Fractional charges in a superstring-derived standardlike model

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(Received 21 October 1991)

We investigate the presence of fractionally charged states in the massless sector of our superstring-derived standardlike model. We show that these states are divided into three classes. The states in two of these classes will receive a Planck mass, at the trilinear level of the superpotential, by giving a vacuum expectation value to a set of neutral singlets in the string model, while the states of the third class are exotic color triplets and are therefore confined and may become superheavy at the nonrenormalizable level of the superpotential.

PACS number(s): 12.10.Gq, 11.17.+y, 11.30.Pb, 14.80.Pb

Superstring theory [1] is a unique candidate for the consistent unification of all the known fundamental interactions. However, this miraculous unification occurs at an energy scale which is much above our observational capability. The only way we know how to connect between the scale of the string and the low-energy scale is by using $N=1$ space-time supersymmetry. However, by using space-time supersymmetry, we are faced with the problem of proton decay. In the supersymmetric standard model we have the trilinear terms in the superpotential that can lead to fast proton decay,

$$\eta_1 u_L^C d_L^C d_L^C + \eta_2 d_L^C Q L \ ,$$

where generation indices are suppressed. If η_1, η_2 are of order 1, the proton will decay in a matter of minutes. These terms can be avoided if the gauge symmetry of the standard model is extended by an additional U(1) gauge symmetry which is a combination of $B-L$, baryon number minus lepton number, and T_{3_R} [2]. Moreover, it may eventually turn out that the only possible way to avoid proton decay in a supersymmetric string model is by having gauged $B-L$ for the following reason: If we extend the gauge symmetry, the matter spectrum is extended to include the right-handed neutrinos, N_L^C. The effective quartic, gauge-invariant terms can then be formed:

$$\eta_1 (u_L^C d_L^C d_L^C N_L^C)\Phi + \eta_2 (d_L^C Q L N_L^C)\Phi \ ,$$

where Φ is a combination of fields that fixes the string selection rules [3] and gets a vacuum expectation value (VEV) of $O(m_{\mathrm{Pl}})$. Because of the nonperturbative nature of the string, one has to check that such terms are not induced to all orders of such nonrenormalizable terms. Even though it is not possible to rule out the existence of a symmetry that will forbid those terms to all orders, it is hard to believe that such terms will not be induced to all orders of such nonrenormalizable terms. Therefore, if we assume that the $B-L$ symmetry is broken by a VEV of the right-handed neutrino, it is seen that the scale of the $B-L$ symmetry breaking controls the rate of proton decay. In Ref. [4], we derived a string model in the four-dimensional free fermionic formulation [5]. The gauge symmetry after the application of the Dine-Seiberg-Witten [6] mechanism is $SU(3)\times SU(2)\times U(1)_{B-L} \times U(1)_{T_{3_R}}$.[1] In Refs. [7,8], motivated by proton decay, we considered the possibility that the additional U(1) symmetry remains unbroken down to low energies. It was found that the extra Z' is expected at $M_{Z'}\sim 1$ TeV [7] and that the τ neutrino is expected at a mass scale of ~ 10 MeV, while the lighter neutrinos can be light enough to avoid any cosmological problems because of the heaviness of the top quark relative to the other quarks [8].

In spite of its appealing phenomenology, our model suffers from the problem of fractional charges in the massless sector of the string model. It has been observed in several papers that the string theory has some difficulties with fractional charges. In Ref. [9] the presence of fractionally charged particles in Calabi-Yau compactification with an E_6 gauge group broken by Wilson lines was discussed. Schellekens has recently argued [10] that their presence is a generic feature of models derived from the superstring with a level-one Kac-Moody current algebra. While many experimental searches for fractional charges have been conducted, no reported observation of a fractionally charged particle has ever been confirmed and there are upper bounds on the abundance of any such particles in the range of $10^{-19}-10^{-26}$ [11] of the nucleon abundance for charges between $\frac{1}{3}$ and 1. This may be a fundamental property of nature or merely an accidental property of the low-energy spectrum that we have been able to observe so far, and indeed fractionally charged particles may exist provided they are sufficiently heavy or sufficiently rate. In this paper we investigate the presence of fractionally charged states in the massless sector of our superstring-derived standardlike model. We show that these states can be divided into three classes and that most of them can be made extremely heavy (m_{Pl}), at the trilinear level of the superpotential, by giving a VEV to some neutral singlets in the string model, while the rest of those states are exotic color triplets and are therefore confined by QCD and may be expected to become superheavy at the unrenormalizable level of the superpotential.

