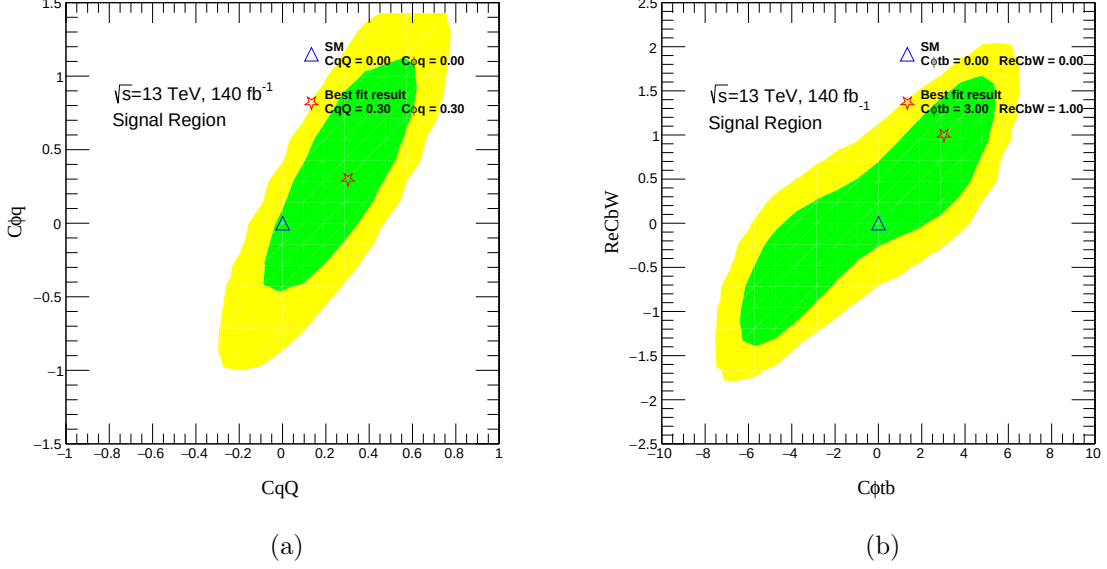


Figure 6.20: Likelihood scan of (a) C_{qQ} and $C_{\phi q}$, and (b) $C_{\phi tb}$ and C_{bW} with all other parameters profiled, obtained from real data, the best-fit results are marked with red stars, and the expected SM points are marked as the teal triangles.

Figure 6.21(b) shows the scan between C_{tW} and C_{bW} . A weak correlation is visible, reflecting their shared impact on the tWb vertex and decay kinematics. However, the difference in the nature of their contributions—dipole vs. helicity flip—ensures their effects remain distinguishable.

The scan of C_{bW} versus C_{qQ} , shown in Figure 6.22(a), exhibits largely uncorrelated behaviour. This is expected, since C_{bW} affects the decay via a dipole-like term, while C_{qQ} contributes through a four-fermion interaction that modifies the production process. The orthogonality of their physical effects is reflected in the shape of the likelihood contours.

Finally, Figure 6.22(b) presents the scan between $C_{\phi tb}$ and the scale factor for W +jets. A moderate correlation is observed, which may stem from overlapping sensitivity in kinematic CRs. Since $C_{\phi tb}$ can shift the helicity structure and event topology, it may affect how well W +jets events are separated and normalised, leading to the observed correlation in the fit.

These two-dimensional scans offer valuable insight into the interplay between various operators and highlight which parameter combinations are most constrained by the current data.

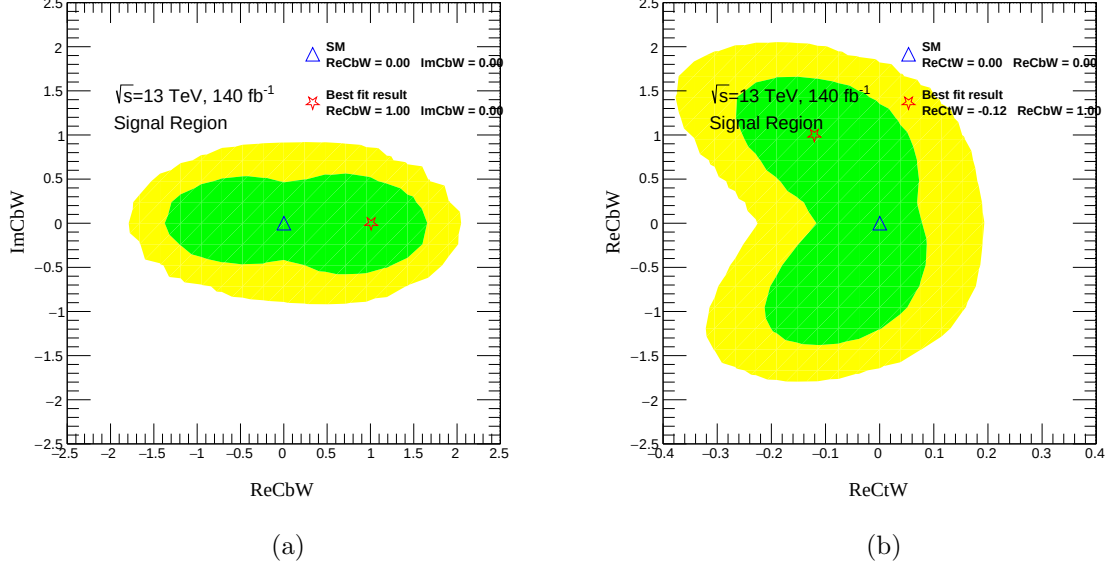


Figure 6.21: Likelihood scan of (a) C_{bW} and C_{ibW} , and (b) C_{tW} and C_{bW} with all other parameters profiled, obtained from real data, the best-fit results are marked with red stars, and the expected SM points are marked as the teal triangles.

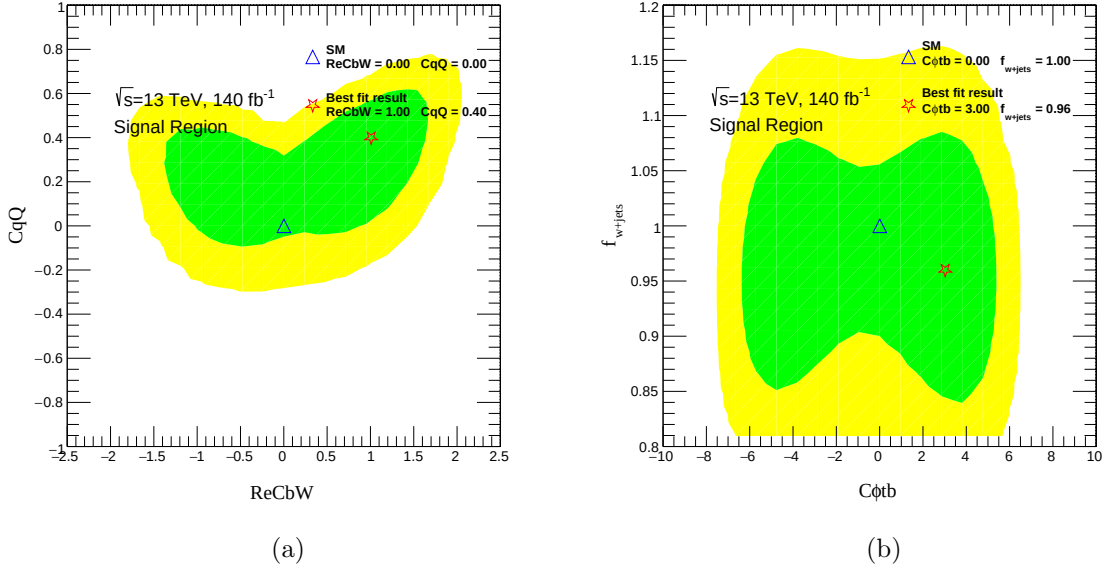


Figure 6.22: Likelihood scan of (a) $C_{\phi tb}$ and scale factor of the W +jets, and (b) C_{itW} and C_{bw} with all other parameters profiled, obtained from real data, the best-fit results are marked with red stars, and the expected SM points are marked as the teal triangles.

Double-minimum likelihood shapes in data: In this analysis, one-dimensional likelihood scans of a POI exhibit a double-minimum structure when using real data, whereas the same scans performed with an Asimov dataset yield a single, smooth minimum.

This indicates the presence of two distinct regions in parameter space that are locally preferred by the data. Such double-well structures often emerge when the observed data fluctuates around the SM prediction in a way that creates two competing solutions that describe different subsets of the data comparably well. Since the Asimov dataset represents an ideal, fluctuation-free expectation, it does not exhibit this behaviour.

Several effects can give rise to such structures:

- **Interference and degeneracy effects:** In EFT analyses, many operators contribute to the signal through both interference (linear in the Wilson coefficient) and quadratic terms. This can create a situation where two different values of the coefficient lead to nearly identical predicted yields or distributions, resulting in two local minima in the likelihood. This degeneracy is especially relevant when the interference term can change sign, and the likelihood lacks sensitivity to the direction of deviation.
- **Non-Gaussian fluctuations or category-wise tension:** In real data, statistical fluctuations in different regions or event categories can favour distinct parameter values. For instance, an excess in one region might prefer a positive coefficient value, while a deficit in another region favours a negative one. The global fit thus compromises between these regions, producing a bi-modal likelihood that reflects non-Gaussian behaviour, unlike the parabolic form expected from simple statistical models.
- **Complex likelihood landscapes:** The appearance of a second minimum may arise from the intrinsic complexity of the likelihood surface, especially in high-dimensional fits involving many NPs. The likelihood may feature multiple valleys or shallow plateaus, making it challenging for the optimiser to identify the true global minimum. This can lead to local minima being misidentified as best-fit points.

The observed double-minima in the likelihood scans for coefficients like C_{bW} in Figure 6.14(b) and $C_{\phi tb}$ in Figure 6.16(b) reflect these effects. The minima are asymmetric, and while sometimes one is slightly favoured, the difference in likelihood is often small enough that both remain statistically viable. This can result in wider or split confidence intervals, and suggests a need for careful interpretation. These features are important diagnostic tools in EFT fits, revealing either subtle physics effects such as interference and phase structure, or pointing to areas where the statistical model may need refinement.

Post-fit pulls and NP constraints: Pull distributions are a standard diagnostic tool used to evaluate the behaviour of NPs in a likelihood fit. Each pull represents the deviation of a fitted NP from its nominal value, normalised by its pre-fit uncertainty. It

is defined as

$$\text{Pull}(\theta_i) = \frac{\hat{\theta}_i - \theta_i^{\text{prefit}}}{\Delta\theta_i^{\text{prefit}}},$$

where $\hat{\theta}_i$ is the post-fit value, θ_i^{prefit} is the nominal (typically zero), and $\Delta\theta_i^{\text{prefit}}$ is the associated pre-fit uncertainty. A pull close to zero indicates that the data are consistent with the prior expectation, while larger deviations suggest tension or strong preference by the data. The post-fit uncertainty (or constraint) reflects how well each NP is constrained: values below one indicate a reduction in uncertainty due to the data, whereas values near or above one suggest little constraint. The post-fit pulls for representative NP groups are shown in Figures 6.23–6.25.

Figure 6.23(a) shows the pulls for the PDF uncertainty group, which includes 30 LHAPDF eigenvector variations as explained in Section 6.5. These NPs are only mildly constrained, with most pull values clustering around zero and symmetric uncertainty bands. This behaviour is expected, as PDF variations tend to affect overall rate predictions and are typically weakly constrained by shape-sensitive observables.

In contrast, the JER uncertainties, shown in Figure 6.23(b), exhibit stronger post-fit constraints. The fit significantly narrows the uncertainties on several effective NP components, especially affecting the most important central jet resolution. This suggests sensitivity in the data to resolution effects, which influence the shapes of kinematic distributions and angular variables used in the fit.

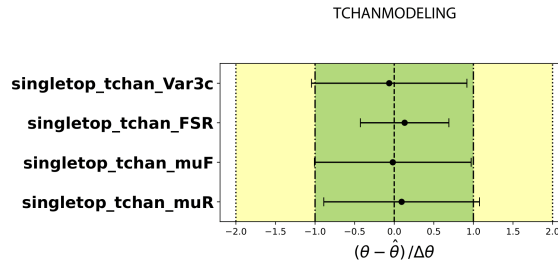


Figure 6.24: Post-fit pulls for selected NP for the t -channel modelling. Each point shows the shift of a NP from its nominal value in units of its pre-fit uncertainty. A pull near zero indicates agreement with the prior, while significant deviations suggest data-driven preference. The uncertainty bars represent post-fit uncertainties: values below one indicate the NP is constrained by the data; values near or above one indicate little constraint.

Pile-up-related uncertainties are presented in Figure 6.23(c). These include systematic variations related to the modelling of average interactions per bunch crossing, jet-area pile-up corrections, and topology-based reweighting. The fit exhibits modest constraint

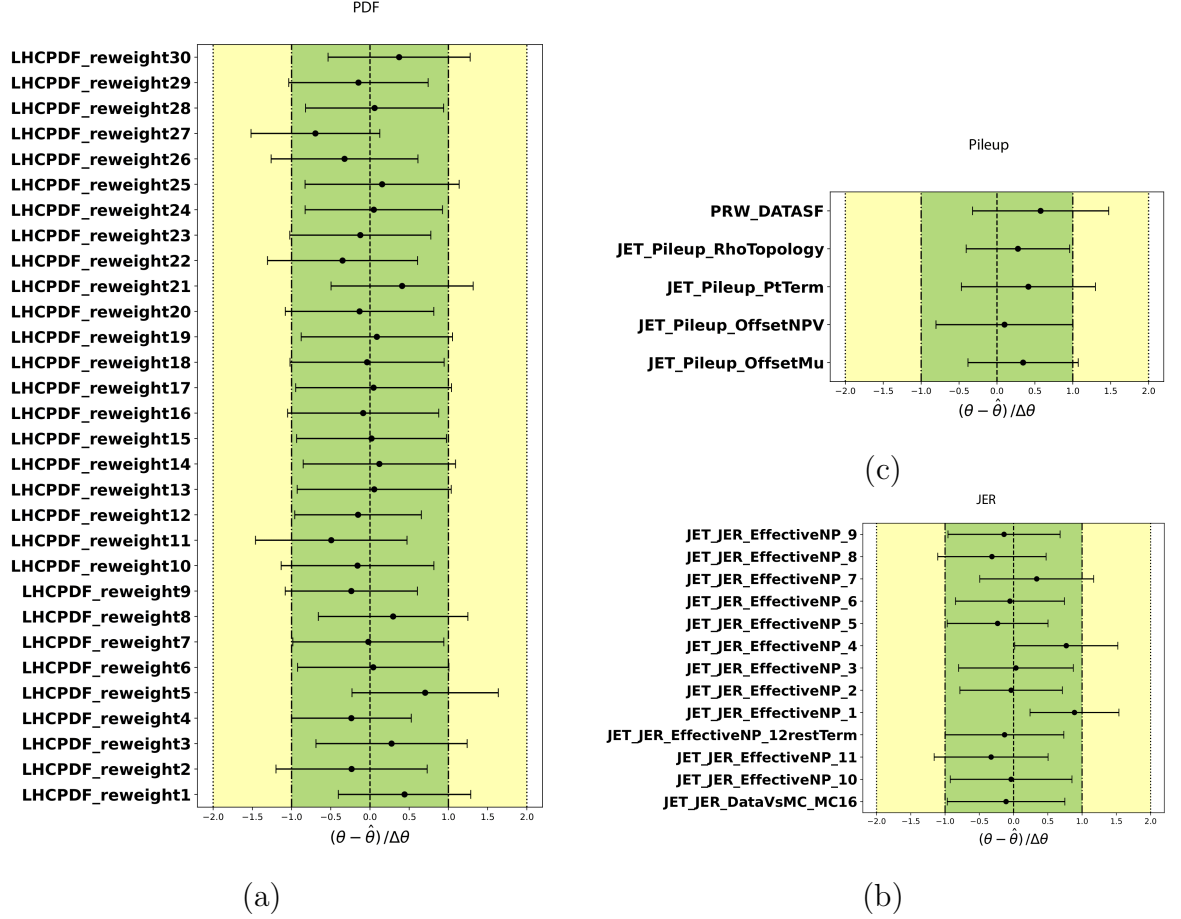


Figure 6.23: Post-fit pulls for selected NP groups: (a) PDF uncertainties; (b) JER systematics; (c) pile-up systematics. Each point shows the shift of a NP from its nominal value in units of its pre-fit uncertainty. A pull near zero indicates agreement with the prior, while significant deviations suggest data-driven preference. The uncertainty bars represent post-fit uncertainties: values below one indicate the NP is constrained by the data; values near or above one indicate little constraint.

on these components, with no significant pull from the nominal, indicating that pile-up effects are well-modelled and not a limiting factor in this analysis.

Figure 6.24 shows the modelling-related NPs for $t\bar{t}$ production. The pulls for variations such as FSR, ISR, and the A14 tune parameter **Var3c** reflect moderate constraint and in some cases slight shifts from the prior. Notably, the FSR component shows a visible pull, likely driven by angular and jet multiplicity observables sensitive to radiation effects.

The final pull summaries are shown in Figure 6.25. Figure 6.25(a) contains the W +jets modelling systematics, including variations of the renormalisation and factorisation scales. These are slightly constrained, reflecting the relatively minor contribution of W +jets backgrounds in the SRs. Figure 6.25 (c) provides a comprehensive view of the JES

systematic components. Some of the JES NPs are pulled and tightly constrained by the fit, particularly the one affecting η -intercalibration. This indicates that the fit is highly sensitive to the JES modelling across different detector regions and jet types.

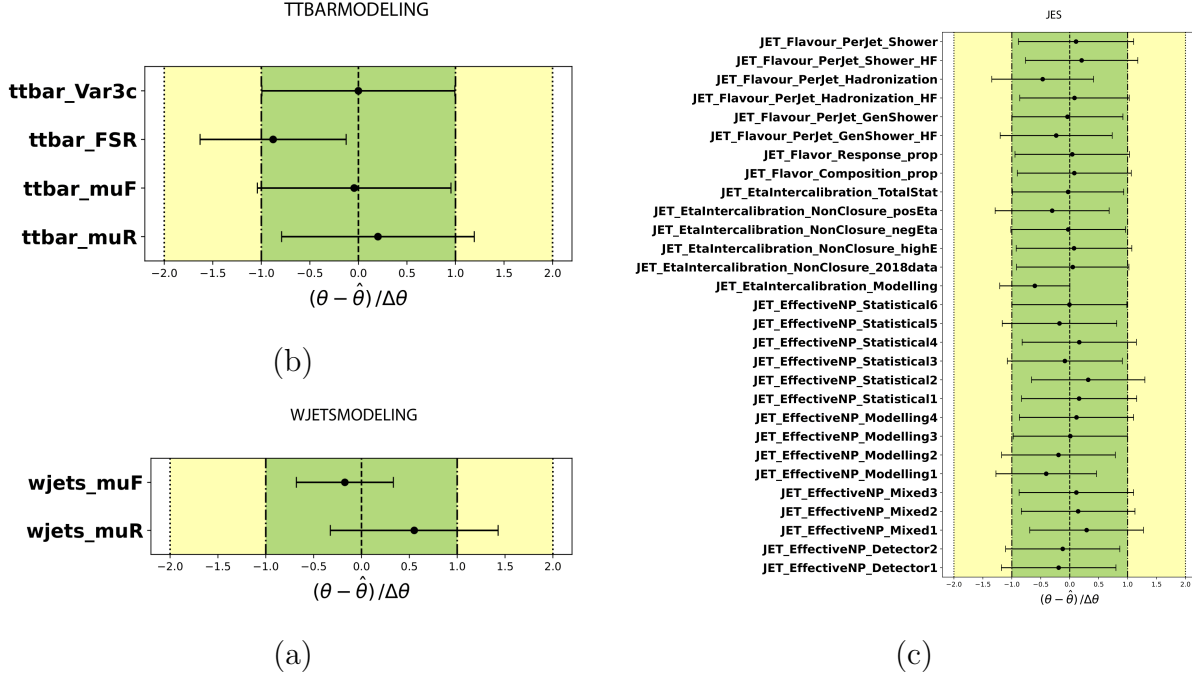


Figure 6.25: Post-fit pulls for NP groups: (a) W +jets modelling; (b) $t\bar{t}$ modelling; and (c) JES systematics. Each point shows the shift of a NP from its nominal value in units of its pre-fit uncertainty. A pull near zero indicates agreement with the prior, while significant deviations suggest data-driven preference. The uncertainty bars represent post-fit uncertainties: values below one indicate the NP is constrained by the data; values near or above one indicate little constraint.

The final pull summaries are shown in Figure 6.25. Figure 6.25(a) contains the W +jets modelling systematics, including variations of the renormalisation and factorisation scales. These are slightly constrained, reflecting the relatively minor contribution of W +jets backgrounds in the SRs. Figure 6.25 (c) provides a comprehensive view of the JES systematic components. Some of the JES NPs are pulled and tightly constrained by the fit, particularly the one affecting η -intercalibration. This indicates that the fit is highly sensitive to the JES modelling across different detector regions and jet types.

Overall, the pull distributions demonstrate that the fit is consistent and quite stable, with the strongest constraints appearing in JES and JER systematics, and modelling NPs exhibiting data-driven shaping in key processes like $t\bar{t}$. A more complete collection of pull distributions covering additional systematic groups is provided in Appendix H.

Ranking of systematic uncertainties: Ranking distributions provide a quantitative

measure of the sensitivity of the fitted Wilson coefficients to individual systematic uncertainties. These distributions are constructed after the fit is performed and reflect how each NP influences the POI, such as C_{tW} and C_{itW} . For each systematic group, the impact is evaluated by shifting the associated NP up and down by its post-fit uncertainty while re-optimising all other parameters in the likelihood. The resulting variation in the POI is shown as a blue box on the upper axis, thereby indicating the contribution of that particular uncertainty to the total uncertainty of the fit.

The lower part of the distribution shows the pulls of each NP, defined as the deviation of the fitted value from its pre-fit expectation, normalised by the pre-fit uncertainty. These are displayed as filled circles, providing insight into how strongly the data constrain each NP. The black uncertainty bars around the pulls correspond to the post-fit uncertainties relative to their nominal (pre-fit) values. A horizontal dotted line at unity serves as a reference, where post-fit uncertainties equal their initial values. Deviations from this line indicate either constraints from the data (below unity) or weakly constrained NPs (above unity). Together, the ranking distributions highlight both the dominant sources of uncertainty and how well each NP is constrained in the fit. Figures 6.26 to 6.29 present the ranking distributions for all EFT parameters considered in this analysis, summarising the relative impact of grouped systematic uncertainties on each fit result.

From the ranking distributions in Figure 6.26, it is evident that the leading uncertainties affecting C_{tW} and C_{itW} arise from JES, JER, and MC statistical uncertainties. In particular, the JES-related components cause noticeable shifts in both real and imaginary parts of C_{tW} , reflecting the sensitivity of angular observables to jet kinematics. Similarly, Figure 6.27 shows that the dominant contributions to the uncertainties in C_{bW} and C_{ibW} also originate from MC statistics, along with most sensitivity to JES and missing ET, respectively. These results are consistent with the subdominant nature of b -quark dipole effects in single-top-quark processes, where precision is largely driven by experimental inputs. For $C_{\phi tb}$ and $C_{\phi Q}$, shown in Figure 6.28, jet-related uncertainties dominate, though in the case of $C_{\phi tb}$ there is a notable impact from JER. The ranking distribution for C_{qQ} (see Figure 6.29) reveals that the coefficient is especially affected by PDF and MC statistics, consistent with its role in scaling the single top-quark production cross-section, which is highly sensitive to the jet acceptance and normalisation.

Together, these ranking distributions provide valuable insight into the dominant sources of systematic uncertainty for each EFT coefficient. They also highlight which NP groups would benefit from further refinement to improve the overall EFT sensitivity.

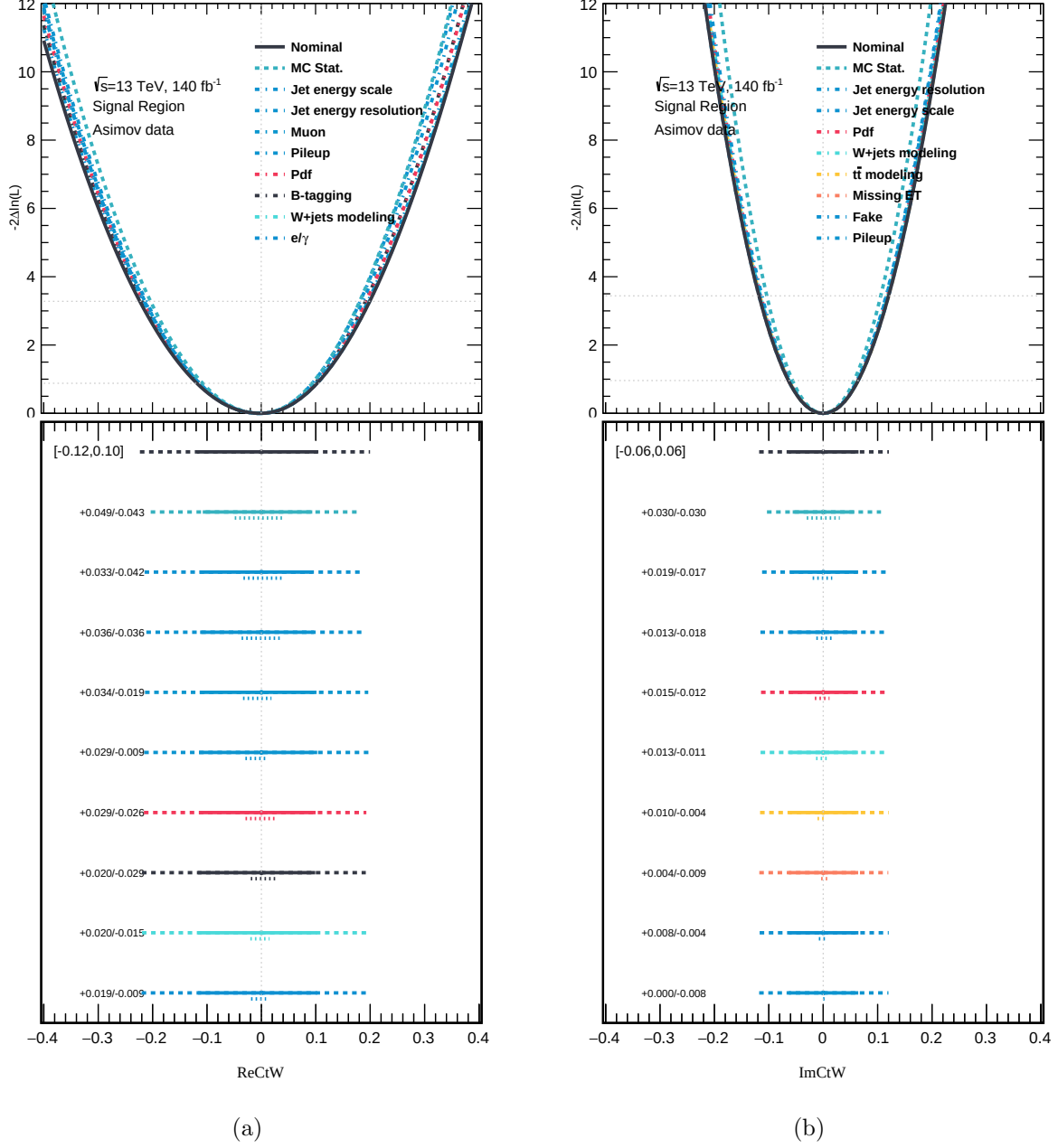


Figure 6.26: Ranking distributions for (a) C_{tW} and (b) C_{itW} , showing the impact of grouped systematic uncertainty categories on the signal strength. The vertical axis lists grouped sources (e.g. JES, PDF, MC stat.), and the horizontal bars indicate the shift in the best-fit value of the POI when each source is varied by $\pm 1\sigma$, while profiling all other parameters.

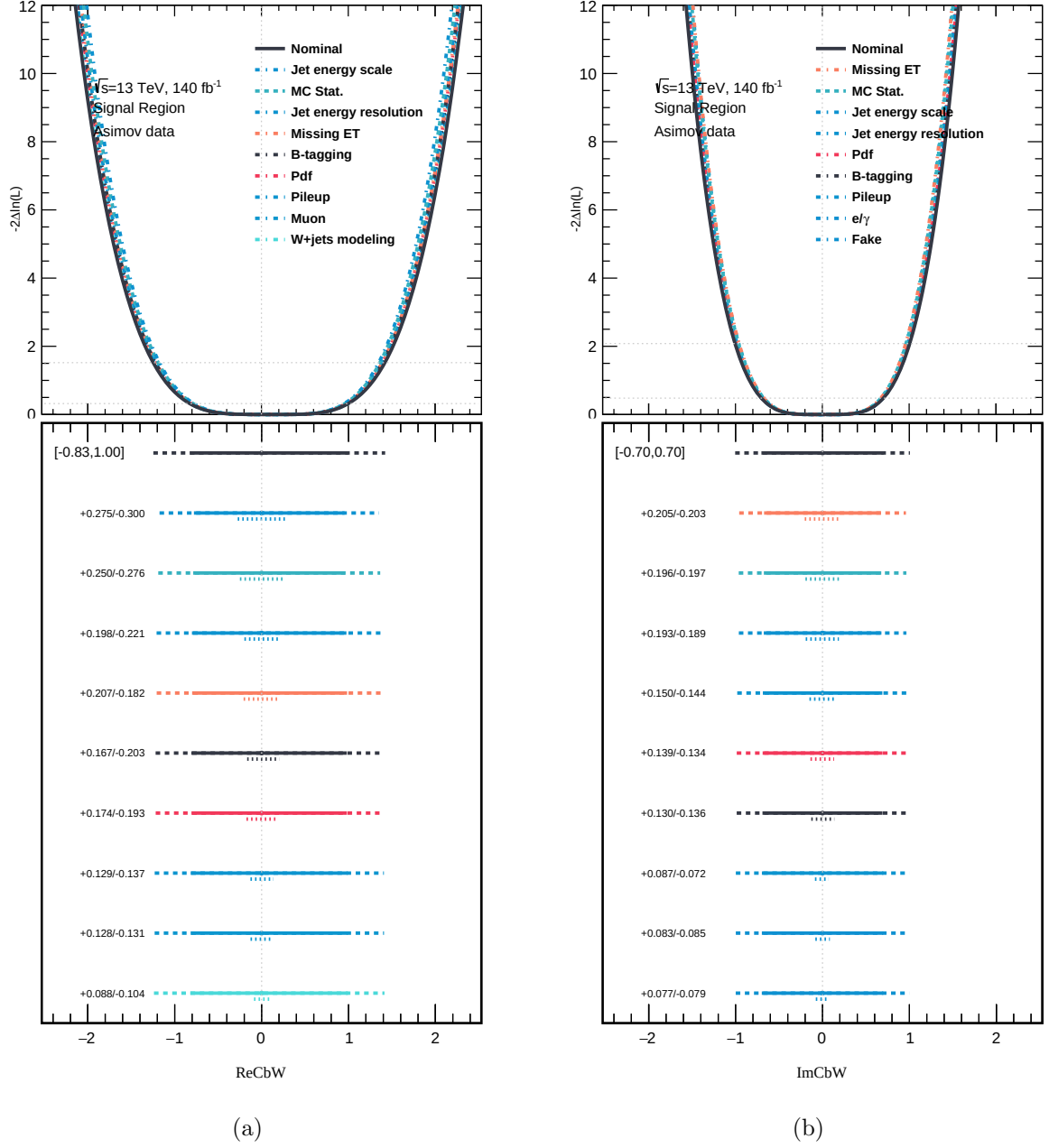


Figure 6.27: Ranking distributions for (a) C_{bW} and (b) C_{ibW} . The vertical axis lists grouped sources (e.g. JES, PDF, MC stat.), and the horizontal bars indicate the shift in the best-fit value of the POI when each source is varied by $\pm 1\sigma$, while profiling all other parameters.

