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Constraining the intrinsic structure of top-quarks

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ABSTRACT

The basic structure of top-quarks as spin-1/2 particles is characterized by the radius R_t and the intrinsic magnetic dipole moment κ_t , both individually associated with gauge interactions. They are predicted to be zero in pointlike theories as the Standard Model. We derive upper limits of these parameters in the color sector from cross sections measured at Tevatron and LHC in top pair production $p\bar{p}/pp \rightarrow t\bar{t}$, and we predict improved limits expected from LHC in the future, especially for analyses exploiting boosted top final states. An additional method for measuring the intrinsic parameters is based on $t\bar{t} + \text{jet}$ final states.

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1. Basic set-up

The top-quark is the heaviest particle in the Standard Model (SM), even if the Higgs particle is included as a contender. This observation led to many approaches in which the top-quark plays the role of portal to physics beyond the Standard Model, see e.g. Refs. [1,2]. Scales characterizing the novel interactions in which the top-quark is identified with the crucial source field, may be realized not far beyond the TeV size. As a consequence, the top-quark may be endowed with intrinsic structure at the TeV scale. This should be contrasted with the pointlike character of all fundamental fields within the Standard Model, extending up to scales close to the Planck scale for low Higgs mass.

The basic non-pointlike structure will manifest itself in a non-zero radius R_t and a non-zero anomalous magnetic dipole moment κ_t in C_P -invariant scenarios, probed in interactions with gauge fields [3]. Due to the high energy available, the LHC will enable us to probe the intrinsic top-quark structure in the colored sector at an unprecedented level [4,5,7,8,6]. Non-pointlike interactions with the gluon field¹ modify the color quark current to [3]

$$\mathcal{J}_\mu = F_t \gamma_\mu + i \frac{\kappa_t}{2m_t} \sigma_{\mu\nu} Q^\nu. \quad (1.1)$$

The current incorporates the form factor

$$F_t = 1 + \frac{1}{6} R_t^2 Q^2, \quad (1.2)$$

with the top-quark radius R_t related by

$$R_t = \sqrt{6}/\Lambda_*, \quad (1.3)$$

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¹ The massless gluon gauge field is assumed intrinsically pointlike in the present analysis. This assumption can be removed, see Ref. [9], at the expense of increasing complexity. Non-pointlike structures of the weak current remain non-effective as long as the top decay is treated inclusively with $BR(t \rightarrow bW)$ very close to unity.

to the new scale parameter Λ_* , and the anomalous chromo-magnetic dipole moment κ_t [beyond the loop value [10]]. To protect fermion masses from acquiring large values, the theory is generally assumed chiral [11], and the breaking of the chiral symmetry by anomalous magnetic moments is suppressed by two powers of the scale Λ_* , in the simplest possible realization:

$$\kappa_t = \rho m_t^2 / \Lambda_*^2, \quad (1.4)$$

where $|\rho|$ is an $\mathcal{O}(1)$ number. The quadratic Λ_* dependence of κ_t is effectively equivalent to the scaling of the form factor. The quadratic dependence in the heavy quark mass singles out the top-quark as unique particle for which κ_t may be accessible experimentally, in contrast to much less sensitive light quarks or leptons. Assuming Λ_* to be of order 1 TeV and beyond, compatible with bounds on contact interactions from Tevatron and LHC [12], κ_t could be expected at the level of several per-cent.

Both the anomalous parameters, color radius and color magnetic dipole moment, can be introduced through effective Lagrangians [13] in an $SU(3)_c$ gauge-invariant and parity-even form²:

$$\mathcal{L}_R = -g_s \frac{R_t^2}{6} \bar{t} \gamma^\mu \mathcal{G}_{\mu\nu} D^\nu t + \text{h.c.}, \quad (1.5)$$

$$\mathcal{L}_\kappa = g_s \frac{\kappa_t}{4m_t} \bar{t} \sigma^{\mu\nu} \mathcal{G}_{\mu\nu} t, \quad (1.6)$$

with the gluon field \mathcal{G}_μ , in octet matrix notation, and the gluon field strength $\mathcal{G}_{\mu\nu} = D_\nu \mathcal{G}_\mu - D_\mu \mathcal{G}_\nu$, while $D^\nu = \partial^\nu + ig_s \mathcal{G}^\nu$ denotes the covariant derivative of QCD. Besides the components generating the anomalous top color current, the Lagrangians are complemented by additional two-gluon and three-gluon top interactions, as demanded by gauge invariance. The effective Lagrangians unambiguously translate the anomalous parameters from scattering to annihilation processes.

The classical method for studying radius and anomalous magnetic dipole moment of the top quark is given by the elastic Rutherford-type scattering of a top quark t with a light quark q [taken pointlike in the present scenario], which is mediated by the exchange of a gluon in $qt \rightarrow qt$. Rutherford-type scattering is also embedded in the process $gq \rightarrow t\bar{t}q$. At very high energies, gluon partons in the protons split into beams of long-lived top-quark pairs traveling parallel to the gluon momentum. Thus, the events of the $t\bar{t}q$ process, characterized by a forward moving t -quark plus a $\bar{t}q$ -pair, with the two partons in the pair balanced in transverse momentum, signal Rutherford qt scattering. [Elastic gluon-top scattering is independent of the radius R_t and cannot be exploited.]

