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Constraining four-heavy-quark operators with top-quark, Higgs, and electroweak precision data

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Abstract: We establish constraints on the dimension-six four-heavy-quark operators in the Standard Model Effective Field Theory (SMEFT) by synthesising LHC measurements of top-quark and single-Higgs production with electroweak precision observables. We scrutinise the choice of the γ_5 scheme in single-Higgs calculations, demonstrating its non-negligible impact on SMEFT fits.

Keywords: SMEFT, higher-loop computations, Monte Carlo, collider physics

Contents

1 Introduction

The absence of direct evidence for new light particles beyond the Standard Model (SM) at the Large Hadron Collider (LHC) has motivated a campaign of indirect searches in the SM Effective Field Theory (SMEFT) framework. In SMEFT, new-physics effects are parametrised through a series of higher-dimensional operators modifying the interactions of the SM particles. Thus, SMEFT provides a systematic and model-agnostic way of probing new physics in the absence of new light states. A key strength of the SMEFT framework lies in its ability to correlate effects across different sectors of the SM interactions. In order to fully exploit its potential in identifying signatures of physics beyond the SM, global analyses and interpretations within the SMEFT paradigm have become essential, prompting significant ongoing efforts [1–7].

Such global interpretations are crucial, as they can reveal potential signs of new physics or, at the very least, place constraints on the energy scale at which new physics could appear by setting bounds on the Wilson Coefficients (WCs) of higher-dimensional operators. Additionally, global interpretations help identify sectors with greater potential for deviations from the SM by highlighting the least constrained operator classes. Notably, operators involving four heavy-quark fields stand out as among the least constrained by current experimental data. These induce contact interactions of four top quarks ($t\bar{t}t\bar{t}$), four bottom quarks ($b\bar{b}b\bar{b}$), and interactions involving a top-quark pair with a bottom-quark pair ($t\bar{t}b\bar{b}$).

Results from global fits indicate that the new-physics scale associated with this class of operators can be as low as a few hundred GeV; see, for example, the recent global analysis by the SMEFiT collaboration [5]. These loose constraints arise because, at tree level, these interactions are predominantly probed by $t\bar{t}t\bar{t}$ or $t\bar{t}b\bar{b}$ production, which suffer from large experimental uncertainties [8–10]. Moreover, these inclusive measurements are not sufficient to distinguish different colour and chirality structures in contact interactions, leading to flat directions that weaken the constraints.

These findings–together with model-building arguments suggesting that new physics might be top-philic [11–14]–have motivated indirect probes of the four-heavy contact operators via their higher-loop contributions to observables that are measured more precisely than multi-top-quark production. In particular, the effects of four-heavy-quark operators on electroweak precision observables (EWPO) [15–17], single-Higgs production in gluon fusion [18, 19], top-quark-pair production [20, 21] and flavour observables [17, 22] have been computed. These studies demonstrate that such indirect probes supply information complementary to four-top-quark production and must be included to obtain tighter limits on the strength of these interactions.

In this work we carry out a fit that combines direct and indirect probes, thereby exploiting their complementarity. Specifically, we include leading-order (LO) contributions to $t\bar{t}t\bar{t}$ and $t\bar{t}b\bar{b}$ production; next-to-leading-order (NLO) contributions to $t\bar{t}$ and $t\bar{t}H$ production; two-loop effects in gluon-fusion Higgs production ($gg \to H$) and Higgs decays, together with the two-loop contributions to EWPO.

The chiral nature of the four-fermion contact interactions in the SMEFT demands particular care in loop computations, because the Dirac algebra necessarily involves γ_5. Since γ_5 is intrinsically four-dimensional, one must define a consistent continuation to $d = 4 - 2\epsilon$ dimensions. In this work, we adopt two distinct continuation schemes: the *naïve dimensional regularisation* (NDR) scheme [23] and the *Breitenlohner–Maison–'t Hooft–Veltman* (BMHV) scheme [24, 25].

Provided that each scheme is implemented self-consistently, any differences in EFT matrix elements can be traced either to the scheme-dependent definition of the WCs [19] or to finite terms specific to the chosen renormalisation prescription [26]. This is due to the fact that NDR and BMHV predictions can be different in loop computations. We demonstrate this explicitly for the two-loop process of single-Higgs production and decay mediated by four-heavy-quark operators: in a dedicated fit, the bounds extracted for the WCs differ between the NDR and BMHV schemes. The scheme dependence has been studied also in the context of di-Higgs production at the LHC [27] and flavour physics [28–32].

The paper is organised as follows: in Section 2, we state our flavour assumptions, introduce the SMEFT operators relevant to this study, and outline the computational setup. Section 3 presents our predictions for the processes under consideration and summarises the analytic expressions for the four–top-quark operators that modify the ggH and $\gamma\gamma H$ couplings in both the NDR and BMHV γ_5 schemes. The scheme dependence of the resulting bounds on the WCs is examined in Section 4. Our fitting method and core results are detailed in Section 5. Finally, Section 6 summarises our findings.

2 Theoretical framework and computation setup

In this section, we discuss the SMEFT theoretical employed, along with the technical details underlying our computations.

2.1 SMEFT framework

A generic SMEFT Lagrangian, including terms up to $\mathcal{O}(\Lambda^{-4})$, can be written as

$$\mathcal{L}_{\text{SMEFT}} = \mathcal{L}_{\text{SM}} + \sum_i \frac{c_i^{(6)}\,\mathcal{O}_i^{(6)}}{\Lambda^2} + \sum_j \frac{c_j^{(8)}\,\mathcal{O}_j^{(8)}}{\Lambda^4} + \mathcal{O}\big(\Lambda^{-6}\big)\,, \tag{2.1}$$

where $c_i^{(D)}$ and $\mathcal{O}_i^{(D)}$ are the WCs and SMEFT operators of mass dimension D, respectively, and Λ denotes the scale of new physics. Restricting to the dimension-six operators, the SMEFT prediction for cross section can be parametrised as

$$\sigma_{\text{SMEFT}} \;=\; \sigma_{\text{SM}} \;+\; \sigma_{\text{int}}^{(i)}\,\frac{c_i}{\Lambda^2} \;+\; \sigma_{\text{quad}}^{(i)}\,\frac{c_i^2}{\Lambda^4} \;+\; \sigma_{\text{cross}}^{(i,j)}\,\frac{c_i\,c_j}{\Lambda^4}\,, \tag{2.2}$$

Here σ_{int} originates from the interference between the SM and dimension-six SMEFT amplitudes scaling as Λ^{-2}, while σ_{quad} and σ_{cross} denote the diagonal (c_i^2) and off-diagonal (c_ic_j) quadratic contributions scaling as Λ^{-4}. An analogous parametrisation will be adopted for the partial widths, Γ. In all our results we set $\Lambda = 1$ TeV.

We use a specific flavour assumption of the SMEFT focused on top-quark interactions:

$$U(3)_l \times U(3)_e \times U(2)_q \times U(2)_u \times U(3)_d \equiv U(2)^2 \times U(3)^3, \tag{2.3}$$

where the subscripts denote the five-fermion representations of the SM. This minimal relaxation of the $U(3)^5$ group allows for top-quark-chirality-flipping interactions, such as dipole interactions and modifications to the top-Yukawa coupling. We adhere to the notation and operator conventions of Refs. [20, 33] and focus on the four-heavy subclass of dimension-six four-fermion operators defined as follows:

$$\begin{aligned}
\mathcal{Q}_{qq}^{1(ijkl)} &= (\bar{q}_i\gamma^\mu q_j)(\bar{q}_k\gamma_\mu q_l), \quad & \mathcal{Q}_{qq}^{3(ijkl)} &= (\bar{q}_i\gamma^\mu\tau^I q_j)(\bar{q}_k\gamma_\mu\tau^I q_l),\\
\mathcal{Q}_{qu}^{1(ijkl)} &= (\bar{q}_i\gamma^\mu q_j)(\bar{u}_k\gamma_\mu u_l), \quad & \mathcal{Q}_{qu}^{8(ijkl)} &= (\bar{q}_i\gamma^\mu T^A q_j)(\bar{u}_k\gamma_\mu T^A u_l),\\
& & \mathcal{Q}_{uu}^{(ijkl)} &= (\bar{u}_i\gamma^\mu u_j)(\bar{u}_k\gamma_\mu u_l),
\end{aligned}$$

where the notation $\mathcal{Q}$ denotes operators written in the original Warsaw basis [34], with corresponding WCs denoted as C_i. However, in this work, we use operators aligned with the `dim6top` [33] and `SMEFTatNLO` [20] conventions, hereafter referred to as the 'top basis'—in which the operators are written as $\mathcal{O}_i$ with Wilson coefficients c_i. The translations of four-heavy coefficients at tree level from the Warsaw basis to the top basis are shown in Eq. (2.4). The corresponding state-of-the-art constraints are reported in the recent global fit of Ref. [5].

$$\begin{aligned} c^1_{QQ} &= 2C^{(1)}_{qq} - \frac{2}{3}C^{(3)}_{qq}, \quad & c^1_{Qt} &= C^{(1)}_{qu}, \\ c^1_{tt} &= C^{(1)}_{uu}, & c^8_{QQ} &= 8C^{(3)}_{qq}, \\ c^8_{Qt} &= C^{(8)}_{qu}. \end{aligned} \tag{2.4}$$

We note here that, in all our computations, we adopt the definition of the four-heavy operator in terms of Warsaw-basis operators, i.e. $\mathcal{O}^8_{QQ} = \mathcal{Q}^{(3)}_{qq}/8 + \mathcal{Q}^{(1)}_{qq}/24$, rather than $\mathcal{O}^8_{QQ} = (\bar{Q}\gamma^\mu T^A Q)(\bar{Q}\gamma_\mu T^A Q)$. The two expressions differ by an evanescent operator, as also discussed in Ref. [17]. Numerical results can differ between the two definitions when the evanescent operator contributes. All our results are consistent with the first definition, and we will comment on this further.

2.2 Computation setup

For all our predictions, we utilise `MadGraph5_aMC@NLO` [35] and the `SMEFT@NLO` [20] package, with the exceptions being $gg \to H$, for which we employ the analytic expression given in Ref. [19] and the EWPO, for which we use the expressions in Eq. (3.1) extracted from [16, 17] – see dedicated discussions below. The parton distribution functions (PDF) set `NNPDF3.1` in the five-flavour scheme at NLO with $\alpha_s(m_Z) = 0.118$ [36] is used as input for all Monte Carlo (MC) simulations through the `LHAPDF` interface [37]. All computations are carried out in the G_F scheme [38, 39], which is recommended for SMEFT analyses by the LHC EFT WG [40]. The Fermi constant value is set to $G_F = 1.16637 \times 10^{-5}\,\text{GeV}^{-2}$. The masses of the Higgs boson and the top quark are set to $m_H = 125\,\text{GeV}$ and $m_t = 172\,\text{GeV}$, respectively.

As will be relevant later, it is important to note that in `SMEFT@NLO` [20], in $d \neq 4$ dimensions, γ_5 is treated as anti-commuting, and the cyclic property of Dirac matrices traces is not maintained following the KKS scheme [41]. The latter is understood to be equivalent to the NDR scheme supplemented with a fixed reading point.

The factorisation and renormalisation scales, μ_F and μ_R, are set to half of the sum of the masses of the final state particles. The scale μ_{EFT} introduced in the counterterms of the WCs is set to μ_R to ensure $\overline{\text{MS}}$ renormalisation for the EFT poles [21]. The effects of renormalisation group equations (RGEs) on the WCs are not considered in this work, and we interpret the WCs at a typical electroweak (EW) scale. We refer to Refs. [42–45] for RGE effects in this context. Considering the EWPO at scale $Q \sim m_Z$ and four-top-quark production at $Q \sim 2m_t$, we adopt scales within a factor of four of each other and thus expect RGE effects to be under control.

3 Studied processes

In this section, we review the characteristic features of the LHC processes analysed in this work: four-top-quark production; top-quark pair production; top-quark pair production in association with a Higgs boson; single-Higgs production with its subsequent decay. Finally, we discuss the EWPO predictions employed in our analysis.

3.1 Four-top-quark production

Our $pp \to t\bar{t}t\bar{t}$ predictions are listed in Table 3 of Appendix B. As demonstrated in Ref. [46], subleading terms stemming from the interference of four-fermion operators with weakly mediated SM amplitudes play a non-negligible role; all such contributions are therefore included in our calculation. It is worth noting that, unlike $t\bar{t}$ production, the richer colour structure of the $pp \to t\bar{t}t\bar{t}$ process allows colour-singlet operators to interfere with the QCD SM amplitudes already at LO.

