

Faculté des sciences

Observables sensitive to 4-quark operators

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Chapter 1

Introduction

One of the most profound questions in science is whether the complexity of the Universe that we live in — from the motion of galaxies down to the behaviour of the smallest known constituents of matter — can be reduced to a small set of fundamental laws. The history of physics is, in many ways, the history of that reduction. The unification of electricity and magnetism by Maxwell in the nineteenth century, the discovery of quantum mechanics in the early twentieth century, and the development of quantum field theory each represented a dramatic simplification of our picture of Nature : fewer ingredients, more explanatory power. By the 1970s, these efforts had culminated in the Standard Model of particle physics (SM), a single quantum field theory that describes three of the four known fundamental forces — the electromagnetic, weak, and strong interactions — and classifies all the elementary particles observed to date.

The construction of large particle colliders was central to establishing and testing this framework. Each new generation of machines brought higher energies allowing physicists to probe shorter and shorter distance scales and confront the predictions of the SM with ever-increasing precision. The discovery of the W and Z bosons [1], the top quark [2], and finally of the Higgs boson [3, 4] at the Large Hadron Collider (LHC) each marked a milestone in this programme, confirming the internal consistency of the SM and demonstrating its remarkable predictive power. Today, after decades of precision measurements across a vast range of energies and processes, the SM stands as one of the most thoroughly tested theories in the history of science.

Yet it is widely understood to be incomplete which points unambiguously to the existence of physics beyond the SM. The LHC was originally conceived as a discovery machine, aiming to produce new massive states directly and identify them as resonances. The absence of any such signal after more than a decade of operation has progressively shifted the focus towards indirect searches, where new physics manifests not as a new particle but as small deviations in the properties of known processes. The Standard Model Effective Field Theory (SMEFT) provides the natural language for this programme : by extending the SM Lagrangian with higher-dimensional operators built from SM fields and consistent with its gauge symmetries,

$$\mathcal{L}_{\text{SMEFT}} = \mathcal{L}_{\text{SM}} + \sum_{d>4} \sum_i \frac{C_i^{(d)}}{\Lambda^{d-4}} \mathcal{O}_i^{(d)},$$

one can parametrise in a model-independent way the low-energy imprints of any heavy new physics at a scale Λ too large to be directly produced. The Wilson coefficients $C_i^{(d)}$ are then constrained by comparing SMEFT predictions to experimental data.

In this master's thesis, we study the ten four-light-quark operators of the Warsaw basis, which directly modify quark-quark scattering amplitudes. We investigate their contributions to the associated production of a photon with two jets, $pp \rightarrow \gamma jj$, exploiting the fact that the photon couples differently to the up-type quarks and the down-type quarks and can thereby constrain certain operators more than others. We derive constraints on the Wilson coefficients by studying differential distributions.

The thesis is organised as follows : Chapter 2 reviews the Standard Model, covering the aspects most relevant to this work. Chapter 3 introduces the SMEFT framework, including the Warsaw basis, power counting, and the need for differential observables. Chapter 4 presents the $pp \rightarrow \gamma jj$ analysis setup and derives individual constraint and two dimensional contours on the 4LQ Wilson coefficients from the photon transverse momentum distribution. Chapter 5 develops the interference revival strategy, identifies the directions in coefficient space for which the total interference vanishes, and search for observables that maximise sensitivity to these directions. Chapter 6 extends the analysis beyond the linear level by including $\mathcal{O}(\Lambda^{-4})$ contributions for a subset of operators. Chapter 7 separates the colour-singlet and colour-octet components of selected operators and derives independent constraints on each. We then conclude this thesis with a final chapter that summarises the main results and outlines directions for future work.

Chapter 2

The standard Model of Particle physics

This section provides a concise overview of the Standard Model (SM) of particle physics, covering the essential theoretical background needed to follow the discussions presented in this thesis. Rather than aiming for completeness, the goal is to introduce the key concepts and building blocks of the SM that will be directly relevant in the chapters that follow. For the interested reader, more can be found in [5, 6, 7].

2.1 The Standard Model :

The Standard Model (SM) of particle physics stands as one of the most successful theoretical frameworks in the history of science, providing an accurate and experimentally well-tested description of the elementary constituents of matter and their fundamental interactions. Formulated within the language of quantum field theory, the SM is a gauge theory based on the gauge group

$$G = SU(3)_c \times SU(2)_L \times U(1)_Y \quad (2.1)$$

where the subscripts refer to colour charge, weak isospin, and hypercharge, respectively. This gauge structure entirely determines the nature of the strong, weak, and electromagnetic interactions, with the corresponding coupling constants extracted from experiment.

The particle content of the SM is summarised in Fig.2.1. Matter is described by spin-1/2 fermions, comprising six quarks and six leptons organised into three generations, while the fundamental interactions are mediated by spin-1 gauge bosons: the gluon, the photon, and the $W^\pm$ and Z bosons. The particle spectrum is completed by the Higgs boson, the unique spin-0 scalar of the SM, whose vacuum expectation value triggers electroweak symmetry breaking (EWSB) and is responsible for generating the masses of all elementary particles through the Higgs mechanism [8].

2.1.1 The Lagrangian of the Standard Model :

The complete Lagrangian for the SM is then formed by adding up the different pieces and reads

$$\mathcal{L}_{SM} = \mathcal{L}_{Gauge} + \mathcal{L}_{Fermions} + \mathcal{L}_{Higgs} + \mathcal{L}_{Yukuwa} \quad (2.2)$$

Strictly speaking it is the Lagrangian density

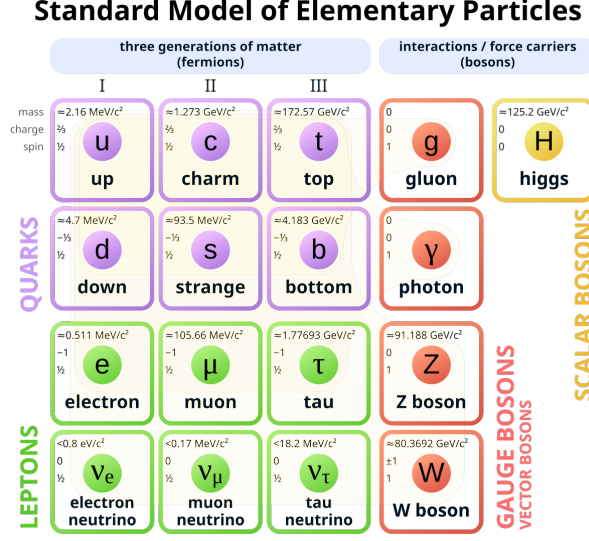


Figure 2.1: Particle content of the SM [9]

which includes the gauge, fermion, Higgs, and Yukawa sectors of the theory. We will then, in the next sections discuss these terms in detail.

2.1.2 The Gauge sector :

The Lagrangian for the gauge sector of the SM is given by

$$\mathcal{L}_{Gauge} = -\frac{1}{4}G_{\mu\nu}^a G_a^{\mu\nu} - \frac{1}{4}W_{\mu\nu}^i W_i^{\mu\nu} - \frac{1}{4}B_{\mu\nu} B^{\mu\nu} \quad (2.3)$$

We have the following definitions for the field strength tensors associated respectively to $SU(3)_c$, $SU(2)_L$ and $U(1)_Y$

$$G_{\mu\nu}^a = \partial_\mu G_\nu^a - \partial_\nu G_\mu^a - g_s f^{abc} G_\mu^b G_\nu^c, \quad a, b, c \in [1, 8], \quad (2.4)$$

$$W_{\mu\nu}^i = \partial_\mu W_\nu^i - \partial_\nu W_\mu^i - g \epsilon^{ijk} W_\mu^j W_\nu^k, \quad i, j, k \in [1, 3], \quad (2.5)$$

$$B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu. \quad (2.6)$$

where g_s and g are the respective gauge coupling constants, the indices $a, b, c = 1, \dots, 8$ label the colour degrees of freedom, and $i, j, k = 1, \dots, 3$ label the weak isospin degrees of freedom. The eight gluon fields G_μ^a and the three weak gauge fields W_μ^i are associated with the generators of $SU(3)_c$ and $SU(2)_L$, defined as $T^a = \lambda^a/2$ and $\tau^i = \sigma^i/2$, respectively, where λ^a are the Gell-Mann matrices and σ^i the Pauli matrices :

$$\lambda^1 = \begin{pmatrix} 0 & 1 & 0 \\ 1 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \quad \lambda^2 = \begin{pmatrix} 0 & -i & 0 \\ i & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \quad \lambda^3 = \begin{pmatrix} 1 & 0 & 0 \\ 0 & -1 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \quad \lambda^4 = \begin{pmatrix} 0 & 0 & 1 \\ 0 & 0 & 0 \\ 1 & 0 & 0 \end{pmatrix}, \quad (2.7)$$

$$\lambda^5 = \begin{pmatrix} 0 & 0 & -i \\ 0 & 0 & 0 \\ i & 0 & 0 \end{pmatrix}, \quad \lambda^6 = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 1 \\ 0 & 1 & 0 \end{pmatrix}, \quad \lambda^7 = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & -i \\ 0 & i & 0 \end{pmatrix}, \quad \lambda^8 = \frac{1}{\sqrt{3}} \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & -2 \end{pmatrix}, \quad (2.8)$$

and

$$\sigma^1 = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}, \quad \sigma^2 = \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix}, \quad \sigma^3 = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}. \quad (2.9)$$

These generators satisfy the following Lie algebra commutation relations :

$$[T^a, T^b] = if^{abc}T^c, \quad [\tau^i, \tau^j] = i\epsilon^{ijk}\tau^k, \quad [Y, Y] = 0, \quad (2.10)$$

where f^{abc} are the $SU(3)_c$ structure constants and ϵ^{ijk} denotes the Levi-Civita tensor. The non-Abelian nature of $SU(3)_c$ and $SU(2)_L$ gives rise to cubic and quartic self-interactions among the gauge fields, a feature that fundamentally distinguishes non-Abelian gauge theories from the Abelian $U(1)$ case.

2.1.3 The Fermion sector :

The matter content of the Standard Model is composed of spin-1/2 fermion fields arranged into three generations. These are divided into quarks and leptons, each generation containing two quarks and two leptons with identical gauge quantum numbers but different masses. The fermion fields transform under the gauge group $SU(3)_c \times SU(2)_L \times U(1)_Y$, and their representation content reflects the chiral nature of the electroweak interaction.

1) Quarks :

The quarks are the fermions that carry colour charge and therefore participate in the strong interaction, which is described by Quantum Chromodynamics (QCD). In the Standard Model, QCD is a non-Abelian gauge theory based on the symmetry group $SU(3)_c$, under which quark fields transform as triplets. This means that each quark comes in three colour states, conventionally labelled red, green, and blue.

An important property of QCD is confinement: coloured particles such as quarks and gluons are not observed as free asymptotic states in nature, but only inside colour-neutral bound states called hadrons, such as mesons and baryons. At the same time, QCD also exhibits asymptotic freedom, meaning that the strong interaction becomes weaker at high energy or short distances. This behaviour is of central importance in high-energy collider physics, where quarks can often be treated as approximately free partons inside the proton.

For each generation $i = 1, 2, 3$, the quark fields are organised as

$$Q_{L_i} = \begin{pmatrix} u_i \\ d_i \end{pmatrix}_L, \quad u_{R_i}, \quad d_{R_i}, \quad (2.11)$$

where Q_{L_i} is a left-handed weak doublet, while u_{R_i} and d_{R_i} are right-handed weak singlets.

2) Leptons :

The leptons are fermions that do not carry a colour charge and are therefore singlets under $SU(3)_c$. As a consequence, they do not couple to gluons and do not participate in QCD interactions. They are nevertheless charged under the electroweak group and therefore interact through the weak and electromagnetic forces, depending on their quantum numbers.

For each generation $i = 1, 2, 3$, the lepton fields are written as

$$L_{L_i} = \begin{pmatrix} \nu_i \\ \ell_i \end{pmatrix}_L, \quad \ell_{R_i}, \quad (2.12)$$

where L_{L_i} is a left-handed weak doublet and ℓ_{R_i} is a right-handed charged-lepton singlet. It is worth noting that in the SM, no right-handed neutrino field is introduced, so neutrinos are massless.

The transformation properties of the fermion fields under $SU(3)_c \times SU(2)_L \times U(1)_Y$ are summarized in Table 2.1

Table 2.1: Gauge quantum numbers of one fermion generation in the Standard Model. Here $i = 1, 2, 3$ is the generation index, Y denotes the weak hypercharge, and the electric charge satisfies $Q = T_3 + Y/2$.

Field	$SU(3)_c$	$SU(2)_L$	$Y/2$	Q
$Q_{L_i} = \begin{pmatrix} u_i \\ d'_i \end{pmatrix}_L$	3	2	1/6	$(\frac{2}{3}, -\frac{1}{3})$
u_{R_i}	3	1	2/3	2/3
d_{R_i}	3	1	-1/3	-1/3
$L_{L_i} = \begin{pmatrix} \nu_i \\ \ell_i \end{pmatrix}_L$	1	2	-1/2	$(0, -1)$
ℓ_{R_i}	1	1	-1	-1

The kinetic part of the fermion Lagrangian is given by

$$\mathcal{L}_{\text{Fermions}} = \sum_{i=1}^3 \left(\bar{Q}_{L_i} i\gamma^\mu D_\mu Q_{L_i} + \bar{u}_{R_i} i\gamma^\mu D_\mu u_{R_i} + \bar{d}_{R_i} i\gamma^\mu D_\mu d_{R_i} + \bar{L}_{L_i} i\gamma^\mu D_\mu L_{L_i} + \bar{\ell}_{R_i} i\gamma^\mu D_\mu \ell_{R_i} \right), \quad (2.13)$$

where the covariant derivative contains the gauge interactions with the $SU(3)_c$, $SU(2)_L$, and $U(1)_Y$ gauge bosons,

$$D_\mu = \partial_\mu + ig_s T^a G_\mu^a + ig \tau^i W_\mu^i + ig' \frac{Y}{2} B_\mu. \quad (2.14)$$

Where g_s, g, g' are gauge coupling constants and the precise form of D_μ depends on the representation of the field on which it acts.

A fundamental feature of the fermion sector is its chiral structure. Only left-handed fermions transform as doublets under $SU(2)_L$, while right-handed fermions are singlets. As a consequence, the charged weak interaction couples only to left-handed fermions and right-handed antifermions, leading to the $V - A$ structure of the weak current.

The weak-interaction eigenstates of the quarks are not, in general, identical to their mass eigenstates. This misalignment is described by the Cabibbo–Kobayashi–Maskawa (CKM) matrix, defined through

$$d'_{i_L} = V_{ij} d_{j_L}, \quad (2.15)$$

where repeated flavour indices are implicitly summed over if up-quarks are in the mass basis. The CKM matrix appears in charged-current interactions and is responsible for flavour mixing in the quark sector.

The chiral gauge structure of the Standard Model therefore fixes the form of the fermionic kinetic terms and gauge interactions, while the generation of fermion masses is deferred to the Yukawa sector discussed below.

2.1.4 The Higgs sector :

In the SM the Higgs doublet is defined as set of complex scalar field put in a $SU(2)$ doublet

$$\phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix} \quad (2.16)$$

whose self-interactions are responsible for generating the mass terms of both gauge and fermion fields through spontaneous symmetry breaking (SSB). This $SU(2)_L$ doublet, gives rise to the only physical scalar in the SM spectrum, namely the Higgs boson. The electroweak quantum numbers of its components ϕ^+ and ϕ^0 are listed in Table 2.2.

	T_3	$Y/2$	Q
ϕ^+	$+1/2$	$1/2$	1
ϕ^0	$-1/2$	$1/2$	0

Table 2.2: Electroweak quantum numbers of the Higgs doublet components. Where T_3 its third component of the weak isospin, Y the hypercharge, and Q the electric charge.

The Higgs sector of the SM Lagrangian reads

$$\mathcal{L}_{\text{Higgs}} = (D^\mu \phi)^\dagger D_\mu \phi - V(\phi), \quad V(\phi) = \mu^2 |\phi|^2 + \lambda |\phi|^4, \quad (2.17)$$

where $V(\phi)$ is the renormalisable Higgs potential illustrated in Fig 2.2, and the covariant derivative acting on the doublet is

$$D_\mu \phi = \left(\partial_\mu + ig\tau^i W_\mu^i + i\frac{g'}{2} B_\mu \right) \phi. \quad (2.18)$$

For $\mu^2 < 0$, the neutral component ϕ^0 acquires a non-vanishing vacuum expectation

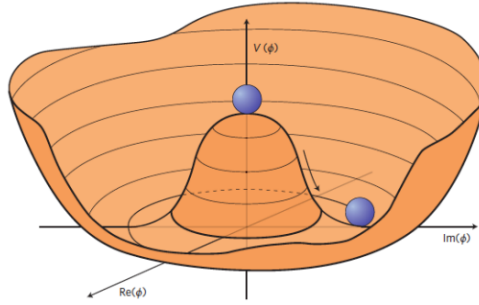


Figure 2.2: The Higgs potential [10]

value (vev) v , defined by

$$v = \sqrt{-\frac{\mu^2}{\lambda}}, \quad \langle \phi \rangle_0 \equiv \langle 0 | \phi | 0 \rangle = \begin{pmatrix} 0 \\ v/\sqrt{2} \end{pmatrix}. \quad (2.19)$$

To isolate the vev, it is convenient to parameterise the Higgs doublet in this way

$$\phi = \exp(i\tau^i \theta^i(x)) \begin{pmatrix} 0 \\ (v + H(x))/\sqrt{2} \end{pmatrix}, \quad (2.20)$$

where $\theta^i(x)$ and $H(x)$ are real scalar fields with vanishing vev. The three Goldstone fields $\theta^i(x)$ can be removed by performing an $SU(2)_L$ gauge transformation, bringing ϕ into the unitary gauge:

$$\phi(x) \longrightarrow \begin{pmatrix} 0 \\ (v + H(x))/\sqrt{2} \end{pmatrix}. \quad (2.21)$$

Applying the covariant derivative to ϕ in the unitary gauge yields

$$D_\mu \phi = \frac{1}{\sqrt{2}} \begin{pmatrix} igW_\mu^+(v + H(x))/\sqrt{2} \\ \partial_\mu H(x) - ig_Z Z_\mu(v + H(x))/2 \end{pmatrix}, \quad (2.22)$$

so that the kinetic term $(D^\mu \phi)^\dagger D_\mu \phi$ expands to

$$\begin{aligned} (D^\mu \phi)^\dagger D_\mu \phi = & \frac{1}{2} \partial^\mu H \partial_\mu H + m_W^2 W_\mu^- W^{+\mu} + \frac{1}{2} m_Z^2 Z^\mu Z_\mu \\ & + gm_W W_\mu^- W^{+\mu} H + g_Z m_Z Z^\mu Z_\mu H + \frac{g^2}{4} W_\mu^- W^{+\mu} H^2 + \frac{g_Z^2}{8} Z^\mu Z_\mu H^2, \end{aligned} \quad (2.23)$$

where $g_Z \equiv g/\cos\theta_W$ is the effective coupling of the Z boson, and where the gauge boson masses are given by

$$m_W = \frac{1}{2}gv, \quad m_Z = \frac{1}{2}g_Z v = \frac{m_W}{\cos\theta_W}. \quad (2.24)$$

where θ_W is the weak mixing angle (Weinberg angle), defined by

$$\sin\theta_W \equiv \frac{g'}{\sqrt{g^2 + g'^2}}, \quad (2.25)$$

with g and g' the $SU(2)_L$ and $U(1)_Y$ gauge couplings respectively. Its measured value at the Z pole is $\sin^2\theta_W \approx 0.231$ [11]. Following their discovery at CERN [1], the $W^\pm$ and Z boson masses have been measured with high precision and being given by [11]

$$m_W = 80.3692 \pm 0.0133 \text{ GeV}, \quad m_Z = 91.1880 \pm 0.0020 \text{ GeV} \quad (2.26)$$

We have that the vev can be related to the Fermi constant G_F through

$$\frac{1}{G_F\sqrt{2}} = \frac{g^2}{8m_W^2} \quad \Rightarrow \quad v = \frac{1}{(G_F\sqrt{2})^{1/2}} \simeq 246 \text{ GeV}. \quad (2.27)$$

The Higgs potential, expanded around the vev, contains both the Higgs mass term and its cubic and quartic self-interactions:

$$\begin{aligned} V(\phi) &= \frac{\mu^2}{2}(v + H)^2 + \frac{\lambda}{4}(v + H)^4 \\ &= \frac{1}{2}m_H^2 H^2 + \sqrt{\frac{\lambda}{2}}m_H H^3 + \frac{\lambda}{4}H^4 - \frac{\lambda v^4}{4}, \end{aligned} \quad (2.28)$$

where the physical Higgs mass is

$$m_H = v\sqrt{2\lambda} = \sqrt{2}|\mu|. \quad (2.29)$$

The Higgs boson was observed by the ATLAS and CMS Collaborations at the CERN Large Hadron Collider in 2012 [3, 4], with a measured mass of $m_H \approx 125 \text{ GeV}$.

2.1.5 The Yukawa sector :

The last ingredient of the Standard Model Lagrangian is the Yukawa sector, which describes the interactions between the fermion fields and the Higgs doublet. These terms are essential because explicit fermion mass terms of the form $m\bar{\psi}\psi$ are not gauge invariant under $SU(2)_L \times U(1)_Y$, since the left-handed and right-handed fermions transform differently under the electroweak group. We therefore generate fermion masses in a gauge invariant way through Yukawa interactions with the Higgs field.

The Yukawa Lagrangian of the Standard Model is given by

$$\mathcal{L}_{\text{Yukawa}} = -y_{ij}^d \bar{Q}_{Li} \phi d_{Rj} - y_{ij}^u \bar{Q}_{Li} \tilde{\phi} u_{Rj} - y_{ij}^\ell \bar{L}_{Li} \phi \ell_{Rj} + \text{h.c.}, \quad (2.30)$$

where $i, j = 1, 2, 3$ are flavour indices, y_{ij}^u , y_{ij}^d and y_{ij}^ℓ are the Yukawa coupling matrices in flavour space, and

$$\tilde{\phi} = i\sigma^2 \phi^* \quad (2.31)$$

is the conjugate Higgs doublet. The use of $\tilde{\phi}$ in the up-type quark term ensures that the interaction is invariant under the gauge symmetry.

Once electroweak symmetry breaking takes place, the Higgs doublet develops a vacuum expectation value,

$$\langle \phi \rangle_0 = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v \end{pmatrix}. \quad (2.32)$$

Substituting this expression into the Yukawa Lagrangian yields fermion mass terms of the form

$$\mathcal{L}_{\text{mass}} = -\frac{v}{\sqrt{2}} y_{ij}^d \bar{d}_{Li} d_{Rj} - \frac{v}{\sqrt{2}} y_{ij}^u \bar{u}_{Li} u_{Rj} - \frac{v}{\sqrt{2}} y_{ij}^\ell \bar{\ell}_{Li} \ell_{Rj} + \text{h.c.} \quad (2.33)$$

From this expression, one identifies the fermion mass matrices as

$$M_u = \frac{v}{\sqrt{2}} y^u, \quad M_d = \frac{v}{\sqrt{2}} y^d, \quad M_\ell = \frac{v}{\sqrt{2}} y^\ell. \quad (2.34)$$

These matrices are in general non-diagonal in flavour space. Passing to the mass basis therefore requires independent unitary rotations of the left-handed and right-handed fermion fields. In the quark sector, the mismatch between the rotations of the up-type and down-type left-handed fields gives rise to the Cabibbo–Kobayashi–Maskawa (CKM) matrix, which appears in the charged-current weak interaction and encodes flavour mixing.

