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Probing the Exotic Particle Content
Beyond the Standard Model

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Abstract

We explore the possible exotic particle content beyond the standard model by examining all its scalar bilinear combinations. We categorize these exotic scalar fields and show that without the suppression of (A) their Yukawa couplings with the known quarks and leptons, and (B) the trilinear couplings among themselves, most are already constrained to be very heavy from the nonobservation of proton decay and neutron-antineutron oscillations, the smallness of $K^0 - \overline{K^0}$, $D^0 - \overline{D^0}$ and $B_d^0 - \overline{B_d^0}$ mixing, as well as the requirement of a nonzero baryon asymmetry of the universe. On the other hand, assumption (B) may be naturally violated in many models, especially in supersymmetry, hence certain exotic scalars are allowed to be below a few TeV in mass and would be easily detectable at planned future hadron colliders. In particular, large cross sections for the distinctive processes like $\bar{p}p \to tt, \bar{t}c$ and $pp \to tt, bb$ would be expected at the Fermilab Tevatron and CERN LHC, respectively.

1. Introduction

The quarks and leptons of the minimal standard model are familiar fixtures of particle physics. Under the standard $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge group, they transform as follows

$$\begin{pmatrix} u_i \\ d_i \end{pmatrix}_L \sim (3, 2, 1/6), \quad u_{iR} \sim (3, 1, 2/3), \quad d_{iR} \sim (3, 1, -1/3); \tag{1}$$

$$\begin{pmatrix} \nu_i \\ l_i \end{pmatrix}_L \sim (1, 2, -1/2), \quad l_{iR} \sim (1, 1, -1). \tag{2}$$

In the above, the index i denotes the known 3 families. Only one scalar bilinear combination of these fermions is required in the minimal model, *i.e.*

$$(\phi^+, \phi^0) \sim (1, 2, 1/2), \tag{3}$$

which couples $\overline{(u_i, d_i)}_L$ to u_{jR} and d_{jR}, as well as $\overline{(\nu_i, l_i)}_L$ to l_{jR}. As ϕ^0 acquires a nonzero vacuum expectation value, $v \equiv \langle \phi^0 \rangle$, the quarks and leptons obtain masses and there is mixing among the quark families, but not among the lepton families. The other possible bilinear combinations may give rise to unobserved phenomenology. Therefore their masses and couplings to fermions have been stringently constrained from low energy and collider processes [1, 2].

However, the scalar bilinear combination which couples $(\nu_i, l_i)_L$ to $(\nu_j, l_j)_L$, *i.e.*,

$$(\xi^{++}, \xi^+, \xi^0) \sim (1, 3, 1), \tag{4}$$

has been shown to have important phenomenological implications [3] beyond these considered in the previous works. The new ingredient is the trilinear coupling of ξ to the standard model Higgs boson which, together with the Yukawa couplings to fermions, may give rise to baryogenesis or, alternatively, to wash away the baryon asymmetry of the universe. Also, for a very large M_ξ, it is natural for ξ^0 to acquire a tiny vacuum expectation value, thereby allowing neutrinos to obtain very small Majorana masses and to mix with one another.

In this paper, we extend the above consideration to all possible scalar bilinear combinations of the quarks and leptons of the minimal standard model including diquark scalars. We categorize these exotic scalar fields and ascertain their various contributions to physics within and beyond the standard model. From a general consideration we can show that most are already constrained by present experimental data to be very heavy. This is based on two assumptions: (A) the Yukawa couplings of these exotic scalars with the known quarks and leptons are all of order unity, and (B) the trilinear couplings of these exotic scalars among themselves are of order the mass scale of the heaviest particle involved. Whereas assumption (A) is a natural one in almost any model, assumption (B) is subject to many other possible qualifications. For example, exact supersymmetry often forbids the existence of trilinear scalar couplings, in which case any such coupling should not exceed the scale of supersymmetry breaking, which may be very small compared to the mass of the heaviest particle involved. Relaxing assumption (B) allows certain exotic scalar particles to be below a few TeV in mass and they would easily be detectable at planned future hadron colliders. In particular, large cross sections for the flavor-changing neutral-current (FCNC) process $\bar{p}p \rightarrow \bar{t}c$, and for the quark flavour violating resonance processes $pp \rightarrow tt, bb$ may be expected at the Fermilab Tevatron and CERN LHC, respectively. Also, the cross section of the resonance process $\bar{p}p \rightarrow tt$ at Tevatron, though sea quark suppressed, may be large enough to provide an observable excess of the same-sign dilepton final states indicating clearly for the new physics.

The outline of the paper is as follows. After the Introduction we classify the exotic scalar bilinears and discuss constraints on their masses and couplings. In Section 3 we present our results on diquark mediated processes at hadron colliders. We conclude in Section 4.

Representation	qq	$\bar{q}\bar{l}$	$q\bar{l}$	ll	Comment
$(1,1,-1)$				X	$n_f \geq 2$
$(1,3,-1)$				X	neutrino masses
$(1,1,-2)$				X	e^-e^- collider
$(3^*,1,1/3)$	X	X			$p \to e^+\pi^0$
$(3^*,3,1/3)$	X	X			$n_f \geq 2$
$(3^*,1,4/3)$	X	X			$n_f \geq 2$
$(3^*,1,-2/3)$	X				$n_f \geq 2$
$(3,2,1/6)$			X		HERA
$(3,2,7/6)$			X		anomaly
$(6,1,-2/3)$	X				neutron-
$(6,1,1/3)$	X				antineutron
$(6,1,4/3)$	X				oscillations
$(6,3,1/3)$	X				$K^0 - \overline{K^0}$ and
$(8,2,1/2)$					$D^0 - \overline{D^0}$ mixing

Table 1: Exotic scalar particles beyond the standard model.

2. Classification and Indirect Constraints

In this section we classify the exotic scalars and constrain their masses and couplings from low-energy phenomenology and from cosmological considerations. In Table 1 we list all possible scalar representations which are products of two known fermion representations. The colour singlet fields have been listed before in [1] and the colour triplets and sextets in [4]. We cite the most stringent constraints on their masses and interactions from previous works and add several new bounds.

2.1. Dileptons

First we consider the three scalar dileptons.

