Probing top-Higgs non-standard interactions at the LHC

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Effective interactions involving both the top quark and the Higgs field are among the least constrained of all possible (gauge invariant) dimension-six operators in the Standard Model. Such a handful of operators, in particular the top quark chromomagnetic dipole moment, might encapsulate signs of the new physics responsible for electroweak symmetry breaking. In this work, we compute the contributions of these operators to inclusive Higgs and $t\bar{t}h$ production. We argue that: i) rather strong constraints on the overall size of these operators can already be obtained from the current limits/evidence on Higgs production at the LHC; ii) $t\bar{t}h$ production will provide further key information that is complementary to $t\bar{t}$ measurements, and the possibility of discriminating among different contributions by performing accurate measurements of total and differential rates.

I. INTRODUCTION

The Standard Model (SM) has been tested with an impressive accuracy and is so far in excellent agreement with the experimental data. The room left for new physics at the TeV scale is therefore getting more and more squeezed, thanks to the LHC. Effective field theory (EFT) provides a model independent parametrization of the potential deviations from the SM while keeping its successes if the new degrees of freedom are heavy. EFT has been intensively used for instance in flavor physics to translate the accuracy of the measurements into strong constraints on the coefficients of the associated operators [1]. Slightly softer constraints on the operators involving weak bosons have also been derived from the electroweak precision measurements [2, 3]. In comparison, the operators involving the top quark are poorly constrained so far [4], especially the chromomagnetic moment operator of the top quark [5, 6], while those involving the Higgs field remain largely unconstrained. However, this status is about to change. In particular, modifications of the top quark interactions can significantly change the main Higgs production mechanism at hadron colliders, which is under scrutiny at the LHC.

In this paper, we focus on operators that involve both the top quark and the Higgs field. Not only they are little tested, but it is also where one might expect new physics associated with electroweak symmetry breaking to show up. First, we compute their contributions to $gg \rightarrow h$ due to a top loop. The only non-trivial contribution due to the chromomagnetic operator is logarithmic divergent and can be as large as the SM one. We then derive the constraints from the experimental bound on the Higgs production rate. Higgs production by gluon fusion alone does not allow to distinguish the new contributions since they are all proportional to the SM amplitude. In Section IV, we argue that $t\bar{t}h$ production can provide complementary information to further constrain and to disentangle the various top-Higgs operators.

II. OPERATORS OF INTEREST

Recently, constraints on effective Higgs interactions from the latest Higgs searches have been derived [7–13], with an emphasis on the $d = 6$ operators derived from the Higgs and SM gauge bosons. These papers display global fits in a large parameter space. While Ref. [8] and Ref. [9], are restricted to a particular UV set-up where only a sub-class of operators are important, Refs. [7, 10] included the modification of all Higgs interactions to the SM particles but considered that only the Yukawa coupling of the fermions were changed, and therefore have not considered the chromomagnetic operator. The spirit of this work is different in that our motivation is to focus only on $d = 6$ operators that involve both the Higgs field and the top quark. We study their effect on Higgs production by gluon fusion and associated with a $t\bar{t}$ pair, assuming in particular that $hWW$ and $hZZ$ tree-level couplings are not affected by new physics. The results of our analysis can easily be updated once the $hWW$ and $hZZ$ couplings are better determined. We start with the effective lagrangian [14–16]

$$\mathcal{L} = \mathcal{L}_{SM} + \sum \frac{c_i}{\Lambda^2} O_i + O\left(\frac{1}{\Lambda^4}\right).$$

