Contributions from dimension six strong flavor changing operators to $t\bar{t}$, $t$ plus gauge boson, and $t$ plus Higgs boson production at the LHC

P.M. Ferreira *, R. Santos †
Centro de Física Teórica e Computacional, Faculdade de Ciências,
Universidade de Lisboa, Avenida Professor Gama Pinto, 2, 1649-003 Lisboa, Portugal

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Abstract. We study the effects of a set of dimension six flavor changing effective operators on several processes of production of top quarks at the LHC. Namely, $t\bar{t}$ production and associated production of a top and a gauge or Higgs boson. Analytical expressions for the cross sections of these processes are derived and presented.

I. INTRODUCTION

The large mass of the top quark [1] makes it a natural laboratory to search for new phenomena beyond those predicted by the Standard Model (SM). One possible avenue of research consists in using effective operators of dimensions larger than four to parameterize the effects of any new physics. One advantage of using this formalism is that one works in a model independent manner. The complete set of dimension five and six operators that preserve the gauge symmetries of the SM is quite large, and was first obtained by Buchmuller and Wyler [2]. This methodology has been used by many authors to study the top quark. For instance, in refs. [3] contributions to top quark physics arising from several types of dimension five and six operators were studied. In ref. [4] a detailed study of the $Wtb$ vertex was undertaken. Analysis of flavor changing neutral currents in supersymmetric theories and models with two Higgs doublets may be found in [5]. For a recent study of single top production in supersymmetric models see [6] and for a study on single top-quark production in flavor-changing $Z'$ models see [7]. NLO and threshold corrections to top quark flavor changing operators were obtained in [8]. The four fermion operator contributions to $t\bar{t}$ production were studied in detail in [9].

Recently [10, 11] we studied the effects on the phenomenology of the top quark of a subset of operators of dimension six - namely, those with contributions to strong flavor changing neutral currents. The set chosen included several operators studied by other authors, but which had not been considered together before. We also benefitted greatly from working in a fully gauge-invariant manner, taking advantage of the equations of motion to eliminate several unknown effective couplings. We considered both gluonic operators and four-fermion ones. A detailed analysis of the contributions of these operators in phenomena such as the top’s width, its rare decays and the cross section for single top production at the LHC was performed. It was shown that the operator set we chose may have large contributions to the single top production cross section at the LHC, and that that channel is an excellent probe into the existence of new physics.

In this paper we wish to analyze the effect of that set of effective operators on other potentially interesting channels of top production at the LHC, namely: top and anti-top production; associated production of a top quark with a single gauge boson (a photon, a $W$ or a $Z$ boson); and associated production of a top quark and a Higgs boson. Our aim, as in references [10, 11], is to produce analytical expressions whenever possible, so that the results of this paper may be used directly by our experimental colleagues in their Monte Carlo simulations. This work is structured as follows: in section II we will review the effective operator formalism and the operators studied in refs. [10, 11]. Namely, we will explain the criteria behind that choice and the role the equations of motion play in how many of them are truly independent. In section III we will study the effect of our operator set in the production of $t\bar{t}$ pairs at the LHC. In section IV we will analyze the processes of associated production of a top and a gauge boson and of a top and a Higgs boson. In section VI we will present numerical results for the cross sections of these processes at the LHC. Finally, we will conclude in section VII with an overall analysis of the results.

* ferreira@cii.fc.ul.pt
† rsantos@cii.fc.ul.pt
II. THE EFFECTIVE OPERATOR APPROACH

A physical system rarely provides us enough information for a complete description of its properties. A way to solve this problem is to parameterize any physical effects not yet observed by introducing an effective lagrangian with a set of new interactions to be determined phenomenologically. This effective lagrangian has the Standard Model as its low energy limit, and can serve to represent the effect of any high-energy theory at a given energy scale $\Lambda$. We write this lagrangian as

$$\mathcal{L} = \mathcal{L}^{SM} + \frac{1}{\Lambda^2} \mathcal{L}^{(5)} + \frac{1}{\Lambda^4} \mathcal{L}^{(6)} + O\left(\frac{1}{\Lambda^6}\right),$$

where $\mathcal{L}^{SM}$ is the SM lagrangian and $\mathcal{L}^{(5)}$ and $\mathcal{L}^{(6)}$ are all of the dimension 5 and 6 operators which, like $\mathcal{L}^{SM}$, are invariant under the gauge symmetries of the SM. The $\mathcal{L}^{(5)}$ terms break baryon and lepton number conservation, and usually are not considered. This leaves us with the $\mathcal{L}^{(6)}$ operators. Some of these, after spontaneous symmetry breaking, generate dimension five terms. The complete list of effective operators was obtained in [2]. Our purpose, in this and previous works [10, 11] is to study flavor changing interactions, restricted to the strong sector of the theory, involving a top quark. Therefore, we choose operators with a single top quark, that do not comprise gauge or Higgs bosons (except for those that arise from covariant derivatives), and that involve some sort of strong flavor changing interactions. Finally, we choose those $\mathcal{L}^{(6)}$ operators that have no sizeable impact on low energy physics (by which we mean below the TeV scale).

Only two gluon operators survive these criteria which, in the notation of ref. [2], are written as

$$\mathcal{O}_{uG} = i \frac{\alpha_t}{\Lambda^2} \left( \bar{u}_R \gamma^\mu D^\mu u_R \right) G_{\mu\nu}^a,$$

$$\mathcal{O}_{uG\phi} = \frac{\beta_{tij}}{\Lambda^2} \left( \bar{g}_L \gamma^\mu \sigma_{\mu\nu} u_R^j \right) \phi G_{\mu\nu}^a.$$