(1,1,−1): This scalar singlet couples to $\nu_i l_j - l_i \nu_j$ where $i \neq j$, hence the number of families n_f must be 2 or greater. It contributes to μ and τ decays leading to the constraints [1, 5]

$$\mu \to e\gamma \quad : \quad \frac{M_X}{(f_{e\tau} f_{\mu\tau})^{1/2}} \gtrsim 1.6 \times 10^4 \text{ GeV}\,;$$

$$G_\tau \quad : \quad \frac{M_X}{f_{e\tau}} \gtrsim 2.2 \times 10^3 \text{ GeV}, \quad \frac{M_X}{f_{\mu\tau}} \gtrsim 3.2 \times 10^3 \text{ GeV}\,; \tag{5}$$

where M_X denotes the scalar mass. Similar bounds can be derived from the precision electroweak data. If it is assumed that this exotic scalar contributes less than 0.1% of the μ decay rate, then its mass divided by $|f_{e\mu}|$ is greater than 1.1×10^4 GeV [6]. It has been also shown to be useful for the radiative generation of neutrino masses [7].

In models with two or more Higgs doublets, there are in general trilinear couplings given by

$$\mathcal{L} = h_1 X^+ (\nu_i l_j - l_i \nu_j) + h_2 M_X X^- (\phi_i^+ \phi_j^0 - \phi_i^0 \phi_j^+)\,. \tag{6}$$

Lepton-number conservation is thus violated and if h_1 and h_2 are unsupressed (which is the

case in most models of radiative neutrino masses), the resulting interactions will erase any lepton asymmetry of the universe before the onset of the electroweak phase transition. This will deprive the anomalous sphaleron-induced processes from converting an existing lepton asymmetry into a baryon asymmetry of the universe [8]. To forestall this eventuality, X must be heavy:

$$\frac{M_X}{h_1^2} \gtrsim 10^{15} \text{ GeV} \quad \text{or} \quad \frac{M_X}{h_2^2} \gtrsim 10^{15} \text{ GeV} \quad \text{and} \quad \frac{M_X}{h_1^2 h_2^2} \gtrsim 10^{16} \text{ GeV}. \tag{7}$$

(1,3,−1): This scalar triplet couples according to $\xi^0 \nu_i \nu_j + \xi^+ (\nu_i l_j + l_i \nu_j)/\sqrt{2} + \xi^{++} l_i l_j$ and allows us to have the most general 3×3 Majorana neutrino mass matrix without right-handed neutrinos. It also couples to the standard Higgs doublet according to $\bar{\xi}^0 \phi^0 \phi^0 + \sqrt{2}\xi^- \phi^+ \phi^0 + \xi^{--} \phi^+ \phi^+$ and it has been shown [3] that a tiny $\langle \xi^0 \rangle$ of a few eV or less is obtained for $M_\xi \gtrsim 10^{13}$ GeV. If there are two such triplets, a successful leptogenesis scenario [3] for the baryon asymmetry of the universe may also be obtained.

Alternatively, it is possible that ξ does not couple to ϕ, in which case lepton number is conserved, but not lepton flavor. In this latter scenario, the most stringent constraints come from $\mu \rightarrow ee\bar{e}$ [1] decay and $\mu - e$ conversion in nuclei [9], which are given by

$$\frac{M_X}{(f_{ee} f_{e\mu})^{1/2}} \gtrsim 2.2 \times 10^5 \text{ GeV} \quad \text{and} \quad \frac{M_X}{(f_{ea} f_{\mu a})^{1/2}} \gtrsim 4 \times 10^4 \text{ GeV}, \tag{8}$$

respectively, where, $a = e, \mu$ or τ. The doubly charged component of the triplet may give rise to unique lepton-flavor-violating signatures at collider experiments [1, 10, 11]. In particular, for certain neutrino mass patterns where the neutrino decay lifetimes are constrained, there are lower bounds on the cross sections of the processes $e^- e^- (\mu^- \mu^-) \rightarrow l^- l^-$ such that they *must* be seen at future lepton colliders [11].

(1,1,−2): This doubly charged scalar singlet couples to $l_{iR} l_{jR}$ symmetrically. It is relevant for $e^- e^-$ colliders [1, 10], but it also contributes to lepton-flavor-changing processes such as

$\mu - e$ conversion, $\mu \rightarrow e\gamma$, and $\mu \rightarrow eee$. The bounds from these processes are the same as those for the (1,3,–1) scalar.

If this exotic scalar X coexists with the previously discussed $(1, 1, -1)$, now call it Y, which is much lighter than X, then there may be lepton-number-violating couplings given by

$$\mathcal{L} = h_1 X^{++} l_R l_R + h_2 M_X X^{++} Y^- Y^-.$$

In this case, X has to be again heavy to satisfy the leptogenesis constraints [8],

$$\frac{M_X}{h_1^2} \gtrsim 10^{15} \text{ GeV} \quad \text{or} \quad \frac{M_X}{h_2^2} \gtrsim 10^{15} \text{ GeV} \quad \text{and} \quad \frac{M_X}{h_1^2 h_2^2} \gtrsim 10^{16} \text{ GeV}. \tag{9}$$

2.2. Triplet Diquarks and Leptoquarks

The decomposition of 3×3 being $3^* + 6$ under SU(3), there are two types of scalar diquarks. The antisymmetric combination 3^* may also couple to an antiquark and an antilepton as shown in Table 1. This means that proton decay is always possible, if not at tree level [4] then in one loop as we show below.

(3*,1,1/3): This exotic scalar singlet couples to $u_{iL}^\alpha d_{jL}^\beta - u_{iL}^\beta d_{jL}^\alpha - d_{iL}^\alpha u_{jL}^\beta + d_{iL}^\beta u_{jL}^\alpha$ where $\alpha, \beta = 1, 2, 3$ are color indices. This combination is antisymmetric under both $SU(3)_C$ and $SU(2)_L$, but symmetric under the interchange of families. It also couples to $u_{iR}^\alpha d_{jR}^\beta - u_{iR}^\beta d_{jR}^\alpha$ as well as $\bar{u}_{iL} \bar{l}_{jL} - \bar{d}_{iL} \bar{\nu}_{jL}$ and $\bar{u}_{iR} \bar{l}_{jR}$. Hence it mediates proton decay with the effective operators $(uudl)$ and $(udd\nu)$. Using $\tau(p \rightarrow \pi^0 e^+) \gtrsim 9 \times 10^{32}$ years, one finds the constraint

$$\frac{M_X}{|f_{ud} f_{ul}|^{1/2}} \gtrsim 10^{16} \text{ GeV}. \tag{10}$$

Although this is the strongest bound, it is applicable only to u, d, and s quarks. However, if other couplings are large, there will be fast baryon-number-violating interactions in the early universe and a baryon asymmetry cannot be maintained. Thus strong bounds exist for

all quarks, namely

$$\frac{M_X}{f^2} \gtrsim 10^{15} \text{ GeV}. \tag{11}$$

On the other hand, if baryon asymmetry is generated after these scalars have decayed away, then these bounds will not be valid.