The chromomagnetic dipole moment operator modifies the interactions between the gluons and the top quark,

$$O_{hg} = (\bar{Q}_L H) \sigma^{\mu\nu} T^a t_R G^a_{\mu\nu} \, ,$$

where $\sigma^{\mu\nu} = \frac{i}{2} [\gamma^{\mu}, \gamma^{\nu}]$ and $T^a$ is such that $\text{Tr}(T^a T^b) = \delta^{ab}/2$. Besides, one operator contains the top density

$$O_{Ht} = H^\dagger H (\bar{Q}_L) t_R$$

and three operators can be built from the top and Higgs currents,

$$O_{Ht} = H^\dagger D_\mu H (\bar{Q}_L)^R \gamma^\mu t_R$$
$$O_{HQ} = H^\dagger D_\mu H \bar{Q}_L \gamma^\mu Q_L$$
$$O_{HQ}^{(3)} = H^\dagger \sigma^I D_\mu H \bar{Q}_L \sigma^I \gamma^\mu Q_L \, .$$
Other operators of dimension 6 play a role in the top-Higgs interaction even though they do not contain both fields. One of them is

\[ O_H = \partial_\mu (H^\dagger H) \partial^\mu (H^\dagger H), \]  

which amounts to an overall renormalization of the Higgs wave function and therefore to a trivial shift of the top-quark Yukawa coupling [17].

The corrections from those operators to Higgs production by gluon fusion are shown in Fig. 1. In the large top mass limit, the contribution of the operators in Eqs. (4,3) can be seen as corrections to the \( O_{HG} \) operator

\[ O_{HG} = \frac{1}{2} H^\dagger H G^a_{\mu\nu} G^a_{\mu\nu} \]  

generated by the scale anomaly. Therefore, we are going to derive the constraints on \( O_{HG} \) from Higgs production, which we will then re-express in terms of limits on a combination of the above operators.

One should remark that not only the Higgs production rate is sensitive to the modifications of the top interactions but also the \( h \rightarrow \gamma \gamma \) decay. The operator \( O_H \) does not change the branching ratios since it multiplies all partial widths by the same factor. However, \( O_{H_H} \) and the electromagnetic version of \( O_{hg} \) induce

\[ O_{H_H} = \frac{1}{2} H^\dagger H F_{\mu\nu} F^{\mu\nu}. \]  

The main effect of this operator will be to relax the constraints from the \( h \rightarrow \gamma \gamma \) channel. We reiterate that we do not consider corrections to \( hhWW \) and \( hZZ \) vertices. New top interactions affect all these channels at one-loop. However, their effects to the loop-induced processes \( h \rightarrow \gamma \gamma \) and \( gg \rightarrow h \) are expected to be relatively larger than for \( h \rightarrow WW \) and \( h \rightarrow ZZ \) because the new operators modify the SM leading order in the first case and the NLO corrections in the second.

III. HIGGS PRODUCTION BY GLUON FUSION

\( O_{HG} \) is the only dimension-six operator inducing Higgs production by gluon fusion at tree-level. Its effect on the partonic cross-section is (see also Refs. [18, 19])

\[ \sigma (gg \rightarrow h) = \sigma (gg \rightarrow h)_{SM} \left( 1 + \frac{c_{HG} 6\pi v^2}{A^2 \alpha_s} \right)^2, \]  

where we have taken the heavy top limit for the SM, i.e., \( m_t > m_H/2 \), and \( v \approx 246 \text{ GeV} \) is the Higgs vacuum expectation value (vev). The contribution from \( O_{HG} \) is quite large compared to the SM one \( (6\pi v^2/\alpha_s \sim 10 \text{ TeV}^2) \) because the latter is only generated at the loop-level. Consequently, the upper limits on the Higgs production cross-section from the Tevatron [20] and the LHC [21, 22] strongly constrain the allowed range for \( c_{HG} \), as shown on Fig. 2. For this figure, we assume that only \( O_{HG} \) is added to the SM Lagrangian, i.e., we neglect the modifications of the other production mechanisms or of the decay widths except for \( h \rightarrow gg \). We used the same NNLO K factor for the contribution of the \( O_{HG} \) as for the SM [24] since both amplitudes are the same up to a global factor. The errors on these limits have been estimated by varying simultaneously the renormalization (\( \mu_R \)) and factorization scales (\( \mu_F \)) for the SM and the \( O_{HG} \) tree-level contributions. Other theoretical errors are much smaller. For \( m_H = 125 \text{ GeV} \), we obtain \(-0.29 \lesssim c_{HG}(\text{TeV}^2/A^2) \lesssim 0.036 \). We also show in yellow how the constraints on \( c_{HG} \) are relaxed when including the effect of \( O_{H_H} \). The exclusion in the plane \( (c_{HG}, c_{H_H}) \) is shown in Fig. 3. Again, Fig. 2 is valid only for SM \( hWW \) and \( hZZ \) couplings but a similar plot can
be drawn once the actual values of $hWW$ and $hZZ$ will be determined.