2. Theoretical groundwork

We will analyze the total cross sections for the production of top-quark pairs

$$p\bar{p}/pp \rightarrow q\bar{q}, \quad gg \rightarrow t\bar{t} \quad (2.1)$$

at Tevatron and LHC for deriving limits on the color radius R_t , the anomalous chromo-magnetic dipole moment κ_t and the Λ_* parameter in practice. Additional constraints can be derived from the angular dependence of the top-quarks, and the correlations between longitudinal spin components of t and \bar{t} [14], which can be measured unperturbed by fragmentation due to the short top lifetime [15]. Related analyses have been discussed in Refs. [18,16,17].

We will assume that the non-pointlike contributions to the observables are small and, correspondingly, we will expand the observables linearly in the analytic formulae. In fact, anomalous chromo-magnetic dipole moment and chromo-radius are the first terms of a multipole expansion including scale parameters beyond the Standard Model. The systematic expansion would continue with higher-order moments the quadratic terms in R_t^2 and κ_t would compete with. An analysis of these contributions is beyond the scope of the present Letter.

The hadron cross sections are built up by the incoherent superposition of quark-antiquark annihilation and gluon fusion to top-antitop pairs. Quark-antiquark annihilation is mediated only by s -channel gluon exchange³, gluon fusion by s -channel gluon and t, u -channel top exchanges.

The anomalous terms of the independent cross sections at the parton level can be summarized as follows [see also references quoted above], using $\beta = \sqrt{1 - 4m_t^2/s}$, where s is the partonic center-of-mass energy:

Quark-antiquark annihilation:

$$\frac{\Delta\sigma}{\sigma_B} = \frac{s}{3} R_t^2 + \frac{6\kappa_t}{3 - \beta^2}, \quad (2.2)$$

$$\frac{\Delta d\sigma/d\cos\theta}{d\sigma_B/d\cos\theta} = \frac{s}{3} R_t^2 + \frac{4\kappa_t}{2 - \beta^2(1 - \cos^2\theta)}, \quad (2.3)$$

$$\frac{\{t_R\bar{t}_R + t_L\bar{t}_L\} - \{t_R\bar{t}_L + t_L\bar{t}_R\}}{\{t_R\bar{t}_R + t_L\bar{t}_L\} + \{t_R\bar{t}_L + t_L\bar{t}_R\}} = -\frac{1 + \beta^2}{3 - \beta^2} + \frac{8\beta^2}{(3 - \beta^2)^2} \kappa_t. \quad (2.4)$$

Gluon fusion:

$$\frac{\Delta\sigma}{\sigma_B} = \frac{(36\beta - 64\tanh^{-1}\beta)\kappa_t}{\beta(59 - 31\beta^2) - 2(33 - 18\beta^2 + \beta^4)\tanh^{-1}\beta}, \quad (2.5)$$

² Electroweak gauge invariance can be ensured by expanding the Lagrangians to the complete third generation and incorporating the Higgs field [13].

³ We neglect electroweak interactions in the following.

$$\frac{\Delta d\sigma/d\cos\theta}{d\sigma_B/d\cos\theta} = \frac{4(1-\beta^2\cos^2\theta)\kappa_t}{2-\beta^4-(1-\beta^2+\beta^2\cos^2\theta)^2}, \quad (2.6)$$

$$\begin{aligned} & \frac{\{t_R\bar{t}_R+t_L\bar{t}_L\}-\{t_R\bar{t}_L+t_L\bar{t}_R\}}{\{t_R\bar{t}_R+t_L\bar{t}_L\}+\{t_R\bar{t}_L+t_L\bar{t}_R\}} \\ &= -\frac{\beta(66-37\beta^2+31\beta^4)-2(33-33\beta^2+17\beta^4-\beta^6)\tanh^{-1}\beta}{\beta^2[\beta(59-31\beta^2)-2(33-18\beta^2+\beta^4)\tanh^{-1}\beta]} \\ &\quad -\frac{4[\beta(33+11\beta^2)-(33-\beta^4)\tanh^{-1}\beta][\beta(41-31\beta^2)-2(17-18\beta^2+\beta^4)\tanh^{-1}\beta]}{\beta^2[\beta(59-31\beta^2)-2(33-18\beta^2+\beta^4)\tanh^{-1}\beta]^2}\kappa_t, \end{aligned} \quad (2.7)$$

in agreement, wherever overlapping, with e.g. [19,18]. Other helicity asymmetries are related by \mathcal{P} and \mathcal{C} invariance. The effective top current Eq. (1.1), generates the same dependence on the anomalous parameters in the $q\bar{q}$ amplitude, so that the top radius and the anomalous magnetic moment R_t, κ_t can indeed be interpreted as gauge-invariant characteristics of the top quark.

The $q\bar{q}$ annihilation channel is modified by both the radius and the chromo-magnetic moment. By contrast, gluon fusion does not depend on the radius to leading order – reminiscent of the Thomson cross section in QED – but only on the anomalous magnetic moment (see Appendix A and e.g. Ref. [18]). Since top production at the Tevatron is driven by $q\bar{q}$ collisions, both parameters can in principle be determined in top measurements at this collider. On the other hand, the LHC, where gluon fusion is by far the dominant inclusive top channel, is highly sensitive to the value of the color anomalous magnetic dipole moment, leading to large bounds on the scale parameter Λ_* . However, at the expense of reduced cross sections, the relative weight of the $q\bar{q}$ channel can be increased by the production of boosted top events at LHC, and $t\bar{t}$ production in this configuration becomes also sensitive to the radius.