(3*,3,1/3): This triplet scalar diquark couples to a symmetric combination of $SU(2)_L$ doublets, hence it must be antisymmetric in its coupling to quark families. It mediates proton decay with the effective operators $(uusl)$ and $(uds\nu)$. From the nonobservation [5] of processes such as $p \to e^+ K^0$ or $p \to \mu^+ K^0$, we find its mass divided by $|f_{us}f_{ul}|^{1/2}$ to be also greater than about 10^{16} GeV.

There will also be constraints from the survival of the baryon asymmetry of the universe in this case, which are of similar magnitude but applicable to all generations, $M_X/f^2 \gtrsim 10^{15}$ GeV.

(3*,1,4/3): As a diquark, this exotic scalar couples only to $u_{iR}u_{jR}$, where $i \neq j$. As a leptoquark, it couples only to $\bar{d}_{iR}\bar{l}_{jR}$. Since the c or t quark must appear in a tree-level effective operator, proton decay here requires a one-loop diagram such as the one depicted in Fig. 1. The resulting effective operator $(uusl)$ is suppressed by the factor $G_F m_c m_s V_{cs} V_{us}/16\pi^2$. We find thus the constraint on the mass of this scalar to be

$$\frac{M_X}{|f_{uc}f_{sl}|^{1/2}} \gtrsim 5.3 \times 10^{11} \text{ GeV}. \tag{12}$$

There will also be constraints from baryogenesis in this case, which are of similar magnitude but applicable to all generations, $M_X/f^2 \gtrsim 10^{15}$ GeV.

(3*,1,–2/3): Because there is no singlet neutrino, this exotic scalar acts as a diquark, but not a leptoquark. It couples only to $d_{iR}d_{jR}$, where $i \neq j$. Hence it does not appear to mediate proton decay. However, the trilinear scalar interaction

$$(3^*, 1, -2/3)(3^*, 1, 1/3)(3^*, 1, 1/3) \tag{13}$$

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is generally allowed and the effective operators $(uds\bar\nu)$ and $(dds\bar l)$ may be obtained in one loop as depicted in Fig. 2. Note that these have the selection rule $\Delta B = -\Delta L$ instead of the usual $\Delta B = \Delta L$. Let the trilinear coupling be μ and the mass of the $(3^*, 1, 1/3)$ scalar be M_Y, then the loop suppression factor for decays such as $n \to e^- K^+$ and $p \to \nu K^+$ is $\mu m_p / 16\pi^2 M_Y^2$, where m_p is the proton mass. Taking the natural value of μ to be equal to $M_Y \simeq 10^{16}$ GeV, we find the mass of the $(3^*, 1, -2/3)$ scalar divided by $|f_{ds} f_{ud} f_{d\nu}|^{1/2}$ to be greater than about

$$\frac{M_X}{|f_{ds} f_{ud} f_{d\nu}|^{1/2}} \gtrsim \frac{(m_p M_Y)^{1/2}}{4\pi} \simeq 7.7 \times 10^6 \text{ GeV.} \tag{14}$$

In addition to this bound, if the scalar $(3^*,1,-2/3)$ has a trilinear coupling as assumed above, then its interaction with two quarks and simultaneously with two $(3,1,-1/3)$ scalars will break baryon number. To prevent an existing baryon asymmetry from being washed out, we would then get a much stronger bound on its mass and couplings,

$$M_X^3/\mu^2 \gtrsim 10^{15} \text{ GeV} \quad \text{or} \quad M_X/f^2 \gtrsim 10^{15} \text{ GeV} \quad \text{and} \quad M_X^3/f^2\mu^2 \gtrsim 10^{16} \text{ GeV,} \tag{15}$$

assuming that $M_Y \ll M_X$. The same consideration applies to the $(3^*,1,4/3)$ scalar if the trilinear coupling $(3^*,1,4/3)$ $(3^*,1,-2/3)$ $(3^*,1,-2/3)$ also exists.

2.3. Leptoquarks

There are two scalar leptoquark representations, both of which are $SU(2)_L$ doublets. They are $q\bar l$ (and not ql) combinations, and may be relevant [12] to the erstwhile HERA anomaly [13] of excess events at large momenta in $e^+ p$ scattering. Reviews on the constraints on leptoquark couplings to fermions can be found in [2].

(3,2,1/6): This scalar leptoquark couples to $(\bar l_i, \bar\nu_i)_L d_{jR}$. It mediates lepton flavor-changing processes such as $K^0 \to e^+ \mu^-$. From the experimental upper bound [5] of the latter, one finds its mass divided by $|f_{de} f_{s\mu}|^{1/2}$ to be greater than about 2.4×10^5 GeV.

(3,2,7/6): This scalar leptoquark couples to $(\bar{l}_i, \bar{\nu}_i)_L u_{jR}$ and $\bar{l}_{iR}(u_j, d_j)_L$. The same constraint applies here as in the previous case.

2.4. Sextet Diquarks

The other type of scalar diquark is an $SU(3)_C$ sextet. It couples symmetrically to two $SU(3)_C$ quark triplets. It has not received much attention in the past, but it is potentially an important hint for physics beyond the standard model. These scalars contribute at tree level to $K^0 - \overline{K^0}$, $D^0 - \overline{D^0}$ and $B_d^0 - \overline{B_d^0}$ mixings, and neutron-antineutron oscillations naturally occur from the trilinear scalar interactions $(6, 1, -2/3)(6, 1, 1/3)^2$, $(6, 1, -2/3)(6, 3, 1/3)^2$ and $(6, 1, 4/3)(6, 1, -2/3)^2$.

(6,1,–2/3): This exotic scalar couples to $d_{iR}d_{jR}$ symmetrically, whereas $(3^*, 1, -2/3)$ does so antisymmetrically as already discussed. Comparing the two cases, we see that a similar loop diagram to Fig. 2 would generate the effective operators $(udd\bar{\nu})$ and $(ddd\bar{l})$. Hence its mass divided by $|f_{dd} f_{ud} f_{d\nu}|^{1/2}$ should also be greater than about 7.7×10^6 GeV. The effective $dd \to ss$ and $dd \to bb$ transitions which induce $K^0 - \overline{K^0}$ and $B_d^0 - \overline{B_d^0}$ mixings give somewhat weaker bounds, 1.5×10^6 GeV and 4.6×10^5 GeV, on its mass over the couplings $|f_{dd} f_{ss}|^{1/2}$ and $|f_{dd} f_{bb}|^{1/2}$, respectively.