In the pure SM, subleading EW contributions almost exactly cancel among themselves [47] making the leading NLO QCD prediction highly reliable. In the EFT, however, such cancellation is not guaranteed because the SMEFT operators alter the kinematic structure of the amplitudes. A fully consistent NLO prediction in the SMEFT would thus require the complete set of NLO QCD *and* EW corrections—an undertaking that is presently beyond the reach of existing automated tools. Therefore, we employ only tree-level predictions for this process.

3.2 Top-quark pair production (and in association with a Higgs)

In $pp \to t\bar{t}(H)$, four-fermion operators that involve a third-generation doublet contribute at tree level through b-quark–initiated amplitudes. The interference of colour-singlet four-fermion operators with the SM vanishes at tree level when only purely QCD-induced amplitudes are considered [48], because the top quarks are always produced in a colour-octet configuration.[1] At NLO in QCD, however, both real and virtual corrections alter the colour flow and can induce a non-zero interference for colour-singlets.

Inclusive predictions for all processes as well as the differential predictions in the Higgs boson transverse momenta, p_T^H, for both the SM and the SMEFT are collected in Tables 2 and 3 of Appendix B.[2] The SMEFT results are separated into linear, $\mathcal{O}(\Lambda^{-2})$, and (off-) diagonal quadratic, $\mathcal{O}(\Lambda^{-4})$, terms. As previously mentioned, at LO, the operators under study contribute only through b-quark–initiated channels so the diagonal-quadratic piece from c_{tt}^1 vanishes. Because c_{tt}^1 first appears at one loop, obtaining its quadratic contribution would require squaring the loop amplitudes rendering it beyond the perturbative order considered in this work.

[1] For weak-mediated Born amplitudes, i.e. $t\bar{t}$ production via an s-channel weak boson, colour-singlets can already interfere at LO. These contributions are generally expected to be subdominant and are not considered here.

[2] Differential predictions in the top-quark-pair invariant mass, $m_{t\bar{t}}$, for $t\bar{t}$ production are omitted here for brevity, but can be provided upon request.

We observe that c^1_{Qt} dominates the SMEFT corrections at NLO, providing by far the largest linear contribution to the cross section. This is the case in both $t\bar{t}$ and $t\bar{t}H$ production. Moreover, the WC of the right-handed four-top-quark operator c^1_{tt} features some strong cancellation between the gluon- and quark-initiated channels in $t\bar{t}H$ production. This cancellation amplifies the scale dependence and results in sizeable QCD uncertainties for the linear contribution, as illustrated in Table 3. The same pattern is visible differentially in one of the p_T^H bins in Table 2. In $t\bar{t}$ production, this effect is absent in the inclusive rate.

Moving to the off-diagonal quadratic terms, at LO, interference between colour-singlet and colour-octet structures—whether between SM and EFT amplitudes or between two different EFT operators—vanishes exactly, and so only singlet–singlet and octet–octet combinations survive, as shown in Table 3. At NLO, real emissions or virtual gluon exchange can mix the colour flows, generating singlet–octet cross terms which are numerically tiny–cross terms consistent with zero at the 2σ level are therefore omitted from the tables for clarity.

Finally, in Fig. 1, the differential distributions of the top–quark-pair invariant mass, $m_{t\bar{t}}$, and the Higgs transverse momentum, p_T^H, are displayed in the left and right panels, respectively. For $\mathcal{O}(1)$ WCs, the four–heavy operators modify the SM prediction by at most

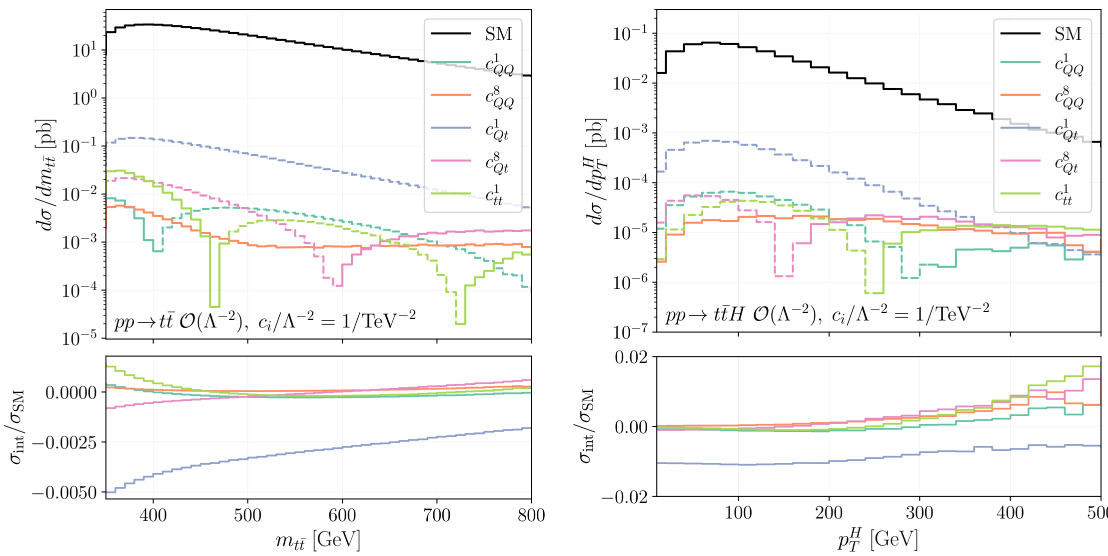


Figure 1: Differential distributions of the top-quark-pair invariant mass in $t\bar{t}$ (left) and of the Higgs transverse momentum in $t\bar{t}H$ (right). The curves show the linear SMEFT contributions of the five four-heavy operators, compared to the NLO SM prediction. The absolute values are plotted, and the dashed lines indicate where the interference becomes destructive—that is, where the contributions are negative.

the percent level. This mild impact is expected, given both the one-loop suppression and the small b-quark parton density driving these contributions. The dominant correction comes from $\mathcal{O}^1_{Qt}$: it enhances the low-invariant-mass region of the $t\bar{t}$ spectrum while inducing an essentially flat shift in the p_T^H distribution. All other operators yield sub-percent effects;

in some kinematic bins their contributions even change sign, as indicated by the dashed curves in the figures.

3.3 Electroweak precision observables

Our EWPO predictions are obtained at the two-loop level leveraging the work of Ref. [17] and reported here in Eq. (3.1). It is important to emphasise that all EWPO results were extracted in the Warsaw basis from Ref. [17] to maintain consistency with our `SMEFT@NLO` calculations, which also adopt this basis. We stress again that this choice is different by an evanescent contribution compared to the contribution of $\mathcal{O}^{8}_{QQ}$ defined with two colour-octet currents, since the mapping in Eq. (2.4) is a tree-level one (cf. Ref. [17]). No other coefficients are affected. We further note that this evanescent term can be numerically significant, and a substantial shift in the c^{8}_{QQ} contribution would be expected had the results in the top basis conventions of Ref. [17] been employed from the outset.

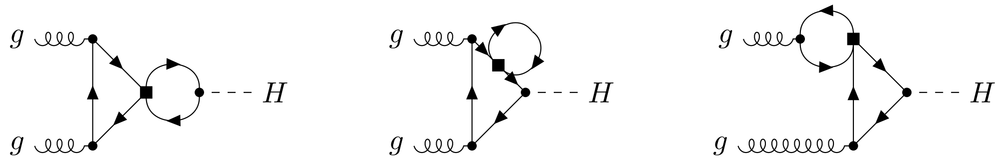
$$
\begin{aligned}
\delta\Gamma_Z^{b\bar{b}} &= 9.5412\times 10^{-4}\, c^{1}_{QQ} + 1.0098\times 10^{-4}\, c^{8}_{QQ} - 1.1409\times 10^{-3}\, c^{1}_{Qt} \\
&\quad + 4.4956\times 10^{-7}\, c^{8}_{Qt} - 3.12\times 10^{-6}\, c^{1}_{tt}, \\
\delta R_c &= -9.699\times 10^{-5}\, c^{1}_{QQ} - 1.0265\times 10^{-5}\, c^{8}_{QQ} + 1.1598\times 10^{-4}\, c^{1}_{Qt} \\
&\quad - 4.5701\times 10^{-8}\, c^{8}_{Qt} + 3.1718\times 10^{-7}\, c^{1}_{tt}, \\
\delta R_l &= 1.1688\times 10^{-2}\, c^{1}_{QQ} + 1.2371\times 10^{-3}\, c^{8}_{QQ} - 1.3977\times 10^{-2}\, c^{1}_{Qt} \\
&\quad + 5.5074\times 10^{-6}\, c^{8}_{Qt} - 3.8222\times 10^{-5}\, c^{1}_{tt}, \\
\delta R_b &= 4.4158\times 10^{-4}\, c^{1}_{QQ} + 4.6736\times 10^{-5}\, c^{8}_{QQ} - 5.2803\times 10^{-4}\, c^{1}_{Qt} \\
&\quad + 2.0806\times 10^{-7}\, c^{8}_{Qt} - 1.444\times 10^{-6}\, c^{1}_{tt}, \\
\delta A_b &= 2.4597\times 10^{-4}\, c^{1}_{QQ} + 3.2227\times 10^{-5}\, c^{8}_{QQ} - 2.9442\times 10^{-4}\, c^{1}_{Qt} \\
&\quad + 5.834\times 10^{-7}\, c^{8}_{Qt} + 5.7326\times 10^{-5}\, c^{1}_{tt}, \\
\delta A_{b,\mathrm{FB}} &= 2.5306\times 10^{-4}\, c^{1}_{QQ} + 8.3078\times 10^{-5}\, c^{8}_{QQ} - 3.0434\times 10^{-4}\, c^{1}_{Qt} \\
&\quad + 5.3495\times 10^{-6}\, c^{8}_{Qt} + 5.2565\times 10^{-4}\, c^{1}_{tt}\,.
\end{aligned}
\tag{3.1}
$$

The definitions of the observables, as well as details of the computation and numerical inputs, are provided in Appendix C. We note that (off-)diagonal quadratic EFT contributions are strongly suppressed relative to the linear contributions shown in Eq. (3.1). Moreover, the contributions of c^{8}_{Qt} and c^{1}_{tt} are purely two-loop induced, as they do not contribute to the EWPO at one loop.

Although the one-loop contributions are not shown here–see Appendix C for said contributions–in comparison, we find the two-loop corrections to be negligible for all observables except for the bottom-quark asymmetry, A_b, and its forward–backward counterpart, $A_{b,\mathrm{FB}}$, where they are significant.

3.4 Single-Higgs production in gluon-fusion and Higgs decays

Higgs production via gluon fusion and its decays into gluons and photons are loop–induced already at LO. Four–top-quark operators contribute for the first time through two-loop



(a) Higgs-top-quark coupling (b) Top-quark propagator (c) Gluon-top-quark vertex

Figure 2: Feynman diagrams illustrating the corrections to $gg \to H$ induced by four-top-quark SMEFT operator insertions.

diagrams, such as those shown in Fig. 2 for the production process.

For single-Higgs production, we adopt the results of Ref. [19]. A key challenge in these computations arise from the presence of γ_5 in the loop amplitudes rendering a delicate treatment necessary when dimensional regularisation is used. It has been shown in Ref. [19] that isolated contributions of four-top-quark operators depend on the continuation scheme of the γ_5 matrix to $d = 4 - 2\epsilon$ dimensions. The reference studied two schemes: the NDR [23] and the BMHV scheme [24, 25]. Whilst the former is algebraically inconsistent in the presence of traces involving six or more γ^μ matrices [41, 49–51], the latter remains consistent but the regulator spuriously breaks chiral symmetries and hence requires symmetry-restoring counterterms [52–58].

The continuation-scheme dependence is expected to cancel upon matching, once a process-specific set of operators at a given loop order is included.[3] We adopt the loop-order definitions of Refs. [59, 60], which requires the assumption of weakly interacting and renormalisable UV models. The two schemes can be connected by a set of relations, available in Ref. [19] for the operators entering single-Higgs production. Scheme-independent results can be achieved by including operators that enter at a lower loop order–see also Appendix A.

