Flavour changing strong interaction effects on top quark physics at the LHC

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Abstract. We perform a model independent analysis of the flavour changing strong interaction vertices relevant to the LHC. In particular, the contribution of dimension six operators to single top production in various production processes is discussed, together with possible hints for identifying signals and setting bounds on physics beyond the standard model.

1 Introduction

The top quark is a fascinating particle. It is the heaviest particle discovered so far, and one of the least known \cite{1}. These two facts make it an ideal laboratory for searches of physics beyond the Standard Model (SM). The theoretical implications of new physics may be studied in two ways: by proposing a new theory and computing its effects on observable quantities; or by parameterizing all possible effects of new physics using effective operators of dimensions higher than four, and analysing their impact on measurable variables. The full set of dimension five and six operators was obtained by Buchmuller and Wyler \cite{2}. There has been much work on the effect of these operators on the properties of the top quark. For instance, the authors of the series of papers listed in reference \cite{3,4,5} studied how several types of dimension five and six operators affected the top quark, both at the Tevatron and at the LHC, as well as other types of physics (such as $b\bar{b}$ production). The authors of \cite{6} undertook a detailed study of the $Wtb$ vertex at LHC conditions. New physics contributions for flavour changing neutral currents were studied in \cite{7}. The contributions from four-fermion operators to $t\bar{t}$ production were considered in ref. \cite{8}.

In this work we will focus on the case of dimension six operators with flavour changing interactions involving the top quark, and the processes for which one has, at the LHC, a single quark $t$ being produced. The set of operators we chose includes, after symmetry breaking, the dimension five ones mentioned above, but also new contributions not considered simultaneously before. Our methodology will also be slightly different from that of previous works in this area: whenever possible, we will present full analytical expressions for our results, so that our experimental colleagues at the Tevatron or LHC may use them in their Monte Carlo generators,
to study the sensibility of the experiments to this new physics. The outline of this paper is
as follows: in section 2 we will describe the operator set we have chosen and the reasons that
motivated that choice. We will discuss some subtleties related to the number of independent
operators and the inclusion of four-fermion terms, as well. In section 3 we will compute the
contributions arising from these operators to the top quark width and also to a number of
physical processes of single top production at the LHC. Analysis of these results and conclusions
will be presented in section 4.

2 Dimension six operators

Following the philosophy and notation of ref. [2], we consider the lagrangean
\[ \mathcal{L} = \mathcal{L}^{SM} + \frac{1}{\Lambda} \mathcal{L}^{(5)} + \frac{1}{\Lambda^2} \mathcal{L}^{(6)} + O \left( \frac{1}{\Lambda^3} \right), \]

where \( \mathcal{L}^{SM} \) is the SM lagrangean, \( \Lambda \) is the energy scale at which new physics makes itself manifest
and \( \mathcal{L}^{(5)} \) and \( \mathcal{L}^{(6)} \) are operators of dimension 5 and 6 respectively which, like \( \mathcal{L}^{SM} \), are invariant
under the SM gauge group and built with its fields. Imposing baryon and lepton number
conservation, the term \( \mathcal{L}^{(5)} \) is eliminated. However, after spontaneous symmetry breaking,
dimension five operators will appear in the lagrangean, arising from \( \mathcal{L}^{(6)} \). The total number
of dimension six operators is quite large (see [2] for the full list). It is clearly not practical to
consider them all when studying the effects of new physics, and so some selection criteria are
needed. We assumed from the start that, whatever new physics contributions appear at the
lagrangean \( \mathcal{L} \) as dimension six operators, low-energy phenomenology will not be substantially altered by them. By “low-energy” we mean any process occurring below the LEP energies.
Another selection criterion was that we are only interested in operators involving a single top quark - we have in mind the study of single top quark production, so operators with more than a top quark are of no interest to us. Finally, we restricted ourselves to a particular type of new physics: flavour changing vertices involving the top quark and the strong interactions, namely, vertices of the form \( g t \bar{c} \) or \( g t \bar{u} \). It is a choice like any other, but motivated by two arguments: first, the LHC environment is ideally suited to study the strong sector of the SM (much like LEP was designed to study its electroweak sector) so, if a deviation from normal QCD vertices exists, it will likely be possible to discover it at the LHC; second, some of the operators we will be considering - the so-called chromomagnetic operators - are generated in the context of the SM at high orders, but their effects are expected to be too small to be measured. However, they appear in more general theories, such as supersymmetric models or theories with multiple Higgs doublets (see, for example, [2]).