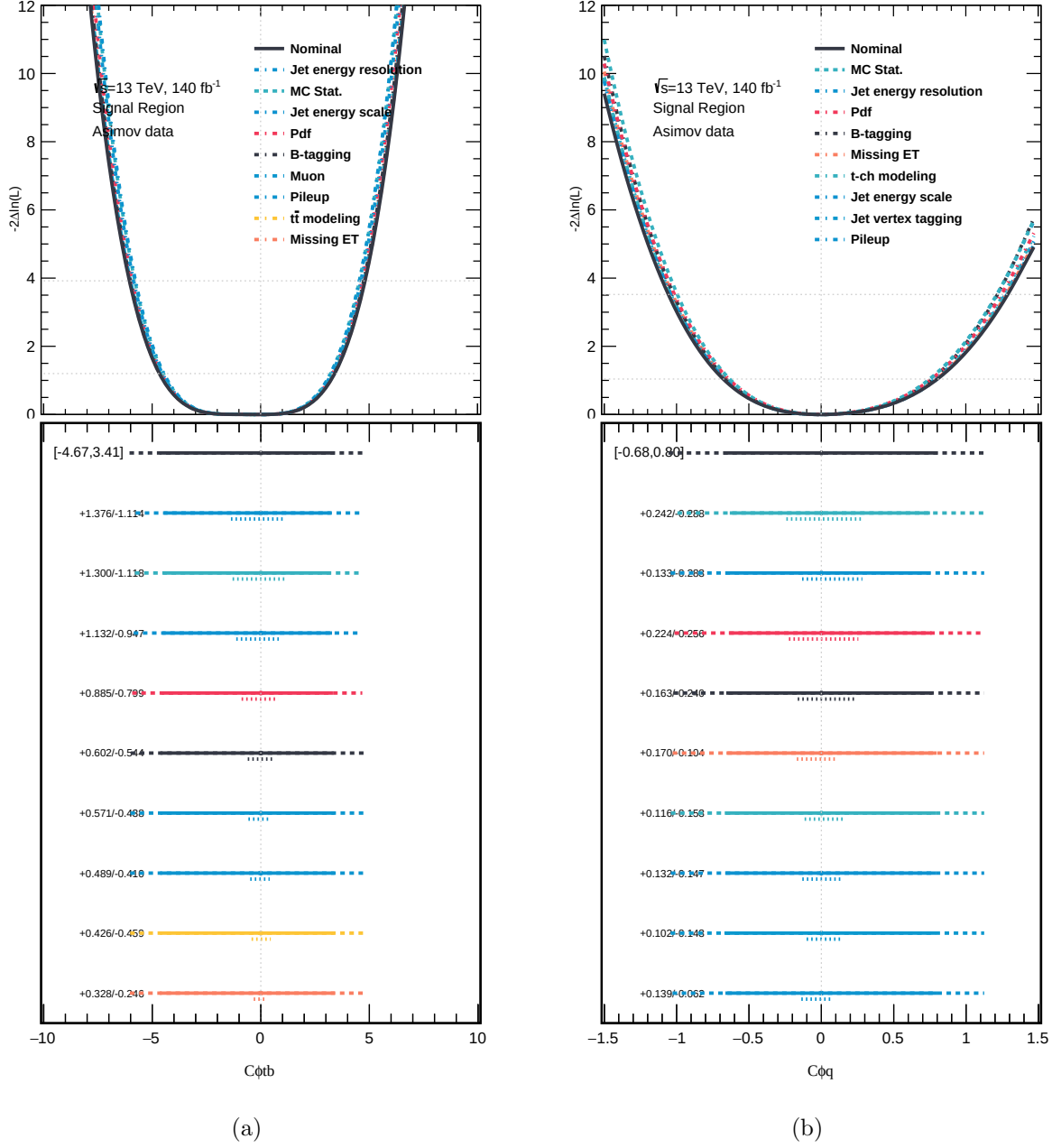


Figure 6.28: Ranking distributions for (a) $C_{\phi tb}$ and (b) $C_{\phi Q}$, illustrating the leading sources of systematic uncertainty for each parameter. The vertical axis lists grouped sources (e.g. JES, PDF, MC stat.), and the horizontal bars indicate the shift in the best-fit value of the POI when each source is varied by $\pm 1\sigma$, while profiling all other parameters.

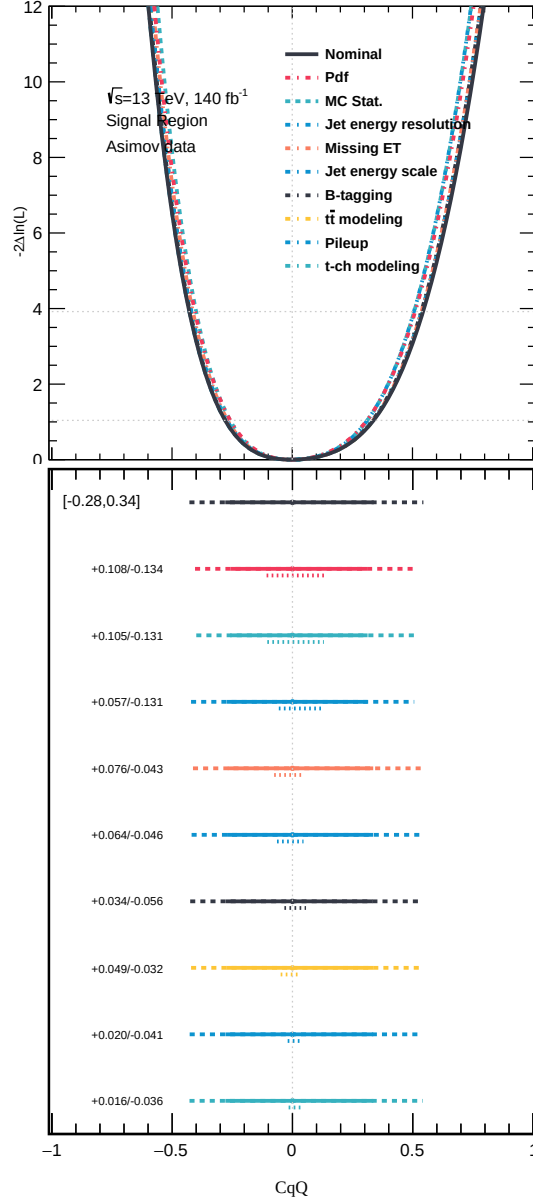


Figure 6.29: Ranking distribution for $C_{QQ}^{(3)}$, showing the impact of dominant experimental and theoretical uncertainties on the extracted coefficient. The vertical axis lists grouped sources (e.g. JES, PDF, MC stat.), and the horizontal bars indicate the shift in the best-fit value of the POI when each source is varied by $\pm 1\sigma$, while profiling all other parameters.

The ranking distributions shown are based on Asimov datasets using the $\mathcal{O}(1/\Lambda^2)$ approximation, rather than real data, as the extraction of final constraints from observed data lies outside the scope of this thesis, being done recently by the ATLAS collaboration.

6.9 Conclusions

The analysis done within this thesis delivers a comprehensive and state-of-the-art study of single top-quark production in the t -channel within the framework of the SMEFT, using the full angular information from the decay of the top quark. By employing a fully differential fit to angular observables, stringent constraints are placed on seven dimension-six Wilson coefficients affecting the tWb vertex ($C_{tw}, C_{itw}, C_{bw}, C_{ibw}, C_{\phi q}, C_{\phi tb}, C_{qQ}$), along with associated NPs and normalisation factors. The methodology incorporated both event rates and shape information, enabling an unprecedented sensitivity to subtle deviations from SM predictions.

This work represents a methodological advancement in SMEFT analyses of top-quark production, notably through the development and application of the M-function formalism for extracting angular coefficients from fully differential decay distributions [247]. By expanding the differential decay rate into a basis of orthonormal angular functions, the analysis enables a simultaneous fit to all angular observables, preserving full kinematic information. The use of the morphing technique to model the dependence of signal templates on EFT coefficients allows for continuous interpolation between different operator values, ensuring smooth and accurate template construction across parameter space. In addition, systematic uncertainties are incorporated through a high-dimensional likelihood with over 200 NPs, profiled simultaneously in the fit. These methodological components collectively enable a robust and model-independent extraction of EFT constraints and establish a blueprint for future high-precision top-quark analyses in the SMEFT framework.

Crucially, the analysis probed the structure of the tWb vertex in both its real and imaginary components. The tight constraint on the dipole coefficient C_{tW} , as well as its CP-violating counterpart C_{itW} , marks a major improvement over previous results. These results reaffirm the SM expectation of a V–A structure with no evidence of new physics, but place its potential deviations under tight scrutiny: the couplings are now constrained at the few-percent level or $\mathcal{O}(10^{-2})$ TeV^{-2} in the EFT expansion. The corresponding numerical limits are summarised in Table 6.3, which presents the one-sigma confidence intervals on all relevant EFT coefficients—including both real and imaginary parts of C_{tW} and C_{bW} , as well as $C_{\phi tb}$, $C_{\phi q}$, and C_{qQ} —across multiple fit configurations accounting for statistical and systematic uncertainties.

These results can be compared to the ATLAS+CMS combination, which as already discussed in the Section 2.3.2 provides 95% confidence intervals on individual EFT coefficients such as $C_{\phi tb} \in [-3.48, 5.16]$ and $C_{bW} \in [-0.96, 0.67]$. These limits primarily

C/Λ^2	MC Stat. + Data Stat.	Data Stat.	Full Syst. + Data Stat.	Full Syst. + Full Stat.
$\text{Re}C_{tW}$	$[-0.18, -0.06]$	$[-0.17, -0.07]$	$[-0.23, -0.02]$	$[-0.22, 0.01]$
$\text{Im}C_{tW}$	$[-0.04, 0.06]$	$[-0.04, 0.05]$	$[-0.06, 0.04]$	$[-0.07, 0.05]$
$\text{Re}C_{bW}$	$[0.65, 0.99]$	$[0.68, 0.99]$	$[0.68, 1.25]$	$[0.61, 1.25]$
$\text{Im}C_{bW}$	$[-0.11, 0.22]$	$[-0.08, 0.20]$	$[-0.30, 0.20]$	$[-0.29, 0.25]$
$C_{\phi tb}$	$[2.18, 3.85] \cup [-5.17, -4.00]$	$[2.34, 3.74] \cup [-5.00, -4.20]$	$[0.91, 4.94] \cup [-5.90, -3.00]$	$[0.34, 5.04] \cup [-6.06, -3.00]$
$C_{\phi q}$	$[0.11, 0.58]$	$[0.16, 0.55]$	$[-0.12, 0.91]$	$[-0.16, 0.95]$
C_{qQ}	$[0.04, 0.55]$	$[0.07, 0.54]$	$[0.26, 0.45]$	$[0.22, 0.46]$

Table 6.3: One sigma confidence interval limits on EFT coefficients in units of C/Λ^2 (TeV^{-2}) for different fit configurations.

constrain one linear direction in the $(C_{\phi tb}, C_{bW})$ plane, aligned with contributions to the right-handed W boson helicity fraction F_R . The global fit is effectively blind in the orthogonal direction due to a lack of complementary observables. In contrast, the two-dimensional likelihood contour obtained (see Figure 6.20(b)) in this analysis shows sensitivity not only along the F_R -dominated axis, but also retains non-negligible resolving power in the orthogonal direction. This is a direct consequence of using angular observables in a multidimensional fit, which access spin correlations and chirality structures beyond inclusive or single-observable measurements. As such, this result provides complementary constraints that help lift flat directions present in global combinations.

The improvement in sensitivity achieved through this analysis is illustrated in Figure 6.30, which compares the allowed regions for the anomalous tWb couplings in the complex plane of $\text{Re}[g_R/V_L]$ vs. $\text{Im}[g_R/V_L]$. The figure shows progressively tighter confidence intervals obtained from earlier ATLAS measurements using two and three angular observables at 7 TeV [95] and 8 TeV [96], respectively, culminating in the present 13 TeV result, which incorporates the full set of angular variables describing the kinematics of the top-quark decay and achieves an improvement of approximately one order of magnitude in sensitivity.

As mentioned previously in Section 2.4.2, it is worth noting that one-loop EW corrections in the SM induce a small but non-zero effective g_R , with a complex phase. When translated into the EFT coefficient C_{tW} using the matching relation in Equation 2.12, this leads to a value of $C_{tW} \approx -0.014$ based on Ref. [255], or up to -0.03 according to other evaluations [256, 257]. Compared to the current statistical uncertainty on C_{tW} in this analysis, this effect is at the level of one-quarter to one-half sigma. Since the imaginary part of g_R enters T-odd observables and appears with opposite sign in top-quark and top-antiquark decays, this contribution largely cancels in CP-even observables. However, given the precision expected in future datasets, especially at HL-LHC, such SM loop-induced contributions may need to be included in the MC baseline or subtracted

analytically to avoid bias in EFT interpretations.

A key distinction of the current result is the simultaneous fit of all six relevant dimension-six EFT operators affecting the tWb vertex, including complex phases. In particular, this fit disentangles the contribution of the four-fermion operator C_{qQ} .

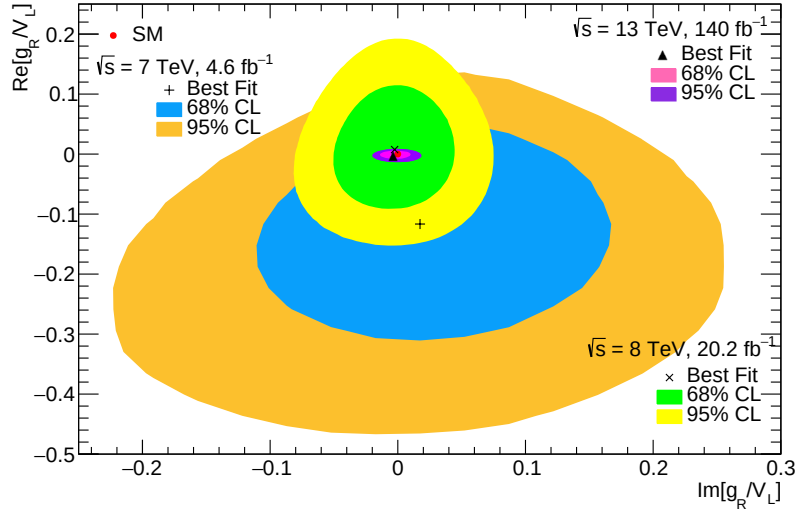


Figure 6.30: Summary of constraints on the anomalous tWb couplings in the complex plane of the ratio $\text{Im}[g_R/V_L]$ vs $\text{Re}[g_R/V_L]$, where V_L is the SM left-handed coupling and g_R is the right-handed coupling. Shown are the 68% (inner) and 95% (outer) confidence regions obtained using two angular observables (blue/orange), three angles (green/yellow), and the full four-angle analysis (pink/purple). These correspond to the progressively tighter constraints from earlier 7 TeV (4.6 fb^{-1}) [95] and 8 TeV (20.2 fb^{-1}) [96] ATLAS measurements to the new 13 TeV result. The markers (+, $\times$, $\blacktriangle$) indicate the best-fit point for each case, and the red dot denotes the SM expectation at (0,0). The dramatic reduction in the allowed region with additional observables highlights how including all four decay angles significantly improves sensitivity to both the real and imaginary parts of g_R .

This thesis demonstrates that fully exploiting the angular structure of top-quark decays, combined with a rigorous SMEFT interpretation, provides an exceptionally sensitive probe of new physics in the EW top-quark sector. The integration of innovative methods—such as the M-function formalism and morphing interpolation—enables a level of precision and generality not previously achieved in top-quark EFT analyses. The constraints obtained for some of the EFT parameters are among the most stringent to date, reinforcing the validity of the SM while significantly narrowing the window for potential deviations.

Looking forward, the techniques developed here pave the way for future analyses

using higher-statistics datasets from Run 3 and from the future runs at the HL-LHC. If deviations from the SM do exist in the charged-current interaction of the top quark, this framework is well-equipped to uncover them. Until then, the results presented in this work stand as a benchmark for precision and completeness in SMEFT-based top-quark analyses.

Appendix

A Sagitta bias effect on opposite signed muons transverse momentum

This appendix section presents comparative distributions of p_T for two muons. The δ_s bias is known to be proportional to the p_T , significantly altering the momentum of muons when a bias is introduced. Figure 31 compares the p_T distributions of two muons in scenarios: one with an original MC event sample and another with an injected sagitta bias. These distributions clearly demonstrate how the sagitta bias influences the p_T of muons.

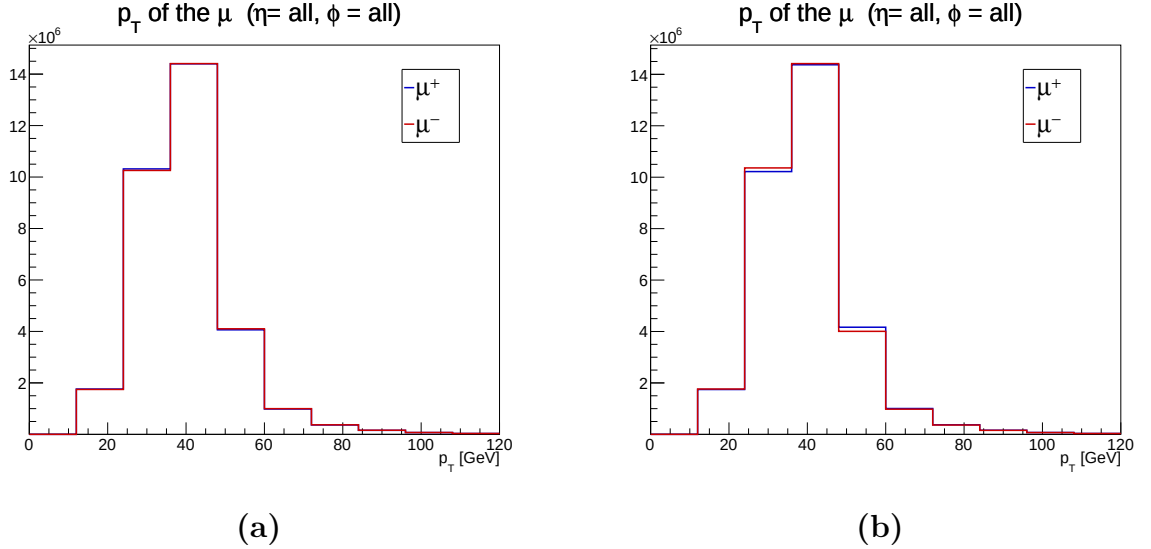


Figure 31: Comparison of the transverse momentum (p_T) distributions for oppositely charged muons. Figure 12(a) shows the p_T distributions for a pure MC generated sample without any injected bias. Figure 12(b) shows the distributions after introducing an artificial sagitta bias. The observed separation between positive and negative muon distributions highlights the dependence of the transverse momentum measurement on sagitta distortions.

The second example explores the impact of injecting varying artificial biases. Figure 32(a) demonstrates how a negative bias shifts the transverse momentum distribution of negatively charged muons to be softer (i.e. to have lower p_T). Conversely, a positive bias, as shown in Figure 32 (b), shifts the distribution in the opposite direction (i.e. to be harder, higher p_T). This illustrates the directional dependence of muon transverse momentum on the nature of the injected bias.

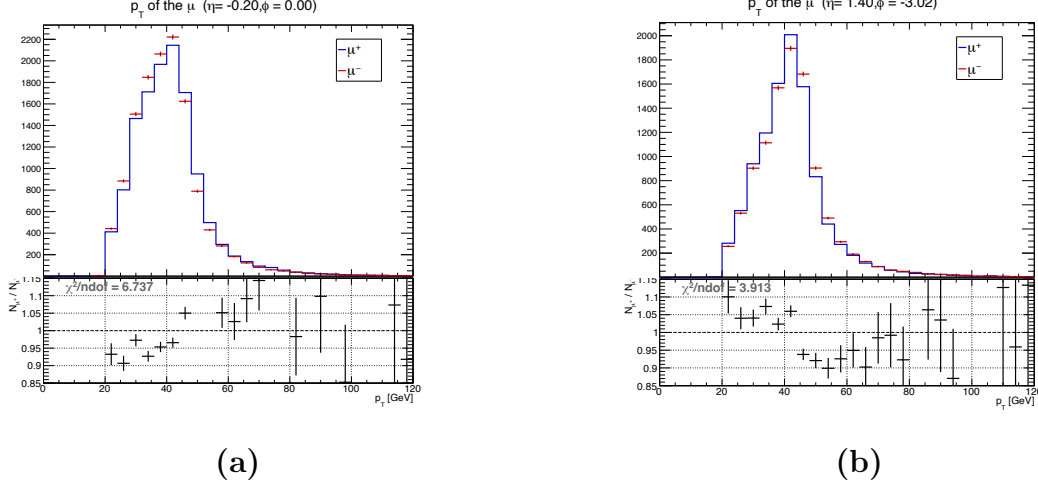


Figure 32: Comparison of the transverse momentum (p_T) distributions for oppositely charged muons within a single (η, ϕ) bin under injected artificial biases. Figure 13(a) shows the result after applying a negative bias (approximately $-0.2 \times 10^{-3} \text{ GeV}^{-1}$), where the transverse momentum distribution of negatively charged muons shifts to lower values. Figure 13 (b) presents the case with a positive injected bias (approximately $+0.15 \times 10^{-3} \text{ GeV}^{-1}$), resulting in a shift of the distribution of the positively charged muons toward higher transverse momentum values.

B Testing of the VarMin method with Y-shaped injected bias

This appendix details the testing of a new method using a “Y” shaped bias, with a strength of 0.4 TeV^{-1} . As noted in the main text, even a pure MC event sample is not devoid of bias, contributing noise to the sagitta measurement map. Here, the subtraction of MC simulated noise from the measurement over several iterations is demonstrated. For instance, Figure 36 illustrates the subtraction of a ten-iteration sagitta correction map from a corresponding ten-iteration MC simulated sample map, while figures Figures 33 to 35, for one, two and three-iterations. It is crucial to subtract the measured sagitta map from the pure MC simulated sagitta map for an equal number of iterations to ensure accurate convergence.

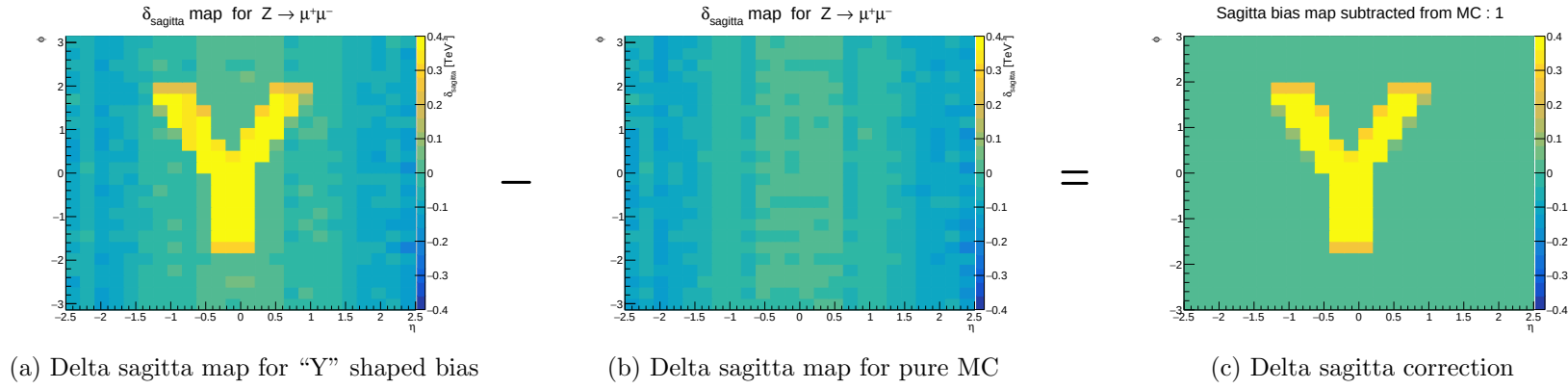


Figure 33: Sagitta bias map for pure MC simulated subtracting from the DS with “Y” shaped bias injected. The Figure 34c is the correction value, which is what was injected in the first place, 0.4 for bins in “Y” shape and flat zero as a background. The maps correspond to one iterations output.

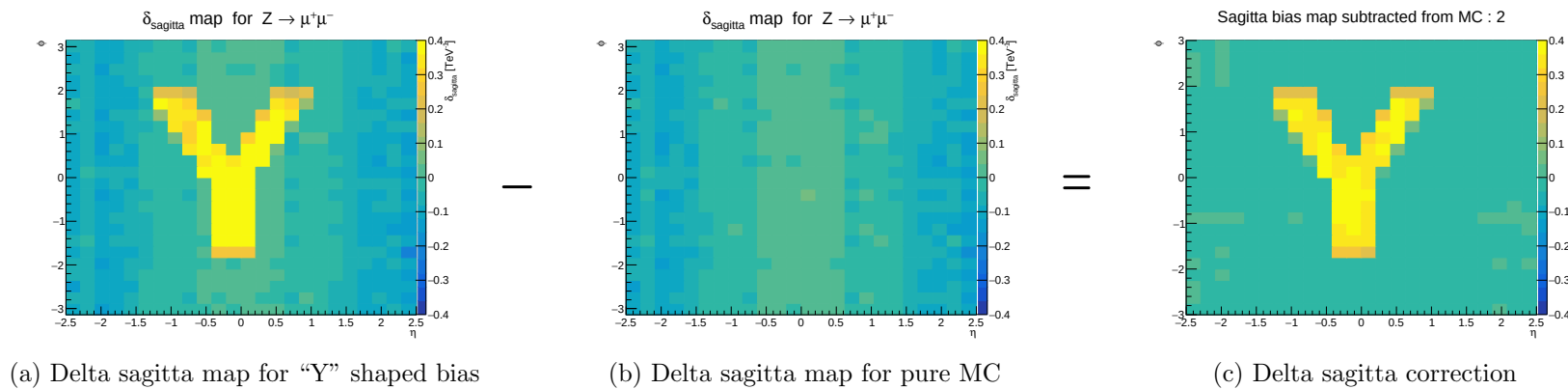


Figure 34: Sagitta bias map for pure MC simulated subtracting from the DS with “Y” shaped bias injected. The Figure (c) is the correction value, which is what was injected in the first place, 0.4 for bins in “Y” shape and flat zero as a background. The maps correspond to 2 iterations output.

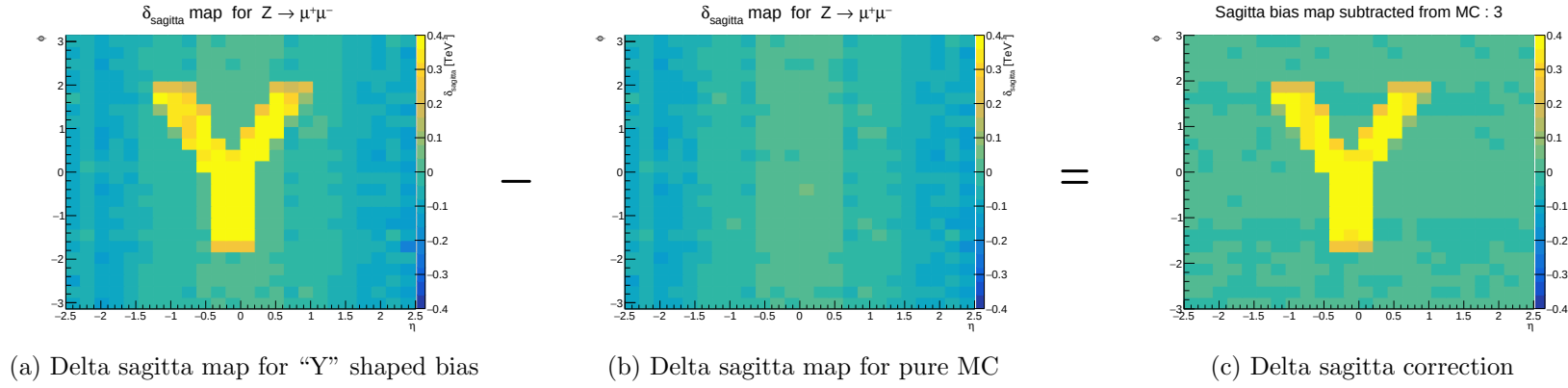


Figure 35: Sagitta bias map for pure MC subtracting from the DS with “Y” shaped bias injected. The Figure (c) is the correction value, which is what was injected in the first place, 0.4 for bins in “Y” shape and flat zero as a background. The maps correspond to 3 iterations output.

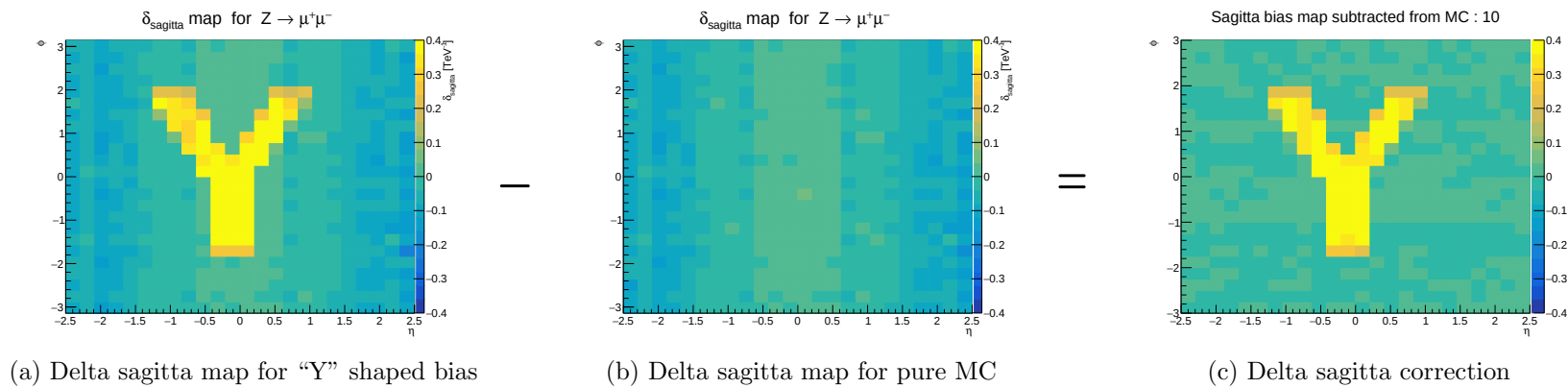


Figure 36: Sagitta bias map for pure MC subtracting from the DS with “Y” shaped bias injected. The Figure (c) is the correction value, which is what was injected in the first place, 0.4 for bins in “Y” shape and flat zero as a background. The maps correspond to 10 iterations.

The final analytical component involves Z projection images of the sagitta bias map before MC simulated subtraction (as shown in Figures 37 and 38). These projections allow observation of the background distribution and the injected bins across one to ten iterations.

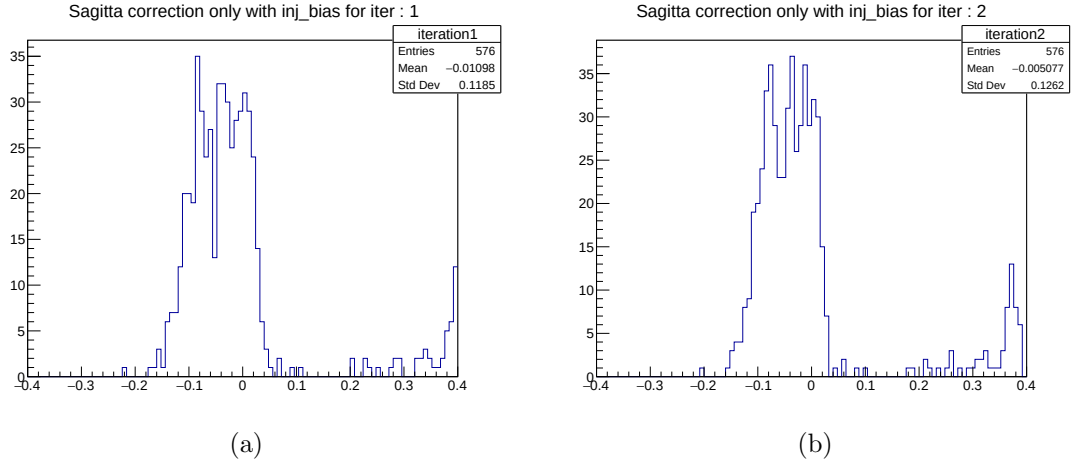


Figure 37: Distributions (a) and (b) are Z projections of the sagitta bias maps from Figures 33(a) and 34(a).

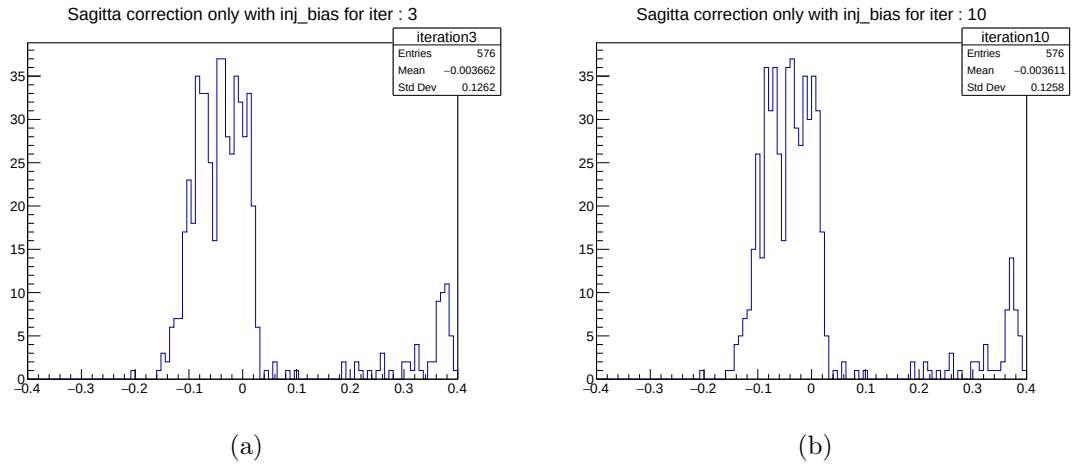


Figure 38: Distributions (a) and (b) are Z projections of the sagitta bias maps from Figures 35(a) and 36(a).

C Morphing functions for the top-quark and anti-top-quark channel

In this appendix, all 23×2 normalized coefficients as a one-dimensional function of the Wilson coefficients are shown.