Here, we consider only the impact of four-top-quark operators, whose matrix element $\mathcal{M}$ reads

$$\mathcal{M}_{\mathrm{OS}}^{ggH} = \left(c_{Qt}^{1} + \left(c_F - \frac{c_A}{2}\right) c_{Qt}^{8}\right) \mathcal{K}_{tG} \frac{1}{\Lambda^2} \mathcal{M}_{tG}^{ggH} + \left(c_{Qt}^{1} + c_F c_{Qt}^{8}\right) \frac{1}{\Lambda^2} \left(B_{ggH} + \mathcal{K}_{t\varphi}\right) \mathcal{M}_{\mathrm{SM}}^{ggH}, \tag{3.2}$$

$$\mathcal{K}_{tG} = \begin{cases} \frac{\sqrt{2} m_t g_s}{16\pi^2 v} & \text{(NDR)} \\ 0 & \text{(BMHV)}, \end{cases} \qquad \mathcal{K}_{t\varphi} = \begin{cases} \frac{m_H^2 - 4m_t^2}{16\pi^2} & \text{(NDR)} \\ 0 & \text{(BMHV)}, \end{cases} \tag{3.3}$$

$$B_{ggH} = \frac{4m_t^2 - m_H^2}{8\pi^2} \left(\beta \log\left(\frac{\beta - 1}{\beta + 1}\right) + 2 + \log\left(\frac{\mu_R^2}{m_t^2}\right)\right), \quad \beta = \sqrt{1 - \frac{4m_t^2}{m_H^2}}. \tag{3.4}$$

The matrix element in Eq. (3.2) has been obtained in the on-shell (OS) renormalisation scheme for the top-quark mass, whilst the operators are renormalised in the $\overline{\mathrm{MS}}$ scheme. We note that this differs from the renormalisation scheme of Ref. [19]. The scheme dependence in Eq. (3.2) and in the following Eq. (3.5) is parametrised by the $\mathcal{K}$ terms in Eq. (3.3). We report the explicit expressions for $\mathcal{M}_{tG}^{ggH}$ and $\mathcal{M}_{\mathrm{SM}}^{ggH}$ in Appendix A.

[3] This applies to scenarios in which UV divergences are absent in the UV model, a feature that is guaranteed for the operators considered here by the superficial degree of divergence [26].

The matrix element for the Higgs-photon coupling ($\gamma\gamma H$) can be obtained by performing the replacements indicated in Appendix A on Eq. (3.2) yielding

$$\mathcal{M}_{\rm OS}^{\gamma\gamma H} = \left(c_{Qt}^{1} + c_F c_{Qt}^{8}\right) \mathcal{K}_{tG} \frac{Q_t e}{g_s} \frac{1}{\Lambda^2} \mathcal{M}_{t\gamma}^{\gamma\gamma H} + \left(c_{Qt}^{1} + c_F c_{Qt}^{8}\right) \frac{1}{\Lambda^2} \left(B_{ggH} + \mathcal{K}_{t\varphi}\right) \mathcal{M}_{\rm SM}^{\gamma\gamma H}. \quad (3.5)$$

We note that different combination of WCs entering in Eq. (3.2) compared to Eq. (3.5)–first term of the former. The phenomenological consequences of this observation are discussed in the following section.

4 Interpretation of SMEFT constraints in different γ_5 schemes

In this section, we perform a simplified (toy) fit to assess the impact of the γ_5 prescription on the bounds on WCs from single–Higgs production. This fit uses the Higgs signal strength and its associated data from Ref. [61]. We restrict our study to the dominant gluon–fusion channel and consider only total production rates, since measurements of the Higgs transverse-momentum spectrum are not yet available. We use the same numerical inputs as in Section 2.2 and we set $\mu_R = \mu_F = m_H/2$ for Higgs production and $\mu_R = m_H$ for the partial widths [62].

In Table 4 of Appendix B we list the numerical results for the single-Higgs production cross section, σ, and the Higgs partial width, Γ, computed in both γ_5 schemes using the formulae derived in the previous section. We omit the $\mathcal{O}(\Lambda^{-4})$ terms, as they enter at one loop order higher than the SM–$\mathcal{O}(\Lambda^{-2})$ interference.

The theoretical signal strengths, $\mu^{\rm Th}$, used in the fit are defined as follows:

$$\mu_X^{\rm Th} = \frac{\sigma_{\rm SMEFT}\,{\rm BR}(H\to X)_{\rm SMEFT}}{\sigma_{\rm SM}\,{\rm BR}(H\to X)_{\rm SM}}, \qquad X \equiv [\gamma\gamma,\ W^+W^-,\ ZZ,\ b\bar{b},\ \tau^+\tau^-,\ \mu^+\mu^-]. \quad (4.1)$$

In Eq. (4.1), the same K-factor is used to account for higher-order corrections both in the SM and linear, $\mathcal{O}\left(\Lambda^{-2}\right)$, EFT contributions to the production cross section and so it cancels out in the signal strengths.

Concerning the branching ratios, ${\rm BR}(H\to X)$, in Eq. (4.1), four-top-quark operators modify only the loop–induced partial widths $H\to gg$ and $H\to\gamma\gamma$–later denoted as Γ_{gg} and $\Gamma_{\gamma\gamma}$, respectively. Consequently the total Higgs width, $\Gamma^{\rm Tot}$, changes and it multiplies all branching ratios by a common factor. Every theoretical signal strength in our fit therefore receives this universal rescaling—except for $\mu_{\gamma\gamma}^{\rm Th}$. Said channel is additionally affected by the process-specific four-top-quark contribution given in Eq. (3.5). In particular, we have

$$\Gamma^{\rm Tot} = \Gamma_{\rm SM}^{\rm Tot} + K_{gg}\Gamma_{\rm int}^{gg} + K_{\gamma\gamma}\Gamma_{\rm int}^{\gamma\gamma}, \quad (4.2)$$

where $\Gamma_{\rm int}^{gg}$ and $\Gamma_{\rm int}^{\gamma\gamma}$ can be read from Table 4, whilst the $K-$factors K_{gg} and $K_{\gamma\gamma}$ are obtained as the ratio between the SM best estimates [63–76] and the SM values in Table 4. The following are the values we employ [62]: $\Gamma_{\rm SM}^{\rm Tot} = 4.088\,{\rm MeV}$, $K_{gg} = 1.707$, $K_{\gamma\gamma} = 0.913$.

We present here the theoretical signal strengths, expanded to linear order in the WCs, in the NDR scheme:

$$
\begin{aligned}
\mu^{\rm Th}_{\gamma\gamma} &= 1 - 0.0159 \times c^1_{Qt} - 0.00239 \times c^8_{Qt} \\
&\quad - 6.60 \times 10^{-5} (c^1_{Qt})^2 - 2.34 \times 10^{-5} (c^8_{Qt})^2 - 9.34 \times 10^{-5} c^1_{Qt} c^8_{Qt}, \\
\mu^{\rm Th}_{\rm Y} &= 1 - 0.0186 \times c^1_{Qt} - 0.00606 \times c^8_{Qt} \\
&\quad - 1.47 \times 10^{-5} (c^1_{Qt})^2 - 1.17 \times 10^{-6} (c^8_{Qt})^2 - 8.39 \times 10^{-6} c^1_{Qt} c^8_{Qt},
\end{aligned}
\tag{4.3}
$$

and the BMHV one:

$$
\begin{aligned}
\mu^{\rm Th}_{\gamma\gamma} &= 1 - 0.00451 \times c^1_{Qt} - 0.00601 \times c^8_{Qt} \\
&\quad - 2.53 \times 10^{-6} (c^1_{Qt})^2 - 4.50 \times 10^{-6} (c^8_{Qt})^2 - 6.75 \times 10^{-6} c^1_{Qt} c^8_{Qt}, \\
\mu^{\rm Th}_{\rm Y} &= 1 - 0.00491 \times c^1_{Qt} - 0.00655 \times c^8_{Qt} \\
&\quad - 5.66 \times 10^{-7} (c^1_{Qt})^2 - 1.01 \times 10^{-6} (c^8_{Qt})^2 - 1.51 \times 10^{-6} c^1_{Qt} c^8_{Qt},
\end{aligned}
\tag{4.4}
$$

where $\rm Y \equiv [\, W^+W^-,\ ZZ,\ b\bar{b},\ \tau^+\tau^-,\ \mu^+\mu^-]$.

In Fig. 3 we display the two-dimensional $\Delta\chi^2$ contours for the single-Higgs production fit in the $(c^1_{Qt},\, c^8_{Qt})$ plane. This result highlights the scheme dependence introduced by the choice of γ_5-continuation prescription (NDR vs. BMHV). The analysis is restricted to four-top-quark operators and the fit retains all linear, quadratic, and mixed (cross) terms in Eqs. (4.3) and (4.4).

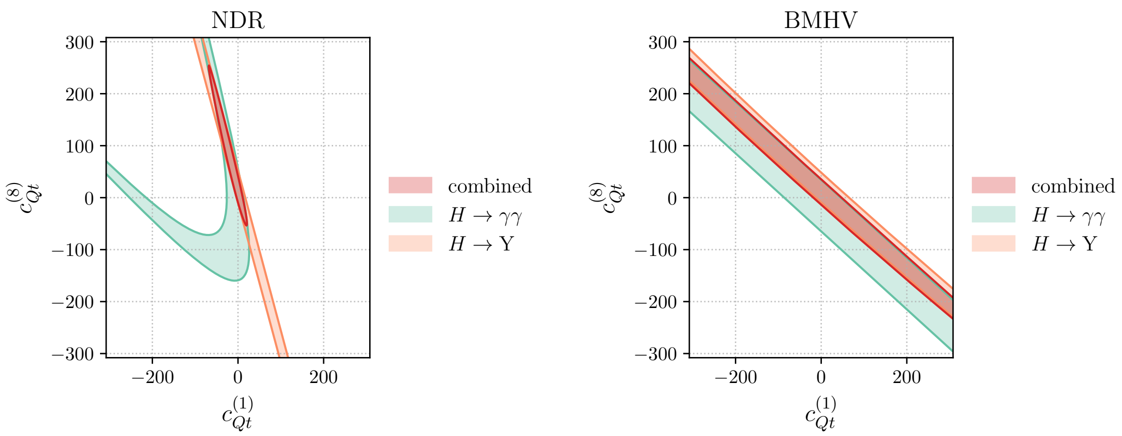


Figure 3: $\Delta\chi^2$ contours for the single-Higgs production fit in the $(c^1_{Qt},\, c^8_{Qt})$ plane, shown for each decay channel, i.e. Y and $\gamma\gamma$ and for their combination, labelled as 'combined'. The left (right) panel corresponds to the NDR (BMHV) γ_5 scheme. A pronounced flat direction emerges in the BMHV fit, whereas no such degeneracy appears in the NDR scheme.

As can be inferred from Fig. 3, the fit can distinguish between c^1_{Qt} and c^8_{Qt} in NDR but finds a flat direction for $c^1_{Qt} + c_F c^8_{Qt}$ in BHMV. This can be understood by inspecting Eq. (3.2): in the BMHV scheme, contributions of the type shown in Fig. 2c vanish, leaving a degeneracy in the WCs. Conversely, in the NDR scheme, they are non-vanishing and are proportional to the linear combination $c^1_{Qt} + \left(c_A - \frac{c_F}{2}\right) c^8_{Qt}$, which lifts the degeneracy.

As detailed in Appendix A, the WCs of operators entering at one-loop order depend on the scheme in such a way that they compensate for the scheme-dependent four-top-quark contribution. This observation highlights that a γ_5-prescription–independent result can be obtained only when a sufficiently complete set of SMEFT operators is included, as demonstrated in Ref. [44]. If those additional operators are omitted, our fit—restricted to four-quark operators—must be interpreted differently in the NDR and BMHV schemes: it effectively probes UV scenarios that generate four-top-quark operators while leaving all other operators absent.

5 Constraining four-quark operators: fit method, inputs and results

5.1 Fit method

We analyse the impact of SMEFT operators on measured inclusive and differential cross sections, $\sigma_{\rm Ex}$, by performing individual and marginalised χ^2 fits. For an operator coefficient c_i, the theoretical cross section, $\sigma_{\rm Th}$, in each bin, can be written as shown in Eq. (2.2)

When statistical and systematic uncertainties are provided separately by the experimental collaborations, the total experimental uncertainty, $\Delta_{\rm Ex}$, in each bin is determined by combining both uncertainties in a quadrature, i.e. assuming no correlation–the total experimental uncertainty provided directly by the collaborations is used when available. The normalisation, $\Delta_{\rm Tot}$, entering the test statistic is the quadrature sum of experimental and theoretical uncertainties, the latter being the QCD scale uncertainties of SM predictions. Conservatively, we choose that as the maximum of the scale uncertainty envelope.

For each bin, the χ^2 contribution is therefore calculated as

$$\chi^2_{\rm bin} = \left(\frac{\sigma_{\rm Ex} - \sigma_{\rm Th}}{\Delta_{\rm Tot}}\right)^2, \tag{5.1}$$

and the total χ^2 is obtained by summing over all bins, as well as over all considered observables and processes:

$$\chi^2 = \sum_{\rm proc.}\sum_{\rm obs.}\sum_{\rm bins}\chi^2_{\rm bin}. \tag{5.2}$$

Observables from different experiments are assumed to be uncorrelated. Finally, for the EWPO fit, we use the correlations between the different observables as quoted in Ref. [77].