(6,1,1/3): This exotic scalar couples to $u_{iL}^\alpha d_{jL}^\beta + u_{iL}^\beta d_{jL}^\alpha - d_{iL}^\alpha u_{jL}^\beta - d_{iL}^\beta u_{jL}^\alpha$, which is antisymmetric under $SU(2)_L$ and the interchange of families. It also couples to $u_{iR}^\alpha d_{jR}^\beta + u_{iR}^\beta d_{jR}^\alpha$. Since the trilinear scalar interaction

$$(6, 1, 1/3)(6, 1, -2/3)(6, 1, 1/3) \tag{16}$$

is generally allowed, the analog of Fig. 2 is also possible here. We find the mass of the $(6, 1, 1/3)$ scalar divided by $|f_{ud} f_{db} f_{b\nu}|^{1/2}$ to be also greater than about 7.7×10^6 GeV.

As discussed before, the coexistence of a trilinear scalar coupling and a Yukawa coupling

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with the quarks will cause fast baryon-number violation. Forbidding it will impose a stronger bound on the mass and couplings of the heavier scalar, $M_X/f^2 \gtrsim 10^{15}$ GeV.

(6,1,4/3): This exotic scalar couples to $u_{iR}u_{jR}$ symmetrically. It does not participate in proton decay. However, from the allowed trilinear scalar interaction

$$(6, 1, 4/3)(6, 1, -2/3)(6, 1, -2/3), \tag{17}$$

we obtain the effective $(udd)^2$ operator which induces neutron-antineutron oscillations. From the present experimental upper bound [14] of the latter, we find the mass of the $(6, 1, 4/3)$ scalar divided by $|f_{dd}||f_{uu}|^{1/2}$ to be greater than about 1.3×10^3 GeV. However, this limit is superceded by considering the effective $uu \to cc$ transition which induces $D^0 - \overline{D^0}$ mixing. From the experimental upper bound [5] $\Delta m_D < 1.4 \times 10^{-10}$ MeV, we find the mass of the $(6, 1, 4/3)$ scalar divided by $|f_{uu}f_{cc}|^{1/2}$ to be greater than about 7.3×10^5 GeV. Again the consideration of baryogenesis applies as before.

(6,3,1/3): This exotic $SU(2)_L$ triplet has all the couplings of the previous 3 singlets. It partcipates in the trilinear scalar interaction

$$(6, 3, 1/3)(6, 3, 1/3)(6, 1, -2/3), \tag{18}$$

which induces $n - \bar{n}$ oscillations. Hence its mass divided by $|f_{ud}|^{1/2}|f_{dd}|^{1/4}$ should be greater than about 1.0×10^5 GeV. The baryogenesis bound will also be similar as in the other cases. Previously quoted bounds coming from the $K^0 - \overline{K^0}$, $D^0 - \overline{D^0}$ and $B_d^0 - \overline{B_d^0}$ mixings apply also in this case..

2.5. Octet

(8,2,1/2): This $SU(3)_C$ octet couples $(\overline{u_i^\alpha, d_i^\alpha})_L$ to u_{jR}^β and d_{jR}^β. It carries neither baryon nor lepton number. Consider the neutral member of this $SU(2)_L$ doublet. Unlike the standard neutral Higgs boson, it has in general nondiagonal couplings to quarks. Hence there is an

effective $d\bar{s} \to s\bar{d}$ transition which induces $K^0 - \overline{K^0}$ mixing. We find thus the mass of this scalar divided by $|f_{ds}f_{sd}|^{1/2}$ to be greater than about 1.5×10^6 GeV. Bounds from the other neutral meson mixings are also valid here.

All the bounds coming from the survival of the baryon asymmetry of the universe may be avoided to some extent, if a baryon asymmetry of the universe is generated after the heavy scalars (whose interactions violate baryon number) have all decayed away. However, in this case, there will still be a bound on the masses of these heavy scalars, which is the scale of baryogenesis. If these scalars are lighter than the scale at which baryon asymmetry of the universe is generated, then their interactions will erase the asymmetry thus generated. Hence the bounds from the constraints of baryogenesis can at most be made milder by generating a baryon asymmetry of the universe at a lower energy scale.

3. Diquarks at Hadron Colliders

In general, the above derived bounds on the couplings and masses of the $SU(3)_C$ triplet and sextet diquarks depend on the attributed flavor indices and, most importantly, on the assumption of the existence of the trilinear couplings. For example, if the latter are absent or strongly suppressed as in some models of supersymmetry, then the most stringent bounds from proton decay, neutron-antineutron oscillations and baryogenesis can be evaded in the case of sextet diquarks. The bounds coming from the measurements of $K^0 - \overline{K^0}$, $D^0 - \overline{D^0}$ and $B_d^0 - \overline{B_d^0}$ mixings still apply, but if we take care of them by suppressing the diquark diagonal couplings involving the first two families, then no bounds at all exist for the others. The diquark couplings to the third family, in particular to the top quark, can only be tested in experiments at Tevatron and LHC. In order to be consistent with the above discussed experimental constraints there are widespread arguments [15] that new flavor-changing interactions are likely to affect mainly the physics of the third family of quarks only.

At Tevatron diquarks may give rise to t-channel processes $u\bar{u} \to \bar{u}_i u_j$, $u\bar{d} \to \bar{d}_i u_j$ and $d\bar{d} \to \bar{d}_i d_j$ due to the proton and anti-proton valence quarks collisions; and s-channel resonance processes $uu \to u_i u_j$, $ud \to d_i u_j$ and $dd \to d_i d_j$ due to the proton valence quarks and anti-proton sea quarks collisions. Among these the most interesting final states are $t\bar{t}$, tt, $t\bar{b}$, tb, $b\bar{b}$, bb and $t\bar{c}$, tc because top, bottom and charm tagging allows one to distinguish them effectively from the standard model background. In particular, the observation of a large rate for the flavor-changing process $uu \to t\bar{c}$ at the Tevatron Runs II and III would indicate unambiguously the existence of new physics [16]. Also, large excess of the same-sign dileptons from the tt final states may indicate for the new physics.

