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No like-sign tops at Tevatron: Constraints on extended models and implications for the $t\bar{t}$ asymmetry

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ABSTRACT

We use a recent upper limit from CDF on like-sign top pair production to place constraints on general new vector bosons and scalars mediating $uu \to tt$. The possible vector bosons comprise neutral colour singlets or octets, and charge 4/3 colour triplets or sextets, whereas the new scalars can be neutral colour singlets or octets and charge 4/3 colour sextets. We also estimate the expected bounds from like-sign top pair production at LHC in the near future. Then, we address the implications of these limits for the forward–backward asymmetry in $t\bar{t}$ production measured at Tevatron. In particular, we find that models explaining the observed asymmetry by the exchange of a single t-channel heavy Z' boson are already excluded. On the other hand, light Z' bosons with a mass $M_{Z'} \simeq 150$ GeV, which could also account for a recent CDF dijet excess in W + jet production, are barely allowed.

1. Introduction

The production of like-sign top quark pairs would be a striking signature of physics beyond the Standard Model (SM). At hadron colliders, charge conservation implies that tt pairs can only be produced from initial up or charm quarks. Hence, proton–proton colliders are better suited for studying this signal, in particular in $uu \to tt$. In fact, like-sign top production is a golden channel for early discoveries at the Large Hadron Collider (LHC) [1,2]. Maybe more surprising is the fact that the high statistics accumulated by Tevatron already allows to extract useful limits on various SM extensions mediating this process.

Recently, the CDF Collaboration has set a limit on like-sign top production at Tevatron, using a luminosity of 6.1 fb^{-1} [3],

$$\sigma(tt + \bar{t}\bar{t}) \times \mathrm{Br}(W \to \ell\nu)^2 < 54\ \mathrm{fb}, \tag{1}$$

with a 95% confidence level (CL). As we shall show, this limit puts significant constraints on the different SM extensions that can potentially give observable contributions to $uu \to tt$. These extensions must contain boson fields mediating this process at tree level. Assuming renormalizable interactions, these fields can be either extra vector bosons or extra scalars. Since the new interactions must fulfil the $\mathrm{SU}(3)_C \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ gauge invariance of the SM, the quantum numbers of the new fields are not arbitrary. It is easy to check that the only possibilities for particles of spin 1 are [4]:

- a neutral colour-singlet Z', which can be in a singlet $\mathrm{SU}(2)_L$ representation (denoted here as $\mathcal{B}_\mu$) or belong to a triplet $\mathcal{W}_\mu$;
- a neutral colour-octet g', either an isosinglet $\mathcal{G}_\mu$ or member of an isotriplet $\mathcal{H}_\mu$;
- the charge 4/3 component of a colour-triplet isodoublet $\mathcal{Q}^5_\mu$;
- the charge 4/3 component of a colour-sextet isodoublet $\mathcal{Y}^5_\mu$.

On the other hand, the possible new particles of spin 0 are

- the neutral scalar of an isodoublet, which can be a colour singlet ϕ or octet Φ;
- the charge 4/3 component of a colour-sextet, either an isosinglet Ω^4 or included in an isotriplet Σ.

Obviously, charge 4/3 and neutral particles are exchanged in s and t channels, respectively. In this Letter, we translate the CDF upper limit in Eq. (1) into constraints on the extended sectors enumerated above. Furthermore, we compare with potential LHC measurements with 2010 data (35 pb^{-1}) and 2011 data (1 fb^{-1}). In particular, we place direct limits on flavour-changing neutral (FCN) couplings of new neutral vector bosons Z', g' as well as neutral scalars. In order to unify the computations for the different fields, we make use of effective field theory in the following manner: we first integrate out the new heavy states and obtain their contribution to $uu \to tt$ in terms of four-fermion operators; then, we obtain the cross section in terms of effective operator coefficients. In an appendix we discuss the range of validity of this approximation and present exact results for light Z' bosons.

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Table 1
Vector bosons and scalar representations mediating $uu \to tt$.