The constraints on $c_{HG}$ of Eq. (2) translate into constraints on a combination of the coefficients of the operators Eqs. (2)−(5). Contrary to the result reported in the first version of our paper as well as in Ref. [25], the one-loop correction from $O_{hg}$ to the operator $O_{HG}$ diverges logarithmically in a way consistent with the general expectations from dimension-six operators. Its one-loop contribution can be written as

$$\delta c_{HG} = \frac{g_s m_t}{4\Lambda^2} \Re c_{hg} \log \left( \frac{\Lambda^2}{m_t^2} \right). \quad (9)$$

The operators $O_{Hy}$ and $O_H$ renormalize the top mass

$$m_t = y_t \frac{v}{\sqrt{2}} - \frac{\Re(c_{Hy})}{2\sqrt{2}} \frac{v^3}{\Lambda^2} \quad (10)$$

and/or the top Yukawa coupling,

$$L^{hht} = \bar{t} \Gamma H \gamma_{\pm} t \propto \left( H^\dagger H \right) \frac{J^\mu \pm J_5^\mu}{2} \quad (14)$$

because the vector current is conserved. Their contributions to Higgs production through the effective operator $H^\dagger H \Gamma_{\mu \nu} \Gamma_{\mu \nu}$, generated by the axial anomaly, vanish in the SM due to parity. This result is consistent with the operator relations derived in Ref. [27].

Taking $m_t = 174.3$ GeV, $m_H = 125$ GeV, $v = 246$ GeV, $\Lambda = 1$ TeV and $g_s = 1.2$, we obtain

$$\delta c_{HG} \approx 0.1 \Re c_{hg} - 0.006 c_y. \quad (15)$$

Even if the effects due to the new interactions of the top quark are loop suppressed, they cannot be neglected. The coefficient $c_y$, probing the relation between the top mass and its Yukawa coupling, is not constrained by any other process than Higgs production (see recent and rather weak constraints on $c = 1 - c_y (v/\Lambda)^2$ in

$$\frac{\delta c_{HG}}{\Lambda^2} = \frac{\alpha_s}{6\pi v^2} \times \left( -c_y \frac{v^2}{\Lambda^2} \right). \quad (13)$$

The other three operators listed in Eq. (4) do not contribute to Higgs production by gluon fusion. In fact, the vertex $htt$ comes from the sum of those operators and of their Hermitian conjugates [10]. The relevant part of the operators can thus be written as

$$\delta c_{HG} \approx 0.1 \Re c_{hg} - 0.006 c_y. \quad (15)$$

Even if the effects due to the new interactions of the top quark are loop suppressed, they cannot be neglected. The coefficient $c_y$, probing the relation between the top mass and its Yukawa coupling, is not constrained by any other process than Higgs production (see recent and rather weak constraints on $c = 1 - c_y (v/\Lambda)^2$ in

