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Subleading operators and γ_5-scheme dependence in SMEFT for Higgs boson pair production

Stefano Di Noi,[a] **Ramona Gröber,**[a] **Gudrun Heinrich,**[b,*] **Jannis Lang,**[b] **Ludovic Scyboz**[c] **and Marco Vitti**[d,e]

[a] *Dipartimento di Fisica e Astronomia "G. Galilei", Università di Padova, Italy, and Instituto Nazionale di Fisica Nucleare, Sezione di Padova, 35131 Padova, Italy*

[b] *Institute for Theoretical Physics, Karlsruhe Institute of Technology (KIT), Wolfgang-Gaede-Str. 1, 76131 Karlsruhe, Germany*

[c] *School of Physics and Astronomy, Monash University, Clayton VIC 3800, Australia*

[d] *Institute for Theoretical Particle Physics, Karlsruhe Institute of Technology (KIT), Wolfgang-Gaede-Str. 1, 76131 Karlsruhe, Germany*

[e] *Institute for Astroparticle Physics, Karlsruhe Institute of Technology (KIT), Hermann-von-Helmholtz-Platz 1, 76344 Eggenstein-Leopoldshafen, Germany*

E-mail: gudrun.heinrich@kit.edu, jannis.lang@kit.edu, stefano.dinoi@phd.unipd.it, ramona.groeber@pd.infn.it, ludovic.scyboz@monash.edu, marco.vitti@kit.edu

The calculation of contributions from the chromomagnetic and four-top-quark-operators within Standard Model Effective Field Theory (SMEFT) to Higgs boson pair production in gluon fusion is presented, in combination with NLO QCD corrections. Here we focus on the γ_5-scheme dependence introduced by the four-top-quark-operators and the interplay with other operators contributing to this process in SMEFT.

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*Speaker

1. Introduction

Higgs boson pair production plays a special role in the LHC program since it is the prime process to constrain the trilinear Higgs coupling. The gluon fusion production mode has the largest cross section, therefore a lot of effort has been put into providing increasingly accurate predictions for this process.

In these proceedings we focus on the description of $gg \to hh$ within Standard Model Effective Field Theory (SMEFT), combining leading and subleading operators with NLO QCD corrections in the Standard Model (SM) as described in more detail in Refs. [1, 2]. The contributions to $gg \to hh$ of the leading operators in SMEFT have been calculated in Ref. [3] and have been implemented in the `Powheg-Box-V2` event generator [4], while the calculation of the operators in the HEFT framework has been presented in Refs. [5, 6], including the NLO QCD corrections calculated in Ref. [7]. The chromomagnetic and the 4-top-operators are suppressed by loop factors compared to the leading operators when the potential UV completion is assumed to be a weakly coupling and renormalisable quantum field theory [8, 9]. We will demonstrate that these operators are intricately related through a γ_5-scheme dependence; the scheme dependence only cancels when they are consistently combined in a renormalised amplitude, as has been shown in Ref. [10] for the case of single Higgs production and in Ref. [1] for double Higgs production.

2. Operators contributing to $gg \to hh$ beyond the leading order

Any bottom-up EFT is defined by its degrees of freedom, its symmetries and a power counting scheme. SMEFT [11–14] builds on the field content and gauge symmetries of the SM and its main power counting, which relies on the counting of the canonical (mass) dimension, expanding in inverse powers of a new physics scale Λ which suppresses operators beyond dimension-4. The dominant contributions are expected to be described by dimension-6 operators, on which we focus here. We also impose a flavour symmetry $U(2)_q \times U(2)_u \times U(3)_d$ in the quark sector, which forbids chirality flipping operators bilinear in light quarks (including b-quarks), such that only 4-top-operators remain. Further, we neglect operators whose contributions involve diagrams with electroweak particles propagating in the loop.