Because collider phenomenology of the diquarks depends on their quantum numbers we have to discuss their couplings in more detail. The form of the Yukawa Lagrangian describing the $SU(2)_L$ singlet diquark X coupling to down type quarks is given by

$$\mathcal{L} = f_{ij} \, [T_{ab}]^m \, \overline{(d_R^c)_i^a} \, d_{Rj}^b \, X_m + h.c. \,, \tag{19}$$

where $\overline{d^c} = -d^T C^T$, $i, j = 1, 2, 3$ are family indices, $a, b = 1, 2, 3$ are colour indices, R denotes the chirality of the quarks and matrices T form the basis of the n dimensional 3×3 matrix representation of the $SU(3)_c$ group. Therefore $m = 1, ..., n$. For colour triplets ($n = 3$) the matrices T_{ab} should be anti-symmetric in a and b and one may identify $[T_{ab}]^m = \varepsilon_{ab}{}^m$, where $\varepsilon_{ab}{}^m$ is the three dimensional totally anti-symmetric tensor. For colour sextets ($n = 6$) the matrices T_{ab} are symmetric in a and b. Analogously to Eq. (19) we have for the $SU(2)_L$ singlet diquark couplings to up and down type quarks

$$\mathcal{L} = f_{ij} \, [T_{ab}]^m \, \overline{(u_R^c)_i^a} \, d_{Rj}^b \, X_m + h.c. \,, \tag{20}$$

and for the couplings to up type quarks

$$\mathcal{L} = f_{ij} \, [T_{ab}]^m \, \overline{(u_R^c)_i^a} \, u_{Rj}^b \, X_m + h.c. \,. \tag{21}$$

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Yukawa interactions of the $SU(2)_L$ triplet diquarks to quarks can be expressed by the Lagrangian

$$\mathcal{L} = f_{ij} \, [T_{ab}]^m \, (\tau_\kappa \varepsilon)^{\alpha\beta} \, \overline{(q_L^c)_{i\alpha}^a} \, q_{Lj\beta}^b \, X_m^\kappa + h.c. \,, \tag{22}$$

where the new indices $\alpha, \beta = 1, 2$ are the $SU(2)_L$ indices, τ^κ are the three Pauli matrices and $\varepsilon_{\alpha\beta}$ is the 2×2 anti-symmetric tensor acting in the $SU(2)_L$ space. In the Lagrangians above, for the $SU(2)_L$ singlet diquarks the coupling matrices f are anti-symmetric in the generation space while for the $SU(2)_L$ triplets f is symmetric. It follows that both singlets and triplets can mediate the t-channel processes at hadron colliders while the s-channel processes can be induced only by the $SU(2)_L$ triplets.

We have calculated the cross sections of the above mentioned processes mediated by the diquarks at Tevatron by convoluting over the default parton structure functions MRS (G) [17] of the CERN Library package PDFLIB [18] version 7.09. The colour factors F_c coming from averaging over the initial state colours and summing over the final state colours depend on the $SU(3)_c$ representation. For the s- and t-channel cross sections we obtain $F_c = 4/3$ and $F_c = 2/3$, and for the decay widths $F_c = 2$ and $F_c = 1$ corresponding to triplets and sextets, respectively. For definiteness we shall in the following consider the colour sextet diquarks; the corresponding cross sections for the colour triplets are *larger* due to the larger colour factors.

Let us first study the s-channel resonance production of dijets at Tevatron. Neglecting the final state fermion masses the diquark partial width is given by

$$\Gamma = \frac{f^2}{8\pi} F_c M_X \,, \tag{23}$$

which implies that the expected resonance is quite narrow. Assuming that the charge 4/3 and -2/3 components of (6,3,1/3) diquark decay 100% to uu and bb final states, respectively, we plot in Fig. 3 the dijet s-channel cross sections as functions of the diquark masses for

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two values of collision energy $\sqrt{s} = 1.8$ TeV and $\sqrt{s} = 2$ TeV. The diquark couplings are taken to be equal to unity. One can see from the behaviour of the cross section curves in Fig. 3 that the narrow width approximation is effective up to diquark masses ~ 700 GeV. For considerably higher masses the width becomes larger according to Eq. (23) and the resonance peak will be smeared. The difference in $p\bar{p} \to uu$ and $p\bar{p} \to bb$ cross sections is only due to the difference in the parton structure functions for the u and d quarks.

The narrow resonances in dijet signal [19], and in the double b-tagged dijet signal separately [20], have been searched for with Tevatron CDF detector. Negative results have resulted in model-independent bounds on the narrow resonance cross sections. Therefore, for the diquark couplings equal to unity the present CDF data allowes us to exclude the charge 4/3 component of the (6,3,1/3) diquark in the mass range

$$270 \text{ GeV} \lesssim M_X \lesssim 500 \text{ GeV} \tag{24}$$

and the charge -2/3 component in the mass range

$$200 \text{ GeV} \lesssim M_X \lesssim 570 \text{ GeV}, \tag{25}$$

provided they decay 100% to uu and bb pairs, respectively.

Next we consider the top quark pair production at Tevatron. In Fig. 4 we plot the cross sections of (6,1,4/3) or (6,3,1/3) diquark mediated t-channel process $p\bar{p} \to t\bar{t}$ and the (6,3,1/3) diquark mediated s-channel resonance process $p\bar{p} \to tt$ as functions of the diquark masses for two values of collision energy, as indicated in the figure. The couplings f are taken to be equal to unity. For the resonance production we have assumed diquark branching ratio to tt pair to be 100%. With these assumptions the resonance cross section exceeds the t-channel cross section for diquark masses below ~ 700 GeV. Above that scale t-channel production becomes dominant.

To derive constraints on the diquark masses and couplings we proceed as follows. The standard-model NLO (next to leading order) $t\bar{t}$ [21] and also $t\bar{b}$ [22] cross sections at the Tevatron and LHC are known within a total error of 15%. If the total cross sections of the new processes exceed 15% of the standard-model cross section, then the new signal will be detectable (such a criterion has recently been used in [23]). Assuming that one of the cross sections dominates we find from the Tevatron Run I ($\sqrt{s} = 1.8$ TeV, L=0.1 fb^{-1}) tt and $t\bar{t}$ cross sections that the masses of (6,3,1/3) and (6,1,4/3) diquarks should exceed 700 and 600 GeV, respectively. At Run II ($\sqrt{s} = 2$ TeV, L=2 fb^{-1}) the corresponding bounds would be 750 and 600 GeV, respectively.