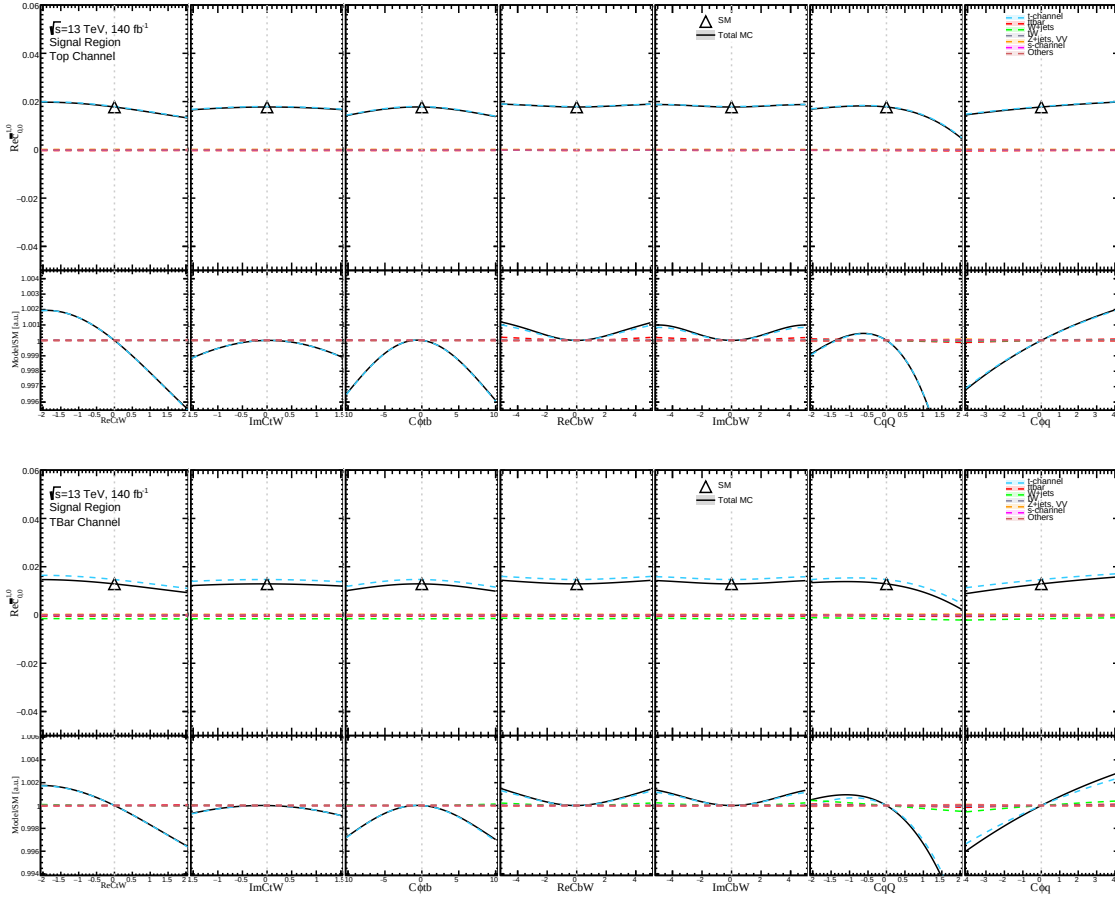


Figure 6.39: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re} c_{0,0}^{1,0} \rangle^{\pm, \text{SR}}(C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses overlap each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

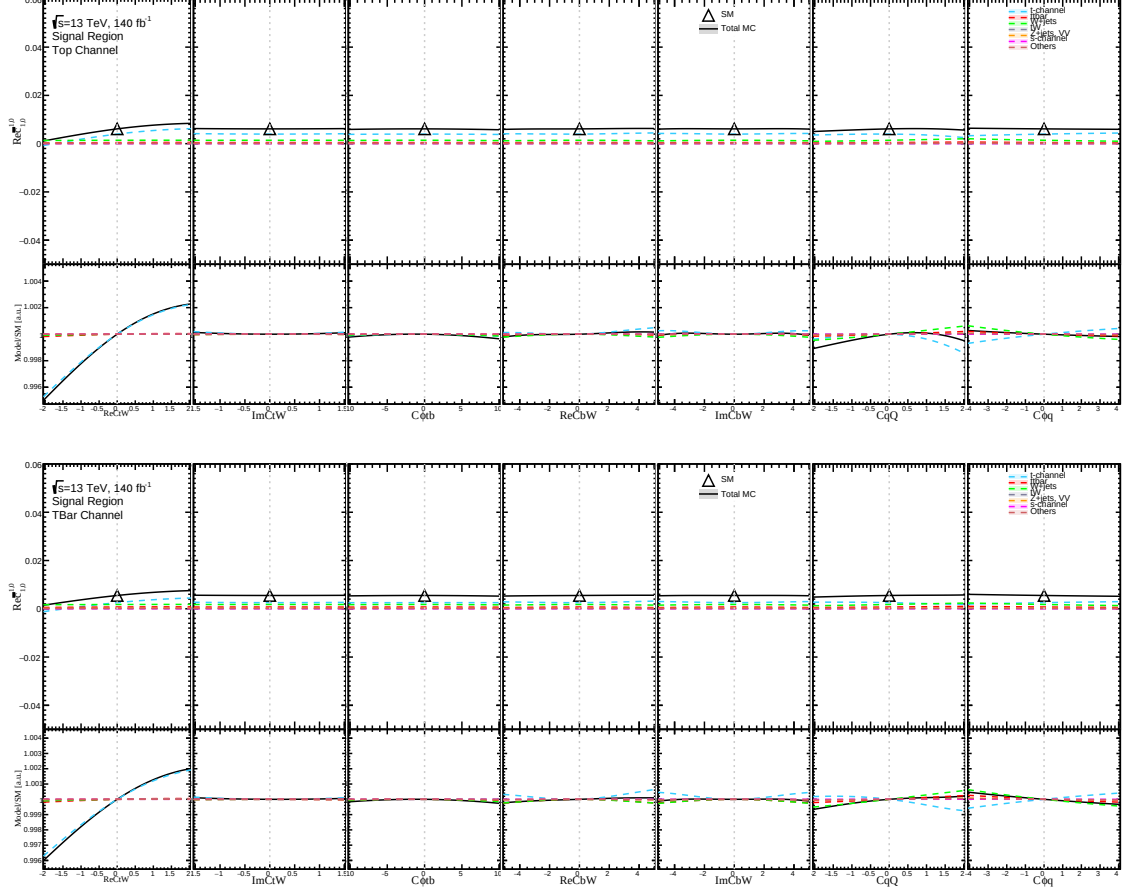


Figure 6.40: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } c_{1,0}^{1,0} \rangle^{\pm, \text{SR}}(C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses overlap each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

D Morphing 2d Plots

In this appendix, all 23×2 normalized coefficients as a two-dimensional function of the Wilson coefficients $(C_{\phi b}, C_{bW})$ and $(C_{\phi b}, C_{ibW})$ are shown.

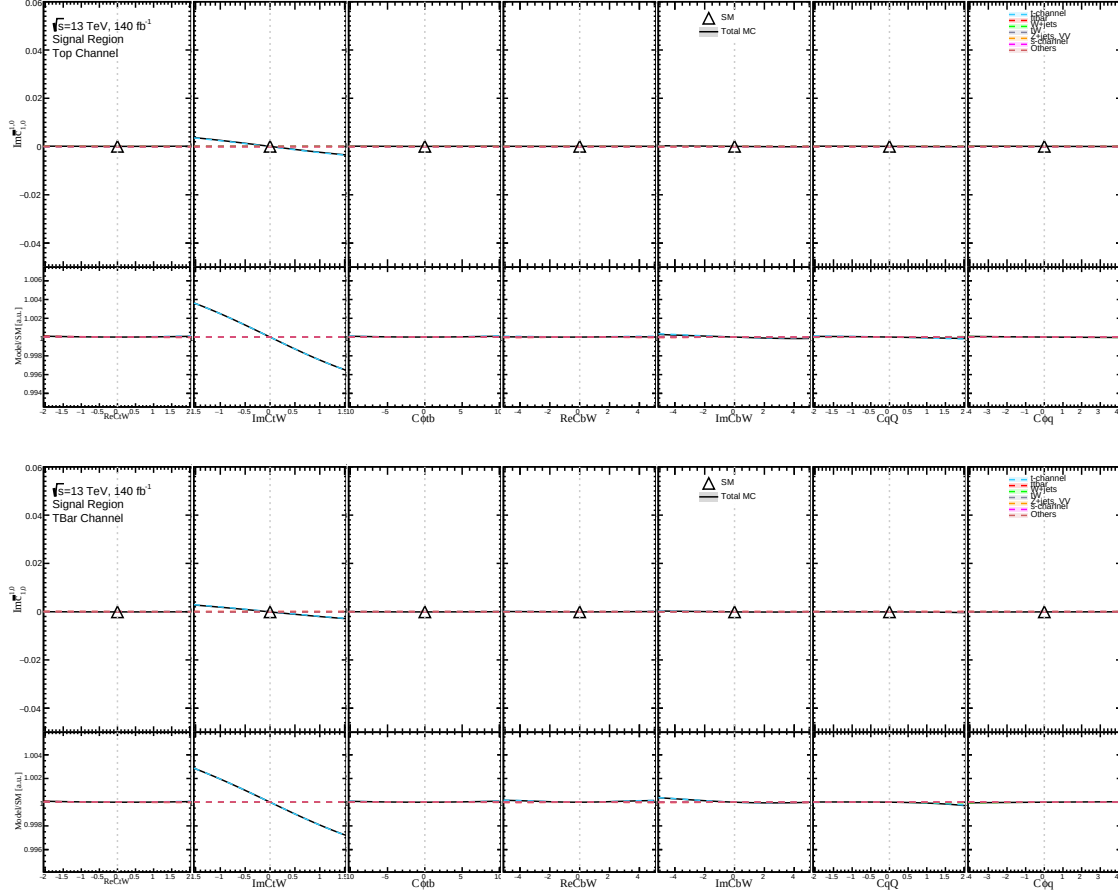


Figure 6.41: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Im } c_{1,0}^{1,0} \rangle^{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses are overlapped on each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

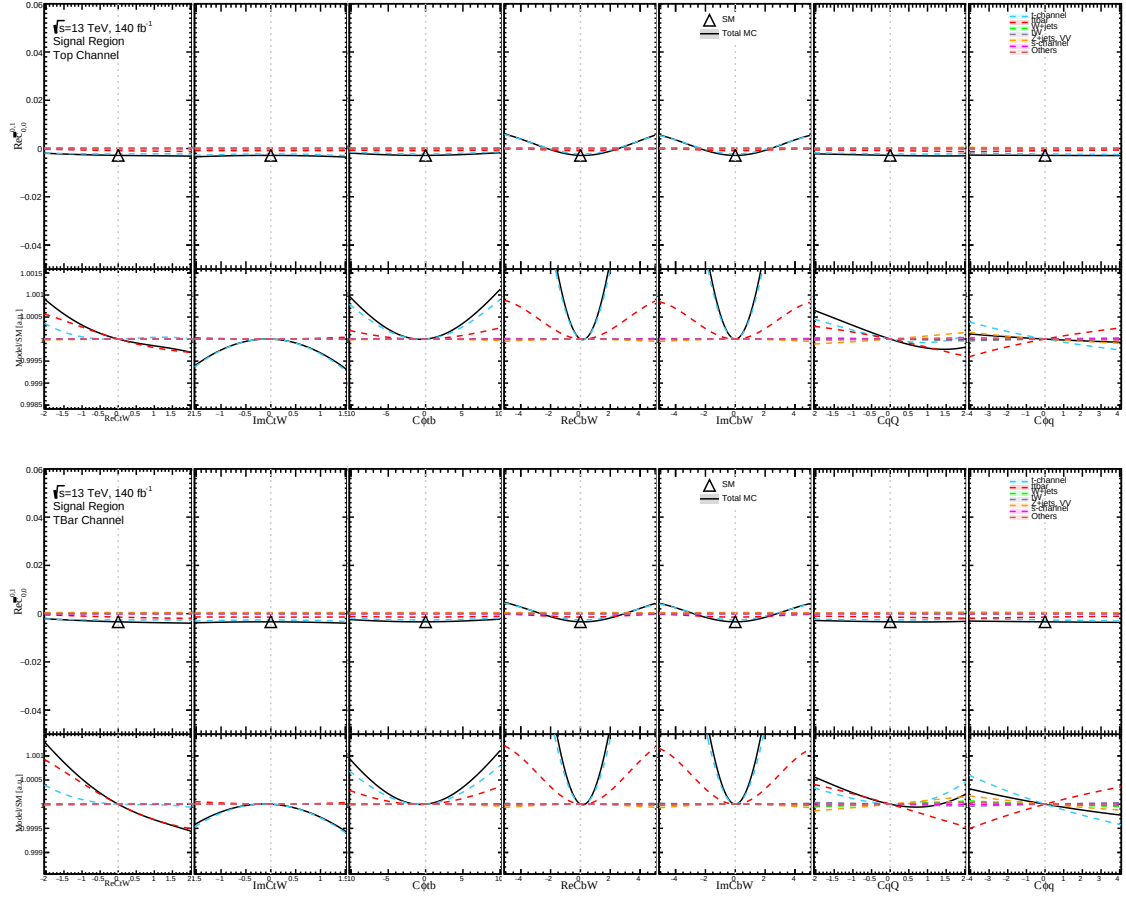


Figure 6.42: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } c_{0,0}^{0,1} \rangle^{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses overlap each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

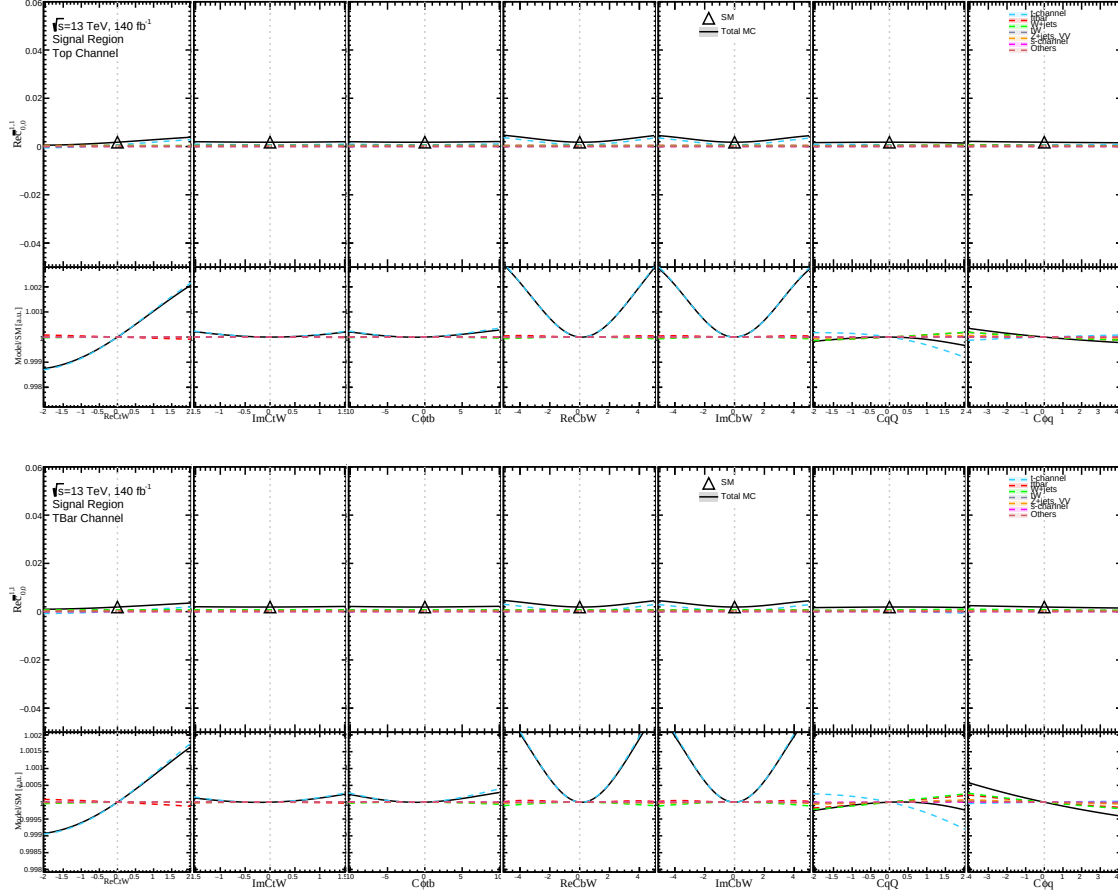


Figure 6.43: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } c_{0,0}^{1,1} \rangle_{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses overlap each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

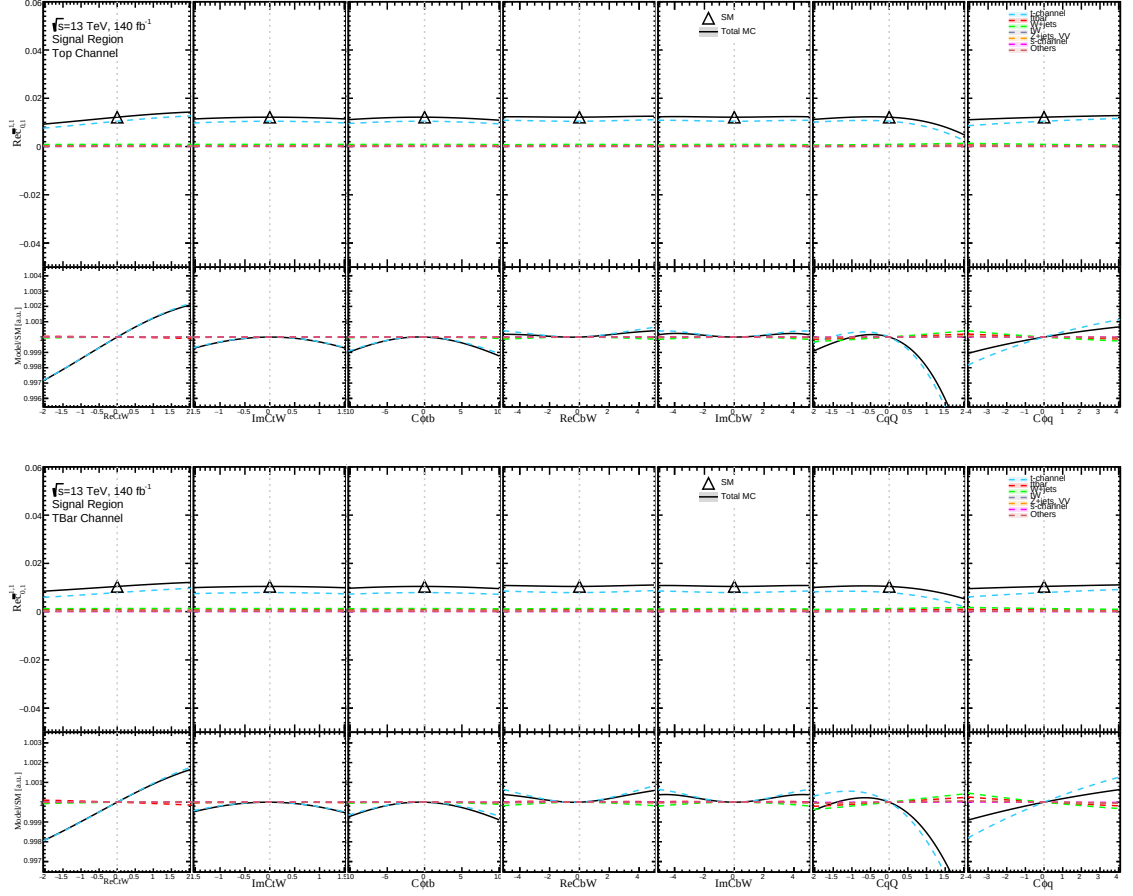


Figure 6.44: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } c_{0,1}^{1,1} \rangle^{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses overlap each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

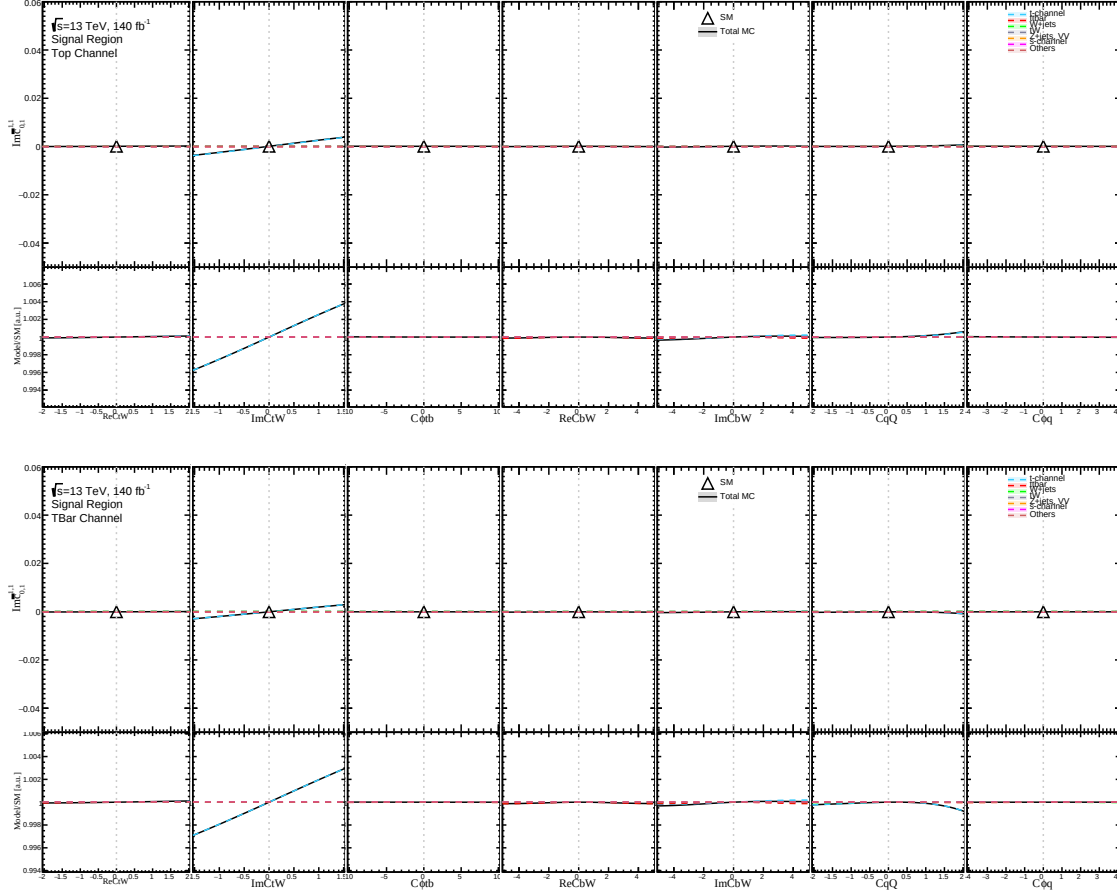


Figure 6.45: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Im } c_{0,1}^{1,1} \rangle_{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses overlap each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes comes from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

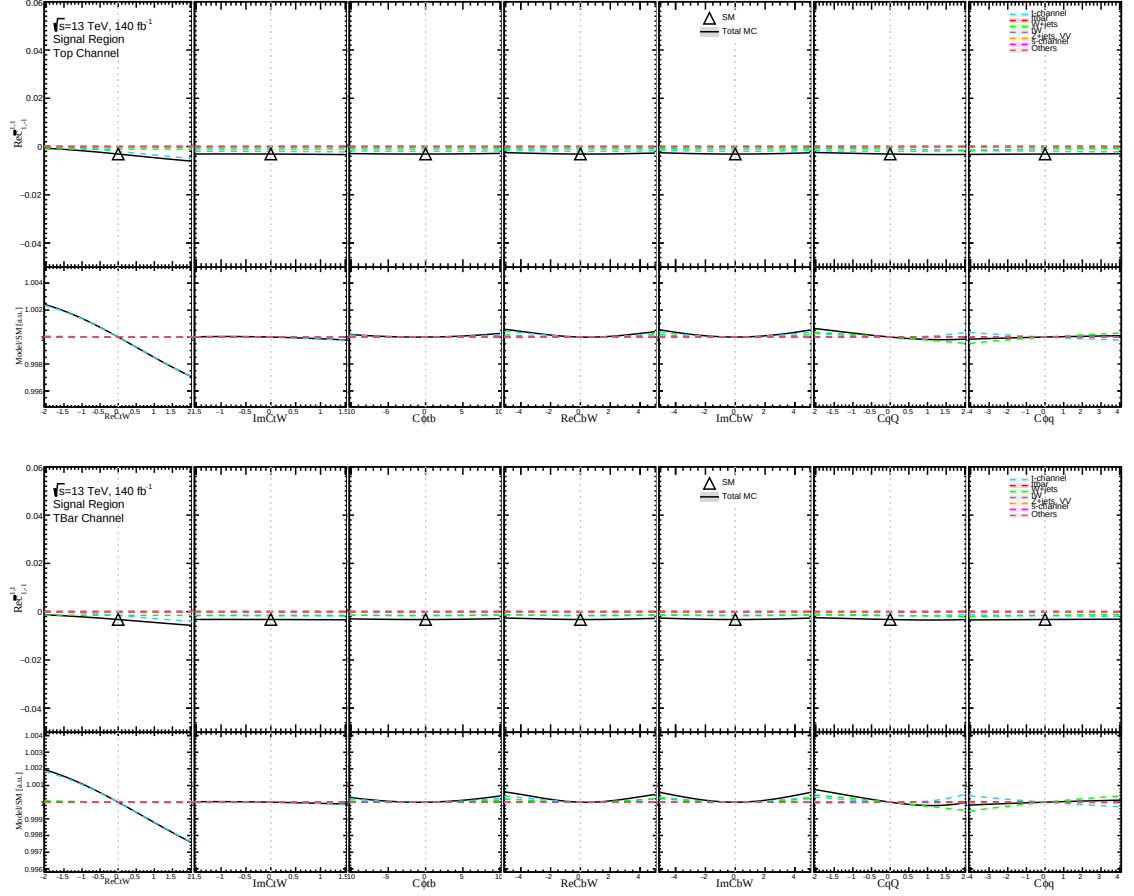


Figure 6.46: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re} c_{1,-1}^{1,1} \rangle^{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses overlap each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

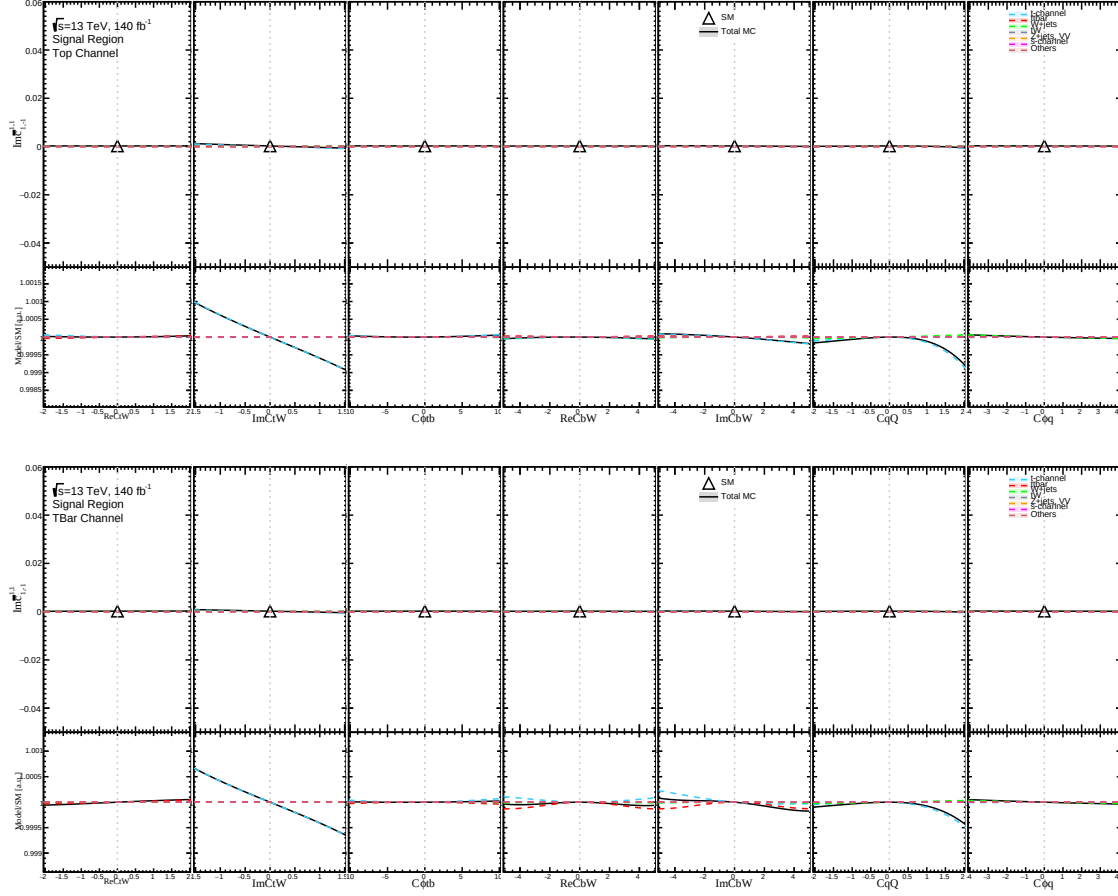


Figure 6.47: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Im } c_{1,-1}^{1,1} \rangle^{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses overlap each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

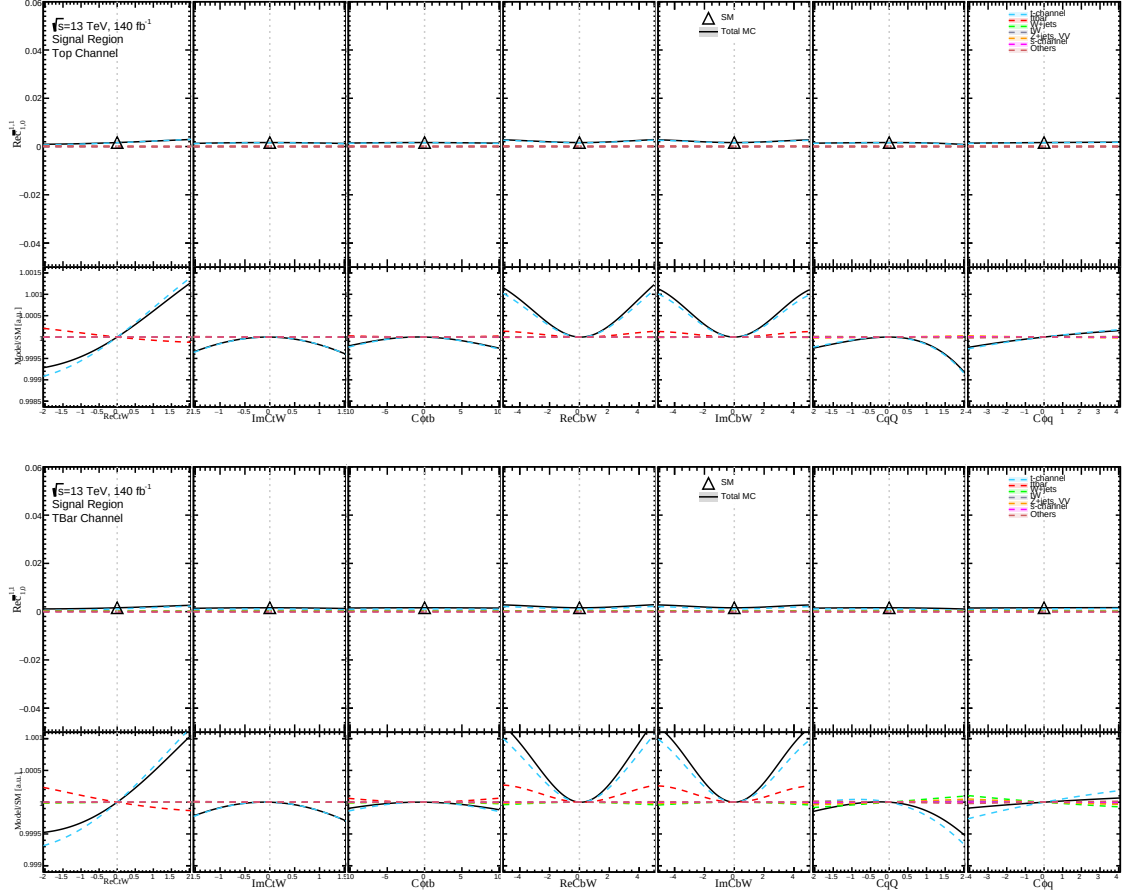


Figure 6.48: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } c_{1,0}^{1,1} \rangle^{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses are overlapped on each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

