TeV Scale Leptoquarks as a Signature of Standard–like Superstring Models

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ABSTRACT

We show that there can be TeV scale scalar and fermionic leptoquarks with very weak Yukawa couplings in a generic standard–like superstring model. Leptoquark–(down–like) quark mixing though present, is not large enough to violate the unitarity bounds on the CKM matrix. The constraints on leptoquark masses and couplings from FCNCs are easily satisfied whereas those from baryon number violation may cause problems. The leptoquarks of the model are compared to the ones in $E_6$ Calabi–Yau and flipped $SU(5) \times U(1)$ models.

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1. Introduction

Superstring theories are [1], to date, the most promising Planck scale theories of particle physics. In spite of their successes, one of the many drawbacks of realistic superstring models is their loose connection to TeV or weak scale physics. TeV (or weak) scale signs or predictions of realistic superstring models are rare even though they can reproduce most of the known low–energy physics. It is important to look for these signs or predictions either to make specific superstring models more plausible or to rule them out.

In this letter, we show that, under certain conditions, there can be TeV scale leptoquarks in a class of standard–like superstring models [3,4]. We find that these leptoquarks have very weak (i.e. $< 10^{-3}$) Yukawa couplings if at all. Supersymmetry (SUSY) constraints in the observable and hidden sectors play an important role in these results. One of the leptoquarks mixes with down–like quarks, but the mixing is small enough to satisfy the unitarity constraints on the CKM matrix. Due to the small leptoquark Yukawa couplings, constraints from flavor changing neutral currents (FCNCs) on leptoquark masses are easily satisfied. Baryon number violation may impose severe constraints on leptoquark masses unless the Yukawa couplings to diquarks are absent up to very high orders. We also compare these leptoquarks with those that arise from $E_6$ Calabi–Yau [1] and flipped $SU(5) \times U(1)$ [2] models and discuss their differences.

The standard–like superstring models that we consider have the following properties [3,4]:

1. $N = 1$ space–time SUSY.

2. A $SU(3)_C \times SU(2)_L \times U(1)^n \times$ hidden gauge group.

3. Three generations of chiral fermions and their superpartners, with the correct quantum numbers under $SU(3)_C \times SU(2)_L \times U(1)_Y$.

4. Higgs doublets that can produce realistic electro–weak symmetry breaking.
5. Anomaly cancellation, apart from a single “anomalous” U(1) which is canceled by application of the Dine–Seiberg–Witten (DSW) mechanism [5].

The superstring standard–like models are constructed in the four dimensional free fermionic formulation [6]. The models are generated by a basis of eight boundary condition vectors for all world–sheet fermions [3,4]. The observable and hidden gauge groups after application of the generalized GSO projections are $SU(3)_C \times U(1)_C \times SU(2)_L \times U(1)_L \times U(1)_R^6$ and $SU(5)_H \times SU(3)_H \times U(1)^2$, respectively. The weak hypercharge is given by $U(1)_Y = \frac{1}{3}U(1)_C + \frac{1}{2}U(1)_L$ and has the standard $SO(10)$ embedding. The orthogonal combination is given by $U(1)_{Z'} = U(1)_C - U(1)_L$. The models have six right–handed and six left–handed horizontal symmetries $U(1)_{r_j} \times U(1)_{\ell_j}$ ($j = 1, \ldots, 6$), which correspond to the right–moving and left–moving world–sheet currents respectively.

A generic standard–like superstring model including the complete massless spectrum with the quantum numbers and the cubic superpotential were presented in Ref. [3] and will not be repeated here. The notation of Ref. [3] is used throughout this letter.

2. SUSY constraints

In order to preserve SUSY at $M_P$, one has to satisfy a set of F and D constraints. The set of F and D constraints is given by the following equations:

$$D_A = \sum_k Q_k^A |\chi_k|^2 = \frac{-g^2 e^{\phi_D}}{192 \pi^2} Tr(Q_A) \frac{1}{2\alpha'}$$  \hspace{1cm} (1a)

$$D^j = \sum_k Q_k^j |\chi_k|^2 = 0 \quad j = 1 \ldots 5$$  \hspace{1cm} (1b)

$$D^j = \sum_k Q_k^j |\chi_k|^2 = 0 \quad j = C, L, 7, 8$$  \hspace{1cm} (1c)

$$W = \frac{\partial W}{\partial \eta_i} = 0$$  \hspace{1cm} (1d)

\* $U(1)_C = \frac{3}{2}U(1)_{B-L}$ and $U(1)_L = 2U(1)_{T_3_R}$. 