These criteria reduce the number of operators involving gluons to just two, namely
\[ \mathcal{O}_{uG} = i \frac{\alpha_{ij}}{\Lambda^2} \left( \bar{u}_R^i \gamma^\mu \gamma^\nu \sigma^\mu \sigma^\nu \bar{u}_R^j \right) G^a_{\mu\nu}, \]
\[ \mathcal{O}_{uG\phi} = i \frac{\beta_{ij}}{\Lambda^2} \left( \bar{q}_L^i \gamma^\mu \gamma^\nu \sigma^\mu \sigma^\nu \bar{q}_L^j \right) \tilde{\phi} G^a_{\mu\nu}, \]

where \( u_R^i \) is an up quark right spinor (an SU(2) singlet, therefore), \( q_L^i \) a left quark doublet spinor, \( \phi \) the charge conjugate of the Higgs doublet, \( G^a_{\mu\nu} \) the gluon tensor. \( \alpha_{ij} \) and \( \beta_{ij} \) are complex couplings and the remaining quantities are well-known. The indices \( (i, j) \) on the spinors are flavour indices, indicating to which generation each one belongs. As an example of the application of our criteria, consider another \textit{a priori} possible operator from [2],
\[ \mathcal{O}_{qG} = i \left( \bar{q}_L^i \gamma^\mu \gamma^\nu q_L^j \right) G^a_{\mu\nu}. \]
In terms of up and down quarks this operator becomes
\[
\mathcal{O}_{qG} = i \left( \bar{u}_L^i \lambda^a \gamma^\mu D^\nu u_L^j + \bar{d}_L^i \lambda^a \gamma^\mu D^\nu d_L^j \right) G^a_{\mu\nu} .
\] (4)

The second term of this operator gives an interaction involving only down type quarks. It would therefore have direct implications on low-energy physics (such as bottom physics, since one of the doublets, per one of our criteria, must belong to the third generation). As such, its size is already immensely constrained by the wealth of existing data. Hence, the relevance of this operator to new phenomena will be very limited and we do not consider it. It might also seem bizarre that we are considering the second of the operators of equation (2), as it has a Higgs field in it, and our criteria did not include phenomenology of scalar fields. However, after spontaneous symmetry breaking, the neutral component of \( \phi \) acquires a non-zero vev (\( \phi_0 \rightarrow \phi_0 + \text{v} \), with \( \text{v} = 246/\sqrt{2} \text{ GeV} \)) and the operator decomposes into several pieces containing Higgs scalars and one without them, namely
\[
\mathcal{O}_{uG\phi} \rightarrow \beta_{ij} \frac{v}{\Lambda^2} \left( \bar{u}_L^i \lambda^a \sigma^\mu\nu u_R^j \right) G^a_{\mu\nu} .
\] (5)

The dimension 6 operator thus becomes identical to a dimension 5 chromomagnetic one.

One of the generation indices on the operators (2) must correspond to the third family, and so our lagrangean for new physics becomes
\[
\mathcal{L} = \alpha_{tu} \mathcal{O}_{tu} + \alpha_{ut} \mathcal{O}_{ut} + \beta_{tu} \mathcal{O}_{tu\phi} + \beta_{ut} \mathcal{O}_{ut\phi} + \text{h.c.}
\]
\[
= \frac{i}{\Lambda^2} \left[ \alpha_{tu} \left( \bar{t}_R^i \lambda^a \gamma^\mu D^\nu u_R^j \right) + \alpha_{ut} \left( \bar{u}_R^i \lambda^a \gamma^\mu D^\nu t_R^j \right) \right] G^a_{\mu\nu} +
\]
\[
+ \frac{v}{\Lambda^2} \left[ \beta_{tu} \left( \bar{t}_L^i \lambda^a \sigma^\mu\nu u_R^j \right) + \beta_{ut} \left( \bar{u}_L^i \lambda^a \sigma^\mu\nu t_R^j \right) \right] G^a_{\mu\nu} + \text{h.c.} ,
\] (6)

