ON POSSIBLE GLOBAL SYMMETRIES IN SUPERSTRING INSPIRED SUPERSYMMETRIC MODELS

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Received 19 February 1987

We investigate phenomenological roles of possible global symmetries in a class of superstring inspired models. Such global symmetries may be present in the full string theory or may incidentally appear in the effective theory of the low-energy sector. We find that if a model has a discrete symmetry, Z_m, which acts in a generation-independent way, it is naturally embedded in a U(1) symmetry of the low-energy lagrangian. Then, such global U(1) symmetries (through suitable redefinition in terms of the local U(1)'s) remain even after the spontaneous breakdown of the gauge symmetries at $O(M_W)$. Some discussion is also presented for generation-dependent symmetries.

It has recently been discovered that superstring theories may present a mathematically consistent description for the unification of all the fundamental particle spieces and their interactions including gravity [1–3]. Then, it is expected that the low-energy phenomenologies of the superstring theories may be simulated by a class of supersymmetric models [2,4–9]. The gauge symmetry of such superstring inspired models may be given by $K\times K'$ which results from the gauge group $E_8\times E_8'$ of the original heterotic string theory through the compactification of the extra six-dimensional space admitting a nontrivial Wilson loop [4,5,10]. The chiral superfields surviving to low energies will be singlet under K' and transform as certain representations of K in the decomposition of 27 or $\overline{27}$ of E_6.

One of the interesting features of the string inspired supersymmetric models is that due to the discrete symmetries and/or the topology of the internal six-dimensional space, some couplings in the superpotential may vanish so as to reproduce successful low-energy phenomenology, e.g., prohibition of the fast proton decay mediated by the scalar quarks [4,11]. Then, the effective low-energy lagrangian below $M_{\rm Compact}\sim M_{\rm Planck}$ including the associated soft supersymmetry breakings may exhibit certain global continuous symmetries (except for explicit breakings of order $M_{\rm Pl}^{-n}$ ($n>0$) induced by integrating out the superheavy sector). For example, in the models considered in ref. [9], a Z_9 symmetry is embedded in a $U(1)_{PQ}$ symmetry and a U(1) of lepton number appears "automatically".

In this letter, we would like to investigate phenomenological roles of such global symmetries in a class of low-energy supersymmetric models derived form the superstring.

We begin with the consideration of a simple model of

$$K=SU(3)_C\times SU(2)_W \times U(1)_Y\times U(1)_I\times U(1)_J\,, \tag{1}$$

with rank 6, where the U(1)'s are embedded in a maximal subgroup $SU(6)\times SU(2)$ of E_6 such that $U(1)_Y\times U(1)_I\subset SU(6)$ and $U(1)_J\subset SU(2)$, respectively, and Y is the usual weak hypercharge. The chiral superfields contained in a 27 of E_6 are identified by the K-quantum numbers (C, W, Y, I, J) as

$$q\sim(3,2,\tfrac{1}{6},-\tfrac{1}{3},0)\,,$$

$$u^c\sim(\bar{3},1,-\tfrac{2}{3},-\tfrac{1}{3},0),\quad d^c\sim(\bar{3},1,\tfrac{1}{3},\tfrac{1}{6},\tfrac{1}{2})\,,$$

$$D^c\sim(\bar{3},1,\tfrac{1}{3},\tfrac{1}{6},-\tfrac{1}{2}),\quad D\sim(3,1,-\tfrac{1}{3},\tfrac{2}{3},0)\,, \tag{2}$$

$$\ell \sim (1, 2, -\tfrac{1}{2}, \tfrac{1}{6}, \tfrac{1}{2}), \quad e^c \sim (1, 1, 1, -\tfrac{1}{3}, 0) ,$$

$$H \sim (1, 2, -\tfrac{1}{2}, \tfrac{1}{6}, -\tfrac{1}{2}), \quad H^c \sim (1, 2, \tfrac{1}{2}, \tfrac{2}{3}, 0) ,$$

$$N \sim (1, 1, 0, -\tfrac{5}{6}, \tfrac{1}{2}), \quad N^c \sim (1, 1, 0, -\tfrac{5}{6}, -\tfrac{1}{2}) . \tag{2 cont'd}$$

The most general K-invariant superpotential is given by

$$W = \lambda_1 u^c q H^c + \lambda_2 d^c q H + \lambda_3 e^c \ell H + \lambda_4 N D^c D + \lambda_5 N H^c H + \lambda_6 N^c H^c \ell + \lambda_7 N^c d^c D + \lambda_8 D^c q \ell + \lambda_9 D u^c e^c + \lambda_{10} q q D + \lambda_{11} u^c d^c D^c , \tag{3}$$

which is obtained by decomposing $27 \cdot 27 \cdot 27$ of E_6. Here, the generation indices are suppressed for simplicity; these couplings should be regarded as, e.g., $(\lambda_1)_{\alpha\beta\gamma} u^c_\alpha q_\beta H^c_\gamma$ (α, β, $\gamma = 1$–3) in the three-generation space.