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where $\chi_k$ and $\eta_i$ are the fields that do and do not get VEVs respectively and $Q^I_k$ are their charges. $(2\alpha')^{-1} = g^2 M_P^2 / 8\pi = M^2 \sim (10^{18} \text{ GeV})^2$ and $W$ is the superpotential. From Eq. (1a) we see that, $SO(10)$ singlet scalars must get VEVs $\sim g^2 M/4\pi \sim M/25$ in order to preserve SUSY at $M_P$.

The set of F constraints in the observable sector has been studied before [7]. One finds that SUSY requires $\langle \Phi_{12} \rangle = \langle \bar{\Phi}_{12} \rangle = \langle \xi_3 \rangle = 0$ even though the number of fields is larger than the number of constraints. Then, one is left with only three F constraints from the observable sector [7]. F constraints in the hidden sector which are derived from the cubic superpotential have also been investigated recently [8]. These lead to conditions on hidden sector VEVs which are particularly strong if one also requires realistic quark and lepton masses. Then, SUSY in the hidden sector (at $M_P$) imposes $\langle H_i \rangle = 0$ where $i = 13, \ldots, 26$ in the notation of Ref. 3 (with at most one pair among these having non–zero VEVs in special cases) [8].

Once SUSY is dynamically broken by the hidden sector condensates, the VEVs which vanish above can become non–zero. For broken SUSY, $\langle F \rangle \sim M^2_{\text{SUSY}}$ and $m_{3/2} \sim M^2_{\text{SUSY}} / M < O(\text{TeV})$, in order to solve the hierarchy problem. For a light gravitino (and light squark and slepton masses), i.e. $m_{3/2} \sim O(100 \text{ GeV})$, we need $M_{\text{SUSY}} \sim 10^{10} \text{ GeV}$ or $\langle F \rangle \sim 10^{20} \text{ GeV}$. As a result, the VEVs which vanished due to SUSY can now be non–zero and up to $O(\text{TeV})$. Note that for a heavy gravitino with $m_{3/2} \sim O(\text{TeV})$ these VEVs can be up to $O(10 \text{ TeV})$.

### 3. Leptoquarks of the model

In the massless $b_1 + b_2 + \alpha + \beta + (S)$ sector of standard–like superstring models there are two color triplet, electroweak singlet states, $D_{45}$ and $\bar{D}_{45}$ [3,4]. Under $SU(3)_C \times SU(2)_L \times U(1)_C \times U(1)_L$, $D_{45}$ and $\bar{D}_{45}$ transform as $\begin{pmatrix} 3, 1, -1, 0 \end{pmatrix}$ and $\begin{pmatrix} \bar{3}, 1, 1, 0 \end{pmatrix}$ respectively. Since $Q_Y = Q_C/3 + Q_L/2$ and $Q_Z' = Q_C - Q_L$ we find that $Q_Y(D_{45}) = Q_{EM}(D_{45}) = -1/3$ and $Q_Z'(D_{45}) = -1$ with $\bar{D}_{45}$ having opposite charges. Another combination of $Q_C$ and $Q_L$ gives $Q_{B-L} = 2Q_C/3$ which is a gauge symmetry in these models. Thus, $Q_{B-L}(D_{45}) = -2/3$ and $Q_{B-L}(\bar{D}_{45}) = 2/3$. 

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From all these quantum numbers we see that $D_{45}$ and $\bar{D}_{45}$ are actually leptoquarks. Note that $D_{45}$ and $\bar{D}_{45}$ are superfields and as a result there are two scalar and two fermionic leptoquarks in this model.