$q_L$ and $u_R$ are spinors (a left quark doublet and up-quark right singlet of $SU(2)$, respectively), $\phi$ is the charge conjugate of the Higgs doublet and $G_{\mu\nu}^a$ is the gluon tensor. $\alpha_{ij}$ and $\beta_{tij}$ are complex dimensionless couplings and the $(i,j)$ are flavor indices. According to the criteria listed above, one of these indices must belong to the third generation. After spontaneous symmetry breaking the neutral component of the field $\phi$ acquires a vev $(\phi_0 \rightarrow \phi_0 + v)$, with $v = 246/\sqrt{2}$ GeV) and the second of these operators generates a dimension five term. The lagrangian for new physics is then given by

$$\mathcal{L} = \alpha_{tu} \mathcal{O}_{tu} + \alpha_{ut} \mathcal{O}_{ut} + \beta_{tu} \mathcal{O}_{tu\phi} + \beta_{ut} \mathcal{O}_{ut\phi} + \text{h.c.}$$

$$= i \frac{\alpha_t}{\Lambda^2} \left[ \alpha_{tu} \left( \bar{t}_R \gamma^\mu D^\mu u_R \right) + \alpha_{ut} \left( \bar{u}_R \gamma^\mu D^\mu t_R \right) \right] G_{\mu\nu}^a +$$

$$\frac{v + h/\sqrt{2}}{\Lambda^2} \left[ \beta_{tij} \left( \bar{i}_L \gamma^\mu \sigma_{\mu\nu} u_R^j \right) + \beta_{ut} \left( \bar{u}_L \gamma^\mu \sigma_{\mu\nu} t_R \right) \right] G_{\mu\nu}^a + \text{h.c.},$$

where $h$ is the SM Higgs boson. This lagrangian describes new vertices such as $g \bar{t} u$, $g \gamma \bar{t} u$, $g Z \bar{t} u$ and $g h \bar{t} u$. There are, of course, analogous vertices involving the top quark, instead of the anti-top one. We will also consider an analogous lagrangian (with new couplings $\alpha_{tc}$, $\beta_{ct}$, ...) for vertices of the form $g \bar{t} c$, $g \gamma \bar{t} c$, $g Z \bar{t} c$ and $g h \bar{t} c$. Notice how the operators with $\beta$ couplings correspond to a chromomagnetic momentum for the $t$ quark. Several extensions of the SM, such as supersymmetry and two Higgs doublet models, may generate contributions to this type of operator [12].

The Feynman rules for these anomalous vertices are shown in figs. [1], [2] and [3], with quark momenta following the arrows and incoming gluon momenta. The only operators that contribute to the Feynman rule in fig. [2] are $\mathcal{O}_{ut}$ and $\mathcal{O}_{tu}$. They generate the vertices $g \gamma \bar{t} u$ and $g Z \bar{t} u$ when we consider the electroweak gauge fields present in the covariant derivatives of eq. [3]. On the other hand, the Feynman rule in fig. [3] comes from the operator $\mathcal{O}_{uG\phi}$, where the vev was replaced by the Higgs field. Of course, we have considered analogous vertices involving the $c$ quark instead of the $u$ one.

In ref. [10] we calculated the effect of these operators on the width of the quark top. They allow for the decay $t \rightarrow u g$ ($t \rightarrow c g$) (which is also possible in the SM, but only at higher orders), and the corresponding width is given by

$$\Gamma(t \rightarrow u g) = \frac{m_t^3}{12 \pi \Lambda^4} \left( m_t^2 |\alpha_{ut} + \alpha_{tu}^*|^2 + 16 v^2 (|\beta_{tu}|^2 + |\beta_{ut}|^2) \right) +$$

$$\Gamma(t \rightarrow c g) = \frac{m_t^3}{12 \pi \Lambda^4} \left( m_t^2 |\alpha_{ct} + \alpha_{tc}^*|^2 + 16 v^2 (|\beta_{ct}|^2 + |\beta_{tc}|^2) \right) +$$

or, in terms of the coefficients

$$\Gamma(t \rightarrow u g) = \frac{m_t^3}{12 \pi \Lambda^4} \left( m_t^2 |\Delta\alpha|^2 + 16 v^2 (|\Delta\beta|^2) \right) +$$

$$\Gamma(t \rightarrow c g) = \frac{m_t^3}{12 \pi \Lambda^4} \left( m_t^2 |\Delta\alpha|^2 + 16 v^2 (|\Delta\beta|^2) \right) +$$
and an analogous expression for $\Gamma(t \to cg)$. In this expression and in all results of this paper we have set all quark masses, except that of the top, equal to zero. We performed the full calculations and verified that the error introduced by this approximation is extremely small. Notice how the top width (4) depends on $\Lambda^{-4}$. There are processes with a $\Lambda^{-2}$ dependence, namely the interference terms between the anomalous operators and the SM diagrams of single top quark production, via the exchange of a $W$ gauge boson - processes like $u\bar{d} \to t\bar{d}$. They were studied in ref. [11] in detail, and we discovered that, due to a strong CKM suppression, the contributions from the anomalous vertices are extremely small.

As was discussed in refs. [10, 11], the operators that compose the lagrangian (3) are not completely independent. If one performs integrations by parts and uses the fermionic equations of motion [2, 13], one obtains the following relations between them:

\begin{align}
O_{ut}^I &= O_{tu} - \frac{i}{2}(\Gamma_u^I O_{ut}^I + \Gamma_u O_{tu}) \\
O_{tu}^I &= O_{tu} - i g_s \tilde{t} \gamma^\mu \gamma_R \lambda^a u \sum_i (\bar{u}^i \gamma^\mu \gamma_R \lambda_a + \bar{d}^i \gamma^\mu \gamma_R \lambda_a d^i)
\end{align}

(5)

where $\Gamma_u$ are the Yukawa couplings of the up quark and $g_s$ the strong coupling constant. In the second equation four-fermion terms appear, which means that they have to be taken into account in these studies. Indeed, their role was of great importance for the processes studied in ref. [11]. In the current paper, however, they will have no bearing in our results. The most interesting thing about eqs. (5), which is in fact a direct consequence of working in a fully gauge invariant manner, is that they tell us that there are two relations between the several operators. This means that we are allowed to set two of the couplings to zero. We have used this in ref. [11] to simplify immensely
All of these diagrams contribute to the process $p p \rightarrow t \bar{t}$ and $q \bar{q} \rightarrow t \bar{t}$, as shown in fig. 3. The expressions obtained, by setting one of the four-fermion couplings to zero, as well as making $\beta_{tu} = \beta_{tc} = 0$. For consistency, we will make the same choice in the current work, to allow for a direct comparison with the results of \[11\].