for the vertices \( g \tilde{t} u \), \( g \tilde{u} t \), and an analogous lagrangean (with different coupling constants, clearly) for \( g \tilde{c} t \) and \( g \tilde{t} \tilde{c} \). Then, the Feynman rules for the anomalous vertices \( g \tilde{t} u \) and \( g \tilde{t} \tilde{u} \) are

\[
\text{1Gluon momenta are defined as incoming, quark momenta follow the sense of the arrows in the figure.}
\]
\[
\lambda^a \frac{1}{\Lambda^2} \left\{ \gamma_{\mu} \gamma_R \left( \alpha_{ij} p_{\mu} + \alpha^*_{ij} p'_{\mu} \right) + v \sigma_{\mu\nu} \left( \beta_{ij} \gamma_R + \beta^*_{ij} \gamma_L \right) \right\} \times \left( k^\mu g^{\nu\alpha} - k'^\mu g^{\mu\alpha} \right)
\] (7)

If one uses the fermionic equations of motion and integrates by parts, one obtains the following relations between operators:

\[
O^\dagger_{ut} = O_{tu} - \frac{i}{\Lambda^2} (\Gamma^\dagger_{ut} O_{ut\phi} + \Gamma_{ut} O_{tu\phi})
\]

\[
O^\dagger_{ut} = O_{tu} - i g_s \bar{t} \gamma_{\mu} \gamma_R \lambda^a u \sum_i (\bar{u}_i \gamma_{\mu} \gamma_R \lambda^a u_i + \bar{d}_i \gamma_{\mu} \gamma_R \lambda^a d_i) ,
\] (8)

where \( \Lambda^a \) are the Yukawa coupling matrices associated with the \( u \) quarks and \( g_s \) is the strong gauge coupling. In the second of these equations we see the appearance of four-fermion operators. Then, as long as one considers these extra contributions, we obtain two relations between \( O_{tu} \), \( O_{ut} \), \( O_{tu\phi} \) and \( O_{ut\phi} \) and so we can set two of the \( \{ \alpha, \beta \} \) constants to zero. If, however, we leave out the four fermion contributions, only the first of eqs. (8) is valid and therefore only one of \( \{ \alpha_{tu}, \alpha_{ut}, \beta_{tu}, \beta_{ut} \} \) may be set to zero. These considerations are important for more elaborate calculations than those that will be undertaken in this paper - for instance, the process \( gg \to t\bar{c} \). In the present work, for reasons to be understood later, the inclusion of the four-fermion operators shown in eq. (8) has no effect on the results and thus we do not consider them. Therefore, we use only one of the equations (8) and reduce the dependence of our expressions to three complex parameters per generation, setting \( \beta_{tu} \) and \( \beta_{tc} \) to zero. Nevertheless, we will present our full results, for completeness. Finally, equation (8) shows the advantage of having chosen this particular operator set. With it, we may take advantage of the gauge invariance of the theory, which is gives us relations between the several operators, allowing us to reduce the number of independent parameters.

3 Effects of the new physics on top quark observables

The vertices \( g\bar{t}c \) and \( g\bar{t}u \) constitute new possible decay channels for the top quark. Hence they will contribute to the top’s width, and a straightforward calculation shows that

\[
\Gamma(t \to ug) = \frac{m_t^3}{12\pi^4 \Lambda^4} \left\{ m_t^2 \left| \alpha_{ut} + \alpha^*_{tu} \right|^2 + 16 v^2 \left( |\beta_{tu}|^2 + |\beta_{ut}|^2 \right) \right. + \left. 8 v m_t \text{Im} \left[ (\alpha_{ut} + \alpha^*_{tu}) \beta_{tu} \right] \right\}
\] (9)

where we assumed that all of the quark masses, except the top’s, are zero. Having performed the full calculation, we verified that this is an excellent approximation. This will allow us to impose limits on the values of the new couplings. The main decay channel for the top quark is \( t \to bW \), the tree level decay width being given, in the SM, by

\[
\Gamma(t \to bW) = \frac{g^2}{64\pi} |V_{tb}|^2 \frac{m_t^3}{M_W^2} \left( 1 \right. + \left. \frac{M_W^2}{m_t^2} - 2 \frac{M_W^2}{m_t^2} \right) \left( 1 \right. - \left. \frac{M_W^2}{m_t^2} \right) ,
\] (10)

assuming \( m_b = 0 \). After QCD corrections and using \( m_t = 175 \text{ GeV} \), this becomes \( \Gamma(t \to bW) = 1.42 |V_{tb}|^2 \text{ GeV} \). (11)