In general one expects that $D_{45}$ and $\bar{D}_{45}$ get large masses (of $O(10^{17} \text{ GeV})$) at the level of the cubic superpotential. Even if this is not the case $D_{45}$ and $\bar{D}_{45}$ can get large masses from higher order terms (i.e. $N > 3$ terms) in the superpotential and decouple from the low–energy spectrum. In fact, in the standard–like model under consideration, there are potential mass terms for $D_{45}$ and $\bar{D}_{45}$ at the cubic level [3,4]

$$W_{D,\bar{D}} = \frac{1}{2}D_{45}\bar{D}_{45}\xi_{3} + \frac{1}{2}D_{45}H_{18}H_{21}$$  (2)

where $H_{21}$ is a hidden sector state which is a $\bar{3}$ of color, $\xi_{3}$ and $H_{18}$ are singlets of $SU(3)_{C} \times SU(2)_{L} \times U(1)_{C} \times U(1)_{L}$. We see that due to the SUSY constraints in the observable and hidden sectors, i.e. since $\langle \xi_{3} \rangle = \langle H_{18} \rangle = 0$, $D_{45}$ and $\bar{D}_{45}$ remain massless at the cubic level of the superpotential.

As noted earlier there may be higher order terms which give large masses to $D_{45}$ and $\bar{D}_{45}$. Higher order ($N > 3$) non–renormalizable contributions to the superpotential are obtained by calculating correlators between vertex operators [9] $A_{N} \sim \langle V_{f}^{i}V_{f}^{j}V_{b}^{k} \cdots V_{b}^{N} \rangle$ where $V_{f}^{i}$ ($V_{b}^{i}$) are the fermionic (bosonic) vertex operators corresponding to different fields. The non–vanishing terms are obtained by applying the rules of Ref. [9]. First, since only $H_{23}, H_{25}$ or $H_{24}, H_{26}$ can get VEVs due to the SUSY constraints in the hidden sector, a mass term containing $H_{21}$ is not possible to any order in $N$. Second, $H_{21}$ gets a large mass from the term $\frac{1}{2}H_{21}H_{22}\xi_{1}$ in the cubic superpotential and decouples from the low–energy spectrum [3,4]. Therefore $D_{45}$ cannot mix with $H_{21}$ at low or intermediate energies. On the other hand, there are $D_{45}\bar{D}_{45}$ mass terms which arise from $N > 3$ terms in the superpotential. At $N = 5$ we find a large number of terms which can be combined to give

$$D_{45}\bar{D}_{45}(\xi_{2} + \xi_{3})\frac{\partial W}{\partial \xi_{3}} + \Phi_{12}\frac{\partial W}{\partial \Phi_{12}} + \bar{\Phi}_{12}\frac{\partial W}{\partial \bar{\Phi}_{12}}$$  (3)
These vanish because of the SUSY constraints, Eq. (2) and Eq. (3), which can be written as

$$\frac{\partial W}{\partial \xi_3} = \frac{\partial W}{\partial \Phi_{12}} = \frac{\partial W}{\partial \bar{\Phi}_{12}} = 0$$

(4)

There are many other higher order ($N > 5$) $D_{45} \bar{D}_{45}$ terms arising from the observable states which are proportional to the F terms in Eq. (8) and therefore vanish. There may be terms which are not proportional to the F terms in Eq. (8) at very high orders. It is difficult to disregard this possibility because the number of terms increases rapidly with the order $N$. Here we assume that if there are such terms they can be made to vanish by an appropriate choice of vanishing VEVs.

When hidden sector states are taken into account, there are $N = 6$ terms

$$D_{45} \bar{D}_{45} T_2 \bar{T}_2 \Phi_{45} \Phi_2^+ (\xi_1 + \xi_3)$$

(5)

which give large masses to $D_{45}$ and $\bar{D}_{45}$ if $\Phi_2^+$ gets a VEV. Here $T_2, \bar{T}_2$ are 5, 5 of the hidden $SU(5)_H$ gauge group. There are no phenomenological constraints from quark and lepton masses or quark mixing on $\langle \Phi_2^+ \rangle$. In addition, the SUSY F and D constraints can be satisfied whether $\langle \Phi_2^+ \rangle$ vanishes or not. If $\langle \Phi_2^+ \rangle \neq 0$, then generically $\langle \Phi_2^+ \rangle \sim M/10 \sim 10^{17}$ GeV and the $\langle T_2 \bar{T}_2 \rangle \sim \Lambda_H^2$ where $\Lambda_H \sim 10^{14}$ GeV is the hidden $SU(5)_H$ condensation scale [10]. This gives $M_{D, \bar{D}} \sim 10^8$ GeV. Potential leptoquark mass terms arising from VEVs of $H_1$ vanish due to the SUSY constraints in the hidden sector, Eqs. (4).