An important consequence of the Yukawa mechanism is that the coupling of the physical Higgs boson to a fermion is proportional to the fermion mass. Indeed, after symmetry breaking one obtains interactions of the form

$$\mathcal{L}_{Hff} = -\frac{m_f}{v} H \bar{f} f, \quad (2.35)$$

showing that heavier fermions couple more strongly to the Higgs boson.

2.2 The need for physics beyond the SM :

Despite its extraordinary predictive power, the SM is not a complete theory of Nature. In the spirit of the scientific method, the role of a physicist is not merely to validate existing frameworks but to identify their limitations and seek more fundamental descriptions of

reality. In this regard, several well-established experimental observations and theoretical considerations point unambiguously to the existence of physics beyond the Standard Model (BSM). A non-exhaustive list of such open problems is the following:

- **Neutrino masses:** The SM predicts massless neutrinos and does not accommodate right-handed neutrino fields. However, the observation of neutrino oscillations in solar, atmospheric and reactor neutrinos [12, 13, 14] provides conclusive evidence that neutrinos are massive [15], implying that their flavour and mass eigenstates are distinct and that a non-trivial mixing matrix must exist.
- **Dark matter:** Astrophysical and cosmological observations consistently indicate that the SM accounts for only approximately 20% of the total matter content of the Universe. Evidence from galactic rotation curves [16], anisotropies in the Cosmic Microwave Background (CMB) [17], and the Bullet Cluster [18] all point to the existence of a stable, electrically neutral, and non-luminous form of matter whose presence is inferred solely through its gravitational effects, and for which the SM offers no viable candidate.
- **Gravity:** While the SM successfully describes the strong, weak, and electromagnetic interactions, it does not incorporate gravity. A consistent quantum field-theoretic description of gravitational interactions remains one of the most profound open problems in theoretical physics. The SM is expected to break down at the Planck scale, $M_{\text{Pl}} \sim 10^{19}$ GeV, where gravitational effects become of order unity and a more fundamental unified description is required.

Taken together, these observations strongly suggest that the SM should be regarded as a low-energy effective description of a deeper and more complete theory. Despite the many theoretical frameworks that have been proposed over the years to address these shortcomings, the absence of direct experimental evidence for any specific BSM scenario has left the question of the correct UV completion of the SM wide open.

Chapter 3

Effective Field Theories

The remarkable success of the SM, combined with the absence of unambiguous evidence for new physics (NP) at the Large Hadron Collider (LHC), presents a fundamental challenge in defining future directions for collider-based exploration. Two complementary strategies have historically driven progress in this regard.

Direct searches rest on the principle that sufficiently energetic collisions can produce new massive states on-shell, detectable as resonances in experimental data — as exemplified by the discovery of the Higgs boson. This strategy is inherently bounded by the collider’s centre-of-mass energy, and the absence of any signal after more than a decade of operation at CERN suggests that NP may lie beyond the current energy frontier.

Indirect searches follow a different logic: rather than producing new particles, they seek to uncover new interactions through high-precision measurements confronted with accurate theoretical predictions. This approach is limited not by collider energy, but by experimental and theoretical precision, making it particularly sensitive to NP effects in the tails of differential distributions as presented in Fig.3.1

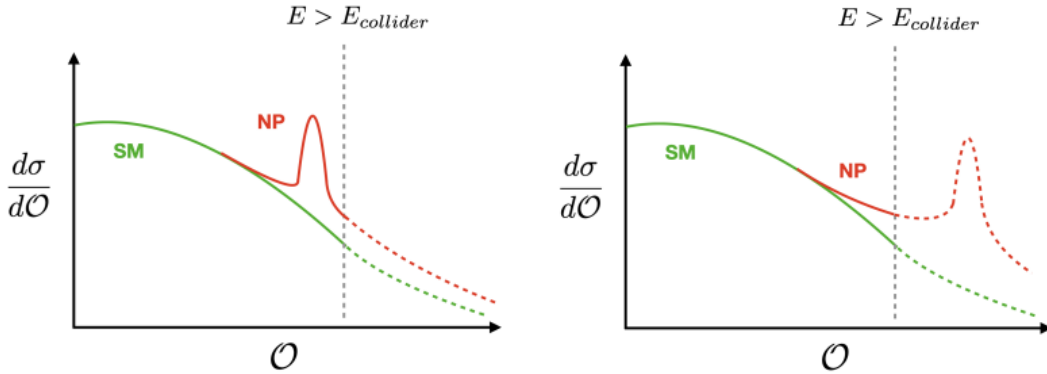


Figure 3.1: The two main strategies for discovering New Physics at colliders are here shown. The observable $\mathcal{O}$ is for example the invariant mass of a final state. On the left, direct search: the resonance is within the reach of the machine and can be directly produced. On the right: the new resonance is beyond the energy frontier and the only way to infer its existence is by scouting the tail of the distribution and find a deviation from the SM [19].

As the field enters an era of unprecedented precision, driven by the High-Luminosity LHC [20] and prospective future lepton colliders [21], this second strategy becomes es-

pecially compelling. In this context, Effective Field Theories (EFTs) provide a natural and model-independent framework for parametrising the low-energy imprints of heavy unknown interactions. Applying this construction to the SM yields the Standard Model Effective Field Theory (SMEFT) [22, 23], which serves as the primary theoretical framework of this thesis.

We will then use this chapter to explain this framework.

3.1 The strategy to build an EFT :

The conceptual foundation of EFTs rests on the *separation of scales*: physical phenomena occurring at a given energy scale can be described self-consistently without requiring detailed knowledge of the dynamics at higher scales. A simple illustration is the fact that Newtonian mechanics governs macroscopic motion without any reference to the underlying quark structure of matter — two problems separated by many orders of magnitude in energy, each admitting its own appropriate theoretical description. More generally, before constructing an effective theory, three ingredients must be identified: the relevant energy scale, the appropriate degrees of freedom (d.o.f) at that scale, and the symmetries that constrain the dynamics. In the context of quantum field theory, these ingredients entirely determine the structure of the Lagrangian, whose operator content is fixed by symmetry requirements and organised systematically through a power-counting expansion.

Two complementary approaches to EFT construction are commonly distinguished:

- **Top-Down:** when the ultraviolet (UV) completion is known, heavy degrees of freedom can be systematically integrated out, yielding a low-energy effective description in which the Wilson coefficients are calculable from the full theory. An example is the Weak Effective Theory (WET) [24].
- **Bottom-Up:** when the UV theory is unknown, an effective description is constructed by augmenting the known low-energy theory with higher-dimensional operators consistent with the observed symmetries. These operators systematically parametrise our ignorance of the underlying dynamics, and the associated Wilson coefficients remain a priori undetermined, to be constrained by experimental data. The SMEFT is the prime example of this approach.

We have in Fig.3.2 a visual illustration of both approaches, i.e top-down and bottom-up.

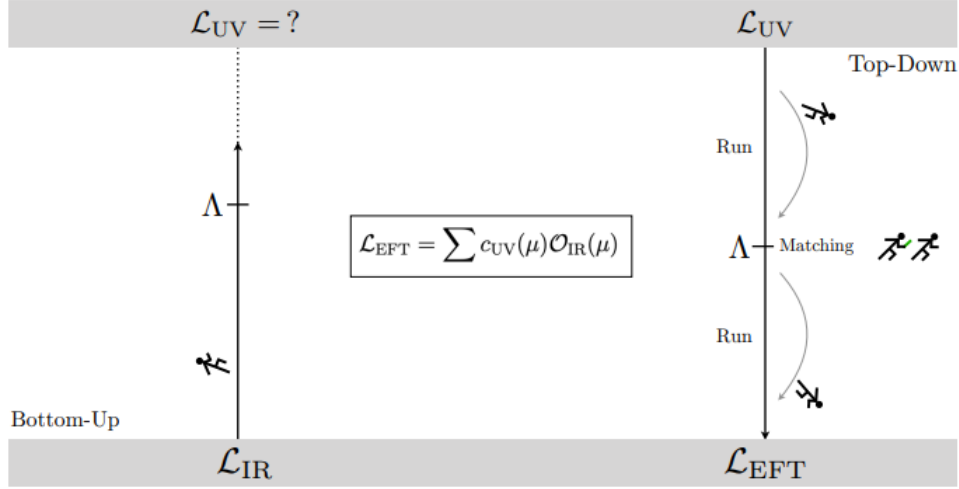
3.2 Power counting :

A systematic way to organize the operator content of an EFT is through power counting, which amounts to tracking the mass dimensions of all terms appearing in the Lagrangian. In natural units, the action is dimensionless and, in four dimensions, takes the form

$$S = \int d^4x \mathcal{L}(x) \quad \Rightarrow \quad [\mathcal{L}(x)] = 4$$

so that, writing the Lagrangian as a sum over local operators,

$$\mathcal{L}(x) = \sum_i c_i \mathcal{O}_i(x),$$

Figure 3.2: Pictorial representation of the different approaches to an EFT [25]

an operator $\mathcal{O}_i$ of mass dimension D must be accompanied by a Wilson coefficient of dimension $4 - D$. The mass dimensions of the fundamental fields are fixed by their kinetic terms: the scalar field obeys the Klein-Gordon equation (two derivatives), the fermion field the Dirac equation (one derivative), and the gauge field kinetic term is built from two contracted field strengths (each carrying one derivative), with $[\partial_\mu] = 1$. Furthermore, noticing that the covariant derivative $D_\mu = \partial_\mu + igA_\mu$ has dimension one fixes the dimension of the gauge coupling $[g]$. Collecting these results,

$$[\psi] = \frac{3}{2}, \quad [\phi] = 1, \quad [A_\mu] = 1, \quad [g] = [y] = 0, \quad [\lambda] = 0,$$

where y and λ denote the Yukawa coupling appearing in $y\phi(\bar{\psi}\psi)$ and the quartic scalar coupling in $\lambda\phi^4$, respectively. Thus, the key building blocks for SMEFT operator construction — bispinors, gauge field strengths $X_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$, covariant derivatives, and scalar fields — carry the following dimensions:

$$[\bar{\psi}\psi] = 3, \quad [X] = 2, \quad [D_\mu] = 1, \quad [\phi] = 1. \quad (3.1)$$

For a dimension-six Lagrangian, one writes

$$\mathcal{L}_{dim-6} = \sum_i \frac{c_i}{\Lambda^2} \mathcal{O}_i, \quad [c_i] = 0, \quad [\Lambda] = 1, \quad [\mathcal{O}_i] = 6,$$

where each operator is expressed as a product of the building blocks,

$$\mathcal{O}_i \sim (\bar{\psi}\psi)^{N_F} (X)^{N_X} (D)^{N_D} (\phi)^{N_\phi}, \quad \Rightarrow \quad 3N_F + 2N_X + N_D + N_\phi = 6. \quad (3.2)$$

This power-counting relation systematically enumerates all operators contributing at dimension six. However, the resulting set is overcomplete : many operators are related to one another through the classical Equations of Motion (EOM) and integration by parts, and are therefore not independent. Identifying a minimal, non-redundant set of operators that spans the space of physical amplitudes is essential for a well-defined EFT parametrisation. In the SMEFT literature, the standard choice is the Warsaw basis [26], which provides such a complete and non-redundant operator basis at dimension six.

3.3 A first example : The Fermi theory

To get familiarised with the idea of an EFT we will explain here the Fermi theory of the weak interaction. It was first developed to explain the beta decay [27] :

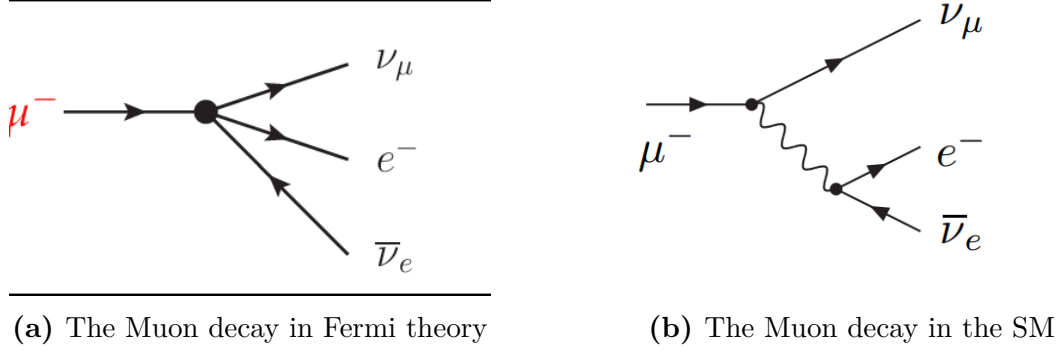


Figure 3.3: Feynman diagrams of the Muon decay in each theory

In the decay process $\mu \rightarrow e \bar{\nu}_e \nu_\mu$, the muon undergoes a transition mediated by an intermediate W gauge boson in the full electroweak theory (see Fig. 4.1b). Since the momentum transfer satisfies $p^2 \ll m_W^2$, the appropriate expansion parameter in the low-energy effective field theory (EFT) is

$$\delta = \frac{p}{m_W}.$$

At energies well below the electroweak scale, the process is accurately described by the four-fermion operator

$$\mathcal{L}_{\text{EFT}} \supset -\frac{4G_F}{\sqrt{2}} (\bar{\nu}_\mu \gamma^\alpha P_L \mu) (\bar{e} \gamma_\alpha P_L \nu_e),$$

where $P_L = \frac{(1-\gamma_5)}{2}$ denotes the left-handed chiral projector and G_F is the Fermi constant.

In the ultraviolet (UV) complete theory, the corresponding four-point amplitude arises from tree-level W exchange and is given by

$$i\mathcal{M}^{\text{UV}} = \bar{u}(p_\nu) \left(\frac{ig}{\sqrt{2}} \gamma^\alpha P_L \right) u(p_\mu) \cdot \frac{ig_{\alpha\beta}}{q^2 - m_W^2} \cdot \bar{u}(k_e) \left(\frac{ig}{\sqrt{2}} \gamma^\beta P_L \right) v(k_\nu).$$

Evaluating the product of couplings yields

$$i\mathcal{M}^{\text{UV}} = \frac{ig^2}{2(q^2 - m_W^2)} [\bar{u}(p_\nu) \gamma^\alpha P_L u(p_\mu)] [\bar{u}(k_e) \gamma_\alpha P_L v(k_\nu)].$$

Expanding the propagator in powers of q^2/m_W^2 and keeping only the leading term gives

$$i\mathcal{M}^{\text{UV}} = -\frac{ig^2}{2m_W^2} [\bar{u}(p_\nu) \gamma^\alpha P_L u(p_\mu)] [\bar{u}(k_e) \gamma_\alpha P_L v(k_\nu)] + \mathcal{O}\left(\frac{q^2}{m_W^2}\right).$$

In the EFT description, the corresponding amplitude is

$$i\mathcal{M}^{\text{EFT}} = -\frac{4iG_F}{\sqrt{2}} [\bar{u}(p_\nu) \gamma^\alpha P_L u(p_\mu)] [\bar{u}(k_e) \gamma_\alpha P_L v(k_\nu)].$$

Matching the UV amplitude to the EFT amplitude at leading order yields

$$i\mathcal{M}^{\text{UV}} = i\mathcal{M}^{\text{EFT}} \quad \Rightarrow \quad G_F = \frac{g^2}{4\sqrt{2}m_W^2}.$$

This is the standard relation expressing the Fermi constant in terms of the electroweak parameters. The dependence on m_W follows directly from dimensional analysis: since the Lagrangian carries mass dimension 4 and each spinor field has mass dimension 3/2, the Fermi constant must have $[G_F] = -2$.

3.4 The Standard Model Effective Field Theory :

In this thesis, possible effects of physics beyond the Standard Model are described within the framework of the Standard Model Effective Field Theory (SMEFT). In this approach, the Standard Model Lagrangian is extended by higher-dimensional operators constructed only from Standard Model fields and required to satisfy Lorentz invariance together with the Standard Model gauge symmetry

$$SU(3)_c \otimes SU(2)_L \otimes U(1)_Y. \quad (3.3)$$

The SMEFT Lagrangian can therefore be written as

$$\mathcal{L}_{\text{SMEFT}} = \mathcal{L}_{\text{SM}} + \sum_{d>4} \sum_i \frac{C_i^{(d)}}{\Lambda^{d-4}} \mathcal{O}_i^{(d)}, \quad (3.4)$$

where Λ denotes the characteristic scale of new physics, $\mathcal{O}_i^{(d)}$ are higher-dimensional operators of canonical dimension d , and $C_i^{(d)}$ are the corresponding Wilson coefficients.

In this work, we restrict ourselves to dimension-six operators only. Under this assumption, the effective Lagrangian reduces to

$$\mathcal{L}_{\text{SMEFT}} = \mathcal{L}_{\text{SM}} + \sum_i \frac{C_i^{(6)}}{\Lambda^2} \mathcal{O}_i^{(6)}. \quad (3.5)$$

The dimension-six operators are organized in the Warsaw basis [26], which provides a complete and non-redundant set of operators after eliminating redundant structures through integration by parts, field redefinitions, and the equations of motion.

The full set of dimension-six operators can be divided into operators without four fermions and four-fermion operators. For convenience, these two classes are listed separately in Tables 3.1 and 3.2. In the presence of three fermion generations, the number of independent flavor configurations becomes very large, especially in the four-fermion sector.

X^3		φ^6 and $\varphi^4 D^2$		$\psi^2 \varphi^3$	
Q_G	$f^{ABC} G_\mu^{A\nu} G_\nu^{B\rho} G_\rho^{C\mu}$	Q_φ	$(\varphi^\dagger \varphi)^3$	$Q_{e\varphi}$	$(\varphi^\dagger \varphi)(\bar{l}_p e_r \varphi)$
$Q_{\tilde{G}}$	$f^{ABC} \tilde{G}_\mu^{A\nu} G_\nu^{B\rho} G_\rho^{C\mu}$	$Q_{\varphi\Box}$	$(\varphi^\dagger \varphi)\Box(\varphi^\dagger \varphi)$	$Q_{u\varphi}$	$(\varphi^\dagger \varphi)(\bar{q}_p u_r \tilde{\varphi})$
Q_W	$\varepsilon^{IJK} W_\mu^{I\nu} W_\nu^{J\rho} W_\rho^{K\mu}$	$Q_{\varphi D}$	$(\varphi^\dagger D^\mu \varphi)^* (\varphi^\dagger D_\mu \varphi)$	$Q_{d\varphi}$	$(\varphi^\dagger \varphi)(\bar{q}_p d_r \varphi)$
$Q_{\tilde{W}}$	$\varepsilon^{IJK} \tilde{W}_\mu^{I\nu} W_\nu^{J\rho} W_\rho^{K\mu}$				
$X^2 \varphi^2$		$\psi^2 X \varphi$		$\psi^2 \varphi^2 D$	
$Q_{\varphi G}$	$\varphi^\dagger \varphi G_{\mu\nu}^A G^{A\mu\nu}$	Q_{eW}	$(\bar{l}_p \sigma^{\mu\nu} e_r) \tau^I \varphi W_{\mu\nu}^I$	$Q_{\varphi l}^{(1)}$	$(\varphi^\dagger i \overleftrightarrow{D}_\mu \varphi)(\bar{l}_p \gamma^\mu l_r)$
$Q_{\varphi \tilde{G}}$	$\varphi^\dagger \varphi \tilde{G}_{\mu\nu}^A G^{A\mu\nu}$	Q_{eB}	$(\bar{l}_p \sigma^{\mu\nu} e_r) \varphi B_{\mu\nu}$	$Q_{\varphi l}^{(3)}$	$(\varphi^\dagger i \overleftrightarrow{D}_\mu^I \varphi)(\bar{l}_p \tau^I \gamma^\mu l_r)$
$Q_{\varphi W}$	$\varphi^\dagger \varphi W_{\mu\nu}^I W^{I\mu\nu}$	Q_{uG}	$(\bar{q}_p \sigma^{\mu\nu} T^A u_r) \tilde{\varphi} G_{\mu\nu}^A$	$Q_{\varphi e}$	$(\varphi^\dagger i \overleftrightarrow{D}_\mu \varphi)(\bar{e}_p \gamma^\mu e_r)$
$Q_{\varphi \tilde{W}}$	$\varphi^\dagger \varphi \tilde{W}_{\mu\nu}^I W^{I\mu\nu}$	Q_{uW}	$(\bar{q}_p \sigma^{\mu\nu} u_r) \tau^I \tilde{\varphi} W_{\mu\nu}^I$	$Q_{\varphi q}^{(1)}$	$(\varphi^\dagger i \overleftrightarrow{D}_\mu \varphi)(\bar{q}_p \gamma^\mu q_r)$
$Q_{\varphi B}$	$\varphi^\dagger \varphi B_{\mu\nu} B^{\mu\nu}$	Q_{uB}	$(\bar{q}_p \sigma^{\mu\nu} u_r) \tilde{\varphi} B_{\mu\nu}$	$Q_{\varphi q}^{(3)}$	$(\varphi^\dagger i \overleftrightarrow{D}_\mu^I \varphi)(\bar{q}_p \tau^I \gamma^\mu q_r)$
$Q_{\varphi \tilde{B}}$	$\varphi^\dagger \varphi \tilde{B}_{\mu\nu} B^{\mu\nu}$	Q_{dG}	$(\bar{q}_p \sigma^{\mu\nu} T^A d_r) \varphi G_{\mu\nu}^A$	$Q_{\varphi u}$	$(\varphi^\dagger i \overleftrightarrow{D}_\mu \varphi)(\bar{u}_p \gamma^\mu u_r)$
$Q_{\varphi WB}$	$\varphi^\dagger \tau^I \varphi W_{\mu\nu}^I B^{\mu\nu}$	Q_{dW}	$(\bar{q}_p \sigma^{\mu\nu} d_r) \tau^I \varphi W_{\mu\nu}^I$	$Q_{\varphi d}$	$(\varphi^\dagger i \overleftrightarrow{D}_\mu \varphi)(\bar{d}_p \gamma^\mu d_r)$
$Q_{\varphi \tilde{W}B}$	$\varphi^\dagger \tau^I \varphi \tilde{W}_{\mu\nu}^I B^{\mu\nu}$	Q_{dB}	$(\bar{q}_p \sigma^{\mu\nu} d_r) \varphi B_{\mu\nu}$	$Q_{\varphi ud}$	$i(\tilde{\varphi}^\dagger D_\mu \varphi)(\bar{u}_p \gamma^\mu d_r)$

Table 3.1: Dimension-six operators other than the four-fermion ones.