5.2 Fit inputs

We discuss here the inputs for the processes included in our fit, which are summarised in Table 1.

$\boldsymbol{pp \to t\bar{t}}$ The SM predictions in the bins of ATLAS$_{t\bar{t}}$ are taken from the corresponding publication [79], where MC simulations were generated using `Powheg-Box` v2 [86–89] interfaced with `Pythia` 8.210 [90]. The cross-section normalisation of these MC samples is set to the NNLO + next-to-next-to-leading-logarithmic (NNLL) QCD prediction [91–97],

[4] https://www.hepdata.net/

Tag	$\sqrt{s}$, $\mathcal{L}$	Final state	Observable	$n_{\rm dat.}$	Ref.(Ex)	Location/HEPData	Ref.(Th)
			$pp \to t\bar{t}$				
CMS$_{t\bar{t}}$	13 TeV, 137 fb^{-1}	lepton+jets	$d\sigma/dm_{t\bar{t}}$	15	[78]	/Tab. 37	[78]
ATLAS$_{t\bar{t}}$	13 TeV, 36 fb^{-1}	lepton+jets	$d\sigma/dm_{t\bar{t}}$	9	[79]	/Tab. 617	[79]
			$pp \to t\bar{t}H$				
ATLAS$_{t\bar{t}H}$	13 TeV, 140 fb^{-1}	$H \to b\bar{b}$	$d\sigma/dp_T^H$	6	[80]	Fig. 3 in the ref./	[81]
			$pp \to t\bar{t}t\bar{t}$				
CMS$_{4t}$	13 TeV, 138 fb^{-1}	multi-leptons	$\sigma^{\rm tot.}_{t\bar{t}t\bar{t}}$	1	[9]	/Fig. 8	[82]
ATLAS$_{4t}$	13 TeV, 140 fb^{-1}	multi-leptons	$\sigma^{\rm tot.}_{t\bar{t}t\bar{t}}$	1	[8]	/Tab. 17	[82]
			$gg \to H$				
CMS$_{ggH}$	13 TeV, 138 fb^{-1}	$W^+W^-, ZZ, b\bar{b}, \tau^+\tau^-, \mu^+\mu^-$	$\mu^{\rm Ex}_{\rm Y}$	1	[61]	/Tab. 12	Eqs. (4.3) and (4.4)
CMS$_{ggH}$	13 TeV, 138 fb^{-1}	$\gamma\gamma$	$\mu^{\rm Ex}_{\gamma\gamma}$	1	[61]	/Tab. 12	Eqs. (4.3) and (4.4)
			$pp \to t\bar{t}b\bar{b}$				
ATLAS$_{t\bar{t}b\bar{b}}$	13 TeV, 36.1 fb^{-1}	lepton+jets	$\sigma^{\rm tot.}_{t\bar{t}b\bar{b}}$	1	[83]	/Tab. 5	as in [5]
CMS$^1_{t\bar{t}b\bar{b}}$	13 TeV, 35.9 fb^{-1}	all-jets	$\sigma^{\rm tot.}_{t\bar{t}b\bar{b}}$	1	[10]	Fig. 3 in the ref./	as in [5]
CMS$^2_{t\bar{t}b\bar{b}}$	13 TeV, 35.9 fb^{-1}	dilepton	$\sigma^{\rm tot.}_{t\bar{t}b\bar{b}}$	1	[84]	Tab. 4 in ref. (FPS)/	as in [5]
CMS$^3_{t\bar{t}b\bar{b}}$	13 TeV, 35.9, fb^{-1}	lepton+jets	$\sigma^{\rm tot.}_{t\bar{t}b\bar{b}}$	1	[84]	Tab. 4 in ref. (FPS)/	as in [5]
			EWPO				
EWPO^1	Z−pole, /	Z decays	$\Gamma_Z^{b\bar{b}}, R_c, R_l, R_b$	1	[77]	Tabs. 7.1 and 8.4 in ref.	[85]
EWPO^2	Z−pole, /	Z decays	$A_b, A_{b,\rm FB}$	1	[77]	Tab. 8.4 in ref.	[77]

Table 1: Summary of the inputs used in the fit. For each dataset we list, from left to right: (i) the dataset label; (ii) the centre-of-mass energy and integrated luminosity; (iii) the measured final state; (iv) the observable under study; (v) the number of data points; (vi) the experimental publication; (vii) the location of the experimental data (either in the publication or via its HEPData entry[4]); and (viii) the reference used for theoretical predictions.

as indicated explicitly in Table 1 of the publication. Predictions and measurements from ATLAS$_{t\bar{t}}$ are normalised to the total cross section of 832^{+20}_{-29} pb, as reported therein.

The SM predictions in the bins of CMS$_{t\bar{t}}$ are also taken from the corresponding publication [78]. The cross-section normalisation of these MC samples is set to the NNLO in the strong coupling constant including the resummation of NNLL soft-gluon terms calculated with TOP++ (version 2.0) [91]. This amounts to a normalisation factor, i.e. the inclusive $t\bar{t}$ production cross section, of 832^{+40}_{-46} pb. The SM theoretical uncertainties are taken from the respective publications.

$\boldsymbol{pp \to t\bar{t}H}$ We use the SM predictions and the associated uncertainties provided in the analysis of Ref. [81]. We use the experimental data from the most recent measurement of Ref. [80] reporting a measured total cross section of $411^{+24\%}_{-22\%}$ fb. The SM differential predictions [62, 98–103] extracted from Ref. [81] are in good agreement with our own results presented in Table 2 of Appendix B.

$\boldsymbol{pp \to t\bar{t}t\bar{t}}$ For both datasets, we adopt the SM cross section of Ref. [82], computed at NLO (QCD+EW)+NLL accuracy, yielding $13.37^{+7.77\%}_{-13.3\%}$ fb.

$\boldsymbol{pp \to t\bar{t}b\bar{b}}$ We extract the experimental measurements and theoretical predictions—for *both* the SM–assigned a 10% theoretical uncertainty– and the SMEFT—from Ref. [5].[5]

[5] The SMEFiT database can be found here: `https://github.com/LHCfitNikhef/smefit_database/`

$\boldsymbol{gg \to H}$ The analysis presented in Ref. [61] reports signal-strength modifiers, $\mu_i^{\rm Ex}$, categorised through their decay modes i, with uncertainties representing the total experimental error, combining both systematic and statistical contributions. These results are summarised as follows:

$$\begin{aligned} \mu_{\mu\mu}^{\rm Ex} &= 0.33^{+0.74}_{-0.70}\,, \quad \mu_{\tau\tau}^{\rm Ex} = 0.66^{+0.21}_{-0.21}\,, \quad \mu_{ZZ}^{\rm Ex} = 0.93^{+0.14}_{-0.13}\,, \\ \mu_{WW}^{\rm Ex} &= 0.90^{+0.11}_{-0.10}\,, \quad \mu_{bb}^{\rm Ex} = 5.31^{+2.97}_{-2.54}\,, \quad \mu_{\gamma\gamma}^{\rm Ex} = 1.08^{+0.12}_{-0.11}\,. \end{aligned} \tag{5.3}$$

The results reported in Eq. (5.3) are of the gluon-fusion production mode and constitute the experimental input for our fit. Given that gluon fusion is directly sensitive to the operators we consider and is the dominant production mode for single-Higgs production, we do not expect significant sensitivity from other production modes.We set the SM prediction to unity.

EWPO Our predictions are obtained at the two-loop level leveraging the work of Ref. [17] as discussed in Section 3.3 and in Appendix C. Experimental measurements and correlations are taken from Ref. [77]. SM predictions for all observables apart from A_b and $A_{b,\rm FB}$ are extracted from Table 2 of Ref. [85]. SM predictions of A_b and $A_{b,\rm FB}$ are taken directly from Table 8.4 in Ref. [77].

5.3 Fit results

Figure 4 shows the two-dimensional $\Delta\chi^2$ contours at 95% CL for the case where only linear EFT terms are included whilst Fig. 5 corresponds to the scenario where quadratic contributions are also taken into account. In both panels, we display the contours for the combined fit (all processes) under two distinct scenarios: one in which only the two WCs of interest are varied (black solid-line contour, labelled as 'comb-2D' in the plots) with all other coefficients fixed to zero, and one in which those two coefficients are scanned while the remaining WCs are profiled (black dashed-line contour, labelled as 'comb-profiled' in the plots). The best-fit point (BFP) is indicated in each case. Finally, in Table 5 of Appendix B we list all the 95% CL bounds obtained on each of the five WCs. We note that all $gg \to H$ results are the ones obtained using the NDR scheme and thus consistent with the NLO $t\bar{t}(H)$ results from `SMEFT@NLO`.

We observe that, in the linear case of Fig. 4, the comb-2D fit yields smaller and differently shaped contours than those obtained from the individual channels. This illustrates that the inclusion of additional production modes provides complementary information and strengthens the overall constraints. By contrast, the comb-profiled contours display significantly weaker constraining power in comparison to the comb-2D ones, underscoring the impact of the profiled coefficients. Examining at individual datasets, we observe that most are plagued by flat directions; four-top-quark production, for example, shows flat directions for all coefficient pairs we consider. Top-quark pair production in association with a Higgs and $t\bar{t}b\bar{b}$ seem to exhibit the least constraining power among the different processes. Inclusive Higgs production probes only a subset of the coefficients and also has no significant impact on the final combination.

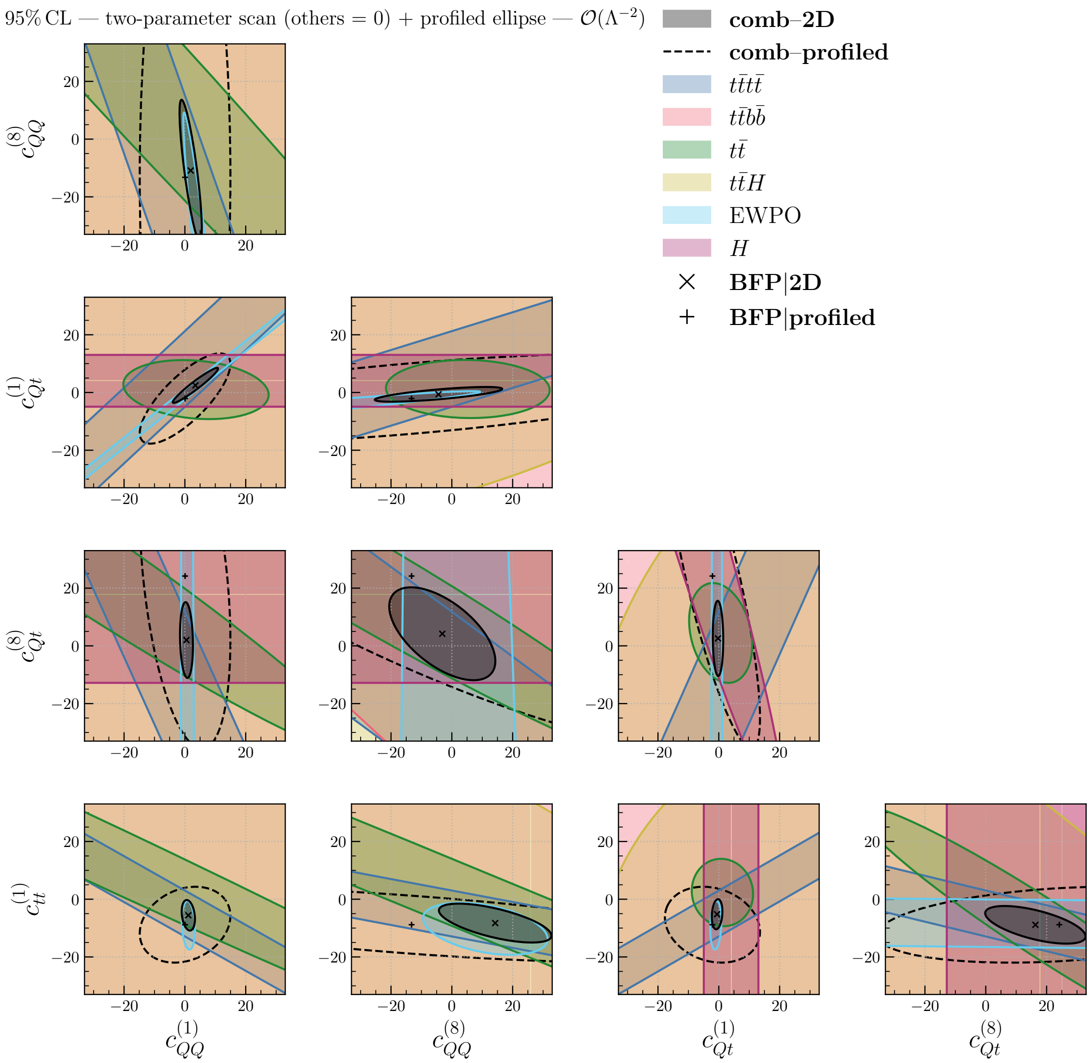


Figure 4: Two-dimensional fits for the four–heavy operator coefficients. Shown are the constraints from each set of observables separately and their combination. Only linear terms, $\mathcal{O}(\Lambda^{-2})$, in the EFT parametrisation are included. The best-fit point (BFP) for the combined fit is indicated for both the two-parameter scan and the profiled scan.