If, on the other hand, $\langle \Phi_2^+ \rangle = 0$, then the $D_{45} \bar{D}_{45}$ mass terms come from the SUSY breaking VEVs. (Once again we assume that if there are $N > 6$ terms similar to Eq. (9), then they can be made to vanish by an appropriate choice of vanishing VEVs.) As mentioned earlier, the VEVs vanishing due to SUSY can become non–zero (and up to the $TeV$ scale) once SUSY is broken. Therefore, when SUSY is broken, $D_{45}$ and $\bar{D}_{45}$ get $TeV$ scale masses from the cubic superpotential, i.e. from the terms in Eq. (6) since now $\langle \xi_3 \rangle \sim O(TeV)$ (for a scenario with light squark and lepton masses). In addition the scalar leptoquarks get contributions
to their masses from soft SUSY breaking terms. These are generally less than a TeV. (For the case we consider, soft SUSY breaking masses for the scalars \(m_0 \sim m_{3/2} \sim O(100 \text{ GeV})\).) Thus, under the conditions given above, there are two scalar and two fermionic leptoquarks with masses around the TeV scale in this model. The lower bound on \(M_{D, \bar{D}}\) from direct leptoquark searches is 45 GeV [12] which is easily satisfied.

The fermionic leptoquarks may mix with down–like quarks. In this model, there is a \(d_3 D_{45}\) mixing term of the form

\[
d_3 D_{45} N_3 \Phi_{13} \Phi_3^+ (\xi_1 + \xi_2 + \xi_3)
\]

Similar mixing terms for the other two down–like quarks may appear at higher orders, \(N > 6\). There are no \(d_i \bar{D}_{45}\) mixing terms, for the left–handed down–like quarks, due to the conservation of \(Q_L\) and \(Q_C\). \(\langle N_{fi}^c \rangle\) appears in non–renormalizable terms which induce dimension four baryon number violating operators [7]. From the proton lifetime, we get the constraint \(\langle N_{fi}^c \rangle \sim O(TeV)\) at most. As a result, the mixing term in Eq. (10) is at most about \(O(\text{GeV})\) and the others are smaller by at least an order of magnitude since they appear at higher orders. (Note that the mixing term in Eq. (10) can be made to vanish by taking \(\langle \Phi_{13} \rangle = 0\) or \(\langle \Phi_3^+ \rangle = 0\).)

Now, the \(3 \times 3\) CKM matrix becomes non–unitary because of the new mixing terms in the \(4 \times 4\) down quark mass matrix. (In standard–like superstring models, CKM matrix arises mainly from the down quark mass matrix [10].) The strongest bounds on the magnitude of the \(d_3 D_{45}\) mixing arise from unitarity of the CKM matrix \((V_{ij})\) which imposes \(|V_{ud}| < 0.07\) [12] and \(|Re V_{id} V_{is}| < 2.4 \times 10^{-5}, \ i = u, c, t\) (from flavor changing Z currents [13]). In our case, when \(d_3 D_{45}\) mixing is much smaller than \(M_D, |V_{ud}| \sim \langle N_{fi}^c \rangle / 10^3 \langle \xi_3 \rangle \sim 10^{-3}\) and \(|V_{cD}| \sim |V_{tD}|\) are (at least) an order of magnitude smaller. With these results the constraints from unitarity are easily satisfied. Conversely, since the \(d_3 D_{45}\) mixing is very small compared to \(M_D\), there will not be an appreciable violation of unitarity in the \(3 \times 3\) CKM matrix.