These bounds based on the cross section estimates indicate the sensitivity of Tevatron to considered processes. Of course, dedicated Monte Carlo studies with appropriate kinematical cuts would allow one to achieve a much better signal-over-background ratio and therefore higher bounds on the diquark masses (in particular at Run III, $\sqrt{s} = 2$ TeV, L=30 fb^{-1}). This is because the diquarks are scalars while the standard model top pair production background is dominantly produced in gluon-gluon and gluon-quark collisions. This statement applies especially to the s-channel tt production. This process is quark flavour violating and the distribution of the final state top quarks is flat. The unambiguous diquark signal in this process would be the same-sign dilepton originating from top decays. This has very little background from the standard model. However, while large excess of the same-sign dileptons would indicate a clear signal of new physics then reconstruction of diquark resonance peak in this channel is impossible due to the missing energy carried away by neutrinos. On the other hand, studying the kinematics of the visible particles in terms of endpoint spectra it should be possible to extract information on the diquark mass also in this channel (the Jacobian peak should be visible). This information can further be combined with the studies of hadronic and semileptonic channels. These kind of dedicated studies are beyond the scope

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of the present work.

To demonstrate the effectiveness of dedicated studies we first plot in Fig. 5 the cross section of the process $p\bar{p} \to t\bar{c}$ and $p\bar{p} \to tc$ against the mass of the (6,3,1/3) diquark M_X for $\sqrt{s} = 1.8$ TeV and $\sqrt{s} = 2$ TeV assuming the couplings to be of order unity. The single-top production at Tevatron due to anomalous chromomagnetic couplings has been studied in [24, 25]. Assuming that in the case of the diquark mediated process the same sensitivity to the $p\bar{p} \to t\bar{c}$ signal can be achieved as in Ref.[25], then comparison of the $p\bar{p} \to t\bar{c}$ cross sections gives for Run I a bound $M_X \gtrsim 750$ GeV and for Run II and III $M_X \gtrsim 1.2$ TeV and $M_X \gtrsim 1.6$ TeV, respectively. These estimates are indicative of the sensitivity of the Tevatron to flavor-changing diquark processes.

For completeness, we obtain from the $t\bar{b}$, tb cross section estimates at Tevatron that at Run I that the (6,3,1/3) diquark mass should exceed 650 GeV and at Run II 720 GeV. Appropriate kinematical cuts and tagging of hard b jets should separate the signal from the background.

Because LHC will be a pp collider, it will be the ideal place to search for scalar diquarks. The most interesting process, $pp \to tt$, can be mediated by the s-channel (6,3,1/3) diquark. Therefore, resonance production is possible. The discussion on detectability of this process at Tevatron applies also here. Another interesting process to study is $pp \to bb$ which may also occur due to the s-channel (6,3,1/3) diquark resonance. The resonance can be detected searching for doubly b-tagged dijet final states. We plot the cross sections of these processes in Fig. 6 as functions of M_X. The couplings are taken equal to unity and only one decay channel is assumed for the diquarks. Considering tt production, the excess of same-sign charged dilepton final states will be clear signal of the new physics. The main experimental error to this signal will come from the misidentification of the lepton charge but this is expected to be small. To estimate which diquark mass scales can be probed at LHC we

assume that about 10^3 events will constitute a discovery. It follows from Fig. 6 that for couplings of order unity, scalar diquarks as heavy as 7 TeV can be discovered already at the low luminosity (L=10 fb^{-1}) run of LHC. With the final luminosity, L=500 fb^{-1}, diquark masses of 13 TeV can be probed. The same conclusions apply also to the process $pp \to tc$ provided that the involved couplings are of order unity.

4. Conclusions

We see from the above discussion that all exotic scalars which are bilinear combinations of two known fermion representations are all very heavy if their couplings are independent of the quark or lepton family and unsuppressed. At closer scrutiny, we see also that the above derived bounds on the couplings and masses of the exotic scalars depend on the flavor indices and on the assumption of the existence of trilinear couplings. If the latter couplings are absent, the most stringent bounds coming from the nonobservation of proton decay and neutron-antineutron oscillations, and from baryogenesis can be evaded in many cases. Certain combinations of the couplings can be tested only in high-energy experiments. The possibility exists for some of the diquark masses to be in the range of a few TeV, in which case they can be discovered at future hadron colliders. At Tevatron the best sensitivity to diquarks can be achieved by studying the processes $\bar{p}p \to \bar{t}c, \bar{t}t$. The resonance process $\bar{p}p \to tt$, though sea quark suppressed, may also provide an observable excess of the same-sign dilepton final states indicating clearly for the new physics. LHC will be an ideal facility for the studies of diquarks because the quark-flavor-violating resonance processes $pp \to tt, bb$ will not be suppressed. Therefore diquark masses as high as 13 TeV can be probed at LHC.

ACKNOWLEDGEMENT

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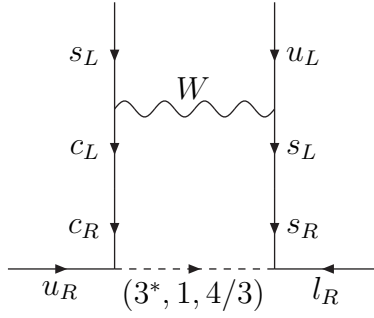
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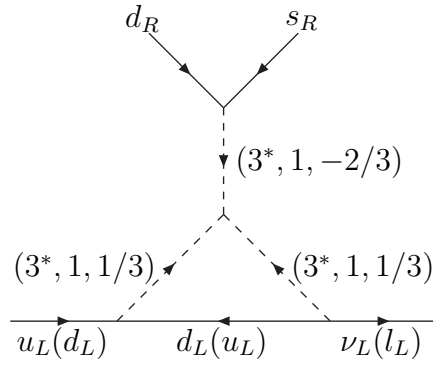**Fig. 1.** One-loop proton decay due to the $(3^*, 1, 4/3)$ scalar.

