Probing Top Anomalous Couplings at the Tevatron and the Large Hadron Collider

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Abstract

Chromomagnetic and chromoelectric dipole interactions of the top quark are studied in a model independent framework. Limits are set on the scale of new physics that might lead to such contributions using available Tevatron data. Prospects at the LHC are reviewed.

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1 Introduction

The Standard Model (SM) embodies our current understanding of the fundamental constituents of the universe and their interactions. This model is well tested up to the energy scale of a few hundred GeVs and experiments have shown many of its predictions to be astounding accuracy. However, in spite of this stupendous success, certain questions remain unanswered. Among them are questions regarding the mechanism responsible for giving masses to fundamental particles.

Within the SM, the generation of masses is explained by the spontaneous breaking of the electroweak symmetry and the Higgs mechanism. However, no such Higgs scalar has been found yet. Furthermore, the SM fails to explain why, even though the underlying mechanism is the same, there is a difference of six orders of magnitude between the masses of the lightest fermion (electron) and the heaviest one (top quark). The Yukawa couplings of the fermions with the Higgs are parameters in the SM and cannot be explained or predicted by the theory.

It is possible to bring about electroweak symmetry breaking (EWSB) without introducing a new fundamental field such as the Higgs \(H\). What is important to note is that, any theory that provides a mechanism for generation of masses must have a large coupling to the top quark. Consequently, even in the absence of an actual observation of the Higgs boson, experiments with the top may be used to probe the EWSB mechanism. Not only is the top

\[1\] We do not consider the neutrinos here as even the nature of their mass term is, as yet, uncertain. Were they to be purely Dirac ones, the hierarchy worsens.
quark the heaviest particle in the SM with a mass $\sim 175$ GeV, its mass differs widely from those of the other fermions (the next heaviest is the $b$-quark with a mass of 4.2 GeV [2]). This prompts us to examine whether the top quark has couplings different from and in addition to those of the other quarks.

Mechanisms for dynamic breaking of electroweak symmetry through the formation of $t\bar{t}$ condensates are discussed in Ref. [3]. Such scenarios require the top quark to have non-QCD ‘strong’ interactions and give rise to interaction terms such as $\bar{t}t\bar{t}$ and $\bar{t}tbb$ which contribute to the higher order corrections to the $ttg$ vertex.

At the Tevatron proton-antiproton collider at Fermilab, the mass and charge of the top quark have been measured reasonably well [4, 5]. But, the high threshold for top production has meant that its couplings are still not well measured. The possibility that the top quark has anomalous couplings is still open. Once the Large Hadron Collider at CERN comes into operation, precise measurements of top couplings and detection of anomalous couplings will become possible.

Various anomalous couplings of the top have been discussed in Ref. [6]. Of these, the ones that pertain to the QCD-sector would be expected to modify the production rates significantly and, thus, to be probed during the early phase of the LHC. On the other hand, modifications of the electroweak couplings would play only a rather sub-dominant role in $t\bar{t}$ production and it is the decay patterns that would be more sensitive to them. Consequently, a search for the latter type would require both a thorough understanding of the detector as well as the accumulation of large statistics. Given this, we concentrate here on the former set.

Large anomalous couplings may arise in a plethora of models, most prominent being the aforementioned scenarios of dynamical EWSB. Contributions may also arise from theories with additional heavy fermions that couple to the top. Examples include, but are not limited to, Little Higgs models [7,8] or models with extra spacetime dimensions [9,10,11]. Another possibility is the SM augmented by color-triplet or color-sextet scalars that have Yukawa couplings with the top-quark [12].