4. Leptoquark interactions
The leptoquarks $D_{45}$ and $\bar{D}_{45}$ carry color, electric and $Z'$ charge and therefore have strong, electromagnetic and $Z'$ gauge interactions. Of these, the $Z'$ interactions will be very weak at the TeV scale if the $Z'$ gauge boson has a large mass (i.e. $M_{Z'} > \sim TeV$). Otherwise all gauge interactions of $D_{45}$ and $\bar{D}_{45}$ are appreciable at the TeV scale. In any case, the leptoquarks will be easy to produce in $e^+e^-$ or $pp$ collisions. Production of leptoquarks has been investigated in Ref. [14] in detail. As final states $D_{45}$ and $\bar{D}_{45}$ will look like new, very massive, $SU(2)_L$ singlet, down–like quarks.

Yukawa couplings of $D_{45}$ and $\bar{D}_{45}$ are more interesting since it is these that allow $D_{45}$ and $\bar{D}_{45}$ to decay into quarks and leptons. In addition, Yukawa couplings are model dependent and therefore useful to distinguish between different models. The Yukawa couplings allowed by $Q_C$ and $Q_L$ conservation are

$$W_Y = L_i Q_i D_{45} \Phi + e_i u_i D_{45} \Phi + N_i d_i D_{45} \Phi$$

(7)

where $i$ is the generation index ($i = 3$ is the lightest generation in the notation of Ref. [3].) and $\Phi$ is a generic string of $SO(10)$ singlet fields which get VEVs. Effective Yukawa couplings for the three terms ($g_{1i}, g_{2i}, g_{3i}$) are obtained from the VEVs of the string of singlets divided by the proper power of $M$. Each term in (11) also has a coefficient which can be calculated exactly and is $O(1)$ [9].

We look for terms which induce effective Yukawa couplings for $D_{45}$ and $\bar{D}_{45}$ at orders $N > 3$. We find that all kinds of couplings given in Eq. (11) are allowed for all generations by the gauge symmetries of the model at $N = 4$ and $N = 5$. All of these terms except one vanish due to the string selection rules as given in Ref. [9]. These selection rules arise from the left–handed $U(1)$ symmetries and the world–sheet sigma model operators that appear in the vertex operators in the world–sheet correlators (after picture changing is taken into account). The term that remains at $N = 5$ is

$$W_Y = L_3 Q_3 D_{45} \Phi_{45} \xi_3$$

(8)

which potentially gives an effective Yukawa coupling of $\langle \Phi_{45} \xi_3 \rangle / M \sim 10^{-2}$. But
\[ \langle \xi_3 \rangle = 0 \text{ due to SUSY, so this term vanishes. (Even after SUSY breaking, } \langle \xi_3 \rangle / M \sim 10^{-15} \text{ and this term is negligible.) We see that the } D_{45} \text{ and } \bar{D}_{45} \text{ Yukawa couplings can only arise from terms at } N > 5 \text{ in this model. As a result they are at most } \sim 10^{-3} \text{ and probably smaller which means that } D_{45} \text{ and } \bar{D}_{45} \text{ have very weak decays into leptons and quarks.}

\]

If the Yukawa couplings of \( D_{45} \) or \( \bar{D}_{45} \) are not “diagonal”, i.e. \( D_{45} \) or \( \bar{D}_{45} \) couple to more than one quark and lepton generation, they induce FCNCs. FCNC processes such as \( K_L \rightarrow e^+e^- \) and \( K^+ \rightarrow \pi^+\nu\bar{\nu} \) give the strongest bounds on \( M_{D,\bar{D}} \) and the Yukawa couplings \( g_{1,2,3}i \). From the analysis of Ref. [15] we get

\[ |g_jg_3| < 5.65 \times 10^{-8} \sin \theta_c M_{D,\bar{D}}^2 \]  

from \( K_L \rightarrow e^+e^- \), and a slightly lower bound from \( K^+ \rightarrow \pi^+\nu\bar{\nu} \). Here \( j \) is the index for the different couplings and \( \theta_c \) is the Cabibbo angle with \( \sin \theta_c \sim 0.2 \).