Our model is constructed by using the free fermionic

[1] In Ref. [4] we have $U(1)_C = \frac{3}{2} U(1)_{B-L}$ and $U(1)_L = 2U(1)_{T_{3_R}}$.

formulation [5]. In this formulation the extra degrees of freedom which are needed to cancel the conformal anomaly are interpreted as internal fermionic degrees of freedom propagating on the string world sheet. Under parallel transport around a noncontractible loop, the fermionic states pick up a phase and the specification of the phases for all the fermionic states around all the noncontractible loops constitute a spin structure of the model. Requiring the partition function to be invariant under modular transformation leads to a set of constraints on the possible spin structures. A model is constructed by choosing a set of boundary conditions that satisfies these constraints. The massless physical spectrum is then obtained by applying the generalized Gliozzi-Scherk-Olive (GSO) projections. In Ref. [4], we constructed such a model which is generated by the following basis of eight vectors (including the sector 1) of boundary conditions for all the world-sheet fermions:

$$S=(\underbrace{1,\ldots,1}_{\psi^\mu,\chi^{1\cdots6}},0,\ldots,0|0,\ldots,0)\,, \tag{1a}$$

$$b_1=(\underbrace{1,\ldots,1}_{\psi^\mu,\chi^{12},y^{3\cdots6},\bar{y}^{3\cdots6}},\quad 0,\ldots,0|\underbrace{1,\ldots,1}_{\bar{\psi}^{1\cdots5},\bar{\eta}^1},0,\ldots,0)\,, \tag{1b}$$

$$b_2=(\underbrace{1,\ldots,1}_{\psi^\mu,\chi^{34},y^1\omega^6,y^2\bar{y}^2,\omega^5\bar{\omega}^5,\bar{y}^1\bar{\omega}^6},\quad 0,\ldots,0|\underbrace{1,\ldots,1}_{\bar{\psi}^{1\cdots5},\bar{\eta}^2},0,\ldots,0)\,, \tag{1c}$$

$$b_3=(\underbrace{1,\ldots,1}_{\psi^\mu,\chi^{56},\omega^1\omega^3,\omega^2\bar{\omega}^2,\bar{\omega}^1\bar{\omega}^3,\omega^4\bar{\omega}^4},\quad 0,\ldots,0|\underbrace{1,\ldots,1}_{\bar{\psi}^{1\cdots5},\bar{\eta}^3},0,\ldots,0)\,, \tag{1d}$$

$$b_4=(\underbrace{1,\ldots,1}_{\psi^\mu,\chi^{12},y^3y^6,\bar{y}^3\bar{y}^6,\omega^4\bar{\omega}^4,\omega^5\bar{\omega}^5},\quad 0,\ldots,0|\underbrace{1,\ldots,1}_{\bar{\psi}^{1\cdots5},\bar{\eta}^1},0,\ldots,0)\,, \tag{1e}$$

$$\alpha=(\underbrace{1,\ldots,1}_{\psi^\mu,\chi^{56},\omega^1\omega^3,\bar{y}^3\bar{y}^6,\bar{y}^1\bar{\omega}^6,\omega^4\bar{\omega}^4,y^2\bar{y}^2},\quad 0,\ldots,0|\underbrace{1,\ldots,1}_{\bar{\psi}^{1\cdots3},\bar{\eta}^1,\bar{\eta}^2}\ \underbrace{0,\ldots,0}_{\psi^{4,5},\bar{\eta}^3},\mathbf{A})\,, \tag{1f}$$