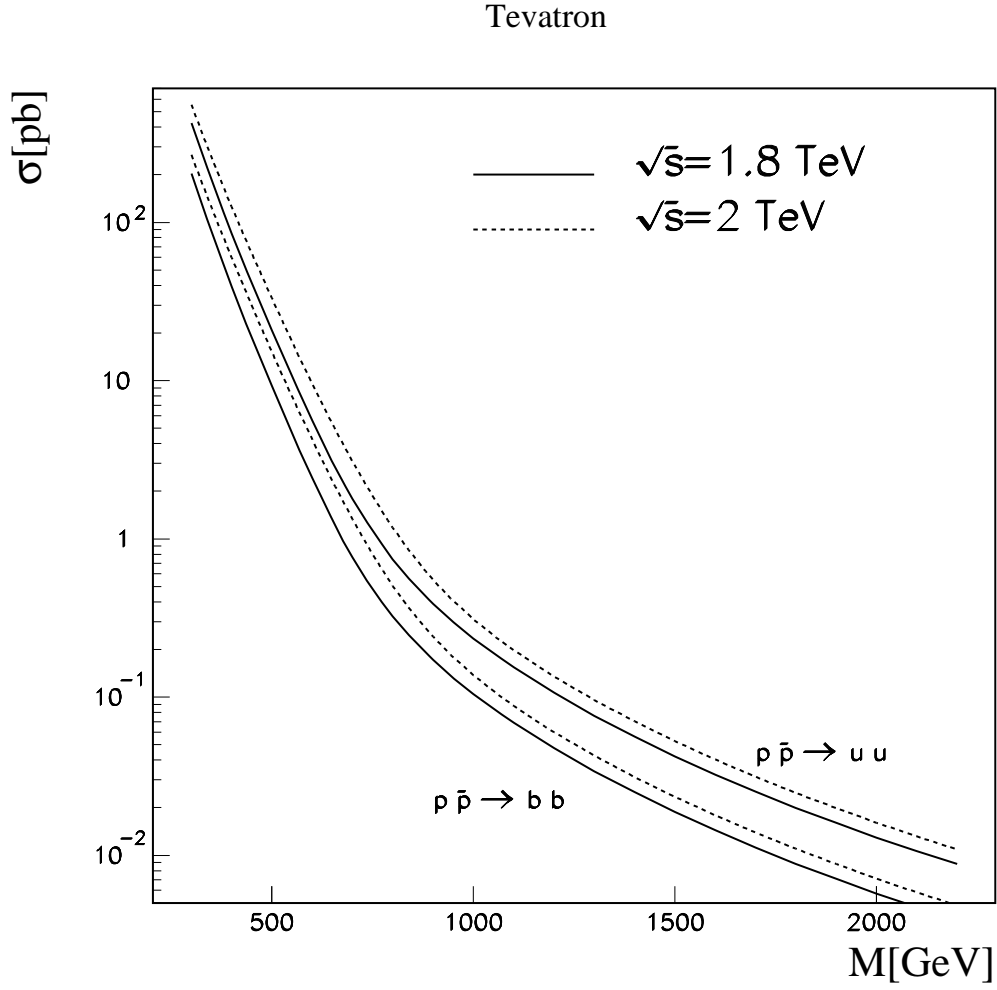**Fig. 2.** One-loop proton decay due to the $(3^*, 1, -2/3)$ scalar.

Tevatron

Fig. 3. Cross sections of the s-channel processes $p\bar{p} \to uu$ and $p\bar{p} \to bb$ at Tevatron as functions of the scalar $(6,3,1/3)$ diquark masses. The couplings are taken to be equal to unity.