With the upper bound of \( \sim 10^{-3} \) that we obtained for the \( D_{45} \) and \( \bar{D}_{45} \) Yukawa couplings above, we find that the lower bound on \( M_{D,\bar{D}} \) from FCNCs is \( M_{D,\bar{D}} > 10 \text{ GeV} \) which is not a constraint at all since the bound from direct searches is \( M_{D,\bar{D}} > 45 \text{ GeV} \). FCNC constraints are severe only for leptoquarks with Yukawa couplings \( \sim O(1) \) to more than one generation. The same bounds and remarks also apply to the squark–quark–fermionic leptoquark Yukawa couplings but since squark masses are very large (and unknown) no useful bounds exist in this case.

\( D_{45} \) and \( \bar{D}_{45} \) can also have diquark couplings such as

\[ W' = u_i^c d_i^c \bar{D}_{45} \Phi + Q_i Q_i D_{45} \Phi \]  

(We will call these effective couplings \( g_{4i} \) and \( g_{5i} \).) Baryon number (B) violation imposes severe constraints on the strength of these couplings if there are non–zero leptoquark couplings like those in Eq. (11). This is because if both leptoquark and diquark couplings are non–zero at the same time, \( D_{45} \) and \( \bar{D}_{45} \) exchange leads to very large B violating processes such as \( Q_i Q_i \rightarrow u_j e_j \) or \( u_i d_i \rightarrow Q_j L_j \) where \( i, j \) are generation indices.
Both kinds of diquark couplings for all generations are allowed by the local symmetries of the model at $N = 5$ and $N = 6$. As for the leptoquark couplings, all these terms except one vanish due to the string selection rules. The term that remains at $N = 5$ is

$$u_3^c d_3^c \bar{D}_{45} \Phi_{45} \xi_3$$

which vanishes due to SUSY ($\langle \xi_3 \rangle = 0$). Proton lifetime constrains the product of leptoquark and diquark couplings such that

$$\frac{|g_{1i} g_{4i}|}{M_{D, \bar{D}}^2} < 10^{-32} \text{ GeV}^{-2}$$

or $|g_{1i} g_{4i}| < 10^{-26}$ for $M_{D, \bar{D}} \sim O(\text{TeV})$ with a similar bound on $|g_{2i} g_{5i}|$. This requires a search up to $N = 11$ for both kinds of terms (and to higher orders for one if the other appears at $N < 11$) which is difficult to do due to the very large number of terms at these orders. As before, we assume that if unwanted terms appear, they can be made to vanish by an appropriate choice of vanishing VEVs. If this cannot be done, one has to give very large masses to $D_{45}$ and $\bar{D}_{45}$ in order to satisfy the constraint, Eq.(16), from B violation.

5. Discussion and Conclusions

We found that, under certain conditions, standard–like superstring models have two scalar and two fermionic leptoquarks with $M_{D, \bar{D}} \sim O(\text{TeV})$ and the quantum numbers given above. $M_{D, \bar{D}} \sim O(\text{TeV})$ only for light ($\sim O(100 \text{ GeV})$) squarks and sleptons. Moreover, $D_{45}$ and $\bar{D}_{45}$ have Yukawa couplings to leptoquarks which are weaker than $\sim 10^{-3}$. As we saw, SUSY constraints in the observable and hidden sectors play an important role in these results. $D_{45}$ also mixes with the down quark (and possibly with s and b). These mixings are small enough (compared to $M_D$) to satisfy the unitarity constraints on the CKM matrix. FCNC constraints on the leptoquark masses can be easily satisfied due to the small Yukawa couplings. On the other hand, baryon number violation may impose severe constraints on $M_{D, \bar{D}}$ if there are Yukawa couplings to both leptoquarks and diquarks at low $N$. 
Leptoquarks also appear in $E_6$ Calabi–Yau (CY) and flipped $SU(5) \times U(1)$ models. We now compare these with the leptoquarks of standard–like model discussed above. In CY models, leptoquarks are in each 27 of $E_6$, i.e. there is a leptoquark pair for each generation. In the $SU(3)^3$ CY model [16], leptoquarks get masses at $\sim 10^{12-14}$ GeV where the gauge symmetry is broken spontaneously and decouple from the spectrum. In CY models with a rank 5 gauge group (e.g. $SU(3) \times SU(2) \times U(1)^2$ [14]) leptoquarks are light because the superpotential contains (for each generation)

