Higgs production at NLO in the Standard Model
Effective Field theory

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Abstract. The Effective Field Theory approach is a fruitful way of constraining new physics in a model independent way. As the Higgs sector is one of the most popular candidate for deviations from the Standard Model, it is particularly important that the constraints extracted from the experimental data on the Higgs boson be as meaningful as possible, which entails making accurate and precise theoretical predictions.

In this proceeding, I discuss a two loop calculation performed to improve the existing Leading Order result for the Higgs gluon-fusion cross section in the Standard Model Effective Field Theory. We present the modern multi-loop calculation techniques employed to obtain this result and highlight the unusual divergence structure of the amplitude.

1. Introduction
The Run I of the LHC has yielded the much awaited discovery of a scalar boson with properties compatible with the Standard Model (SM) Higgs boson [1, 2]. Observing the Higgs boson at ATLAS and CMS in different production channels and decay modes has opened the possibility to probe its couplings to other SM particles, showing a good agreement between predictions and observations [3].

Despite the current success of the SM in predicting Higgs boson properties and other LHC phenomena, we know that our theory for fundamental interactions is incomplete. The Standard Model Effective Field Theory (SMEFT) provides a consistent framework to parametrize possible deviations in a model independent way by encapsulating new physics effects in higher-dimensional operators involving the SM fields [4–6]. This approach is very powerful as it can be systematically related to any renormalizable theory beyond the SM (BSM) in which all the new fields are much heavier than LHC energies.

Recent years have seen a rapid progress in the development of precise predictions in the SMEFT, in particular in the form of a number of next-to-leading-order (NLO) results in top physics [8–11] and in Higgs boson physics [12–15].

Our work is part of this effort to provide precise SMEFT predictions for key LHC processes. We wish to improve the existing leading-order (LO) prediction for Higgs-boson production through gluon fusion [16] by including the NLO QCD radiative corrections. We focus on operators that modify the couplings of the top quark and the Higgs boson.
Figure 1: Sample diagrams in the leading order contribution to gluon fusion for $O_1$, $O_2$, and $O_3$. The blobs indicate the dimension-6 operators.

This proceeding focuses on the calculation of the virtual corrections to gluon fusion in the SMEFT and is organized as follows: section 2 details how we set up the calculation in the SMEFT and discusses the relevant operators. section 3 discusses the structure of the amplitude. Finally, section 4 describes the multi-loop techniques used for the evaluation of the virtual corrections.

2. EFT Setup

The SMEFT is the extension of the SM by the tower of higher dimensional operators involving SM fields:

$$\mathcal{L}_{\text{SMEFT}} = \mathcal{L}_{\text{SM}} + \sum_{i=1}^{\infty} \sum_{j=0}^{N_i} \frac{C^{(i)}_j}{\Lambda^i} O^{(i)}_{ij},$$

where the operators $O^{(i)}_{ij}$ have an energy dimension of $4+i$, $\Lambda$ is an energy scale larger than typical LHC energies and $C^{(i)}_j$ are dimensionless couplings. While non-renormalizable in general, the SMEFT Lagrangian can be safely truncated at some order $1/\Lambda^n$ and the resulting set of operators is stable under renormalization as long as observables are also truncated at that order. For most practical purposes it is sufficient to work at the order $1/\Lambda^2$: it is the first non-trivial order which features operators that affect virtually all aspects of LHC physics. A dimension 5 operator exists but is related to neutrino physics and bears little effect on collider observables [4].

Three CP-conserving operators affect gluon fusion by modifying the top loop [11]:

$$O_1 = \bar{Q}_L \phi t_R (\phi^\dagger \phi - \frac{v^2}{2}) + \text{h.c.},$$

$$O_2 = g_s^2 G^a_{\mu\nu} G^{\mu\nu}_a (\phi^\dagger \phi - \frac{v^2}{2}) + \text{h.c.},$$

$$O_3 = g_s \bar{Q}_L T^a \sigma^{\mu\nu} l_R G^{\mu\nu}_a \phi + \text{h.c.}.$$

Their leading-order contribution to this process is shown in figure 1. The one-loop renormalization of the SMEFT has been studied in detail [7] and we can use these results to obtain the required counterterms. We work in dimensional regularization with a number of dimensions $d = 4 - 2\epsilon$. We use the same conventions as [11], adding to each operator its Hermitian conjugate even if it is already Hermitian, and subtracting vacuum expectation values from operators to avoid shifting SM definitions of fields and couplings. This leads to including counterterms for the strong coupling and the top quark mass proportional to EFT couplings.
The renormalization of the relevant parameters is the following:

\[ Z_g = \mu^{-2\epsilon} \left( 1 + \alpha_s \frac{\delta Z_g}{\epsilon} + \alpha_s C_3 \frac{\delta Z_g}{\epsilon} \right), \quad \alpha_s^B = \mu^{2\epsilon} S_\epsilon^{-1} \alpha_s (1 + \alpha_s \frac{\delta Z_a}{\epsilon} + \alpha_s C_3 \frac{\delta Z_a}{\epsilon}), \]

\[ m_t^B = m_t (1 + \alpha_s \frac{\delta Z_a}{\epsilon} + \alpha_s C_3 \frac{\delta Z_g}{\epsilon}), \quad C_1^B = \mu^{3\epsilon} \left( C_1 (1 + \alpha_s \frac{\delta Z_1}{\epsilon}) + C_3 \alpha_s \frac{\delta Z_1}{\epsilon} \right), \quad C_2^B = C_2 + C_3 \frac{\delta Z_2}{\epsilon} + \alpha_s C_3 \frac{\delta Z_2}{\epsilon} \left( \frac{\delta Z_2^2}{\epsilon^2} + \frac{\delta Z_1^2}{\epsilon} \right). \]

We denote bare parameters with a superscript \( B \) and renormalized parameters without a superscript. The renormalization scale is \( \mu \), \( Z_g \) is the gluon field renormalization constant and the \( \delta Z_X \) are the counterterms (the superscript \( C \) indicates counterterms to SM parameters generated by the EFT). Finally, \( S_\epsilon = (4\pi)^\epsilon \exp(-\gamma \epsilon) \) and \( \gamma \) is the Euler-Mascheroni constant.

These counterterms highlight a recurrent theme in EFT renormalization: operator mixing. The bare SMEFT couplings are expressed not only in terms of their renormalized counterpart, but also have contributions coming from other SMEFT couplings. This behavior does not only impact on counterterms, but also the running of these operators, which mix together when the renormalization scale changes. All the counterterms required for the calculation exist in the literature except for \( \delta Z_2^3 \), which is a two-loop counterterm that will be fixed by our calculation.

3. Structure of the amplitude

The SM gluon fusion process, as well as the contributions generated by \( O_1 \) and \( O_3 \) are loop-induced processes and their NLO virtual corrections will therefore consist of two-loop diagrams. We will first focus on the LO calculation, which has unusual features, and then move on to discussing the NLO corrections.

3.1. Leading-order amplitude

The amplitude \( A \) at leading order is the sum of four contributions, which are exemplified in figure 1:

\[ A = \alpha_s^B A_{SM}^{(0)}(m_t^B) + \alpha_s^B C_1 (m_t^B) + \alpha_s^B C_2^B (m_t^B) + C_3 A_3^{(0)}(m_t^B), \]

where \( A_{SM}^{(0)} \) is the sum of LO SM diagrams and the \( A_i^{(0)} \) are the sums of diagrams containing operators \( O_i \). It is interesting to note that the diagrams in \( A_3^{(0)} \) have UV divergences already at leading order. This divergence has already been noted in previous work [16, 17] and the LO amplitude is renormalized by the leading-order mixing of \( C_2 = C_2^B = C_2 + C_3 (\delta Z_2^3/\epsilon) + \mathcal{O}(\alpha_s) \). As a result the expression of the amplitude in terms of renormalized quantities is

\[ A = \alpha_s A_{SM}^{(0)}(m_t) + \alpha_s C_1 A_1^{(0)}(m_t) + \alpha_s C_2 A_2^{(0)}(m_t) + \alpha_s C_3 \left( A_3^{(0)}(m_t) + S_\epsilon^{-1} \frac{\delta Z_2^3}{\epsilon} A_2^{(0)}(m_t) \right) + \mathcal{O}(\alpha_s^2). \]