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\]
Taking $V_{tb} \sim 1$ this will be, to very good approximation, the total width of the top quark. Then,

$$\frac{\Gamma(t \rightarrow ug)}{\Gamma_t} = \frac{10^{-3}}{\Lambda^4} \left\{ 3.07 |\alpha_{ut} + \alpha_{tu}^*|^2 + 48.53 \left( |\beta_{tu}|^2 + |\beta_{ut}|^2 \right) \right. +
$$

$$24.41 \text{Im} \left[ (\alpha_{ut} + \alpha_{tu}^*) \beta_{tu} \right] \right\}, \quad (12)$$

with $\Lambda$ expressed in TeV. Let us now assume that these flavour changing decays are not observed at the LHC. If $L$ is the degree of precision with which the top quark width will be known at the LHC then we will have $\Gamma(t \rightarrow qg)/\Gamma_t \leq L$, where $q = u, c$. Since $\beta_{tu}$ and $\beta_{tc}$ were set to zero, upper bounds on the $\{\alpha, \beta\}$ couplings are obtained as $|\alpha_{qt} + \alpha_{qg}^*|/\Lambda^2 \leq 18.05 \sqrt{L}$ TeV$^{-2}$ and $|\beta_{qt}|/\Lambda^2 \leq 4.51 \sqrt{L}$ TeV$^{-2}$. Assuming that it will be possible to measure $\Gamma_t$ with 10% precision at the LHC (which seems reasonable, considering preliminary studies already done \[11\]), these upper limits become, respectively, 5.71 and 1.44 TeV$^{-2}$. The authors of ref. \[12\] used CDF data \[12\] to impose the limit $BR(t \rightarrow cg) < 0.45$ and obtaining from this $|\beta_{ct}|/\Lambda^2 \leq 3.1$ TeV$^{-2}$, which is in agreement with the result we obtained above.

The flavour changing vertices $g \bar{t} c$, $g \bar{t} u$, \ldots
duce several new channels for single top production. The first, and as we will see, the most important, is direct top production, in which a gluon and a quark $u/c$ are extracted from the colliding protons, as is illustrated in fig. 1. With a single particle in the final state and the colliding protons having large parton density functions (pdf’s), this channel is favoured over all others. Due to the one-particle final state, the cross section calculation is very simple, and we obtain

$$\sigma(pp \rightarrow t) = \sum_{q = u,c} \frac{\pi m_t^2}{12 \Lambda^4} \left\{ m_t^2 |\alpha_{qt} + \alpha_{qg}^*|^2 + 16 v^2 \left( |\beta_{qt}|^2 + |\beta_{qg}|^2 \right) \right. +
$$

$$8 v m_t \text{Im} \left[ (\alpha_{qt} + \alpha_{qg}^*) \beta_{qt} \right] \left\} \int_{m_t^2/E_{CM}^2}^{1} \frac{2 m_t}{E_{CM}^2 x_1} f_g(x_1) f_q(m_t^2/(E_{CM}^2 x_1)) \, dx_1 \right. \quad (13)$$

where $E_{CM}$ is the proton-proton center of mass energy (14 TeV for the LHC), $f_g$ and $f_q$ are the parton density functions for the gluon and the quark $q$ (an up or charm quark). We see that the partonic cross section factorizes from the pdf integration and, using the results from eq. (9), we obtain

$$\sigma(pp \rightarrow t) = \sum_{q = u,c} \Gamma(t \rightarrow qg) \frac{\pi^2}{m_t^4} \int_{m_t^2/E_{CM}^2}^{1} \frac{2 m_t}{E_{CM}^2 x_1} f_g(x_1) f_q(m_t^2/(E_{CM}^2 x_1)) \, dx_1 . \quad (14)$$