Naturally assuming universality for the light quarks $q = u, d$, the bounds on R_t can be transcribed easily to the scale of the standard color-octet vectorial contact interactions $\mathcal{L}_{ct} = g_{ct}^2/\Lambda_{ct}^2(\bar{q}\gamma^\mu T^A q)(\bar{t}\gamma_\mu T^A t)$, where $T^A = \lambda^A/2$ denote the SU(3) generators of QCD, expressed by the Gell-Mann matrices λ^A . After inserting the effective contact coupling $g_{ct}^2 = 4\pi$ of the two quark currents, as generally defined, the contact scale Λ_{ct} is related to the compositeness scale Λ_* by

$$\Lambda_{ct} \sim 1/\sqrt{\alpha_s} \Lambda_*. \quad (2.8)$$

The different coupling strengths boost the contact scale to a value half an order of magnitude above the compositeness scale. Current constraints limit the octet contact scale to $\Lambda_{ct} \gtrsim 2.8$ TeV, see below and Ref. [18]. The singlet contact scale of general chiral quark interactions has been constrained to $\gtrsim 3.4$ TeV at LHC [12]; this bound may be compared with $\sqrt{3/\sqrt{2}} \Lambda_{ct} \simeq 4.1$ TeV for top interactions if the singlet energy density is identified, hypothetically, with the octet density.

3. Numerical evaluation

The determination of the anomalous parameters by three independent measurements of cross sections at three different energies and different superpositions of the parton subprocesses at Tevatron and LHC is over-constrained. Leaving the exhaustive evaluation to experimental analyses proper we focus in this theoretical study on the total cross sections at Tevatron and LHC. Combining the cross sections of both colliders the different weight of $q\bar{q}$ and gg events allows us to separate the parameters R_t and κ_t . We will also investigate the cross section for boosted final-state tops, which are well-accessible at the LHC with 14 TeV center-of-mass energy, again initiated by $q\bar{q}$ and gg parton compositions different from inclusive cross sections. These experimental observables are well documented by both the collaborations at the two colliders [20–22].

The $t\bar{t} + X$ cross section follows from the modified Born-level $t\bar{t}$ amplitudes $\mathcal{M} = \mathcal{M}_{SM} + \mathcal{M}(\kappa_t, R_t)$ for the partonic subprocesses $ab = q\bar{q}, gg$, where q denotes the light quark flavors, so that

$$\begin{aligned} \Delta\sigma &= \sum_{ab \in \{q\bar{q}, gg\}} \iiint dx_1 dx_2 dLIPS f_a(x_2, \mu_F^2) f_b(x_2, \mu_F^2) \{|\mathcal{M}_{ab}|^2 - |\mathcal{M}_{SM ab}|^2\} \\ &= \sum_{ab \in \{q\bar{q}, gg\}} \iiint dx_1 dx_2 dLIPS f_a(x_2, \mu_F^2) f_b(x_2, \mu_F^2) 2\text{Re}\{\mathcal{M}_{SM ab}^* \mathcal{M}_{ab}(\kappa_t, R_t)\} + \mathcal{O}\left(\frac{1}{\Lambda_*^4}\right). \end{aligned} \quad (3.1)$$

For the remainder of this analysis we choose the CTEQ6L1 parton distribution set [23].

We have implemented the parton-level cross section of Eq. (3.1) in a fully flexible numerical program based on the VBFNLO framework [24]. The calculation of the matrix elements is performed with a set of custom-built HELAS routines [25] which facilitate the numerical evaluation of the three- and four-point contributions of Eq. (1.5). The R_t -dependent terms in the on-shell gluon-induced subprocess drop out, as discussed earlier. This cancellation persists in the gluon initiated subprocess of $p(p) \rightarrow t\bar{t} + \text{jet}$ production, which incorporates off-shell gluons by emitting an additional jet (see below). However, this process is still worth studying since the quark-gluon initiated channel first enters at this order and thus offers new ways of probing R_t . The numerical implementation for $p(p) \rightarrow t\bar{t} + \text{jet}$ is set up analogously to $p(p) \rightarrow t\bar{t}$, supplementing the relevant five-point interactions following from Eq. (3.1). We have checked all contributing matrix elements for gauge invariance and we have validated our phase space integration against MADEVENT [26] and SHERPA [27].