$$\beta=(\underbrace{1,\ldots,1}_{\psi^\mu,\chi^{34},y^5\bar{y}^5,\bar{y}^3\bar{y}^6,y^1\omega^6,\bar{y}^1\bar{\omega}^6,\omega^2\bar{\omega}^2},\quad 0,\ldots,0|\underbrace{\tfrac{1}{2},\ldots,\tfrac{1}{2}}_{\bar{\psi}^{1\cdots5},\bar{\eta}^{1,2,3}},\mathbf{B})\,, \tag{1g}$$

with

$$\mathbf{A}=(\underbrace{1,\ldots,1}_{\bar{\phi}^{1\cdots4}},\underbrace{0,\ldots,0}_{\bar{\phi}^{5\cdots8}}) \tag{2a}$$

and

$$\mathbf{B}=(\underbrace{1,\ldots,1}_{\bar{\phi}^{3,4}},\underbrace{\tfrac{1}{2},\ldots,\tfrac{1}{2}}_{\bar{\phi}^1,\bar{\phi}^{5,6,7}},\underbrace{0,\ldots,0}_{\bar{\phi}^2,\bar{\phi}^8})\,. \tag{2b}$$

In this notation 1 stands for periodic fermions, 0 for antiperiodic, and $\frac{1}{2}$ for those twisted by a phase $-i$. The vertical line separates real from complex fermions.

We have chosen a basis in which all left movers ($\psi^\mu,\chi^i,y^i,\omega^i$: $i=1,\ldots,6$) are real, among which world-sheet supersymmetry is realized nonlinearly, 12 right movers are real ($\bar{y}^i,\bar{\omega}^i$) and 16 right movers ($\bar{\psi}^{1\cdots5},\bar{\eta}^1,\bar{\eta}^2,\bar{\eta}^3,\bar{\phi}^{1\cdots8}$) are complex. Pairs of two real fermions that have the same boundary conditions in all

TABLE I. Fractionally charged states of class (A) with charges $\pm\frac{1}{2}$.

F	Sector	$SU(3)_C\times SU(2)_L$	Q_C	Q_L	Q_1	Q_2	Q_3	Q_4	Q_5	Q_6	SO(4)×SU(3)	Q_{11}	Q_7	Q_8	Q_9
V_{41}	$1+b_1+\alpha+2\beta$	(1,1)	0	1	0	$\frac{1}{2}$	0	$\frac{1}{2}$	0	$-\frac{1}{2}$	(1,1)	0	$\frac{1}{2}$	0	$-\frac{1}{2}$
V_{42}		(1,1)	0	-1	0	$-\frac{1}{2}$	0	$\frac{1}{2}$	0	$\frac{1}{2}$	(1,1)	0	$-\frac{1}{2}$	0	$\frac{1}{2}$
V_{43}		(1,1)	0	1	0	$-\frac{1}{2}$	0	$\frac{1}{2}$	0	$-\frac{1}{2}$	(1,1)	0	$-\frac{1}{2}$	0	$\frac{1}{2}$
V_{44}		(1,1)	0	-1	0	$\frac{1}{2}$	0	$\frac{1}{2}$	0	$\frac{1}{2}$	(1,1)	0	$\frac{1}{2}$	0	$-\frac{1}{2}$
V_{45}		(1,2)	0	0	0	$-\frac{1}{2}$	0	$-\frac{1}{2}$	0	$\frac{1}{2}$	(1,1)	0	$\frac{1}{2}$	0	$-\frac{1}{2}$
V_{46}		(1,2)	0	0	0	$\frac{1}{2}$	0	$-\frac{1}{2}$	0	$-\frac{1}{2}$	(1,1)	0	$-\frac{1}{2}$	0	$\frac{1}{2}$
V_{47}	$I+\alpha+2\beta$	(1,1)	0	1	0	0	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	0	(1,1)	0	$-\frac{1}{2}$	0	$\frac{1}{2}$
V_{48}		(1,1)	0	-1	0	0	$-\frac{1}{2}$	$\frac{1}{2}$	$-\frac{1}{2}$	0	(1,1)	0	$\frac{1}{2}$	0	$-\frac{1}{2}$
V_{49}		(1,1)	0	1	0	0	$-\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	0	(1,1)	0	$\frac{1}{2}$	0	$-\frac{1}{2}$
V_{50}		(1,1)	0	-1	0	0	$\frac{1}{2}$	$\frac{1}{2}$	$-\frac{1}{2}$	0	(1,1)	0	$-\frac{1}{2}$	0	$\frac{1}{2}$
V_{51}		(1,2)	0	0	0	0	$-\frac{1}{2}$	$-\frac{1}{2}$	$\frac{1}{2}$	0	(1,1)	0	$-\frac{1}{2}$	0	$\frac{1}{2}$
V_{52}		(1,2)	0	0	0	0	$\frac{1}{2}$	$-\frac{1}{2}$	$-\frac{1}{2}$	0	(1,1)	0	$\frac{1}{2}$	0	$-\frac{1}{2}$