$(\bar{L}L)(\bar{L}L)$		$(\bar{R}R)(\bar{R}R)$		$(\bar{L}L)(\bar{R}R)$	
Q_{ll}	$(\bar{l}_p \gamma_\mu l_r)(\bar{l}_s \gamma^\mu l_t)$	Q_{ee}	$(\bar{e}_p \gamma_\mu e_r)(\bar{e}_s \gamma^\mu e_t)$	Q_{le}	$(\bar{l}_p \gamma_\mu l_r)(\bar{e}_s \gamma^\mu e_t)$
$Q_{qq}^{(1)}$	$(\bar{q}_p \gamma_\mu q_r)(\bar{q}_s \gamma^\mu q_t)$	Q_{uu}	$(\bar{u}_p \gamma_\mu u_r)(\bar{u}_s \gamma^\mu u_t)$	Q_{lu}	$(\bar{l}_p \gamma_\mu l_r)(\bar{u}_s \gamma^\mu u_t)$
$Q_{qq}^{(3)}$	$(\bar{q}_p \gamma_\mu \tau^I q_r)(\bar{q}_s \gamma^\mu \tau^I q_t)$	Q_{dd}	$(\bar{d}_p \gamma_\mu d_r)(\bar{d}_s \gamma^\mu d_t)$	Q_{ld}	$(\bar{l}_p \gamma_\mu l_r)(\bar{d}_s \gamma^\mu d_t)$
$Q_{lq}^{(1)}$	$(\bar{l}_p \gamma_\mu l_r)(\bar{q}_s \gamma^\mu q_t)$	Q_{eu}	$(\bar{e}_p \gamma_\mu e_r)(\bar{u}_s \gamma^\mu u_t)$	Q_{qe}	$(\bar{q}_p \gamma_\mu q_r)(\bar{e}_s \gamma^\mu e_t)$
$Q_{lq}^{(3)}$	$(\bar{l}_p \gamma_\mu \tau^I l_r)(\bar{q}_s \gamma^\mu \tau^I q_t)$	Q_{ed}	$(\bar{e}_p \gamma_\mu e_r)(\bar{d}_s \gamma^\mu d_t)$	$Q_{qu}^{(1)}$	$(\bar{q}_p \gamma_\mu q_r)(\bar{u}_s \gamma^\mu u_t)$
		$Q_{ud}^{(1)}$	$(\bar{u}_p \gamma_\mu u_r)(\bar{d}_s \gamma^\mu d_t)$	$Q_{qu}^{(8)}$	$(\bar{q}_p \gamma_\mu T^A q_r)(\bar{u}_s \gamma^\mu T^A u_t)$
		$Q_{ud}^{(8)}$	$(\bar{u}_p \gamma_\mu T^A u_r)(\bar{d}_s \gamma^\mu T^A d_t)$	$Q_{qd}^{(1)}$	$(\bar{q}_p \gamma_\mu q_r)(\bar{d}_s \gamma^\mu d_t)$
				$Q_{qd}^{(8)}$	$(\bar{q}_p \gamma_\mu T^A q_r)(\bar{d}_s \gamma^\mu T^A d_t)$
$(\bar{L}R)(\bar{R}L)$ and $(\bar{L}R)(\bar{L}R)$		B -violating			
Q_{ledq}	$(\bar{l}_p^j e_r)(\bar{d}_s^k q_t^j)$	Q_{duq}	$\varepsilon^{\alpha\beta\gamma} \varepsilon_{jk} [(d_p^\alpha)^T C u_r^\beta] [(q_s^{\gamma j})^T C l_t^k]$		
$Q_{quqd}^{(1)}$	$(\bar{q}_p^j u_r) \varepsilon_{jk} (\bar{q}_s^k d_t)$	Q_{qqu}	$\varepsilon^{\alpha\beta\gamma} \varepsilon_{jk} [(q_p^{\alpha j})^T C q_r^{\beta k}] [(u_s^\gamma)^T C e_t]$		
$Q_{quqd}^{(8)}$	$(\bar{q}_p^j T^A u_r) \varepsilon_{jk} (\bar{q}_s^k T^A d_t)$	Q_{qqq}	$\varepsilon^{\alpha\beta\gamma} \varepsilon_{jn} \varepsilon_{km} [(q_p^{\alpha j})^T C q_r^{\beta k}] [(q_s^{\gamma m})^T C l_t^n]$		
$Q_{lequ}^{(1)}$	$(\bar{l}_p^j e_r) \varepsilon_{jk} (\bar{q}_s^k u_t)$	Q_{duu}	$\varepsilon^{\alpha\beta\gamma} [(d_p^\alpha)^T C u_r^\beta] [(u_s^\gamma)^T C e_t]$		
$Q_{lequ}^{(3)}$	$(\bar{l}_p^j \sigma_{\mu\nu} e_r) \varepsilon_{jk} (\bar{q}_s^k \sigma^{\mu\nu} u_t)$				

Table 3.2: Four-fermion operators.

In the present analysis, we do not consider the full dimension-six operator space. Instead, we restrict ourselves to the ten four-quark operators listed in Table 3.3. This choice is motivated by the fact that these operators directly modify quark–quark scattering amplitudes.

Table 3.3: List of 4LQ operators considered in this study. q denotes the left-handed quark doublets of all three generations, while u and d are the right-handed up-type and down-type quark fields, respectively. σ^I , with $I = \{1, 2, 3\}$, are the Pauli matrices, while T^A , with $A = \{1, \dots, 8\}$, are the $\text{SU}(3)_C$ generators. p, r, s, t are flavour indices; spin and colour indices are omitted.

Operator	Coefficient	Definition
$\mathcal{O}_{qq}^{(1)}$	$C_{qq}^{(1)}$	$(\bar{q}_p \gamma_\mu q_r)(\bar{q}_s \gamma^\mu q_t)$
$\mathcal{O}_{qq}^{(3)}$	$C_{qq}^{(3)}$	$(\bar{q}_p \gamma_\mu \sigma^I q_r)(\bar{q}_s \gamma^\mu \sigma^I q_t)$
$\mathcal{O}_{uu}$	C_{uu}	$(\bar{u}_p \gamma_\mu u_r)(\bar{u}_s \gamma^\mu u_t)$
$\mathcal{O}_{dd}$	C_{dd}	$(\bar{d}_p \gamma_\mu d_r)(\bar{d}_s \gamma^\mu d_t)$
$\mathcal{O}_{ud}^{(1)}$	$C_{ud}^{(1)}$	$(\bar{u}_p \gamma_\mu u_r)(\bar{d}_s \gamma^\mu d_t)$
$\mathcal{O}_{ud}^{(8)}$	$C_{ud}^{(8)}$	$(\bar{u}_p \gamma_\mu T^A u_r)(\bar{d}_s \gamma^\mu T^A d_t)$
$\mathcal{O}_{qu}^{(1)}$	$C_{qu}^{(1)}$	$(\bar{q}_p \gamma_\mu q_r)(\bar{u}_s \gamma^\mu u_t)$
$\mathcal{O}_{qu}^{(8)}$	$C_{qu}^{(8)}$	$(\bar{q}_p \gamma_\mu T^A q_r)(\bar{u}_s \gamma^\mu T^A u_t)$
$\mathcal{O}_{qd}^{(1)}$	$C_{qd}^{(1)}$	$(\bar{q}_p \gamma_\mu q_r)(\bar{d}_s \gamma^\mu d_t)$
$\mathcal{O}_{qd}^{(8)}$	$C_{qd}^{(8)}$	$(\bar{q}_p \gamma_\mu T^A q_r)(\bar{d}_s \gamma^\mu T^A d_t)$

The Wilson coefficients associated with these operators are treated as free parameters and encode possible deviations from the Standard Model induced by heavy new physics.

3.5 The need for observables and differential shapes

As already explained in the SMEFT, the effects of heavy new physics are parametrized through higher-dimensional operators added to the Standard Model Lagrangian, such that collider observables admit an expansion in powers of Λ^{-1} . For a generic observable X , one may write

$$\frac{d\sigma}{dX} = \frac{d\sigma_{\text{SM}}}{dX} + \sum_i \frac{C_i}{\Lambda^2} \frac{d\sigma_i^{\text{int}}}{dX} + \mathcal{O}(\Lambda^{-4}), \quad (3.6)$$

where the coefficients C_i encode the strength of the effective interactions. Since Wilson coefficients are not directly observable, EFT constraints at the LHC must be extracted from experimentally measurable quantities such as cross sections, invariant-mass spectra, transverse-momentum distributions, angular observables, and event-shape variables.

In this context, differential observables are particularly valuable because they retain the kinematic information of the event sample, whereas a cross section reduces the full phase space to a single number. This distinction is especially important in our case, since the linear contribution at order Λ^{-2} arises from the interference between the Standard Model and EFT amplitudes, and this interference can vary in sign across phase space which will

be studied later in this thesis.

As a result, sizable effects may partially cancel after integration, making the inclusive rate significantly less sensitive than the corresponding differential distribution. By contrast, suitably chosen shapes can isolate the regions in which the EFT contribution is enhanced, recover sensitivity to the sign and magnitude of the Wilson coefficients, and help disentangle operators that would otherwise remain degenerate in a cross-section only analysis.

This is one of the main reasons why observables, shape distributions, and high-energy tails play a central role in modern global SMEFT fits at the LHC, although the enhanced sensitivity of these regions must always be balanced against the question of EFT validity and the size of higher-order corrections.

3.6 Some definitions : CMS coordinate system and angular variables

We will finally end the introductory part of this thesis by defining some widely used quantities in collider physics, we consider the two incoming particles to move along opposite directions on the z axis, which we call the beam axis. The plane perpendicular to this axis is called the transverse plane. In the CMS coordinate system, the x -axis points toward the centre of the LHC ring, the y -axis points upward, and the azimuthal angle ϕ is measured in the transverse (x, y) plane. We define the transverse momentum $\vec{p}_T$ the projection of $\vec{p}$ on the xy plane, orthogonal to the beam direction.

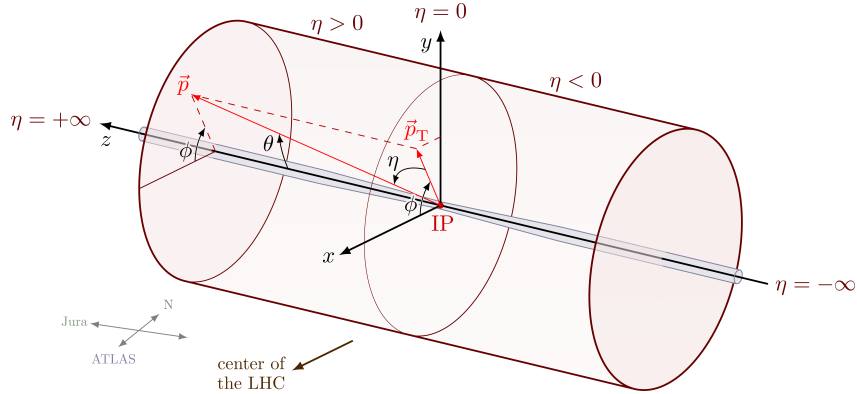


Figure 3.4: Coordinate system for the CMS experiment in the Large Hadron Collider [28]

The direction of a final-state object is described by the polar angle θ with respect to the beam axis. In collider physics, it is more convenient to use the pseudorapidity η , defined as

$$\eta = -\ln \left(\tan \frac{\theta}{2} \right). \quad (3.7)$$

For two jets, one may define the difference in azimuthal angle as

$$\Delta\phi_{jj} = |\phi_{j1} - \phi_{j2}|, \quad (3.8)$$

and the difference in pseudorapidity as

$$\Delta\eta_{jj} = \eta_{j_1} - \eta_{j_2}. \quad (3.9)$$

The angular distance between the two jets in the (η, ϕ) plane is then defined by

$$\Delta R_{jj} = \sqrt{(\Delta\eta_{jj})^2 + (\Delta\phi_{jj})^2}. \quad (3.10)$$

This can then be generalized to any other particle.

Chapter 4

Constraining the coefficients

Dijet production at the LHC provides a powerful probe of four-quark operators; however, it suffers from important limitations. For example, because all quark flavours and chiralities look the same once in jets, these measurements are inherently insensitive to both the flavour and chirality structure of the underlying four-light-quark interactions. As a result, dijet measurements alone are insufficient to disentangle the full set of 4LQ Wilson coefficients.

The associated production of a photon with two jets can address these limitations. On one hand, the presence of the additional photon opens up a broader set of differential observables, such as photon-jet angular distribution and invariant-mass distributions of both the dijet and γ +dijet systems. On the other hand, since up- and down-type quarks carry different electric charges, their couplings to the photon differ, making γ +jets observables sensitive to the flavour structure of the four-light-quark operator space. This provides access to flavour-specific directions in the Wilson-coefficient space that dijet measurements alone cannot probe. We therefore consider the process $pp \rightarrow \gamma jj$ as our channel of investigation. All the 4LQ can take part in this process and we present a feynman diagram of one contribution in Fig4.1

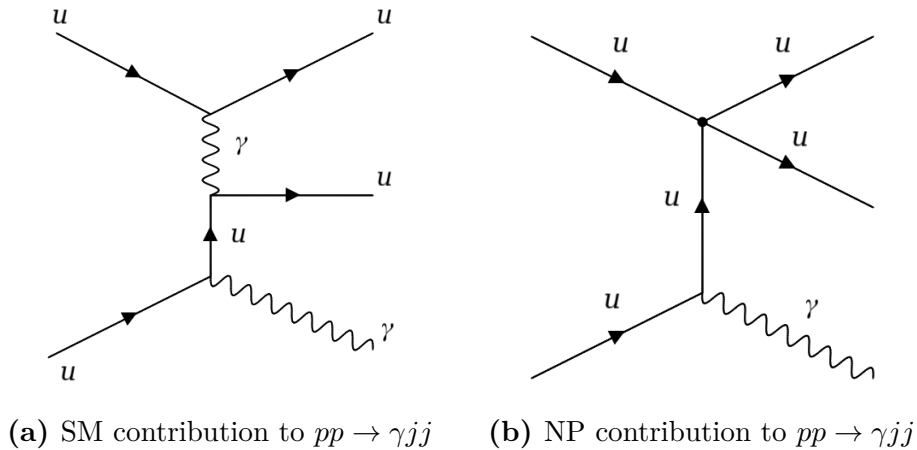


Figure 4.1: Feynman diagrams

4.1 Framework :

All events were generated with MADGRAPH5 aMC@NLO [29] from an UFO [30] model containing the 4LQ operators, which are defined in Tab. 3.3, on top of the SM. It is worth noting that : $\mathcal{O}_{qq}^{(1)}$, $\mathcal{O}_{qq}^{(3)}$, $\mathcal{O}_{uu}$ and $\mathcal{O}_{dd}$ present two fermionic currents with the same chirality, left-handed for the first two and righthanded for the others. For this reason, there are two possible ways to contract their flavour indices : inside the fermion bilinears or between them these two contractions carry respectively a colour singlet and octet structure, so they can be constrained separately. In the UFO model used, the two contractions are summed together and associated to the same Wilson coefficient [31]. The top quark is included as well, but we excluded it when generating the processes. Throughout this work, we restrict ourselves to the leading-order (LO) and partonic level. NNPDF3.0 [32] was used as parton distribution function (PDF) set. The leptons and all the quarks but the top are taken as massless and the CKM matrix is assumed to be diagonal. The total transverse energy of the event divided by 2, $H_T/2$, was chosen as dynamical scale. We have used the following experimental cuts [33], they are summarized in Table 4.1. We presents results for the LHC at 13 TeV.

Table 4.1: Photon and jet selection requirements, and event categorization.

Requirements on photon	$E_T^\gamma > 150 \text{ GeV}, \quad \eta^\gamma < 2.37$
Requirements on jets	at least two jets with $R = 0.4$ $p_T^{\text{jet}} > 100 \text{ GeV}, \quad y^{\text{jet}} < 2.5, \quad \Delta R_{\gamma\text{-jet}} > 0.8$

The total cross section for the SM and the linear terms are reported in Table 4.2

$\gamma + \text{jets}$ (pb)	
SM	$18.98 \pm 0.14\%$
$\mathcal{O}(1/\Lambda^2)$	
$\mathcal{O}_{qq}^{(1)}$	$-0.579 \pm 0.17\%$
$\mathcal{O}_{qq}^{(3)}$	$-2.327 \pm 0.12\%$
$\mathcal{O}_{uu}$	$-0.458 \pm 0.12\%$
$\mathcal{O}_{dd}$	$-0.051 \pm 0.42\%$
$\mathcal{O}_{ud}^{(1)}$	$0.015 \pm 0.54\%$
$\mathcal{O}_{ud}^{(8)}$	$-0.143 \pm 0.46\%$
$\mathcal{O}_{qu}^{(1)}$	$-0.016 \pm 0.45\%$
$\mathcal{O}_{qu}^{(8)}$	$-0.286 \pm 0.45\%$
$\mathcal{O}_{qd}^{(1)}$	$0.003 \pm 0.52\%$
$\mathcal{O}_{qd}^{(8)}$	$-0.115 \pm 0.53\%$

Table 4.2: Cross sections (in fb) for $\gamma + \text{jets}$ production for the SM and the 4LQ operators at $\mathcal{O}(1/\Lambda^2)$ whose coefficients are set to 1 TeV^{-2} .

For the Standard Model prediction, the LO cross section is obtained via MADGRAPH5_AMC@NLO.

4.2 Constraints at the cross-section level :

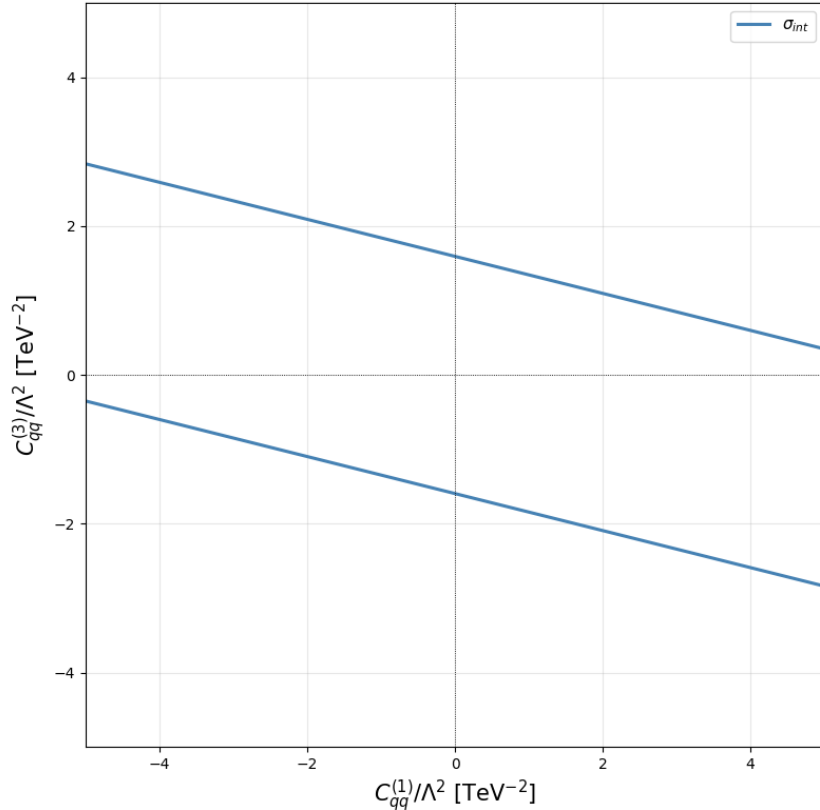
We begin by examining the constraints that can be extracted from the cross section. The allowed region is defined by the condition :

$$\left| \sum_{i=1}^{operators} \frac{c_i}{\Lambda^2} \sigma_i^{1/\Lambda^2} \right| \leq \pm \Delta, \quad (4.1)$$

where σ_i^{1/Λ^2} is the interference-level cross section associated with the i -th operator of the four-light-quark (4LQ) basis and $\frac{c_i}{\Lambda^2}$ are the corresponding Wilson coefficients, and $\Delta = 10\% \times \sigma_{\text{SM}}$ is the maximal EFT-induced deviation relative to the Standard Model prediction that is considered acceptable.

Restricting equation 4.1 to two operators, the resulting allowed band in the $C_{qq}^{(1)}$ vs $C_{qq}^{(3)}$ plane is shown in Figure 4.2. As can be readily observed, the cross-section constraint is highly asymmetric : one direction in coefficient space is strongly bounded, while the orthogonal direction remains unconstrained. This is a direct consequence of the limited sensitivity of the inclusive cross section to the full structure of the operator basis. The primary goal of this thesis is therefore to exploit the additional information encoded in differential distributions in order to lift this degeneracy and obtain tighter bounds on the Wilson coefficients of the 4LQ operators

Figure 4.2: Cross-section constraints in the $C_{qq}^{(1)}$ vs $C_{qq}^{(3)}$ plane for $\gamma + \text{jets}$ production.



4.3 Analysis and differential distribution :

It was shown in [31] that among the different observables considered in their analysis, such as the dijet invariant mass M_{jj} , the γjj invariant mass $M_{\gamma jj}$, the azimuthal separation $\Delta\phi_{jj}$, etc, the photon transverse momentum, p_T^γ , was found to show the largest 4LQ effect relative to the Standard Model prediction. For this reason, we chosed to focus on this observable to obtain bounds.

The differential distribution for p_T^γ is shown in Figure 4.3, where we applied the binning [150, 175, 200, 250, 300, 350, 400, 470, 550, 650, 750, 900, 1100, 1500] GeV with the last one including the overflow. We notice that, for this observable, the 4LQ linear shapes are different from the SM ones, but no remarkable distinctions can be identified among the operators. We also note that the photon transverse momentum presents larger 4LQ effects compared to the SM ones, by at least a factor four in the tail of its distribution. For this reason, we used p_T^γ to extract bounds over the coefficients.