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FIG. 3: The dashed blue and solid red lines are the limits from $h \to (WW, ZZ)$ and $h \to \gamma \gamma$ respectively. The $WW/ZZ$ constraints on $c_{HG}$ are stronger only when the branching ratio to $\gamma \gamma$ goes below $10^{-3}$ (SM value), corresponding to $0 \lesssim c_{HG} \lesssim 0.1$. For larger branching ratio, the $\gamma \gamma$ constraints are stronger and do not allow for large values of $c_{HG}$. Note that the allowed region is symmetric along the dotted black lines where $\sigma(gg \to h) = 0$ and $\Gamma(h \to \gamma \gamma) = 0$. We have checked that a more refined analysis combining all the channels along the lines of Ref. [24] gives qualitatively similar results, although slightly more constraining of course.
Refs. [9, 10]. Similarly, the present constraints on $c_{HG}$ due to top pair production [5] including the latest ATLAS combination [29], i.e., $-0.75 \lesssim c_{HG} (\text{TeV}/\Lambda)^2 \lesssim 3$ at 1σ and $-1.2 \lesssim c_{HG} (\text{TeV}/\Lambda)^2 \lesssim 3.5$ at 2σ, still allow the contribution from the chromomagnetic operator to have a noticeable effect on the allowed range for $c_{HG}$ as will be illustrated in the summary plots of Sec. IV.

The next question concerns the typical expectation for the size of the coefficient $c_{HG}$. For example, the one-loop contributions from $O_H$ and $O_{HG}$ have been shown to be as large as the $O_{HG}$ contribution in little Higgs models [20]. The reason is that those operators $O_H$ and $O_{HG}$ can be induced by the tree-level exchange of a heavy particle while $O_{HG}$ is only generated at the loop-level in a perturbative UV completion of the SM (see Fig. 4). The operator $O_H$ is also enhanced compared to $O_{HG}$ in strongly interacting Higgs models [17].

On the contrary, the chromomagnetic operator can hardly be enhanced. It is also generated only at the loop-level (see Fig. 4) in perturbation theory and thus for $O_{hg}$ to be the dominant new physics effects requires $O_{HG}$ to be relatively suppressed. While the diagram of Fig. 4 can be obtained by using twice the lower part of the second diagram in Fig. 5, the first diagram in Fig. 5 with $c = 1$ does not imply the presence of $O_{HG}$. As a consequence, it is possible to generate the chromomagnetic operator, $O_{hg}$, at one-loop and not the operator $O_{HG}$. An explicit example is given in Appendix A. While dominant in this example, the effects from the chromomagnetic operator are too small to be observed. Alternatively, the hierarchy may come from strongly coupled theories and can be estimated with the help of Naive Dimensional Analysis [30, 31]. If only the right-handed top is strongly coupled, the dominant operator involves four top quarks yet does not contribute even at two-loop [32]. In that case, the coefficient of the chromomagnetic operator is only suppressed by one power of the strong coupling compared to two for $c_{HG}$ and both operators can have similar contribution when the strong coupling approaches $4\pi$. However, its effects may again be too small to be observed. So, let us now move to study the effect of these operators on $t\bar{t}h$ production.

IV. $t\bar{t}h$ Production

While both Higgs direct coupling to the gluons and new top interactions significantly affect Higgs production, they cannot be distinguished using this process only. Contrary to Higgs production by gluon fusion, the four operators $O_{HG}$, $O_{hg}$, $O_{H}$ and $O_{HG}$ all contribute to $t\bar{t}h$ at the tree-level (see Fig. 6). Again, the three operators in Eq. (4) have no contribution for this process due to parity. There is only one additional operator affecting this process,

$$O_G = f^{ABC} C^A \mu^A G^{B\bar{B}} G^{C\mu}.$$  

And this relation holds at NLO (at least in the flavor universal limit). The total cross-section at 14 TeV is given by

$$\sigma(pp \to t\bar{t}h) = \sigma(pp \to t\bar{t}h)_{SM} \left(1 - c_y \frac{\beta^2}{\Lambda^2}\right)^2$$