The lowest-dimensional anomalous coupling of the top with the gluon can be parametrized by extra terms in the interaction Lagrangian of the form

$$\mathcal{L}_{\text{int}} \supset \frac{g_8}{\Lambda} F_a^{\mu\nu} i\sigma_{\mu\nu} (\rho + i \rho' \gamma_5) T_a t$$

where $\Lambda$ denotes the scale of the effective theory. While $\rho$ represents the anomalous chromomagnetic dipole moment of the top, $\rho'$ indicates the presence of a $(CP$-violating) chromoelectric dipole moment. Within the SM, $\rho'$ is non-zero only at the three-loop level and is, thus, tiny. $\rho$, on the other hand, receives a contribution at the one-loop level and is $\mathcal{O}(\alpha_s/\pi)$ for $\Lambda \sim m_t$. The evidence for a larger $\rho$ or $\rho'$ would thus be a strong indicator of new physics lurking nearby. Whereas both $\rho$ and $\rho'$ can, in general, be complex, note that any imaginary part thereof denotes absorptive contributions and would render the Lagrangian non-Hermitian. We desist from considering such a possibility.

The phenomenological consequences of such anomalous couplings have been considered earlier in Ref. [13]. However, we reopen the issue in light of the improved measurements of top quark mass and $t\bar{t}$ cross-section.
2 Analytic Calculation

The inclusion of a chromomagnetic moment term leads to a modification of the vertex factor for the usual $ttg$ interaction to $ig_s[\gamma^\alpha + (2 i \rho/\Lambda) \sigma^{\alpha\beta} k_{\mu}] T^\alpha$, where $k$ is the momentum of the gluon coming into the vertex. An additional quartic interaction involving two top quarks and two gluons is also generated with the corresponding vertex factor being $(2 i g_s^2 \rho/\Lambda) f_{abc} \sigma^{\alpha\beta} T^c$. The changes in the presence of the chromoelectric dipole moment term are analogous, with $\rho$ above being replaced by $(i \rho' \gamma_5)$.

At a hadronic collider, the leading order contributions to $t\bar{t}$ production come from the $q\bar{q}$ to $t\bar{t}$ and $gg$ to $t\bar{t}$ sub-processes. Summing (averaging) over spin and color degrees of freedom and defining $\Theta_{\pm} = 1 \pm \beta^2 \cos^2 \theta$ where $\beta = \sqrt{1 - 4m_t^2/s}$ and $\theta$ are, respectively, the velocity and scattering angle of the top in the parton center-of-mass frame, the differential cross-sections can be expressed as

\[
\left( \frac{2\hat{s}}{\pi\alpha_s^2\beta} \right) \frac{d\hat{\sigma}_{q\bar{q}}}{d\cos\theta} = \frac{2}{9} \Theta_+ + \frac{8 m_t^2}{9 \hat{s}} \left( 32 \rho m_t \left( \frac{8 t \Theta_+ + 4 m_t^2}{9 \Lambda^2} \right) + \frac{8 \rho^2}{9 \Lambda^2} \left( \frac{\hat{s} \Theta_- + 4 m_t^2}{9 \Lambda^2} \right) \right),
\]

\[
\left( \frac{2\hat{s}}{\pi\alpha_s^2\beta} \right) \frac{d\hat{\sigma}_{gg}}{d\cos\theta} = \frac{2}{3\Theta_-} \left( 1 + 4 \frac{m_t^2}{s} + \frac{m_t^2}{s^2} \right) - \left( \frac{1}{3} + \frac{3}{16} \Theta_+ + \frac{3 m_t^2}{2s} + \frac{16 m_t^2}{3s} \Theta_+ \right)
\]

\[
+ \frac{\rho m_t}{\Lambda} \left( -3 + \frac{16}{3\Theta_-} \right) + \frac{\rho^2}{\Lambda^2} \left[ \frac{7}{3} \hat{s} + m_t^2 \right] \left\{ -6 + \frac{34}{3\Theta_-} \right\}
\]

\[
+ \frac{\rho^2}{\Lambda^2} \left[ \frac{7}{3} \hat{s} + \frac{2m_t^2}{\Theta_-} \right]
\]

\[
+ \frac{\rho}{\Lambda} \left( \frac{\rho^2}{\Lambda^2} + \frac{\rho^2}{\Lambda^2} \right) m_t \left( \frac{28}{3} \hat{s} - \frac{20}{3m_t^2} \hat{s} \right)
\]

\[
+ \frac{4}{3} \left( \frac{\rho^2}{\Lambda^2} + \frac{\rho^2}{\Lambda^2} \right)^2 \left( \frac{\hat{s}^2 \Theta_- - m_t^2 \hat{s} + \frac{4}{\Theta_-} m_t^2}{} \right).
\]