The gauge group K is spontaneously broken at $O(M_W)$ as $K \to SU(3)_C \times U(1)_{EM}$ by the VEV's of H, H^c, N and N^c. Here, it should be remarked that N^c may develop a nonzero VEV of $O(M_W)$ due to the D^2 terms in the scalar potential if the gauge coupling constants and the soflt supersymmetry breaking scalar mass terms satisfy $(g_I/g_J)^2 < [3(1+r)/5(1-r)] < \frac{3}{5}$ and $-1 < r < 0$ where $r \equiv (m^2_{N^c}/m^2_N)$ with $m^2_N < 0$. (The argument in ref. [12] to rule out $\langle N^c \rangle \neq 0$ is valid for $g_I = g_J$. However, it is possible in the superstring theory that $g_I \neq g_J$ [13].) We do not consider the possibility of gauge symmetry breaking at an intermediate scale. Some variants of our model, e.g., a rank-five model, Higgs multiplets coming from an incomplete set in $(27 + \overline{27})$, etc., will be discussed later.

We first examine possible abelian symmetries, U(1) and Z_n, in the superpotential (3) which act in a *generation-independent* way; the chiral multiplets with the same K-quantum numbers transform in the same way. (Such symmetries are also preserved in the effective low-energy lagrangian if certain couplings in (3) together with the associated soft supersymmetry breakings are vanishing.) The charge X associated with an abelian symmetry may be identified by solving linear algebraic equations for nonvanishing couplings λ_a,

$$A_{ai} X_i = n \cdot k_a \leftrightarrow \lambda_a \neq 0 , \tag{4}$$

where the k_a are certain integers without an overall common factor, and the A_{ai} are non-negative integers characterizing the relevant couplings λ_a with the chiral fields $\phi_i = q$, u^c, ..., e.g., nonzero entries are $A_{1i} = 1$ for $\phi_i = u^c$, q, H^c in the coupling λ_1, etc. The charges X_i for the chiral fields should be determined modulo n as integers without an overall common factor. Each linearly independent solution of eq. (4) realizes a U(1) ($n = 0$) or a Z_n (generated by $\exp[i(2\pi/n)X]$).

All the possible 11 equations in eq. (4) with 11 fields cannot be linearly independent with respect to X_i since W obviously preserves the local U(1)'s in $K \subset E_6$. In fact, we find that the 11×11 matrix A_{ai} has rank 8. Hence, if all the 11 couplings are nonvanishing, there is no room for additional U(1) symmetries in W other than the local $U(1)_Y \times U(1)_I \times U(1)_J$.

Some couplings in W and the associated soft supersymmetry breaking terms, however, may vanish due to the discrete symmetries and/or the topology of the internal manifold [4,11]. Then, the effective low-energy theory may incidentally exhibit larger global symmetries than the original ones in the whole sector of the superstring theory, (except for breaking terms of order M_{Pl}^{-n} ($n > 0$)). For example, suppose that a certain internal manifold results in a discrete symmetry [14]

$$P: (D, D^c, N^c) \to (D, D^c, N^c) , \tag{5}$$

so that

$$\lambda_6, \lambda_8\text{–}\lambda_{11} = 0 . \tag{6}$$

Then, besides the local U(1)'s, two global U(1) symmetries appear, corresponding to N_B (baryon number) and $N_{(D-N)^c)} \equiv N_D - N_{D^c} - N_{N^c}$ ($D - D^c - N^c$ number). It is interesting to observe that the discrete symmetry looks "promoted" to a continuous symmetry $U(1)_{(D-N^c)}$ such that $P = \exp(i\pi N_{(d-N^c)})$. (See also ref. [9] for this kind of phenomenon.)

We may proceed to find possible U(1) symmetries by considering certain phenomenological constraints. Such symmetries will characterize superstring inspired modles.

We first note that λ_1–λ_4 must be nonvanishing for the quark and lepton mass generation and that λ_5 is also inevitable to obtain $\langle \mathrm{H} \rangle \neq 0$ and $\langle \mathrm{H}^c \rangle \neq 0$. The remaining couplings λ_6–λ_{11} are optional. We here simply assume that

$$\lambda_6 = 0 , \tag{7}$$

to avoid a Dirac neutrino mass of $O(M_W)$.