Tevatron
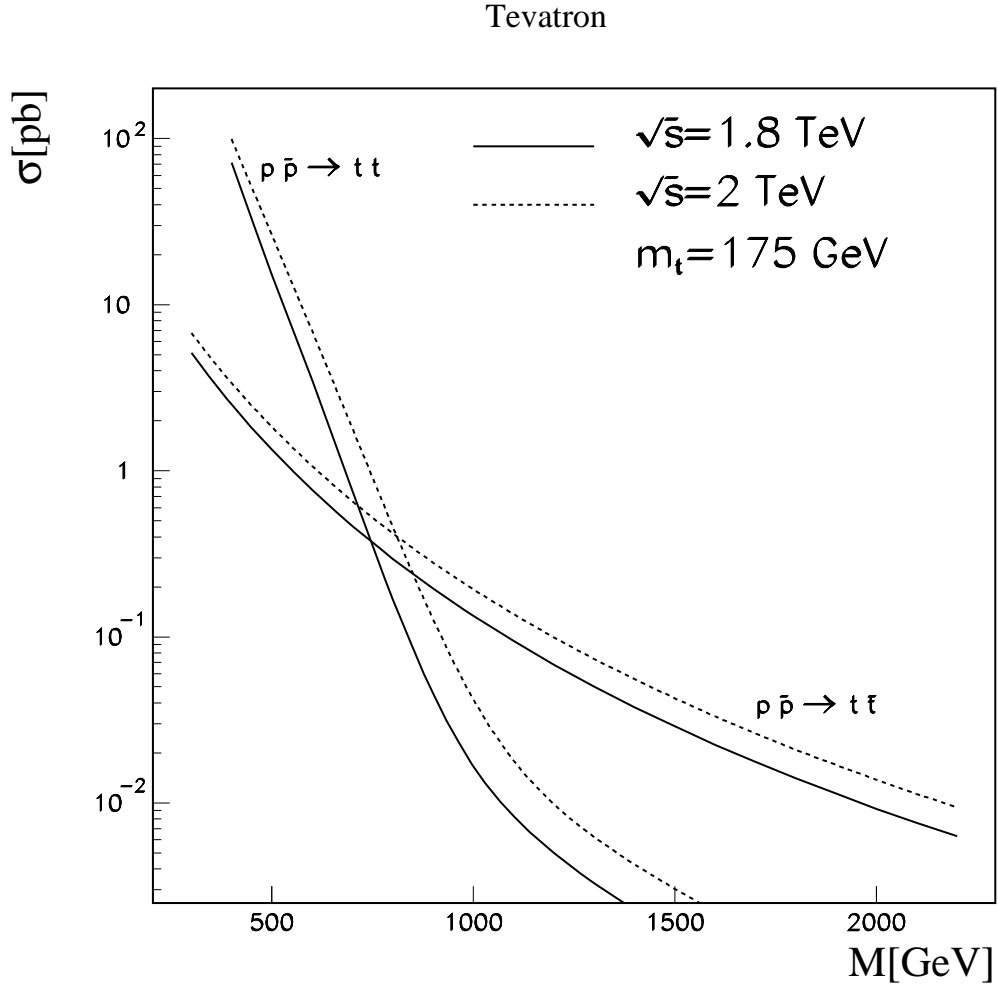

Fig. 4. Cross sections of the processes $p\bar{p} \to t\bar{t}$ and $p\bar{p} \to tt$ at Tevatron against the (6,3,1/3) scalar diquark mass M_X. The couplings are taken to be equal to unity.
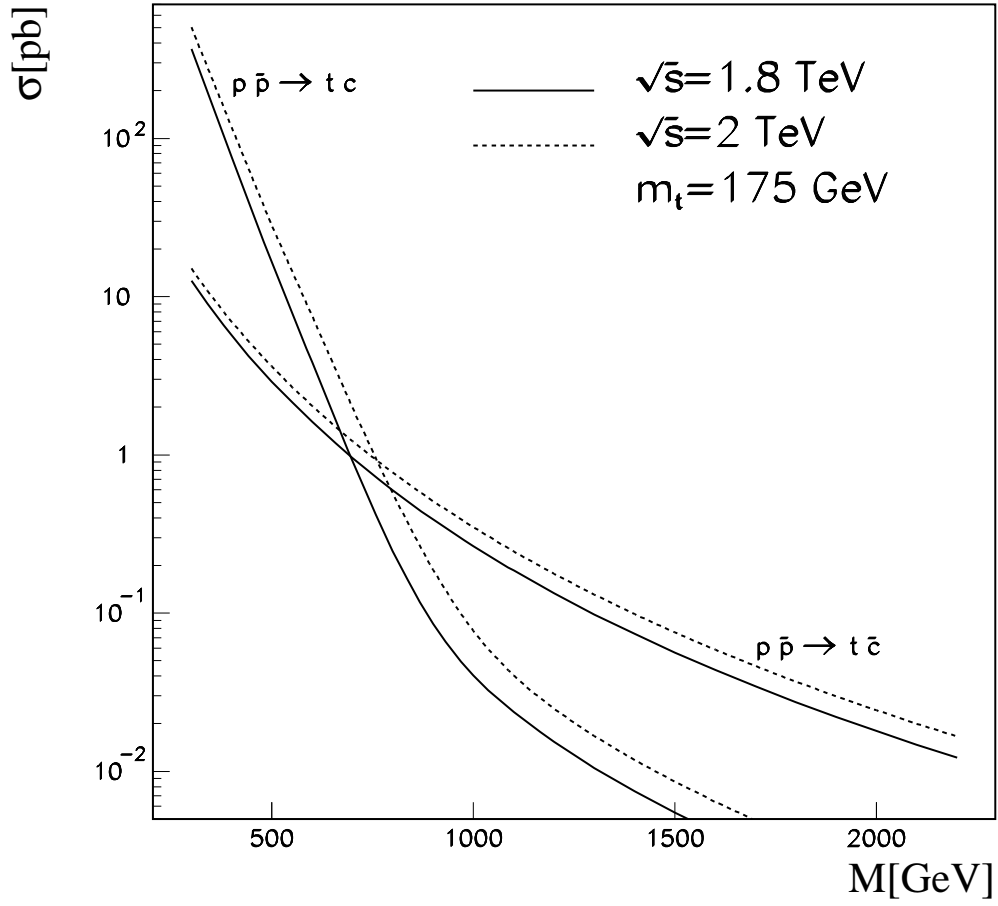
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Fig. 5. Cross sections of the FCNC processes $p\bar{p} \to t\bar{c}$ and $p\bar{p} \to tc$ at Tevatron against the (6,3,1/3) scalar diquark mass M_X. The couplings are taken to be equal to unity.
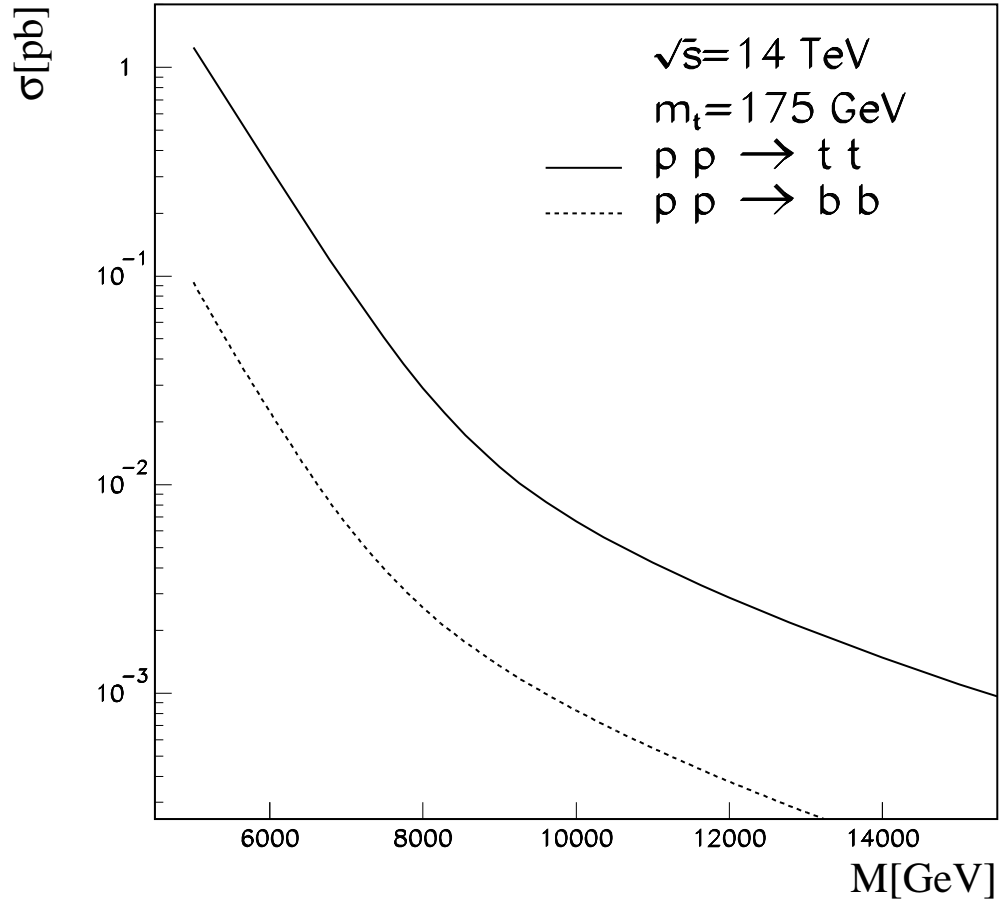
LHC

Fig. 6. Cross sections of the *s*-channel process $pp \to tt$ and $pp \to bb$ at LHC against the scalar (6,3,1/3) diquark masses. The couplings are taken to be equal to unity.