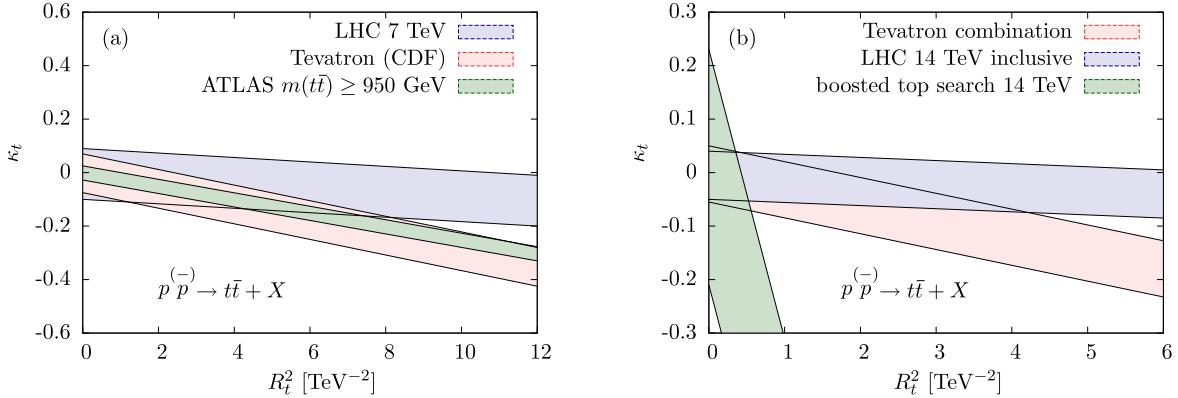


Fig. 1. (a) Bands allowed in R_t, κ_t space by $t\bar{t}$ production at Tevatron and LHC for 7 TeV, available data; (b) The same for LHC at 14 TeV, theoretical expectation of inclusive cross sections and boosted top events. We also show the region that is allowed by the high top pair invariant mass bin as reported in a recent ATLAS investigation of the differential $t\bar{t}$ cross section [32], which is largely equivalent in sensitivity to the CDF analysis. We use $\mu_R = \mu_F = m_t$ and $\mu_F = \mu_R = \bar{m}_T$ for the boosted search, where \bar{m}_T denotes the average transverse mass of the top quarks.

Table 1

Upper bounds on t radius and magnetic moment after combining Tevatron and LHC data/future expectations for $t\bar{t}$ production – inclusive and boosted top measurements at LHC.

	R_t	$ \kappa_t $
Tevatron \oplus LHC[7 TeV]	$2.9 \text{ TeV}^{-1} \sim 0.57 \times 10^{-16} \text{ cm}$	0.17
Tevatron \oplus LHC[14 TeV]	$2.1 \text{ TeV}^{-1} \sim 0.41 \times 10^{-16} \text{ cm}$	0.07
LHC[14 TeV]: inclusive \oplus boosted top	$0.7 \text{ TeV}^{-1} \sim 0.14 \times 10^{-16} \text{ cm}$	0.05

The upper limits $\Delta\sigma(t\bar{t} + X)$ as generated⁴ by the anomalous top parameters, color radius R_t and magnetic dipole moment κ_t , are identified with the difference between the measured and the theoretically predicted SM cross sections, both including errors (for details see below):

$$\text{Tevatron (CDF) [20]: } \sigma(t\bar{t} + X) = 7.5 \pm 0.31 \text{ (stat)} \pm 0.34 \text{ (syst)} \pm 0.15 \text{ (lumi) pb}, \quad (3.2)$$

$$\text{LHC, } \sqrt{s} = 7 \text{ TeV [21]: } \sigma(t\bar{t} + X) = 177 \pm 3 \text{ (stat)}^{+8}_{-7} \text{ (syst)} \pm 7 \text{ (lumi) pb}. \quad (3.3)$$

We have adopted the theoretical expectations of Ref. [28]

$$\text{Tevatron: } \sigma(t\bar{t} + X) = 7.13^{+0.30}_{-0.40} \text{ (scale)}^{+0.17}_{-0.12} \text{ (pdf) pb}, \quad (3.4)$$

$$\text{LHC, } \sqrt{s} = 7 \text{ TeV: } \sigma(t\bar{t} + X) = 164.3^{+3.3}_{-9.2} \text{ (scale)}^{+4.4}_{-4.5} \text{ (pdf) pb}, \quad (3.5)$$

$$\text{LHC, } \sqrt{s} = 14 \text{ TeV: } \sigma(t\bar{t} + X) = 908.3^{+9.8}_{-40.5} \text{ (scale)}^{+15.2}_{-16.7} \text{ (pdf) pb} \quad (3.6)$$

as representative figures of the inclusive $t\bar{t}$ cross sections [29–31]. It has been shown in Refs. [30,31] that the perturbative evolution up to the full NNLO precision result for $t\bar{t}$ production at the Tevatron reduces the renormalization scale uncertainty by $\mathcal{O}(30\%)$ and a similar improvement is expected for LHC predictions. We include the theoretical uncertainty due to variations of the renormalization scale and errors of the parton densities by adding it to the previously mentioned experimental error in quadrature; the differences of the theoretical and experimentally expected mean values are added equivalently.

This procedure gives rise to a band of viable values in the $\{R_t, \kappa_t\}$ -parameter plane from each of the two colliders, Fig. 1. The crossing of the bands allows us to determine the upper limits of the two parameters separately, resulting in the conservative upper bounds collected in Table 1. At the LHC, the inclusive $t\bar{t}$ cross section is driven by the gluon-fusion channel, which has no dependence on R_t , see Eq. (2.5). This makes it difficult to obtain stringent bounds on R_t , in contrast to the Tevatron where the quark-antiquark channel is dominant, see Eq. (2.2).