Furthermore, we observe a sizeable impact of the two-loop corrections in our EWPO fit—particularly when compared with the corresponding one-loop contours of Ref. [21]. These effects are driven mainly by the non-negligible two-loop contributions to the asymmetry observables A_b and $A_{b,\mathrm{FB}}$; see Appendix C for a detailed comparison between the one- and two-loop results for these quantities. We find that the combination of EWPO, top-quark pair and four-top-quark production significantly reduces the allowed parameter space. For some pairs we note that the contours extracted from the EWPO do not include the SM at the 95% CL. This is related to the well-known discrepancy between the measurement of $A_{b,\mathrm{FB}}$ and its SM prediction [77].

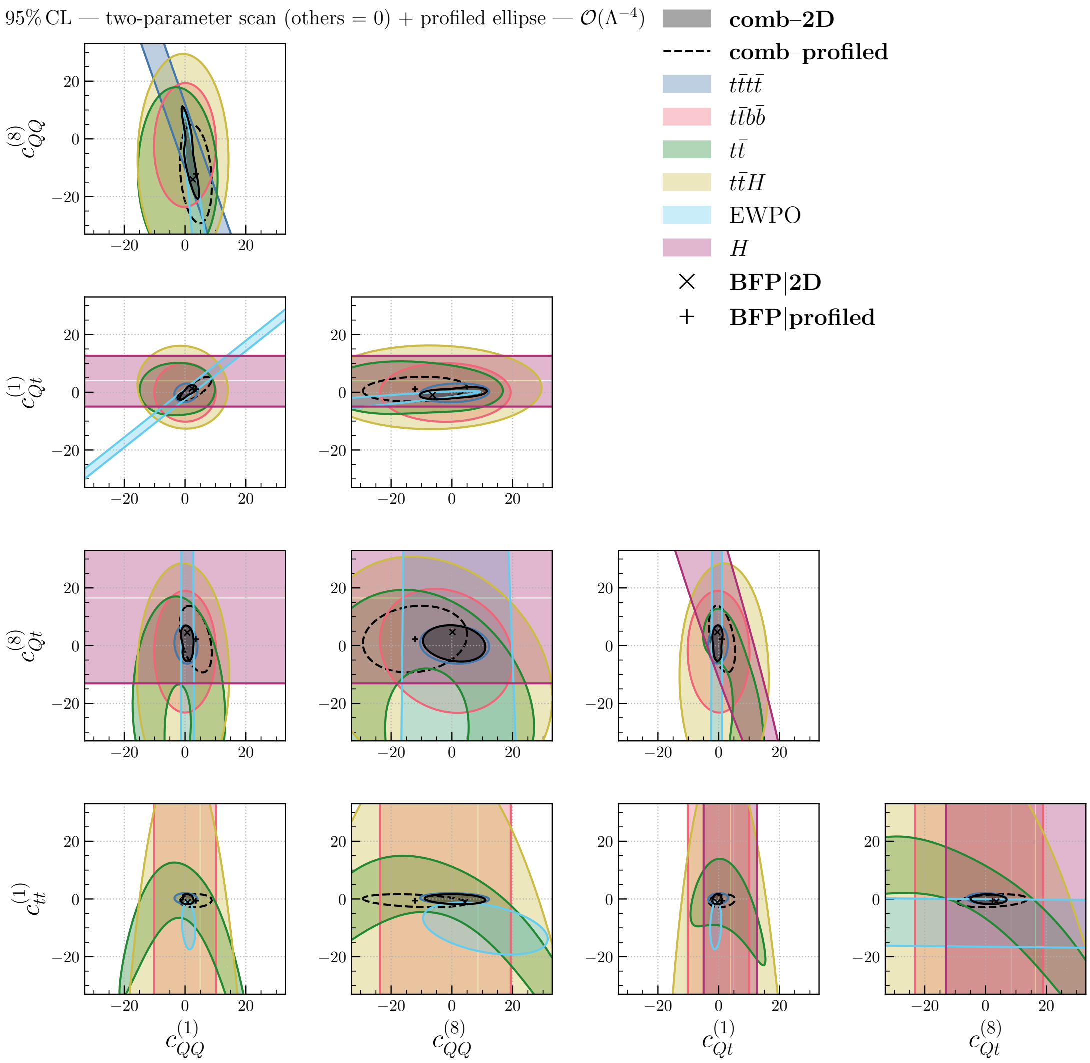


Figure 5: Same as Fig. 4, but including quadratic terms, $\mathcal{O}(\Lambda^{-4})$, in the EFT parametrisation.

At the quadratic level of Fig. 5, the combined fit yields markedly tighter bounds, driven predominantly by the four–top-quark measurements. The comb-2D contour lies very close to the four–top-quark one, with modest additional tightening from the EWPO. Moreover, the discrepancy between the comb-profiled and comb-2D contours is less pronounced than in the linear case. The most significant reduction in the allowed parameter space appears in the ($c_{QQ}^{(8)}$–$c_{QQ}^{(1)}$) plane, where the four–top-quark channel alone exhibits an almost blind direction; here the EWPO combination is essential to lift the degeneracy.

6 Conclusion

In this work we explored how various classes of observables can shed light on the dimension-six four-heavy-quark operators. These operators are notoriously difficult to constrain:

global EFT fits typically leave them essentially undetermined at the linear level. To overcome this, several studies have examined loop-induced effects in specific observables, which offer complementary sensitivity to these otherwise elusive interactions.

In particular, we have explored the tree-level contributions of the four-heavy operators to four-top-quark production at the LHC; their one-loop contributions to top-quark pair production and to associated Higgs production–both inclusive and differentially; the two-loop contributions to single-Higgs production in gluon fusion and its subsequent decays; and the one- and two-loop contributions to EWPO. We presented results for all these channels at linear and, where available, quadratic terms in the EFT expansion.

In the case of single-Higgs production via gluon fusion and its subsequent decays, we scrutinised the dependence of the results on the choice of the prescription for the continuation of γ_5 to $d = 4 - 2\epsilon$. We show that this scheme dependence numerically propagates into the bounds on the WCs. Therefore, one must exercise caution when interpreting these limits. Indeed, by restricting the analysis to four-top-quark operators, this assumption leads to a different interpretation of the fit in the BMHV scheme compared with NDR. Establishing a coherent correspondence between the two schemes therefore necessitates the inclusion of additional operators.

Finally, we performed a fit combining experimental data from the LHC and LEP, incorporating all relevant theoretical and experimental uncertainties. This allowed us to identify the most sensitive observables and to elucidate the complementarity between tree- and loop-level probes. In the linear fit, the synergy between heavy-quark production and EWPO is key to reducing the allowed parameter space. In the quadratic fit, the strongest constraints arise from four-top-quark production, with the EWPO further reducing the allowed parameter space and lifting degeneracies where present. Moreover, we observed non-negligible two-loop effects in the EWPO fit, predominantly driven by the sizeable two-loop corrections to the asymmetry observables.

Finally, we note that a comprehensive assessment of the impact of the observables considered in this work requires their inclusion in a global fit together with the complete set of relevant operators. It is also worthwhile to investigate UV-complete theories which, upon matching onto the EFT, reproduce the operator basis examined here.

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A Matrix elements for ggH and $\gamma\gamma H$

In this appendix we provide the explicit expression for the auxiliary quantities appearing in Eqs. (3.2) and (3.5). We also include the contribution from the operators

$$\begin{aligned} \mathcal{O}_{t\varphi} \equiv \mathcal{Q}_{u\varphi}^{(33)} = (\bar{q}_3\tilde{\varphi}u_3)\left(\varphi^\dagger\varphi\right), \quad \mathcal{O}_{tG} \equiv \mathcal{Q}_{uG}^{(33)} = \bar{q}_3\tilde{\varphi}\sigma^{\mu\nu}T^A u_3 G^A_{\mu\nu}, \\ \mathcal{O}_{\varphi G} \equiv \mathcal{Q}_{\varphi G} = G^{A,\mu\nu}G^A_{\mu\nu}\left(\varphi^\dagger\varphi\right), \end{aligned} \tag{A.1}$$

which are required to obtain a result that is independent of the γ_5 continuation scheme. We note that the operator $\mathcal{O}_{\varphi G}$ is introduced to renormalise the one-loop contribution from $\mathcal{O}_{tG}$. We note that the results presented here employ the $\overline{\text{MS}}$ renormalisation scheme for the WCs and the on-shell scheme for the top-quark mass, in contrast to Ref. [19], where all parameters were renormalised in the MS scheme.

The matrix element for the process $gg \to H$ (or, equivalently, $H \to gg$) reads

$$\begin{aligned} \mathcal{M}_{\text{OS}}^{ggH} =& \frac{c_{\varphi G}}{\Lambda^2}\mathcal{M}_{\varphi G}^{ggH} + \left[c_{tG} + \left(c_{Qt}^1 + \left(c_F - \frac{c_A}{2}\right)c_{Qt}^8\right)\mathcal{K}_{tG}\right]\frac{1}{\Lambda^2}\mathcal{M}_{tG}^{ggH} \\ &+ \left[1 + \left(c_{Qt}^1 + c_F c_{Qt}^8\right)\frac{1}{\Lambda^2}\left(B_{ggH} + \mathcal{K}_{t\varphi}\right) - \frac{v^3}{\sqrt{2}m_t\Lambda^2}c_{t\varphi}\right]\mathcal{M}_{\text{SM}}^{ggH}, \end{aligned} \tag{A.2}$$

where c_A and c_F are the SU(N) Casimir invariants in the adjoint and fundamental representations, respectively. For SU(3)$_\text{C}$ they take the values $c_A = 3$ and $c_F = 4/3$. The factor B_{ggH} is defined in Eq. (3.4).

The matrix elements entering Eq. (A.2) are

$$\begin{aligned} \mathcal{M}_{tG}^{ggH} =& -T_F\frac{g_s m_t}{\sqrt{2}\pi^2}L^{\mu_1\mu_2}\epsilon_{\mu_1}(p_1)\epsilon_{\mu_2}(p_2)\delta^{A_1A_2}\times \\ &\left(\frac{m_t^2}{m_H^2}\log^2\left(\frac{\beta-1}{\beta+1}\right) + \sqrt{1-\frac{4m_t^2}{m_H^2}}\log\left(\frac{\beta-1}{\beta+1}\right) + 2\log\left(\frac{\mu_R^2}{m_t^2}\right) + 1\right), \end{aligned} \tag{A.3}$$

$$\mathcal{M}_{\text{SM}}^{ggH} = T_F\frac{g_s^2 m_t^2}{4\pi^2 v m_H^2}L^{\mu_1\mu_2}\epsilon_{\mu_1}(p_1)\epsilon_{\mu_2}(p_2)\delta^{A_1A_2}\times\left(\frac{m_H^2-4m_t^2}{m_H^2}\log^2\left(\frac{\beta-1}{\beta+1}\right) - 4\right), \tag{A.4}$$

$$\mathcal{M}_{\varphi G}^{ggH} = -4ivL^{\mu_1\mu_2}\epsilon_{\mu_1}(p_1)\epsilon_{\mu_2}(p_2)\delta^{A_1A_2}. \tag{A.5}$$

Here T_F is the Dynkin index of the fundamental representation of SU(N). The Lorentz structure of the amplitude is

$$L^{\mu_1\mu_2} = (m_H^2/2\; g^{\mu_1\mu_2} - p_1^{\mu_2}p_2^{\mu_1}), \tag{A.6}$$

with p_1 and p_2 being the gluon momenta. The indices A_1 and A_2 denote the gluon colour indices.