In refs. \[10\] \[11\] we considered the contributions from the anomalous flavor changing operators to single top production. In particular, we calculated all processes with a top quark in the final state, with a jet or isolated, stemming from either direct top production or associated production with a gluon or a light quark. We determined that the single top channel is an excellent one for detection of new physics, as our calculations demonstrated that one could obtain a significant increase in the cross section of single top quark via the anomalous couplings relative to the SM predicted production and the associated production of a top and a gauge or Higgs boson. These are all channels of great physical interest. In the former case, the LHC is expected to be a veritable top-anti-top factory, with an estimated production of around eight million top quark pairs per year and per experiment. With such high statistics, a deviation from the SM prediction has, a priori, a good chance of being detected. In the latter case, the final state presents a very clear signal for experiments. As such it should be easy to isolate it from the LHC backgrounds.

At this point we must emphasize one important aspect: we are not considering the most general set of operators for associated top plus gauge or Higgs boson production processes. As mentioned before, there are contributions from the electroweak dimension six operators that we could have considered as well. However, that is not our goal: we have established that the operator set we have chosen, which corresponds to a specific type of possible new physics (strong flavor changing interactions), may have a sizeable impact on the single top channel. We now wish to verify if the same operators might have important contributions to other interesting channels. This is an important verification, to ensure the consistency of the operator set chosen. If it predicts significant increases on several physical processes but only some of those are observed, then we will have a powerful clue that the operators we chose do not tell the whole story of the new top physics at the LHC. If, however, the observations are according to the predictions arising from these operators, that will constitute good evidence that they parameterize well whatever new physics lies beyond the SM.

III. CROSS SECTIONS FOR $g g \rightarrow t \bar{t}$ AND $q \bar{q} \rightarrow t \bar{t}$

There are three Feynman diagrams contributing to the partonic cross section of $t \bar{t}$ production, as is shown in fig. 3. All of these diagrams contribute to the process $p p \rightarrow t \bar{t}$ and interfere with the (analogous) SM tree-level diagrams for the same processes. Notice that, since each of these diagrams includes two anomalous vertices, they will generate amplitudes of the order $1/\Lambda^4$. With the Feynman rules shown in fig. 4 and the corresponding SM Feynman rules, it is easy to obtain the interference cross section for $g g \rightarrow t \bar{t}$, given by

$$\frac{d \sigma(g g \rightarrow t \bar{t})}{dt} = -\frac{g_s^2}{1536 \pi \Lambda^4} \frac{F_{gg}^1 [\alpha_{ut} + \alpha_{tu}]^2 + F_{gg}^2 [|\beta_{ut}|^2 + |\beta_{tu}|^2] + F_{gg}^3 Im[\alpha_{ut} \beta_{tu} - \alpha_{tu} \beta_{tu}^*]}{s^3 (m_t^2 - t) (m_t^2 - u) u},$$

where

$$F_{gg}^1 = 7 m_t^{12} t - 23 m_t^{10} t^2 + 16 m_t^{12} u - 20 m_t^{10} t u + 51 m_t^{10} u^2 - 73 m_t^8 t^3 u + 37 m_t^8 t^4 u^2 + 2 t^5 u^2 + 16 m_t^8 u^3 - 73 m_t^6 t u^3 - 31 m_t^4 t^2 u^3 + 42 m_t^2 t^3 u^3 - 16 t^4 u^3 + 37 m_t^4 t u^4 +$$

$$F_{gg}^2 = \frac{1}{s} \frac{2 m_t^8 t + 51 m_t^8 u - 73 m_t^6 t u + 37 m_t^6 t^2 u}{s} \frac{t^2 + 16 m_t^8 u^3 - 73 m_t^6 t u^3 - 31 m_t^4 t^2 u^3 + 42 m_t^2 t^3 u^3 - 16 t^4 u^3 + 37 m_t^4 t u^4 +}{s^2}$$

$$F_{gg}^3 = \frac{1}{s^2} \frac{2 m_t^4 t + 51 m_t^4 u - 73 m_t^2 t u + 37 m_t^2 t^2 u}{s} \frac{t^2 + 16 m_t^4 u^3 - 73 m_t^2 t u^3 - 31 m_t^4 t^2 u^3 + 42 m_t^2 t^3 u^3 - 16 t^4 u^3 + 37 m_t^4 t u^4 +}{s^2}.$$
occur in the SM at tree level. This is, thus, an order 1/detection. However, we must recall that the final states of these channels will also include jets, stemming either from the production of a top quark along with a gauge or Higgs boson, hopefully a much “cleaner” signal for experimental motion, we are allowed to set

\[
F_{gg}^2 = 16 v^2 t u (7 m^6 t - 15 m^4 t^2 + 8 m^2 t^3 + 7 m^6 u - 26 m^4 t u + 20 m^2 t^2 u + t^3 u - 15 m^4 u^2 + 20 m^2 t u^2 - 16 t^2 u^2 + 8 m^6 u^3 + t u^3)
\]

\[
F_{gg}^{3} = -2 v m_t (7 m^6 t - 23 m^4 t^2 + 16 m^2 t^3 + 7 m^6 u - 52 m^4 t u + 145 m^6 t^2 u - 158 m^4 t^3 u + 60 m^2 t^4 u - 23 m^8 u^2 + 145 m^6 t u^2 - 258 m^4 t^2 u^2 + 136 m^6 t^3 u^2 + 4 t^4 u^2 + 16 m^6 u^3 - 158 m^4 t u^3 + 136 m^2 t^2 u^3 - 64 t^3 u^3 + 60 m^2 t u^4 + 4 t^2 u^4)
\]

(7)