4.4 Finding bounds on the 4LQ Operators :

4.4.1 The χ^2 distribution :

For this analysis we have chosen to assume that the experimental data follows the SM one generated by us, resulting in the individual χ^2 expression for a given observable X

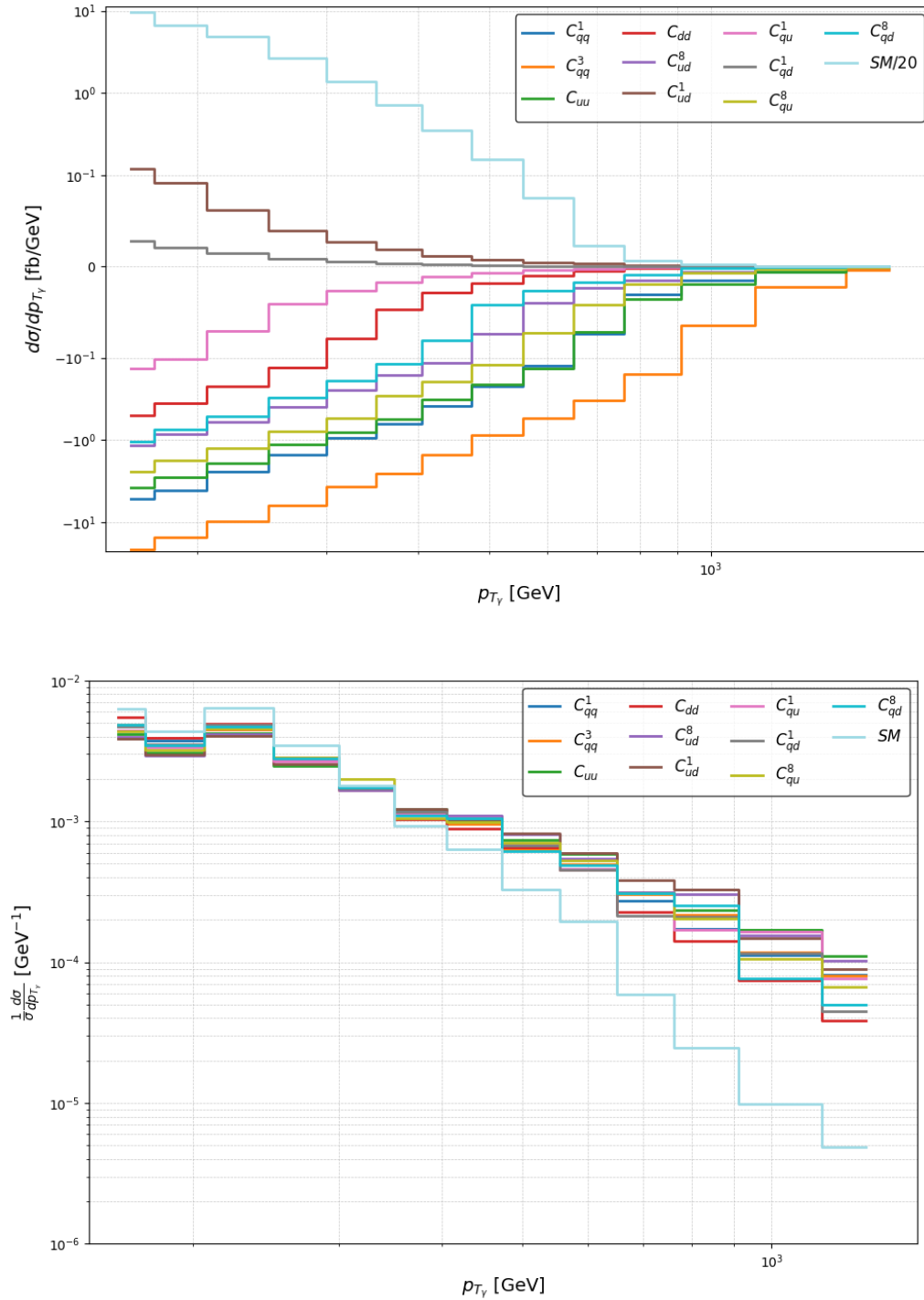
$$\chi^2(\vec{c}) = \sum_{i=1}^{N_{bins}} \frac{1}{\Delta_i^2} \left(\frac{d\sigma_i^{\text{exp}}}{dX} - \frac{d\sigma_i^{\text{SM}}}{dX} - \sum_j^{\text{operators}} \frac{c_j}{\Lambda^2} \frac{d\sigma_{j,i}^{1/\Lambda^2}}{dX} \right)^2 \quad (4.2)$$

$$= \sum_i^{N_{bins}} \frac{1}{\Delta_i^2} \left(\sum_j^{\text{operators}} \frac{c_j}{\Lambda^2} \frac{d\sigma_{j,i}^{1/\Lambda^2}}{dX} \right)^2 \quad (4.3)$$

With here, the c_j are the component of $\vec{c}^T = (C_{qq}^{(1)}, C_{qq}^{(3)}, \dots, C_{qd}^{(8)})$ which is the vector of ten 4LQ operators. The uncertainty at the denominator, Δ_i , has to contain uncertainties from PDFs and scale variations, but also experimental errors, for simplicity, and to facilitate direct comparison with the results obtained at the cross-section level, was assigned to 10% of the SM value in each bin.

This is, however, only an estimation; since we are analysing some simulations, we do not have experimental uncertainties to involve, while they would be taken into account when studying real data. One can see that equation 4.3 is minimized at $\vec{c}^T = \vec{0}$, since the analysis is defined relative to deviations from the Standard Model. In real data, however, the minimum is expected to be slightly nonzero. Therefore, in our simulation, bounds closer to $\vec{c}^T = \vec{0}$ indicate potentially a more effective procedure for constraining $\vec{c}^T$.

Figure 4.3: Differential (top) and normalised (bottom) distributions of the photon transverse momentum in $\gamma + \text{jets}$ production, for the SM (light blue) and the interference of the ten operators included in this analysis at LO, with all the coefficients set to $C_i/\Lambda^2 = 1 \text{ TeV}^{-2}$



4.4.2 The bounds on the coefficients :

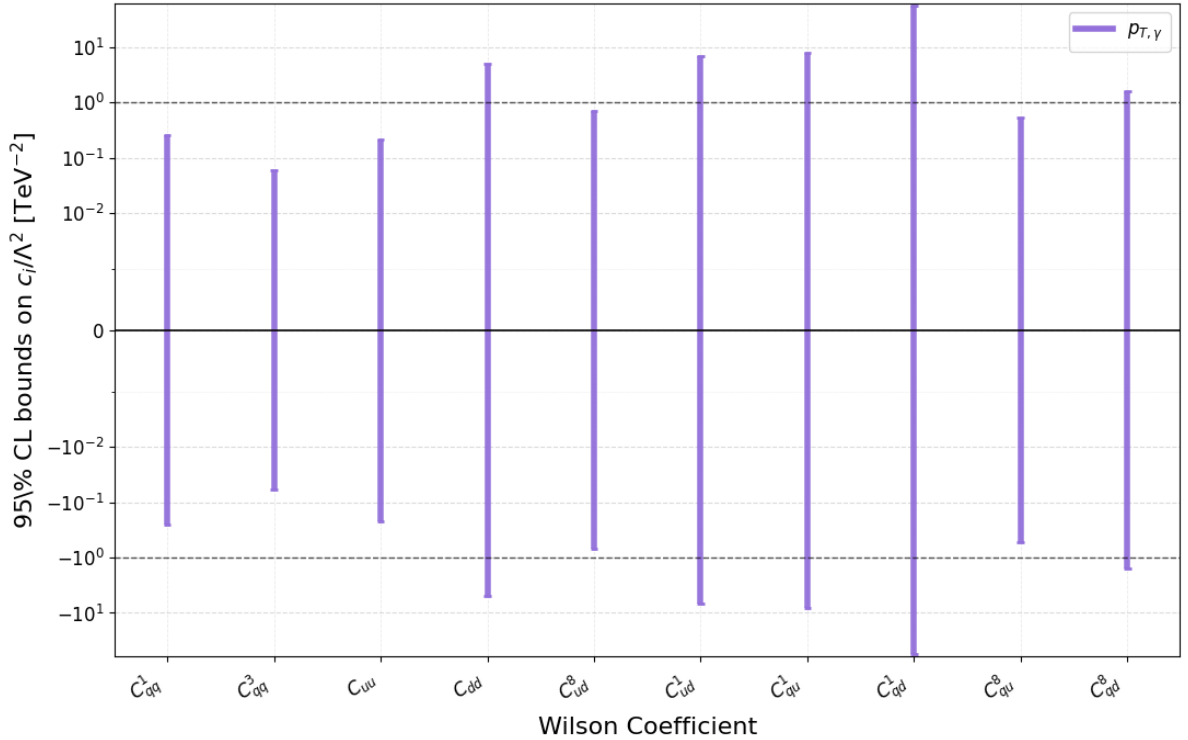
We have obtained limits on the 4LQ coefficients through the χ^2 in equation 4.3 : during this work we will perform individual bounds meaning that all the coefficients but one are set to zero.

The individual bounds obtained from the photon transverse momentum are both shown on Fig. 4.4 and Table 4.3.

Table 4.3: Individual limits, at 95% CL, on the 4LQ operators from $\gamma + \text{jets}$ production with p_T^γ .

Operator	95% CL Interval
$C_{qq}^{(1)}$	$[-0.2523, 0.2523]$
C_{qq}^3	$[-0.0598, 0.0598]$
C_{uu}	$[-0.2175, 0.2175]$
C_{dd}	$[-4.9147, 4.9147]$
C_{ud}^8	$[-0.7011, 0.7011]$
C_{ud}^1	$[-6.9088, 6.9088]$
C_{qu}^1	$[-8.0209, 8.0209]$
C_{qd}^1	$[-55.7959, 55.7959]$
C_{qu}^8	$[-0.5372, 0.5372]$
C_{qd}^8	$[-1.5784, 1.5784]$

Figure 4.4: Individual limits, at 95% CL, on the 4LQ operators from $\gamma + \text{jets}$ production with p_T^γ



We can then see that the stringest bounds are found on the $C_{qq}^{(3)}$ coefficient.

4.5 Global constraints :

In order to constrain a combination of Wilson coefficients simultaneously, we construct a global binned χ^2 function using the linear SMEFT contribution of each operator. For each bin a of the distribution, the prediction is written as

$$\sigma_{\text{th}}^a(\vec{c}) = \sigma_{\text{SM}}^a + \sum_{i=1}^{10} \frac{c_i}{\Lambda^2} \sigma_i^a, \quad (4.4)$$

where σ_{SM}^a denotes the Standard Model prediction in bin a , and σ_i^a is the linear interference contribution associated with operator i .

Assuming again that the experimental distribution is given by the SM prediction generated by us, and assigning an uncertainty of 10% of the SM value in each bin, the χ^2 function can be written as

$$\chi^2(\vec{c}) = \sum_{a=1}^{N_{\text{bins}}} \left[\frac{\sum_{i=1}^{10} c_i \sigma_i^a}{0.1 \sigma_{\text{SM}}^a} \right]^2. \quad (4.5)$$

This expression is quadratic in the Wilson coefficients and can therefore be rewritten in matrix form as

$$\chi^2(\vec{c}) = \sum_{i,j=1}^{10} c_i M_{ij} c_j = \vec{c}^T M \vec{c}, \quad (4.6)$$

with

$$M_{ij} = \sum_{a=1}^{N_{\text{bins}}} \frac{\sigma_i^a \sigma_j^a}{(0.1 \sigma_{\text{SM}}^a)^2}. \quad (4.7)$$

The matrix M summarizes the sensitivity of the observable to the different operators and their correlations. Since it is built from a sum of products, it is symmetric,

$$M_{ij} = M_{ji}, \quad (4.8)$$

and it is also positive semi-definite. Indeed, for any vector $\vec{x}$ in coefficient space, one has

$$\vec{x}^T M \vec{x} = \sum_{a=1}^{N_{\text{bins}}} \left[\frac{\sum_i x_i \sigma_i^a}{0.1 \sigma_{\text{SM}}^a} \right]^2 \geq 0. \quad (4.9)$$

As a consequence, all eigenvalues of M are real and non-negative.

In the present setup, the matrix M entering the quadratic form of the χ^2 coincides with the Fisher matrix F . Indeed, for Gaussian and uncorrelated uncertainties, the Fisher matrix is defined as [34]

$$F_{ij} = \sum_{a=1}^{N_{\text{bins}}} \frac{1}{(\Delta^a)^2} \frac{\partial \sigma_{th}^a}{\partial c_i} \frac{\partial \sigma_{th}^a}{\partial c_j}, \quad (4.10)$$

Since

$$\frac{\partial \sigma_{th}^a}{\partial c_i} = \sigma_i^a, \quad (4.11)$$

and since in our analysis the uncertainty in each bin is taken to be

$$\Delta^a = 0.1 \sigma_{\text{SM}}^a, \quad (4.12)$$

it follows that

$$F_{ij} = \sum_{a=1}^{N_{\text{bins}}} \frac{\sigma_i^a \sigma_j^a}{(0.1 \sigma_{\text{SM}}^a)^2}. \quad (4.13)$$

Therefore,

$$F_{ij} = M_{ij}, \quad (4.14)$$

One can note that this matrix is independent of the Wilson coefficients and therefore of the fit results. Because F is a real symmetric matrix, it can be diagonalized as

$$F = R \text{diag}(\lambda_1, \lambda_2, \dots, \lambda_{10}) R^T, \quad (4.15)$$

where the columns of R are the eigenvectors of F , and the λ_k are the corresponding eigenvalues. The eigenvectors define independent linear combinations of the original Wilson coefficients, that is, the principal directions in the ten-dimensional parameter space.

Under usual assumptions, the Fisher matrix is the inverse of the covariance matrix [34]. One therefore has

$$\text{Cov} = F^{-1} = R \text{diag}\left(\frac{1}{\lambda_1}, \frac{1}{\lambda_2}, \dots, \frac{1}{\lambda_{10}}\right) R^T. \quad (4.16)$$

This shows that, in the basis of Fisher eigenvectors, the variance associated with the k -th principal direction is

$$\text{Var}(d_k) = \frac{1}{\lambda_k}, \quad (4.17)$$

and the corresponding uncertainty scales as

$$\Delta d_k \sim \frac{1}{\sqrt{\lambda_k}}. \quad (4.18)$$

This provides a direct physical interpretation of the eigenvalues. A large eigenvalue corresponds to a small variance, and therefore to a small uncertainty on the associated linear combination of Wilson coefficients. Such a direction is thus strongly constrained by the observable. Conversely, a small eigenvalue corresponds to a large variance and therefore to a weakly constrained direction.

An exactly flat direction corresponds to the case $\lambda_k = 0$. In that situation, the variance along the corresponding eigendirection formally diverges, which means that the associated linear combination of Wilson coefficients is unconstrained by the observable. Equivalently, there exists a non-zero vector $\vec{d}$ such that

$$F\vec{d} = M\vec{d} = 0, \quad (4.19)$$

or, bin by bin,

$$\sum_i d_i \sigma_i^a = 0 \quad \text{for every bin } a. \quad (4.20)$$

Hence, a flat direction requires an exact cancellation of the interference contribution in all bins of the considered distribution.

In general, one expects approximate flat directions rather than exact ones. These correspond to very small but non-zero eigenvalues, and therefore to very large uncertainties

on the associated combinations of Wilson coefficients.

In practice, the eigenvector associated with the largest eigenvalue defines the best-constrained linear combination of Wilson coefficients, since it corresponds to the smallest variance. Conversely, the eigenvector associated with the smallest eigenvalue defines the least-constrained combination, namely the direction to which the observable is least sensitive. We summarize in Table 5.10 the value of the Wilson coefficient corresponding to the best-constrained direction, together with the associated eigenvalue. The eigenvectors are defined with a unitary norm and showing a positive scalar product with $(1, 1, \dots, 1)$.

Table 4.4: Best constrained direction in coefficient space and its associated eigenvalue, for $\gamma + \text{jets}$ production by the photon transverse momentum, p_T^γ .

Operator	Coefficient
C_{qq}^1	+0.22
C_{qq}^3	+0.93
C_{uu}	+0.26
C_{dd}	+0.01
C_{ud}^8	+0.08
C_{ud}^1	-0.01
C_{qu}^1	+0.01
C_{qd}^1	-0.00
C_{qu}^8	+0.10
C_{qd}^8	+0.03
λ	1235.65

We can also show how the bounds at the cross-section level are versus the ones obtained with p_T^γ as illustrated in Fig 4.5. We can see that the most constrained direction is mostly aligned with $C_{qq}^{(3)}$, which can be explained by the fact that $C_{qq}^{(3)}$ is the most constrained coefficient at the individual level. We then show the individual contours associated to some two-dimensional coefficient plane. Examining the two-dimensional contour in the plane containing $C_{qq}^{(3)}$, we observe that the ellipse is tilted with respect to the parameter axes, indicating a correlation between the two Wilson coefficients for the observables considered. This suggests that the observable is predominantly sensitive to a specific linear combination of the two coefficients rather than to each independently.

One can also see that compared to the constraints obtained at the cross-section level, the observable further shrinks the allowed parameter space to a narrower region, indicating improved sensitivity but along essentially the same correlated direction in Wilson-coefficient space. That means p_T^γ does not change the main correlation direction very much; instead, it mainly sharpens the constraint along the same degeneracy already present in the cross-section fit.

Figure 4.5: Individual comparison of the cross-section constraints versus the ones obtained by p_T^γ in the $C_{qq}^{(1)}$ vs $C_{qq}^{(3)}$ plane (*top left*), in the C_{uu} vs $C_{qq}^{(3)}$ one (*top right*), in the C_{dd} vs $C_{qq}^{(3)}$ one (*bottom left*) and in the $C_{qu}^{(8)}$ vs $C_{qq}^{(3)}$ one (*bottom right*) for $\gamma + \text{jets}$ production.

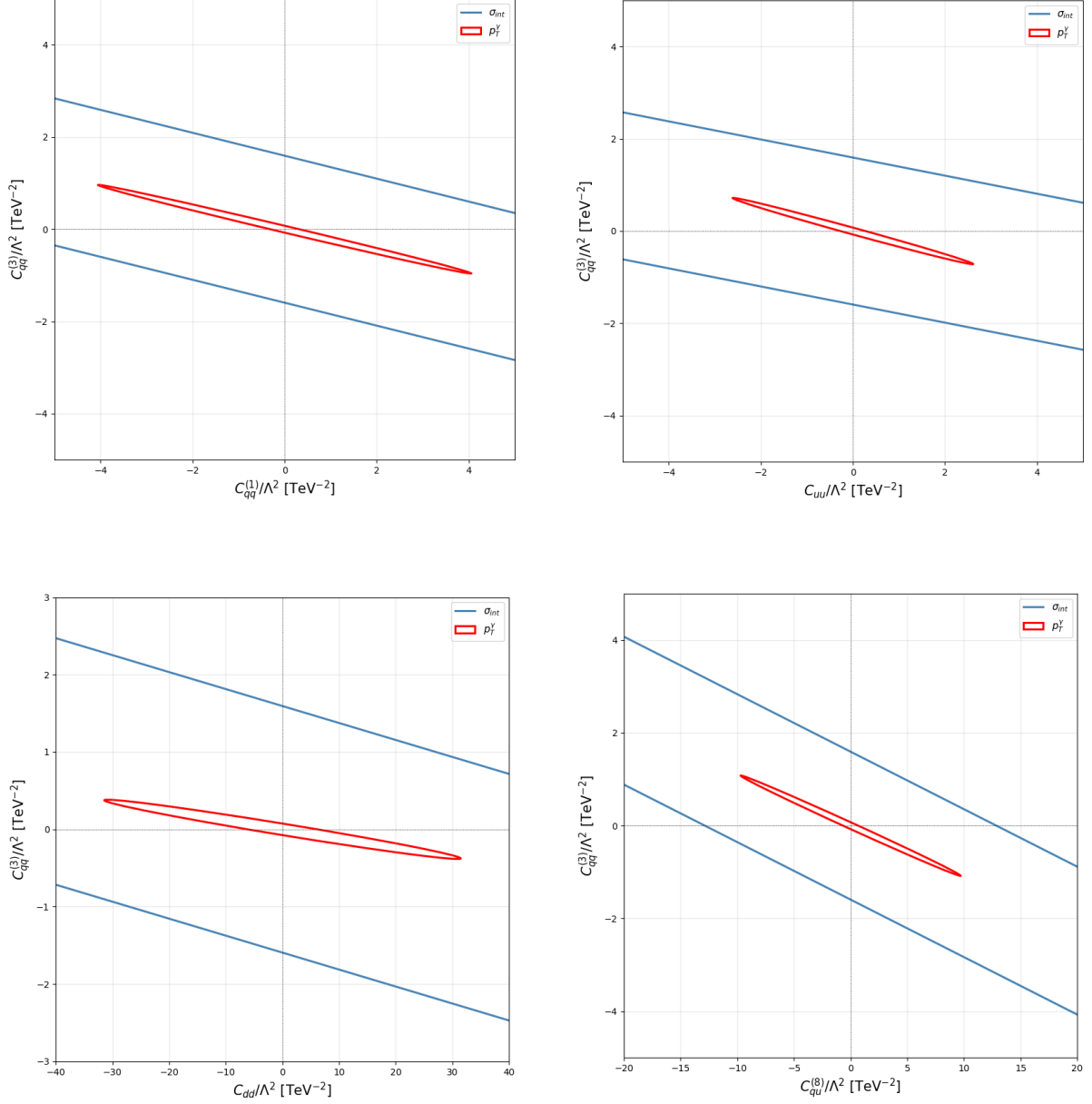
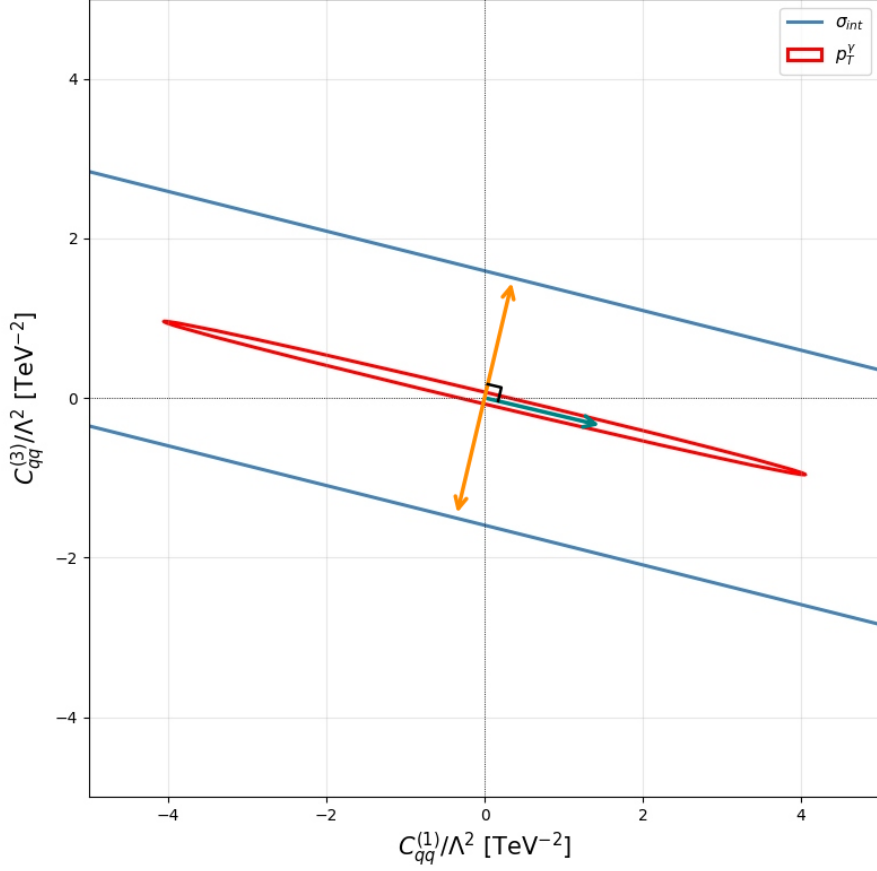


Figure 4.6: Illustration of the directions in the $C_{qq}^{(1)}$ vs $C_{qq}^{(3)}$ plane for $\gamma + \text{jets}$ production.

In Figure 4.6, we display the directions corresponding to the most tightly constrained and to the weakly constrained combinations of Wilson coefficients. The short axis of the ellipse identifies the best-constrained direction, while the long axis corresponds to a weakly constrained direction along which the observables provide essentially no information.