(17)
at 8 TeV by

$$
\frac{\sigma(pp \rightarrow t\bar{t}h)}{fb} = 128 + [94\Re c_{hg} - 9.7c_G + 27c_{HG} \\
- 15c_y \left( \frac{\text{TeV}}{\Lambda} \right)^2 + [53.9(\Re c_{hg})^2 + 137c_G^2 \\
+ 9.6c_{HG} + 0.4c_y^2 + 19.3\Re c_{cHG} \\
- 9.6\Re c_{hg}c_y - 1.2\Re c_{HG}c_{HG} \\
- 0.7c_{cHG}] \left( \frac{1\text{ TeV}}{\Lambda} \right)^{-4},
$$

(19)

and at 7 TeV by

$$
\frac{\sigma(pp \rightarrow t\bar{t}h)}{fb} = 86.3^{+10}_{-15} + [63^{+20}_{-14}\Re c_{hg} + 22.3^{+8}_{-HG} \\
- 10.2^{+4}_{-2.5}c_y - 5.6c_G] \left( \frac{\text{TeV}}{\Lambda} \right)^2 \\
+ \left[ \frac{43.6^{+17}_{-12}(\Re c_{hg})^2 + 78.3c_G^2 + 8.6^{+1}_{-3}c_{HG}}{c_y} \right] \\
+ 0.3c_y^2 + 21^{+6}_{-2}\Re c_{hg}c_{HG} - 7.2^{+1}_{-1.1}\Re c_{hg}c_y \\
- 1.5^{+1}_{-1}\Re c_{HG}c_{HG} - 1.1^{+1}_{-1}\Re c_{HG}c_{HG}] \left( \frac{\text{TeV}}{\Lambda} \right)^{-4},
$$

(20)

for $m_H = 125$ GeV. We included $c_G$ and $c_y^2$ terms for indication (but not $c_Gc_y$ terms), however, as mentioned earlier, we will set $c_G = 0$ in the rest of the analysis.

The same factorization and renormalization scales as for top pair production, i.e., $\mu_F = \mu_R = m_t$ have been used since we have only considered a light Higgs boson. The cross-section will slightly decrease if higher values taking into account the Higgs mass are chosen. The errors are again obtained by varying the factorization and renormalization scales simultaneously from $\mu_F = \mu_R = m_t/2$ to $\mu_F = \mu_R = 2 m_t$, except for the last two terms $\Re(c_{HG})c_{HG}$ and $c_{HG}c_{HG}$ for which the numerical errors are larger. Results have been obtained via the FeynRules-MadGraph 5 simulation chain [34–37].

The new physics has been computed at the tree-level and the SM contribution at NLO [24–26]. The $O(1/\Lambda^4)$ terms have been computed to check the $1/\Lambda$ expansion and only take into account the operators that contribute also at the $1/\Lambda^2$ order, i.e., contain either squares of the operators $O(1/\Lambda^2)$ or the interference of the SM with an amplitude involving two new vertices. Additional contributions from the operators in Eq. (1) or dimension-eight operators and proportional to the imaginary part of $c_{HG}$ or $c_{hg}$ are not included. The values of the $1/\Lambda^4$ coefficients tell us that the $1/\Lambda$ expansion breaks down around the TeV for $c_i = 1$. This lower value compared to top pair production [3] is expected due to the higher energy required for this final state. While Eqs. (18)–(20) have been obtained only for a particular value of the Higgs mass, the ratios of the new physics contributions over the SM do not change drastically with the Higgs mass as shown on Fig. 7.

As shown by Eqs. (18)–(20) and Fig. 7 $t\bar{t}$ associated Higgs production can mainly be affected by the chromo-magnetic operator. As a consequence, the constraints from a measurement of the $t\bar{t}h$ cross-section would complement those from Higgs production as illustrated in Fig. [8] which displays the $c_{hg}$ range allowed by the present measurements of the $t\bar{t}$ cross section. By the time the $t\bar{t}h$ cross section will be measured, the improved constraints from $t\bar{t}$ measurements will also help in reducing further the allowed range for $c_{hg}$ according to Ref. [5]:

$$
\delta\sigma_{pp \rightarrow t\bar{t}H} = 144 c_{hg} \left( \frac{\text{TeV}}{\Lambda} \right)^2 + 22.5 c_{hg}^2 \left( \frac{\text{TeV}}{\Lambda} \right)^4.
$$