$$W = \lambda_1 H_1 H_2 N + \lambda_2 D \bar{D} N$$  \hspace{1cm} (13)

in addition to leptoquark and diquark couplings of the form given in Eqs. (11) and (14). Here $N$ is a $SO(10)$ singlet whose VEV gives the Higgs mixing as well as the leptoquark masses. Since Higgs mixing has to be $< O(TeV)$ in order to get weak symmetry breaking, $M_{D, \bar{D}} < O(TeV)$ as in our case. Note that in the CY model the scale of leptoquark masses is correlated with the Higgs mixing. If Higgs mixing is small, then $M_{D, \bar{D}}$ can be smaller than $O(TeV)$. There is no such connection between Higgs mixing and leptoquark masses in standard–like models. Instead, the correlation is between the scale of the sparticle spectrum and leptoquark masses via the gravitino mass.

In CY models, either the leptoquark or the diquark Yukawa couplings (or both) are eliminated by the discrete symmetries $(-1)^{3B}$ or $(-1)^L$ respectively. (B and L are baryon and lepton numbers respectively.) The one which is not eliminated is, in general, $\sim O(1)$. In our case, the Yukawa couplings are at most $\sim 10^{-3}$ and probably smaller. Also the discrete symmetries of CY models do not exist in standard–like models since B and L are not good quantum numbers but only (local) $B - L$ is. The explicit terms in Eqs. (12) and (15) are counterexamples to these discrete symmetries.

Flipped $SU(5) \times U(1)$ models [2] must also have leptoquarks (called vector–like heavy quarks) in 5 or 10 representations of $SU(5)$ for the gauge coupling unification
scale to be about the string unification scale, $10^{18} \text{ GeV}$ [17]. In the minimal case, only one pair of these is needed but cases with more than one pair are also possible. The leptoquark mass, in this case, can be computed from the requirement of gauge coupling unification and is given by [17]

$$m_D = 12.4 \left( \frac{\text{TeV}}{\tilde{m}} \right)^{37/12} \times (11.04^{+1}) \text{GeV}$$

(14)

where $\tilde{m}$ is the gaugino or squark mass. Thus, $m_D$ can easily be around or less than the $\text{TeV}$ scale as in our case. It has been noted that, in these models, squark masses decrease as the leptoquark masses increase which is exactly the opposite of what happens in standard–like models. This point may be instrumental in distinguishing between them. Much cannot be said about (leptoquark induced) B and L violation in the flipped $SU(5)$ case since the Yukawa couplings of leptoquarks have not been calculated.

If $\text{TeV}$ scale leptoquarks are observed, one can think of a number of scenarios in which it would be possible to distinguish between the different superstring models. For example if a) more than one pair of leptoquarks or b) leptoquarks with Yukawa couplings of $O(1)$ or c) sparticle masses much larger than $O(100 \text{ GeV})$ are observed, standard–like models will probably be ruled out. If no leptoquarks are observed at the $\text{TeV}$ scale, both standard–like models and $SU(3)^3 \text{ CY}$ models are possible but the rank 5 CY and flipped $SU(5)$ models are not. Finally, if the amount of Higgs mixing turns out to be different than leptoquark masses rank 5 CY models are ruled out. Standard–like and flipped $SU(5)$ leptoquarks are very similar to each other. It seems that the only way to distinguish between them, until Yukawa couplings of the latter are known in more detail, is by considering the sparticle mass scales in these models. Then one can use the correlation (anti–correlation) between sparticle and leptoquark masses as a possible signature of standard–like (flipped $SU(5)$) models.

The leptoquarks of standard–like models are also interesting because of their mixing with the down–like quarks. Since the CKM matrix is determined mainly by
$M_d$ in these models [10], one may try to obtain the quark mixing only from these terms. In fact, this idea has been explored in flipped $SU(5)$ models in a qualitative manner [18]. In addition, this mixing may realize the Nelson–Barr mechanism which is a possible solution to the strong CP problem, naturally. These issues are currently under investigation and will be reported in the future.

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