With these restrictions, the dimension-6 CP even operators that contribute to $gg \to hh$, after electroweak symmetry breaking and in the unitary gauge, are given by

$$
\begin{aligned}
\mathcal{L}_{\rm SMEFT} \supset & -\left(\frac{m_t}{v}\left(1+v^2\frac{C_{H,\rm kin}}{\Lambda^2}\right)-\frac{v^2}{\sqrt{2}}\frac{C_{tH}}{\Lambda^2}\right) h\,\bar{t}\,t-\left(m_t\frac{C_{H,\rm kin}}{\Lambda^2}-\frac{3v}{2\sqrt{2}}\frac{C_{tH}}{\Lambda^2}\right) h^2\,\bar{t}\,t\\
& -\left(\frac{m_h^2}{2v}\left(1+3v^2\frac{C_{H,\rm kin}}{\Lambda^2}\right)-v^3\frac{C_H}{\Lambda^2}\right) h^3+\frac{C_{HG}}{\Lambda^2}\left(v\,h+\frac{1}{2}\,h^2\right) G^a_{\mu\nu}G^{a,\mu\nu}\\
& +g_s\bar{t}\,\gamma^\mu T^a\,t\,G^a_\mu+\frac{C_{tG}}{\Lambda^2}\sqrt{2}\,(h+v)\left(\bar{t}\,\sigma^{\mu\nu}T^a\,t\,G^a_{\mu\nu}\right)\\
& +\frac{C^{(1)}_{Qt}}{\Lambda^2}\bar{t}_L\gamma^\mu t_L\bar{t}_R\gamma_\mu t_R+\frac{C^{(8)}_{Qt}}{\Lambda^2}\bar{t}_L\gamma^\mu T^a t_L\bar{t}_R\gamma_\mu T^a t_R\\
& +\frac{C^{(1)}_{QQ}}{\Lambda^2}\bar{t}_L\gamma^\mu t_L\bar{t}_L\gamma_\mu t_L+\frac{C^{(8)}_{QQ}}{\Lambda^2}\bar{t}_L\gamma^\mu T^a t_L\bar{t}_L\gamma_\mu T^a t_L\\
& +\frac{C_{tt}}{\Lambda^2}\bar{t}_R\gamma^\mu t_R\bar{t}_R\gamma_\mu t_R\ ,
\end{aligned}
\tag{1}
$$

where $\sigma^{\mu\nu} = \frac{i}{2}\,[\gamma^\mu,\gamma^\nu]$ and $\tilde{\phi} = i\sigma_2\phi$ is the charge conjugate of the Higgs doublet. The first two lines in Eq. (1) contain the leading EFT operators, which have been studied in Ref. [3]. The remaining lines contain the chromomagnetic operator and the 4-top operators, where we use $O^{(1),(8)}_{QQ}$, related to the corresponding operators in the Warsaw basis [12] by

$$
C^{(1)}_{QQ} = C^{(1)\,3333}_{qq,\,\rm Warsaw}-\frac{1}{3}C^{(3)\,3333}_{qq,\,\rm Warsaw}\ ,\quad C^{(8)}_{QQ} = 4\,C^{(3)\,3333}_{qq,\,\rm Warsaw}\ . \tag{2}
$$

We emphasize that v denotes the full vacuum expectation value including a higher dimensional contribution of C_H/Λ^2 and the relation between the top-Yukawa parameter y_t of the SM Lagrangian and the top quark mass is given by

$$
m_t = \frac{v}{\sqrt{2}}\left(y_t-\frac{v^2}{2}\frac{C_{tH}}{\Lambda^2}\right)\ . \tag{3}
$$