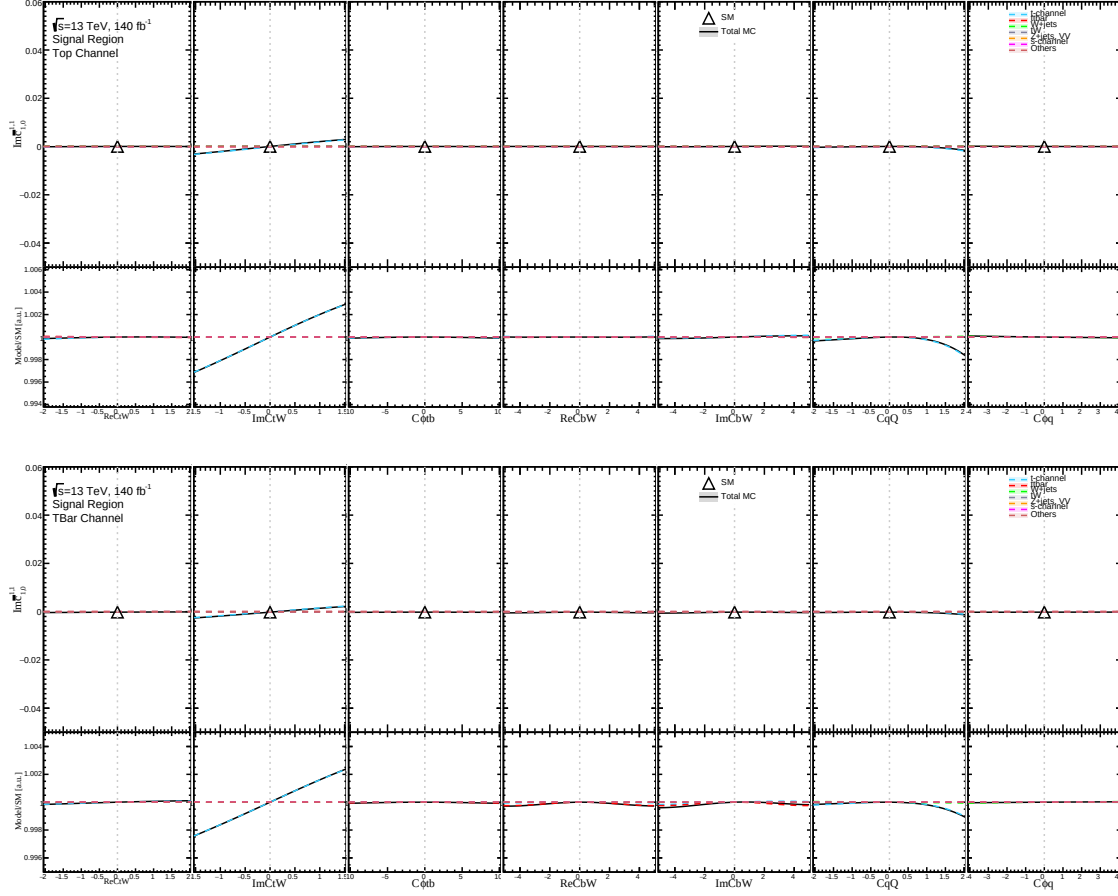


Figure 6.49: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Im } c_{1,0}^{1,1} \rangle^{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses overlap each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

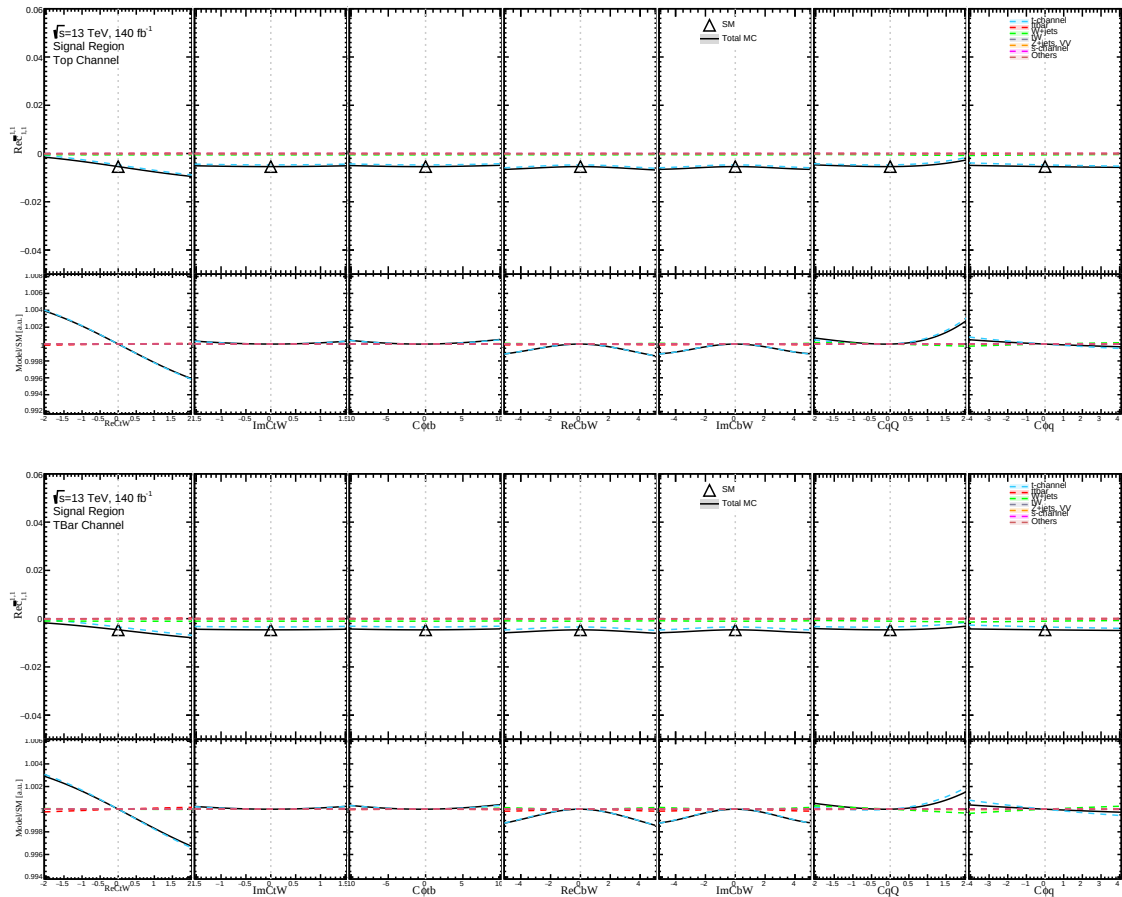


Figure 6.50: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re} c_{1,1}^{1,1} \rangle^{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses are overlapped on each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

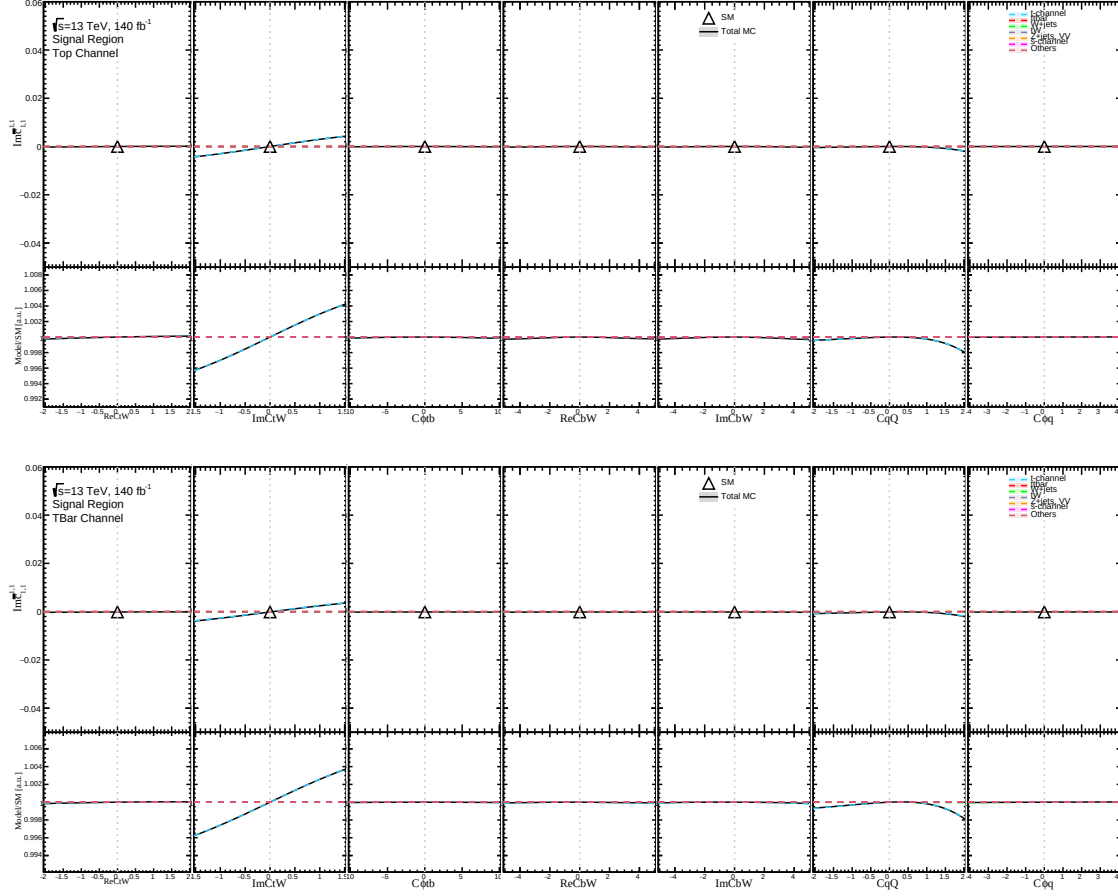


Figure 6.51: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Im } c_{1,1}^{1,1} \rangle_{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocess are overlapped on each other, the solid black line is corresponding to the linear combination of each sub process, and the triangle point is corresponding to the SM. A MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

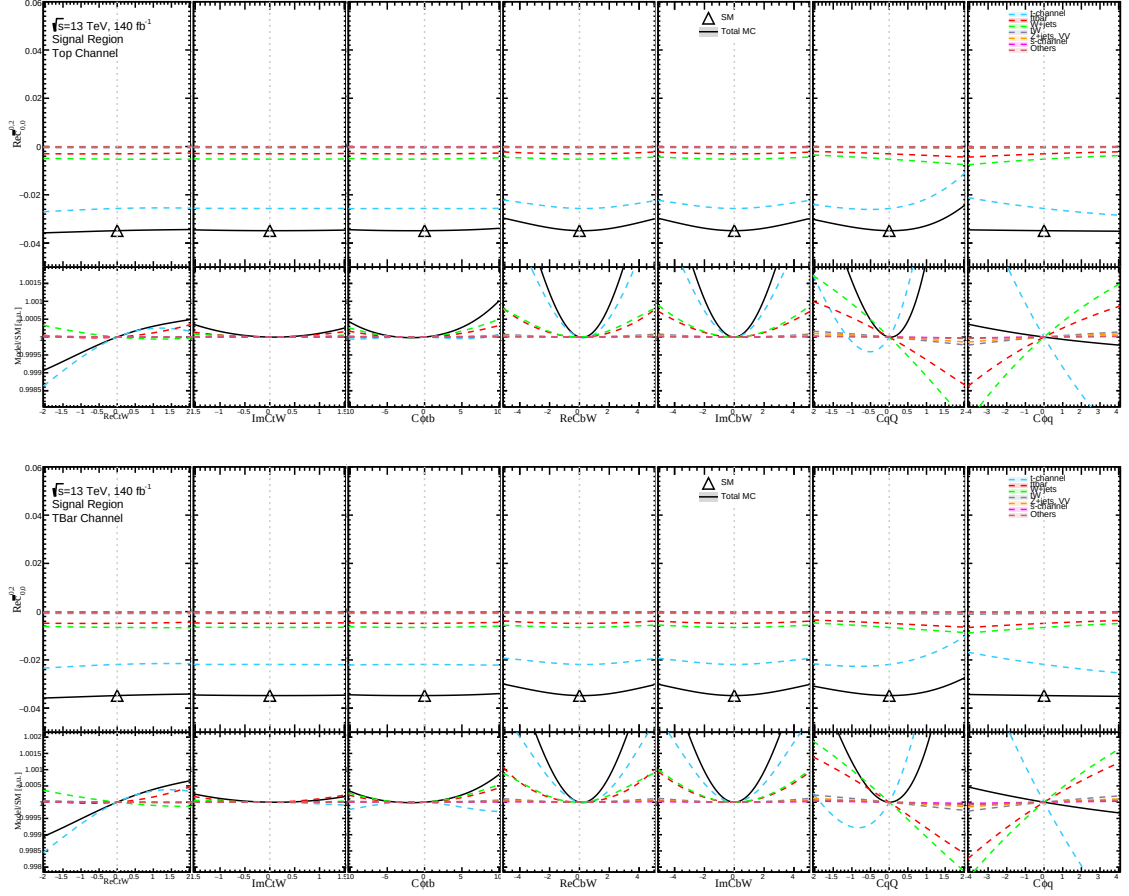


Figure 6.52: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } c_{0,0}^{0,2} \rangle^{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses overlap each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. A MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

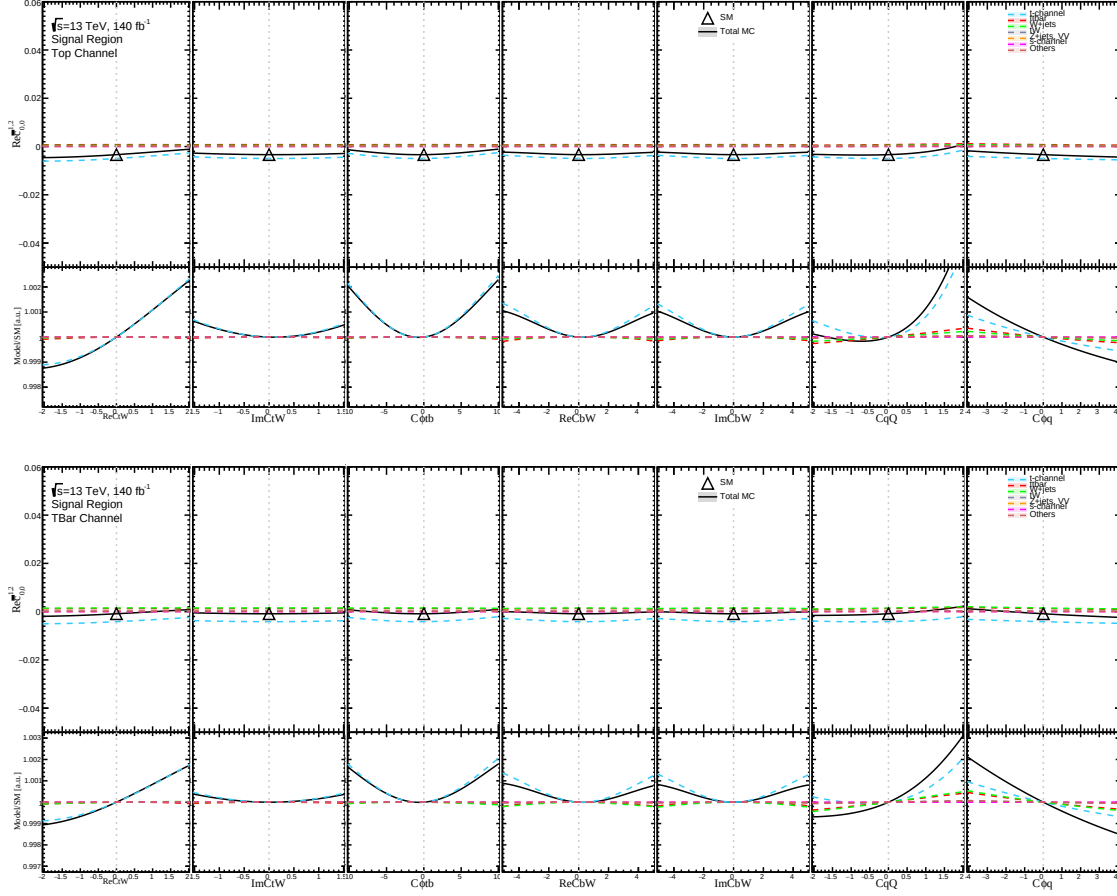


Figure 6.53: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re} c_{0,0}^{1,2} \rangle_{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses overlap each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

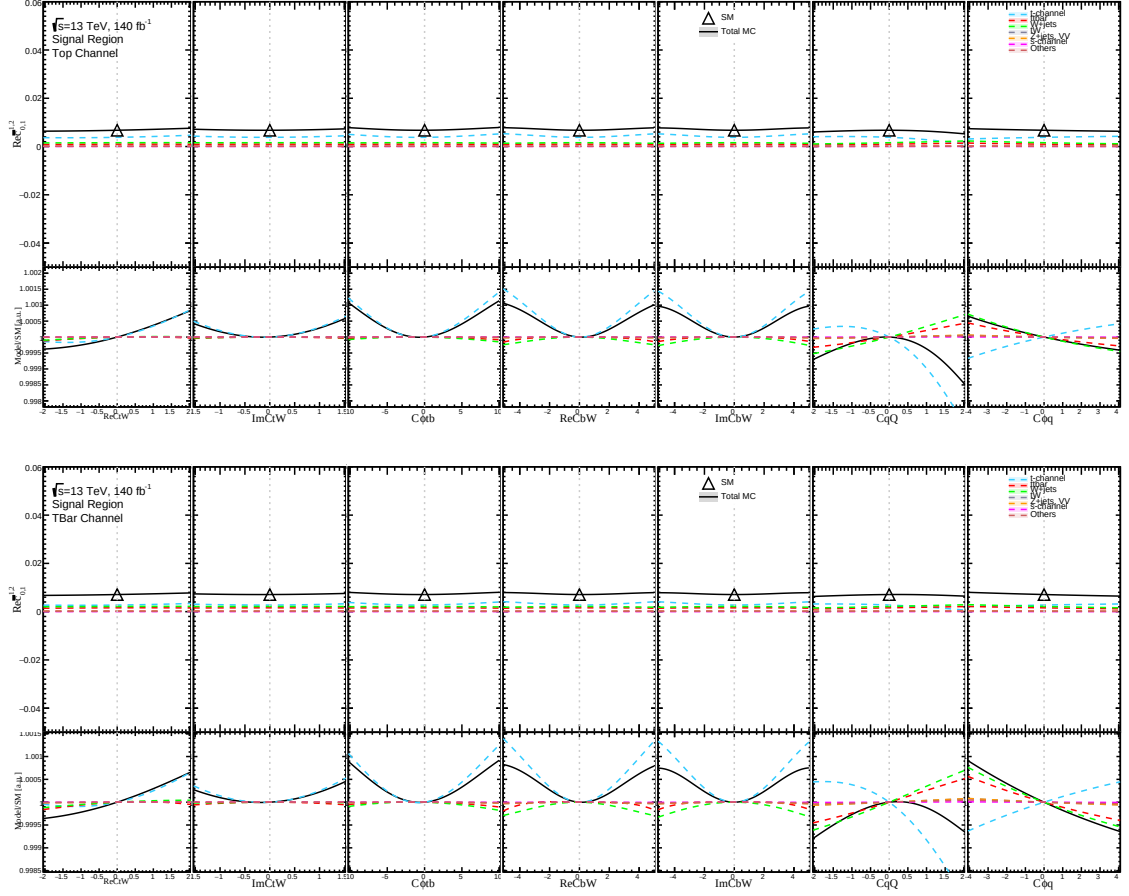


Figure 6.54: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } c_{0,1}^{1,2} \rangle^{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses overlap each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

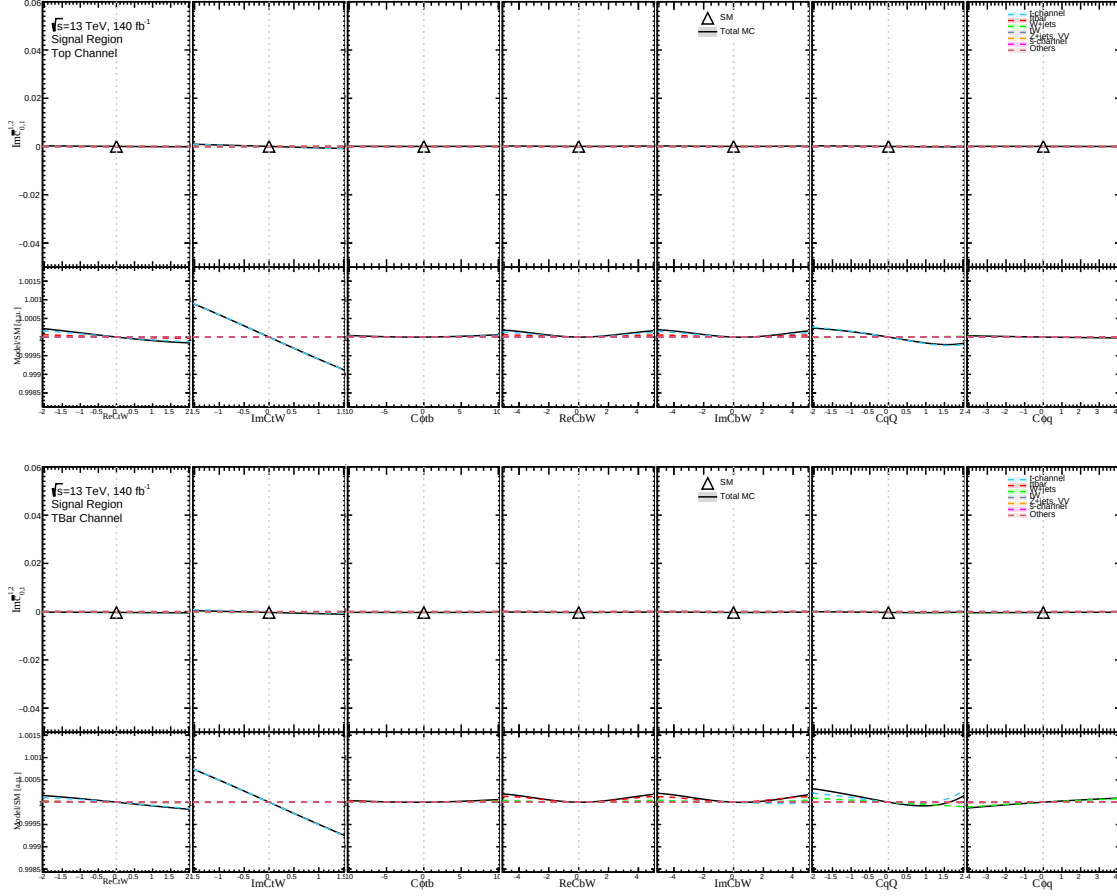


Figure 6.55: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Im } c_{0,1}^{1,2} \rangle_{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses are overlapped on each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

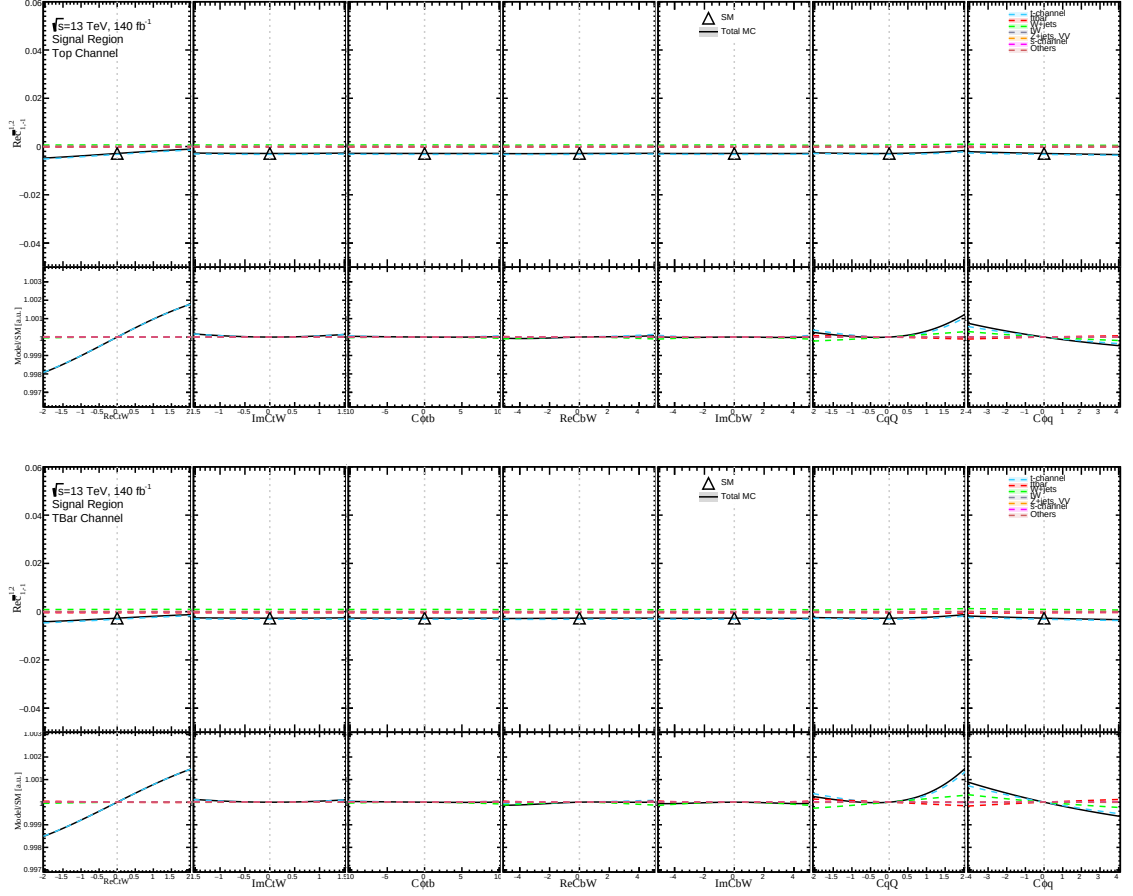


Figure 6.56: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re} C_{1,-1}^{1,2} \rangle^{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses overlap each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

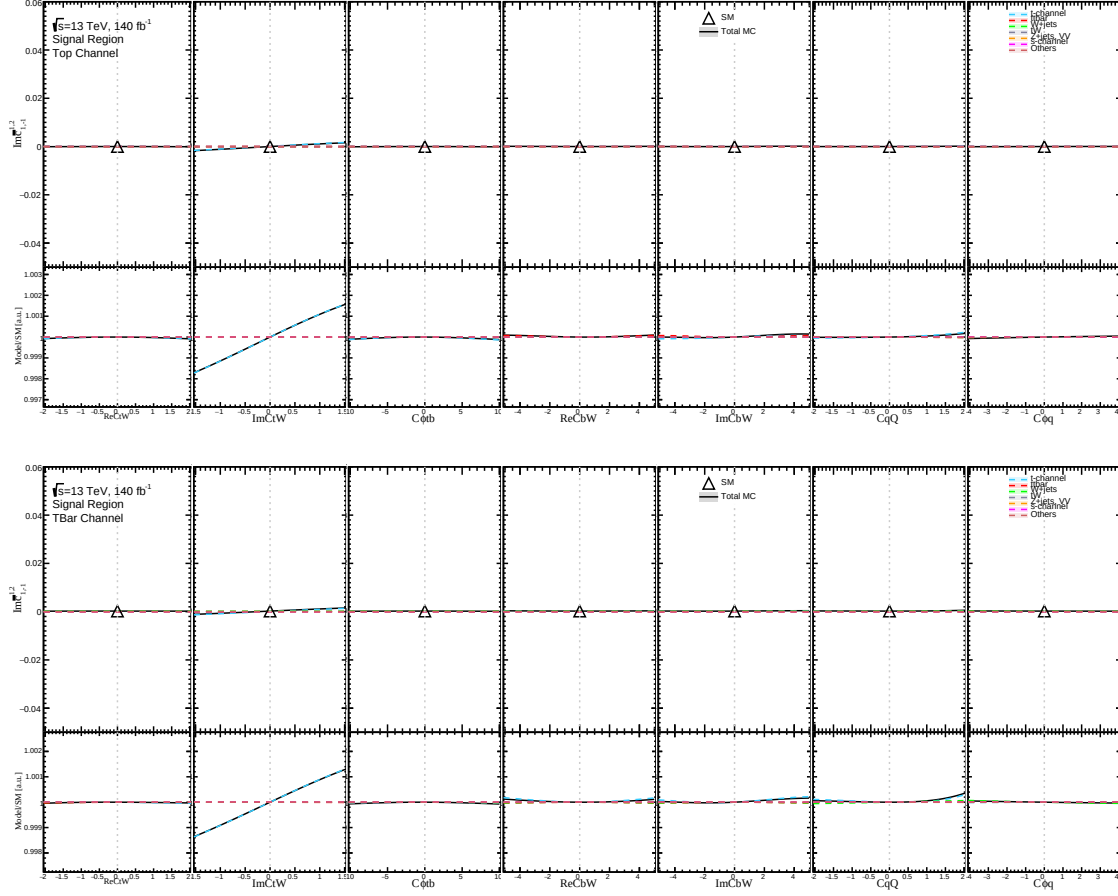


Figure 6.57: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Im } c_{1,-1}^{1,2} \rangle_{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses overlap each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

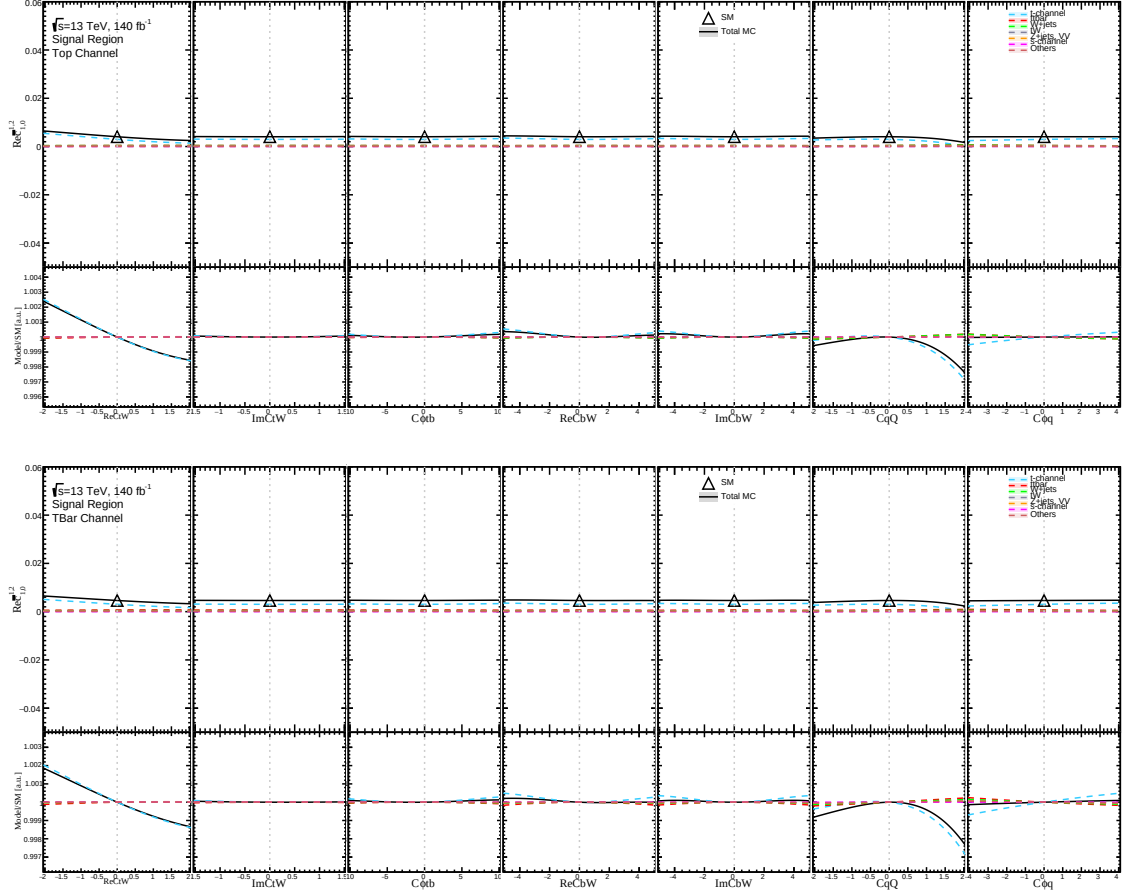


Figure 6.58: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re} c_{1,0}^{1,2} \rangle^{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses overlap each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

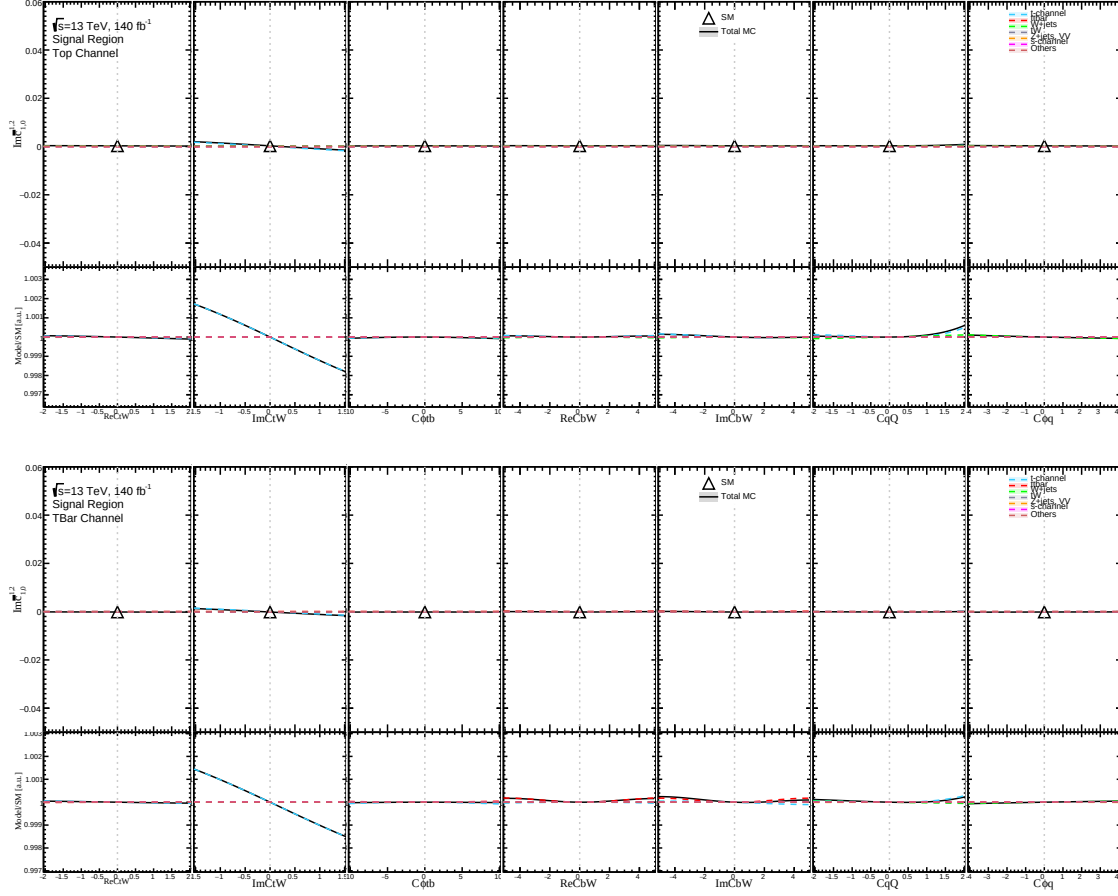


Figure 6.59: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Im } c_{1,0}^{1,2} \rangle^{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses overlap each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