The scheme dependence arising from the two-loop contributions of four-top-quark operators, parametrised by $\mathcal{K}_{tG}$, and $\mathcal{K}_{t\varphi}$ in Eq. (3.3), can be compensated by assuming the WCs of operators entering at one loop order are scheme dependent. In particular, the relations

$$c_{t\varphi}^{\rm NDR} = c_{t\varphi}^{\rm BMHV} + \left(c_{Qt}^{1} + c_F c_{Qt}^{8}\right) \frac{y_t(\lambda - y_t^2)}{8\pi^2}, \tag{A.7}$$

$$c_{tG}^{\rm NDR} = c_{tG}^{\rm BMHV} - \left(c_{Qt}^{1} + \left(c_F - \frac{c_A}{2}\right) c_{Qt}^{8}\right) \frac{y_t g_s}{16\pi^2}, \tag{A.8}$$

render the prediction in Eq. (A.2) scheme independent. In these expressions, y_t is the top-quark Yukawa coupling and $\lambda = m_H^2/(2v^2)$. We note that some of the shifts depend on the strong coupling constant, g_s, and thus on the renormalisation scale. This scale dependence must be accounted for when a dynamical scale choice is employed in the calculation.

We now present the matrix element for the Higgs-photon coupling. We use $F^{\mu\nu}$ to denote the photon field strength tensor and introduce the operators $\mathcal{O}_{t\gamma} = (\bar{t}_L \sigma^{\mu\nu} t_R) \frac{H+v}{\sqrt{2}} F_{\mu\nu}$ and $\mathcal{O}_{\varphi\gamma} = H v F^{\mu\nu} F_{\mu\nu}$. These operators are not part of the Warsaw basis, as they are defined directly in terms of the physical fields in the broken phase. Their expression in Warsaw-basis operators can be found in Refs. [104, 105]. We finally obtain

$$\begin{aligned} \mathcal{M}_{\rm OS}^{\gamma\gamma H} =& \mathcal{M}_{{\rm SM},W}^{\gamma\gamma H} + \frac{c_{\varphi\gamma}}{\Lambda^2}\mathcal{M}_{\varphi\gamma} + \left[c_{t\gamma} + \left(c_{Qt}^{1} + c_F c_{Qt}^{8}\right) \mathcal{K}_{tG} \frac{Q_t e}{g_s}\right] \frac{1}{\Lambda^2} \mathcal{M}_{t\gamma}^{\gamma\gamma H} \\ &+ \left[1 + \left(c_{Qt}^{1} + c_F c_{Qt}^{8}\right) \frac{1}{\Lambda^2} \left(B_{ggH} + \mathcal{K}_{t\varphi}\right) - \frac{v^3}{\sqrt{2} m_t \Lambda^2} c_{t\varphi}\right] \mathcal{M}_{\rm SM}^{\gamma\gamma H}. \end{aligned} \tag{A.9}$$

The one-loop matrix element for $H \to \gamma\gamma$ can be obtained from that for $H \to gg$ by making the substitutions $g_s \to eQ_t$ and $T_F \delta^{A_1 A_2} \to N_C$ in $\mathcal{M}_{tG}^{ggH}$, $\mathcal{M}_{\rm SM}^{ggH}$, where e is the electric charge of the electron and $Q_t = 2/3$ is the quantised charge of the top quark. We denote the matrix elements of the Higgs boson decay into photons as $\mathcal{M}_{t\gamma}^{\gamma\gamma H}$, $\mathcal{M}_{\rm SM}^{\gamma\gamma H}$. The tree-level insertion of $\mathcal{O}_{\varphi\gamma}$ is given by $\mathcal{M}_{\varphi\gamma}^{\gamma\gamma H} = -4ivL^{\mu_1\mu_2}\epsilon_{\mu_1}(p_1)\epsilon_{\mu_2}(p_2)$. Regarding the two-loop matrix elements, each four-top-quark operator insertion generates two different colour contractions. We are able to obtain our result from the Higgs-gluon coupling since in the diagrams in Fig. 2, the only non-vanishing term features a single Dirac trace, allowing the colour structures to be identified unambiguously.

For completeness, the SM contribution to $H \to \gamma\gamma$ with W-boson loops reads

$$\mathcal{M}_{{\rm SM},W}^{\gamma\gamma H} = \frac{e^2}{4\pi^2 v}\left(\frac{6m_W^2}{m_H^2} + \left(\frac{6m_W^4}{m_H^4} - \frac{3m_W^2}{m_H^2}\right)\log^2\left(\frac{\beta_W - 1}{\beta_W + 1}\right) + 1\right), \tag{A.10}$$

where $\beta_W = \sqrt{1 - \frac{4m_W^2}{m_H^2}}$.

To render the Higgs decay into photons scheme–independent, we must employ Eq. (A.7) with the relation analogous to Eq. (A.8), namely

$$c_{t\gamma}^{\rm NDR} = c_{t\gamma}^{\rm BMHV} - eQ_t \frac{y_t\left(c_{Qt}^{1} + c_F c_{Qt}^{8}\right)}{8\pi^2}. \tag{A.11}$$

B Numerical predictions

Table 2 presents the differential p_T^H spectrum in the ATLAS $t\bar{t}H$ bins, whilst Table 3 lists the inclusive cross sections for all top-quark processes. SMEFT results are split into interference, quadratic, and cross terms, with total-rate K-factors given. All WCs are set to unity and $\Lambda = 1$ TeV. Predictions are quoted within their QCD scale uncertainties and MC statistical errors–predictions which are compatible with zero within a MC error of 2σ or greater are omitted.

Table 4 lists our numerical predictions for single–Higgs production and for Higgs decays into gluons and photons. The quoted SM value is the LO result with only top- and W-boson loops included. Quadratic SMEFT terms are strongly suppressed relative to the linear ones and are therefore not shown.

	$d\sigma_{\mathrm{NLO}}/dp_T^H\,[\mathrm{pb}]$					
	$p_T^H < 60\,\mathrm{GeV}$	$60 \leq p_T^H < 120\,\mathrm{GeV}$	$120 \leq p_T^H < 200\,\mathrm{GeV}$	$200 \leq p_T^H < 300\,\mathrm{GeV}$	$300 \leq p_T^H < 450\,\mathrm{GeV}$	$p_T^H \geq 450\,\mathrm{GeV}$
SM	$1.197\mathrm{e}{-1} \pm 0.054\%^{+6\%}_{-9.08\%}$	$1.785\mathrm{e}{-1} \pm 0.052\%^{+5.94\%}_{-9.17\%}$	$1.258\mathrm{e}{-1} \pm 0.065\%^{+5.97\%}_{-9.43\%}$	$5.203\mathrm{e}{-2} \pm 0.096\%^{+5.47\%}_{-9.61\%}$	$1.888\mathrm{e}{-2} \pm 0.141\%^{+3.36\%}_{-9.24\%}$	$5.179\mathrm{e}{-3} \pm 0.28\%^{+1.76\%}_{-6.8\%}$
	$\mathcal{O}(c_i/\Lambda^2)$					
c_{tt}^1	$-3.693\mathrm{e}{-5} \pm 1.878\%^{+55.11\%}_{-34.6\%}$	$-1.173\mathrm{e}{-4} \pm 0.867\%^{+49.39\%}_{-31.81\%}$	$-1.223\mathrm{e}{-4} \pm 1.11\%^{+56.16\%}_{-35.46\%}$	$1.134\mathrm{e}{-6} \pm 139\%^{+1632\%}_{-2106\%}$	$9.791\mathrm{e}{-5} \pm 1.872\%^{+21.88\%}_{-18.05\%}$	$1.567\mathrm{e}{-4} \pm 1.387\%^{+31.86\%}_{-22.53\%}$
c_{QQ}^1	$-9.997\mathrm{e}{-5} \pm 1.227\%^{+36.9\%}_{-25.25\%}$	$-1.911\mathrm{e}{-4} \pm 0.96\%^{+38.41\%}_{-26.06\%}$	$-1.63\mathrm{e}{-4} \pm 1.896\%^{+42.95\%}_{-28.45\%}$	$-4.519\mathrm{e}{-5} \pm 6.025\%^{+64.81\%}_{-39.8\%}$	$3.056\mathrm{e}{-5} \pm 9.49\%^{+32.56\%}_{-28.18\%}$	$7.125\mathrm{e}{-5} \pm 5.559\%^{+31.13\%}_{-22.19\%}$
c_{QQ}^8	$2.735\mathrm{e}{-5} \pm 3.993\%^{+11.41\%}_{-14.46\%}$	$5.53\mathrm{e}{-5} \pm 3.214\%^{+9.25\%}_{-13.25\%}$	$8.329\mathrm{e}{-5} \pm 2.527\%^{+13.33\%}_{-13.98\%}$	$8.667\mathrm{e}{-5} \pm 2.701\%^{+15.39\%}_{-12.49\%}$	$8.205\mathrm{e}{-5} \pm 3.134\%^{+24.33\%}_{-17.51\%}$	$7.729\mathrm{e}{-5} \pm 4.074\%^{+32.33\%}_{-22.3\%}$
c_{Qt}^1	$-1.262\mathrm{e}{-3} \pm 0.246\%^{+31.97\%}_{-22.67\%}$	$-1.93\mathrm{e}{-3} \pm 0.203\%^{+32.61\%}_{-22.98\%}$	$-1.338\mathrm{e}{-3} \pm 0.258\%^{+33.74\%}_{-23.52\%}$	$-4.837\mathrm{e}{-4} \pm 0.553\%^{+34.96\%}_{-24.05\%}$	$-1.277\mathrm{e}{-4} \pm 1.555\%^{+34.69\%}_{-23.78\%}$	$-3.17\mathrm{e}{-5} \pm 6.399\%^{+31.59\%}_{-22.27\%}$
c_{Qt}^8	$-1.149\mathrm{e}{-4} \pm 1.008\%^{+39.12\%}_{-28.69\%}$	$-1.266\mathrm{e}{-4} \pm 1.521\%^{+45.7\%}_{-33.33\%}$	$3.364\mathrm{e}{-6} \pm 74.484\%^{+389\%}_{-535\%}$	$9.672\mathrm{e}{-5} \pm 2.479\%^{+20.53\%}_{-14.5\%}$	$1.129\mathrm{e}{-4} \pm 2.405\%^{+28.3\%}_{-19.84\%}$	$1.088\mathrm{e}{-4} \pm 3.605\%^{+33.57\%}_{-23.02\%}$
	$\mathcal{O}(c_i^2/\Lambda^4)$					
c_{tt}^1	$\times$	$\times$	$\times$	$\times$	$\times$	$\times$
c_{QQ}^1	$2.89\mathrm{e}{-5} \pm 2.048\%^{+4.14\%}_{-5.08\%}$	$5.913\mathrm{e}{-5} \pm 1.867\%^{+4.61\%}_{-4.27\%}$	$6.764\mathrm{e}{-5} \pm 1.804\%^{+4.94\%}_{-3.84\%}$	$5.253\mathrm{e}{-5} \pm 2.259\%^{+6.82\%}_{-4.48\%}$	$3.687\mathrm{e}{-5} \pm 3.338\%^{+8.42\%}_{-5.56\%}$	$2.779\mathrm{e}{-5} \pm 4.385\%^{+8.37\%}_{-5.23\%}$
c_{QQ}^8	$4.415\mathrm{e}{-6} \pm 3.202\%^{+8.53\%}_{-5.52\%}$	$8.902\mathrm{e}{-6} \pm 2.504\%^{+7.99\%}_{-4.62\%}$	$1.034\mathrm{e}{-5} \pm 2.034\%^{+9.56\%}_{-5.3\%}$	$8.285\mathrm{e}{-6} \pm 2.89\%^{+10.59\%}_{-5.96\%}$	$5.831\mathrm{e}{-6} \pm 4.934\%^{+12\%}_{-6.83\%}$	$4.937\mathrm{e}{-6} \pm 6.78\%^{+14.48\%}_{-8.38\%}$
c_{Qt}^1	$2.833\mathrm{e}{-5} \pm 1.938\%^{+4.06\%}_{-5.02\%}$	$6.055\mathrm{e}{-5} \pm 1.45\%^{+4.17\%}_{-5.06\%}$	$6.545\mathrm{e}{-5} \pm 1.927\%^{+5.11\%}_{-3.51\%}$	$5.113\mathrm{e}{-5} \pm 2.355\%^{+6.15\%}_{-3.97\%}$	$3.639\mathrm{e}{-5} \pm 3.024\%^{+7.48\%}_{-4.79\%}$	$3.042\mathrm{e}{-5} \pm 4.135\%^{+12.52\%}_{-8.65\%}$
c_{Qt}^8	$4.303\mathrm{e}{-6} \pm 3.461\%^{+8.26\%}_{-5.78\%}$	$8.2\mathrm{e}{-6} \pm 2.807\%^{+8.72\%}_{-6.28\%}$	$1.01\mathrm{e}{-5} \pm 2.563\%^{+9.22\%}_{-5.09\%}$	$7.497\mathrm{e}{-6} \pm 7.655\%^{+11.93\%}_{-7.69\%}$	$4.598\mathrm{e}{-6} \pm 11.933\%^{+15.1\%}_{-13.59\%}$	$4.561\mathrm{e}{-6} \pm 8.879\%^{+12.56\%}_{-6.43\%}$
	$\mathcal{O}(c_i c_j/\Lambda^4)$					
$c_{QQ}^1 c_{Qt}^1$	$1.49\mathrm{e}{-5} \pm 8.241\%^{+8.56\%}_{-6.91\%}$	$2.817\mathrm{e}{-5} \pm 7.635\%^{+9.7\%}_{-8.02\%}$	$2.928\mathrm{e}{-5} \pm 9.17\%^{+10.63\%}_{-6.63\%}$	$1.762\mathrm{e}{-5} \pm 14.089\%^{+10.58\%}_{-6.63\%}$	$1.086\mathrm{e}{-5} \pm 21.595\%^{+8.73\%}_{-8.76\%}$	$-2.153\mathrm{e}{-6} \pm 118\%^{+62.83\%}_{-68.99\%}$
$c_{QQ}^8 c_{Qt}^8$	$3.253\mathrm{e}{-6} \pm 11.559\%^{+2.76\%}_{-5.77\%}$	$5.878\mathrm{e}{-6} \pm 10.101\%^{+10.34\%}_{-13.75\%}$	$7.997\mathrm{e}{-6} \pm 7.513\%^{+4.73\%}_{-6.71\%}$	$4.398\mathrm{e}{-6} \pm 19.689\%^{+9.78\%}_{-12.84\%}$	$4.704\mathrm{e}{-6} \pm 18.524\%^{+5.68\%}_{-1.91\%}$	$6.259\mathrm{e}{-7} \pm 123\%^{+109\%}_{-122\%}$

Table 2: Differential p_T^H predictions in the SM and SMEFT in the $t\bar{t}H$ process.