In each case, the first line refers to the SM result (we do not exhibit the electroweak contribution for the $q\bar{q}$-initiated process as it remains unaltered) and the rest encapsulate the consequences of the anomalous dipole moments.

It should be noted that there are neither terms with an odd power of $\rho'$ nor do the expressions show any possibility of a forward-backward asymmetry, even in the presence of a non-zero $\rho'$. The two facts are inter-related. Clearly, terms odd in $\cos \theta$ can appear only in terms odd in $\rho'$. On the other hand, with the chromoelectric dipole moment being a $CP$-violating one, an odd power of $\rho'$ would denote a $CP$-odd (and $T$-odd) observable. It is easy to see that no such observable can be constructed out of the four momenta alone. Had we the ability to measure the polarizations, that possibility would open up too.

Similarly, the presence of absorptive parts of $\rho$ (or $\rho'$) would have allowed for the existence of such terms (essentially by changing the properties of the operator under time-reversal).

A further feature is the growth of the cross-sections with energy as is expected in a theory with dimension-five (or higher) operators. While this may seem unacceptable on account of

\footnote{Note that, while this vertex has occasionally been dropped or modified in literature, its inclusion is necessary for the $gg \to t\bar{t}$ amplitude to be a gauge invariant one.}
a potential loss of unitarity, one should realize that the theory of eqn. (11) is only an effective one and is expected to be superseded beyond the scale \( \Lambda \). While unitarity may be restored by promoting \( \rho (\rho') \) from constants to form-factors with appropriate powers of \((1 + \hat{s}/\Lambda^2)\), this is an *ad-hoc* measure as the mechanism of unitarity restoration is intricately related to the precise nature of the ultraviolet completion. We desist from doing this with the *a posteriori* justification that the limits of sensitivity for \( \rho \) (as described in the next section) are far beyond the typical subprocess energies \((\hat{s}/\Lambda \ll 1)\). Furthermore, note that the terms of \( \mathcal{O}(\rho) \) do respect partial wave unitarity. This is but a consequence of the fact that such terms appear as a result of interference between pure QCD and the dipole contributions, and owing to the different chirality structures of the operators, have to be proportional to \( m_t \).

### 3 Numerical Analysis

Armed with the results of the previous section, we compute the expected \( t\bar{t} \) cross-section at the Tevatron as well as the LHC. To this end, we use the CTEQ6L1 parton distribution sets [15] with \( m_t \) as the scale for both factorization as well as renormalization. To be consistent with the cross-section measurement reported by the CDF collaboration [16], we use \( m_t = 172.5 \) GeV for the Tevatron analysis. For the LHC analysis, though, we use the updated value of \( m_t = 173.1 \) GeV, obtained as a result of the combined CDF+DØ analysis [4]. To incorporate the higher order corrections absent in our leading order results, we use the \( K \)-factors at the NLO+NLL level [5] calculated by Cacciari et. al. [17]. Once this is done, the theoretical errors in the calculation owing to the choice of PDFs and scale are approximately 7-8% for the Tevatron and 9-10% for the LHC [17]. However, the estimates reported for the LHC operating at 7 TeV, are only leading order ones since NLO calculations for these energies are, so far, unavailable.