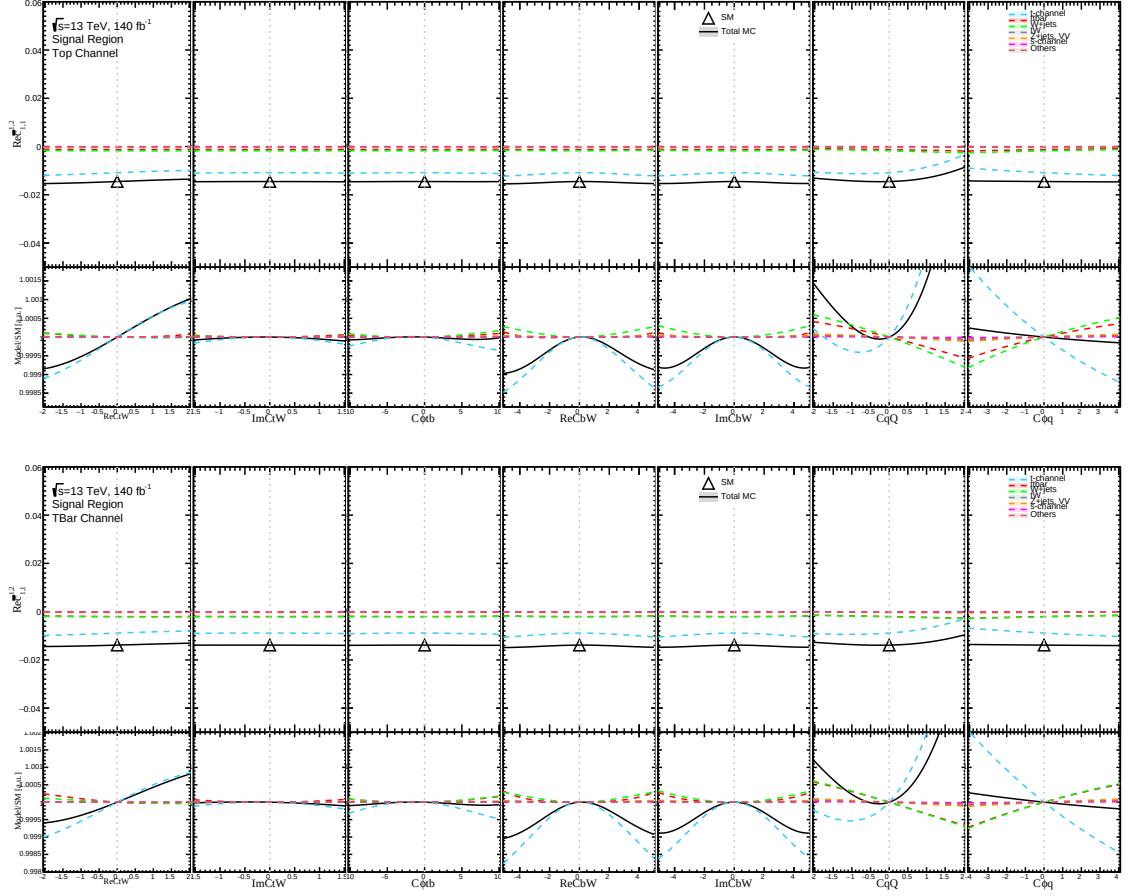


Figure 6.60: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re} c_{1,1}^{1,2} \rangle^{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses overlap each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

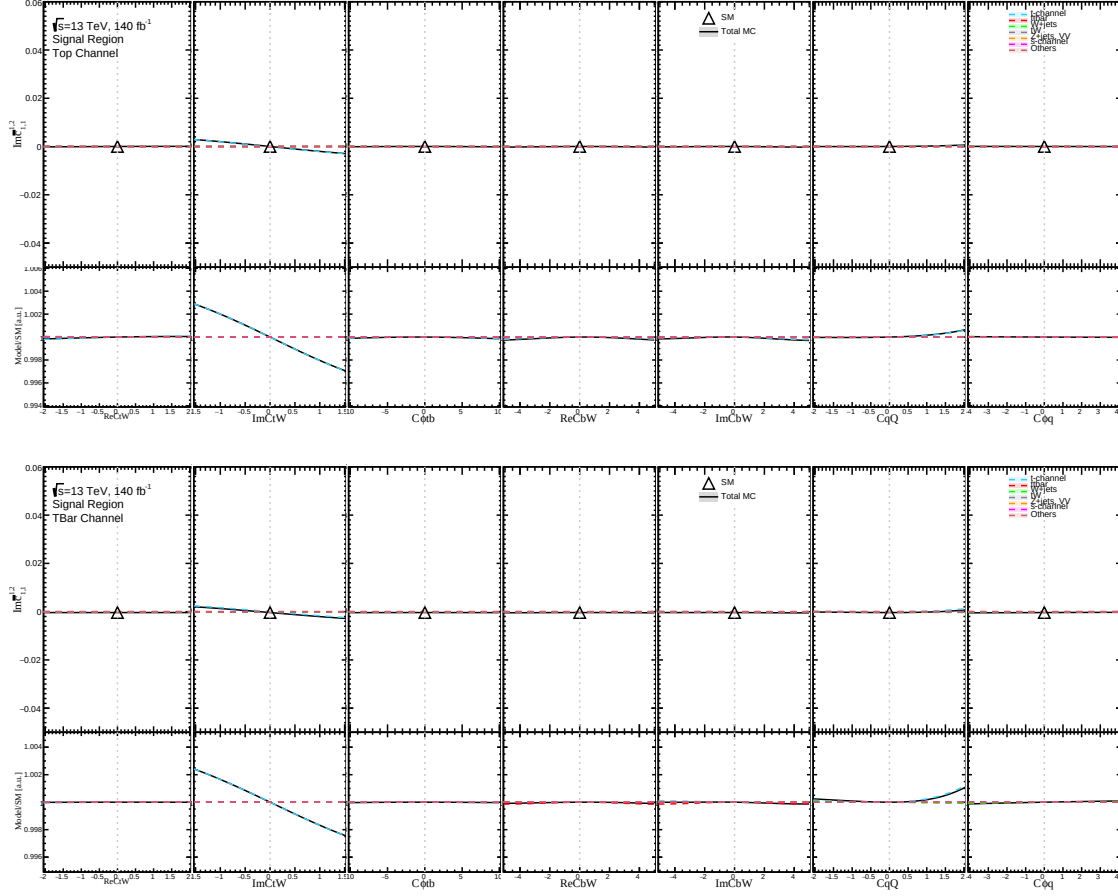


Figure 6.61: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Im } c_{1,1}^{1,2} \rangle^{\pm, \text{SR}} (C_i | \text{others} = 0)$. Each subplot corresponds to the morphed event counts as a function of one of the Wilson Coefficients, with all others fixed at zero. All subprocesses overlap each other, the solid black line corresponds to the linear combination of each subprocess, and the triangle point corresponds to the SM. An MC dataset corresponding to the LHC full run 2 data is being used but only t -channel and $t\bar{t}$ processes are being morphed. The Wilson Coefficient dependency for the other processes are coming from the effect of the morphed event count. Zoomed views are shown below by taking the ratio to the SM point.

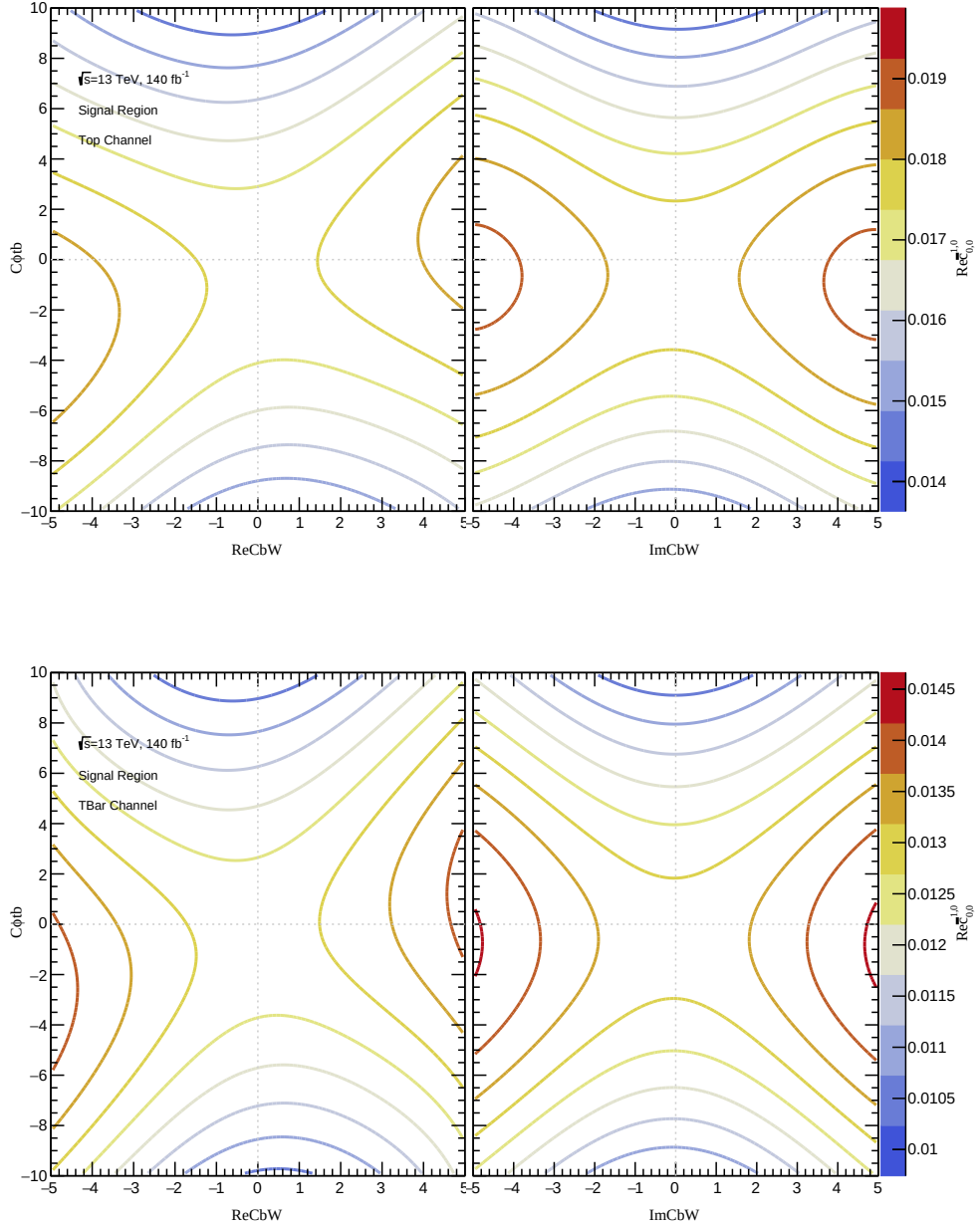


Figure 6.62: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re} C_{0,0}^{1,0} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $1 \vec{S} = \{1, 1, 1\}$.

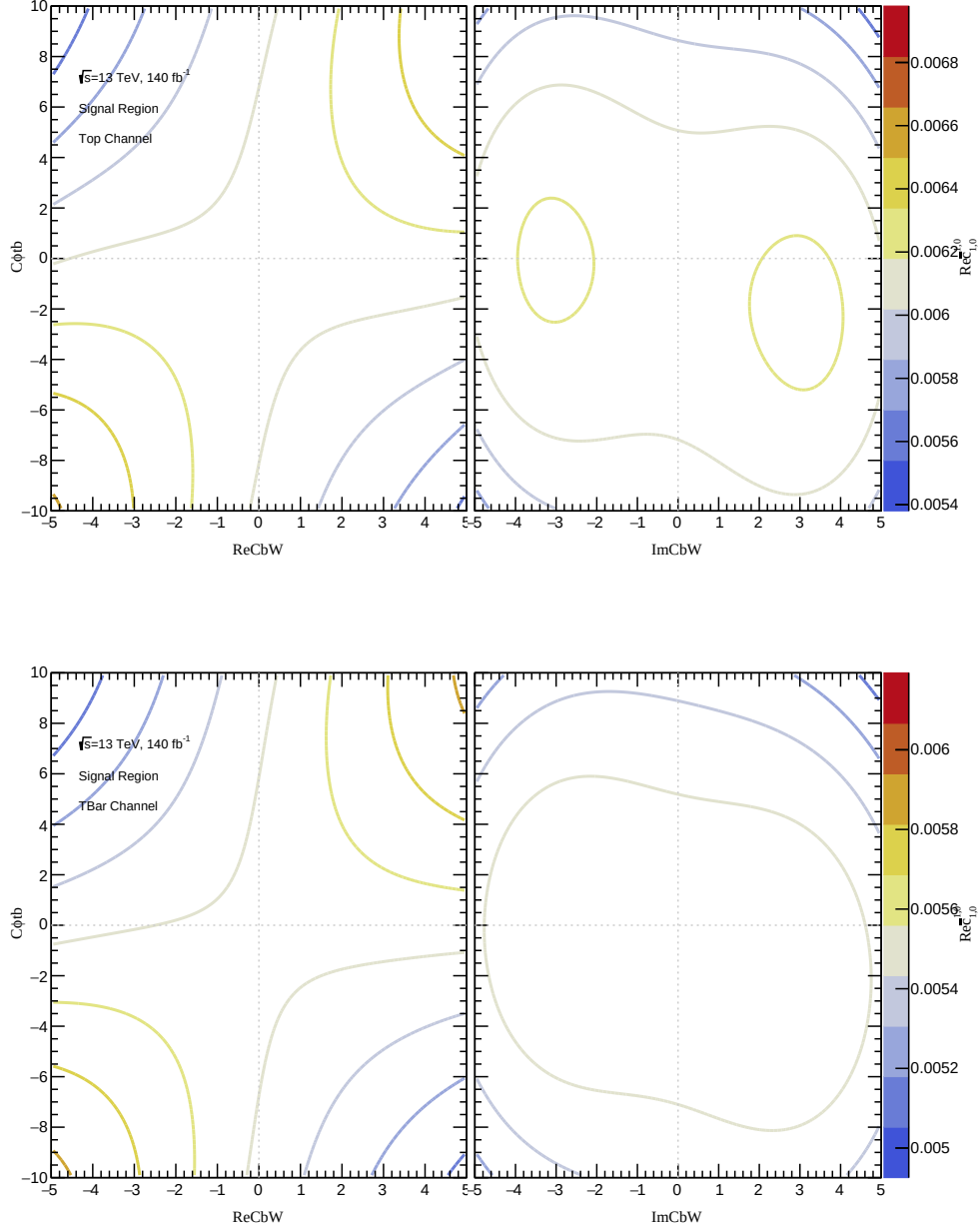


Figure 6.63: Morched event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } C_{1,0}^{1,0} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $\vec{S} = \{1, 1, 1\}$.

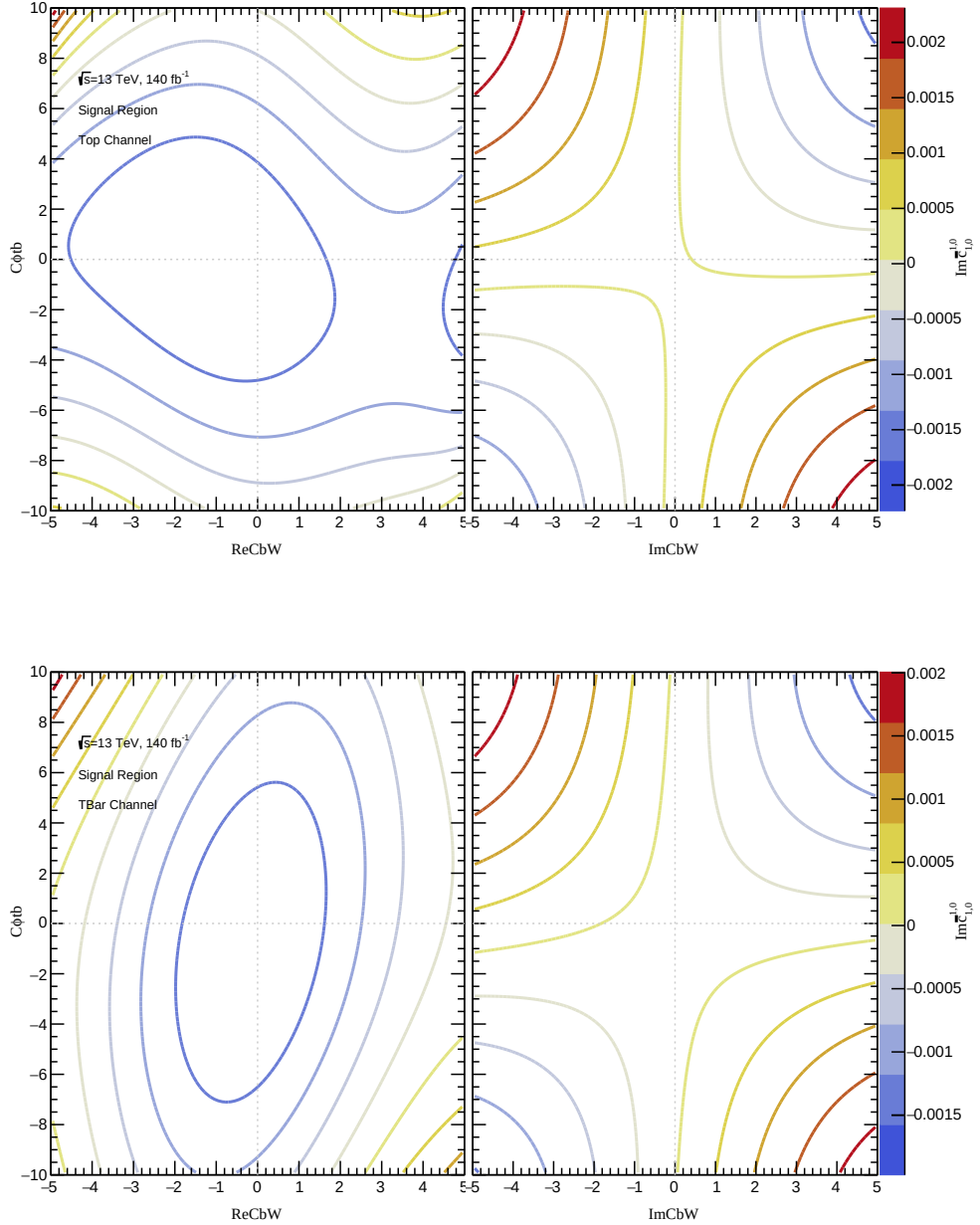


Figure 6.64: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Im } C_{0,0}^{1,0} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $1 \tilde{S} = \{1, 1, 1\}$.

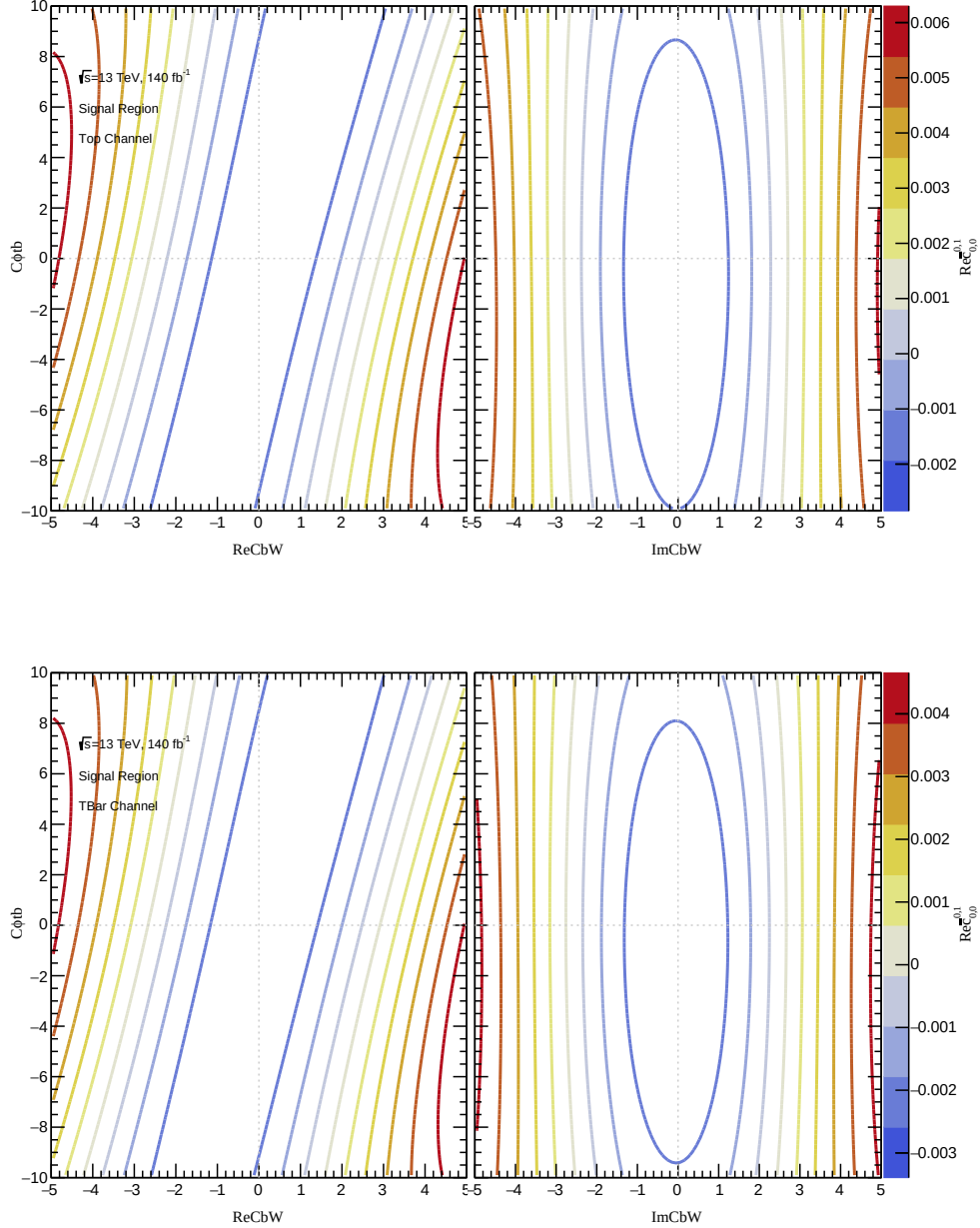


Figure 6.65: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } C_{0,0}^{0,1} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $1 \vec{S} = \{1, 1, 1\}$.

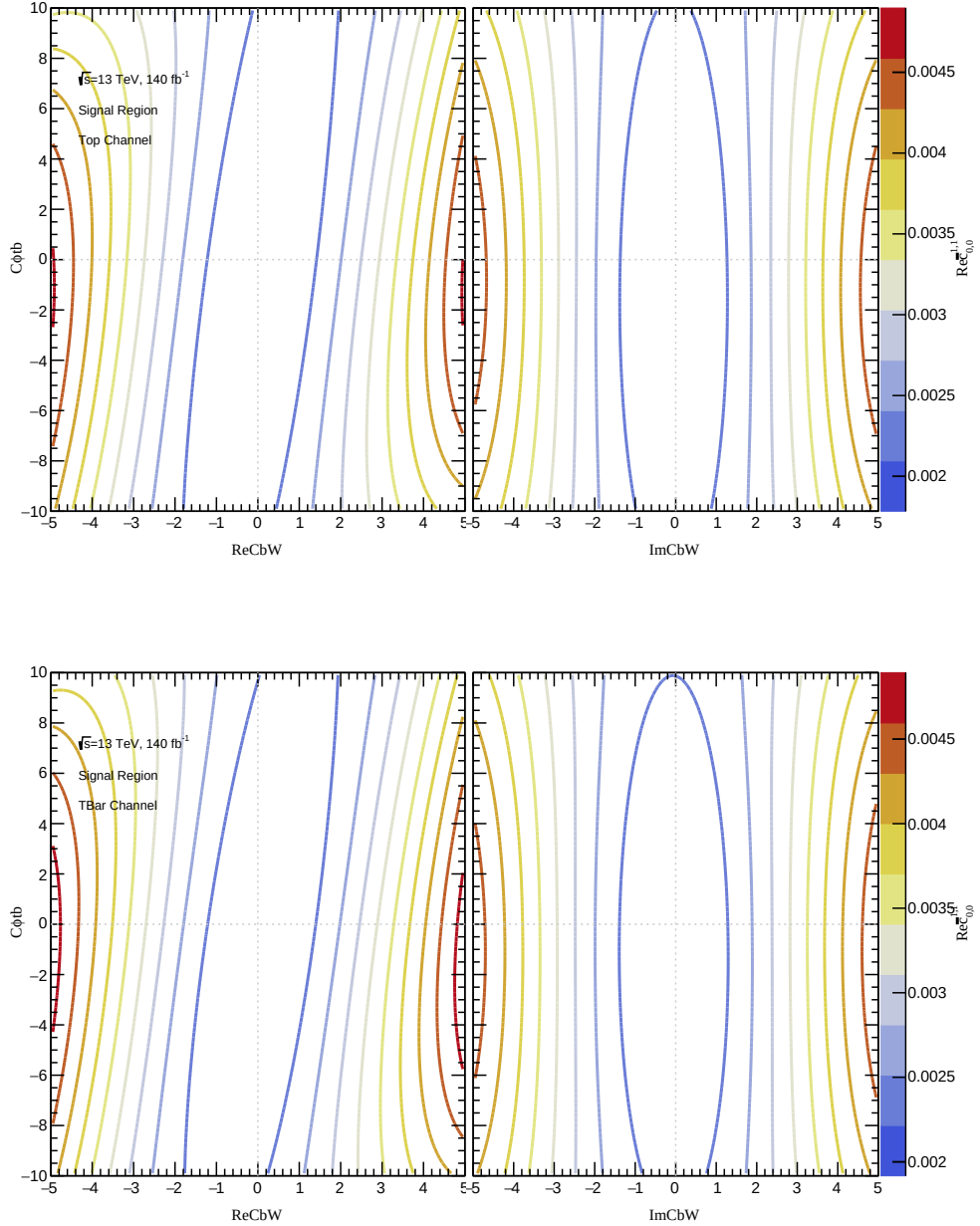


Figure 6.66: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re} C_{0,0}^{1,1} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $1 \tilde{S} = \{1, 1, 1\}$.

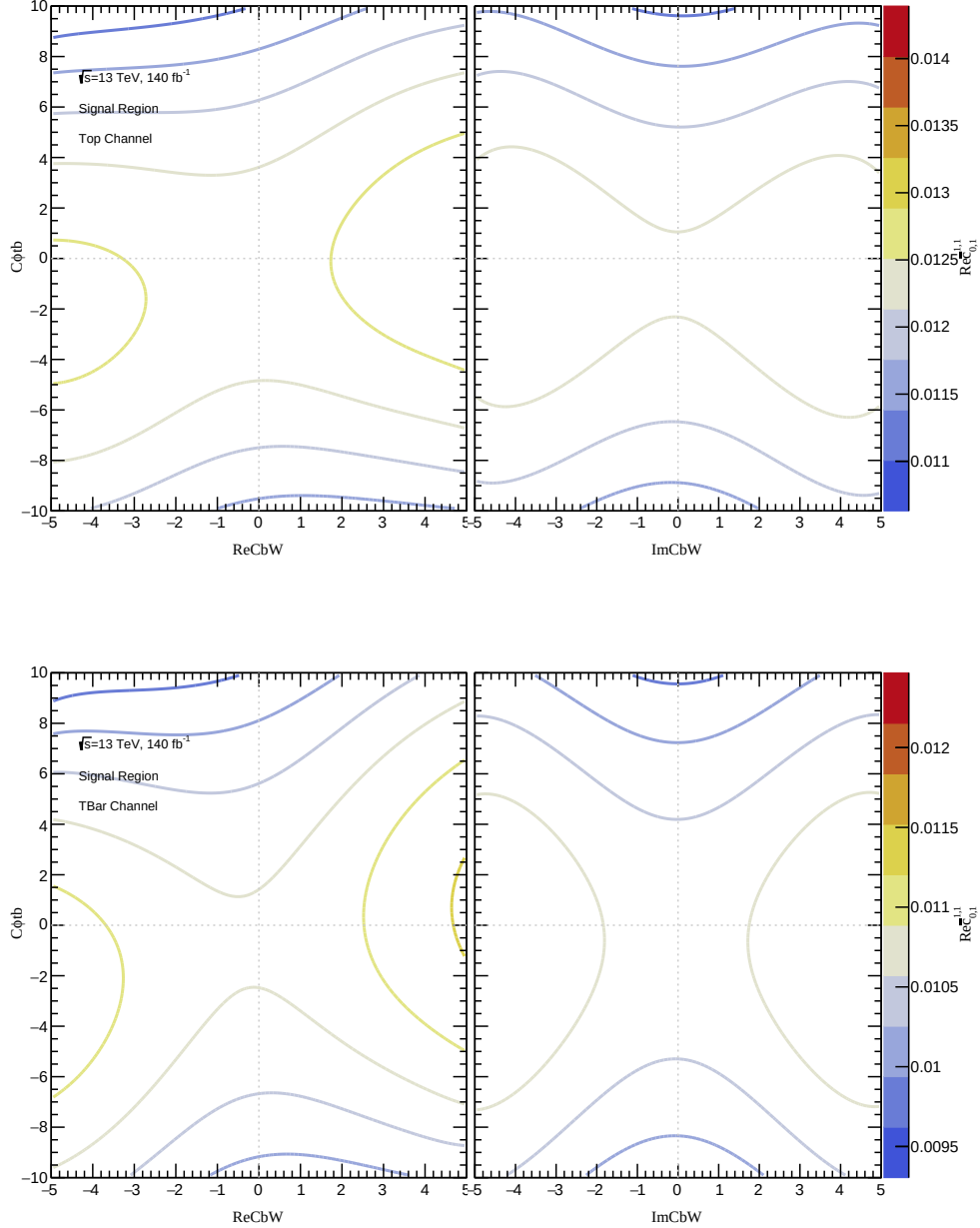


Figure 6.67: Morched event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } C_{0,1}^{1,1} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $1 \vec{S} = \{1, 1, 1\}$.

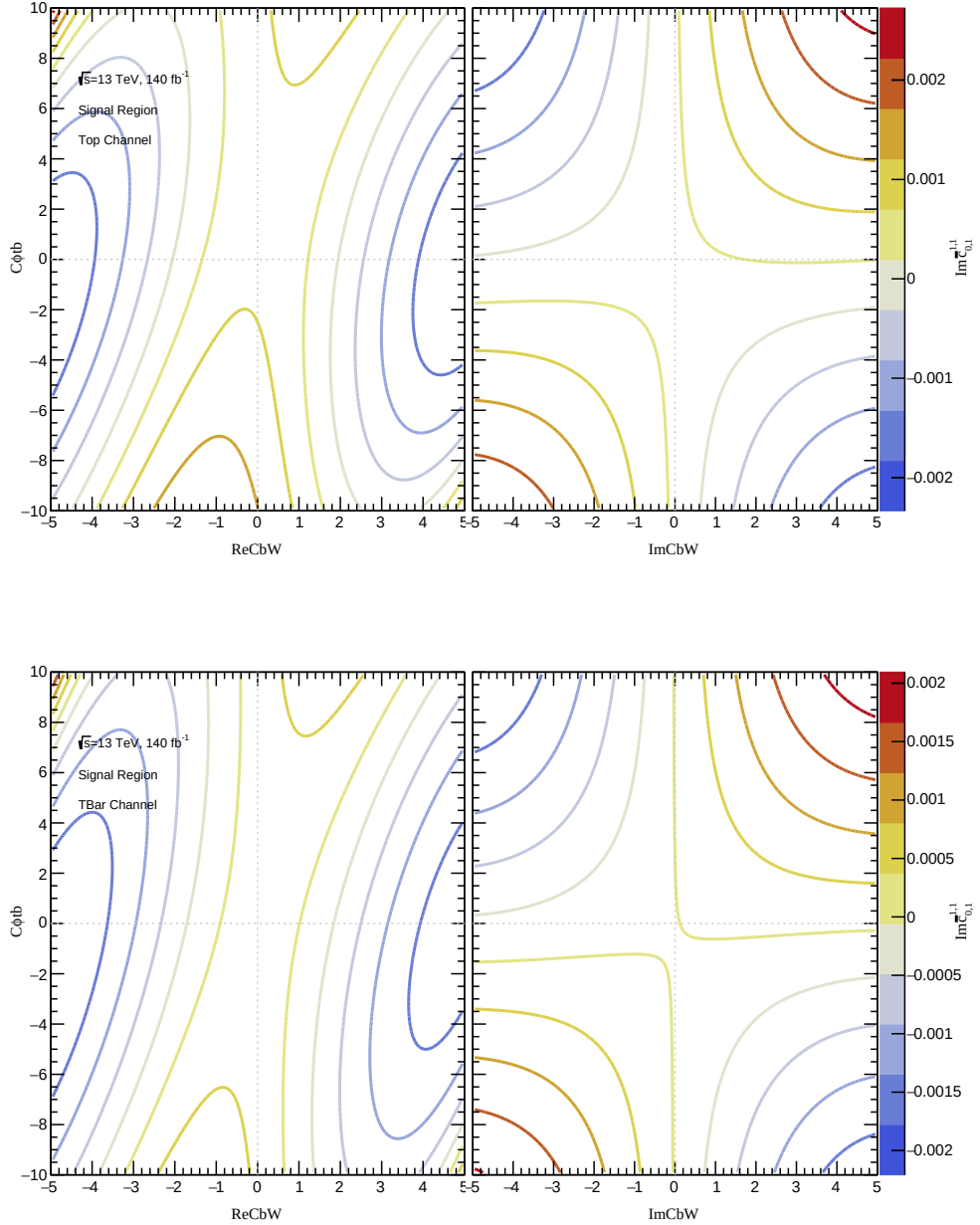


Figure 6.68: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Im } C_{0,0}^{1,1} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $1 \tilde{S} = \{1, 1, 1\}$.