TABLE II. Fractionally charged states of class (B) with charges $\pm\frac{1}{2}$.

F	Sector	$SU(3)_C\times SU(2)_L$	Q_C	Q_L	Q_1	Q_2	Q_3	Q_4	Q_5	Q_6	$SO(4)\times SU(3)$	Q_H	Q_7	Q_8	Q_9
H_1	$\pm\beta$	(1,1)	$\frac{3}{4}$	$\frac{1}{2}$	$\frac{1}{4}$	$\frac{1}{4}$	$\frac{1}{4}$	$-\frac{1}{2}$	$\frac{1}{2}$	0	(2,1)	$\frac{3}{4}$	$\frac{1}{4}$	0	0
H_2		(1,1)	$-\frac{3}{4}$	$-\frac{1}{2}$	$-\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{2}$	$-\frac{1}{2}$	0	(2,1)	$-\frac{3}{4}$	$-\frac{1}{4}$	0	0
H_3	$I\pm\beta$	(1,1)	$\frac{3}{4}$	$\frac{1}{2}$	$\frac{1}{4}$	$\frac{1}{4}$	$\frac{1}{4}$	$\frac{1}{2}$	$\frac{1}{2}$	0	(1,1)	$-\frac{3}{4}$	$-\frac{1}{4}$	$-\frac{1}{2}$	$\frac{1}{2}$
H_4		(1,1	$-\frac{3}{4}$	$-\frac{1}{2}$	$-\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{4}$	$\frac{1}{2}$	$-\frac{1}{2}$	0	(1,1)	$\frac{3}{4}$	$\frac{1}{4}$	$\frac{1}{2}$	$-\frac{1}{2}$
H_5		(1,1)	$\frac{3}{4}$	$\frac{1}{2}$	$\frac{1}{4}$	$\frac{1}{4}$	$\frac{1}{4}$	$\frac{1}{2}$	$-\frac{1}{2}$	0	(1,1)	$-\frac{3}{4}$	$-\frac{1}{4}$	$\frac{1}{2}$	$-\frac{1}{2}$
H_6		(1,1)	$-\frac{3}{4}$	$-\frac{1}{2}$	$-\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{4}$	$\frac{1}{2}$	$\frac{1}{2}$	0	(1,1)	$\frac{3}{4}$	$\frac{1}{4}$	$-\frac{1}{2}$	$\frac{1}{2}$
H_7	$1+b_4\pm\beta$	(1,1)	$\frac{3}{4}$	$\frac{1}{2}$	$\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{2}$	$\frac{1}{2}$	0	(1,1)	$-\frac{3}{4}$	$-\frac{1}{4}$	$\frac{1}{2}$	$\frac{1}{2}$
H_8		(1,1)	$-\frac{3}{4}$	$-\frac{1}{2}$	$-\frac{1}{4}$	$\frac{1}{4}$	$\frac{1}{4}$	$-\frac{1}{2}$	0	$-\frac{1}{2}$	(1,1)	$\frac{3}{4}$	$\frac{1}{4}$	$-\frac{1}{2}$	$-\frac{1}{2}$
H_9		(1,1)	$\frac{3}{4}$	$\frac{1}{2}$	$\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{2}$	0	$-\frac{1}{2}$	(1,1)	$-\frac{3}{4}$	$-\frac{1}{4}$	$-\frac{1}{2}$	$-\frac{1}{2}$
H_{10}		(1,1)	$-\frac{3}{4}$	$-\frac{1}{2}$	$-\frac{1}{4}$	$\frac{1}{4}$	$\frac{1}{4}$	$-\frac{1}{2}$	0	$\frac{1}{2}$	(1,1)	$\frac{3}{4}$	$\frac{1}{4}$	$\frac{1}{2}$	$\frac{1}{2}$
H_{11}	$I+1+b_4\pm\beta$	(1,1)	$\frac{3}{4}$	$\frac{1}{2}$	$\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{4}$	$\frac{1}{2}$	0	$\frac{1}{2}$	(1,1)	$\frac{3}{4}$	$\frac{1}{4}$	0	0
H_{12}		(1,1)	$\frac{3}{4}$	$\frac{1}{2}$	$\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{4}$	$\frac{1}{2}$	0	$\frac{1}{2}$	(1,1)	$\frac{3}{4}$	$\frac{1}{4}$	0	0
H_{13}		(1,1)	$-\frac{3}{4}$	$-\frac{1}{2}$	$-\frac{1}{4}$	$\frac{1}{4}$	$\frac{1}{4}$	$\frac{1}{2}$	0	$-\frac{1}{2}$	(1,1)	$-\frac{3}{4}$	$-\frac{1}{4}$	0	0
H_{14}		(1,1)	$-\frac{3}{4}$	$-\frac{1}{2}$	$-\frac{1}{4}$	$\frac{1}{4}$	$\frac{1}{4}$	$\frac{1}{2}$	0	$-\frac{1}{2}$	(1,1)	$-\frac{3}{4}$	$-\frac{1}{4}$	0	0