If we assume that the branching ratio $BR(t \rightarrow bW)$ is approximately 100% then, using the above results for the top width at the SM, we may express $\Gamma(t \rightarrow qg)$ as $1.42 |V_{tb}|^2 BR(t \rightarrow qg)$. We numerically performed the integration in eq. (14) using the CTEQ6M parton density functions \[13\] setting the factorization scale equal to $m_t^2$. The total cross section is

$$\sigma(pp \rightarrow t) = [10.5 BR(t \rightarrow ug) + 1.7 BR(t \rightarrow cg)] |V_{tb}|^2 10^4 \text{ pb} . \quad (15)$$

The bigger contribution from the $u$ quark channel stems from the fact the pdf for that quark is larger than the $c$ quark’s. For anti-top production, we obtain

$$\sigma(pp \rightarrow \bar{t}) = [2.8 BR(\bar{t} \rightarrow \bar{u}g) + 1.7 BR(\bar{t} \rightarrow \bar{c}g)] |V_{tb}|^2 10^4 \text{ pb} . \quad (16)$$

This choice of factorization scale produces smaller values for the cross section than setting $\mu_F$ equal to the partonic center-of-mass energy (see, for instance, reference \[15\]). For comparison between SM and anomalous cross sections, though, this has no bearing on our conclusions.
In terms of the couplings $\alpha$ and $\beta$, the cross sections become

$$\sigma^{(u)}(pp \to t) = \frac{1}{\Lambda^4} \left\{ 321 |\alpha_{ut} + \alpha_{tu}^*|^2 + 5080 \left( |\beta_{tu}|^2 + |\beta_{ut}|^2 \right) + 2556 \operatorname{Im} [(\alpha_{ut} + \alpha_{tu}^*) \beta_{tu}] \right\} \text{ pb}$$

for top production via an $u$ quark, and, for the $c$ quark contribution,

$$\sigma^{(c)}(pp \to t) = \frac{1}{\Lambda^4} \left\{ 50 |\alpha_{ct} + \alpha_{tc}^*|^2 + 796 \left( |\beta_{tc}|^2 + |\beta_{ct}|^2 \right) + 400 \operatorname{Im} [(\alpha_{ct} + \alpha_{tc}^*) \beta_{tc}] \right\} \text{ pb}$$

Using the expressions [15] and [16] one arrives, in a straightforward manner, to the expressions for anti-top production cross sections. A crude upper bound on the single top cross sections may be obtained if we consider the most unfavourable case in which the decays $t \to u g$, etc, have no visible effect on the top width. Using the estimates of the previous paragraphs and considering both branching ratios identical, the cross section is bounded as $|\sigma(pp \to t) + \sigma(pp \to \bar{t})| \leq 1.7 \times 10^5 L \text{ pb}$. With 5% precision on the measurement of the top’s width at the LHC, this would give an upper bound of 8350 pb for the cross section of this process. Single top quark production has been extensively studied [14]. The value we obtained above is quite large, if compared with the expected value for single top production at the LHC at NLO within the SM framework, roughly $319.7 \pm 19.3 \text{ pb}$ [15]. Alternatively, we can ask ourselves what the values of the couplings $\{\alpha, \beta\}$ should be so that direct top production via flavour changing gluon interactions is not observed. This means that the anomalous contribution is such that $\sigma_{AN} \leq \operatorname{Max}\{\Delta_T, \Delta_E\}$, where $\Delta_T$ and $\Delta_E$ are the theoretical and experimental uncertainties on the cross section. Bounds on the couplings are then obtained considering each one separately and, using the above mentioned value $\Delta_T = 19.3 \text{ pb}$, we get (in units of $\text{TeV}^{-2}$)

$$\frac{|\alpha_{ut} + \alpha_{tu}^*|}{\Lambda^2} \leq \operatorname{Max} \left( 0.22, 0.049 \sqrt{\Delta_E} \right), \quad \frac{|\beta_{ut}|}{\Lambda^2} \leq \operatorname{Max} \left( 0.055, 0.013 \sqrt{\Delta_E} \right)$$

$$\frac{|\alpha_{ct} + \alpha_{tc}^*|}{\Lambda^2} \leq \operatorname{Max} \left( 0.44, 0.10 \sqrt{\Delta_E} \right), \quad \frac{|\beta_{ct}|}{\Lambda^2} \leq \operatorname{Max} \left( 0.11, 0.025 \sqrt{\Delta_E} \right).$$