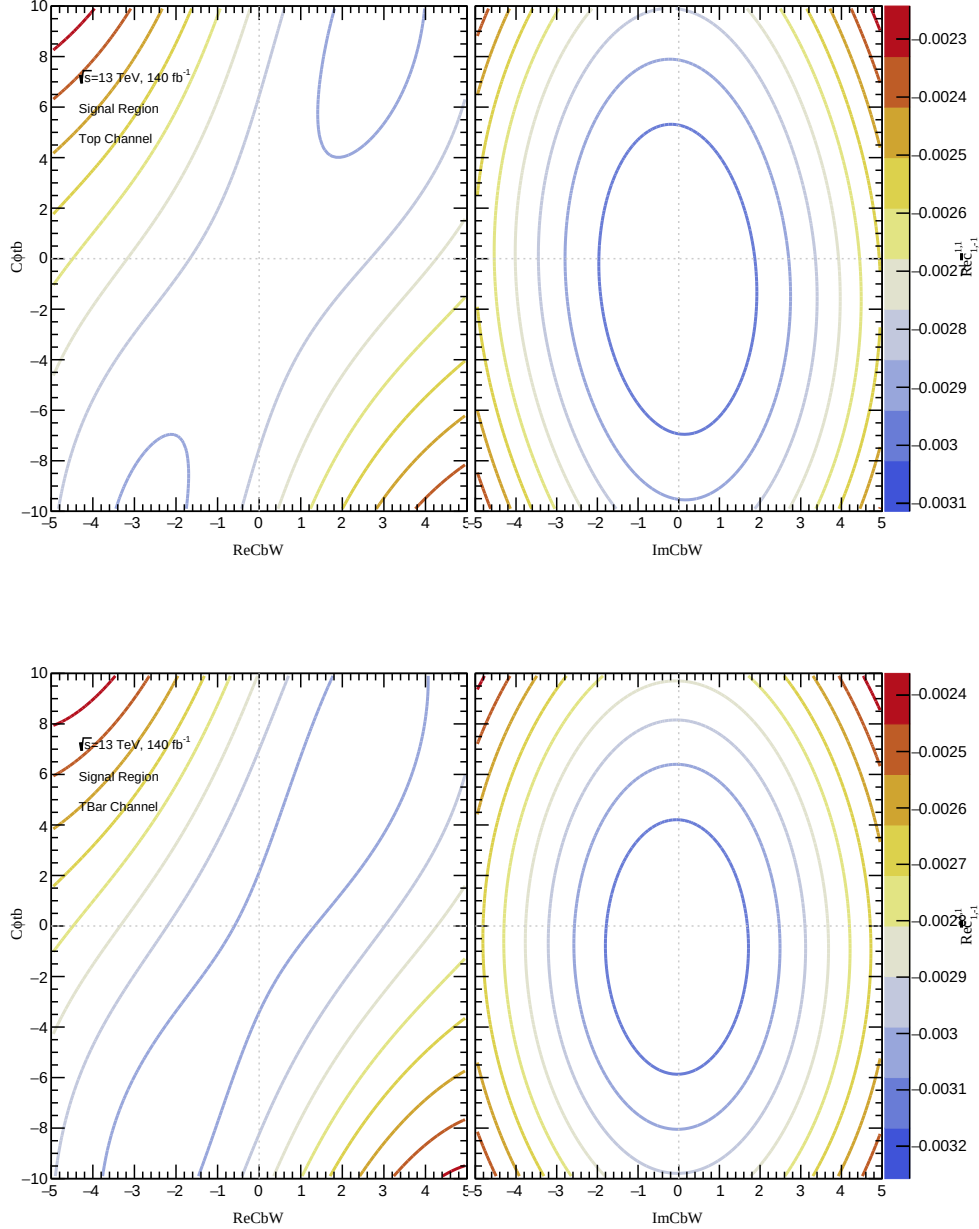


Figure 6.69: Morphe event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } C_{1,-1}^{1,1} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $\vec{S} = \{1, 1, 1\}$.

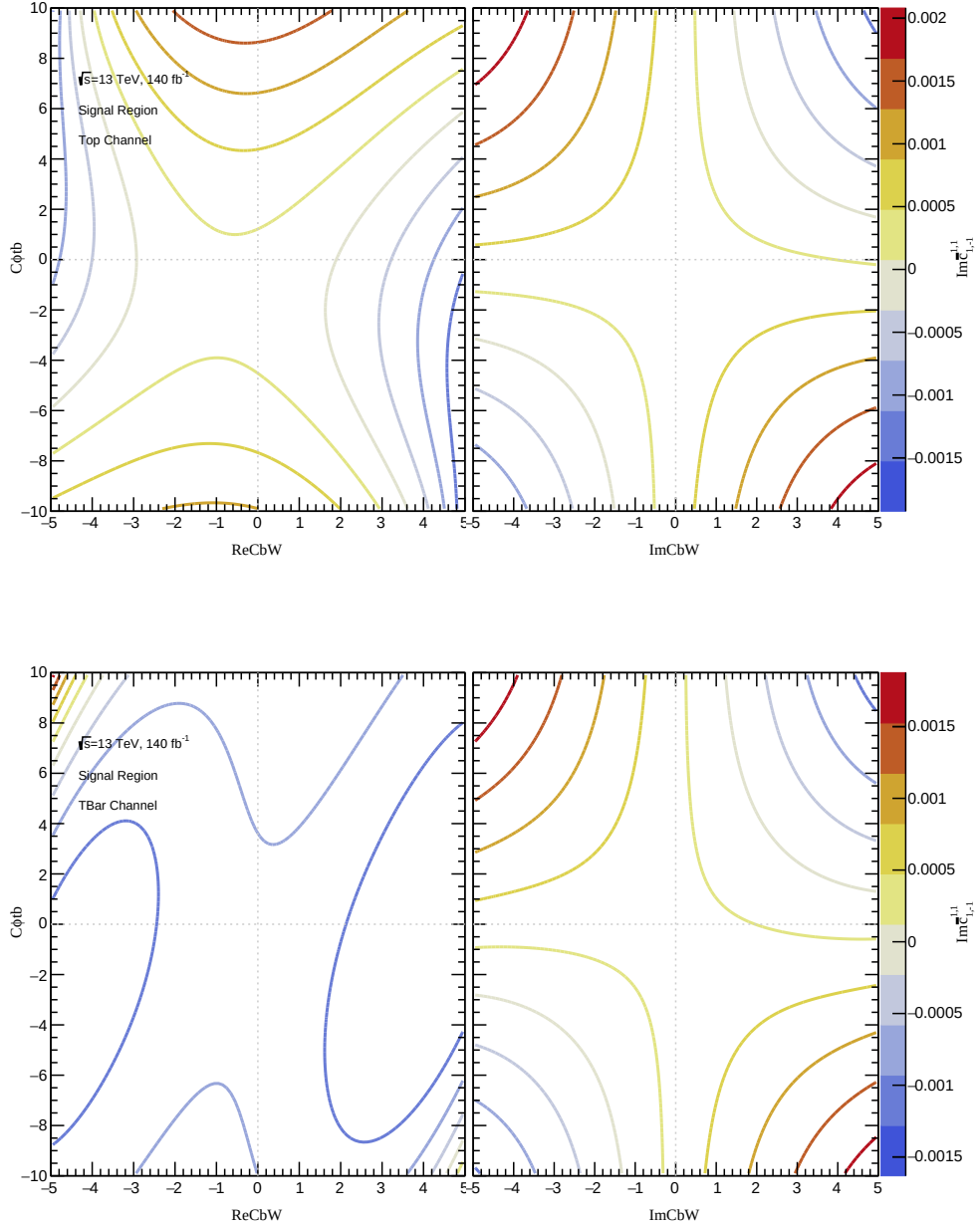


Figure 6.70: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } C_{1,-1}^{1,1} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $1 \ \tilde{S} = \{1, 1, 1\}$.

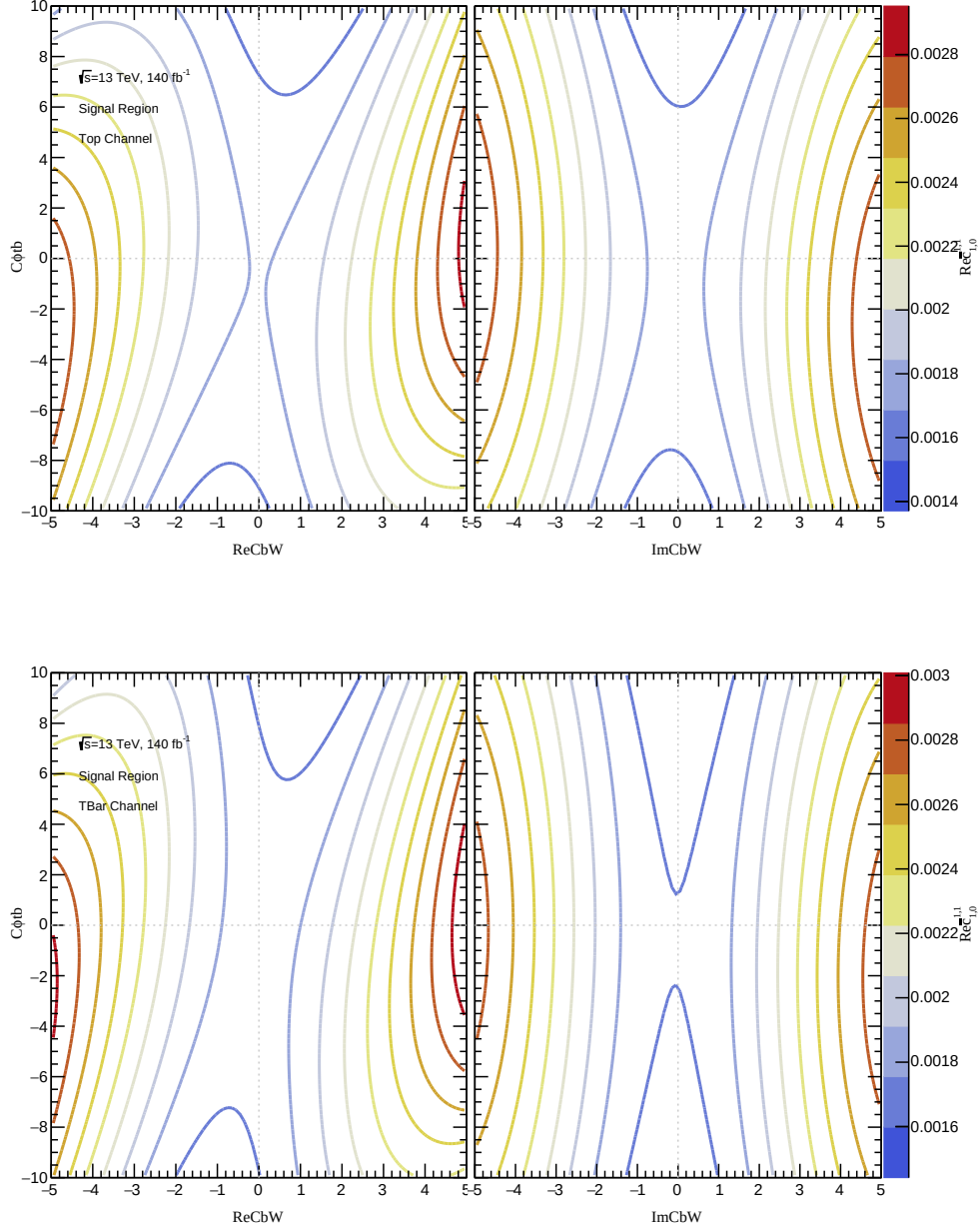


Figure 6.71: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } C_{1,0}^{1,1} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $\vec{S} = \{1, 1, 1\}$.

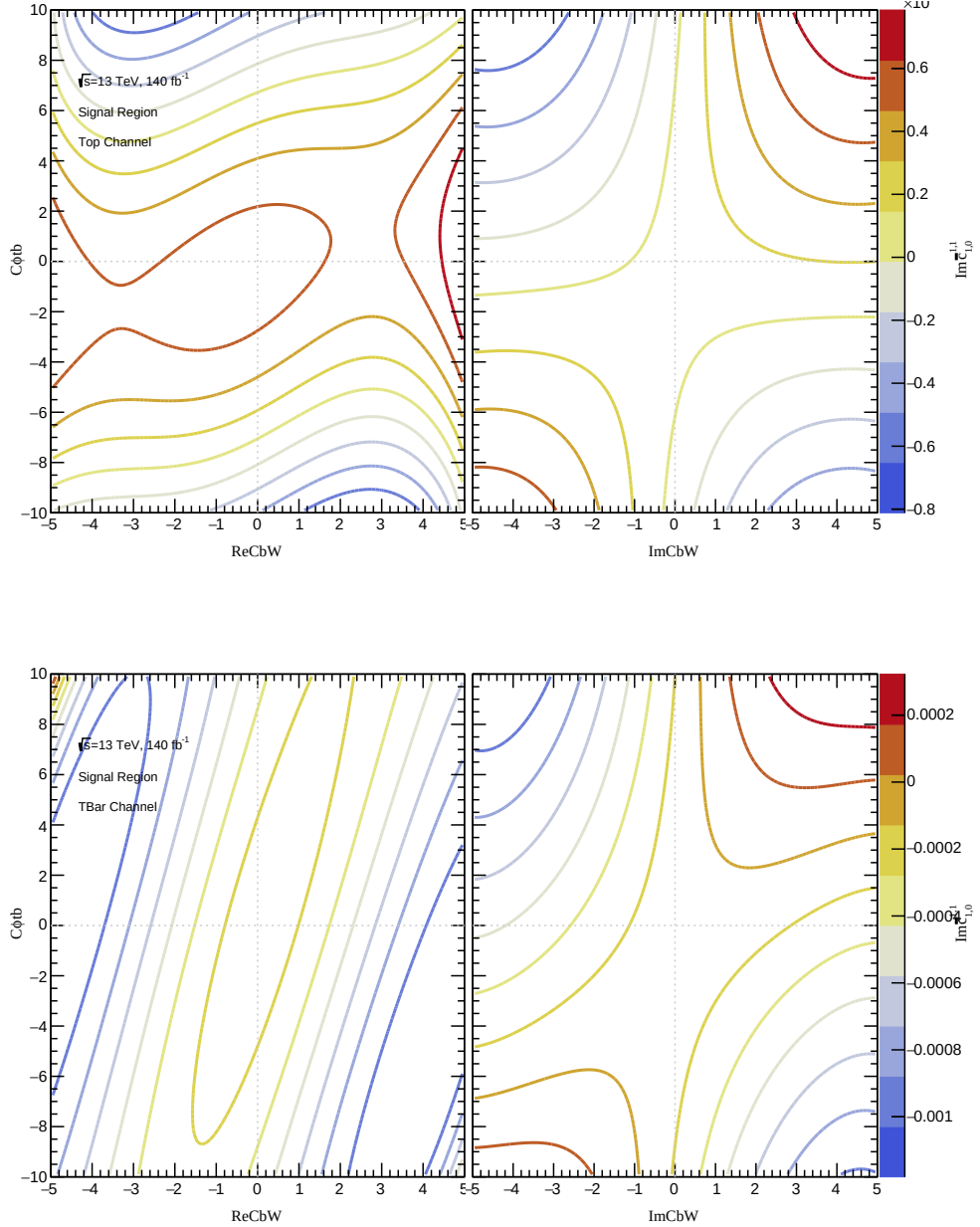


Figure 6.72: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Im } C_{0,0}^{1,0} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $1 \ \tilde{S} = \{1, 1, 1\}$.

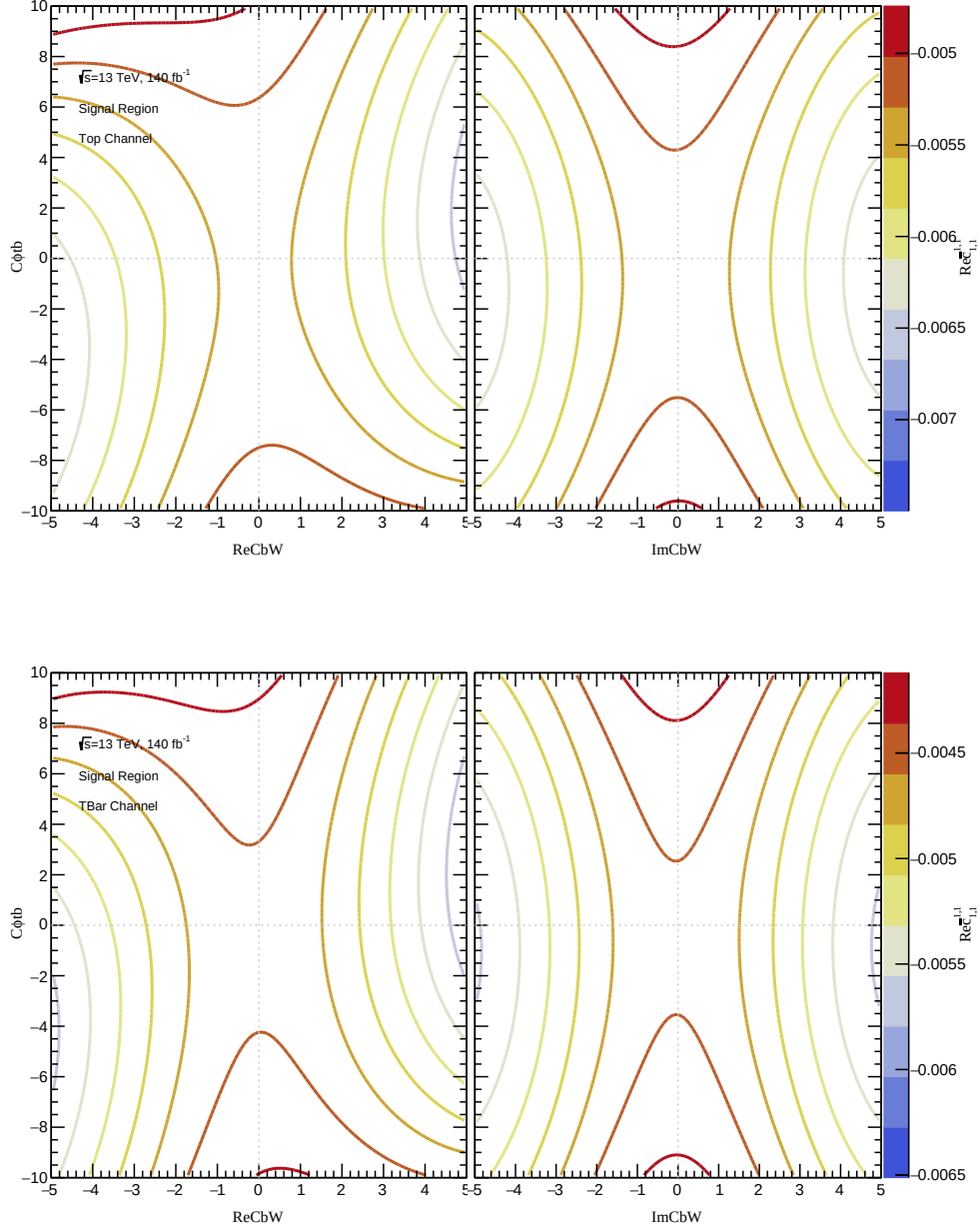


Figure 6.73: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } C_{1,1}^{1,1} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{i\phi b} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{i\phi b} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $\vec{S} = \{1, 1, 1\}$.

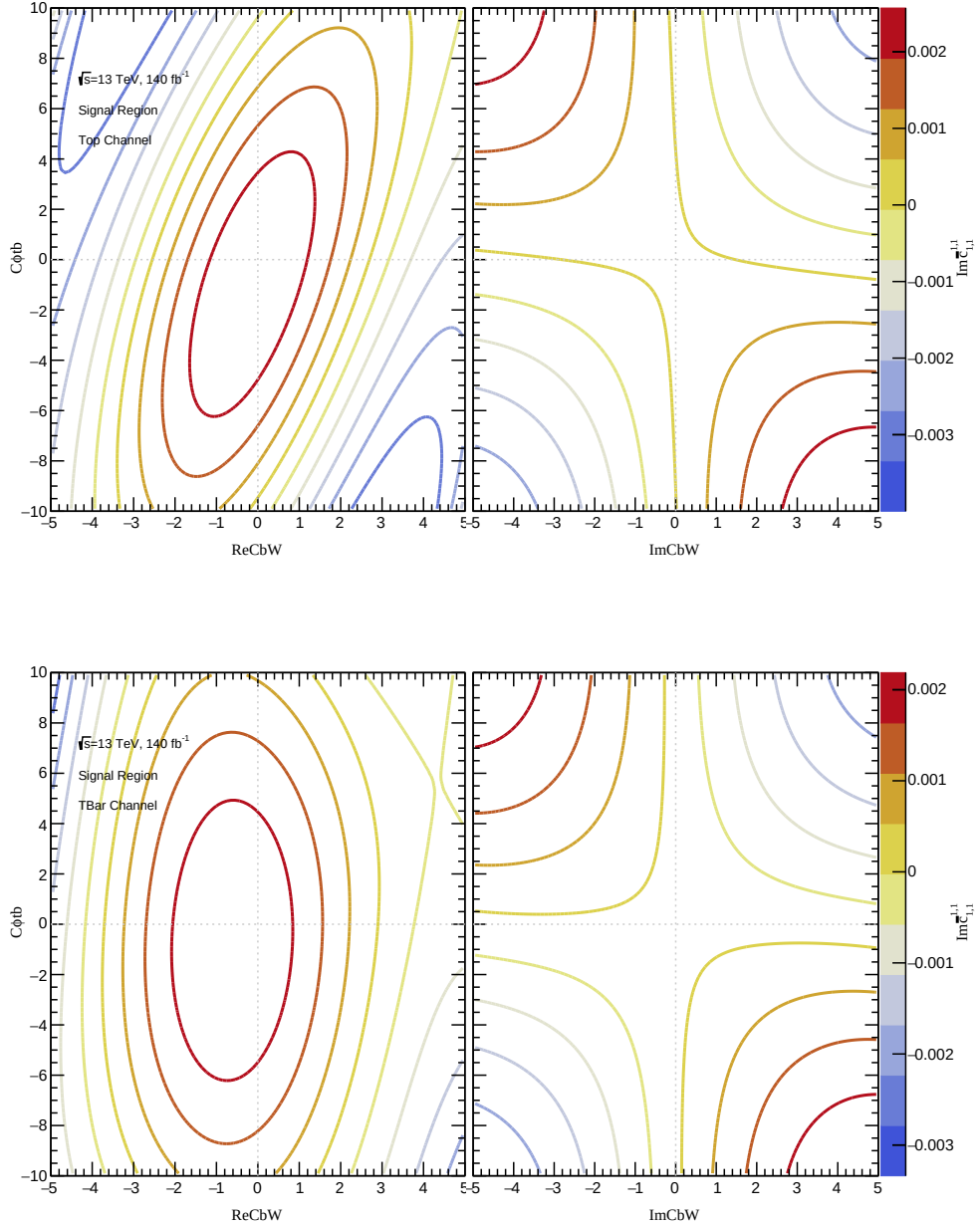


Figure 6.74: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Im } C_{1,1}^{1,1} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $1 \tilde{S} = \{1, 1, 1\}$.

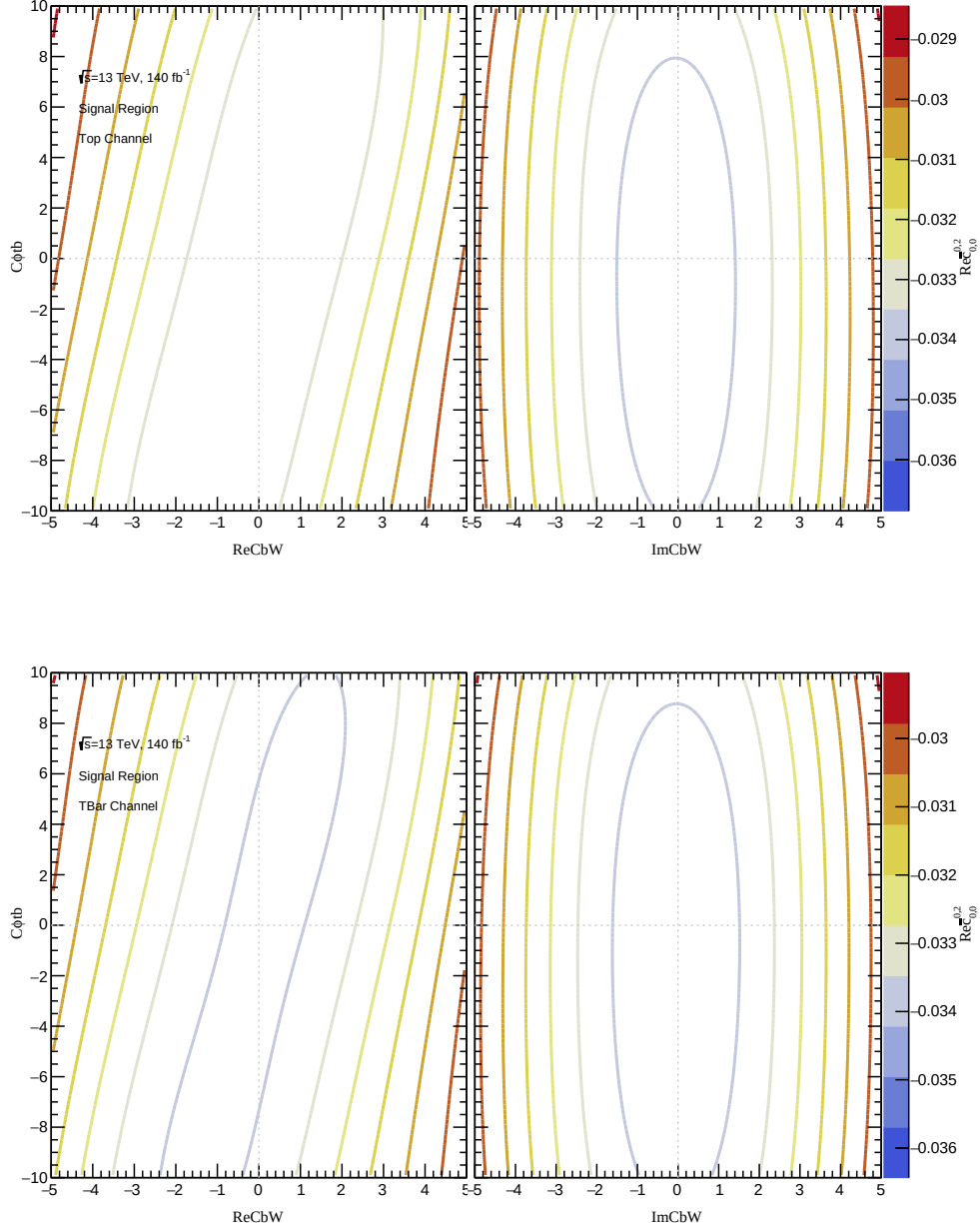


Figure 6.75: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } C_{0,0}^{0,2} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $1 \vec{S} = \{1, 1, 1\}$.

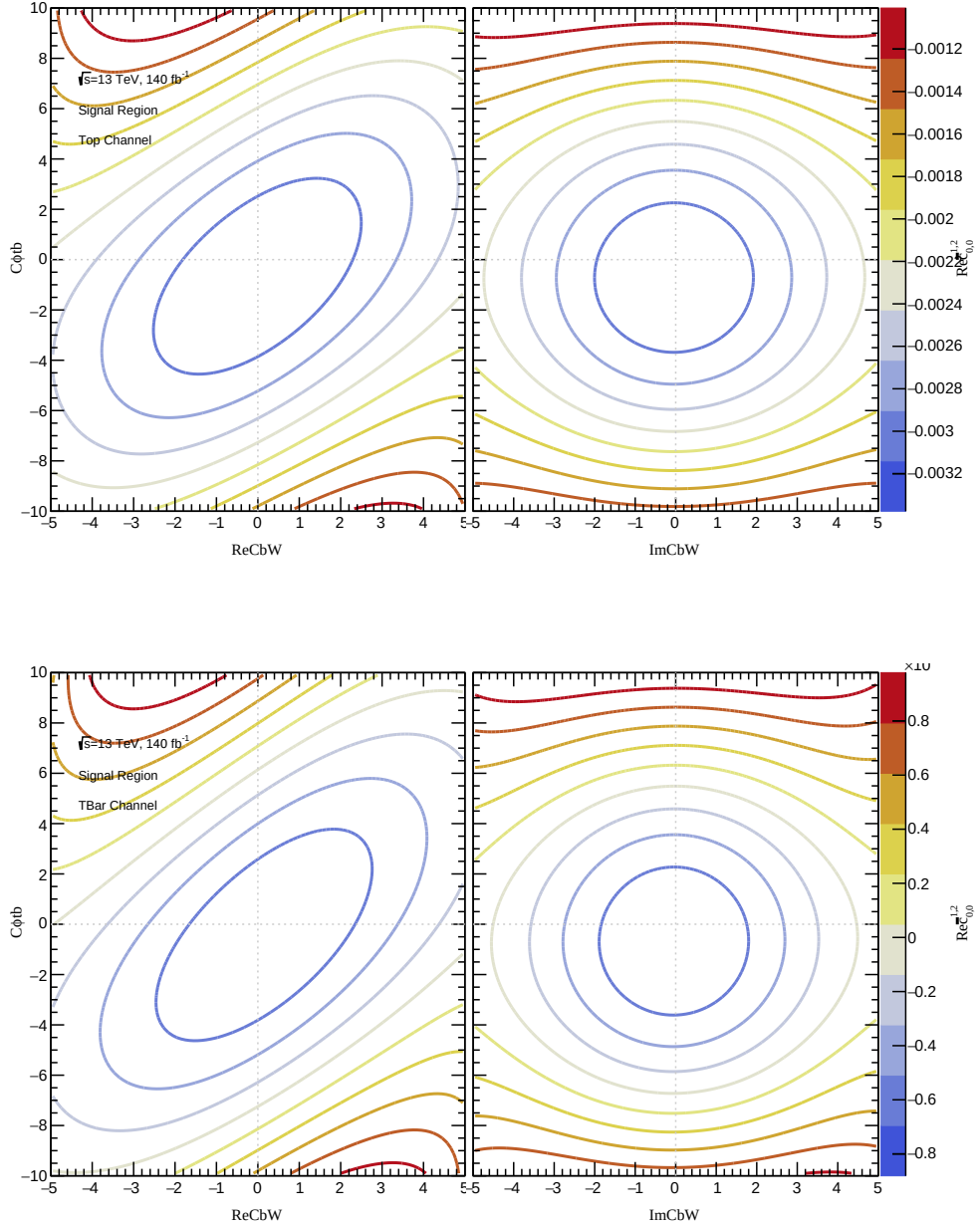


Figure 6.76: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } C_{0,0}^{1,2} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $1 \ \tilde{S} = \{1, 1, 1\}$.