3.2. Next-to-leading-order amplitude

The next-to-leading order correction to the amplitude is the sum of two contributions: virtual corrections, which are corrections of the LO arising through the exchange of a virtual gluon, and real emissions, in which an extra parton is emitted. In the final cross section the square of the real emissions and the interference between the Born-level diagrams and the renormalized virtual...
corrections contribute to the same order and constitute the NLO corrections to the amplitude as shown in Figure 2. The real emissions are at most one-loop diagrams and can therefore be handled by MadGraph5_aMC@NLO [18]. We will focus on discussing the virtual corrections, which include two-loop diagrams and exhibit more unusual behavior related to the leading-order renormalization. In terms of bare quantities, the amplitude at NLO is expressed as

$$A = Z_g \left( \alpha_s A_{SM}^{(0)}(m_t^B) + \alpha_s^2 C_1 A_1^{(0)}(m_t^B) + \alpha_s C_2 A_2^{(0)}(m_t^B) + \epsilon A_3^{(0)}(m_t^B) \right) + \left( \alpha_s^2 C_1 A_1^{(1)}(m_t^B) + \alpha_s^2 C_2 A_2^{(1)}(m_t^B) + \alpha_s C_3 A_3^{(1)}(m_t^B) \right),$$

where $A_{SM}^{(0)}$ is the sum of NLO SM diagrams and the $A_i^{(1)}$ are the sum of NLO diagrams generated by operator $O_i$. We can then express our amplitude in terms of renormalized quantities

$$A = \alpha_s A_{SM}^{(0)}(m_t) + \alpha_s^2 S^{-1} \left( A_{SM}^{(1)}(m_t) + \frac{\delta Z_\alpha + \delta Z_\delta}{\epsilon} A_{SM}^{(0)}(m_t) + m_t \frac{\delta Z_m}{\epsilon} \frac{\partial A_{SM}^{(0)}(m_t)}{\partial m_t} \right)$$

$$+ \alpha_s C_1 A_1^{(0)}(m_t) + \alpha_s^2 C_1 S^{-1} \left( A_1^{(1)}(m_t) + \frac{\delta Z_\alpha + \delta Z_\delta + \delta Z_{11}}{\epsilon} A_1^{(0)}(m_t) + m_t \frac{\delta Z_m}{\epsilon} \frac{\partial A_1^{(0)}(m_t)}{\partial m_t} \right)$$

$$+ \alpha_s C_2 A_2^{(0)}(m_t) + \alpha_s^2 C_2 S^{-1} \left( A_2^{(1)}(m_t) + \frac{\delta Z_\alpha + \delta Z_{23}}{\epsilon} A_2^{(0)}(m_t) + m_t \frac{\delta Z_m}{\epsilon} \frac{\partial A_2^{(0)}(m_t)}{\partial m_t} \right)$$

$$+ \alpha_s C_3 \left( A_3^{(0)}(m_t) + \frac{\delta Z_{13}}{\epsilon} A_3^{(0)}(m_t) \right) + \alpha_s^2 C_3 S^{-1} \left[ A_3^{(1)}(m_t) + \frac{\delta Z_\alpha + \delta Z_2 + \delta Z_{23}}{\epsilon} A_3^{(0)}(m_t) \right]$$

$$+ \frac{\delta Z_{13}}{\epsilon} A_1^{(0)}(m_t) + \frac{\delta Z_{23}}{\epsilon} A_2^{(0)}(m_t) + \left( \frac{\delta Z_{13}}{\epsilon} + \frac{\delta Z_{23}}{\epsilon} \right) A_2^{(0)}(m_t) + \frac{\delta Z_{13}^C}{\epsilon} A_3^{(0)}(m_t)$$

$$+ m_t \frac{\delta Z_m}{\epsilon} \frac{\partial A_{SM}^{(0)}(m_t)}{\partial m_t} + m_t \frac{\delta Z_m}{\epsilon} \frac{\partial A_{SM}^{(0)}(m_t)}{\partial m_t},$$

which makes apparent the UV divergences of the amplitude.