(21)

Like for top pair production, the theoretical uncertainty is responsible for a sizable part of the allowed region. Since those errors mainly affect the overall normalization, this issue could be solved by measuring the shapes of the distributions. Additionally, shape effects could also lift the remaining degeneracy between the four operators. While the contributions of the operators $O_{H_G}$ and $O_H$ have the same shapes as the SM ones, the operators $O_{hg}$ and $O_{HG}$ can induce shape distortions. However, only the contribution of the chromomagnetic operator might have a higher energy dependence than the SM. If the Higgs leg is attached to the effective vertex, the diagrams contain only one chirality flip such that no other chirality flip is needed to interfere with the SM amplitude (Fig. [6c]). Moreover, the vertex is not proportional to the Higgs vev like for top pair production. Those advantages are lost if the Higgs is attached to the top line or to a gluon (Figs. [6b] and [d]). For those diagrams, the amplitude is proportional to $m_t$ and $v$ and no room is left for extra powers of the energy of the process.

The distributions of the transverse momentum of the Higgs, the total $H_T$ and the invariant mass of the Higgs-top system are displayed on Fig. [9]. The shapes of the
1/Λ^4 contributions are also shown for comparison. They are clearly stretched to high energy while the interference and the SM contributions have a very similar behavior. The interference with the diagrams in which the Higgs is connected at the effective vertex do not vanish but are apparently suppressed. The shape effects are only expected if the new physics scale Λ is close to the maximal energy probed because they are due only to the 1/Λ^4 contributions. The plots on the right show how the distributions can differ with respect to the SM in the case c_{hq}(TeV^2/Λ^2) = 1.

Finally, spin correlations could exhibit some dependence on c_{hg}. In the case of tH production, the deviations due to c_{hg} would be of the order of a few percents \cite{5}. For tth, the measurement will be much more challenging and we therefore do not compute the associated spin correlations here but might return to them in due time.

V. CONCLUSION

Only one dimension-six operator, O_{HG}, generates a tree level coupling between the Higgs boson and the gluons. This operator has the largest contribution to Higgs production. Nevertheless, the three operators modifying the contribution from the top loop also have sizable effects compared to the SM one and, in a large class of models, can be comparable to the effect of O_{HG} due to the hierarchy between their coefficients. All those operators are already constrained by the present limits on Higgs production at hadron colliders. However, Higgs production by gluon fusion only constrains a linear combination of these operators and cannot discriminate between them. Interestingly, a light Higgs makes real the possibility of partially solving this issue by using Higgs production in association with a pair of top quarks. Contrary to Higgs production, the leading contribution in this process comes from the chromomagnetic operator O_{hg}, which can therefore be further constrained from the measurement of the total tth cross-section. Shape effects do not come from the interference terms and are dominated by the square of the amplitude involving an effective vertex and could thus be observable for large c_{hg} values only.

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**Appendix A: Explicit example with $c_{HC} \ll c_{hg}$**

In this appendix, we provide a toy model in which the diagrams of Fig. 5 are generated while the diagram of Fig. 4 is not. The new sector is given by

\[ T_{L,R} \sim (3, 1, Y) \]
\[ \Phi \sim (1, 2, Y - 1/6) \]
\[ S \sim (1, 1, Y - 2/3) \]  