the sectors are paired to form complex fermions. The specific choice of the projections coefficient

$$c\begin{pmatrix} b_i \\ b_j \end{pmatrix},$$

which were made in our string model, is summarized in the following matrix form:

$$\begin{array}{c} \\ 1 \\ S \\ b_1 \\ b_2 \\ b_3 \\ b_4 \\ \alpha \\ \beta \end{array}\begin{array}{cccccccc} 1 & S & b_1 & b_2 & b_3 & b_4 & \alpha & \beta \end{array}\hspace{-25em}\begin{array}{c}\\ \end{array}$$

$$\begin{array}{c|cccccccc|} & 1 & S & b_1 & b_2 & b_3 & b_4 & \alpha & \beta \\ \hline 1 & -1 & 1 & -1 & -1 & -1 & 1 & -1 & i \\ S & 1 & 1 & 1 & 1 & 1 & 1 & 1 & 1 \\ b_1 & -1 & -1 & -1 & -1 & -1 & -1 & -1 & i \\ b_2 & -1 & -1 & -1 & -1 & -1 & -1 & -1 & i \\ b_3 & -1 & -1 & -1 & -1 & -1 & -1 & -1 & -1 \\ b_4 & 1 & -1 & -1 & -1 & -1 & 1 & -1 & i \\ \alpha & -1 & -1 & -1 & -1 & -1 & -1 & -1 & 1 \\ \beta & 1 & -1 & 1 & -1 & 1 & 1 & -1 & 1 \end{array} \quad . \tag{3}$$

The model generated by this basis of boundary conditions has the following properties: (1) $N=1$ space-time supersymmetry; (2) three generations of chiral fermions; (3) the gauge group is $SU(3)_C\times SU(2)_L\times U(1)^n$, where we argued that n reduces to 1 or 2 after applying the Dine-Seiberg-Witten mechanism [6]; (4) there are enough scalar doublets and singlets to break the gauge symmetry in a realistic way; and (5) all trilinear terms that give fast proton decay are absent from the superpotential because of the extra U(1) symmetries and the absence of Higgs triplets from the Neveu-Schwarz sector.

The full massless spectrum was derived in Ref. [4]; here we list in Tables I–III only the states with fractional charges. The massless fractional charged states are divided into three classes.

(A) Sectors of the form $\gamma=\alpha+\sum_j k_j b_j$, $k_j=0,1$; $j=1,\ldots,4$ with $(\gamma\cdot\gamma)_R=8$ and γ $(\bar\psi_{1,\ldots,3}=0$ and $\bar\psi_{45}=1)$. In these sectors all the fermionic states have only periodic or antiperiodic boundary conditions. These states are shown in Table I and carry electric charges $\pm\frac{1}{2}$.