	$\sigma^{\text{proc.}}_{\text{(N)LO}}$[pb]						
	$\sigma^{t\bar{t}}_{\text{LO}}$	$\sigma^{t\bar{t}}_{\text{NLO}}$	$K_{t\bar{t}}$	$\sigma^{t\bar{t}H}_{\text{LO}}$	$\sigma^{t\bar{t}H}_{\text{NLO}}$	$K_{t\bar{t}H}$	$\sigma^{t\bar{t}t\bar{t}}_{\text{LO}}$
SM	$5.028\text{e}{+2} \pm 0.003\%^{+29.67\%}_{-21.44\%}$	$7.532\text{e}{+2} \pm 0.004\%^{+11.75\%}_{-11.9\%}$	1.498	$4.005\text{e}{-1} \pm 0.005\%^{+31.45\%}_{-22.25\%}$	$5.003\text{e}{-1} \pm 0.009\%^{+5.71\%}_{-9.24\%}$	1.249	$6.754\text{e}{-3} \pm 0.016\%^{+62.66\%}_{-35.48\%}$
				$\mathcal{O}(c_i/\Lambda^2)$			
c^1_{tt}	×	$2.376\text{e}{-1} \pm 0.461\%^{+24.03\%}_{-18.41\%}$	×	×	$-2.089\text{e}{-5} \pm 17.867\%^{+494\%}_{-339\%}$	×	$-1.185\text{e}{-3} \pm 0.026\%^{+36.07\%}_{-24.93\%}$
c^1_{QQ}	×	$-3.323\text{e}{-2} \pm 1.731\%^{+53.94\%}_{-34.92\%}$	×	×	$-3.975\text{e}{-4} \pm 1.688\%^{+46.88\%}_{-30.48\%}$	×	$-7.053\text{e}{-4} \pm 0.024\%^{+37.35\%}_{-25.48\%}$
c^8_{QQ}	$5.961\text{e}{-2} \pm 0.123\%^{+27.97\%}_{-25.65\%}$	$1.209\text{e}{-1} \pm 0.406\%^{+10.94\%}_{-12.13\%}$	2.028	$1.708\text{e}{-4} \pm 0.393\%^{+22.3\%}_{-21.22\%}$	$4.119\text{e}{-4} \pm 1.3\%^{+12.4\%}_{-11.31\%}$	2.411	$-2.351\text{e}{-4} \pm 0.027\%^{+37.39\%}_{-25.50\%}$
c^1_{Qt}	×	$-2.538 \pm 0.053\%^{+29.22\%}_{-21.29\%}$	×	×	$-5.175\text{e}{-3} \pm 0.138\%^{+33.01\%}_{-23.16\%}$	×	$7.191\text{e}{-4} \pm 0.024\%^{+41.42\%}_{-27.19\%}$
c^8_{Qt}	$5.951\text{e}{-2} \pm 0.135\%^{+27.97\%}_{-25.65\%}$	$-1.116\text{e}{-1} \pm 0.48\%^{+37.42\%}_{-30.59\%}$	1.875	$1.702\text{e}{-4} \pm 0.464\%^{+22.28\%}_{-21.21\%}$	$8.029\text{e}{-5} \pm 7.662\%^{+33.34\%}_{-40.45\%}$	0.471	$-3.026\text{e}{-4} \pm 0.029\%^{+29.98\%}_{-22.35\%}$
				$\mathcal{O}(c_i^2/\Lambda^4)$			
c^1_{tt}	×	×	×	×	×	×	$4.342\text{e}{-3} \pm 0.021\%^{+46.06\%}_{-29.23\%}$
c^1_{QQ}	$2.739\text{e}{-2} \pm 0.233\%^{+13.52\%}_{-16.73\%}$	$5.863\text{e}{-2} \pm 0.378\%^{+6.51\%}_{-7.68\%}$	2.14	$1.608\text{e}{-4} \pm 0.686\%^{+9.34\%}_{-12.45\%}$	$2.728\text{e}{-4} \pm 0.84\%^{+5.9\%}_{-3.77\%}$	1.696	$1.085\text{e}{-3} \pm 0.029\%^{+46\%}_{-29.19\%}$
c^8_{QQ}	$6.061\text{e}{-3} \pm 0.3\%^{+13.53\%}_{-16.74\%}$	$1.029\text{e}{-2} \pm 0.365\%^{+6.82\%}_{-4.65\%}$	1.697	$3.579\text{e}{-5} \pm 0.826\%^{+9.33\%}_{-12.44\%}$	$4.271\text{e}{-5} \pm 1.275\%^{+10.23\%}_{-5.71\%}$	1.193	$1.208\text{e}{-4} \pm 0.045\%^{+45.98\%}_{-29.18\%}$
c^1_{Qt}	$2.73\text{e}{-2} \pm 0.266\%^{+13.54\%}_{-16.75\%}$	$5.877\text{e}{-2} \pm 0.378\%^{+6.68\%}_{-7.78\%}$	2.152	$1.616\text{e}{-4} \pm 0.614\%^{+9.31\%}_{-12.42\%}$	$2.722\text{e}{-4} \pm 0.789\%^{+5.96\%}_{-3.8\%}$	1.684	$1.47\text{e}{-3} \pm 0.048\%^{+46.21\%}_{-29.28\%}$
c^8_{Qt}	$6.047\text{e}{-3} \pm 0.32\%^{+13.53\%}_{-16.74\%}$	$9.186\text{e}{-3} \pm 0.77\%^{+9.22\%}_{-6.16\%}$	1.519	$3.597\text{e}{-5} \pm 0.982\%^{+9.29\%}_{-12.41\%}$	$3.926\text{e}{-5} \pm 2.326\%^{+10.61\%}_{-6.52\%}$	1.091	$3.545\text{e}{-4} \pm 0.055\%^{+45.89\%}_{-29.15\%}$
				$\mathcal{O}(c_ic_j/\Lambda^4)$			
$c^1_{QQ}c^1_{Qt}$	$1.41\text{e}{-2} \pm 0.778\%^{+15.62\%}_{-18.47\%}$	$1.128\text{e}{-2} \pm 3.178\%^{+11.8\%}_{-10.75\%}$	0.8	$1.585\text{e}{-4} \pm 1.231\%^{+10.1\%}_{-13.13\%}$	$9.87\text{e}{-5} \pm 4.431\%^{+11.21\%}_{-7.92\%}$	0.622	$-3.51\text{e}{-4} \pm 0.261\%^{+43.39\%}_{-28.12\%}$
$c^1_{QQ}c^1_{tt}$	–	–	×	–	–	×	$3.566\text{e}{-4} \pm 0.365\%^{+41.6\%}_{-27.33\%}$
$c^1_{QQ}c^8_{QQ}$	–	$1.066\text{e}{-3} \pm 23.367\%^{+24.19\%}_{-23.49\%}$	×	–	–	×	$7.241\text{e}{-4} \pm 0.072\%^{+46.07\%}_{-29.23\%}$
$c^1_{QQ}c^8_{Qt}$	–	–	×	–	–	×	$1.314\text{e}{-4} \pm 0.381\%^{+42.77\%}_{-27.84\%}$
$c^1_{Qt}c^1_{tt}$	–	$5.545\text{e}{-4} \pm 45.484\%^{+14.2\%}_{-16.01\%}$	×	–	–	×	$-7.075\text{e}{-4} \pm 0.202\%^{+43.31\%}_{-28.07\%}$
$c^1_{Qt}c^8_{QQ}$	–	–	×	–	–	×	$-1.18\text{e}{-4} \pm 0.723\%^{+43.31\%}_{-28.06\%}$
$c^1_{Qt}c^8_{Qt}$	–	$-2.956\text{e}{-3} \pm 9.905\%^{+28.22\%}_{-26.3\%}$	×	–	$-1.689\text{e}{-5} \pm 19.597\%^{+23.96\%}_{-23.82\%}$	×	$-8.382\text{e}{-5} \pm 1.078\%^{+41.67\%}_{-27.38\%}$
$c^1_{tt}c^8_{QQ}$	–	–	×	–	$1.931\text{e}{-6} \pm 37.788\%^{+26.67\%}_{-25.98\%}$	×	$1.181\text{e}{-4} \pm 1.002\%^{+41.56\%}_{-27.31\%}$
$c^1_{tt}c^8_{Qt}$	–	$4.615\text{e}{-4} \pm 16.652\%^{+30.32\%}_{-27.57\%}$	×	–	–	×	$2.603\text{e}{-4} \pm 0.461\%^{+42.62\%}_{-27.77\%}$
$c^8_{QQ}c^8_{Qt}$	$3.177\text{e}{-3} \pm 0.965\%^{+15.56\%}_{-18.42\%}$	$3.907\text{e}{-3} \pm 2.285\%^{+2.25\%}_{-3.99\%}$	1.229	$3.45\text{e}{-5} \pm 2.034\%^{+10.22\%}_{-13.24\%}$	$2.685\text{e}{-5} \pm 5.113\%^{+7.94\%}_{-10.45\%}$	0.778	$4.391\text{e}{-5} \pm 0.584\%^{+42.65\%}_{-27.77\%}$

Table 3: Inclusive predictions for top-quark processes in the SM and SMEFT.

	σ_{int} [pb]
$\mathcal{O}(c_i/\Lambda^2)$	
c^1_{Qt} (NDR)	-0.3203
c^8_{Qt} (NDR)	-0.1033
c^1_{Qt} (BMHV)	-0.0830
c^8_{Qt} (BMHV)	-0.1106
σ_{SM}	16.51 pb

(a) $gg \to H$ cross section

	Γ_{int} [MeV]
$\mathcal{O}(c_i/\Lambda^2)$	
c^1_{Qt} (NDR)	-1.912×10^{-3}
c^8_{Qt} (NDR)	-4.810×10^{-4}
c^1_{Qt} (BMHV)	-2.781×10^{-4}
c^8_{Qt} (BMHV)	-3.708×10^{-4}
Γ_{SM}	0.1960 MeV

(b) $H \to gg$ partial width

	Γ_{int} [MeV]
$\mathcal{O}(c_i/\Lambda^2)$	
c^1_{Qt} (NDR)	2.798×10^{-5}
c^8_{Qt} (NDR)	3.731×10^{-5}
c^1_{Qt} (BMHV)	4.069×10^{-6}
c^8_{Qt} (BMHV)	5.426×10^{-6}
Γ_{SM}	1.016×10^{-2} MeV

(c) $H \to \gamma\gamma$ partial width

Table 4: Linear EFT contributions and SM values for single-Higgs inclusive production cross section in gluon-fusion channel (Table 4a) and Higgs decays (Tables 4b and 4c), parametrised as in Eq. (2.2). WCs are set to unity with $\Lambda = 1\,\text{TeV}$. The SM values correspond to the LO one-loop result including only the dominant top- and W-loops. Results are presented in both the NDR and the BMHV schemes.