In the following, we will explain the notions of 'leading' and 'subleading' operators we have used above. In SMEFT, the operators are ordered by their canonical dimension, i.e. the expansion relies on powers in E/Λ. However, in a perturbative expansion, in particular in the combination of an EFT expansion with expansions in a SM coupling, loop suppression factors also play a role. Therefore, a classification of operators into *potentially* tree-level induced and loop-generated operators [8, 14] can be a powerful criterion to identify the relative importance of dimension-6 operators in SMEFT. Loop-generated operators carry an implicit loop factor $\mathbf{L} = \left(16\pi^2\right)^{-1}$, they are typically given by operators containing at least one field strength tensor. We use a boldface notation for the loop factors that are not SM-induced. The loop factors can be derived by supplementing the SMEFT expansion by a chiral counting of operators [9], see also [15]. Such a classification cannot be derived without making some minimal UV assumptions, which are however quite generic, assuming renormalisability and weak coupling of the underlying UV complete theory. Under these assumptions, and if the Wilson coefficients C_i in the SMEFT expansion are considered to be of similar magnitude, it makes sense to expand in

$$
\frac{C_i}{\Lambda^a}\times 1/(16\pi^2)^b\ . \tag{4}
$$

Fixing $a = 2$ (dimension-6 operators), we call the operator contributions with $b = 0$ 'leading' and those with $b > 0$ 'subleading'. The above factors are to be combined with *explicit* loop factors $L = 1/(16\pi^2)^c$ from the SM perturbative expansion. We will see below that this classification is corroborated by observations from renormalisation and the cancellation of scheme-dependent terms [10]. Applying those rules to the Born contributions and associating loop factors of QCD origin with powers of g_s leads to $\mathcal{M}_{\rm Born} \sim O\left((g_s^2 L)\Lambda^{-2}\right)$.

2.1 Chromomagnetic operator insertions

The contribution of the chromomagnetic operator to the amplitude leads to the diagram types shown in Fig. 1. At first sight, the diagrams are at one-loop order. However, taking into account

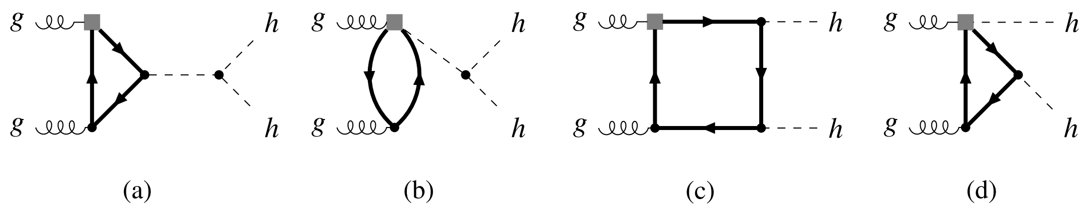


Figure 1: Feynman diagrams involving insertions of the chromomagnetic operator. The gray squares denote insertions of the chromomagnetic operator.

that the chromomagnetic operator belongs to the class of operators that, in renormalisable UV completions, can only be generated at loop level, the order in the power counting is $\mathcal{M}_{tG} \sim O\left((g_s^2 L)\mathbf{L}\,\Lambda^{-2}\right)$, which contains an additional factor $\mathbf{L} = 1/(16\pi^2)$ relative to the leading Born diagrams.

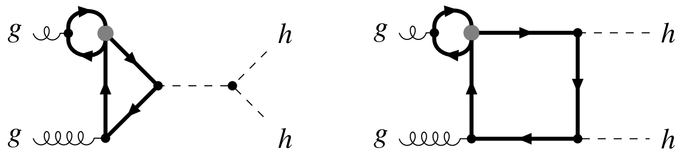
The diagrams of type (a), (b) and (d) are UV divergent even though they constitute the leading order contribution of C_{tG} to the gluon fusion process. This behaviour is well known [16–18] and leads to a renormalisation of $C^0_{HG} = \mu^{2\epsilon}\left(C_{HG} + \delta^{C_i}_{C_{HG}}\right)$, which in the $\overline{\rm MS}$ scheme takes the form [17–19]

$$\delta^{C_{tG}}_{C_{HG}} = \frac{(4\pi e^{-\gamma_E})^\epsilon}{16\pi^2\epsilon}\frac{4\sqrt{2}g_s m_t}{v} T_F\, C_{tG} \;. \tag{5}$$