The bound on $|\kappa|$ from the combination ‘Tevatron \oplus LHC[7 TeV]’ of the presently available data would shrink to $|\kappa| < 0.06$ if the top radius is set to zero. Comparing this value with appropriate values in the literature based on analyses of chromo-magnetic and chromo-electric dipole moments [5,16], they agree within errors of 30%.

For the $t\bar{t}$ cross section at the Tevatron there are statistical improvements upon combining the data sets of DØ and CD [33]. Similar improvements can be expected at the LHC for the 14 TeV run, when more data will become available. We show a projection of this situation in Fig. 1(b), where we scale the CDF error of Eq. (3.2) by a factor $1/\sqrt{2}$, and the LHC systematic uncertainty is saturated at 5% [34]. This shows that we can indeed expect a significant improvement on the limits of $\{R_t, \kappa_t\}$ at the level of inclusive searches.

Even though we cannot carry out a rigorous analysis of quadratic effects in the multipole expansion, nevertheless for LHC[7 TeV], as a typical example, we may quote a rough estimate at what level the linear term may penetrate the quadratic term. With $\langle s \rangle \sim 1/4 \text{ TeV}^2$

⁴ The corresponding code for $\Delta\sigma$ is available upon request from the authors.

the linear correction induced by the radius amounts to $2\langle s \rangle / \Lambda_*^2 \sim 0.7 < 1$, where however it should be noted that this number is still reduced significantly by a negative contribution of the magnetic term – this destructive interference being one of the crucial elements in our analysis. Taking the estimate above at face value, the quadratic term inferred from the interference term is less than about 10%. Thus, the rough estimates signal internal consistency of our analysis.

However, despite its much larger collision energy, the sensitivity of LHC to the anomalous top couplings is improved only moderately compared to the Tevatron, as a result of the prevalence of the R_t -insensitive gluon-fusion component in the total hadronic cross section. A way to eliminate this obstacle is to consider boosted top final states [35]. By restricting ourselves to large momentum transfers we probe the incoming protons at large momentum fractions, thus naturally shifting towards the $q\bar{q}$ contribution, which is more sensitive to R_t . This improvement more than compensates for the significant reduction of the hadronic $t\bar{t}$ production from imposing this cut. We include this search channel in Fig. 1(b), where we choose $p_{T,t} \geq 1$ TeV, for which we expect $\sigma_{SM} \simeq 50$ fb and a 30% measurement uncertainty. This error estimate should be understood figuratively as a dedicated calculation in this phase space region analogous to Refs. [28,30,31,29] is currently not available. Recent analyses of the differential $t\bar{t}$ cross section [32] however suggest that this is roughly the uncertainty that can be expected. It is clear that by increasing the transverse momentum selection, we probe larger partonic center-of-mass energy, which in turn yields a larger sensitivity to the anomalous parameters. This stems from probing predominantly quark-induced subprocesses at large partonic momentum fractions. We can expect that the flat background distribution qualitatively behaves $\sim p_{T,cut}^{-2}$ so that a background fluctuation is parametrically described by $p_{T,cut}^{-1}$. On the other hand the signal cross section in the dominant quark channels for the boosted selection behave $\sim \Lambda_*^2$. Hence the sensitive region for the boosted selection is characterized by $\Lambda_*^2 \lesssim p_{T,cut}$ until the rate at a given luminosity is too small to efficiently reconstruct the $t\bar{t}$ system. Thus, the sensitivity will increase with the cut on the transverse momentum until the error in the cross section becomes overwhelming. By the same reason, the precise value of the involved uncertainties is not too important for the qualitative success of constraining the anomalous top interactions using a boosted selection. It should also be noted that these SM errors are expected to be also the main errors in the part of the cross sections describing anomalous contributions. Central sources for errors like the scales in the QCD coupling and the parton densities are not significantly different from the SM, i.e. they cancel out from observables like the radius, operatively defined in a ratio of cross sections. A combination of either inclusive LHC cross sections together with finalized Tevatron results, or inclusive cross sections and boosted searches solely at the LHC provide good prospects to sharpen the bounds on anomalous top interactions.

The anomalous parameters R_t, κ_t can be translated to the scale parameters Λ_* and $\Lambda_*/\sqrt{|\rho|}$ [as denoted in Eqs. (1.3) and (1.4)]. Using the estimated bounds on the radius R_t from the Tevatron and the LHC experiments, one obtains

$$\text{Tevatron} \oplus \text{LHC}[7 \text{ TeV}]: \quad \Lambda_* \gtrsim 0.84 \text{ TeV}, \quad (3.7)$$

while the bound from the anomalous magnetic moment, for the characteristic choice $\rho = 1$, is weaker by a factor of 2. Identifying $g_S^2/\Lambda_*^2 \rightarrow c_V/(2\Lambda^2)$, this number is in agreement with Ref. [18] when taking into account the different conventions, and the bounds can be improved by fitting the di-top invariant mass distribution [18].

ATLAS has already published results on centrally produced, high invariant-mass top pairs, $m_{t\bar{t}} \geq 950$ GeV [32]. We find that limits obtained from this result are compatible with the combined analysis ‘Tevatron \oplus LHC[7 TeV]’.