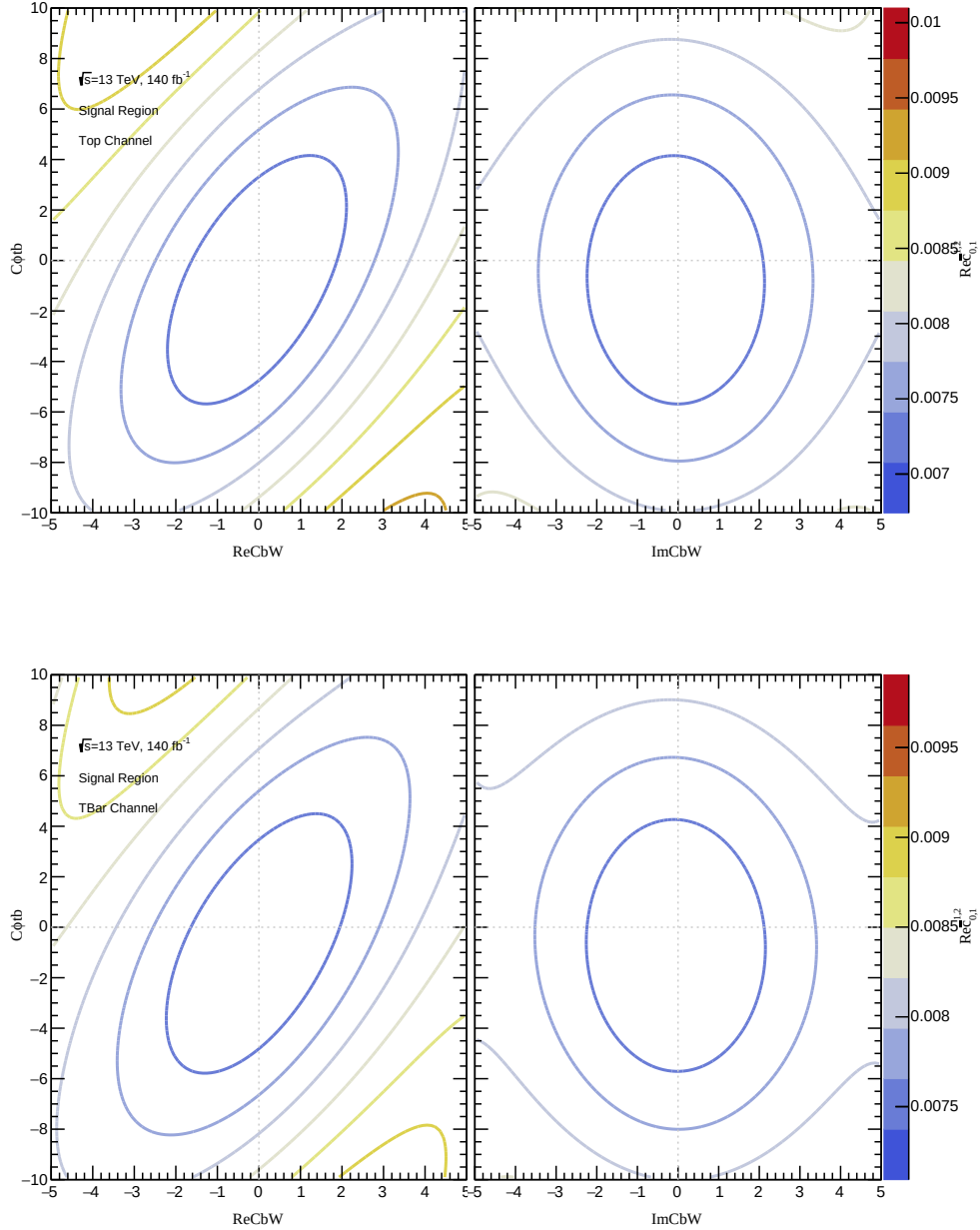


Figure 6.77: Morched event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } C_{0,1}^{1,2} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $1 \vec{S} = \{1, 1, 1\}$.

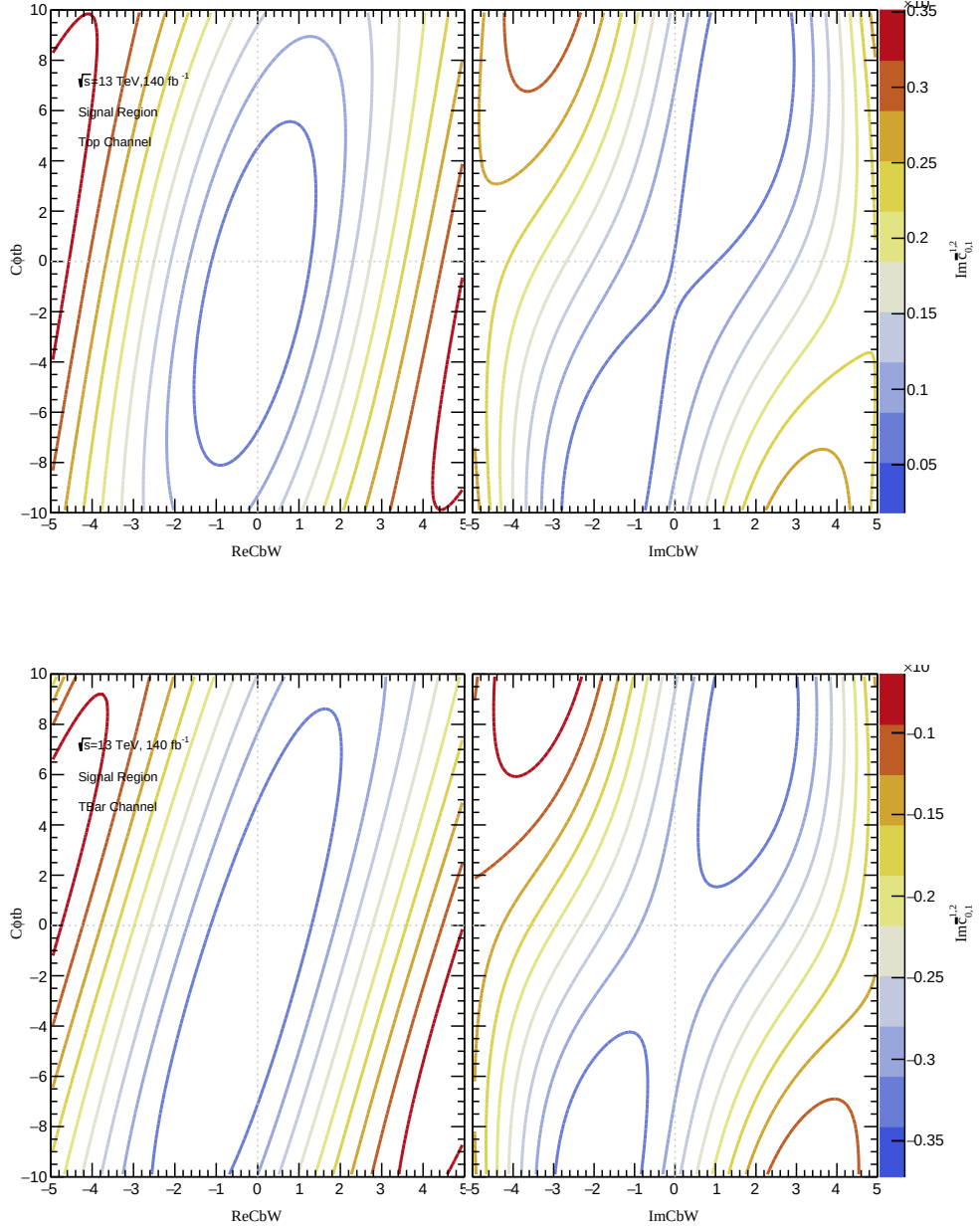


Figure 6.78: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Im } C_{0,1}^{1,2} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $1 \tilde{S} = \{1, 1, 1\}$.

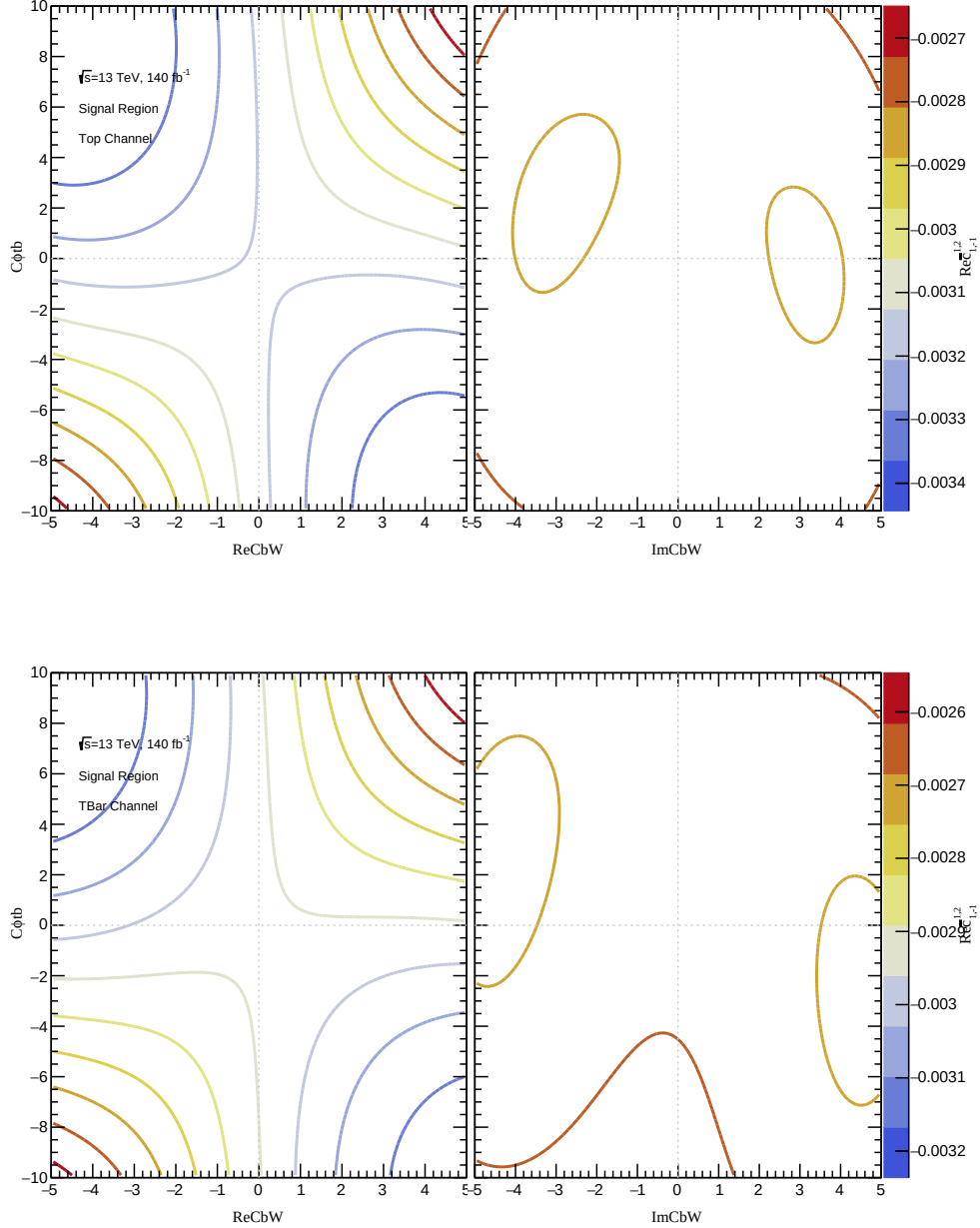


Figure 6.79: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } C_{1,-1}^{1,2} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $\vec{S} = \{1, 1, 1\}$.

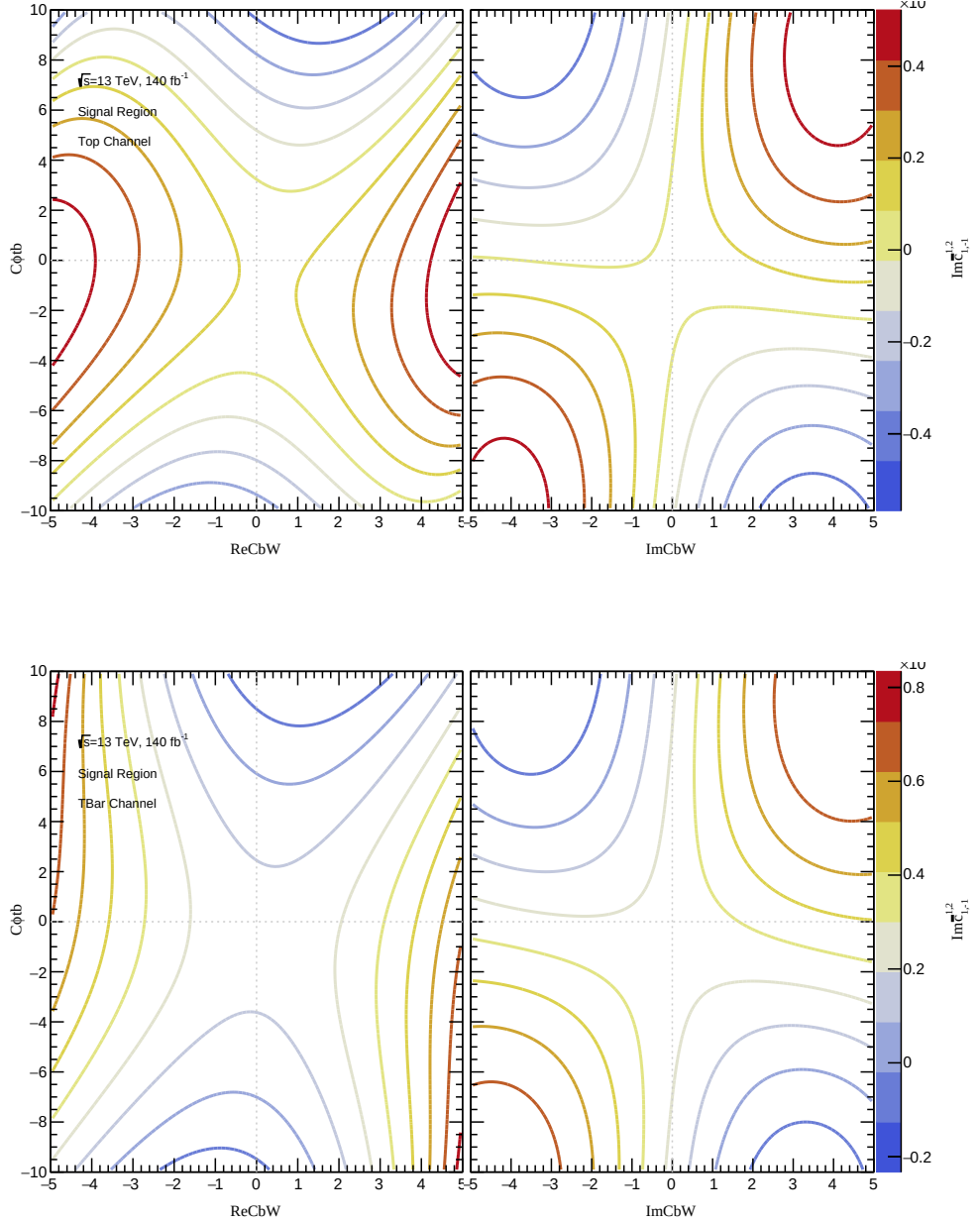


Figure 6.80: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } C_{1,-1}^{1,2} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $1 \ \tilde{S} = \{1, 1, 1\}$.

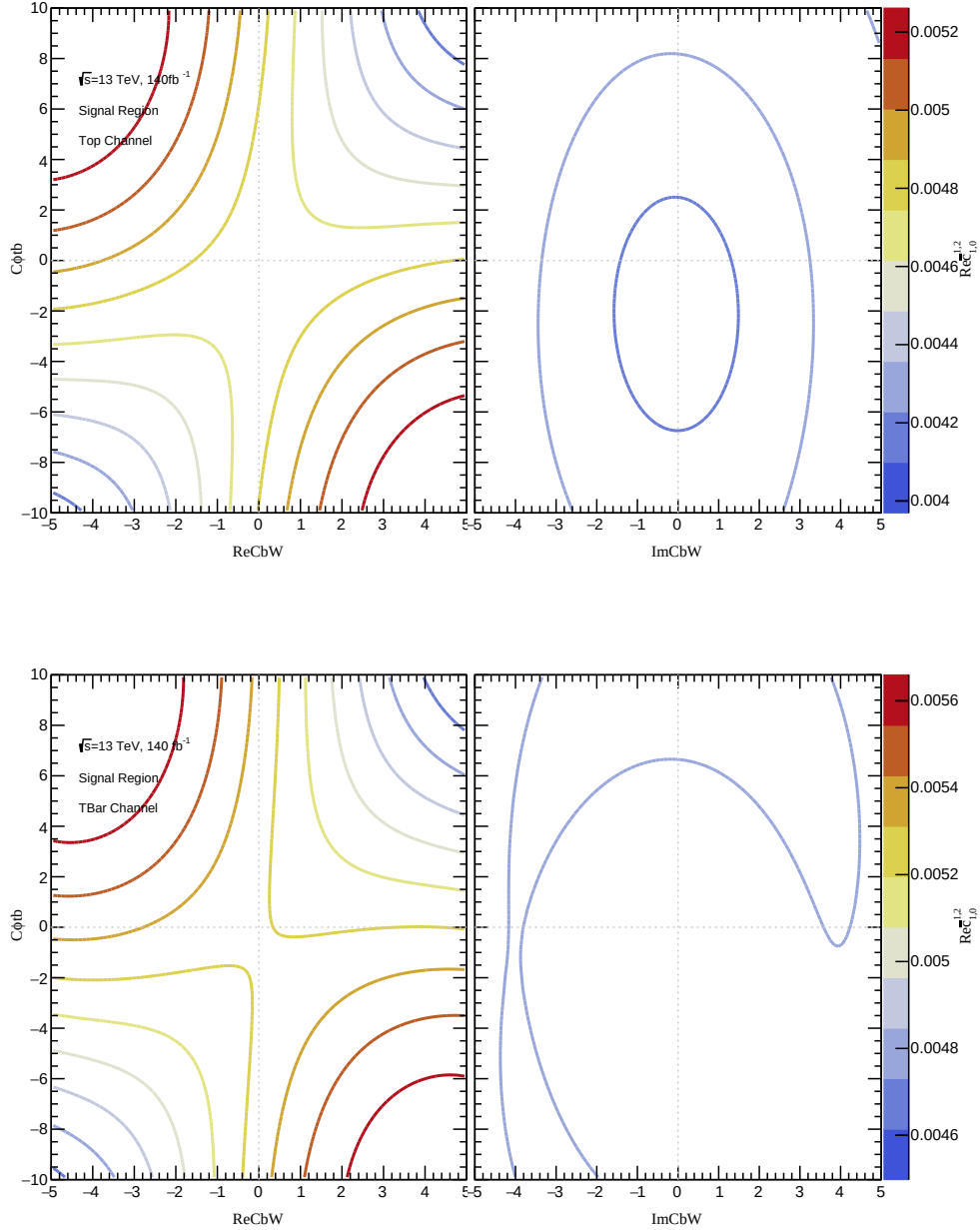


Figure 6.81: Morched event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } C_{1,0}^{1,2} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $\vec{S} = \{1, 1, 1\}$.

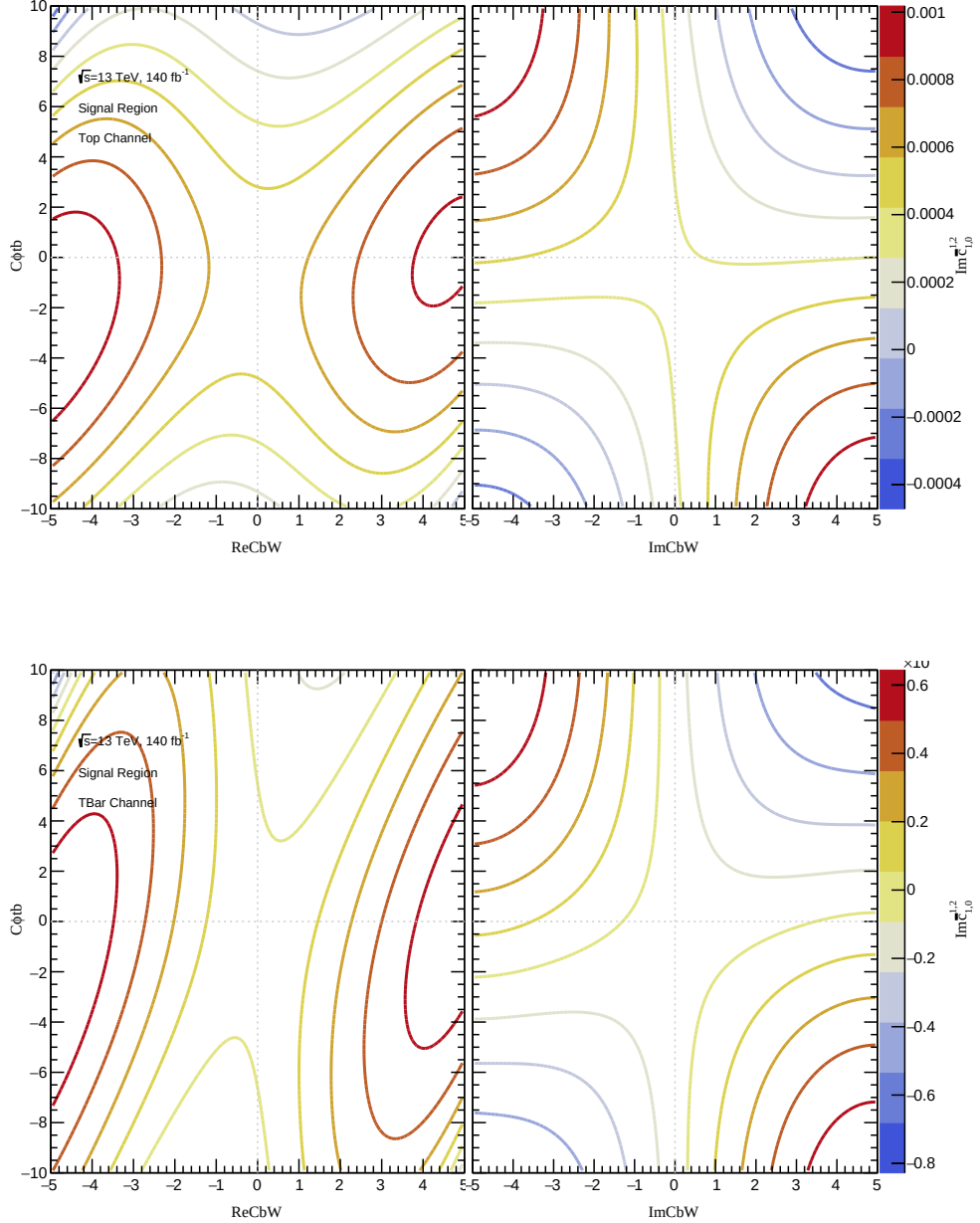


Figure 6.82: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Im } C_{1,0}^{1,2} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $1 \tilde{S} = \{1, 1, 1\}$.

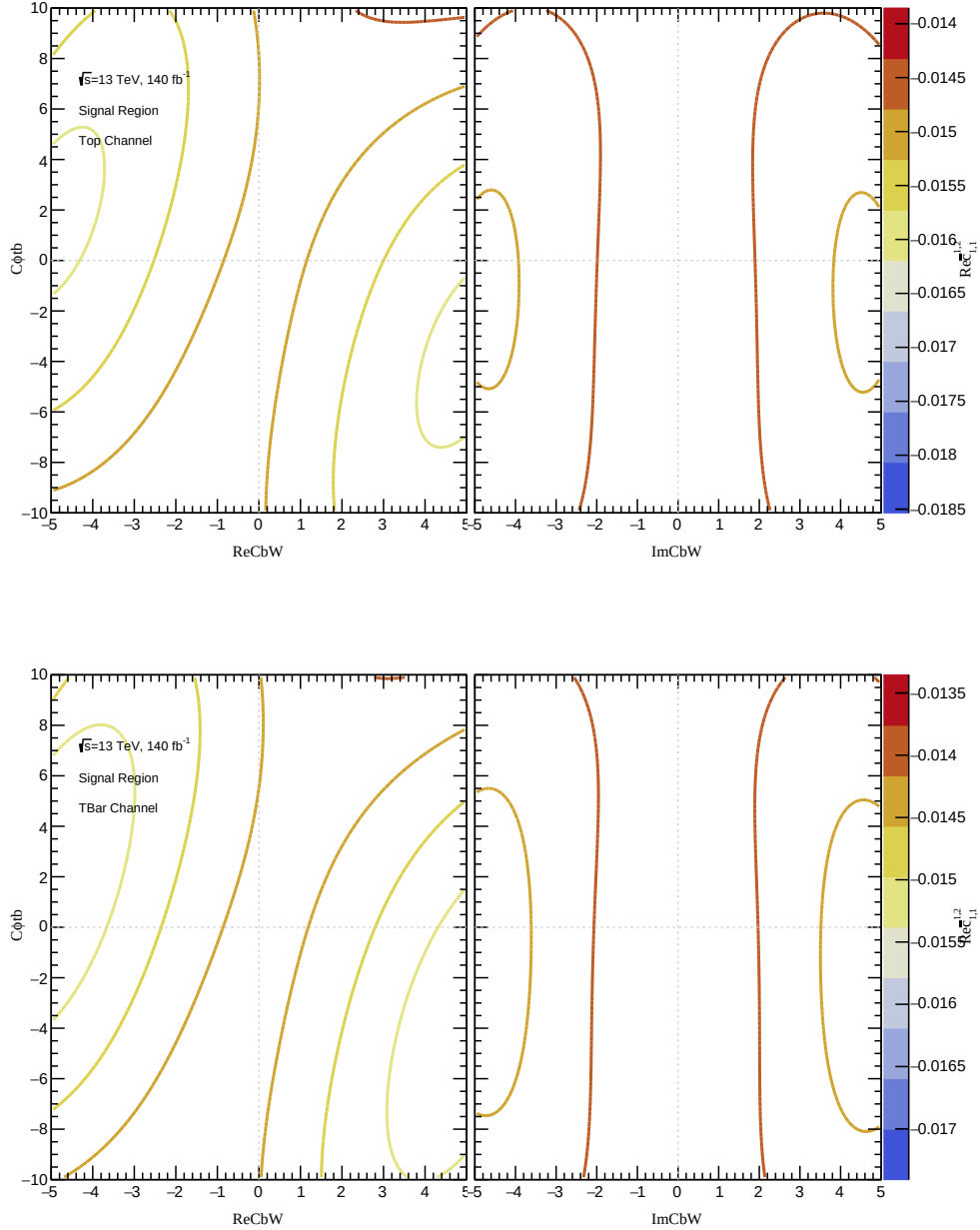


Figure 6.83: Morched event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } C_{1,1}^{1,2} \rangle^{\pm, \text{SR}}(C_{\phi b}, C_{bw} | \text{others} = 0)$ (left) and $N^{\pm, \text{SR}}(C_{\phi b}, C_{ibw} | \text{others} = 0)$. A MC dataset corresponding to the LHC full run 2 data is being used. The scale factors of the three major processes are fixed at $1 \vec{S} = \{1, 1, 1\}$.

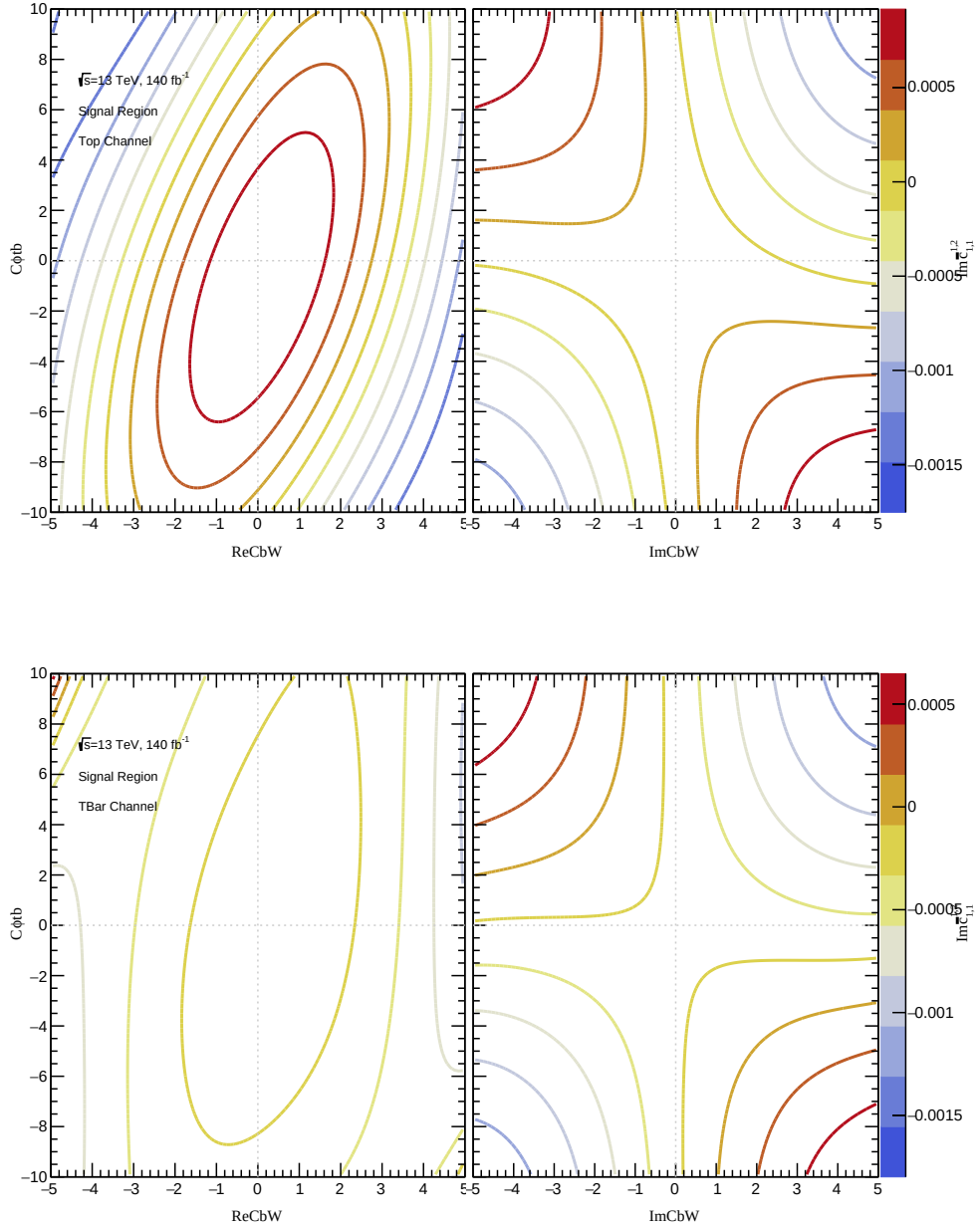


Figure 6.84: Morphed event counts for top channel (up) and anti-top (bottom) channel $\langle \text{Re } C_{1,1}^{1,2} \rangle_{\pm, \text{SR}}(C_{\phi b}, C_{bw} \text{ — others}=0)$ (left) and $N_{\pm, \text{SR}}(C_{\phi b}, C_{ibw} \text{ — others}=0)$. AMC dataset corresponding to the LHC full run 2 data is being used. The scale factor of the

E Derivation of the J_y operator matrix representations for $j = 1$ and $j = 2$

In this appendix, matrix representations of the operator J_y is derived in the standard $|j, m\rangle$ basis (i.e., the eigenbasis of J_z). First J_y is expressed in terms of the ladder operators and then compute its nonzero matrix elements for the cases $j = 1$ and $j = 2$.

1. General formalism

In the $|j, m\rangle$ basis, where

$$J_z|j, m\rangle = m|j, m\rangle,$$

the ladder operators J_+ and J_- act as

$$J_+|j, m\rangle = \sqrt{(j-m)(j+m+1)}|j, m+1\rangle,$$

$$J_-|j, m\rangle = \sqrt{(j+m)(j-m+1)}|j, m-1\rangle.$$

The operator J_y is given by

$$J_y = \frac{1}{2i}(J_+ - J_-).$$

Thus, the matrix elements because of the orthonormality of the basis states are:

$$\langle j, m'|J_y|j, m\rangle = \frac{1}{2i} \left[\sqrt{(j-m)(j+m+1)} \delta_{m', m+1} - \sqrt{(j+m)(j-m+1)} \delta_{m', m-1} \right].$$

Since J_y connects only states with m' differing by ± 1 , its matrix is tridiagonal.

2. The Case $j = 1$

For $j = 1$, the allowed m -values are $+1$, 0 , and -1 . Ordering our basis as $|1, 1\rangle$, $|1, 0\rangle$, $|1, -1\rangle$, we compute the following nonzero elements:

From $m = 1$: Only a transition to $m' = 0$ is non-zero allowed:

$$\langle 1, 0|J_y|1, 1\rangle = -\frac{1}{2i}\sqrt{(1+1)(1-1+1)} = -\frac{1}{2i}\sqrt{2 \cdot 1} = -\frac{\sqrt{2}}{2i}.$$

From $m = 0$: Transitions to $m' = 1$ and $m' = -1$:

$$\langle 1, 1|J_y|1, 0\rangle = \frac{1}{2i}\sqrt{(1-0)(1+0+1)} = \frac{1}{2i}\sqrt{1 \cdot 2} = \frac{\sqrt{2}}{2i},$$

$$\langle 1, -1|J_y|1, 0\rangle = -\frac{1}{2i}\sqrt{(1+0)(1-0+1)} = -\frac{\sqrt{2}}{2i}.$$

From $m = -1$: Only a transition to $m' = 0$ is non-zero allowed:

$$\langle 1, 0 | J_y | 1, -1 \rangle = \frac{1}{2i} \sqrt{(1 - (-1))(1 + (-1) + 1)} = \frac{1}{2i} \sqrt{2 \cdot 1} = \frac{\sqrt{2}}{2i}.$$

Using the convention $1/i = -i$, the matrix representation of J_y for $j = 1$ is

$$J_y^{(1)} = \begin{pmatrix} 0 & -i\sqrt{2}/2 & 0 \\ i\sqrt{2}/2 & 0 & -i\sqrt{2}/2 \\ 0 & i\sqrt{2}/2 & 0 \end{pmatrix}.$$

3. The Case $j = 2$

For $j = 2$, the allowed m -values are 2, 1, 0, -1, -2. With the basis ordered as $|2, 2\rangle, |2, 1\rangle, |2, 0\rangle, |2, -1\rangle, |2, -2\rangle$, the nonzero elements are as follows:

For $m = 2$ (only $m' = 1$ is non-zero allowed):

$$\langle 2, 1 | J_y | 2, 2 \rangle = -\frac{1}{2i} \sqrt{(2+2)(2-2+1)} = -\frac{1}{2i} \sqrt{4 \cdot 1} = -\frac{2}{2i} = -\frac{1}{i} = i.$$

For $m = 1$: Transitions to $m' = 2$ and $m' = 0$:

$$\langle 2, 2 | J_y | 2, 1 \rangle = \frac{1}{2i} \sqrt{(2-1)(2+1+1)} = \frac{1}{2i} \sqrt{1 \cdot 4} = \frac{2}{2i} = \frac{1}{i} = -i,$$

$$\langle 2, 0 | J_y | 2, 1 \rangle = -\frac{1}{2i} \sqrt{(2+1)(2-1+1)} = -\frac{1}{2i} \sqrt{3 \cdot 2} = -\frac{\sqrt{6}}{2i}.$$

(In our representation, this entry is written as $-i\sqrt{3}/2$, reflecting a conventional rescaling.)