2.2 Amplitude structure involving four-top operators

Four-top operators appear first at two-loop order in Higgs- or di-Higgs production in gluon-fusion, where the two loops are explicit. Taking into account the loop-generated nature of the chromomagnetic operator, their contribution is of the same order in the power counting as the chromomagnetic operator, i.e. $\mathcal{M}_{\text{4-top}} \sim O\left((g_s^2 L)\mathbf{L}\,\Lambda^{-2}\right)$. The complete set of diagrams involving 4-top-operators in $gg \to hh$ can be found in Ref. [1], here we only show in Fig. 2 those where a contraction of a one-loop subdiagram leads to topologies of Fig. 1.

In the following we sketch the relation between those classes of diagrams, focusing on the γ_5-scheme dependence, which first has been investigated in this context in Ref. [10]. The four-top operators contain chiral projection operators $(\mathbb{I} \pm \gamma_5)/2$. It is well-known that the treatment of γ_5 in dimensional regularisation is highly non-trivial, as γ_5 is an intrinsically four-dimensional object, see e.g. Refs. [20–23]. We will consider two different schemes for the continuation of γ_5 to $D = 4 - 2\epsilon$



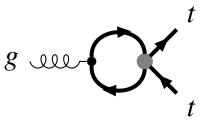
(a) (b)

Figure 2: Selected Feynman diagrams involving insertions of 4-top operators. The gray dots denote insertions of 4-top operators.

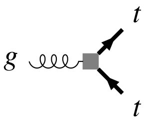
dimensions: naïve dimensional regularisation (NDR) [24] and the Breitenlohner-Maison-t'Hooft-Veltman (BMHV) [25, 26] scheme.

In our calculation, the treatment of γ_5 in the two schemes differs only by $O(\epsilon)$ parts of the Dirac algebra in D dimensions. Therefore, the renormalised result in the limit $D \to 4$ differs between the two schemes only by terms stemming from the ϵ-dependent parts of the Dirac algebra multiplying a pole of a loop integral.

The contributions to the gauge interactions from the diagrams in Fig. 2 for the case of an on-shell external gluon evaluate to



$$= \frac{C_{Qt}^{(1)} + \left(c_F - \frac{c_A}{2}\right) C_{Qt}^{(8)}}{C_{tG}} K_{tG} \times \tag{6}$$



,

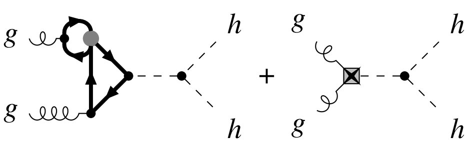
where we find

$$K_{tG} = \begin{cases} -\frac{\sqrt{2} m_t g_s}{16\pi^2 v} & \text{(NDR)} \\ 0 & \text{(BMHV).} \end{cases} \tag{7}$$

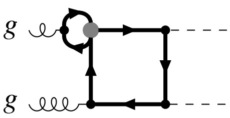
Since the Lorentz structure of the correction to the gauge vertex is similar to the insertion of a chromomagnetic operator, the diagrams in Fig. 2 acquire a UV divergence which, analogous to the case of the chromomagnetic operator, can be absorbed by a (now two-loop) counterterm of C_{HG}. In the $\overline{\text{MS}}$ scheme its explicit form is

$$\delta_{C_{HG}}^{\text{4-top}} = \frac{1}{\epsilon} \frac{(4\pi e^{-\gamma_E})^{2\epsilon}}{(16\pi^2)^2} \frac{(-4) g_s^2 m_t^2}{v^2} T_F \left(C_{Qt}^{(1)} + \left(c_F - \frac{c_A}{2} \right) C_{Qt}^{(8)} \right) . \tag{8}$$