The renormalized NLO virtual amplitude still has divergences that will be canceled against divergences that appear in the infrared and collinear regions of the phase space of the extra
parton in the real emission contributions [19]. The structure of infrared and collinear divergences is universal: for the virtual contributions, we know that the divergences of the renormalized amplitude $A = \alpha_s/(2\pi)A^{(0)}(\alpha_s/(2\pi))^2 A^{(1)}$ appear in a factorized form: $A^{(1)} = A^{(1)\text{finite}} + I_1 A^{(0)}$, where $I_1$ is a universal operator acting on the colors of the external particles. In the case of $gg \rightarrow H$, the color structure is trivial, which makes $I_1$ take the simple form
\[
I_1^{gg \rightarrow H} = -\frac{\epsilon^\gamma}{\Gamma(1-\epsilon)} \left( \frac{C_A}{\epsilon^2} + \frac{\beta_0}{\epsilon} \right) \left( \frac{\mu^2}{-s} \right)^\epsilon. \tag{10}
\]
For example, if we look at the part of our amplitude proportional to $C_2$, we therefore know that
\[
A^{(1)}_i(m_i) + \frac{\delta Z_a + \delta Z_a}{\epsilon} A^{(0)}_2(m_i) + m_i \frac{\delta Z_m}{\epsilon} \frac{\partial A^{(0)}_2(m_i)}{\partial m} = A^{(1)\text{finite}} + I_1 A^{(0)}. \tag{11}
\]
By subtracting its infrared poles from the complete amplitude, we obtain an expression that we know to be finite. We can therefore use this constraint to fix the value of $\delta Z_{23}$ by imposing that all the poles in the amplitude cancel.

4. Multi-loop calculation techniques

The virtual corrections to the gluon fusion process consist of a combination of one- and two-loop diagrams. In this section we discuss the techniques we used to evaluate them.

The first step in such a calculation is to obtain expressions for the diagrams. To this end we use QGRAF to generate diagrams that are then translated to FORM expressions using a private Python program. These expressions are manipulated in FORM to evaluate Dirac matrix traces and index contractions and then imported using Mathematica. It is much easier to work with scalar objects, which we can do by exploiting the simple Lorenz structure of the amplitude:
\[
A = \epsilon_\mu(p_1)\epsilon_\nu(p_2)\delta_{ab}(p_1 \cdot p_2 g^{\mu\nu} - p^\mu_1 p^\nu_2 - p^\mu_2 p^\nu_1)M,
\tag{12}
\]
where $a, b$ are the color indices. The two-loop part of this form factor $M$ is a linear combination of scalar integrals:
\[
M|_{2L} = \sum a_i \int \frac{d^dk_1}{(2\pi)^d} \frac{d^dk_2}{(2\pi)^d} \frac{N_i}{D_{i1} \ldots D_{iq}}, \tag{13}
\]
where the numerators $N$ are products of contracted momenta and the denominators are the usual $D_{ij} = (Q_{ij}^2 - m_{ij}^2)$ where $Q_{ij}$ is some linear combination of momenta and $m_{ij}$ some internal mass (here the top mass or 0). These integrals number in the hundreds and it would not be wise to evaluate them one by one; an organizational principle is required. Integrals are regrouped into “families”, also called “topologies”, which are sets of integrals that share a common list of denominators with arbitrary powers. In our case we defined three families, described in figure 3, which can account for all integrals in our amplitude. Numerators are also accounted for using negative powers of the denominators. For example,
\[
\int \frac{d^dk_1}{(2\pi)^d} \frac{d^dk_2}{(2\pi)^d} \frac{k_2 \cdot k_1}{D_{11} D_{12} D_{17}} = \int \frac{d^dk_1}{(2\pi)^d} \frac{d^dk_2}{(2\pi)^d} \frac{m_i^2}{D_{11} D_{12} D_{17}} + \int \frac{d^dk_1}{(2\pi)^d} \frac{d^dk_2}{(2\pi)^d} \frac{1}{D_{11} D_{12} (D_{10})^{-1} D_{17}} \tag{14}
\]
\[
= m_i^2 I_1(1,1,0,0,0,0,0,1) + I_1(1,1,0,0,0,-1,1).
\]
As a result we can specify any integral in a given family by giving the exponent of each integral, and each can be seen as a point in $\mathbb{Z}^n$ ($n = 7$ here).
Integration-By-Parts Identities (IBP), which exploit the fact that integrals of total derivatives know how to obtain it using a method called the Laporta algorithm [30]. This method relies on integrals within a family can always be expressed as a linear combination of a finite basis of is the same as the denominators appearing in each diagram.

Figure 3: The three families used for our calculations. The set of denominators for each family is the same as the denominators appearing in each diagram.