Improvements of the bound are expected for LHC[14 TeV], particularly if boosted top analyses are exploited:

$$\text{Tevatron} \oplus \text{LHC}[14 \text{ TeV}]: \quad \Lambda_* \gtrsim 1.17 \text{ TeV}, \quad (3.8)$$

$$\text{LHC}[14 \text{ TeV}; \text{inclusive} \oplus \text{boosted top}]: \quad \Lambda_* \gtrsim 3.5 \text{ TeV}, \quad (3.9)$$

dominated again within a factor of at least 2 by the bound on the radius. Boosted strategies are not applicable at the Tevatron due the limited data set and the small available center-of-mass energy compared to LHC[14 TeV]. The LHC[7 TeV] data sample is also too small, but first results can be expected from the LHC[8 TeV] run.

Including the measurements of angular distributions and spin correlations in the experimental analyses will lift these limits to still higher values.

As mentioned above, the scale parameter Λ_* can also be transcribed to octet contact interactions, lifting the scale parameter by half an order of magnitude, cf. Eq. (2.8). Presently a bound of $\Lambda_{ct} \gtrsim 2.8$ TeV has been reached. The bound will improve significantly at 14-TeV LHC,

$$\Lambda_{ct} \gtrsim 11.7 \text{ TeV}, \quad (3.10)$$

in the near future, corresponding for singlet currents even to an estimated 17.0 TeV for the singlet energy density identified, hypothetically, with the color averaged octet density.

4. Jet emission

Earlier we argued that the classical Rutherford process $qt \rightarrow qt$ can be exploited for measuring the radius of the t -quark while the Thomson analogue $gt \rightarrow gt$ does not depend on the radius to leading order. These rules are also effective in the crossed channels $q\bar{q} \rightarrow t\bar{t}$ and $gg \rightarrow t\bar{t}$ applied in practice to measure the t radius and the magnetic moment. Adding a gluon jet to the final state in $\sigma[gg \rightarrow t\bar{t}g]$, the gluon-fusion process still has no dependence on the radius. This is obvious for the logarithmically enhanced splitting process $g \rightarrow gg$ followed by $gg \rightarrow t\bar{t}$, but it remains true also for the non-logarithmic part. This is a consequence of cancellations among the modified three-point ($g\bar{t}t$) vertex and novel four- and five-point ($gg\bar{t}t$, $gggt\bar{t}$) vertex contributions to $gg \rightarrow t\bar{t}g$ [resulting from Eq. (1.5) and (1.6)], which do not only serve to enforce the QCD Ward-identities but also eliminate the R_t -dependent terms. In Fig. 2, it is demonstrated numerically that indeed $\sigma[gg \rightarrow t\bar{t}g]$ is independent of R_t . A cut, $p_{T,j} \geq 100$ GeV, has been imposed on the transverse momentum of the jet. By contrast, the subprocess $gq \rightarrow t\bar{t}q$ depends, weakly though, on the top radius already to logarithmic accuracy through the gluon splitting channel $g \rightarrow q\bar{q}$ followed by $q\bar{q} \rightarrow t\bar{t}$, supplemented by additional non-logarithmic contributions.

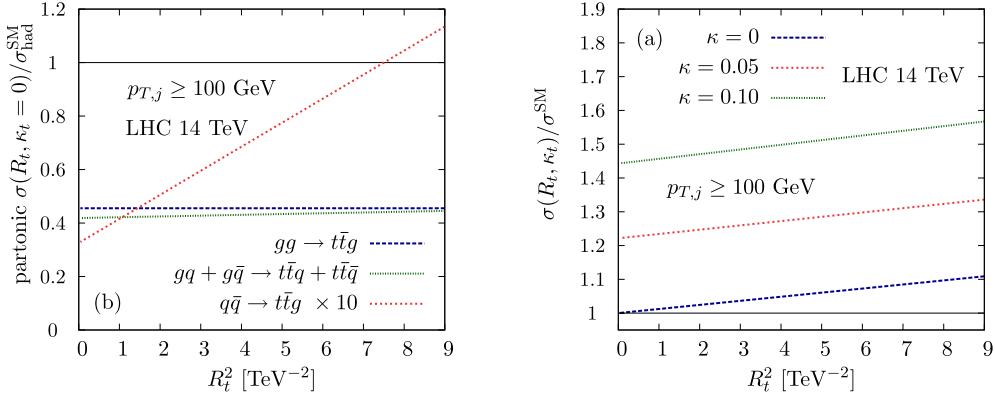


Fig. 2. Dependence of the jet cross sections $\sigma[t\bar{t}j]$ on R_t, κ_t . Left panel: gluon fusion, (anti-)quark-gluon scattering, and quark-antiquark annihilation subprocesses compared with the original Born cross sections at the LHC $\sqrt{s} = 14$ TeV; Right panel: LHC cross sections for various values of κ_t . We use again $\mu_R = \mu_F = \bar{m}_T$.

However, this subchannel is dominated by gluon radiation $q \rightarrow qg$ and $gg \rightarrow t\bar{t}$, which depends on the top radius only beyond the leading logarithmic order. The strongest R_t dependence is predicted for the annihilation channel $q\bar{q} \rightarrow t\bar{l}g$ via two steps, $q \rightarrow qg$ and $q\bar{q} \rightarrow t\bar{l}$, with logarithmic enhancement. The equivalent hadron cross section is shown in the right panel of Fig. 2.