Schematically, we therefore find



$$= \frac{C_{Qt}^{(1)} + \left(c_F - \frac{c_A}{2}\right) C_{Qt}^{(8)}}{C_{tG}} K_{tG} \left(\mathcal{M}_{tG}^{(a)} + \mathcal{M}_{tG}^{(b)} \right)$$



$$= \frac{C_{Qt}^{(1)} + \left(c_F - \frac{c_A}{2}\right) C_{Qt}^{(8)}}{C_{tG}} K_{tG} \mathcal{M}_{tG}^{(c)} , \tag{9}$$

where $\mathcal{M}_{tG}^{(a/b/c)}$ denote the amplitude of diagram types (a), (b) and (c) of Fig. 1, respectively.

3. Relations between Wilson coefficients in different γ_5-schemes

Scheme dependent contributions such as eq. (7) also arise in the corrections to the top-quark propagator and to the top-Higgs coupling. This scheme dependence has the same structure as the one in the process $gg \to h$ which is described in detail in Ref. [10]. The differences in the NDR and BMHV schemes originating from the mixing of four-fermion operators with chiral structure $(\bar{L}L)(\bar{R}R)$ into the chromomagnetic operator are well known in the context of flavour physics, where is was found that his effect can induce a scheme-dependent anomalous dimension matrix [27–31]. The strategy proposed in [27, 29, 32] was to perform a finite renormalisation of the chromomagnetic operator to ensure a scheme-independent anomalous dimension matrix. However, when calculating a physical amplitude, the scheme dependence involving $C_{Qt}^{(1)}$ and $C_{Qt}^{(8)}$ must be compensated by scheme dependent values for the other parameters of the Lagrangian, resulting in an overall scheme independence of the EFT prediction. The γ_5 schemes hence represent equivalent parameterisations of the new physics effects and a translation between the two schemes can be achieved by means of finite shifts of the Lagrangian parameters. The explicit form of the translation relation between the NDR and the BMHV scheme in terms of parameter shifts, derived in Refs. [1, 10] is the following (with the top quark mass renormalisation in the on-shell scheme):

$$\begin{aligned}
\delta m_t^{\text{4-top; BMHV}} &= \delta m_t^{\text{4-top; NDR}} - \frac{m_t^3}{8\pi^2\Lambda^2}\left(C_{Qt}^{(1)} + c_F C_{Qt}^{(8)}\right) \\
C_{tH}^{\text{BMHV}} &= C_{tH}^{\text{NDR}} + \frac{\sqrt{2}m_t\left(4m_t^2 - m_h^2\right)}{16\pi^2 v^3}\left(C_{Qt}^{(1)} + c_F C_{Qt}^{(8)}\right) \\
C_{tG}^{\text{BMHV}} &= C_{tG}^{\text{NDR}} - \frac{\sqrt{2}m_t g_s}{16\pi^2 v}\left(C_{Qt}^{(1)} + \left(c_F - \frac{c_A}{2}\right)C_{Qt}^{(8)}\right) \, .
\end{aligned} \tag{10}$$

Eq. (10) describes a translation scheme, rather than suggesting parameter combinations in which the scheme dependence is absorbed, as the latter would require to define a 'canonical scheme'. In order to avoid such an arbitrary choice in physical predictions within SMEFT, combinations of Wilson coefficients which allow to cancel the scheme dependence at a given order should be considered. When matching to concrete models such relations are automatically fulfilled [10].

4. Phenomenological results

The results presented in the following were obtained for a centre-of-mass energy of $\sqrt{s}$ = 13.6 TeV using the PDF4LHC15_nlo_30_pdfas [33] parton distribution functions, interfaced to our code via LHAPDF [34], along with the corresponding value for α_s. We used $m_h = 125$ GeV, the top quark mass has been fixed to $m_t = 173$ GeV to be coherent with the virtual two-loop amplitude calculated numerically. We set the central renormalisation and factorisation scales to $\mu_R = \mu_F = m_{hh}/2$ and use 3-point scale variations unless specified otherwise.