Symbol	Representation	Interaction Lagrangian	Symmetry
$\mathcal{B}_\mu$	$(1,1)_0$	$-(g^q_{ij}\bar{q}_{Li}\gamma^\mu q_{Lj} + g^u_{ij}\bar{u}_{Ri}\gamma^\mu u_{Rj} + g^d_{ij}\bar{d}_{Ri}\gamma^\mu d_{Rj})\mathcal{B}_\mu$	$g = g^\dagger$
$\mathcal{W}_\mu$	$(1,\mathrm{Adj})_0$	$-g_{ij}\bar{q}_{Li}\gamma^\mu \tau^I q_{Lj}\mathcal{W}^I_\mu$	$g = g^\dagger$
$\mathcal{G}_\mu$	$(\mathrm{Adj},1)_0$	$-(g^q_{ij}\bar{q}_{Li}\gamma^\mu \frac{\lambda^a}{2} q_{Lj} + g^u_{ij}\bar{u}_{Ri}\gamma^\mu \frac{\lambda^a}{2} u_{Rj} + g^d_{ij}\bar{d}_{Ri}\gamma^\mu \frac{\lambda^a}{2} d_{Rj})\mathcal{G}^a_\mu$	$g = g^\dagger$
$\mathcal{H}_\mu$	$(\mathrm{Adj},\mathrm{Adj})_0$	$-g_{ij}\bar{q}_{Li}\gamma^\mu \tau^I \frac{\lambda^a}{2} q_{Lj}\,\mathcal{H}^{aI}_\mu$	$g = g^\dagger$
$\mathcal{Q}^5_\mu$	$(3,2)_{-\frac{5}{6}}$	$-g_{ij}\varepsilon_{abc}\bar{u}_{Rib}\gamma^\mu \epsilon q^c_{Ljc}\,\mathcal{Q}^{5a\dagger}_\mu + \text{h.c.}$	–
$\mathcal{Y}^5_\mu$	$(\bar{6},2)_{-\frac{5}{6}}$	$-g_{ij}\frac{1}{2}[\bar{u}_{Ria}\gamma^\mu \epsilon q^c_{Ljb} + \bar{u}_{Rib}\gamma^\mu \epsilon q^c_{Lja}]\mathcal{Y}^{5ab\dagger}_\mu + \text{h.c.}$	–
ϕ	$(1,2)_{-\frac{1}{2}}$	$-g^u_{ij}\bar{q}_{Li}u_{Rj}\phi - g^d_{ij}\bar{q}_{Li}d_{Rj}\tilde{\phi} + \text{h.c.}$	–
Φ	$(\mathrm{Adj},2)_{-\frac{1}{2}}$	$-g^u_{ij}\bar{q}_{Li}\frac{\lambda^a}{2}u_{Rj}\Phi^a - g^d_{ij}\bar{q}_{Li}\frac{\lambda^a}{2}d_{Rj}\tilde{\Phi}^a + \text{h.c.}$	–
Ω^4	$(\bar{6},1)_{-\frac{4}{3}}$	$-g_{ij}\frac{1}{2}[\bar{u}_{Ria}u^c_{Rjb} + \bar{u}_{Rib}u^c_{Rja}]\Omega^{4ab\dagger} + \text{h.c.}$	$g = g^T$
Σ	$(\bar{6},\mathrm{Adj})_{-\frac{1}{3}}$	$-g_{ij}\frac{1}{2}[\bar{q}_{Lia}\tau^I \epsilon q^c_{Ljb} + \bar{q}_{Lib}\tau^I \epsilon q^c_{Lja}]\Sigma^{Iab\dagger} + \text{h.c.}$	$g = g^T$

We devote special attention to the implications of these limits for the $t\bar{t}$ forward–backward (FB) asymmetry at Tevatron. One mechanism that has been suggested to enhance this asymmetry and thus explain the measured values, in particular $A_{\mathrm{FB}} = 0.475 \pm 0.114$ for $m_{t\bar{t}} > 450$ GeV [5], is the exchange of a flavour-violating Z' boson in the t channel [6–11]. It is well known that, for a single real Z', this automatically implies like-sign tt production. The relation between the two processes can easily be understood, without precise knowledge of the details to be given below, by taking the CP conjugate of one of the two vertices (the fact that the vector boson is real is crucial). Here, we show that the non-observation of like-sign tops at Tevatron already rules out these models as the sole explanation of the Tevatron $t\bar{t}$ asymmetry, except for very light Z' masses which, interestingly, are consistent with a recent CDF dijet excess [12]. We also note that this direct relation between tt production and the value of A_{FB} does not hold any longer when more than one Z' boson is present, as in the model in Ref. [13]. Among the other new particles that can enhance the FB asymmetry, some belong to the list above. Therefore, the limits from tt production reduce the parameter space of extensions with these particles.

2. New bosons and tt production

The process $uu \to tt$ is absent in the SM at the tree level but it can be mediated by new vector bosons in six possible $\mathrm{SU}(3)_C \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ (irreducible) representations [4], or by scalars in four possible representations.[1] They are all collected in Table 1, where in the first column we write the symbol used to label them. The relevant interaction Lagrangian is included as well. In the last column we display the symmetry properties, if any, of the coupling matrices g_{ij}. We use standard notation with left-handed doublets q_{Li} and right-handed singlets u_{Ri}, d_{Ri}; τ^I are the Pauli matrices, λ^a the Gell-Mann matrices normalised to $\mathrm{tr}(\lambda^a\lambda^b) = 2\delta_{ab}$ and $\tilde{\phi} = \epsilon\phi$, $\psi^c = C\bar{\psi}^T$, with $\epsilon = i\tau^2$ and C the charge conjugation matrix. The subindices a, b, c denote colour. The bosons $\mathcal{Q}^5_\mu$, $\mathcal{Y}^5_\mu$, Ω^4, Σ are created by uu fusion and exchanged in the s channel, while the rest are exchanged in the t (and u) channels.