For the \( q \bar{q} \rightarrow t \bar{t} \) partonic channel we obtain

\[
\frac{d\sigma(q \bar{q} \rightarrow t \bar{t})}{dt} = \frac{g^2}{16 \pi \Lambda^4 s^3} \left( F_{qq}^1 |\alpha_{tu}|^2 + F_{qq}^2 |\alpha_{ut}|^2 + F_{qq}^3 Re[\alpha_{ut} \alpha_{tu}] + F_{qq}^4 (|\beta_{ut}|^2 + |\beta_{tu}|^2) + F_{qq}^5 Im[\alpha_{ut} \beta_{tu}] + F_{qq}^6 Im[\alpha_{tu} \beta_{tu}^*] \right),
\]

(8)

where

\[
F_{qq}^1 = (m^2 - t) u^2;
F_{qq}^2 = -(m^6 + t m^4 - 4 u m^4 - 2 t^2 m^2 + u^2 m^2 - 6 t u m^2 + t u^2);
F_{qq}^3 = -(m^6 - t m^4 - 4 u m^4 + 2 t u m^2 - 2 t u^2);
F_{qq}^4 = 8 (3 u m^2 + s^2 - u^2) u^2;
F_{qq}^5 = 2 m_t v (m^2 - 5 t m^2 + 4 u m^2 + 4 t^2 + 12 t u);
F_{qq}^6 = 2 m_t v (m^2 - t m^2 - 6 u m^2 + 4 t u).
\]

(9)

Despite the rather long expressions, notice that the dependence on the \{\alpha, \beta\} anomalous couplings is quite simple. We have kept all couplings in these expressions, but recall that, due to the freedom allowed by the equations of motion, we are allowed to set \( \beta_{tu} = 0 \). Finally, there are identical expressions for the partonic cross sections involving the charm anomalous couplings.

IV. CROSS SECTIONS FOR ASSOCIATED \( t \gamma, t Z, t W \) AND \( t h \) PRODUCTION

For the processes \( q \bar{q} \rightarrow t \gamma, t Z, t W \) there are once more contributions from three diagrams. This process does not occur in the SM at tree level. This is, thus, an order 1/\( \Lambda^4 \) process. In our previous papers, the top quark was produced alongside with a quark or a gluon, and therefore detected through a final state of \( t + \text{jets} \). Here, we have the production of a top quark along with a gauge or Higgs boson, hopefully a much “cleaner” signal for experimental detection. However, we must recall that the final states of these channels will also include jets, stemming either from
from figs. (1) and (2) we may obtain the cross section for associated top plus photon production, given by

\[
\frac{d \sigma(g q \rightarrow t \gamma)}{dt} = \frac{e^2}{18 m_t^2 s^2 t (t + u)^2} \left( (m_t^6 - t m_t^4 + s^2 m_t^2 + 3 st m_t^2 - 2 s^2 t) u \right) \Gamma(t \rightarrow q g) .
\]  

where \(\Gamma(t \rightarrow q g)\) stands for the decay width of a top quark into a light up quark and a gluon. This result is remarkably simple, and quite elegant. Similar expressions had been obtained in refs. \cite{11} for single top production in the direct, gluon-gluon and gluon-quark channels. In fact, we verified that every time there was a gluon in the initial or final states the differential cross section was always proportional to a partial decay width of the top. It is interesting to see the same thing happening when a gluon is replaced by a photon. Eq. \cite{13} then establishes a very powerful link between the rare decays of the top quark and the cross section for associated top plus photon production.

Let us now consider the associated production of a top quark and a \(Z^0\) gauge boson. The calculation is similar to the top plus photon channel, modulus the obvious kinematic differences, and the different Feynman rules. We obtain, for the differential cross section, the expression

\[
\frac{d \sigma(g u \rightarrow t Z)}{dt} = \frac{e^2 m_t^2}{1728 \pi \Lambda^4 S_u^2} \left( F_{1Z}^1 |\alpha_{ut} + \alpha_{gu}^*|^2 + F_{1Z}^2 Im[(\alpha_{ut} + \alpha_{gu}^*) \beta_{tu}] + F_{1Z}^3 |\beta_{tu}|^2 + F_{1Z}^4 |\beta_{ut}|^2 \right) ,
\]  

where

\[
F_{1Z}^1 = 18 m_t^2 m_z^2 s^2 t + 48 m_t^2 m_z^2 s^2 S_w^2 t + 9 s^2 t^3 + 32 m_t^6 m_z^2 S_u^4 u +
32 m_t^2 m_z^2 s^2 S_u^4 u - 18 m_t^2 s^2 t u + 48 m_t^2 s^2 S_w^2 t u -
32 m_t^2 S_w^2 t u + 96 m_z^2 S_u^2 t u - 64 m_z^2 s^2 S_w^2 t u + 9 s^2 t^2 u
\]

\[
F_{1Z}^2 = \frac{4}{m_t} \left( 18 m_t^2 m_z^2 s^2 t + 48 m_t^2 m_z^2 s^2 S_w^2 t + 9 s^2 t^3 + 32 m_t^6 m_z^2 S_u^4 u +
32 m_t^2 m_z^2 s^2 S_u^4 u - 18 m_t^2 s^2 t u + 48 m_t^2 s^2 S_w^2 t u -
32 m_t^2 S_w^2 t u + 96 m_z^2 m_z^2 S_u^4 t u - 64 m_z^2 s^2 S_w^2 t u + 9 s^2 t^2 u \right)
\]

\[
F_{1Z}^3 = \frac{16}{m_t^2} \left( 54 m_t^6 m_z^2 s - 54 m_t^4 m_z^2 s^2 - 48 m_t^4 m_z^2 s^2 S_w^2 + 48 m_t^4 m_z^2 s^3 S_w^2 -
54 m_t^4 m_z^2 s t + 54 m_t^2 m_z^2 s^2 t + 9 s^2 t^3 + 18 m_t^6 m_z^2 u - 54 m_t^4 m_z^2 s u +
18 m_t^2 m_z^2 s^2 u - 192 m_t^4 m_z^2 s^2 S_w^2 u + 192 m_t^2 m_z^2 s^2 S_w^2 u - 48 m_t^2 m_z^2 S_w^2 u +
32 m_t^6 m_z^2 S_w^2 u + 32 m_t^4 m_z^2 s^2 S_w^2 u - 18 m_t^4 m_z^2 s^2 t u + 54 m_t^2 m_z^2 s^2 t u -
18 m_t^2 m_z^2 s^2 t u - 32 m_t^4 m_z^2 S_u^4 t u + 96 m_t^2 m_z^2 S_u^4 t u - 64 m_t^2 m_z^2 S_w^2 t u +
9 s^2 t^2 u - 48 m_t^4 m_z^2 S_w^2 u^2 + 144 m_t^2 m_z^2 S_w^2 u^2 - 48 m_t^2 m_z^2 S_w^2 u^2 \right)
\]