The expected experimental error is $\Delta_E = 16 \text{ pb}$ (5% of the total cross section [11]), therefore the current theoretical error dominates. These results are to be compared with the limits predicted in ref. [3] for the $\beta$ couplings: from a study involving the simulation of the possible LHC backgrounds and for 10 $fb^{-1}$ of data, they obtained, when converted into our notation, $|\beta_{ut}|/\Lambda^2 \leq 0.021 \text{ TeV}^{-2}$ and $|\beta_{ct}|/\Lambda^2 \leq 0.046 \text{ TeV}^{-2}$.

All observables we have considered thus far depend on $\Lambda^{-4}$. For a large energy scale, then, the effects of new physics will be very small. There are however processes for which the scale dependence is $\Lambda^{-2}$, namely the interference terms between the operators [2] and the SM process for which a single top quark is produced. In the SM at tree-level, the Feynman diagrams for single top production are shown in figures [2] and [3] for the reactions $pp \to t j$ and $pp \to t W$. Most of these diagrams will be suppressed by elements of the CKM matrix, and also by small pdf contributions from the incoming quarks. Using the CTEQ6M structure functions and considering all possible combinations of incoming partons we obtain, for the tree-level cross sections, the values listed in table 1. The SM results are different from those quoted in references [15], though not by much. This is due to the size of the QCD corrections and also to the relatively
low factorization scale we chose. It is easy, with our anomalous vertices, to build diagrams that will interfere with the SM processes; namely, those of figures 4 and 5. Notice how, due to the standard gluon vertex, the down-type quarks always belong to the same generation. The interference with the SM diagrams of fig. 2 will thus be restricted to that case. The calculation of the interference terms gives, for the differential partonic cross sections, \((q = u, c)\)

\[
\frac{d\sigma_{1j}^{ij}}{dt} = K \frac{s + t - m_t^2}{s^2(t - m_W^2)} \frac{\Re(\beta qt)}{\Lambda^2}
\]

\[
\frac{d\sigma_{2j}^{ij}}{dt} = \frac{s(s + t + m_W^2 - m_t^2)}{s^2} \frac{\Re(\beta qt)}{\Lambda^2}
\]

\[
\frac{d\sigma_{3j}^{ij}}{dt} = K \frac{s + t - m_t^2}{s^2} \frac{\Re(\beta qt)}{\Lambda^2}
\]

\[
\frac{d\sigma_1^{W}}{dt} = \frac{3K}{32m_W^2} \frac{(t - 2m_W^2)}{t(t - m_t^2)s^2} \frac{\Re(\beta qt)}{\Lambda^2}
\]

\[
\frac{d\sigma_2^{W}}{dt} = \frac{3K}{32m_W^2} \frac{2m_t^4 - t(m_W^2 - 2s - t) + 2m_t^2(t - 2m_W^2)}{t^2} \frac{\Re(\beta qt)}{\Lambda^2}
\]

where \(K\) is a factor including all couplings, spin and colour factors and other constants, given by \(^3\)

\[
K = \frac{8}{9} \sqrt{\frac{2\alpha_S}{\pi}} G_F m_t m_W^2 v |V_{qi} V_{ti}|,
\]

and we have used, for the CKM matrix elements \(V_{ij}\), the central values of the PDG data \(^{16}\). Notice how these results depend solely on one of the couplings - this is due to the particular chiral structure of the anomalous operators \(^2\), but stems also from the fact that we are considering all quark masses to be zero, except the top's. If we relax this approximation (which means, in effect, taking into account the diagrams in fig. 2 with Goldstone bosons \(G^\pm\)), terms in \(\alpha_{ut}\), \(\alpha_{tu}\) and \(\beta_{tu}\) will appear in \(^{20}\). However, direct calculation of those terms shows that they are much smaller than those shown in \(^{20}\) (at least 50 times smaller, but typically much inferior to that). It is interesting to observe that these particular processes probe one single anomalous coupling, thus allowing us, in theory, to determine it independently of the others. Summing all possible contributions from the incoming partons and performing the integrations on both \(t\) and the pdf’s, we obtain the results shown in table 1. As we can see, the anomalous cross sections are typically much smaller than the expected SM values. The largest deviation from the SM predictions occurs for the production of a top and bottom quarks, for which we have