Sensitivity to the radius in $t\bar{t} +$ jet production is largely driven by the $q\bar{q} \rightarrow t\bar{l}g$ subprocess, but some additional sensitivity arises from gq scattering (and the charge-conjugated $g\bar{q}$ channel). As discussed above, the latter originates from two contributions: (i) splitting of an initial-state gluon into a quark-antiquark pair, $g \rightarrow q\bar{q}^*$, followed by the R_t -dependent subprocess $q\bar{q}^* \rightarrow t\bar{l}$; (ii) radiation of an off-shell gluon from the incident quark, $q \rightarrow qg^*$, followed by $gg^* \rightarrow t\bar{l}$. To leading logarithmic order, i.e. for nearly on-shell gluons, the second process is independent of R_t as argued earlier. However, if sufficiently high- p_T jets are observed, the events are pushed out of the DGLAP regime. Since the intermediate gluon in this case is *off-shell*, the contribution of the operator (1.5) is not forbidden by Ward identities. This situation corresponds to the Rutherford-type scattering discussed on page 262. One finds for the R_t -dependent contribution to the cross section (ii) in linear approximation, the transverse part singled out for the sake of transparent illustration:

$$\frac{d\Delta\sigma[gq \rightarrow t\bar{l}q]}{d\cos\theta^*dq^2} \propto \frac{R_t^2}{6} [(1-\xi)(23+50\xi-9\xi^2) + 4(1-\xi)(37+27\xi)m_t^2/s - 64(7+9\xi)m_t^4/s^2]/(1-\xi)^2 + \dots \quad (4.1)$$

[while absent in $gg \rightarrow t\bar{l}g$ due to the anomalous $(g)gg\bar{t}\bar{t}$ vertices according to Eq. (1.5)]. Here $q^2 < 0$ is the 4-momentum transfer in the quark-line, i.e. the virtuality of the off-shell gluon, and $\xi = \beta^2 \cos^2 \theta^*$, where θ^* is the angle of the top-quark in the $t\bar{t}$ rest frame. As follows from scaling, the size of the anomalous part of the cross section is determined by the radius and it is independent of q^2 for $|q^2| \ll R_t^{-2}$, fulfilled in all realistic configurations. For the sake of clarity, the chromo-magnetic moment κ_t has not been included in (4.1), but the results in Fig. 2 are based on both the operators and all relevant diagrams with exact kinematics.

Therefore, with sufficient experimental precision, the top radius R_t can also be probed in $t\bar{t} +$ jet in a unique way that will help to discriminate its effect from the anomalous top magnetic moment. The $t\bar{t} +$ jet process has a smaller cross section $\sigma_{\text{NLO}}^{\text{SM}}(p_{T,j} \geq 50 \text{ GeV}) = 375 \text{ pb}$ [36] as compared to inclusive $t\bar{t}$ production. A measurement of $t\bar{t} +$ jet is also more involved from an experimental systematics point of view, and currently there is no dedicated analysis available at the LHC that targets the high p_T regime [see Ref. [37] for a first measurement of inclusive $t\bar{t} + j$] Nonetheless, sensitivity can also be gained in this channel using similar strategies as discussed in Section 3.

5. Summary

The intrinsic structure of the top-quark can sensitively be probed at the Tevatron and LHC by setting bounds on the color radius and the color magnetic dipole moment of the particle. Values of

$$R_t \lesssim 1.4 \times 10^{-17} \text{ cm} \quad \text{and} \quad |\kappa_t| \lesssim 0.05 \quad (5.1)$$

can be expected from LHC running in the near future. Present bounds, combined with Tevatron results will improve by factors of 4 and 3, respectively. These values can be mapped into effective scale parameters

$$\Lambda_* \gtrsim 3.5 \text{ TeV} \quad \text{and} \quad \Lambda_{ct} \gtrsim 11.7 \text{ TeV} \quad (5.2)$$

[and potentially even 17 TeV for singlet currents], strongly constraining the pointlike character of the top quark.

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Appendix A

A few examples should illustrate the dependence of relevant helicity amplitudes on the anomalous top parameters. Other helicity amplitudes are related by \mathcal{P} and \mathcal{C} symmetries, and the exchange $\cos\theta \rightarrow -\cos\theta$. Dimensional parameters are scaled in $E = \sqrt{s}/2$.

Quark-antiquark annihilation:

$$\begin{aligned} \mathcal{M}(q_R^i \bar{q}_L^j \rightarrow t_L^k \bar{t}_L^l) &= \frac{K_q}{6m_t \sqrt{s}} [2m_t^2(6 + sR_t^2) + 3s\kappa_t] \sin\theta \\ &= -\mathcal{M}(q_L^i \bar{q}_R^j \rightarrow t_L^k \bar{t}_L^l) = -\mathcal{M}(q_R^i \bar{q}_L^j \rightarrow t_R^k \bar{t}_R^l) = \mathcal{M}(q_L^i \bar{q}_R^j \rightarrow t_R^k \bar{t}_L^l), \end{aligned} \quad (\text{A.1})$$

$$\mathcal{M}(q_L^i \bar{q}_R^j \rightarrow t_L^k \bar{t}_R^l) = \frac{K_q}{6} (1 + \cos\theta) [6 + sR_t^2 + 6\kappa_t] = -\mathcal{M}(q_R^i \bar{q}_L^j \rightarrow t_R^k \bar{t}_L^l), \quad (\text{A.2})$$

$$\mathcal{M}(q_R^i \bar{q}_L^j \rightarrow t_L^k \bar{t}_R^l) = -\mathcal{M}(q_L^i \bar{q}_R^j \rightarrow t_R^k \bar{t}_L^l) = \mathcal{M}(q_L^i \bar{q}_R^j \rightarrow t_L^k \bar{t}_R^l)[\cos\theta \leftrightarrow -\cos\theta]. \quad (\text{A.3})$$