Integral families are not only a bookkeeping device, we know that the infinite number of integrals within a family can always be expressed as a linear combination of a finite basis of integrals [29], often called master integrals. Not only do we know a finite basis exists, but we also know how to obtain it using a method called the Laporta algorithm [30]. This method relies on Integration-By-Parts Identities (IBP), which exploit the fact that integrals of total derivatives are 0 in dimensional regularization and that integral families are stable under differentiation,

$$\int \frac{d^dk_1}{(2\pi)^d} \frac{d^dk_2}{(2\pi)^d} \frac{\partial}{\partial k_1^\mu} \frac{\partial}{\partial k_2^\mu} v^\mu D_{11}^{(a_1)} \cdots D_{17}^{(a_2)} = 0. \quad (15)$$

The derivative shifts powers in the denominators and generates numerators, meaning that the left hand side of equation 15 is a linear combination of integrals in the topology. Let us see in a simple example how this can be used to find a basis for a family of master integrals. Our example will be the family of one-loop bubble integrals: $I_b(a_1, a_2) = \int \frac{d^dk}{(k^2)^{a_1}((k+p)^2)^{a_2}}.$

We can generate two independent IBP relations (differentiating with respect to $k_1$ or $k_2$) for this family:

$$I_b(a_1, a_2) = \frac{a_1 + a_2 - d - 1}{p^2(a_2 - 1)} I(a_1, a_2 - 1) + \frac{1}{p^2} I(a_1 - 1, a_2) \quad (16)$$

$$= \frac{a_1 + a_2 - d - 1}{p^2(a_1 - 1)} I(a_1 - 1, a_2) + \frac{1}{p^2} I(a_1, a_2 - 1) \quad (17)$$

These have been set up to highlight how we use these relations: to reduce the powers of the denominators. If we need an integral with a given set of $(a_1, a_2)$, we can express it using these relations recursively in terms of integrals with lower powers. Fortunately, the recursion ends when either power is 0 because the resulting integral is scaleless: $I_b(a_1, 0) = I_b(0, a_2) = 0.$

As shown in figure 4, this results in any integral being expressed in terms of $I(1,1).$ In more complicated cases, not all integrals can be expressed in terms of a single master, but a linear
Figure 4: Reduction of $I_b(3,2)$ to the master integral $I_b(1,1)$ by recursive application of the IBP identities for the bubble topology illustrated in the plane $(a_1,a_2)$. The recursion ends on the lines $a_1 = 0$ and $a_2 = 0$ because all integrals on these are scaleless.

combination of several integrals, which can be obtained from others by recursively applying IBP relations.

In our case the 1969 integrals we need to evaluate can be expressed as linear combinations of only 17 master integrals, which are the same as those required for the SM calculation of Higgs production at two loops [31, 32]. We performed this reduction with LiteRed [24] and cross-checked it with FIRE [25]. All 17 of them have been calculated in previous work [31–35] and we could use these results to evaluate $M$.

The final amplitude is a linear combination of multiple polylogarithms (MPL) depending on the variable

$$x = \frac{\sqrt{4m_t^2 - s} - \sqrt{-s}}{\sqrt{4m_t^2 - s} + \sqrt{-s}}$$

This class of function appears in a large number of loop calculations. MPLs are expressed as iterated integrals over a logarithmic kernel:

$$G(a_1, \ldots, a_n; x) = \int_0^x \frac{dt}{t-a_1}G(a_2, \ldots, a_n, t)$$

$$G(a; x) = \log(1 - x/a).$$

Using these expressions for the scalar integrals, we can evaluate the amplitude in a closed form. We checked our implementation of the published integrals by numerically evaluating their expressions in terms of polylogarithms with Ginac [28], and comparing with the direct numerical evaluation of the integrals in FIESTA [26] and SecDec [27]. The explicit form of the amplitude is of little interest for this proceeding, but its evaluation allowed us to determine the value of the new counterterm that appears in the calculation, which we will report in an upcoming paper.

5. Summary and conclusion

We have reported on the progress achieved in the calculation of the NLO virtual amplitude for the Higgs boson production through gluon fusion. We have shown that several dimension-six operators can modify this process in the SMEFT and discussed how modern multi-loop techniques have enabled us to compute the amplitude for virtual corrections to this process. We are now poised to combine this result with the automatic calculation of the real emission contribution in MadGraph5_aMC@NLO and to provide a more detailed discussion of our result as well as to study the NLO phenomenology of this process in a forthcoming paper.

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