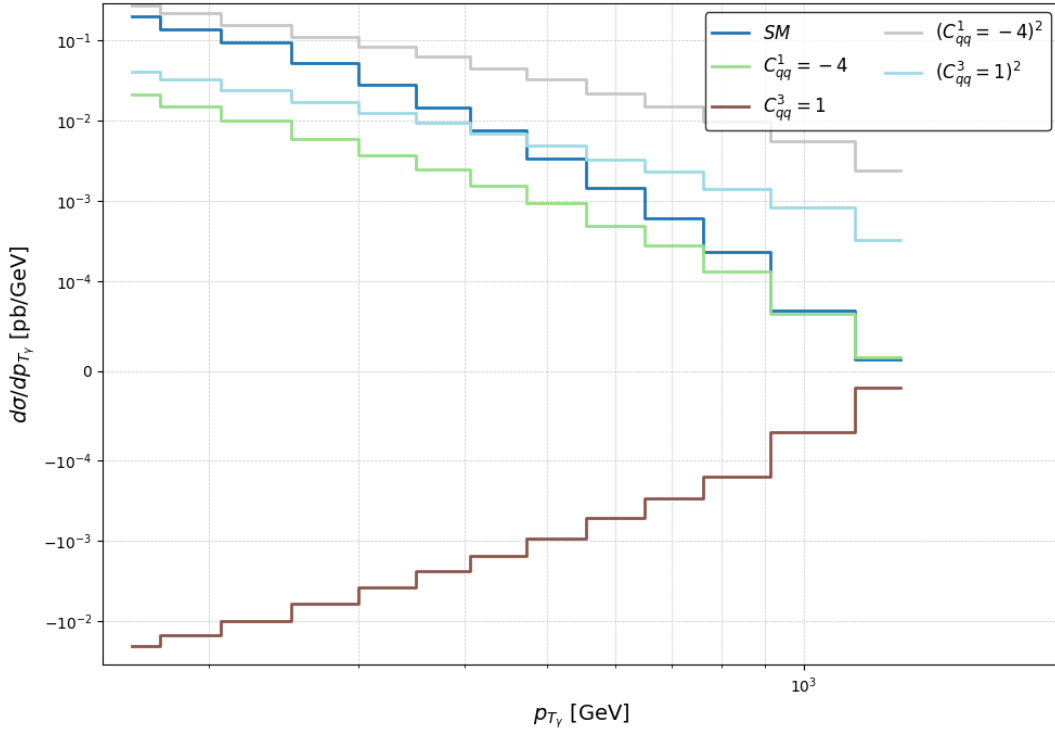
To illustrate the impact of the EFT validity region, we consider a point lying on the edge of the ellipse, namely $(C_{qq}^{(1)}/\Lambda^2, C_{qq}^{(3)}/\Lambda^2) = (-4, 1) \text{ TeV}^{-2}$. The $p_{T,\gamma}$ distribution is shown in Figure 4.7 for the two coefficient choices $(C_{qq}^{(1)}/\Lambda^2, C_{qq}^{(3)}/\Lambda^2) = (-4, 0)$ and $(0, 1) \text{ TeV}^{-2}$, each contributing an interference and a dimension-6 squared curve. The cross-section contributions are reported in Table 4.5.

One can observe that the two interference curves cancel each other, which is expected since this parameter point lies in a direction orthogonal to the cross-section in the $(C_{qq}^{(1)}, C_{qq}^{(3)})$ plane. Furthermore, at this parameter point, the dimension-6 squared contribution exceeds both the SM and interference terms, indicating that the EFT expansion is no longer under control and that this point lies outside the validity region of the effective theory.

Table 4.5: Cross sections for $\gamma + \text{jets}$ production for the SM and the 4LQ operators at $\mathcal{O}(1/\Lambda^4)$ and $\mathcal{O}(1/\Lambda^2)$ whose coefficients are set to -4 and 1 TeV^{-2} for respectively for $C_{qq}^{(1)}$, $C_{qq}^{(3)}$.

$\gamma + \text{jets}$ (pb)	
SM	$18.98 \pm 0.14\%$
$\mathcal{O}(1/\Lambda^2)$	
$\mathcal{O}_{qq}^{(1)}$	$2.32 \pm 0.15\%$
$\mathcal{O}_{qq}^{(3)}$	$-2.327 \pm 0.12\%$
$\mathcal{O}(1/\Lambda^4)$	
$\mathcal{O}_{qq}^{(1)}$	$1.695 \pm 0.21\%$
$\mathcal{O}_{qq}^{(3)}$	$46.64 \pm 0.09\%$

Figure 4.7: Differential distribution of the photon transverse momentum in $\gamma + \text{jets}$ production, for the SM (hard blue), the interference of $\mathcal{O}_{qq}^{(1)}$, $\mathcal{O}_{qq}^{(3)}$ at LO with their $\mathcal{O}(\Lambda^{-4})$ contribution included, with the coefficients set to -4 and 1 TeV^{-2} for respectively for $C_{qq}^{(1)}$, $C_{qq}^{(3)}$



The main objective of the rest of this thesis is to find observables that tighten these bounds and, ideally, render additional directions in coefficient space significantly constrained. To this end, we will investigate directions orthogonal to the interference contribution to the cross section and search for observables that are particularly sensitive to them. Our strategy will exploit the interference revival technique, originally developed to revive otherwise suppressed SMEFT interference effects [35, 36].

Chapter 5

Strategy to find new observables:

Since we want to probe orthogonal directions in the coefficient space compared to the one constrained by the total cross-section, we focus on linear combinations of operators for which the interference contribution to the total cross-section vanishes, *i.e.* the second term in the EFT expansion,

$$\sigma = \sigma_{\text{SM}} + \sum_i \frac{C_i}{\Lambda^2} \sigma_i^{\text{int}} + \mathcal{O}(1/\Lambda^4),$$

The remaining contribution is therefore purely differential with regions that contribute positively and negatively over the full phase space and exactly add up to zero in the total cross-section. First, we look for the normalised coefficients combination that gives the largest contribution in absolute value over the phase space as it should be the one that we can constrain the most. Secondly, we look for differential observables that can separate the positive and negative contributions in the most efficient way as they will give the strongest constraints since the cancellation is minimal [35].

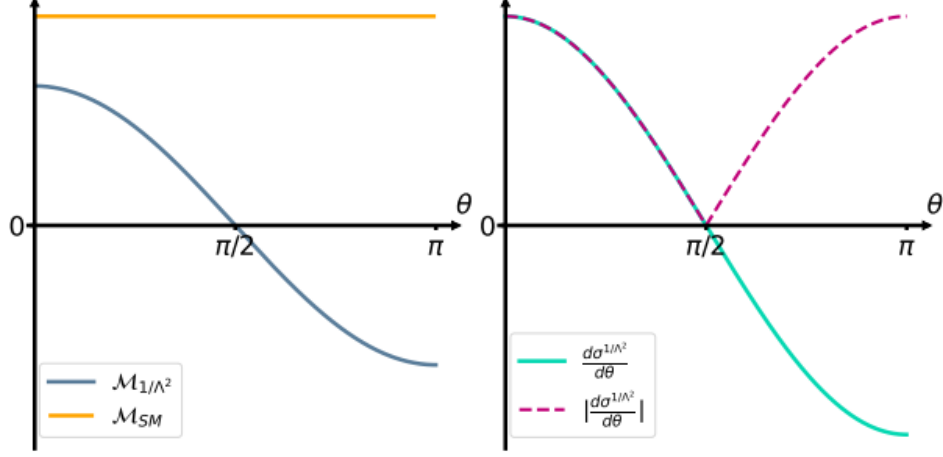
5.1 Large cancellations :

As explained in [31], for example in a $2 \rightarrow 2$ scattering process, Figure 5.1, consider the situation where the SM amplitude is approximately independent of the scattering angle in the centre-of-mass (CoM) frame, while the $\mathcal{O}(1/\Lambda^2)$ contribution exhibits a non-trivial angular dependence. Since the linear (interference) contribution to the differential cross section is proportional to the product of the two amplitudes, $d\sigma_{\text{int}}/d\Omega$ alternates in sign across the angular phase space and therefore oscillates around zero. As a result, when integrating over the scattering angle to obtain the total interference cross section, the positive and negative regions compensate each other, yielding a vanishing net contribution. In contrast, the integral of the absolute interference rate, $\int d\Omega |d\sigma_{\text{int}}/d\Omega|$, is significantly larger, indicating that the null result for the total interference is driven by strong cancellations rather than by a uniformly suppressed differential interference term.

5.2 Benchmark :

In Monte Carlo event generators (e.g. MADGRAPH), each generated event is assigned an event weight w . After imposing generation-level selection cuts, the linear (interference)

Figure 5.1: The plot on the left shows examples of a SM and a $\mathcal{O}(1/\Lambda^2)$ amplitude as functions of the centre-of-mass scattering angle θ in a $2 \rightarrow 2$ process. The plot on the right sketches the corresponding differential interference cross section, signed and in absolute value, as a function of the same angle [31].



cross section can be estimated by summing the weights of the accepted events. In the large-sample limit, this Monte Carlo estimator converges to the corresponding phase-space integral,

$$\sigma^{1/\Lambda^2} = \int d\Phi \frac{d\sigma^{1/\Lambda^2}}{d\Phi} = \lim_{N \rightarrow \infty} \sum_{i=1}^N w_i. \quad (5.1)$$

Since for the interference the weights w_i are not positive definite, cancellations between positive- and negative-weight contributions can significantly suppress σ^{1/Λ^2} , particularly when the two are comparable in value. It is therefore useful to also consider the absolute (integrable) cross section, defined as

$$\sigma^{|int|} = \int d\Phi \left| \frac{d\sigma^{1/\Lambda^2}}{d\Phi} \right| = \lim_{N \rightarrow \infty} \sum_{i=1}^N |w_i|. \quad (5.2)$$

If, for the linear contribution to a given process, this quantity is significantly larger than the total cross section σ^{1/Λ^2} , this may be interpreted as evidence of a phase-space cancellation between two large interference contributions of opposite sign.

It should be noted, however, that although $\sigma^{|int|}$ can be straightforwardly computed from simulated events, it is not generally directly accessible experimentally. Its evaluation requires complete knowledge of the kinematics and quantum numbers of all particles in the event, which is not available at the LHC. In particular, the partonic initial state is unknown on an event-by-event basis, and the presence of neutrinos prevents a full reconstruction of the event. Moreover, final-state particles must be inferred from detector signals, so that both their four-momenta and their identities can only be reconstructed approximately.

5.3 Strategy used :

In practice, we generate Wilson-coefficient configurations such that the total cross section vanishes, or remains compatible with zero within the estimated uncertainty. For simplicity, we restrict ourselves to coefficient-space vectors of the form

$$\vec{c} = \begin{pmatrix} \alpha \\ \beta \\ \gamma \\ 0 \\ 0 \\ 0 \\ 0 \\ 0 \\ 0 \\ 0 \end{pmatrix}, \quad (5.3)$$

which implies that only the coefficients $C_{qq}^{(1)}$, $C_{qq}^{(3)}$, and C_{uu} are taken to be non-zero, whereas all others are set to zero. The value of those coefficient is such that the vector in coefficient space is normalised. The cross section is first computed for a large sample of coefficient choices using a relatively small event sample. The coefficient choices that maximize $\sigma^{|\text{int}|}$ are then selected and generated again with a much larger number of events. Finally, we analyze the differential distribution for a set of observables to determine which ones achieve the best separation between negative- and positive-weight events. Table 5.1 gives the direction chosen and the SM, along with their corresponding value of the cross section and the integratable cross section.

Table 5.1: LO total cross sections for the SM, linear contribution for different direction in coefficients space to $\gamma + \text{jets}$ production, in pb. For the interference, the values of the integrable cross section are also shown.

<i>Direction</i>	$pp \rightarrow \gamma jj$ LO $\mathcal{O}(1/\Lambda^2)$	
	σ	$\sigma^{ \text{int} }$
SM	$18.98 \pm 0.14\%$	$\times$
$(-4.34 \cdot 10^{-1}, -6.86 \cdot 10^{-2}, 8.97 \cdot 10^{-1})$	$-1.648 \cdot 10^{-3} \pm 5\%$	$9.1 \cdot 10^4$

A wide range of observables is available for this analysis, including the invariant mass of the $\gamma + jj$ system, $M_{\gamma jj}$, the dijet invariant mass, M_{jj} , the azimuthal separation between the two jets, $\Delta\phi_{jj}$, and, following the analysis of Chapter 4, the photon transverse momentum, p_T^γ . In addition to these standard kinematic variables, we also investigated other observables, which are defined as follows :

We define the **transverse Sphericity** [37], Sph_T , as

$$\text{Sph}_T = \frac{2\lambda_2}{\lambda_1 + \lambda_2}. \quad (5.4)$$

where $\lambda_1 > \lambda_2$ are the eigenvalues of the transverse-momentum tensor

$$M_{xy} = \frac{1}{\sum_{i=1}^{N_{\text{jets}}} |p_{T,i}|} \sum_{i=1}^{N_{\text{jets}}} \frac{1}{|p_{T,i}|} \begin{pmatrix} p_{x,i}^2 & p_{x,i}p_{y,i} \\ p_{x,i}p_{y,i} & p_{y,i}^2 \end{pmatrix}, \quad (5.5)$$

It represents a useful quantity to differentiate pencil-like events ($\text{Sph}_T \sim 0$) from more isotropic ones ($\text{Sph}_T \sim 1$).

We also define the **Zeppenfeld variable** [38] for the photon as :

$$Z_\gamma = \left| \eta_\gamma - \frac{\eta_{j1} - \eta_{j2}}{2} \right| \quad (5.6)$$

It helps differentiate whether the photon was produced at the middle of the two jets in rapidity or close to one.

For each of these observables, we defined an asymmetry by summing the bin contents of their linear differential distributions in absolute value, and then we compared the results against $\sigma^{|int|}$ to determine which ones are more efficient in separating the positive- and negative- weighted events in our samples. The results for some of those are shown in Table 5.2 . Some variables, like the transverse sphericity, can achieve an efficiency of about 17%. As a comparison, the invariant mass of the dijet system can only revive 8% of the integrable one.

Some of the observables are plotted in Figure 5.2 and Figure 5.3 with their associated numerical uncertainties. These uncertainties are estimated separately for the positive- and negative-weight event distributions, then summed to obtain an estimate of the total uncertainty on each distribution. It should be noted that this approach serves only as an order-of-magnitude indication; a more rigorous treatment might need to check their correlations.

Table 5.2: Comparison for some observables of their power to revive the interference for $\gamma + \text{jets}$ production at LO.

$pp \rightarrow \gamma jj$ LO	
Obs.	% of $\sigma^{ int }$
$\sigma^{ int }$	100
Z_γ	34.4
p_T^γ	18
$\Delta\phi_{jj}$	17.8
Sph_T	17.4
$M_{\gamma jj}$	13.4
M_{jj}	8.6
σ^{1/Λ^2}	0

The best observables to revive the interference turned out to be Z_γ , p_T^γ and $\Delta\phi_{jj}$.

Figure 5.2: LO differential distributions for the direction presenting the higher value of σ^{int} of the photon transverse momentum (top) and the azimuthal separation of the jets (bottom) in $\gamma + \text{jets}$ production, in the linear (black) order. The SM is also shown (dashed green curve). The positive- and negative-weighted contributions to the linear term are shown separately through the shaded histograms. The numerical uncertainties are represented with shaded bands.

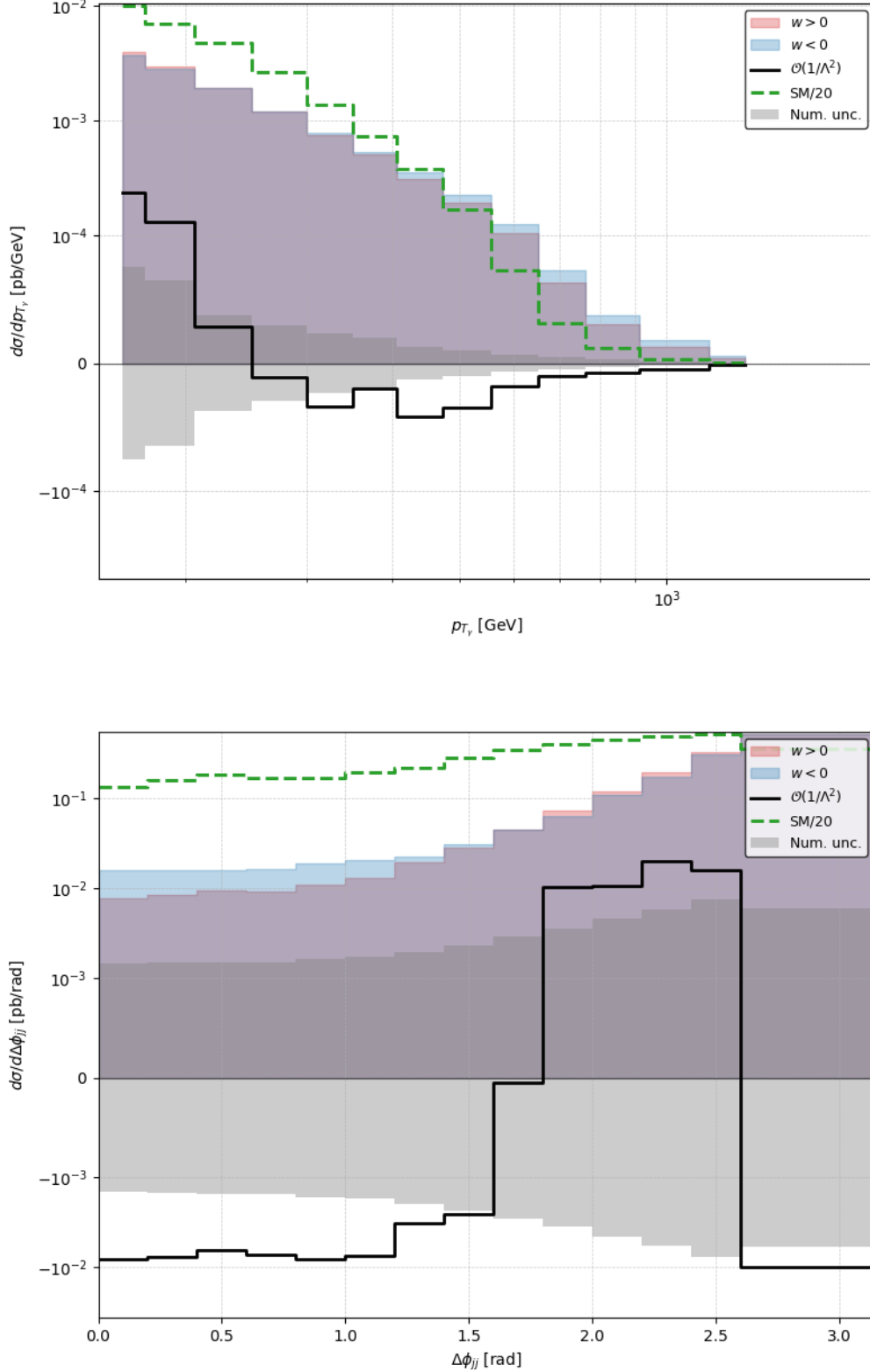
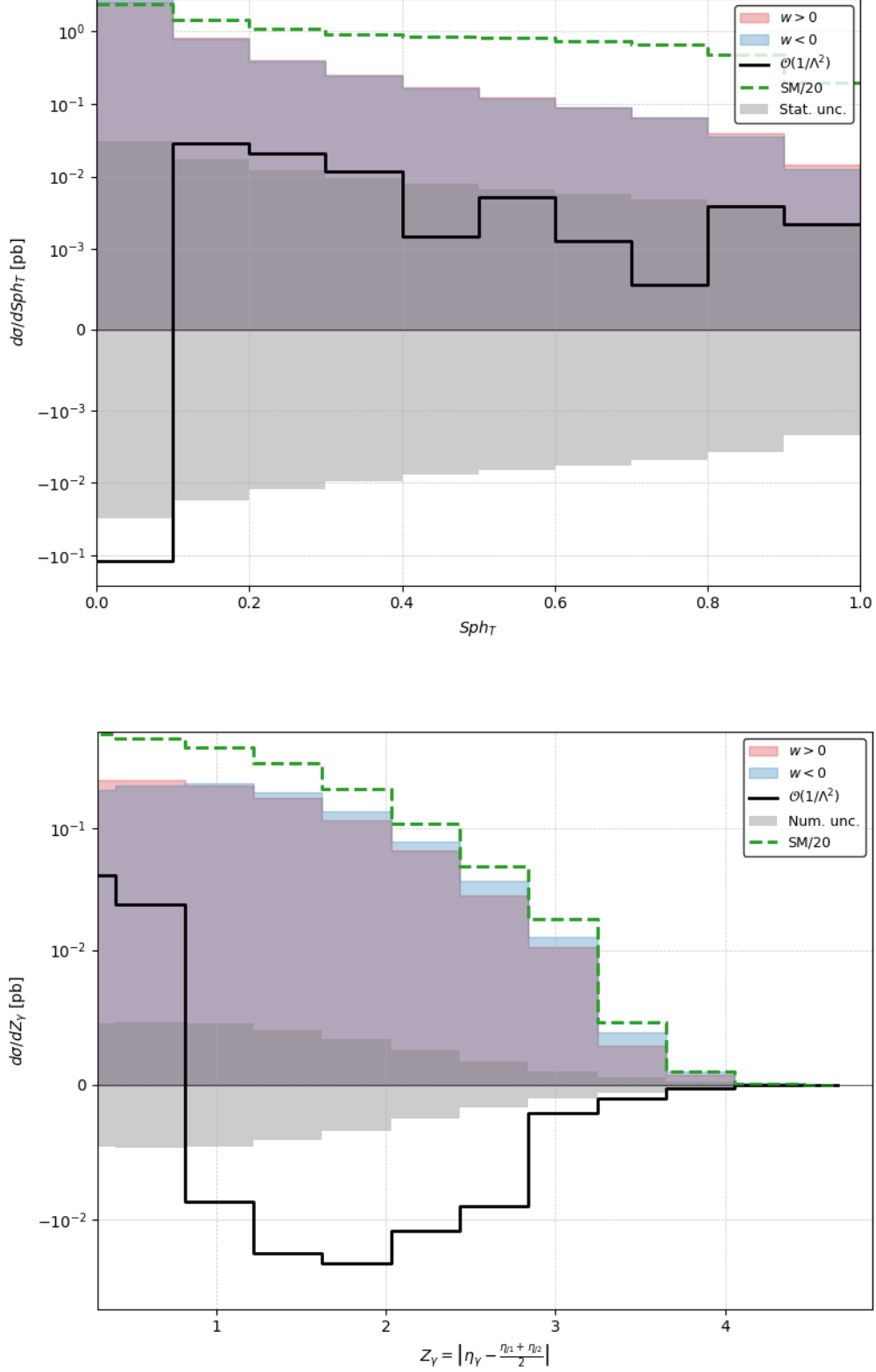


Figure 5.3: LO differential distributions for the direction presenting the higher value of σ^{int} of the transverse sphericity (top) and the zeppenfeld variable for the photon (bottom) in $\gamma + \text{jets}$ production, in the linear (black) order. The SM is also shown (dashed green curve). The positive- and negative-weighted contributions to the linear term are shown separately through the shaded histograms. The numerical uncertainties are represented with shaded bands.



5.4 Another direction :

Having identified a direction in coefficient space that is orthogonal to the interference cross section, we now turn to the direction orthogonal to both the interference cross section and the direction that we have found. The corresponding values of the total cross section and the integrable cross section are listed in Table 5.3.

Table 5.3: LO total cross section for linear contribution for the orthogonal direction in coefficients space to $\gamma + \text{jets}$ production, in pb. The value of the integrable cross section is also shown.

Direction	$pp \rightarrow \gamma jj$ LO $\mathcal{O}(1/\Lambda^2)$	
	σ	σ^{int}
$(-8.686 \cdot 10^{-1}, 2.947 \cdot 10^{-1}, -3.983 \cdot 10^{-1})$	$-1.648 \cdot 10^{-3} \pm 2.6\%$	$1.2 \cdot 10^4$

We have then studied the invariant mass of the $\gamma + jj$ system, $M_{\gamma jj}$, the dijet invariant mass, M_{jj} , the azimuthal separation between the two jets, $\Delta\phi_{jj}$, the photon transverse momentum, p_T^γ , the transverse sphericity, Sph_T , and the Zeppenfeld variable for the photon, Z_γ . Some of them are plotted with the numerical uncertainty shown as before these uncertainties are estimated separately for the positive-and negative-weight event distributions, then summed to obtain an estimate of the total uncertainty on each distribution. Their power to revive the interference is given in Table 5.4.