To demonstrate the effect of different γ_5-schemes on an individual, scheme dependent Wilson coefficient, we show the Higgs boson pair invariant mass distribution, m_{hh}, where we only include $C_{Qt}^{(1)}$ on top of the SM contribution in Fig. 3. We vary $C_{Qt}^{(1)}$ in the interval $-190 \leq C_{Qt}^{(1)} \leq 190$, a range that is inspired by marginalised fits described in Ref. [35]. The grey band denotes the SM scale uncertainties.

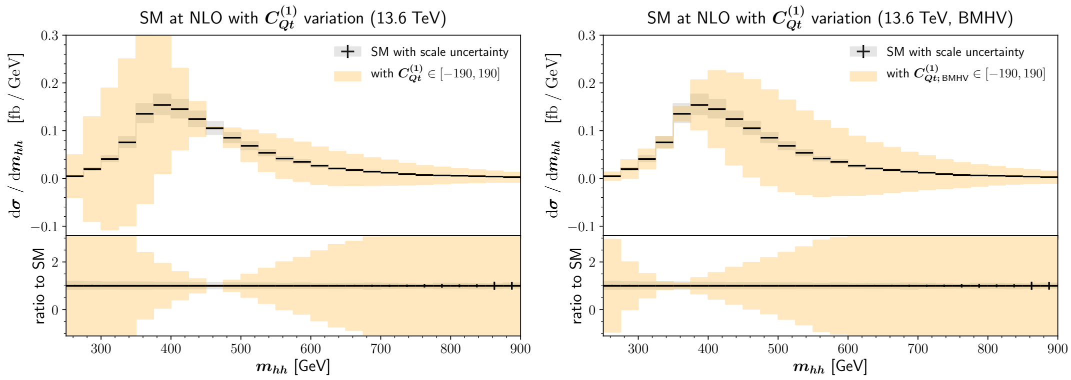


Figure 3: Effects of $C_{Qt}^{(1)}$-variations on m_{hh}-distributions comparing γ_5-schemes. Left: NDR scheme, right: BMHV scheme. The interval for $C_{Qt}^{(1)}$ is oriented at $O\left(\Lambda^{-2}\right)$ constraints from Ref. [35].

Even though the variation range is debatable due to the lack of tight constraints, it is obvious that the scheme differences can be very large for individual Wilson coefficients. For the case of $C_{Qt}^{(1)}$, in NDR (left), the low m_{hh}-regions exhibits a very large effect way beyond the SM scale uncertainties, with unphysical cross sections at very low m_{hh} values and a sign change around $m_{hh} \sim 460$ TeV. This behaviour changes significantly in BMHV (right): there are much weaker effects in the low m_{hh}-region, the sign change occurs around $m_{hh} \sim 360$ TeV and the deviation in the high m_{hh}-region is more pronounced.

We would like to point out that we have combined these operators with the leading SMEFT operators including NLO QCD corrections as described in Refs. [1, 3]. This combination is provided as an extension to the public `ggHH_SMEFT` code as part of the `POWHEG-Box-V2` [4].

5. Conclusions

We have discussed the calculation of contributions from the chromomagnetic operator and 4-top operators to Higgs boson pair production in gluon fusion and argued that these operators both appear at the same order in a power counting scheme that takes into account whether dimension-6 SMEFT operators are loop-generated or (potentially) tree-generated. We have shown that the contributions of those Wilson coefficients, when considered individually, depend on the chosen γ_5-scheme, and we have provided relations that allow a translation between the NDR and BMHV schemes. The explicit example of the 4-top-operator $C_{Qt}^{(1)}$ illustrates that the differences induced by a scheme change can be larger than the SM scale uncertainties. To obtain meaningful results for constraints on such Wilson coefficients, it is therefore recommended not to study those coefficients which are connected through scheme translations in isolation, as only their combination is a scheme independent parametrisation of BSM physics at the considered order in the power counting.

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