#### 3.1 Tevatron Results

At the Tevatron, the dominant contribution accrues from the \( q\bar{q} \) initial states, even on the inclusion of the dipole moments. While the \( \mathcal{O}(\rho^2, \rho'^2) \) terms in \( d\sigma/d\cos\theta \) are always positive (see eqn. 2), the flat \( \mathcal{O}(\rho) \) term can flip sign with \( \rho \). This implies that for \( \rho > 0 \), the change in the cross-section \( \delta\sigma > 0 \), thus severely constraining any deviation of the anomalous couplings in that direction. On the other hand, for \( \rho < 0 \), large cancellations may occur between various pieces of the cross-section. Consequently, substantial negative \( \rho \) could be admitted, but correlated with a substantial \( \rho' \). This is exhibited by Fig. 1a which displays the parameter space that is still allowed by the Tevatron data, namely [16] 

\[
\sigma_{t\bar{t}}(m_t = 172.5 \text{ GeV}) = (7.50 \pm 0.48) \text{ pb}. 
\]

The near elliptical shape of the contours is but a reflection of the fact that, at the Tevatron, the \( q\bar{q} \) contribution far supersedes the \( gg \) one.

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3 Note that \( \rho \) and \( \rho' \) lead to identical unitarity-breaking terms, a consequence of the fact that the difference between them necessarily has to be translated to subdominant terms.

4 And, similarly, those of \( \mathcal{O}(\rho^3, \rho\rho'^2) \).

5 In the absence of a similar calculation incorporating anomalous dipole moments, we use the same \( K \)-factor as obtained for the SM case. While this is not entirely accurate, given the fact that the color structure is similar and drawing from experience with analogous calculations for higher dimensional operators [18], the error associated with this approximation is not expected to be large.
Having seen the extent to which cancellations may, in principle, be responsible for hiding the presence of substantial dipole moments, we now restrict ourselves to the case where only one of $\rho$ and $\rho'$ may be non-zero. While this might seem a gross simplification, it is not really so. For one, with the chromomagnetic moment manifesting itself at $O(\Lambda^{-1})$ and the chromoelectric moment appearing in the cross-sections only at $O(\Lambda^{-2})$, it is obvious that, for large $\Lambda$, the former would, typically, leave a larger imprint. Secondly, it is extremely unlikely that the couplings conspire to be just so that large cancellations take place. This is particularly true because, for a generic underlying ultraviolet completion, one would expect the chromoelectric moment operator to appear at a higher order of perturbation than the chromomagnetic one. On the other hand, the situation could be reversed if there is an underlying symmetry (à la the symmetry proposed in Ref. [19] to account for neutrino magnetic moments) that prevents $\rho$ from appearing while allowing a non-zero $\rho'$.

If only one of the two couplings are to be non-zero, we may rescale $\rho, \rho' = 0, \pm 1$ and, thus, reduce the parameter space to one dimension ($\Lambda$). Of course, $\rho' = \pm 1$ are equivalent. Fig.1b exhibits the corresponding dependence of the total cross-section at the Tevatron on $\Lambda$ for various combinations of ($\rho, \rho'$). For $\rho = +1, -1$, the near-monotonic dependence on $\Lambda$ is reflective of the dominance of the $O(\rho/\Lambda)$ term. This is particularly true for $\Lambda/\rho > 3$ TeV.

The low sensitivity to the chromoelectric moment is understandable in view of the fact that the corresponding contribution is suppressed by at least $\Lambda^2$. Furthermore, unlike in the case of the chromomagnetic moment, the $q\bar{q} \rightarrow t\bar{t}$ cross-section in this case suffers an additional cancellation owing to the chirality structure (see eqn.2). With $\hat{s}$ at the Tevatron being only slightly greater than $4m_t^2$, this cancellation is quite significant.