Table 5.4: Comparison for some observables of their power to revive the interference for $\gamma + \text{jets}$ production at LO.

$pp \rightarrow \gamma jj$ LO	
Obs.	% of σ^{int}
σ^{int}	100
Z_γ	15
Sph_T	13
$\Delta\phi_{jj}$	12
p_T^γ	10
$M_{\gamma jj}$	5
M_{jj}	3
σ^{1/Λ^2}	0

We can thus notice that the power to revive the interference in this direction is greatly reduced than in the other direction. This could be explained by the smaller value of the σ^{int} suggesting that this direction is harder to constrain.

Figure 5.4: LO differential distributions in the orthogonal direction to the direction presenting the higher value of σ^{int} of the photon transverse momentum (top) and the azimuthal separation of the jets (bottom) in $\gamma + \text{jets}$ production, in the linear (black) order. The SM is also shown (green curve). The positive- and negative-weighted contributions to the linear term are shown separately through the shaded histograms the numerical uncertainty is shown in shaded band.

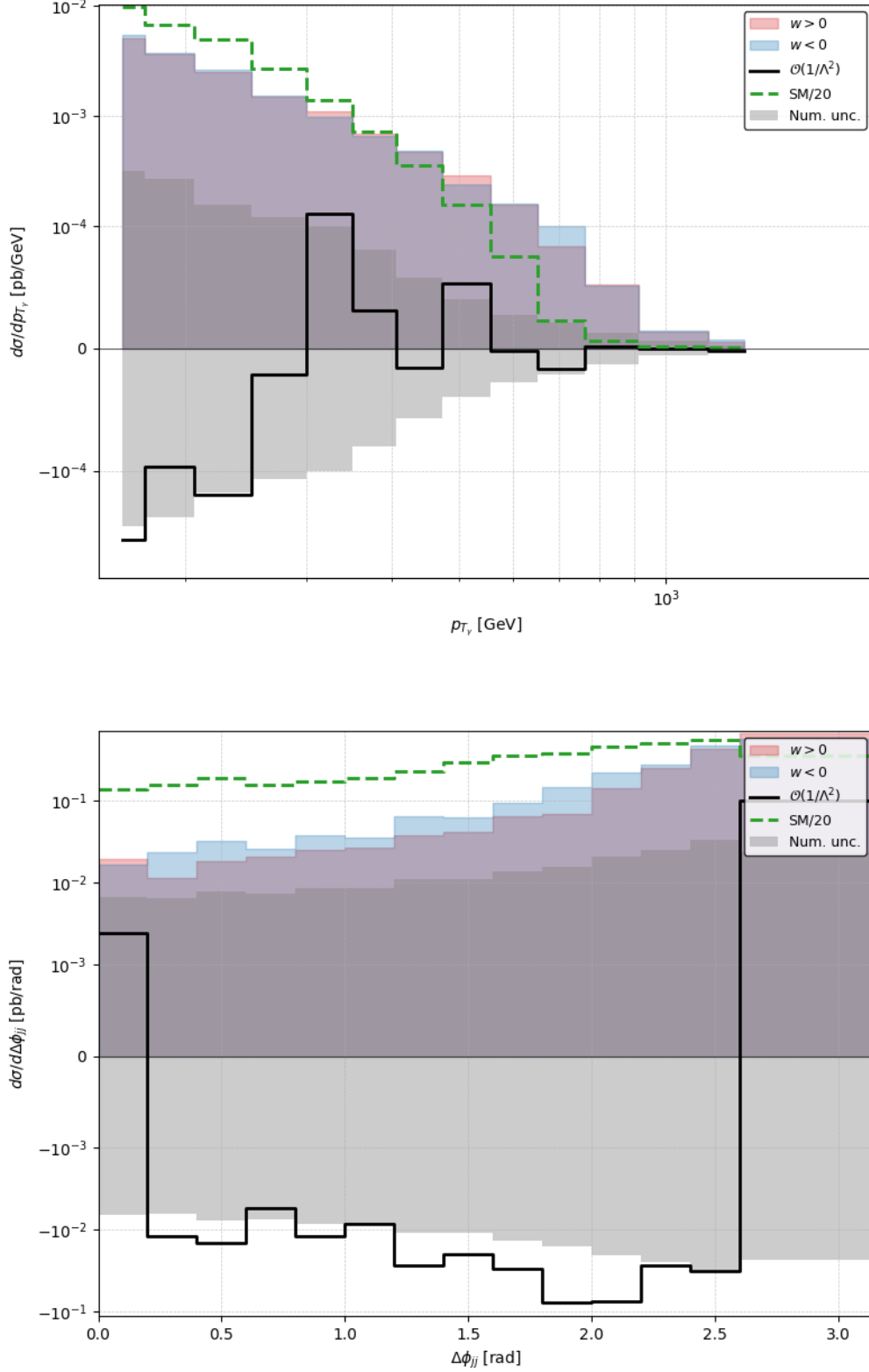
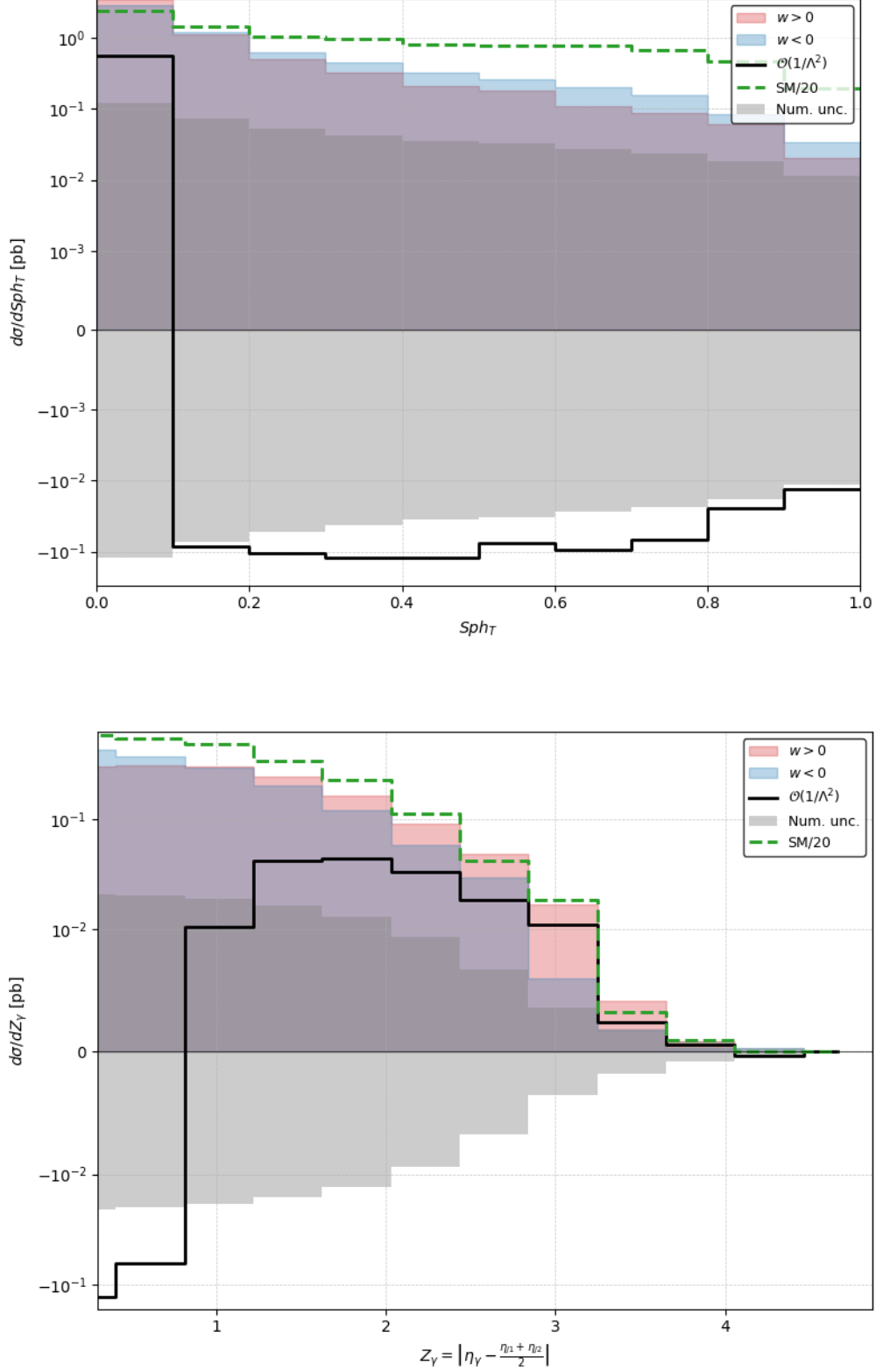


Figure 5.5: LO differential distributions in the orthogonal direction to the direction presenting the higher value of σ^{int} of the transverse sphericity (top) and the zeppenfeld variable for the photon (bottom) in $\gamma + \text{jets}$ production, in the linear (black) order. The SM is also shown (green curve). The positive- and negative-weighted contributions to the linear term are shown separately through the shaded histograms. The numerical uncertainties are represented with shaded bands



5.5 Bounds on the directions :

We can now derive bounds in both directions studied. To do so, we again perform a χ -square analysis, defined by the expression

$$\chi^2 = \sum_{i=1}^{N_{bins}} \frac{1}{\Delta_i^2} \left(\frac{d\sigma_i^{\text{exp}}}{dX} - \frac{d\sigma_i^{\text{SM}}}{dX} - A \frac{d\sigma_i^{1/\Lambda^2}}{dX} \right)^2 \quad (5.7)$$

$$= \sum_{i=1}^{N_{bins}} \frac{1}{\Delta_i^2} \left(A \frac{d\sigma_i^{1/\Lambda^2}}{dX} \right)^2 \quad (5.8)$$

We have here assumed that the experimental data follows the SM distribution that we have generated and Δ_i was assigned to 10% of the SM value in each bin. The binnings are chosen in order to guarantee the 10% accuracy. These numbers of events are calculated, from the differential cross-section in each bin, taking the value of $\mathcal{L} = 139 \text{ fb}^{-1}$ which is constant with the value of the CMS/Atlas experiment during Run 2 [39]. The bounds are placed on A and can thus be translated into bounds on the directions. We show respectively in Table 5.5 and in Table 5.6 the bounds for the direction presenting the highest value of $\sigma^{|int|}$ and the direction orthogonal to it.

Table 5.5: 95% CL bounds on the direction presenting the highest value of $\sigma^{|int|}$ for a list of observables, obtained at the linear level for $\gamma + \text{jets}$ production at LO.

Distribution	$\mathcal{O}(1/\Lambda^2)$ 95% bounds
p_T^γ	[-1.51, 1.51]
Z_γ	[-1.21, 1.21]
Sph_T	[-3.05, 3.05]
$\Delta\phi_{jj}$	[-3.51, 3.51]

We notice that for this direction Z_γ gives the strongest bounds.

Table 5.6: 95% CL bounds on the direction orthogonal to the direction presenting the highest value of $\sigma^{|int|}$ for a list of observables, obtained at the linear level for $\gamma + \text{jets}$ production at LO.

Distribution	$\mathcal{O}(1/\Lambda^2)$ 95% bounds
p_T^γ	[-2.21, 2.21]
Z_γ	[-5.42, 5.42]
Sph_T	[-6.33, 6.33]
$\Delta\phi_{jj}$	[-10.40, 10.40]

We notice that for this direction p_T^γ gives the strongest bounds. And that the first direction presents the tightest bounds overall.

5.6 Double distributions

We also tried double-differential distributions to set bounds on the directions. The results for the two directions are summarized in table 5.7 and 5.8. And as we can see, the first direction posses again the best bounds overall and we also have that for the two directions the best observable is p_T^γ vs $\Delta\phi_{jj}$.

Table 5.7: 95% CL bounds on the direction presenting the highest value of $\sigma^{|int|}$ for a list of observables, obtained at the linear level for γ + jets production at LO.

Distribution	$\mathcal{O}(1/\Lambda^2)$ 95% bounds
p_T^γ vs $\Delta\phi_{jj}$	[-0.06, 0.06]
Z_γ vs p_T^γ	[-1.23, 1.23]
Z_γ vs $\Delta\phi_{jj}$	[-2.19, 2.19]

Table 5.8: 95% CL bounds on the direction orthogonal to the direction presenting the highest value of $\sigma^{|int|}$ for a list of observables, obtained at the linear level for γ + jets production at LO.

Distribution	$\mathcal{O}(1/\Lambda^2)$ 95% bounds
p_T^γ vs $\Delta\phi_{jj}$	[-0.43, 0.43]
Z_γ vs p_T^γ	[-1.29, 1.29]
Z_γ vs $\Delta\phi_{jj}$	[-5.69, 5.69]

In Fig. 5.7 we show the distributions for direction associated to the largest value for $\sigma^{|int|}$ in the p_T^γ vs. $\Delta\phi_{jj}$, $\Delta\phi_{jj}$ vs. Z_γ , and p_T^γ vs. Z_γ . However, for p_T^γ vs. $\Delta\phi_{jj}$ we do not see a clear separation between the negative and positive regions. Therefore, this distribution does not give much additional information, but the other two, namely, $\Delta\phi_{jj}$ vs. Z_γ , and p_T^γ vs. Z_γ seem more promising. In each case, regions of near-zero cross section are observed. A non-rectangular binning could improve the analysis by merging these low-sensitivity bins into a single bin while providing finer resolution in the more relevant kinematic regions.

5.7 Bounds on the Wilson coefficients :

Following this analysis, we now have at our disposal a set of observables that are of interest for further study. We therefore proceed to determine the individual bounds associated with the double distributions presented in Table 5.9 using a χ^2 analysis. We note that even if we do not use the same number of bins for the all observables considered, the differences remain small and the bin counts are sufficiently large in all cases. Therefore, even if not all the distributions have the same number of bins, we can compare their results, as they would change slightly with a different binning. From this, we observe that the distribution providing the strongest bounds is p_T^γ vs $\Delta\phi_{jj}$. We therefore use it to perform our global χ^2 analysis, from which we extract the directions, summarized in Table 5.10. We observe that the most constrained direction remains close to $C_{qq}^{(3)}$.

Table 5.9: 95% CL bounds on the wilson coefficients, obtained at the linear level for γ + jets production at LO. The number of bins is also shown.

Distribution	N_{bins}	$\mathcal{O}(1/\Lambda^2)$ 95% CL bounds		
		$C_{qq}^{(1)}$	$C_{qq}^{(3)}$	C_{uu}
p_T^γ vs $\Delta\phi_{jj}$	34	[-0.047, 0.047]	[-0.014, 0.014]	[-0.044, 0.044]
Z_γ vs $\Delta\phi_{jj}$	28	[-0.928, 0.928]	[-0.272, 0.272]	[-1.158, 1.158]
Z_γ vs p_T^γ	28	[-0.363, 0.363]	[-0.095, 0.095]	[-0.346, 0.346]

Table 5.10: Best constrained directions in coefficient space and their associated eigenvalues, for γ + jets production by the photon transverse momentum, p_T^γ .

Operator	Coefficient	Operator	Coefficient	Operator	Coefficient
C_{qq}^1	0.28	C_{qq}^1	0.55	C_{qq}^1	0.78
C_{qq}^3	0.91	C_{qq}^3	-0.40	C_{qq}^3	-0.04
C_{uu}	0.30	C_{uu}	0.72	C_{uu}	-0.62
λ	21653.37	λ	13.66	λ	4.73

We show in Figure 5.6 the distribution for the first three 4LQ operators, and we notice that, again, for the $\Delta\phi_{jj}$ vs p_T^γ observable, no remarkable distinctions can be identified among the different 4LQ operators.

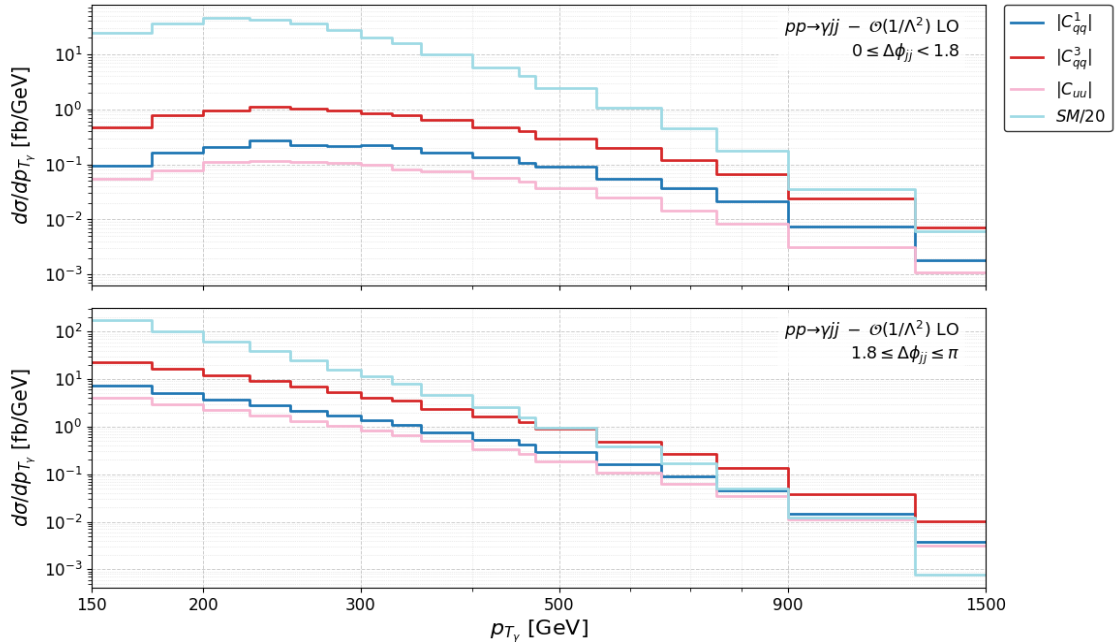
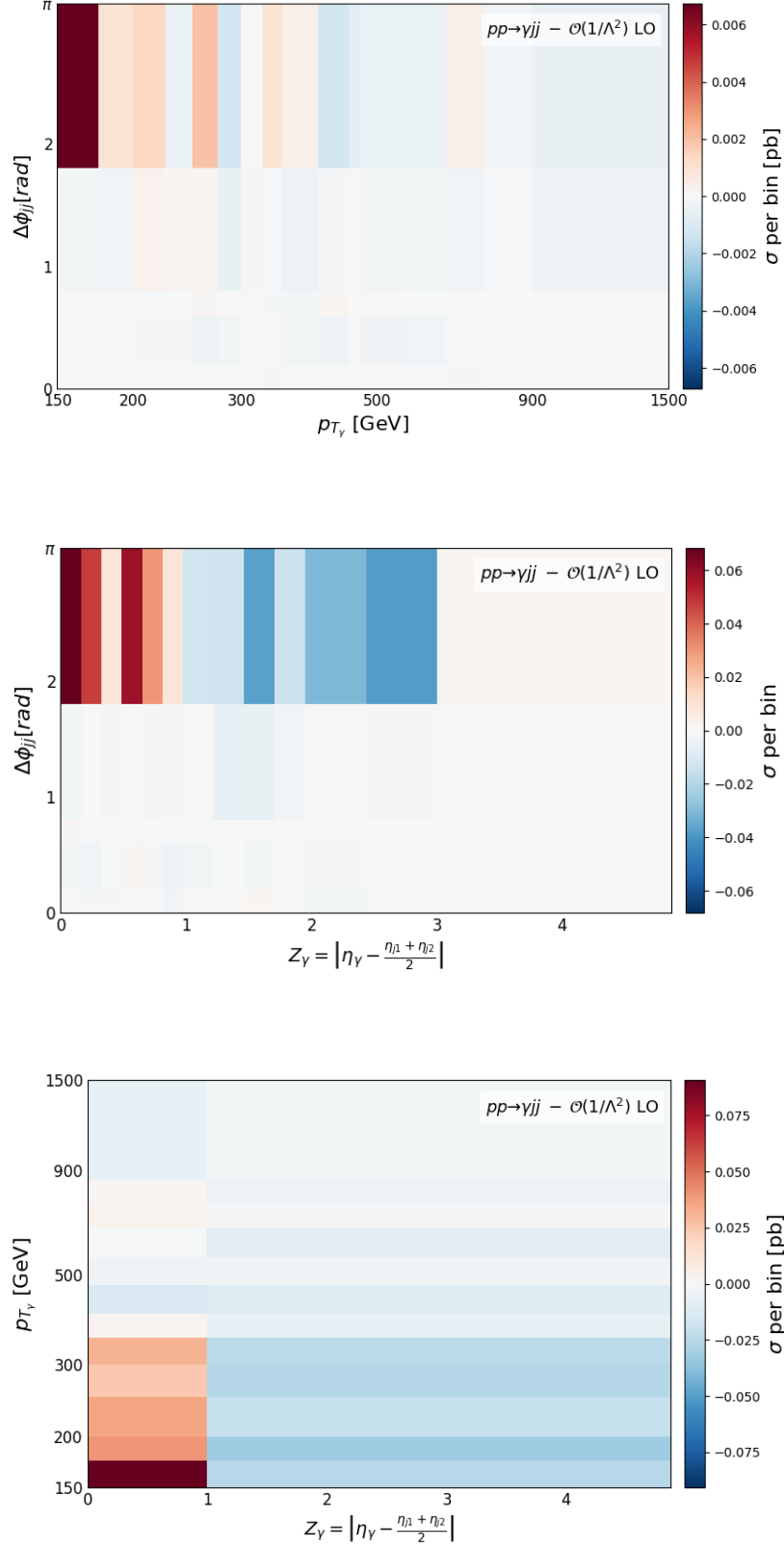
Figure 5.6: Differential distribution of the photon transverse momentum in γ + jets production separated in $\Delta\phi_{jj}$ regions, for the SM (light blue) and the interference of the three first 4LQ operators, at LO, with all coefficients set to $C_i/\Lambda^2 = 1 \text{ TeV}^{-2}$.

Figure 5.7: $\gamma + \text{jets}$ production interference cross section per bin at LO, as a function of p_T^γ and $\Delta\phi_{jj}$ (top), of $\Delta\phi_{jj}$ and Z_γ (middle), and of p_T^γ and Z_γ (bottom). Red (blue) areas mark where the cross section is positive (negative), as the positive- (negative-) weighted contribution dominates there.



We also plot in Figure 5.8 the two dimensional contours and we notice that compared to the constraints obtained with p_T^γ only, $\Delta\phi_{jj}$ vs p_T^γ further shrinks the allowed parameter space to a narrower region, indicating improved sensitivity but unfortunately along essentially the same correlated direction in Wilson-coefficient space. That means that $\Delta\phi_{jj}$ vs p_T^γ does not change the main correlation direction very much; instead, it mainly sharpens the constraint along the same degeneracy already present in the p_T^γ fit. We also have seen that the other two distributions considered namely Z_γ vs $\Delta\phi_{jj}$ and Z_γ vs p_T^γ unfortunately showcased the same behavior as shown in Figure 5.9. To close this chapter, we present the contours in Figure 5.9 obtained with three sets of observables, building from the cross-section up to double differential distributions : the total cross-section, the single differential distribution in $p_{T,\gamma}$, and the double differential distributions in p_T^γ versus $\Delta\phi_{jj}$, Z_γ versus $\Delta\phi_{jj}$, and Z_γ versus p_T^γ .