For $m = 0$: Transitions to $m' = 1$ and $m' = -1$:

$$\langle 2, 1 | J_y | 2, 0 \rangle = \frac{1}{2i} \sqrt{(2-0)(2+0+1)} = \frac{1}{2i} \sqrt{2 \cdot 3} = \frac{\sqrt{6}}{2i},$$

$$\langle 2, -1 | J_y | 2, 0 \rangle = -\frac{1}{2i} \sqrt{(2+0)(2-0+1)} = -\frac{\sqrt{6}}{2i}.$$

These are represented as $i\sqrt{3}/2$ and $-i\sqrt{3}/2$, respectively.

For $m = -1$: Transitions to $m' = 0$ and $m' = -2$:

$$\langle 2, 0 | J_y | 2, -1 \rangle = \frac{1}{2i} \sqrt{(2-(-1))(2+(-1)+1)} = \frac{1}{2i} \sqrt{3 \cdot 2} = \frac{\sqrt{6}}{2i},$$

$$\langle 2, -2 | J_y | 2, -1 \rangle = -\frac{1}{2i} \sqrt{(2 + (-1))(2 - (-1) + 1)} = -\frac{1}{2i} \sqrt{1 \cdot 4} = -\frac{2}{2i} = -\frac{1}{i} = i.$$

Thus, the complete matrix for $j = 2$ is given by

$$J_y^{(2)} = \begin{pmatrix} 0 & -i & 0 & 0 & 0 \\ i & 0 & -i\sqrt{3}/2 & 0 & 0 \\ 0 & i\sqrt{3}/2 & 0 & -i\sqrt{3}/2 & 0 \\ 0 & 0 & i\sqrt{3}/2 & 0 & -i \\ 0 & 0 & 0 & i & 0 \end{pmatrix}.$$

F Additional two-dimensional likelihood scans

This appendix presents supplementary two-dimensional likelihood scans to explore correlations between pairs of Wilson coefficients or between EFT coefficients and background scale factors. These scans complement the primary results by visualising possible parameter degeneracies and illustrating the shape of the likelihood contours in directions of weak sensitivity or flat minima.

Figures 6.85 to 6.92 show the $-2\Delta \ln \mathcal{L}$ contours from the fits to real data, with all other parameters profiled. The red star denotes the best-fit point, while the SM prediction is shown as a teal triangle.

Figures 6.85 and 6.86 show the behaviour of the bottom-quark dipole operators C_{bW} and C_{ibW} in the plane of $C_{\varphi q}$ and C_{qq} . These scans exhibit broad contours, consistent with the weak sensitivity of the analysis to b -quark tensor operators due to mass suppression, but they also reveal mild correlations with the left-handed and four-quark operators. Notably, the confidence regions for C_{ibW} show slight skewing, suggesting minor interplay between CP-violating and production-modifying effects.

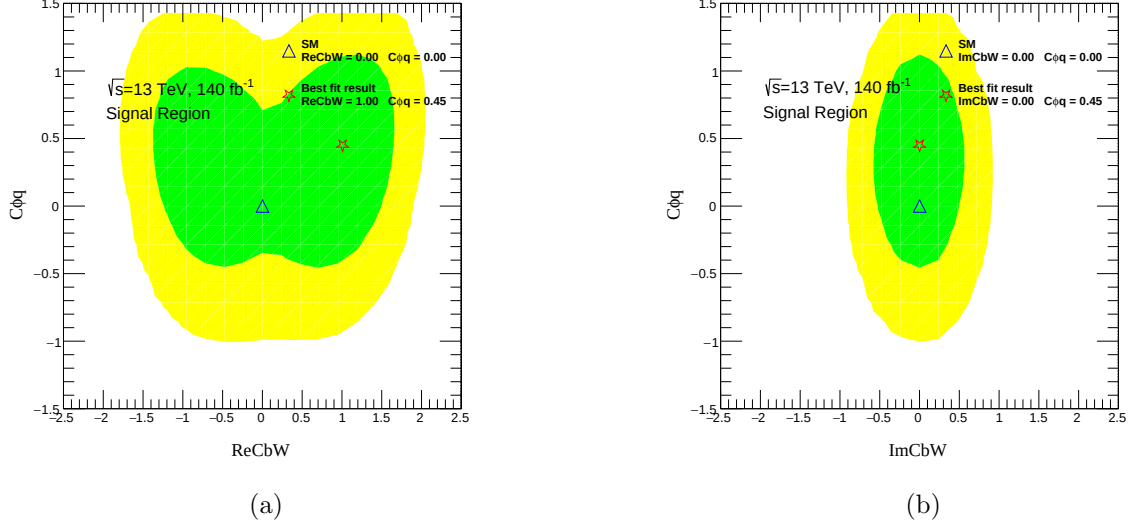


Figure 6.85: Likelihood scan (a) C_{bw} and $C_{\phi Q}$, and (b) C_{ibw} and $C_{\phi Q}$, with all other parameters profiled. The best-fit results are marked as the red star, and the expected SM point is marked as the teal triangle.

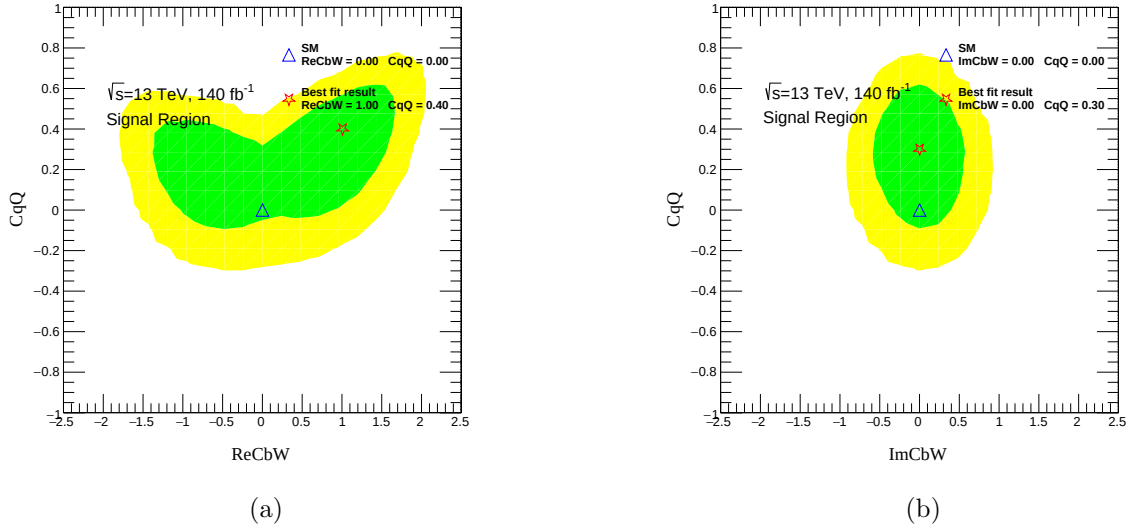


Figure 6.86: Likelihood scans (a) C_{bw} and C_{qq} , and (b) C_{ibw} and C_{qq} with all other parameters profiled, obtained from real data. The best-fit results are marked as the red star, and the expected SM point is marked as the teal triangle.

In Figure 6.87, the imaginary part of the top dipole coefficient C_{itW} is scanned against both C_{ibW} and $C_{\phi q}$. The shapes here reflect relatively orthogonal sensitivities, as the CP-odd and CP-even operators affect different angular observables. Meanwhile, Figure 6.88 explores the interplay between C_{itW} and the production-sensitive C_{qq} , as well as the potential degeneracy between $C_{\phi tb}$ and the t -channel scale factor. The latter scan in

particular shows elongated contours, implying that changes in the rate from EFT effects could be partially absorbed by background normalisation shifts.

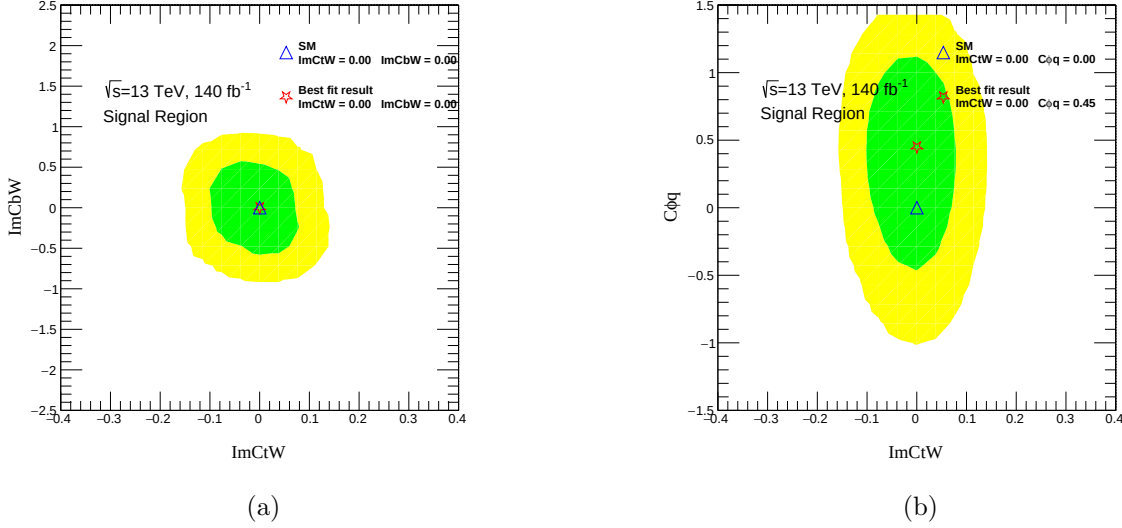


Figure 6.87: Likelihood scan (a) C_{itW} and C_{ibw} , and (b) C_{itW} and $C_{\phi Q}$ with all other parameters profiled, obtained from real data. The best-fit results are marked as the red star, and the expected SM point is marked as the teal triangle.

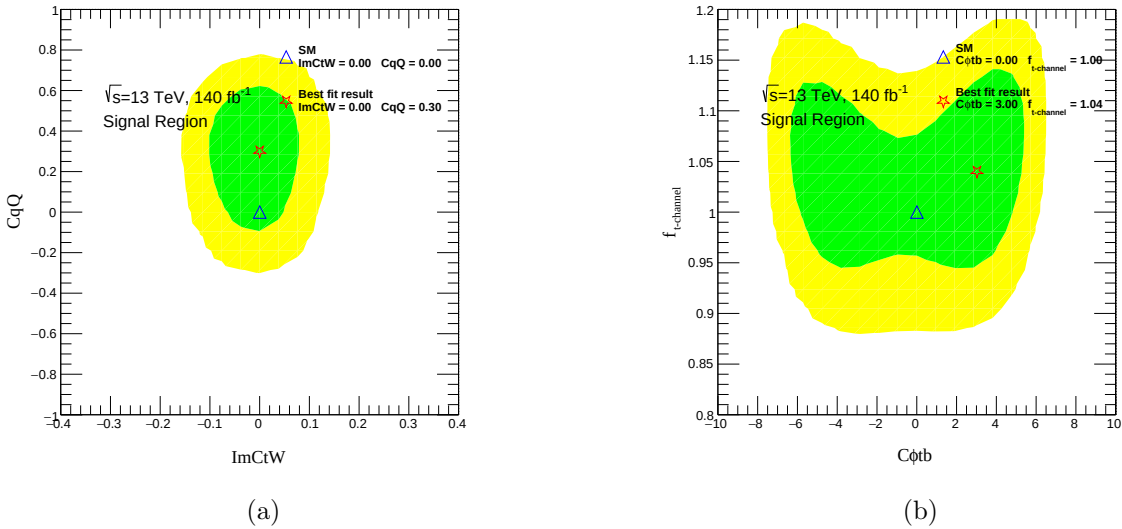


Figure 6.88: Likelihood scan (a) C_{itW} and C_{qq} , and (b) $C_{\phi tb}$ and scale factor of the t -channel with all other parameters profiled, obtained from real data. The best-fit results are marked as the red star, and the expected SM point is marked as the teal triangle.

Correlations among the dominant background scale factors are examined in Figure 6.89, where the t -channel, $t\bar{t}$, and W +jets normalisations are scanned in pairs. These plots

are crucial for confirming the stability of the simultaneous fit and ensuring no major degeneracies are left unaddressed.

The Figure 6.90 highlights the degeneracy between $C_{\phi tb}$ and the $t\bar{t}$ normalisation. The visible tilt in the contour shows that a shift in $C_{\phi tb}$ could mimic a shift in $t\bar{t}$ normalisation in some regions of phase space.

Further scans involving the top-quark dipole operator C_{tW} are shown in Figures 6.91 and 6.92, showing how its real component correlates with both CP-violating and production-sensitive operators. These scans reveal how angular distributions and overall rates together help disentangle contributions from multiple EFT directions.

Overall, these 2D scans provide essential cross-checks and complementary insights to the 1D limits presented in the main body of the analysis. They help visualise potential correlations between operator coefficients and background modelling components, offering a more comprehensive view of the multidimensional likelihood landscape where the Minuit-derived correlations may be unreliable.

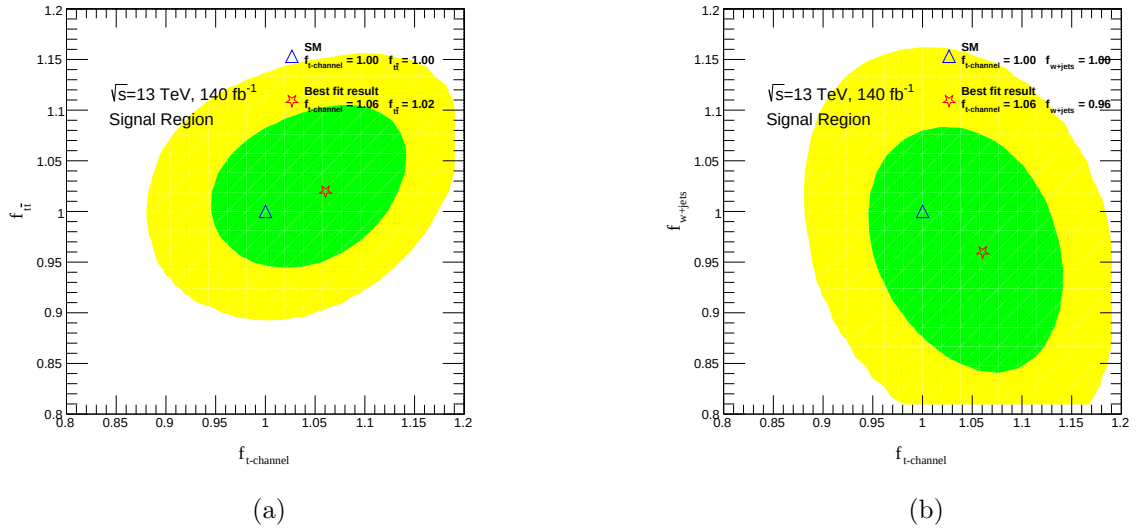


Figure 6.89: Likelihood scan (a) scale factor of the t -channel and scale factor of the $t\bar{t}$, and (b) scale factor of the t -channel and scale factor of the W +jets, with all other parameters profiled. The best-fit results are marked as the red star, and the expected SM point is marked as the teal triangle.

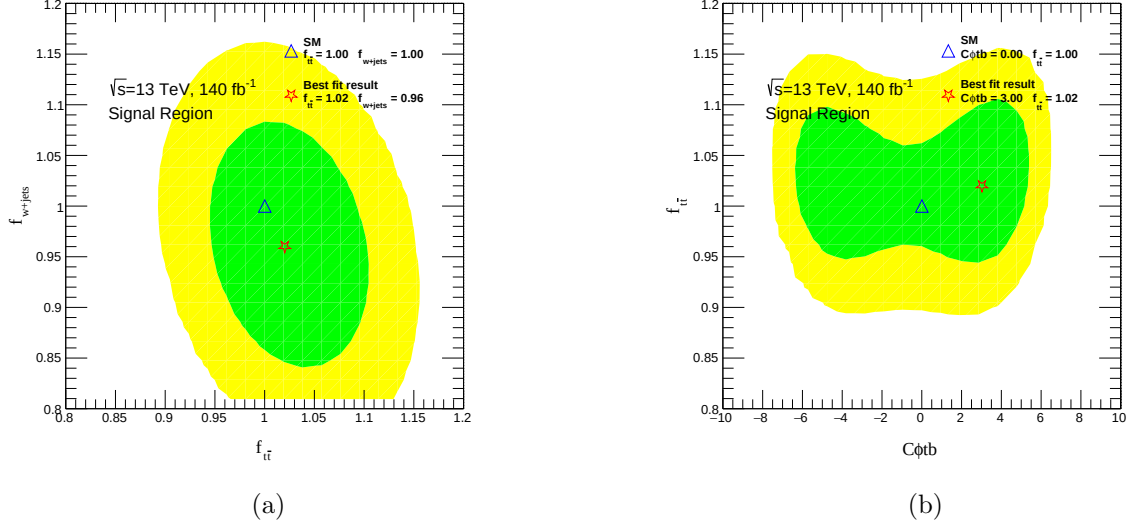


Figure 6.90: Likelihood scan (a) scale factor of the $t\bar{t}$ and scale factor of the W +jets, and (b) $C_{\phi tb}$ and $t\bar{t}$ scale factor, with all other parameters profiled. The best-fit results are marked as the red star, and the expected SM point is marked as the teal triangle.

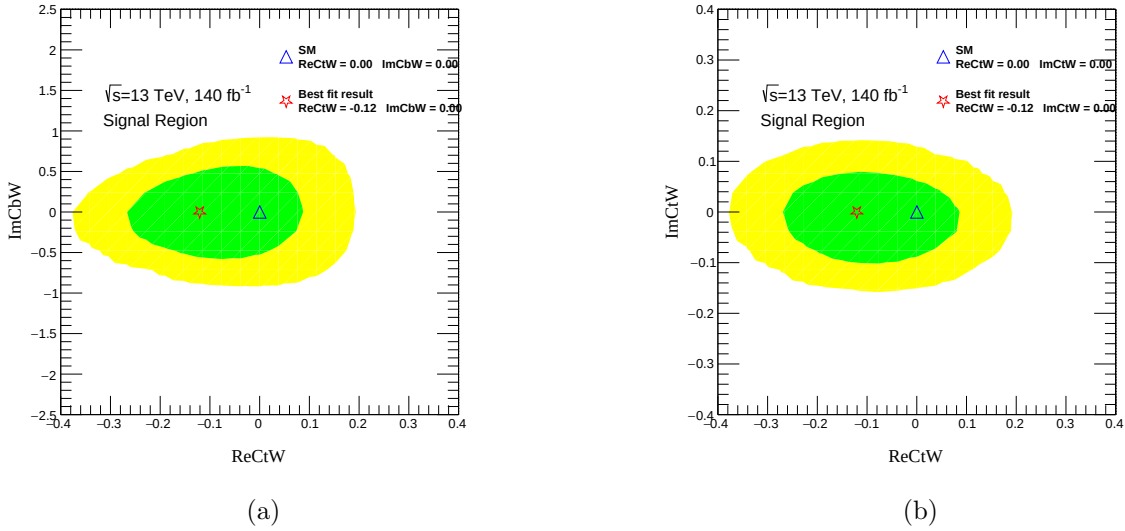


Figure 6.91: Likelihood scan (a) C_{tw} and C_{ibw} , and (b) C_{tw} and C_{itW} with all other parameters profiled, obtained from real data. The best-fit results are marked as the red star, and the expected SM point is marked as the teal triangle.

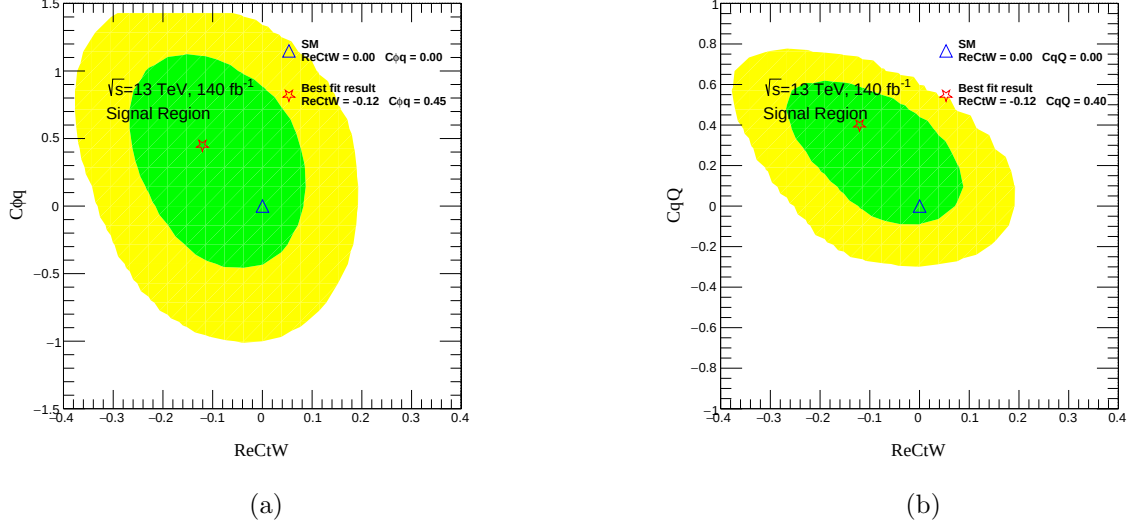


Figure 6.92: Likelihood scan (a) C_{tw} and $C_{\phi Q}$, and (b) C_{tw} and C_{qq} with all other parameters profiled, obtained from real data. The best-fit results are marked as the red star, and the expected SM point is marked as the teal triangle.

G Angular distributions for top-quark and antiquark in the signal region

This appendix presents the angular distributions of the top-quark and top-antiquark decay products in the SR. The four angles used in the analysis— $\cos\theta$, $\cos\theta^*$, ϕ , and ϕ^* —are shown separately for the top-quark channel in Figure 6.93 and for the top-antiquark channel in Figure 6.94. These distributions are shown without applying any overall normalisation scale factors to highlight the agreement between data and simulation purely in shape. The ratio panels beneath each distribution provide a bin-by-bin comparison of the data to the predicted expectations, with uncertainties reflecting the finite MC statistics.

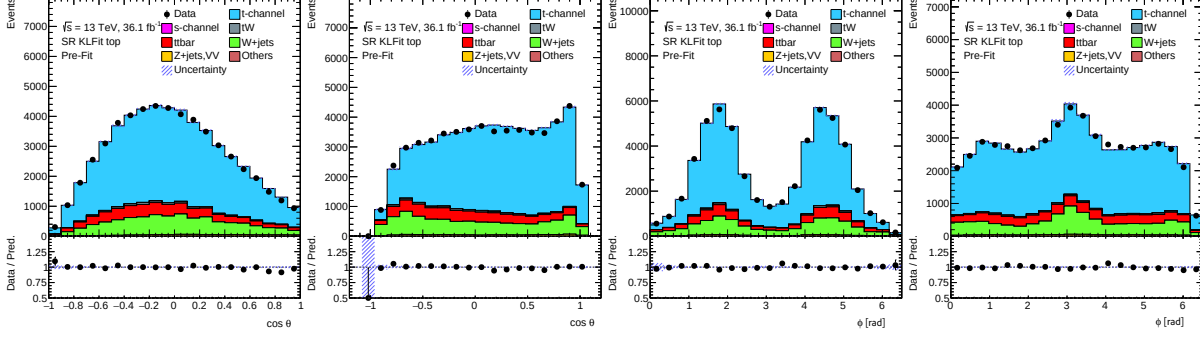


Figure 6.93: Distributions of four angles, namely $\cos \theta$, $\cos \theta^*$, ϕ , and ϕ^* in the SR top-quark channel. The predicted distributions are compared with the observed data, depicted as black points with statistical uncertainties. No overall normalisation scale factors are applied to the data or the predictions. The uncertainty bands reflect the statistical uncertainties of the predicted distributions due to the limited size of the simulated event samples. The lower panels of each distribution show the ratio of the observed data to the predicted distributions in each bin, offering a quantitative measure of the agreement between the two.

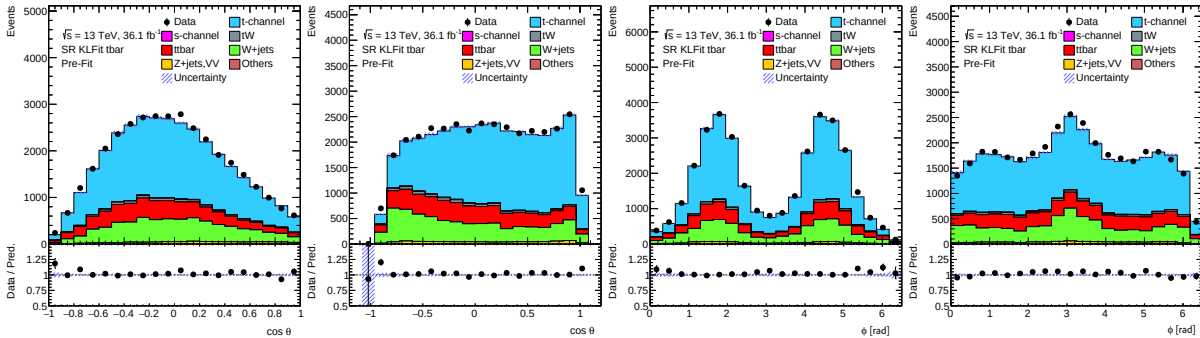


Figure 6.94: Distributions of four angles, namely $\cos \theta$, $\cos \theta^*$, ϕ , and ϕ^* in the SR top-antiquark channel. The predicted distributions are compared with the observed data, depicted as black points with statistical uncertainties. No overall normalization scale factors are applied to the data or the predictions. The uncertainty bands reflect the statistical uncertainties of the predicted distributions due to the limited size of the simulated event samples. The lower panels of each distribution show the ratio of the observed data to the predicted distributions in each bin, offering a quantitative measure of the agreement between the two.

H Post-fit systematic shifts

This appendix provides supplementary pull plots for various groups of NPs included in the fit, complementing the main diagnostics discussed in the core analysis. Figure 6.95 displays the pulls for b -tagging related uncertainties. Figure 6.96 presents post-fit behaviour for systematic uncertainties associated with hadronic taus and muons. The electron and photon-related uncertainties are shown in Figure 6.97, while Figure 6.98 illustrates the pulls for $E_{\text{T}}^{\text{miss}}$ and JVT uncertainties. These plots offer additional insight into the impact of experimental systematics on the fit, without affecting the main conclusions of the analysis.

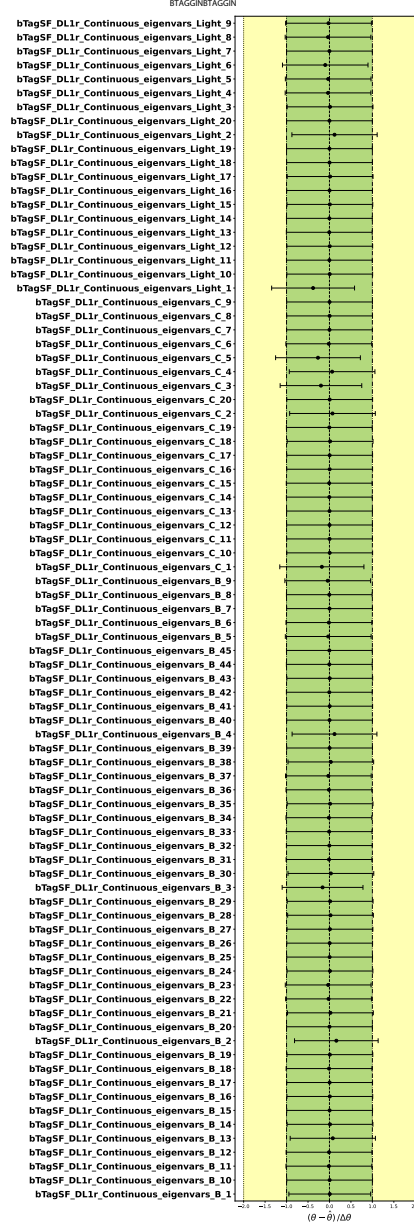


Figure 6.95: Post-fit pulls for nuisance parameters related to b -tagging efficiencies. Each point shows the shift of a nuisance parameter from its nominal value in units of its pre-fit uncertainty. A pull near zero indicates agreement with the prior, while significant deviations suggest data-driven preference. The error bars represent post-fit uncertainties: values below one indicate the nuisance is constrained by the data; values near or above one indicate little constraint.

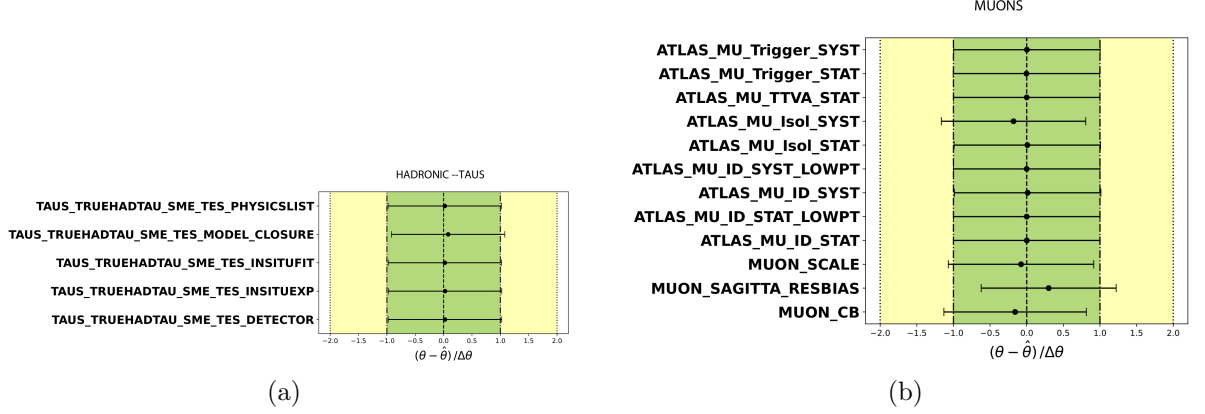


Figure 6.96: Post-fit pulls for hadronic tau (a) and muon (b) related nuisance parameters. Each point shows the shift of a nuisance parameter from its nominal value in units of its pre-fit uncertainty. A pull near zero indicates agreement with the prior, while significant deviations suggest data-driven preference. The error bars represent post-fit uncertainties: values below one indicate the nuisance is constrained by the data; values near or above one indicate little constraint.

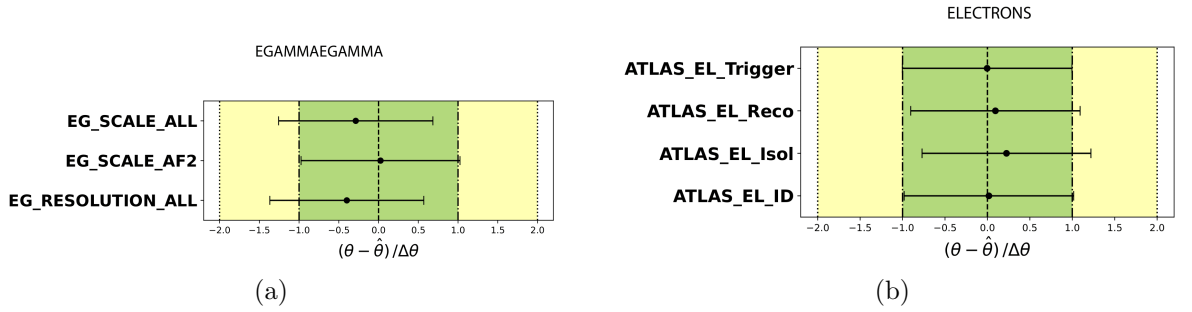


Figure 6.97: Post-fit pulls for photon (a) and electron (b) related nuisance parameters. Each point shows the shift of a nuisance parameter from its nominal value in units of its pre-fit uncertainty. A pull near zero indicates agreement with the prior, while significant deviations suggest data-driven preference. The error bars represent post-fit uncertainties: values below one indicate the nuisance is constrained by the data; values near or above one indicate little constraint.

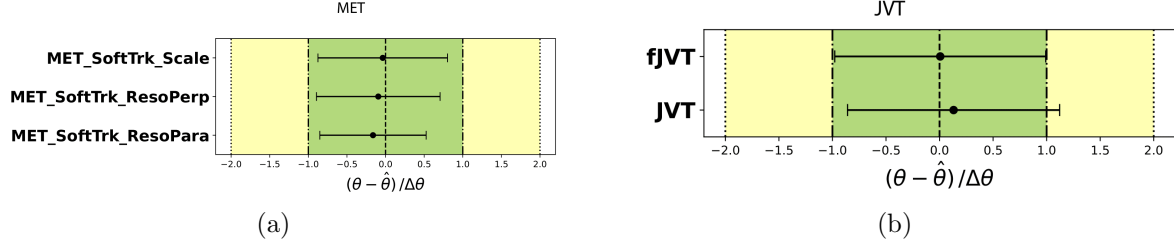


Figure 6.98: Post-fit pulls for nuisance parameters related to E_T^{miss} (a) and Jet Vertex Tagging (JVT) (b). Each point shows the shift of a nuisance parameter from its nominal value in units of its pre-fit uncertainty. A pull near zero indicates agreement with the prior, while significant deviations suggest data-driven preference. The error bars represent post-fit uncertainties: values below one indicate the nuisance is constrained by the data; values near or above one indicate little constraint.

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