\[
F_{1Z}^4 = \frac{16}{m_t^2} \left( 18 m_t^2 m_z^2 s^2 t + 48 m_t^2 m_z^2 s^2 S_w^2 t + 9 s^2 t^3 +
32 m_t^6 m_z^2 S_w^2 u + 32 m_t^4 m_z^2 s^2 S_w^2 u - 18 m_t^2 m_z^2 s^2 t u + 48 m_t^2 m_z^2 S_w^2 t u -
\right)
\]
where we have used, for short, the convention $S_\omega = \sin(\theta_W)$ and $S_{2\omega} = \sin(2\theta_W)$. The nice proportion to the top decay width is destroyed in this expression, caused, no doubt, by the fact that the $Z^0$ boson, unlike the gluon or photon, is not massless.

We now turn to $tW$ production. Contrary to the $tZ$ and $t\gamma$ processes, this one occurs at tree level in the SM. A discussion about the problems related with $tW$ production and detection can be found in \[14\]. In fig. 6 we show the SM Feynman diagrams contributing to this process, along with the single anomalous diagram that also contributes to it. There is only one diagram because there are no contributions of the type of fig. 2. The reason is simple: it is the covariant derivative terms that give rise to the “four-legged” diagram of fig. 2, and that derivative is acting on an $SU(2)$ gauge singlet. Therefore, only the hypercharge $U(1)$ gauge field, $B_\mu$, contributes to the vertex. This means that the operators $O_{uG}$ do not give rise to any diagram of the type of fig. 2 for the $tW$ channels.

A simple calculation yields the SM tree level cross section, which reads

$$d\sigma^{SM} (g d \to tW^-) = \frac{\alpha^2 g_s^2 |V_{td}|^2}{384 m_w^2 \pi s^3 S_w^2 (m_t^2 - t)^2} (m_t^2 u (-t u + m_t^2 (2 t + u)) + 2 m_w^2 (2 m_t^2 u + (s + u) (2 s^2 + 2 s u + u (-2 m_t^2 + u)))) .$$

The interference terms between the anomalous diagram and the SM ones (for an internal $q$ quark, with $q = u, c$) were computed in ref. 11, and are given by

$$d\sigma^{INT} (g d \to tW^-) = -\frac{\alpha^2 g_s |V_{td}| |V_{qd}| m_t v (-m_t^2 t u + 2 m_w^2 (s^2 + (m_t^2 + s) u))}{48 m_w^2 \pi s^2 S_w^2 (m_t^2 - t) t} Re[\beta_{q\gamma}] \frac{|\beta_{q\gamma}|^2}{\Lambda^2} .$$

Finally, the new anomalous term, which corresponds to the squared amplitude of the anomalous diagram in fig. 6, is given by

$$d\sigma^{NEW} (g d \to tW^-) = \frac{\alpha^2 |V_{q\gamma}|^2 (m_t^2 - t) (s t + 2 m_w^2 u) v^2}{24 m_w^2 \pi s^2 S_w^2 t} \frac{|\beta_{q\gamma}|^2}{\Lambda^4} .$$

Notice the dependence, in eqs. 13 and 14, on the CKM matrix elements. We concluded, in 10, that they were of vital importance for the final results.

Finally, we consider the production of a top quark associated with a SM Higgs boson. Like the $tZ$ and $t\gamma$ processes, this process does not occur at tree level in the SM. The Feynman diagrams are the same as for the top plus $\gamma$ or $Z$ processes. There are only two differences. First, we have to use the four point Feynman rule in fig. 3 instead of the one in fig. 2. The “four-legged” Feynman rule involves now only the $\beta$ couplings instead of the $\alpha$ ones, which is natural, when you remember that the $O_{uG}$ operators do not involve the Higgs boson in any way, whereas the $O_{uG\phi}$ ones do. The second difference is the fact that the diagram for top plus Higgs production analogous to the second one in figure 5 does not appear in these calculations. The reason is very simple - the Higgs-quark vertex is proportional to the mass of the quark in question and we have taken, as explained earlier, $m_u = m_c = 0$. 
The cross section for $t\bar{t}$ production can then be written as

$$\frac{d\sigma(g u \rightarrow t\bar{t})}{dt} = \{F_{th}^1|\alpha_{ut} + \alpha_{\bar{t}u}|^2 + F_{th}^2 Im[(\alpha_{ut} + \alpha_{\bar{t}u}^*)\beta_{tu}] + F_{th}^3 (|\beta_{tu}|^2 + |\beta_{ut}|^2)\} \frac{1}{\Lambda^4} ,$$

where

$$F_{th}^1 = \frac{e^2 m_t^2 s^2 (-(m_t^2 m_w^2) - s t + m_h^2 (4 s + t))}{48 m_w^2 (m_t^2 - s)^2 S_w^2}$$

$$F_{th}^2 = \frac{-e m_t^2 s^2 (-\sqrt{2} m_w (m_t^2 - s) S_w (2 m_t^2 - t)) + e m_t^2 (-m_t^2 + 2 (m_t^2 + s)) v)}{12 m_w^2 (m_t^2 - s)^2 S_w^2}$$

$$F_{th}^3 = \frac{-e m_t^2 s (s (-m_t^2 + s) t (2 m_w^2 (-m_t^2 + s) S_w^2 + 2 \sqrt{2} e m_t^2 m_w S_w v - e^2 m_t^2 v^2)) + m_t^2 s (2 m_w^2 (m_t^2 - s)^2 S_w^2 - 2 \sqrt{2} e m_w (m_t^2 - s)^2 S_w^2 v + e^2 (-(m_t^2 s) + (m_t^2 + s)^2) v^2)}{12 m_w^2 (m_t^2 - s)^2 S_w^2} .$$