Figure 5.8: Individual contours in the $C_{qq}^{(1)}$ vs $C_{qq}^{(3)}$ plane (*top left*), in the C_{uu} vs $C_{qq}^{(1)}$ one (*top right*), and in the C_{uu} vs $C_{qq}^{(3)}$ one (*bottom middle*) for γ + jets production by $p_T^\gamma, \Delta\phi_{jj}$.

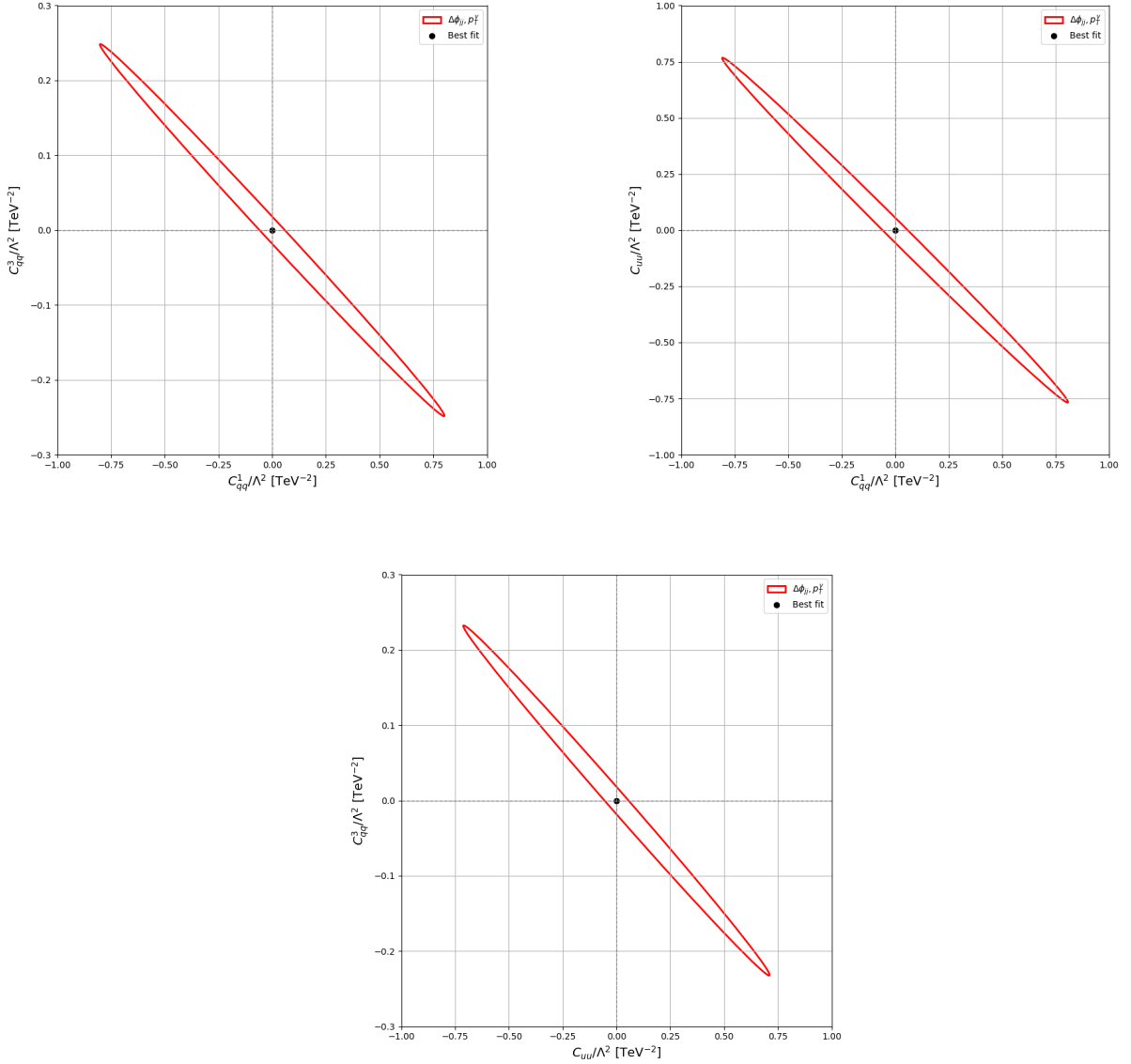
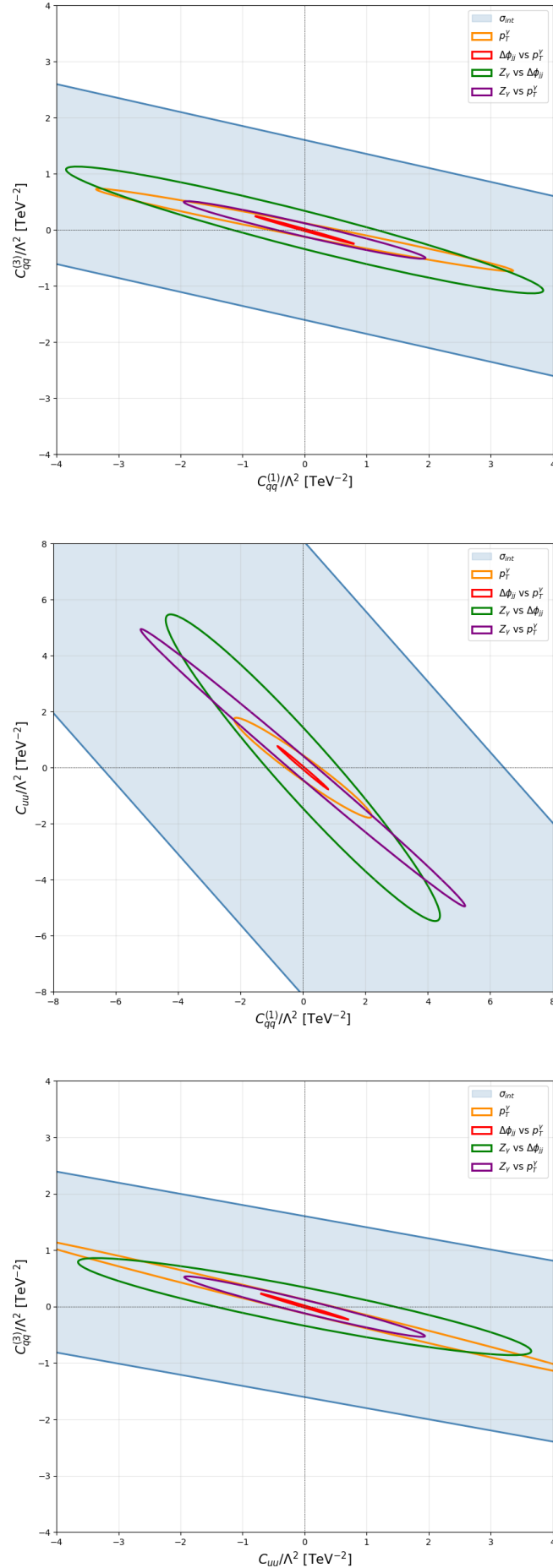


Figure 5.9: Individual contours in the $C_{qq}^{(1)}$ vs $C_{qq}^{(3)}$ plane (*top left*), in the C_{uu} vs $C_{qq}^{(1)}$ one (*middle*), and in the C_{uu} vs $C_{qq}^{(3)}$ one (*bottom*) for $\gamma + \text{jets}$ production by all the observables considered in this study.



Chapter 6

Going beyond the linear level :

Up to now, our analysis has been restricted to the linear SMEFT contributions to the cross section, which represents only a leading-order approximation. To extend this framework, we include the quadratic dimension-6 terms, $\mathcal{O}(\Lambda^{-4})$. To keep the analysis simple, we restrict our focus to a subset of three 4LQ operators: $\mathcal{O}_{qq}^{(1)}$, $\mathcal{O}_{qq}^{(3)}$, and $\mathcal{O}_{uu}$. Consequently, the cross section includes the squared contributions of these operators along with their mutual interference terms. As in the previous chapters, the interference among dimension-8 operators and the SM was not included, due to the large number of terms that should be taken into account [40].

6.1 Analysis :

In the massless fermion limit, chirality (defined by the projection operators P_L and P_R) and helicity (the projection of spin along the momentum axis) are strictly equivalent. Because $\mathcal{O}_{qq}^{(1,3)}$ and $\mathcal{O}_{uu}$ couple to fermions of differing chiralities, they generate transition amplitudes occupying orthogonal helicity sectors. Given these distinct external helicity configurations, their amplitudes do not interfere, and the corresponding cross terms identically vanish. Therefore, under the massless assumption, the only surviving interference term within this is between $\mathcal{O}_{qq}^{(1)}$ and $\mathcal{O}_{qq}^{(3)}$. It is worth noting that the reintroduction of the fermion masses breaks the strict equivalence between chirality and helicity, thus allowing left- and right-chiral structures to mix. In such a scenario, the interference between differing chiral operators is no longer strictly zero, although it will be highly suppressed.

Table 6.1: Cross sections for $\gamma + \text{jets}$ production for the SM and the 4LQ operators considered at $\mathcal{O}(1/\Lambda^4)$ whose coefficients are set to 1 TeV^{-2} .

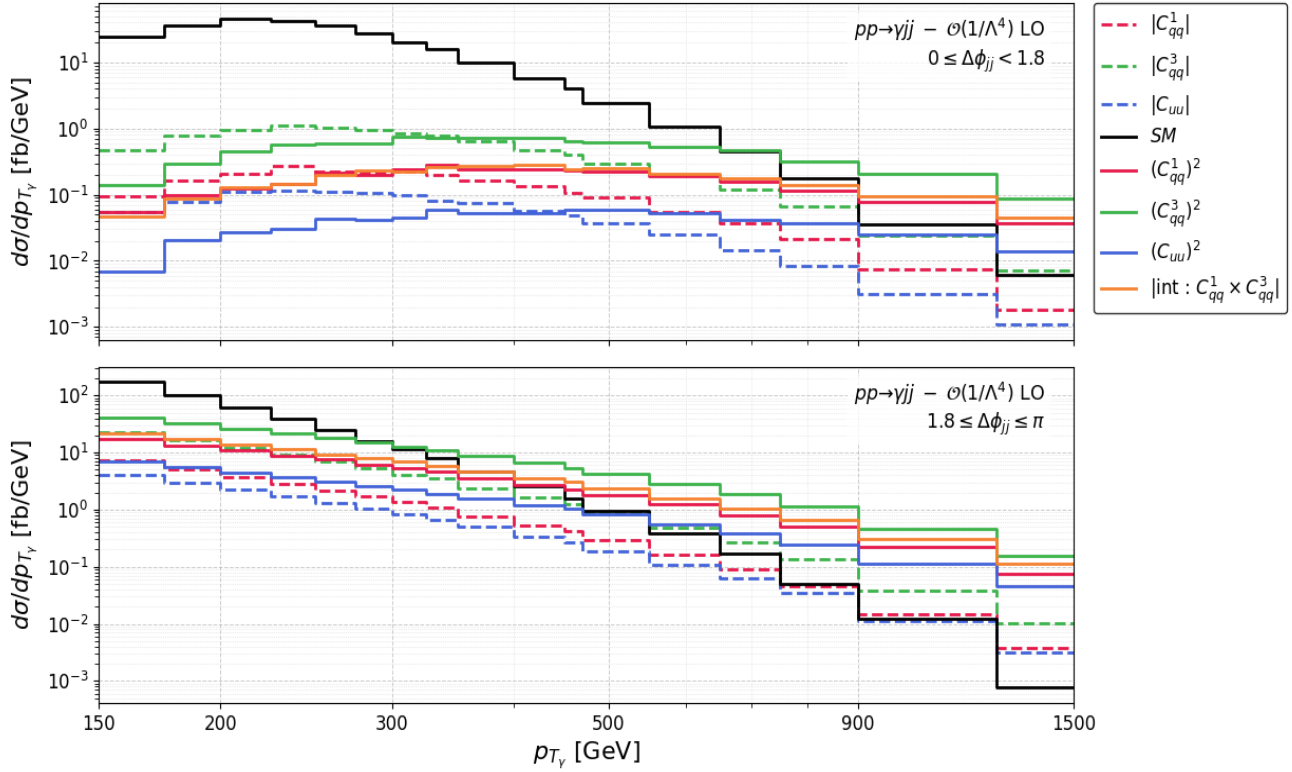
	$\gamma + \text{jets}$ (pb)
SM	$18.98 \pm 0.14\%$
$\mathcal{O}(1/\Lambda^4)$	
$\mathcal{O}_{qq}^{(1)}$	$2.916 \pm 0.09\%$
$\mathcal{O}_{qq}^{(3)}$	$7.036 \pm 0.08\%$
$\mathcal{O}_{uu}$	$1.221 \pm 0.10\%$
Interference term $\mathcal{O}(1/\Lambda^4)$	
$\mathcal{O}_{qq}^{(1)}, \mathcal{O}_{qq}^{(3)}$	$1.695 \pm 0.21\%$

We present in Table 6.1 the cross sections for the SM, the pure squared $\mathcal{O}(\Lambda^{-4})$ contributions, and the interference $\mathcal{O}(\Lambda^{-2})$ terms between the dimension-six operators.

6.2 Observable :

Following the same approach as in the previous chapter, we derive the bounds on the Wilson coefficients by performing a two-dimensional analysis for $(\Delta\phi_{jj}, p_T^\gamma)$. The corresponding differential distributions, binned in two regions of $\Delta\phi_{jj}$, are presented in Fig. 6.1. In the high- p_T^γ tail, the $\mathcal{O}(\Lambda^{-4})$ contributions are found to dominate the $\mathcal{O}(\Lambda^{-2})$ interference terms by at least a factor of three, raising concerns about the validity of the linear approximation and suggesting that higher-order terms in the $1/\Lambda$ expansion may not be negligible in this kinematic regime.

Figure 6.1: Differential distribution of the photon transverse momentum in $\gamma + \text{jets}$ production separated in $\Delta\phi_{jj}$ regions, for the SM (black), the absolute value of the interference of the three first 4LQ operators and the $\mathcal{O}(\Lambda^{-4})$ terms, at LO, with all the coefficients set to $C_i/\Lambda^2 = 1 \text{ TeV}^{-2}$



6.3 Bounds :

With this observable at hand, we proceed to extract bounds on the Wilson coefficients via a χ^2 analysis defined as:

$$\chi^2 = \sum_i^{N_{bins}} \frac{1}{\Delta_i^2} \left(\sum_j^{operators} \frac{c_j}{\Lambda^2} \frac{d\sigma_{j,i}^{1/\Lambda^2}}{dX} + \frac{c_j^2}{\Lambda^4} \frac{d\sigma_{j,i}^{1/\Lambda^4}}{dX} + \frac{c_{ij}}{\Lambda^4} \frac{d\sigma_{mix}^{1/\Lambda^4}}{dX} \right)^2 \quad (6.1)$$

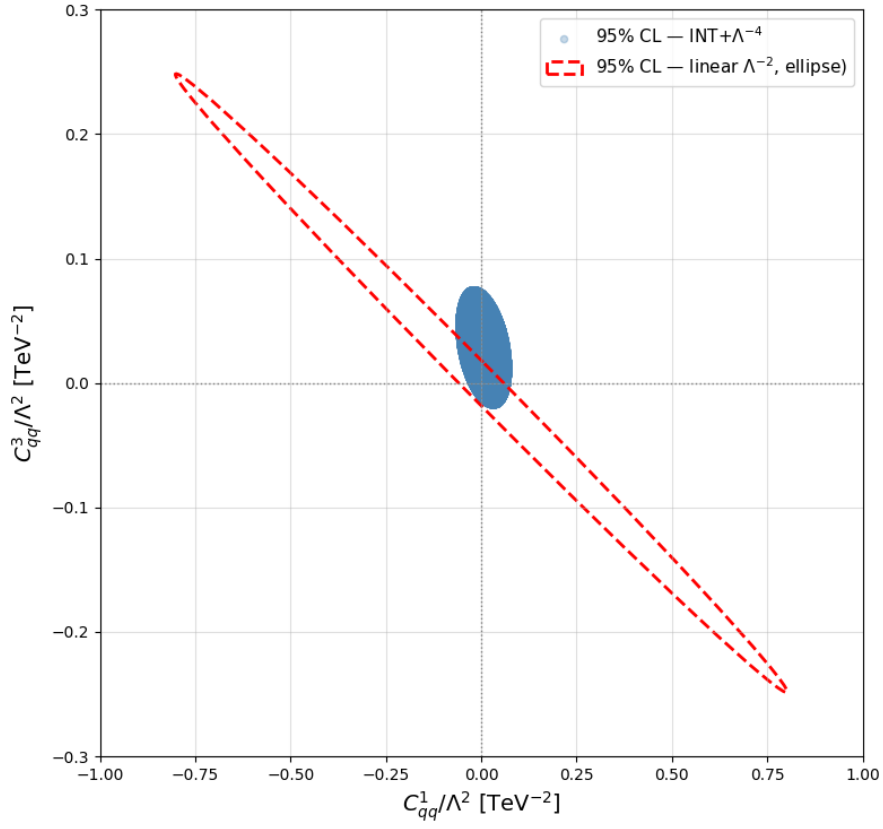
The presence of $\mathcal{O}(\Lambda^{-4})$ terms renders this expression a polynomial of degree four in the Wilson coefficients, which does not admit a straightforward analytical minimization. To circumvent this difficulty, we adopt a numerical approach inspired by Monte Carlo integration techniques, namely the Hit-and-Miss method. Having obtained the bounds from the $\mathcal{O}(\Lambda^{-2})$ analysis, we already have an indication of the region in which the true bounds are expected to lie. We can therefore define a small parameter space (a box) within this region, into which we can systematically substitute the values of the coefficients.

6.3.1 Results :

The resulting countour is presented in Fig. 6.2, where a reduction in the allowed parameter space is observed with respect to the dimension-six truncation. This behavior can be understood directly from the structure of the distributions : the non-negligible $\mathcal{O}(\Lambda^{-4})$ contributions in the high- p_T^γ tail indicate that the linear approximation alone provides an incomplete description of the EFT expansion in this kinematic regime. However, a consistent treatment of $\mathcal{O}(\Lambda^{-4})$ effects would in principle require the inclusion of dimension-eight operators, whose operator basis is significantly larger and makes a systematic analysis considerably more involved. These results should therefore be interpreted with care, and the bounds derived under the dimension-six truncation must be regarded as approximate in the high-energy tail of the distributions. We can note that the loss of symmetry of the ellipse with respect to zero is expected and arises from the fact that the χ^2 function now contains a cubic term in the coefficients.

$$\chi^2 \supset \underbrace{\left(\frac{c_j}{\Lambda^2} \frac{d\sigma_{j,i}^{1/\Lambda^2}}{dX} \right)^2}_{symmetric} + 2 \underbrace{\left(\frac{c_j}{\Lambda^2} \frac{d\sigma_{j,i}^{1/\Lambda^2}}{dX} \right) * \frac{c_j^2}{\Lambda^4} \frac{d\sigma_{j,i}^{1/\Lambda^4}}{dX}}_{not\ symmetric} \quad (6.2)$$

Figure 6.2: Individual contours in the $C_{qq}^{(1)}$ vs $C_{qq}^{(3)}$ plane for $\gamma + \text{jets}$ production by $\Delta\phi_{jj}$ vs p_T^γ including the $\mathcal{O}(\Lambda^{-4})$ for the operators considered



Chapter 7

Disentangling Singlet and Octet Contributions

Up to this point, the color-singlet and color-octet components of the 4LQ operators have been treated as part of the same operator and multiplied by a common Wilson coefficient. In the following, however, these two color structures will be regarded as independent operators, each associated with its own Wilson coefficient.

7.1 Color decomposition into singlet and octet structures

We have thus separated this color structures. We have limited ourselves to the first three operators for simplicity. We then showcase an example of this color singlet/octet separation for the operator $\mathcal{O}_{qq}^{(1)}$.

Colour decomposition of $\mathcal{O}_{qq}^{(1)}$

The operator $\mathcal{O}_{qq}^{(1)} = (\bar{q}_p \gamma^\mu q_r)(\bar{q}_s \gamma_\mu q_t)$ admits two independent flavour-index contractions, which we analyse separately. In the following, i, j, k, l denote $SU(3)_c$ fundamental colour indices.

Contraction A : $\delta_{pr}\delta_{st}$ Both generation indices are contracted *inside* each bilinear. The colour structure factorises trivially:

$$(\bar{q}_p^i \gamma^\mu q_r^i) (\bar{q}_s^j \gamma_\mu q_t^j) \delta_{pr} \delta_{st} \propto \delta^{ij} \delta^{ij} = (\bar{q}^i \gamma^\mu q^i) (\bar{q}^j \gamma_\mu q^j), \quad (7.1)$$

which is manifestly a **colour singlet**: each bilinear is individually colour-neutral, $(\bar{q}\gamma^\mu q) \sim \mathbf{1}_c$.

Contraction B : $\delta_{pt}\delta_{rs}$ Flavour indices are contracted *between* the two bilinears, effectively exchanging the quarks $r \leftrightarrow t$. In colour space, this exchange also swaps the colour lines $j \leftrightarrow i$, producing a colour structure $\delta^{il}\delta^{kj}$ (where k, l are the colour indices of q_r and q_t respectively):

$$(\bar{q}_p^i \gamma^\mu q_t^j) (\bar{q}_s^k \gamma_\mu q_r^l) \delta_{pt} \delta_{rs} \propto \delta^{il} \delta^{kj}. \quad (7.2)$$

We now apply the $SU(3)$ Fierz identity for colour indices:

$$\delta^{il} \delta^{kj} = \frac{1}{N_c} \delta^{ij} \delta^{kl} + 2 (T^A)^{ij} (T^A)^{kl}, \quad (7.3)$$

where $N_c = 3$ and T^A are the generators of $SU(3)_c$ in the fundamental representation. Substituting back into the operator:

$$\begin{aligned} \delta^{il} \delta^{kj} &= \frac{1}{3} \delta^{ij} \delta^{kl} + 2 (T^A)^{ij} (T^A)^{kl} \\ \Rightarrow (\bar{q}^i \gamma^\mu q^j) (\bar{q}^k \gamma_\mu q^l) \delta^{il} \delta^{kj} &= \frac{1}{3} \underbrace{(\bar{q}^i \gamma^\mu q^i) (\bar{q}^j \gamma_\mu q^j)}_{\text{colour singlet } \mathbf{1}_c} + 2 \underbrace{(\bar{q}^i \gamma^\mu T_{ij}^A q^j) (\bar{q}^k \gamma_\mu T_{kl}^A q^l)}_{\text{colour octet } \mathbf{8}_c}. \end{aligned} \quad (7.4)$$

Result : The full operator $\mathcal{O}_{qq}^{(1)}$, therefore contains *both* a colour-singlet and a colour-octet component:

$$\boxed{\mathcal{O}_{qq}^{(1)} \supset \underbrace{(\bar{q} \gamma^\mu q) (\bar{q} \gamma_\mu q)}_{\mathbf{1}_c} \& \underbrace{(\bar{q} \gamma^\mu T^A q) (\bar{q} \gamma_\mu T^A q)}_{\mathbf{8}_c}}. \quad (7.5)$$

This is why $\mathcal{O}_{qq}^{(1)}$ interferes with *both* EW-boson and gluon-mediated diagrams in the $pp \rightarrow \gamma jj$ process.