(A1)
where $Y \neq 2/3$ to avoid the mixing of $T$ with the SM top and $Y \neq -1/3$ to avoid the mixing between $\Phi$ and the Higgs doublet. The extra piece of the Lagrangian is given by

$$\mathcal{L}^{NP} = i\bar{T}PT - MTT - \kappa (\bar{T}_R Q_L \Phi + \bar{Q}_L T_R \Phi^\dagger) - \beta (i\bar{T}_R L^S \bar{S}^I + \bar{T}_L L^S R^S) + D_\mu S^I D^\mu S - M_2 S^I S + \lambda_1 (S^I)^2 + \lambda_2 S^I S H^I + \Lambda_3 S^I S \Phi^\dagger \Phi + D_\mu \Phi^\dagger \Phi^\dagger D^\mu \Phi - M_3 \Phi^\dagger \Phi + \lambda_3 ([\Phi^\dagger \Phi]^2)$$

where the parameters $M$, $M_S$, $M_\Phi$ and $M_3$ are around or above the TeV scale. The model has an accidental $Z_2$ symmetry under which all the SM model particles are even while the new ones are odd. This symmetry prevents any tree-level generation of the higher dimensional operators when the heavy particles are integrated out. The operator $\mathcal{O}_{HC}$ cannot be generated at one-loop since the colored particle does not couple to the Higgs. On the contrary, the equivalent operator for the photon cannot be avoided. Indeed, even if the fermions can be chosen to be neutral, all the new scalars cannot be simultaneously neutral. The constraints from gluon fusion in the low mass will change with the branching ratio to two photons. Nevertheless, the chromomagnetic operator is induced at one-loop and its coefficient given by

$$c_{\rho g} = \frac{\kappa \beta g_s M_3}{4(4\pi)^2 M^3} \left[ \frac{R^2_\Phi (1 - 3 R^2_\Phi)}{R^2_\Phi - 1} + \frac{1}{R^2_\Phi - 1} \right]$$

where $R_S = \frac{M_\Phi}{M}$ and $R_\Phi = \frac{M_\Phi}{M_3}$. 

[1] M. Bona et al. (UTfit Collaboration), JHEP 0803, 049 (2008), 0707.0636.
[2] Z. Han and W. Skiba, Phys.Rev. D71, 075009 (2005), hep-ph/0412166.
[3] R. Barbieri, A. Pomarol, R. Rattazzi, and A. Strumia, Nucl.Phys. B703, 127 (2004), hep-ph/0405040.
[4] C. Zhang, N. Greiner, and S. Willenbrock (2012), 1201.6670.
[5] C. Grazzini, M. Papucci, and S. Willenbrock (2012), 1201.6670.
[6] J. F. Kamenik, M. Papucci, and A. Weiler (2011), 1107.3143.
[7] D. Carmi, A. Falkowski, E. Kuflik, and T. Volansky (2012), 1202.3415.
[8] A. Azatov, R. Contino, and J. Galloway (2012), 1202.3415.
[9] J. Espinosa, C. Grojean, M. Muhlleitner, and M. Trott (2012), 1202.3697.
[10] P. P. Giardino, K. Kannike, M. Raidal, and A. Strumia (2012), 1203.4254.
[11] J. Ellis and T. You (2012), 1204.0464.
[12] A. Azatov, R. Contino, D. Del Re, J. Galloway, M. Grassi, et al. (2012), 1204.4817.
[13] M. Farina, C. Grojean, E. Salvioni (2012), 1205.0011.
[14] W. Buchmuller and D. Wyler, Nucl.Phys. B268, 621 (1986).
[15] B. Grzadkowski, M. Iskrzynski, M. Misiak, and J. Rosiek, JHEP 10, 085 (2010), 1008.4884.
[16] G. Buchalla and O. Cata (2012), 1203.6510.
[17] G. Giudice, C. Grojean, A. Pomarol, and R. Rattazzi, JHEP 0706, 045 (2007), hep-ph/0703164.
[18] A. V. Manohar and M. B. Wise, Phys.Lett. B636, 107 (2006), hep-ph/0601212.
[19] A. Pierce, J. Thaler, and L.-T. Wang, JHEP 0705, 070 (2007), hep-ph/0609049.
[20] T. Aaltonen et al. (CDF and D0 Collaboration) (2011), 1103.3233.
[21] G. Aad et al. (ATLAS Collaboration) (2011), ATLAS-CONF-2011-163.