Note that, for $\rho' = 0$, while Fig.1a shows a second range (close to $\rho/\Lambda \sim -2.2$) consistent with experimental observation, the existence of the same is not apparent in Fig.1b. As can be easily appreciated, for smaller $\Lambda$, the $O(\rho^2/\Lambda^2)$ term gets progressively more important, leading to a rapid growth in the total cross-section for $\rho > 0$ and a cancellation between the two leading terms for $\rho < 0$. Consequently, for smaller $\Lambda$, the lowest curve in Fig.1a would suffer a turnaround, rendering it consistent with the measurements. However, one should not be lead on too far by this. It is the $O(\Lambda^{-2})$ contributions that are largely responsible for this second region of consistency. On the other hand, the Lagrangian considered in eqn.1 contains...
only the lowest dimensional anomalous operators of an effective theory. Higher dimensional operators \( \mathcal{O}(\Lambda^{-2}) \), if included in the Lagrangian, could change the behaviour of the cross-sections. Were we to neglect \( \mathcal{O}(\Lambda^{-2}) \) terms in eqn.2, the shape of the curves would indeed change considerably\( ^6 \), but the limits on \( \Lambda \) for either of \( \rho = \pm 1 \) would hardly alter (see Fig.2). In other words, the sensitivity limits are overwhelmingly dominated by the \( \mathcal{O}(\rho/\Lambda) \) terms, thus making them quite robust. In fact, the change in the limits from inclusion of higher-order terms are well below the theoretical errors from sources such as the dependence on the factorization/renormalization scales, choice of PDF etc. To be quantitative, \( \Lambda \lesssim 7400 \text{GeV} \) can be ruled out at 99% confidence level for the \( \rho = +1 \) case. For \( \rho = -1 \) on the other hand, \( \Lambda \lesssim 9000 \text{GeV} \) can be ruled out at the same confidence level. One expects similar sensitivity for \( \rho = +1 \) and \( \rho = -1 \). The difference essentially owes its origin to the slight discrepancy between the SM expectations (as computed with our choices) and the experimental central value. Of course, restricting to \( \mathcal{O}(\Lambda^{-1}) \) eliminates \( \rho' \) altogether. However, sensitivity to \( \rho' \) may still be obtained by including absorptive pieces and/or by considering polarised scattering.

\[ \begin{align*}
\text{Figure 2: Comparison of production rates obtained at the Tevatron with truncated cross-sections (up to } \mathcal{O}(\Lambda^{-1}) \text{; denoted by subscript } \Lambda \text{ in the key) and full cross-sections(all orders in } \Lambda \text{).}
\end{align*} \]

### 3.2 LHC Sensitivity

At the LHC, it is the \( gg \) flux that rules the roost, especially at smaller \( \hat{s} \) values. Moreover, at high center-of-mass energies, the gluon-initiated cross-sections grow as \( \hat{s}/\Lambda^4 \), whereas the \( q\bar{q} \)-initiated cross-sections remain, at best, constant with \( \hat{s} \). Consequently, it is fair to say that the \( gg \to t\bar{t} \) subprocess dominates throughout. In Fig.3 we present the corresponding cross-sections at the LHC as a function of \( \Lambda \) for various values of the proton-proton center-of-mass energy \( \sqrt{s} \). In the absence of any data, we can only compare these with the SM expectations and the estimated errors. Experimental errors due to systematic and statistical uncertainties are expected to be between 20 and 30 percent for an integrated luminosity of 20 pb\(^{-1} \) at \( \sqrt{s} = 10 \text{ TeV} \) \( ^{21} \) (the errors for other values of \( \sqrt{s} \) are similar) and dominate the theoretical errors quoted earlier. Since experimental errors are expected to decrease with better calibration of the detectors and increase in statistics, we choose to display 10% (optimistic) and 20% error bars for comparison.

\( ^6 \)While it may be argued that, in principle, some as yet unknown symmetry could render such higher order terms in eqn.1 to be very small, we feel that such an eventuality would be a very artificial one.
Figure 3: \( t\bar{t} \) production rates for the LHC as a function of the new physics scale \( \Lambda \). Panels from left to right correspond to \( \sqrt{s} = 7, 10, 14 \) TeV. The horizontal lines show the SM expectation and the 10% and 20% intervals as estimates of errors in the measurement \([21]\).