(B) Sectors of the form $\gamma=\beta+\sum_j k_j b_j$, $k_j=0,1$; $j=1,\ldots,4$ with $(\gamma\cdot\gamma)_R=8$. In these sectors some of the fermionic states have rational $\frac{1}{2}$ boundary conditions. These states are shown in Table II and carry electric charges $\pm\frac{1}{2}$.

(C) Sectors of the form $\gamma=\beta+\alpha+\sum_j k_j b_j$, $k_j=0,1$; $j=1,\ldots,4$ with $(\gamma\cdot\gamma)_R=6$. In this case the fractionally charged states are obtained by acting on the vacuum with $\bar\psi_{1,2,3}$ and $\bar\psi^*_{1,2,3}$ to give electric charges $\pm\frac{1}{3}$. These states are shown in Table III.

The trilinear superpotential of our string model is straightforwardly calculated following the rules given in Ref. [3] and is given by $W=\sqrt{2}g\,[W_1+W_2]$, where

$$W_1=\{(u^c_{L_1}Q_1\bar h_1+N^c_{L_1}L_1\bar h_1+d^c_{L_2}Q_2h_2+e^c_{L_2}L_2h_2+e^c_{L_3}L_3h_3+d^c_{L_3}Q_3h_3)+h_2\bar h_1\bar\phi_{12}+\bar h_2h_1\phi_{12}$$
$$+h_3\bar h_1\phi_{12}+\bar h_3h_1\phi_{13}+h_3\bar h_2\phi_{23}+\bar h_3h_2\phi_{23}+\phi_{12}\bar\phi_{13}\bar\phi_{23}+\bar\phi_{12}\phi_{13}\phi_{23}+(\phi_4\bar\phi_4'+\bar\phi_4\phi_4')\phi_1\}\ , \tag{4a}$$

TABLE III. Fractionally charged states of class (C) with charges $\pm\frac{1}{3}$.

F	Sector	$SU(3)_C\times SU(2)_L$	Q_C	Q_L	Q_1	Q_2	Q_3	Q_4	Q_5	Q_6	$SO(4)\times SU(3)$	Q_H	Q_7	Q_8	Q_9
H_{33}	$b_3+\alpha\pm\beta$	(3,1)	$-\frac{1}{4}$	$-\frac{1}{2}$	$-\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{4}$	0	0	$-\frac{1}{2}$	(1,1)	$\frac{3}{4}$	$-\frac{1}{4}$	$\frac{1}{2}$	0
H_{40}	$b_1+b_2+b_4+\alpha$	$(\bar 3,1)$	$\frac{1}{4}$	$\frac{1}{2}$	$\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{4}$	0	$-\frac{1}{2}$	0	(1,1)	$-\frac{3}{4}$	$\frac{1}{4}$	$\frac{1}{2}$	0

$$W_2=\frac{1}{\sqrt{2}}\{H_1H_2\phi_4+H_3H_4\bar{\phi}_4+H_5H_6\bar{\phi}_4+(H_7H_8+H_9H_{10})\phi_4'+(H_{11}+H_{12})(H_{13}+H_{14})\bar{\phi}_4'+V_{41}V_{42}\bar{\phi}_4+V_{43}V_{44}\bar{\phi}_4+V_{45}V_{46}\phi_4+(V_{47}V_{48}+V_{49}V_{50})\bar{\phi}_4'+V_{51}V_{52}\phi_4'\}+[H_{15}H_{16}\phi_{56}'+H_{17}H_{18}\bar{\phi}_{56}+H_{19}H_{20}\bar{\phi}_{56}'+H_{21}H_{22}\bar{\phi}_{56}+(V_{11}V_{12}+V_{13}V_{14})\phi_{13}+(V_{15}+V_{16})(V_{17}+V_{18})\phi_{13}+V_{19}V_{20}\phi_{13}+V_{21}V_{22}\phi_{12}+V_{23}V_{24}\phi_{12}+(V_{25}+V_{26})(V_{27}+V_{28})\phi_{12}+V_{29}V_{30}\phi_{12}+V_{31}V_{32}\bar{\phi}_{23}+V_{33}V_{34}\phi_{23}+H_{29}H_{30}\bar{\phi}_{13}+H_{36}H_{37}\phi_{12}]\ . \quad (4b)$$

W_1 includes states in the observable sector only and W_2 has couplings between neutral singlets in the observable sector and hidden states, most of them singles as well.