C EWPO

We employ the relations in Eq. (C.1) where the WCs C_{HD} and C_{HWB} are expressed in terms of the shifts to the oblique parameters S and T [16] and in terms of the shifts of the Weinberg angle, s_θ, and the effective couplings [85]. The two-loop contributions to the oblique parameters ΔS and ΔT have been computed in and are taken from Ref. [17].

$$\begin{aligned} C_{HD} &= -\frac{2\,\alpha\,\Delta T}{v^2}, \quad C_{HWB} = \frac{\alpha\,\Delta S}{4\,c_\theta\,s_\theta\,v^2}, \\ \delta s_\theta^2 &= \frac{m_W^2\,C_{HD}}{2\sqrt{2}\,G_F\,m_Z^2} + \frac{m_W\,C_{HWB}}{\sqrt{2}\,G_F\,m_Z}\sqrt{1-\frac{m_W^2}{m_Z^2}}, \quad \delta g_Z = -\frac{C_{HD}}{4\sqrt{2}\,G_F}. \end{aligned} \tag{C.1}$$

Using Eq. (C.1) and substituting into the shifts in vector and axial-vector couplings shown in Eq. (C.2)—which can be found in Ref. [85]—we obtain the modified vector and axial–vector couplings due to the effective operators:

$$\begin{aligned} \delta g_V^f &= \delta g_Z\, g_V^f + Q^f\,\delta s_\theta^2\,, \\ \delta g_A^f &= \delta g_Z\, g_A^f\,, \end{aligned} \tag{C.2}$$

where $g_V^f = T_3/2 - Q^f s_\theta^2$ and $g_A^f = T_3/2$, where T_3 is weak isospin and Q^f is the electric charge. We adopt the conventions of Ref. [85] and use the following definitions of the EWPO:

$$\begin{aligned} \Gamma_i &= \frac{\sqrt{2}\,G_F\,m_Z^3\,N_c}{3\pi}\left(|g_V^i|^2 + |g_A^i|^2\right), \\ \Gamma_{\text{had}} &= \sum_{q=u,d,c,s,b}\Gamma_q, \qquad R_c = \frac{\Gamma_c}{\Gamma_{\text{had}}}, \quad R_b = \frac{\Gamma_b}{\Gamma_{\text{had}}}, \quad R_\ell = \frac{\Gamma_{\text{had}}}{\Gamma_\ell}. \end{aligned} \tag{C.3}$$

For the numerical analysis we adopt the following input parameters:

$$\begin{aligned}
G_F &= 1.166379\times 10^{-5}\ \text{GeV}^{-2}, & m_W &= 80.379\ \text{GeV}, & m_Z &= 91.1876\ \text{GeV},\\
v &= 246.22\ \text{GeV}, & \alpha &= 1/132.184, & s_\theta^2 &= 0.2230,\\
c_\theta^2 &= 1-s_\theta^2, & m_t &= 172.5\ \text{GeV}, & \Lambda &= 1\ \text{TeV},\\
\mu_R &= m_Z, & \Gamma_Z^{\text{SM}} &= 2.4941\ \text{GeV}, & \Gamma_{\text{had}}^{\text{SM}} &= 1.6944\ \text{GeV},\\
R_b^{\text{SM}} &= 0.21582, & A_b^{\text{SM}} &= 0.9347, & A_{b,\text{FB}}^{\text{SM}} &= 0.1029\,.
\end{aligned} \tag{C.4}$$

Corrections to Γ_b

Adopting the one-loop result of Ref. [17] and including the two-loop contributions as described above, the relative shift in the $Z\to b\bar{b}$ partial width, $\delta\Gamma_b$, reads

$$\begin{aligned}
\delta\Gamma_b^{1\text{L}}/\Gamma_Z^{\text{SM}} &= (3.8320\, c_{QQ}^1 + 0.4065\, c_{QQ}^8 - 4.5839\, c_{Qt}^1)\times 10^{-4}\\
\delta\Gamma_b^{1\text{L}+2\text{L}}/\Gamma_Z^{\text{SM}} &= (3.8255\, c_{QQ}^1 + 0.4049\, c_{QQ}^8 - 4.5745\, c_{Qt}^1 + 0.0018\, c_{Qt}^8 - 0.0125\, c_{tt}^1)\times 10^{-4},
\end{aligned} \tag{C.5}$$

where $\delta\Gamma_b^{1\text{L}}$ and $\delta\Gamma_b^{1\text{L}+2\text{L}}$ denote the one-loop and the combined one- and two-loop contributions, respectively.

Corrections to R_c, R_ℓ, and R_b

Implementing the above definitions and expanding to first order in $\delta\Gamma_b$, we obtain

$$\bar{R}_c = \frac{\Gamma_c^{\text{SM}}}{\Gamma_{\text{had}}^{\text{SM}}+\delta\Gamma_b} \simeq R_c^{\text{SM}}\left(1-\frac{\delta\Gamma_b}{\Gamma_{\text{had}}^{\text{SM}}}\right), \quad \bar{R}_\ell = \frac{\Gamma_{\text{had}}^{\text{SM}}+\delta\Gamma_b}{\Gamma_\ell^{\text{SM}}} \simeq R_\ell^{\text{SM}}\left(1+\frac{\delta\Gamma_b}{\Gamma_{\text{had}}^{\text{SM}}}\right),$$

$$\bar{R}_b = \frac{\Gamma_b^{\text{SM}}+\delta\Gamma_b}{\Gamma_{\text{had}}^{\text{SM}}+\delta\Gamma_b} \simeq R_b^{\text{SM}} + \left(1-R_b^{\text{SM}}\right)\frac{\delta\Gamma_b}{\Gamma_{\text{had}}^{\text{SM}}}\,. \tag{C.6}$$

Replacing $\delta\Gamma_b \to \delta\Gamma_b^{1\text{L}}$ or $\delta\Gamma_b^{1\text{L}+2\text{L}}$ gives

$$\begin{aligned}
\delta R_c^{1\text{L}}/R_c^{\text{SM}} &= (-5.6404\, c_{QQ}^1 - 0.5983\, c_{QQ}^8 + 6.7472\, c_{Qt}^1)\times 10^{-4}\\
\delta R_c^{1\text{L}+2\text{L}}/R_c^{\text{SM}} &= (-5.6308\, c_{QQ}^1 - 0.5959\, c_{QQ}^8 + 6.7333\, c_{Qt}^1 - 0.0026\, c_{Qt}^8 + 0.0184\, c_{tt}^1)\times 10^{-4},\\
\delta R_b^{1\text{L}}/R_b^{\text{SM}} &= (20.494\, c_{QQ}^1 + 2.1742\, c_{QQ}^8 - 24.516\, c_{Qt}^1)\times 10^{-4}\\
\delta R_b^{1\text{L}+2\text{L}}/R_b^{\text{SM}} &= (20.459\, c_{QQ}^1 + 2.1654\, c_{QQ}^8 - 24.465\, c_{Qt}^1 + 0.0096\, c_{Qt}^8 - 0.0669\, c_{tt}^1)\times 10^{-4}.
\end{aligned} \tag{C.7}$$

Corrections to A_b, $A_{b,\text{FB}}$

Using the one-loop result of Ref. [17], the one-loop shifts read

$$\begin{aligned}
\delta A_b^{1\text{L}}/A_b^{\text{SM}} &= (2.3648\, c_{QQ}^1 + 0.2508\, c_{QQ}^8 - 2.8288\, c_{Qt}^1)\times 10^{-4}\\
\delta A_{b,\text{FB}}^{1\text{L}}/A_{b,\text{FB}}^{\text{SM}} &= (2.3682\, c_{QQ}^1 + 0.2512\, c_{QQ}^8 - 2.8329\, c_{Qt}^1)\times 10^{-4}
\end{aligned} \tag{C.8}$$

where we have used the relation $A_{b,\mathrm{FB}} = {}^3\!/_4 A_b A_e$ [106]. Including the two-loop shifts using the following relations [106]:

$$A_e \;=\; 2\,\frac{g_V^\ell\, g_A^\ell}{(g_V^\ell)^2 + (g_A^\ell)^2}\,, \qquad A_f \;=\; 2\,\frac{g_V^f\, g_A^f}{(g_V^f)^2 + (g_A^f)^2}\,, \tag{C.9}$$

we obtain the total corrections for the asymmetry observables:

$$\begin{aligned}
\delta A_b^{\mathrm{1L+2L}}/A_b^{\mathrm{SM}} &= (2.6316\, c^1_{QQ} + 0.3447\, c^8_{QQ} - 3.1499\, c^1_{Qt} + 0.0062\, c^8_{Qt} + 0.6133\, c^1_{tt}) \times 10^{-4},\\
\delta A_{b,\mathrm{FB}}^{\mathrm{1L+2L}}/A_{b,\mathrm{FB}}^{\mathrm{SM}} &= (24.593\, c^1_{QQ} + 8.0737\, c^8_{QQ} - 29.576\, c^1_{Qt} + 0.5198\, c^8_{Qt} + 51.083\, c^1_{tt}) \times 10^{-4}.
\end{aligned} \tag{C.10}$$

D Additional fit results

We report here the 95% CL bounds on each of the five WCs from each process, summarised in Table 5. The corresponding individual and marginalised limits from the combined fit are shown in Table 6. All bounds are quoted at both linear and quadratic order in the EFT expansion.

	Order	$t\bar{t}H$	$t\bar{t}t\bar{t}$	$t\bar{t}b\bar{b}$	$gg\to H$	$t\bar{t}$	EWPO
c^1_{tt}	$\mathcal{O}(\Lambda^{-2})$	[-79.98, 37.36]	[-11.31, 1.42]	–	–	[-6.94, 11.72]	[-14.92, -1.68]
	$\mathcal{O}(\Lambda^{-4})$	[-79.98, 37.36]	[-1.62, 1.89]	–	–	[-6.94, 11.72]	[-14.92, -1.68]
c^1_{QQ}	$\mathcal{O}(\Lambda^{-2})$	–	[-19.00, 2.40]	–	–	[-14.75, 22.89]	[-0.94, 2.49]
	$\mathcal{O}(\Lambda^{-4})$	[-13.20, 11.92]	[-3.20, 3.85]	[-9.31, 9.32]	–	[-14.06, 9.36]	[-0.94, 2.49]
c^8_{QQ}	$\mathcal{O}(\Lambda^{-2})$	–	[-57.01, 7.20]	[-39.53, 91.23]	–	[-16.18, 26.89]	[-13.31, 16.79]
	$\mathcal{O}(\Lambda^{-4})$	[-38.30, 24.08]	[-9.60, 11.55]	[-21.91, 17.67]	–	[-41.86, 17.41]	[-13.31, 16.79]
c^1_{Qt}	$\mathcal{O}(\Lambda^{-2})$	[-26.19, 68.73]	[-2.35, 18.64]	–	[-3.14, 11.23]	[-6.77, 9.34]	[-2.08, 0.78]
	$\mathcal{O}(\Lambda^{-4})$	[-10.40, 13.92]	[-3.27, 2.78]	[-9.39, 9.26]	[-3.16, 10.93]	[-7.27, 9.69]	[-2.08, 0.78]
c^8_{Qt}	$\mathcal{O}(\Lambda^{-2})$	[-117.00, 39.40]	[-44.30, 5.59]	[-38.41, 88.66]	[-6.71, 42.30]	[-8.69, 18.14]	–
	$\mathcal{O}(\Lambda^{-4})$	[-44.52, 22.92]	[-5.73, 6.59]	[-21.56, 17.34]	[-7.06, 39.16]	[-56.11, 19.70]	–

Table 5: 95% CL individual bounds from each process.

	Individual		Marginalised	
	$\mathcal{O}(\Lambda^{-2})$	$\mathcal{O}(\Lambda^{-4})$	$\mathcal{O}(\Lambda^{-2})$	$\mathcal{O}(\Lambda^{-4})$
c^1_{tt}	[-9.12, -0.89]	[-1.66, 1.49]	[-16.30, -1.45]	[-1.50, 1.52]
c^1_{QQ}	[-0.99, 2.41]	[-0.84, 3.03]	[-8.42, 8.50]	[-0.97, 5.94]
c^8_{QQ}	[-11.39, 11.46]	[-8.60, 10.84]	[-75.66, 48.77]	[-19.43, 5.85]
c^1_{Qt}	[-1.74, 1.03]	[-2.30, 1.17]	[-10.99, 6.71]	[-1.37, 2.69]
c^8_{Qt}	[-7.93, 12.88]	[-4.39, 6.63]	[-9.50, 58.36]	[-4.09, 6.20]

Table 6: 95% CL individual and marginalised bounds from the combined fit.

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