Setting bounds on the chromoelectric dipole moment now becomes possible. Much of this is due to the fact that the \( gg \rightarrow t\bar{t} \) amplitude is not as chirally suppressed as the \( q\bar{q} \rightarrow t\bar{t} \) one (see eqn[2]). For \( \rho = 0 \), even an early run of the LHC with \( \sqrt{s} = 7 \) TeV (Fig. 3a) would be sensitive to \( \Lambda \lesssim 2700 \) GeV. Unfortunately, the improvement of the sensitivity with the machine energy is marginal at best. As for the chromomagnetic moment, the story is more complicated. For \( \rho = +1 \), naively a sensitivity up to about \( \Lambda \sim 10 \) TeV could be expected. For the \( \rho = -1 \) case though, the aforementioned cancellation between various orders reappears in a more complicated guise even for the \( gg \rightarrow t\bar{t} \) case. On account of this, it appears that the best that the LHC can do is to rule out (for \( \rho = -1 \)) \( \Lambda \sim < 8 \) TeV. This, however, should be compared with the Tevatron results which have already ruled out \( \Lambda \sim < 9 \) TeV.

While the discussion above was based on the full cross-sections as listed in eqn[2], the situation changes significantly if one were to truncate contributions beyond \( O(\Lambda^{-1}) \). In Table[1], we display the bounds on \( \Lambda \) that may be reached, with and without such a truncation, for the three different stages of LHC operation. In reaching these bounds, we have assumed that a 20% deviation constitutes a discernible shift. Several interesting features are apparent now. To begin with, the use of the full cross-section implies that a higher operative energy for the LHC renders it more suitable to the chromomagnetic moment as long as \( \rho = +1 \). The growth in sensitivity is a reflection of the growing importance of the higher order (in \( \rho/\Lambda \)) terms as the energy is increased \([3] \). For \( \rho = -1 \), on the other hand, the situation is reversed. The contrasting behaviour is easy to understand in terms of the constructive (destructive) interferences with the SM amplitude in the two cases. The tiny magnitude of the change is but a reflection of the fact that, owing to the nature of the \( gg \) and \( q\bar{q} \)–fluxes, most of the cross-section accrues from relatively smaller values of \( \hat{s} \). For the truncated case as well, the small increase in sensitivity with the \( pp \) center-of-mass energy can be understood by looking at eqn[2]. The SM as well as the \( O(\rho/\Lambda) \) terms in the cross-section fall with \( \hat{s} \). At lower values of \( \hat{s} \), the SM piece falls faster. However, in the higher \( \hat{s} \) regime, both have a similar behaviour and thus the increase in sensitivity that can be obtained by increasing the center-of-mass energy is only marginal. In fact, even with a several fold increase in the

\[ \text{These pieces in the cross-section do not fall off with } \hat{s} \]
LHC energy, the sensitivity would not increase by much. It is amusing to note that, were LHC a \( p\bar{p} \) collider instead, the use of the full matrix-element-squared would have entailed a substantial increase in the sensitivity with energy, although the situation for the truncated case would have remained quite similar.

If the LHC indeed measures a \( t\bar{t} \) cross-section that is in excess of the expected SM cross-section, it would be desirable to be able to identify the new physics responsible for it. For example, for \( pp \) collisions at \( \sqrt{s} = 10 \) TeV, various combinations \((\rho, \rho') = (1, 0), (0, \pm 1)\) may give rise to positive deviations of the order of, say 15–20\% in the total cross-section, albeit for wildly different values of\( \Lambda \) (Fig. 3b). Could differential distributions be used to discriminate between them or enhance the sensitivity to new physics? To this end, we begin with the invariant mass distribution and consider the full set of expressions of eqn.2 rather than the truncated ones. As Fig. 4a shows, the distributions do indeed diverge significantly for large \( m_{t\bar{t}} \). The two anomalous cross-sections depicted are roughly equal and deviate by approximately 20\% from the SM one. One might argue though that small differences in the spectrum could \((a)\) rise from various effects within the SM and/or experimental resolutions and \((b)\) get washed away as a result of poor statistics. The second objection is countered by the observation that significant deviations are associated not with very tiny cross-sections, but with a rather sizable event rate, even for a moderate value of the integrated luminosity. This deviation is emphasized further if one considers the ratio