V. RESULTS FOR THE INTEGRATED CROSS SECTIONS

A. $t\bar{t}$ production

The cross section for the gluon-gluon channel is identical for the processes with anomalous couplings of the $u$ or $c$ quarks. For the quark–antiquark cross section via a $c$ quark, the numerical results are extremely small, and we do not present them. We have used throughout this work the CTEQ6 part on density functions \[15\] and included a cut of 15 GeV on the transverse momentum of the partons in the final state. This will allow a direct comparison with the results of reference \[11\], where a similar cut was considered to help remove collinear and soft singularities in the gluon-quark processes. Finally, for this particular process, we have taken the factorization scale $\mu_F$ equal to twice the mass of the top quark. As was mentioned in refs. \[11,14\], this will produce smaller values of the cross sections than we would obtain if we had, for instance, set $\mu_F$ equal to the partonic center-of-mass energy. With these specifications, the results we obtain are

$$\sigma_{pp \rightarrow g g \rightarrow t\bar{t}} = \{-0.4 |\alpha_{ut} + \alpha_{\bar{t}u}|^2 + 7.6 (|\beta_{ut}|^2 + |\beta_{tu}|^2) + 9.1 Im[\alpha_{ut} \beta_{tu} - \alpha_{\bar{t}u} \beta_{\bar{t}u}^*]\} \frac{1}{\Lambda^4} \text{ pb}$$

$$\sigma_{pp \rightarrow g \bar{q} \rightarrow t\bar{t}} = \{-0.2 |\alpha_{ut}|^2 - 0.4 |\alpha_{\bar{t}u}|^2 + 0.5 Re[\alpha_{ut} \alpha_{tu}] - 0.5 (|\beta_{ut}|^2 + |\beta_{tu}|^2) - 0.6 Im[\alpha_{ut} \beta_{tu}] - 0.1 Im[\alpha_{tu} \beta_{\bar{t}u}^*]\} \frac{1}{\Lambda^4} \text{ pb} .$$

So the total results for the proton-proton cross sections are

$$\sigma_{pp \rightarrow t\bar{t}}^{(u)} = \{-0.6 |\alpha_{ut}|^2 - 0.8 |\alpha_{\bar{t}u}|^2 - 0.3 Re[\alpha_{ut} \alpha_{tu}] + 7.1 (|\beta_{ut}|^2 + |\beta_{tu}|^2) + 8.5 Im[\alpha_{ut} \beta_{tu}] + 9.0 Im[\alpha_{tu} \beta_{\bar{t}u}^*]\} \frac{1}{\Lambda^4} \text{ pb}$$

$$\sigma_{pp \rightarrow t\bar{t}}^{(c)} = \{-0.4 |\alpha_{ct} + \alpha_{\bar{t}c}|^2 + 7.6 (|\beta_{ct}|^2 + |\beta_{\bar{t}c}|^2) + 9.1 Im[\alpha_{ct} \beta_{\bar{t}c} - \alpha_{\bar{t}c} \beta_{ct}^*]\} \frac{1}{\Lambda^4} \text{ pb} .$$

B. $t\gamma$ and $tZ$ production

The results for top plus $\gamma$ production are particularly simple, as the cross sections are proportional to the top decay width to an up-type quark plus a gluon. The pdf suppression of the $c$ quarks, however, makes the corresponding contributions to the cross section extremely small, which is why we only present the $u$ quark terms. For this channel, we chose $\mu_F = m_t$. The proton-proton cross sections for top plus gamma production are then given by

$$\sigma_{pp \rightarrow t\gamma}^{(u)} = 228 \Gamma(t \rightarrow u g) |V_{tb}|^2 \text{ pb} ;$$
\[ \sigma_{pp \rightarrow t \gamma}^{(u)} = 32.6 \Gamma(t \rightarrow u g) |V_{ub}|^2 \text{ pb} . \] (20)

We have also presented the cross section for anti-top plus gamma production. To obtain this quantity we simply used the differential cross section for the \( t + \gamma \) channel, eq. (10), since there is no difference, in terms of effective operators, between both processes. The different numbers in eqs. (20), then, arise solely from different pdf contributions (namely the \( u \) and \( \bar{u} \) quarks).

For the \( tZ \) processes the expressions we obtain, after integrating on the pdfs (with \( \mu_F = m_t + m_Z \) ), are

\[ \sigma_{pp \rightarrow tZ}^{(u)} = \begin{cases} 4.0 |\alpha_{ut} + \alpha_{tu}^*|^2 + 32.1 Im[(\alpha_{ut} + \alpha_{tu}^*) \beta_{tu}] + 63.8 |\beta_{tu}|^2 + 65.3 |\beta_{ut}|^2 \end{cases} \frac{1}{\Lambda^4} \text{ pb} ; \\
\sigma_{pp \rightarrow tZ}^{(c)} = \begin{cases} 0.4 |\alpha_{ct} + \alpha_{tc}^*|^2 + 3.4 Im[(\alpha_{ct} + \alpha_{tc}^*) \beta_{tu}] + 6.7 |\beta_{tu}|^2 + 7.0 |\beta_{ut}|^2 \end{cases} \frac{1}{\Lambda^4} \text{ pb} ; \\
\sigma_{pp \rightarrow tZ}^{(c)} = \sigma_{pp \rightarrow tZ}^{(c)} = \begin{cases} 0.2 |\alpha_{ct} + \alpha_{tc}^*|^2 + 1.6 Im[(\alpha_{ct} + \alpha_{tc}^*) \beta_{tc}] + 3.2 |\beta_{tc}|^2 + 3.4 |\beta_{ct}|^2 \end{cases} \frac{1}{\Lambda^4} \text{ pb} . \] (21)