7.2 Framework:

We continue our analysis by considering the same process as throughout this work, namely $pp \rightarrow \gamma jj$. All events were generated with MADGRAPH5_aMC@NLO using a UFO model that includes the three selected operators, with the color-singlet and color-octet contributions treated separately, in addition to the SM contribution. The cuts are identical to those used in the previous analysis. As in the rest of this study, we restrict ourselves to leading order and parton level. The total Standard Model cross section is recalled in Table 7.1, together with the contribution of each operator.

7.3 Analysis and differential distribution :

To see the different effects of high and low energy region we have made the choice to separate the observables into two m_{jj} regions. To see the different effects of the high- and low-energy regions, we separate the observables into two m_{jj} regions. The boundary is taken at $m_{jj} = 800$ GeV, which approximately marks the transition between the peak of the SM m_{jj} spectrum and the rapidly falling high-mass tail, as visible in Fig. 7.1. In the spirit of what was done during this thesis, we have looked at two distributions namely p_T^γ and $\Delta\Phi_{jj}$ and the differential distribution for them separated by m_{jj} region are shown in Figure 7.2, where we applied the binning [150, 175, 200, 250, 300, 350, 400, 470, 550, 650, 750, 900, 1100, 1500] GeV with the last one including the overflow for p_T^γ as for $\Delta\Phi_{jj}$ used we have used a binning of 0.2 bins for [0, 2.8] and a last bin from 2.8 to π .

Figure 7.1: Differential distribution of the invariant mass of the dijet system, m_{jj} , in γ + jets production, for the SM (light blue) and the absolute value of the interference of the 6 operators included in this analysis at LO, with all coefficients set to $C_i/\Lambda^2 = 1 \text{ TeV}^{-2}$.

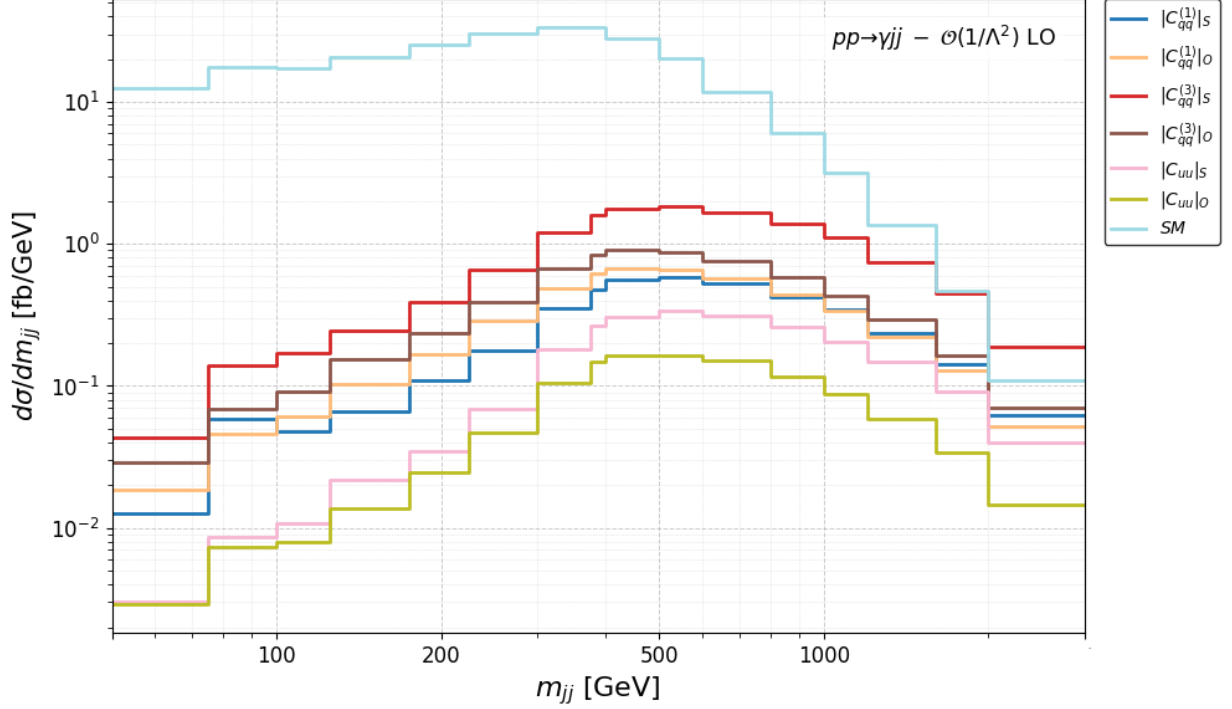
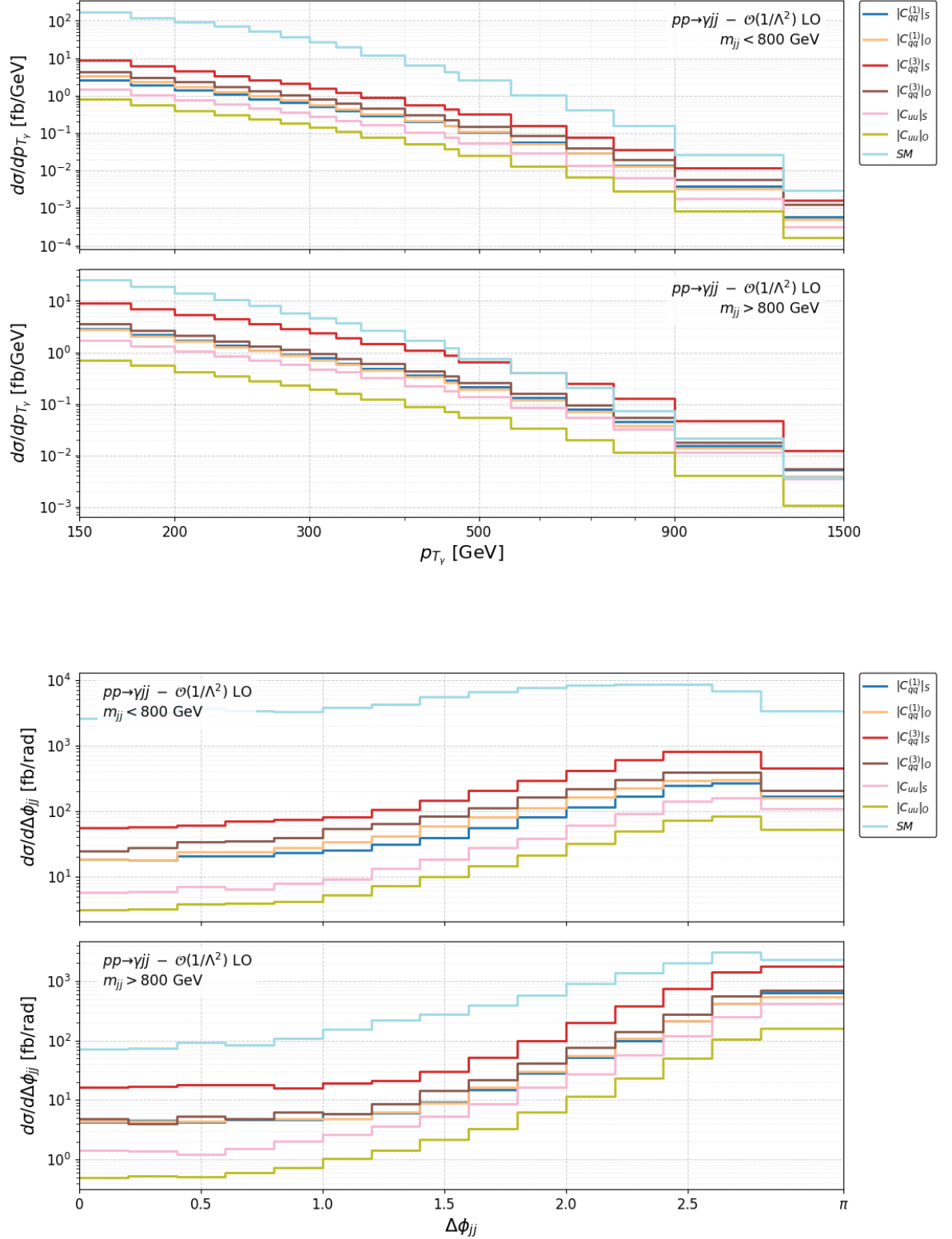


Table 7.1: Cross sections (in pb) for γ + jets production for the SM and the studied operators at $\mathcal{O}(1/\Lambda^2)$ whose coefficients are set to 1 TeV^{-2} .

γ + jets (pb)	
SM	$18.98 \pm 0.14\%$
$\mathcal{O}(1/\Lambda^2)$	
$\mathcal{O}_{qq}^{(1),1}$	$-0.464 \pm 0.16\%$
$\mathcal{O}_{qq}^{(1),8}$	$-0.600 \pm 0.12\%$
$\mathcal{O}_{qq}^{(3),1}$	$-1.974 \pm 0.11\%$
$\mathcal{O}_{qq}^{(3),8}$	$0.415 \pm 0.23\%$
$\mathcal{O}_{uu}^1$	$-0.394 \pm 0.1\%$
$\mathcal{O}_{uu}^8$	$-0.160 \pm 0.12\%$

For both the observables, the 4LQ linear shapes are different from the SM ones, but no remarkable distinctions can be identified among the operators. We still notice that the difference in azimuthal angle between jets seems to be able to present larger 4LQ effects compared to the SM ones, by at least a factor four in the tail of its distribution. For this reason, we used $\Delta\phi_{jj}$ to extract bounds over the coefficients.

Figure 7.2: Differential distribution of the photon transverse momentum (top) and the azimuthal separation of the jets (bottom), separated by m_{jj} region, in $\gamma + \text{jets}$ production, for the SM (light blue) and the absolute value of the interference of the 6 operators included in this analysis at LO, with all coefficients set to $C_i/\Lambda^2 = 1 \text{ TeV}^{-2}$.



7.4 Bounds on the coefficients

To derive bounds on the Wilson coefficients, we perform a χ^2 analysis defined as

$$\chi^2(\vec{c}) = \sum_{i=1}^{N_{\text{bins}}} \frac{1}{\Delta_i^2} \left(\frac{d\sigma_i^{\text{exp}}}{dX} - \frac{d\sigma_i^{\text{SM}}}{dX} - \sum_j \frac{c_j}{\Lambda^2} \frac{d\sigma_{j,i}^{1/\Lambda^2}}{dX} \right)^2, \quad (7.6)$$

The resulting bounds are displayed in Figure 7.3, together with the most constrained direction in Wilson coefficient space for each m_{jj} region in Table 7.2. In both regions, this direction is found to be predominantly aligned with the operator $\mathcal{O}_{qq}^{(3),1}$, which consequently receives the tightest individual bound among all operators considered. Furthermore, the high-invariant-mass region $m_{jj} > 800$ GeV consistently yields more stringent constraints than the complementary region across all observables considered.

These results must, however, be interpreted with care. While the high- m_{jj} region provides the strongest bounds, it is also the kinematic regime in which the validity of the EFT expansion is most vulnerable : the condition $m_{jj} \ll \Lambda$ may no longer be well satisfied, potentially rendering the truncation of the EFT series at $\mathcal{O}(\Lambda^{-2})$ unreliable. The validity of the EFT expansion in this region should therefore be checked before drawing definitive conclusions from these constraints. Individual contours in various planes of Wilson coefficient space are shown in Figure 7.4.

Figure 7.3: Individual limits, at 95% CL, on the 4LQ operators considered in the study from $\gamma + \text{jets}$ production with $\Delta\phi_{jj}$ separated in m_{jj} regions.

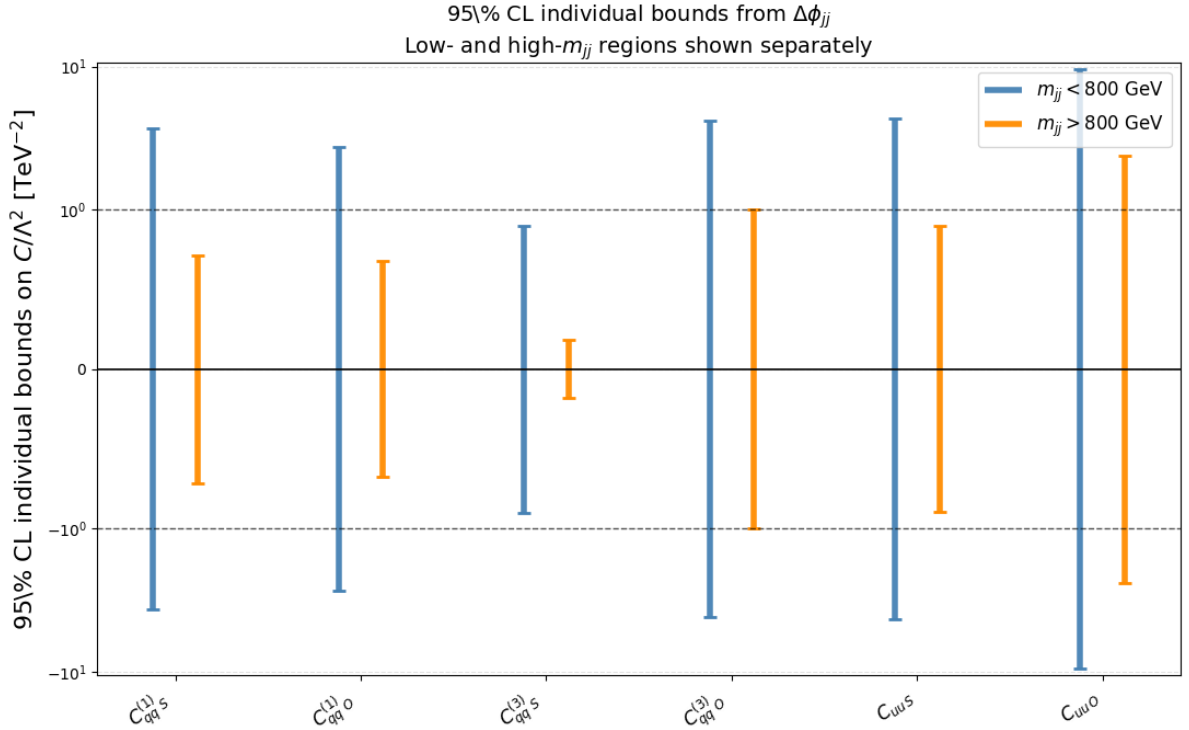
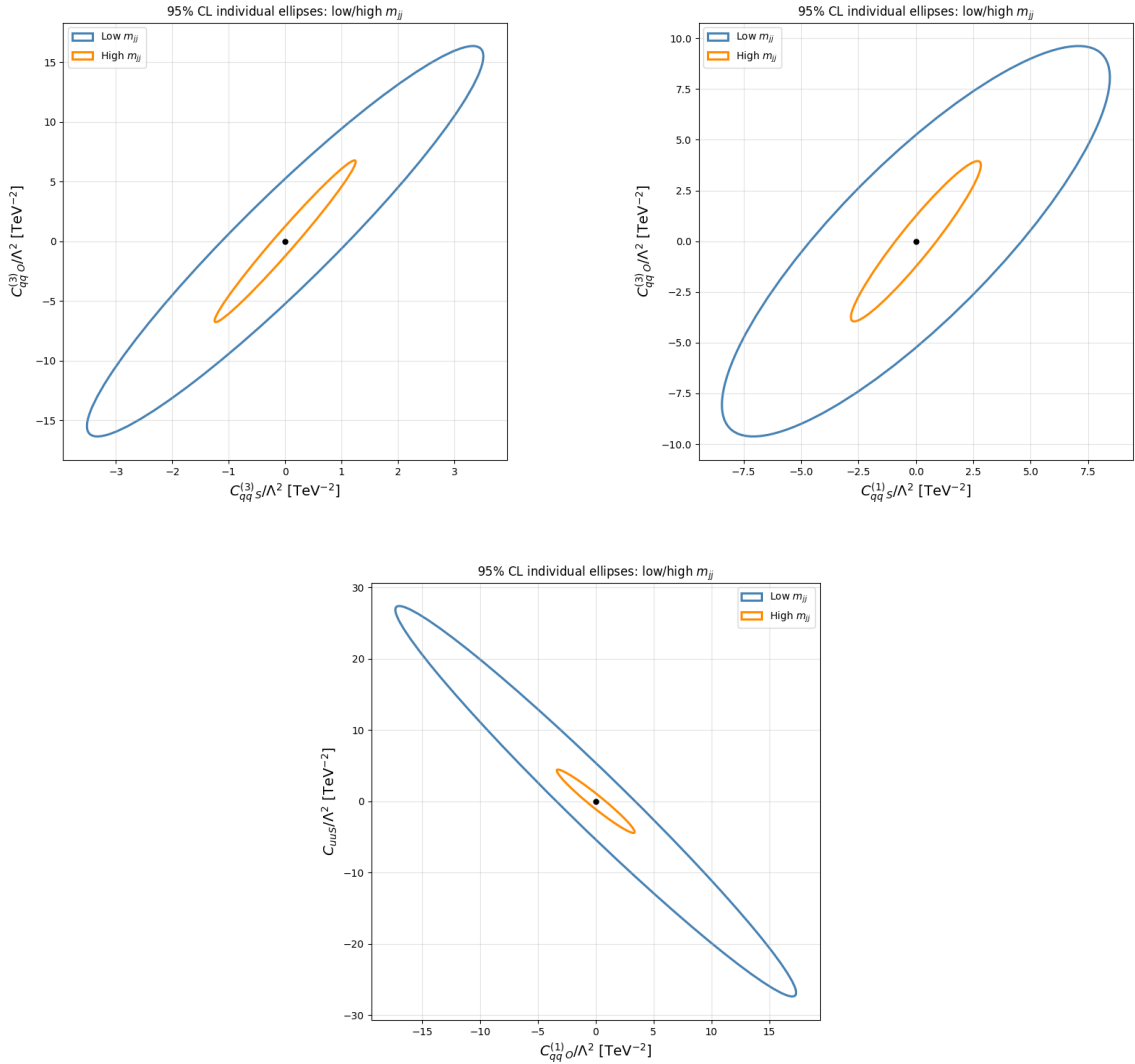


Table 7.2: Best constrained direction in coefficient space and its associated eigenvalue, for $\gamma + \text{jets}$ production by the difference in azimuthal angle between jets, $\Delta\phi_{jj}$.

	$\gamma + \text{jets}$	
	$m_{jj} < 800 \text{ GeV}$	$m_{jj} > 800 \text{ GeV}$
$\mathcal{O}_{qq}^{(1),1}$	0.21	0.23
$\mathcal{O}_{qq}^{(1),8}$	0.29	0.24
$\mathcal{O}_{qq}^{(3),1}$	0.89	0.90
$\mathcal{O}_{qq}^{(3),8}$	-0.18	-0.16
$\mathcal{O}_{uu}^1$	0.18	0.17
$\mathcal{O}_{uu}^8$	0.08	0.06
λ	5.94	136.74

Figure 7.4: Individual contours in the $C_{qq}^{(3),1}$ vs $C_{qq}^{(3),8}$ plane (*top left*), in the $C_{qq}^{(1),1}$ vs $C_{qq}^{(3),8}$ one (*top right*) and in the $C_{qq}^{(1),1}$ vs C_{uu}^1 one (*bottom middle*) for $\gamma + \text{jets}$ production by $\Delta\phi_{jj}$



Chapter 8

Conclusion and Outlook :

This thesis investigated the sensitivity of the process $pp \rightarrow \gamma jj$ to the ten four-light-quark operators of the Warsaw basis within the SMEFT. Working at leading order and parton level with events generated using MADGRAPH5_AMC@NLO, we derived constraints on the Wilson coefficients from the photon transverse momentum distribution and studied several observables to improve said bounds.

The first set of bounds was obtained at the cross-section level, where the allowed region in coefficient space already reveals a strong asymmetry : one direction is tightly constrained while the orthogonal one remains largely unconstrained, illustrating the limited discriminating power of the cross-section. Moving to the differential p_T^γ distribution and constructing a binned χ^2 function allowed us to significantly improve upon this. Among all operators considered, $\mathcal{O}_{qq}^{(3)}$ consistently receives the tightest individual constraint. This can be understood in part to the normalisation convention adopted for this operator : the use of Pauli matrices σ^I rather than $\sigma^I/2$ enhances its contributions relative to $\mathcal{O}_{qq}^{(1)}$ and $\mathcal{O}_{uu}$ by a numerical factor. The global analysis confirmed that the most constrained direction in the ten-dimensional coefficient space is predominantly aligned with $C_{qq}^{(3)}$, while several orthogonal directions remain essentially unconstrained.

To probe these directions, we applied the interference revival strategy, identifying combinations of operators for which the total interference cross section vanishes and constructing differential observables capable of separating the positive- and negative-weight contributions across phase space. Among the observables studied namely p_T^γ , $\Delta\phi_{jj}$, M_{jj} , $M_{\gamma jj}$ the transverse sphericity Sph_T , and the Zeppenfeld variable Z_γ . We saw that Z_γ , the azimuthal separation $\Delta\phi_{jj}$ and p_T^γ showed the largest revival power. Double-differential distributions in p_T^γ vs $\Delta\phi_{jj}$ further sharpened the individual bounds. Exploring additional observables, such as triple products of final-state momenta, could be a natural next step since these quantities are odd under certain discrete symmetries and may be sensitive to directions in coefficient space that standard kinematic variables cannot resolve.

Going beyond the linear level, we included the quadratic $\mathcal{O}(\Lambda^{-4})$ contributions for the three operators $\mathcal{O}_{qq}^{(1)}$, $\mathcal{O}_{qq}^{(3)}$, and $\mathcal{O}_{uu}$. In the high- p_T^γ tail, the squared dimension-six terms dominate over the interference by at least a factor of three, signalling that the linear approximation becomes unreliable in this kinematic regime. The resulting contours are asymmetric and strictly narrower than those from the linear analysis, though a consistent treatment at this order would in principle require including dimension-8 operators [40],

which lie beyond the scope of this work.

Finally, we separated the colour-singlet and colour-octet components of $\mathcal{O}_{qq}^{(1)}$, $\mathcal{O}_{qq}^{(3)}$, and $\mathcal{O}_{uu}$ using SU(3) Fierz identities and derived independent constraints on each. The two colour structures yield qualitatively similar bounds, with the high- m_{jj} region consistently providing the strongest sensitivity. The validity of the EFT expansion in that region should nevertheless be checked carefully before comparing these results to experimental data.

Several extensions of this work would be worth pursuing. The analysis was carried out at parton level without NLO corrections, parton shower, or detector simulation, so the bounds presented here should be regarded as idealised sensitivities. An immediate step would be to include those NLO corrections. Including a more realistic uncertainty model with actual PDF and scale uncertainties and confronting the results with ATLAS or CMS measurements would be the necessary steps towards a fully realistic analysis. The analysis developed in this work could furthermore be extended in several directions. First, the analysis could be applied to other processes: replacing the final-state photon by a W boson is a natural candidate, since the W couples exclusively to left-handed fermions and would therefore probe a complementary set of operator directions. Second, the revival strategy introduced here could be extended to a larger subspace of the full set of ten four-light-quark operators, rather than the restricted subset considered in this thesis.

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