\[
\left( \frac{1}{\sigma} \frac{d\sigma}{dm_{t\bar{t}}} \right) / \left( \frac{1}{\sigma_{SM}} \frac{d\sigma_{SM}}{dm_{t\bar{t}}} \right).
\]

This observable has the benefit of using normalized quantities so that some of the systematic errors such as those due to luminosity measurements or lack of precise knowledge of the parton densities are largely removed. As a perusal of Fig. 4b shows, the qualitative differences between the cases stands out starkly. It should be noted that the fact of the normalized distribution for the \( \rho = +1 \) case being very closely aligned with the SM one is not accidental, but just a consequence of the fact that it corresponds to a larger value of \( \Lambda \) compared to the other case. Consequently, the new physics contribution is dominated by the \( \mathcal{O}(\rho/\Lambda) \) piece, the corresponding integrated subprocess cross-section for which has a \( m_{t\bar{t}} \) behaviour quite similar to that for the SM piece. The other anomalous coupling point corresponds to a much smaller value of \( \Lambda \) and the \( \mathcal{O}(\rho^2/\Lambda^2) \) piece, which now plays an important role, has a very different \( m_{t\bar{t}} \) dependence.

Such distinctions can also be made using other kinematic variables such as transverse momentum and difference in the rapidities of \( t \) and \( \bar{t} \). However, they do not prove to be any more sensitive than the \( m_{t\bar{t}} \) distribution.
Figure 4: (a) The $m_{t\bar{t}}$ spectrum for the LHC at $\sqrt{s} = 10$ TeV along with the 1-$\sigma$ Gaussian error bar for the SM. (b) The ratio of the normalized $m_{t\bar{t}}$ spectra (see eqn. 7). In each case, the two anomalous sets refer to $(\rho, \rho', \Lambda) = (+1, 0, 11$ TeV) and $(0, \pm1, 3$ TeV) respectively. An integrated luminosity of 300 pb$^{-1}$ has been assumed.

4 Summary

The presence of anomalous chromomagnetic and chromoelectric dipole moments can produce sizable deviations in $t\bar{t}$ production cross-sections at hadron colliders. We obtain the range of parameter values allowed by the Tevatron measurements. Thereafter we examine the chromomagnetic and chromoelectric dipole interactions separately and find that the current Tevatron data rules out interactions giving rise to anomalous chromomagnetic moment of the top up to a scale of 7 TeV. At the LHC it will be possible to obtain sensitivity for such interactions up to 10 TeV. One will also be able to probe for anomalous chromoelectric dipole interactions manifest at scales $\sim 2$ TeV. Further, it will be possible to distinguish between the two dipole moments using invariant mass distributions. It is also seen that if the one retains only contributions up to $\mathcal{O}(\Lambda^{-1})$ the change in the limits on chromomagnetic moment remain within experimental error and cannot be distinguished.

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Additional Note

As this paper was being finalised Ref. [22] appeared as a preprint. Though we initially disagreed with some of their expressions, it was due to a misinterpretation, on our part, of
their conventions. We regret this error and thank the authors of Ref. [22] for pointing out this as well as some other typographical errors in our expressions. Our expressions are now in agreement with those in Ref. [22]. Bearing in mind that we use different computational methods, have different choices of parton densities and use different methods to account for NLO effects, our projected cross-sections can also be said to be in reasonable agreement.

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