C. \( tW \) production

With the SM diagrams of figure (13), we obtain the expected tree-level result for \( \sigma_{pp \rightarrow tW^-} \), which is about 30 pb. Unlike the previous cases of production of a top quark alongside with a gauge boson, we now have interference terms between the SM diagrams and the anomalous ones, which are of order \( \Lambda^{-2} \). We obtained these interference cross sections in ref. [10] and present them here again. For the \( u \) quark anomalous couplings we have \( (\mu_F = m_t + m_W) \)

\[ \sigma^{INT}(pp \rightarrow tW^-) = 0.031 Re[\beta_{ut}] \frac{1}{\Lambda^2} \text{ pb} ; \]
\[ \sigma^{INT}(pp \rightarrow tW^+) = 0.022 Re[\beta_{ut}] \frac{1}{\Lambda^2} \text{ pb} , \] (22)

whereas for the \( c \) quark couplings the interference cross sections are

\[ \sigma^{INT}(pp \rightarrow tW^-) = 0.065 Re[\beta_{ct}] \frac{1}{\Lambda^2} \text{ pb} ; \]
\[ \sigma^{INT}(pp \rightarrow tW^+) = 0.063 Re[\beta_{ct}] \frac{1}{\Lambda^2} \text{ pb} . \] (23)

These cross sections have extremely small coefficients, which is due to a double CKM cancellation, as we have discussed in ref. [10]. This cancellation makes the contribution to the cross section arising from the square of the anomalous diagram the largest, as long as \( \Lambda/\sqrt{|\beta_{ut}|} < 45 \text{ GeV} \).

The new contributions to \( tW \) production are then given by, for the \( u \) couplings,

\[ \sigma^{NEW}(pp \rightarrow tW^-) = 63.4 |\beta_{ut}|^2 \frac{1}{\Lambda^2} \text{ pb} \]
\[ \sigma^{NEW}(pp \rightarrow tW^+) = 18.0 |\beta_{ut}|^2 \frac{1}{\Lambda^2} \text{ pb} , \] (24)

and, for the \( c \) quark couplings,

\[ \sigma^{NEW}(pp \rightarrow tW^-) = 14.2 |\beta_{ct}|^2 \frac{1}{\Lambda^2} \text{ pb} \]
\[ \sigma^{NEW}(pp \rightarrow tW^+) = 11.9 |\beta_{ct}|^2 \frac{1}{\Lambda^2} \text{ pb} . \] (25)

It is interesting to notice that because the SM process occurs mostly through a \( gb \) initial state, and the pdfs for a \( b \) quark and a \( \bar{b} \) quark are essentially identical, there is almost no difference in \( t \) and \( \bar{t} \) production. That is, \( \sigma^{SM}(tW^-) - \sigma^{SM}(\bar{t}W^+) \approx 0 \). However, the interference terms and the new ones receive contributions from all quarks leading to a difference

\[ \sigma^{INT}(tW^-) - \sigma^{INT}(\bar{t}W^+) = 0.09 Re[\beta_{ub}] \frac{1}{\Lambda^2} \text{ pb}; \]
\[ \sigma^{NEW}(tW^-) - \sigma^{NEW}(tW^+) = 45.4 |\beta_{ut}|^2 \frac{1}{\Lambda^4} \text{ pb.} \]  

(26)

Therefore, this asymmetry could be a clear sign of new physics. Moreover, it depends on only one of the anomalous couplings and, if no asymmetry of this kind is observed, a stringent bound could be set on \( \beta_{ut} \).

D. \( th \) production

Finally, we consider the numerical results for associated top plus Higgs boson production. The cross sections depend, of course, on the unknown value of the Higgs mass. As we will observe, though, that dependence is not a strong one. We will consider two values for the Higgs mass, \( m_h = 120 \text{ GeV} \) (the preferred value from the current experimental bounds, and a typical value of a supersymmetric Higgs mass) and \( m_h = 300 \text{ GeV} \). Once again, the results we obtained for production of \( th \) via the \( c \) quark are too small, and we do not show them. Likewise, the pdf suppression of the anti-up and anti-charm quark heavily suppress the production of an anti-top quark and a Higgs boson. We are left with \( th \) production via the \( u \) quark, which, for \( m_h = 120 \text{ GeV} \), reads

\[ \frac{d\sigma(\bar{u}t \rightarrow th)}{dt} = \{5.9 |\alpha_{ut} + \alpha^*_{tu}|^2 + 23.6 \text{Im}[(\alpha_{ut} + \alpha^*_{tu}) \beta_{tu}] + 95.2 (|\beta_{tu}|^2 + |\beta_{ut}|^2)\} \frac{1}{\Lambda^4}, \]  

(27)

and, for \( m_h = 300 \text{ GeV} \), we have

\[ \frac{d\sigma(\bar{u}t \rightarrow th)}{dt} = \{3.2 |\alpha_{ut} + \alpha^*_{tu}|^2 + 25.8 \text{Im}[(\alpha_{ut} + \alpha^*_{tu}) \beta_{tu}] + 52.1 (|\beta_{tu}|^2 + |\beta_{ut}|^2)\} \frac{1}{\Lambda^4}. \]  

(28)

We observe some variation of the coefficients of the \( \alpha \) and \( \beta \) coefficients, the cross section for the larger Higgs mass being, naturally, smaller. However, that variation is not a dramatic one, which is perhaps due to the extremely high center-of-mass energy available in the proton-proton collisions at the LHC.

VI. CONCLUSIONS

In this section we perform a joint analysis of the results obtained so far in this paper and those from our previous works [10, 11]. We have calculated all possible one and two body decays, at the partonic level, originating from a set of strong flavor changing operators satisfying our predefined criteria. We now wish to observe if there is any correlation between the cross sections of top quark production we computed. In what follows we have used the liberty afforded to us by the equations of motion to set \( \beta_{tu} = \beta_{tc} = 0 \), to simplify our calculations.