If the new particles are heavy, their contribution to $uu \to tt$ can be described by an effective low-energy Lagrangian. There are only five independent four-fermion operators contributing to $uu \to tt$ [15], which can be taken as O^{1313}_{qq}, $O^{1313}_{qq'}$, O^{1313}_{uu}, O^{1313}_{qu} and $O^{1313}_{qu'}$ (see Appendix A). The coefficients of these operators are given in Table 2 for each of the vector boson and scalar representations in Table 1.[2] The new physics scale Λ equals the mass of the new boson or multiplet in each case.

Table 2
Effective operator coefficients involved in like-sign tt production for each vector boson and scalar representation. The new physics scale Λ equals the mass of the new particle or multiplet.

	C^{1313}_{qq}	$C^{1313}_{qq'}$	C^{1313}_{uu}	C^{1313}_{qu}	$C^{1313}_{qu'}$
$\mathcal{B}_\mu$	$-(g^q_{13})^2$	–	$-(g^u_{13})^2$	–	$2g^q_{13}g^u_{13}$
$\mathcal{W}_\mu$	g^2_{13}	$-2g^2_{13}$	–	–	–
$\mathcal{G}_\mu$	$\frac{1}{6}(g^q_{13})^2$	$-\frac{1}{2}(g^q_{13})^2$	$-\frac{1}{3}(g^u_{13})^2$	$g^q_{13}g^u_{13}$	$-\frac{1}{3}g^q_{13}g^u_{13}$
$\mathcal{H}_\mu$	$-\frac{7}{6}g^2_{13}$	$\frac{5}{6}g^2_{13}$	–	–	–
$\mathcal{Q}^5_\mu$	–	–	–	$2g_{11}g^*_{33}$	$-2g_{11}g^*_{33}$
$\mathcal{Y}^5_\mu$	–	–	–	$-g_{11}g^*_{33}$	$-g_{11}g^*_{33}$
ϕ	–	–	–	$g^u_{13}g^{u*}_{31}$	–
Φ	–	–	–	$-\frac{1}{6}g^u_{13}g^{u*}_{31}$	$\frac{1}{2}g^u_{13}g^{u*}_{31}$
Ω^4	–	–	$g_{11}g^*_{33}$	–	–
Σ	$g_{11}g^*_{33}$	$g_{11}g^*_{33}$	–	–	–

Within this model-independent approach, the cross section for $uu \to tt$ can be compactly written in terms of five effective operator coefficients, the new physics scale Λ and numerical constants E_{1-3} that result from phase space integration and convolution with parton density functions (PDFs),

$$\sigma(tt) = \frac{E_1}{\Lambda^4}\left[\left|C^{1313}_{qq} + C^{1313}_{qq'}\right|^2 + \left|C^{1313}_{uu}\right|^2\right] + \frac{E_2}{\Lambda^4}\left[\left|C^{1313}_{qu'}\right|^2 + \left|C^{1313}_{qu}\right|^2 + \frac{2}{3}\,\mathrm{Re}\,C^{1313}_{qu'}C^{1313*}_{qu}\right] + \frac{E_3}{\Lambda^4}\left\{\mathrm{Re}\,C^{1313}_{qu'}C^{1313*}_{qu} + \frac{1}{6}\left[\left|C^{1313}_{qu'}\right|^2 + \left|C^{1313}_{qu}\right|^2\right]\right\}. \quad (2)$$

The lowest order contributions arise at order $1/\Lambda^4$, since the SM amplitudes vanish [16]. Clearly, higher-order operators can be neglected as long as the extra particles are heavy enough. In Appendix B we discuss in more detail the range of validity of this approximation.

[1] Notice that for scalars mixing up-type quarks there are two additional $\mathrm{SU}(3)_C$ triplet representations [14]: isotriplets coupling to $q_{Li}q^c_{Li}$ and isosinglets coupling to $u_{Ri}u^c_{Rj}$. However, their coupling matrices are anti-symmetric and diagonal couplings to uu^c and tt^c vanish.

[2] For $\mathcal{W}_\mu$ and $\mathcal{H}_\mu$ the normalisation in the Lagrangian differs from Ref. [4] by a factor of two, to simplify the presentation of the limits.