By examining the fractionally charged states and the trilinear superpotential, it is seen that all the states from class (A) and (B) receive a Planck mass by giving a VEV to the neutral singlets $\bar{\phi}_4$, $\bar{\phi}_4'$, ϕ_4, and ϕ_4' which impose the additional F flatness constraint $(\phi_4\bar{\phi}_4'+\bar{\phi}_4\phi_4')=0$.

The states in class (C) come in pairs, so for each state of this form, its complex conjugate exists in the massless spectrum as indeed can be seen by examining Table III. The states of class (C) are exotic color triplets and are therefore confined by QCD. However, because these states come from different sectors in the additive group, they carry different quantum numbers under the additional horizontal U(1) symmetries and therefore they can be expected to receive a Planck mass only at the nonrenormalizable level. In our model such a nonrenormalizable term, which satisfies all the string selection rules [3], appears at the quintic level

$$W=H_{33}H_{40}H_{31}H_{38}\phi_{23}\ , \quad (5)$$

where H_{31}, H_{38}, and ϕ_{23} are neutral under the standard-model gauge group. Giving a VEV to these standard-model singlets makes the exotic quark triplets heavy. These VEV's are constrained by the nonobservation of exotic quark states at CERN e^+e^- collider LEP and the Collider Detector at Fermilab (CDF). A possible numerical scenario is given by taking $\langle\bar{\phi}_{23}\rangle=\langle\phi_{23}\rangle\sim M_{\rm Pl}$, along an F and D flat direction, which constrains $\langle H_{31}\rangle\langle H_{38}\rangle>$ to be above $\sim10^{10}$ GeV. A complete determination of these VEV's can only be obtained by resolving the problem of supersymmetry breaking in superstring models and is beyond the scope of this paper. The additional horizontal U(1) charges also make these states harmless from the point of view of proton decay. It is interesting to note that the states from class (C) are associated with the symmetry breaking down to SU(3)$\times$SU(2) at the string level, as can be observed by examining fractionally charged states of other models in the same formulation. The fractional charged states of class (A) are associated with the symmetry breaking SO(10)$\rightarrow$SO(6)$\times$SO(4) and they are found, for example, in the model of Ref. [12] in which the gauge symmetry at the string level is broken to SO(6)$\times$SO(4) and is not further broken to $SU(3)\times SU(2)\times U(1)_{B-L}\times U(1)_{T_{3_R}}$. The states of class (B) are found in the model of Ref. 13 and it was suggested that, in this model, all those fractionally charged states will be confined by a gauge group of the hidden sector [13,14]. They result because of the use of the rational $\frac{1}{2}$ boundary conditions to break the O(n) gauge symmetry to a U(n) gauge symmetry. The sectors that give rise to the fractional charges of class (C) are found in the model of Ref. [13] as well; however, there, because of the orthogonality of the b_4 and b_5 boundary conditions, those states are projected out while in the case of our model the b_4 and α projections are not orthogonal with respect to those states because α breaks the SO(10) symmetry to SO(6)$\times$SO(4). These states are therefore a unique characteristic of the superstring-derived standardlike models. Their phenomenological implications will be investigated in a future publication.

To conclude, in this paper we have investigated the presence of fractionally charged states in our superstring-derived standardlike model. We have shown that, by giving a VEV to a set of neutral singlets in the string model, most of the fractionally charged states become superheavy while the rest are confined by QCD to give integrally charged states. We showed that these exotic states become heavy at the quintic level of the superpotential. We therefore conclude that, in the superstring standardlike models, the problem with fractionally charged states can be evaded by decoupling of these states from the massless spectrum.

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