Gluon fusion:

$$\begin{aligned} \mathcal{M}(g_L^a g_L^b \rightarrow t_L^k \bar{t}_L^l) &= \frac{1}{2m_t \sqrt{s}} \left[K_t \left[\frac{1 + \cos\theta}{1 - \beta \cos\theta} \{4m_t^2(1 + \beta - \beta \cos\theta) + s(1 - \beta)(2 + \beta - \beta \cos\theta)\kappa_t\} \right] \right. \\ &\quad \left. + K_u [\cos\theta \leftrightarrow -\cos\theta] - K_s \cos\theta [4m_t^2 + s(1 - \beta)\kappa_t] \right] \\ &= -\mathcal{M}(g_R^a g_R^b \rightarrow t_R^k \bar{t}_R^l), \end{aligned} \quad (\text{A.4})$$

$$\mathcal{M}(g_L^a g_L^b \rightarrow t_R^k \bar{t}_L^l) = -\mathcal{M}(g_R^a g_R^b \rightarrow t_L^k \bar{t}_L^l) = \mathcal{M}(g_L^a g_L^b \rightarrow t_L^k \bar{t}_L^l)[\cos\theta \leftrightarrow -\cos\theta, \beta \leftrightarrow -\beta], \quad (\text{A.5})$$

$$\begin{aligned} \mathcal{M}(g_L^a g_R^b \rightarrow t_L^k \bar{t}_L^l) &= \frac{\beta}{2m_t \sqrt{s}} \left[\frac{K_t}{1 - \beta \cos\theta} + \frac{K_u}{1 + \beta \cos\theta} \right] \sin^2\theta [4m_t^2 + s\kappa_t] \\ &= \mathcal{M}(g_R^a g_L^b \rightarrow t_L^k \bar{t}_L^l) = -\mathcal{M}(g_L^a g_R^b \rightarrow t_R^k \bar{t}_R^l) = -\mathcal{M}(g_R^a g_L^b \rightarrow t_R^k \bar{t}_L^l), \end{aligned} \quad (\text{A.6})$$

$$\begin{aligned} \mathcal{M}(g_L^a g_L^b \rightarrow t_L^k \bar{t}_R^l) &= \sin\theta \left[K_t \left(1 + \kappa_t \frac{2 - \beta \cos\theta}{1 - \beta \cos\theta} \right) - K_u (\cos\theta \leftrightarrow -\cos\theta) - K_s [1 + \kappa_t] \right] \\ &= \mathcal{M}(g_R^a g_R^b \rightarrow t_L^k \bar{t}_R^l) = \mathcal{M}(g_L^a g_L^b \rightarrow t_R^k \bar{t}_L^l) = \mathcal{M}(g_R^a g_R^b \rightarrow t_R^k \bar{t}_L^l), \end{aligned} \quad (\text{A.7}) \quad (\text{A.8})$$

$$\mathcal{M}(g_L^a g_R^b \rightarrow t_L^k \bar{t}_R^l) = -\beta \left[\frac{K_t}{1 - \beta \cos\theta} + \frac{K_u}{1 + \beta \cos\theta} \right] \sin\theta (1 + \cos\theta) [1 + \kappa_t] = \mathcal{M}(g_R^a g_L^b \rightarrow t_R^k \bar{t}_L^l), \quad (\text{A.9})$$

$$\mathcal{M}(g_R^a g_L^b \rightarrow t_L^k \bar{t}_R^l) = \mathcal{M}(g_L^a g_R^b \rightarrow t_L^k \bar{t}_L^l) = \mathcal{M}(g_L^a g_R^b \rightarrow t_L^k \bar{t}_R^l)[\cos\theta \leftrightarrow -\cos\theta, \beta \leftrightarrow -\beta]. \quad (\text{A.10})$$

The color factors are defined as $K_q = g_s^2 T_{ij}^a * T_{kl}^a$, $K_s = ig_s^2 f^{abc} T_{kl}^c$, $K_t = g_s^2 T_{ki}^a T_{il}^b$, and $K_u = g_s^2 T_{il}^a T_{ki}^b$, the evaluation of their products follows the standard SU(3) rules, resulting in $K_q^* K_q = g_s^4 N_c C_F/2$, $K_s^* K_s = g_s^4 N_c^2 C_F$, $K_t^* K_t = K_u^* K_u = g_s^4 N_c C_F^2$, and $K_s^* K_t = -K_s^* K_u = g_s^4 N_c^2 C_F/2$, $K_t^* K_u = -g_s^4 C_F/2$.

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