\[
\frac{\sigma^{AN}(pp \to tb)}{\sigma^{SM}(pp \to tb)} = \left[ 1.0 \frac{\Re(\beta_{ud})}{\Lambda^2} + 2.3 \frac{\Re(\beta_{cd})}{\Lambda^2} \right] \%,
\]

Even in this more favourable case, we see that the deviation from the SM is at the most a few percent of the values of the couplings. We would expect that QCD corrections to these calculations would not affect significantly this ratio, as both the SM and anomalous cross sections should have similar contributions from higher order corrections. The reason for the smallness of these numbers is easy to understand: the diagrams for which one has interference are suppressed by off-diagonal CKM matrix elements. There are two ways for this suppression to occur: for instance, the process \(d\bar{d} \to t\bar{u}\) is suppressed quite heavily by the CKM product \(V_{td} V_{ud}\) - even though \(V_{ud}\) is a diagonal CKM element, \(V_{td}\) represents a two family “jump” and as such is miniscule. The second possibility is, for example, the process \(u s \to t s\), which is suppressed by \(V_{ts} V_{us}\) - the product of these two CKM elements, each of them representing a one family “jump”, is also extremely small.

\(^3\)We left out the conversion factor \(0.389 \times 10^9\) pb TeV\(^3\)
| Process | SM cross section (pb) | Anomalous cross section (pb) |
|---------|------------------------|-----------------------------|
| $pp \rightarrow tj$ | 151.0 | $0.80 \frac{\text{Re}(\beta_{ul})}{\Lambda^2} + 0.10 \frac{\text{Re}(\beta_{ct})}{\Lambda^2}$ |
| $pp \rightarrow tb$ | 4.8 | $0.07 \frac{\text{Re}(\beta_{ul})}{\Lambda^2} + 0.09 \frac{\text{Re}(\beta_{ct})}{\Lambda^2}$ |
| $pp \rightarrow \bar{t}j$ | 89.7 | $0.08 \frac{\text{Re}(\beta_{ul})}{\Lambda^2} + 0.08 \frac{\text{Re}(\beta_{ct})}{\Lambda^2}$ |
| $pp \rightarrow \bar{t}b$ | 3.0 | $0.01 \frac{\text{Re}(\beta_{ul})}{\Lambda^2} + 0.09 \frac{\text{Re}(\beta_{ct})}{\Lambda^2}$ |
| $pp \rightarrow tW$ | 31.1 | $0.04 \frac{\text{Re}(\beta_{ul})}{\Lambda^2} + 0.08 \frac{\text{Re}(\beta_{ct})}{\Lambda^2}$ |
| $pp \rightarrow \bar{t}W$ | 31.1 | $0.03 \frac{\text{Re}(\beta_{ul})}{\Lambda^2} + 0.08 \frac{\text{Re}(\beta_{ct})}{\Lambda^2}$ |

Table 1: Results for the interference cross sections. The processes $pp \rightarrow tj, pp \rightarrow \bar{t}j$ refer to production of non-bottom quarks alongside the top.

As for the four-fermion operators, besides those shown in equation (8), there are a priori two other types that we could consider in this calculation [2]. Using the same criteria applied to the gluonic operators, we can classify them as:

- **Type 1**,\(\mathcal{O}_{u1} = \frac{g_s \gamma_{u1}}{\Lambda^2} \left(\bar{t} \gamma^\mu \gamma_R u \right) \left(\bar{q} \gamma_\mu \gamma_R q \right) + \text{h.c.} \), \hspace{1cm} (23)

  where $q$ is any given quark, other than the top; these are the operators that appear in equation (8).