To investigate the dependence of the cross sections on the values of the anomalous couplings, we generated random values for these, and plotted the cross sections against the branching ratios of the top quark for the decays \( t \rightarrow gu \) and \( t \rightarrow gc \). Our rationale for doing this is a simple one: as was discussed in refs. [12, 16], the top quark branching ratios for these decays may vary by as much as eight orders of magnitude, from \( \sim 10^{-12} \) in the SM to \( \sim 10^{-4} \) for some supersymmetric models. This quantity, then, is a good measure of whether any physics beyond that of the standard model exists. In fig. 4 we show the plot of the cross sections for the processes \( pp \rightarrow t + \text{jet} \) and \( pp \rightarrow t + W \) via a \( u \) quark versus the branching ratio \( \text{BR}(t \rightarrow gu) \). This plot was obtained by varying the constants \( \alpha \) and \( \beta \) in a random way. Each combination of \( \alpha \) and \( \beta \) originates a given branching ratio and a particular value for each cross section. Obviously, another set of points may generate the same value for the branching ratio but a different value for the cross section, which justifies the distribution of values of \( \sigma(pp \rightarrow t + \text{jet}) \) and \( \sigma(pp \rightarrow t + W) \). We chose values of \( \alpha \) and \( \beta \) for which the branching ratio varies between its SM value and the maximum theoretical supersymmetric value it may assume. In fig. 8 we show a similar plot, but for top plus jet and top plus a \( W \) boson production via a \( c \) quark.

It is obvious from both fig. 4 and fig. 8 that if the branching ratio is close to its SM value there is no chance to observe new strong flavor changing physics at the LHC. However, as we approach the larger values of \( 10^{-4} \), the cross section for single top becomes visible and the \( Wt \) cross section approaches 0.1 pb. Notice that the \( tw \) cross section is proportional to only one of the couplings, which makes it a very attractive observable - it may allow us to impose constraints on a single anomalous coupling.

It should be noted that single top production depends also on the contributions of the four fermion operators. Hence, even if the branching ratios \( \text{BR}(t \rightarrow gu(c)) \) are very small, there is still the possibility of having a large single
FIG. 7: Cross sections for the processes $pp \rightarrow t + jet$ (crosses) and $pp \rightarrow t + W$ (stars) via an $u$ quark, as a function of the branching ratio $BR(t \rightarrow g u)$.

FIG. 8: $pp \rightarrow t + jet$ (+) and $pp \rightarrow t + W$ (*) via a $c$ quark as a function of the branching ratio $BR(t \rightarrow g c)$.

top cross section with origin in the four fermion couplings. In figs. 7 and 8 we did not consider this possibility, setting the four-fermion couplings to zero. For a discussion on the four-fermion couplings do see [11].

In fig. 9 we plot the cross sections for $pp \rightarrow t + Z$ and $pp \rightarrow t + \gamma$ via a $u$ quark, versus the branching ratio $BR(t \rightarrow g u)$. The equivalent plot with an internal $c$ quark is similar, but the values for the cross section are much smaller. In this plot we can see that both cross sections are very small in the range of $\{\alpha \beta\}$ considered. These results
imply that their contribution will hardly be seen at the LHC, unless the values for the branching ratio are peculiarly large.

\[ \frac{BR(t \rightarrow g u)}{10^{-10}} \]

\[ \frac{BR(t \rightarrow g u)}{10^{-9}} \]

\[ \frac{BR(t \rightarrow g u)}{10^{-8}} \]

\[ \frac{BR(t \rightarrow g u)}{10^{-7}} \]

\[ \frac{BR(t \rightarrow g u)}{10^{-6}} \]

\[ \frac{BR(t \rightarrow g u)}{10^{-5}} \]

\[ \frac{BR(t \rightarrow g u)}{10^{-4}} \]

\[ \frac{BR(t \rightarrow g u)}{10^{-3}} \]

\[ \frac{BR(t \rightarrow g u)}{10^{-2}} \]

\[ \frac{BR(t \rightarrow g u)}{10^{-1}} \]

The same, in fact, could be said for \( p p \rightarrow t + h \). In fig. 10 we present a plot for this cross section, again as a function of the branching ratio of \( t \rightarrow g u \), for two values of the Higgs mass. We readily see that, even for the smallest allowed SM Higgs mass, the values are very small. The same holds true for the process involving the anomalous couplings of the \( c \) quark.

The smallness of the effects of these operators in the several cross sections holds true, as well, for the top–anti-top channel. In this case, even for a branching ratio \( BR(t \rightarrow g u) \approx 10^{-4} \), the contributions to the cross section \( \sigma(pp \rightarrow t t) \) do not exceed, in absolute value, one picobarn. They may be positive or negative, but always extremely small.

In conclusion, the effective operators we have considered in this paper and references \[ \text{[10, 11]} \] are extremely constrained in their impact on the several channels of top quark production. Namely, figs. 7 through 10 illustrate that, with the exception of the cross section for production of a single top plus a jet, the other channels are expected to have anomalous contributions which are probably too small to be observed at the LHC. Thus, if there are indeed strong flavor changing neutral current effects on the decays of the top quark, the results of this paper show that their impact will be restricted to a single channel, single top plus jet production. It is entirely possible, according to our results, to have an excess in the cross section \( \sigma(pp \rightarrow t + \text{jet}) \) arising from new physics described by the operators we have considered here, at the same time obtaining results for the production of a top quark alongside a gauge and Higgs boson, or for \( tt \) production, which are entirely in agreement with the SM predictions. This reinforces the conclusion of reference \[ \text{[11]} \]: that the cross section for single top plus jet production is an excellent probe for the existence of new physics beyond that of the SM. It is a channel extremely sensitive to the presence of that new physics, and boasts a significant excess in its cross section, whereas many other channels involving the top quark remain unchanged. Nevertheless, we are encouraged by the fact that it may still be possible to use some of these unchanged channels, such as top plus \( W \) production, to constrain the \( \beta \) parameters, through the study of asymmetries such as \( \sigma(pp \rightarrow tW^-) - \sigma(pp \rightarrow tW^+) \).

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FIG. 10: Cross section of the process $p p \rightarrow t + h$ via a $u$ quark versus the branching ratio $BR(t \rightarrow g u)$ for $m_h = 120$ GeV and $m_h = 300$ GeV.

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