- **Type 2**,\(\mathcal{O}_{u2} = \frac{g_s \gamma_{u2}}{\Lambda^2} \left[\left(\bar{t} \gamma^\mu \gamma_L u \right) \left(\bar{d}'' \gamma^\mu \gamma_R d \right) + \left(\bar{t} \gamma^\mu \gamma_L d \right) \left(\bar{d}'' \gamma^\mu \gamma_R u \right) \right] + \text{h.c.} \), \hspace{1cm} (24)

  with down and up quarks from several possible generations, excluding the top once more;

- **Type 3**,\(\mathcal{O}_{u3} = \frac{g_s \gamma_{u3}}{\Lambda^2} \left[\left(\bar{t} \gamma^\mu \gamma_R u \right) \left(\bar{b} \gamma^\mu \gamma_R d \right) - \left(\bar{t} \gamma^\mu \gamma_R d \right) \left(\bar{b} \gamma^\mu \gamma_R u \right) \right] + \text{h.c.} \), \hspace{1cm} (25)

  and also,\(\frac{g_s \gamma^a}{\Lambda^2} \left[\left(\bar{t} \gamma^\mu \gamma_L u \right) \left(\bar{d}'' \gamma^\mu \gamma_L d \right) - \left(\bar{t} \gamma^\mu \gamma_L d \right) \left(\bar{d}'' \gamma^\mu \gamma_L u \right) \right] + \text{h.c.} \). \hspace{1cm} (26)

All of these operators could, in principle, interfere with the SM processes of single top production shown in fig. 2. However, due to their chiral structure and the fact that we considered all of the quarks other than the top to have zero mass, a straightforward calculation shows that their interference terms with the SM diagrams are zero.
4 Conclusions

In this paper we analysed the simpler processes affected by our choice of dimension six operators, which modeled, in an effective way, the possibility of strong interactions provoking flavour violations involving the top quark. These induce new decay possibilities for the top quark, namely \( t \to c g \) or \( t \to u g \). Precision measurements of the top quark width at the LHC will be able to set stringent bounds on the new anomalous couplings. The existence of these new vertices gives rise to a new process for production of a single top quark - direct top production - a process for which the calculated cross section is, \textit{a priori}, quite large. At this point, detailed simulations of this process at the LHC are necessary to clarify whether it is possible to distinguish it from other processes of single top production. We anticipate a possible complication, the fact that the top quark is produced with virtually no transverse momentum and a great longitudinal momentum component. The results of its decay, which will then interact with the detectors, might therefore have also a strong \( p_L \) component and concentrate on the small angle region, making their observation more difficult.

There are interference terms between the Feynman diagrams that describe single top production in the SM with two quarks in the initial state and those with a single anomalous gluon coupling. The resulting cross sections depend on \( \Lambda^{-2} \), instead of \( \Lambda^{-4} \) as in the previous examples. These should therefore be a privileged source of bounds on the new operators, specially given that these quantities depend only on one of the new couplings. However, the values obtained for the interference cross sections are very small compared with the expected SM results, or indeed with the other anomalous cross sections obtained, which depend on \( \Lambda^{-4} \). The reason for this, we found, was a strong CKM suppression. Again, a full simulation of this process at the LHC is needed, to investigate whether some kinematical cuts might be capable of extricating the interference terms from the remaining cross section, but that seems unlikely at this stage. The four-fermion operators did not have any importance to the results presented here. They have no contribution to the top’s decay into a light quark and a gluon, nor to the process of direct top production. Further, their interference with the SM processes of single top production is zero.

Finally, the interest of the operator set we considered in this work is not exhausted on the simpler processes of single top production we studied here. They have a strong impact on processes such as \( g g \to t \bar{u}, g u \to t g \) which, due to the large gluon pdf, one can expect to be very significant at the LHC. In ref. \cite{17} we will study these processes and there, the four-fermion operators will have an important role to play as well.

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Figure 1: Direct top production at the LHC via flavour changing gluon vertices

Figure 2: \( t \) and \( s \) channels for single top production according to the Standard Model, at tree-level.
Figure 3: Associated single top production according to the Standard Model, at tree-level.

Figure 4: Diagrams with an anomalous gluon vertex that interfere with those of fig. 2.

Figure 5: Diagram with an anomalous gluon vertex that interfere with those of fig. 3.