The number of possible global U(1)'s is determined by the number of independent conditions (4) for *nonvanishing* λ_a. Let Γ_a indicate a linear combination of X_i; $\Gamma_a \equiv A_{ai} X_i$ (LHS of eq. (4)). Since the conditions $\Gamma_a = 0$ $(a = 1$–$5)$ for λ_1–λ_5 are linearly independent, we find

$$\#(\text{global U(1)'s}) = 3 - \#(\text{independent } \Gamma_a = 0 \text{ for nonzero } \lambda_7\text{–}\lambda_{11}) , \tag{8}$$

where the maximal number of global U(1)'s is obtained by $3 = 11$ (fields) $- 5(\lambda_1$–$\lambda_5 \neq 0) - 3$(local U(1)'s). Identity relations

$$\Gamma_6 - \Gamma_7 \simeq \Gamma_8 \simeq -\Gamma_9 \quad (\text{modulo } \Gamma_1\text{–}\Gamma_5) ,$$

$$\Gamma_{10} \simeq -\Gamma_{11} \quad (\text{modulo } \Gamma_1\text{–}\Gamma_5) , \tag{9}$$

may be useful to count the number of independent $\Gamma_a = 0$ to identify possible symmetries, where $\Gamma_a \simeq \Gamma_b$ (modulo Γ_1–Γ_5) means an identity $\Gamma_a - \Gamma_b \equiv \sum_c B_c \Gamma_c$ $(c = 1-5)$ with suitable integers B_c. These relations imply, for example, that $\Gamma_8 = 0$ for $\lambda_8 \neq 0$ automatically requires $\Gamma_9 = 0$ if Γ_1–$\Gamma_5 = 0$, i.e., as long as λ_1–$\lambda_5 \neq 0$, the survival of λ_8 and λ_9 does not give independent constraints for identifying possible global symmetries. Furthermore, since Γ_6 is not required to be zero if $\lambda_6 = 0$, $\Gamma_7 = 0$ for $\lambda_7 \neq 0$ serves as an independent condition in (8) to reduce the number of possible U(1)'s.

It should be also noticed that the couplings λ_7–λ_{11} may or may not vanish in accordance with the baryon number assignment for the extra D quarks. Two cases are available to reproduce the baryon number conservation for preventing rapid proton decay [8]:

$$N_B^{(1)}(\mathrm{D}) = \tfrac{1}{3} \leftrightarrow \lambda_{10} = \lambda_{11} = 0 , \tag{10}$$

or

$$N_B^{(2)}(\mathrm{D}) = -\tfrac{2}{3} \leftrightarrow \lambda_7 = \lambda_8 = \lambda_9 = 0 . \tag{11}$$

(We cannot assign $N_B^{(2)}(\mathrm{N}^c) = 1$ to admit nonzero λ_7 for the second case since $\langle \mathrm{N}^c \rangle$ must be nonvanishing to break the extra local U(1)'s).

We now find from the conditions (8)–(11) that possible models and the corresponding global U(1) charges are classified as follows, depending on the vanishing (denoted by "$\times$") or nonvanishing ("$\bigcirc$") of the relevant couplings:

$$\begin{aligned}
(\lambda_7, \lambda_{8,9}, \lambda_{10,11}) &= (\bigcirc, \bigcirc, \times) \to N_B^{(1)} \\
&= (\bigcirc, \times, \times) \to N_B^{(1)} \oplus N_{(\mathrm{d}-\mathrm{N}^c)} \\
&= (\times, \bigcirc, \times) \to N_B^{(1)} \oplus N_{\mathrm{N}^c} \\
&= (\times, \times, \bigcirc) \to N_B^{(2)} \oplus N_{\mathrm{N}^c} \\
&= (\times, \times, \times) \to N_B^{(1)} \oplus N_{\mathrm{D}} \oplus N_{\mathrm{N}^c}
\end{aligned} \tag{12}$$

(The cases $(\bigcirc, \bigcirc, \bigcirc)$, $(\bigcirc, \times, \bigcirc)$ and $(\times, \bigcirc, \bigcirc)$ are excluded by (11).)

One may also be interested in finding possible Z_n symmetries. We can, however, show that they are always embedded in the U(1) symmetries found so far. Suppose that there is a nontrivial solution of eq. (4) for a Z_n. Then, we can redefine the X_i charges without affecting the Z_n action on the chiral fields such that

$$X_i' \equiv X_i - n \cdot \Delta_i , \tag{13}$$

where the Δ_i are arbitrary integers. By suitably arranging the Δ_i, eq. (4) for X_i may be rewritten as

$$A_{ai} X_i' = 0 , \tag{14}$$

so that the X_i' represent a U(1) symmetry, i.e., Z_n is embedded in a U(1). For instance, for a model $(\bigcirc, \bigcirc, \times)$ in (12) with $\lambda_7 \neq 0$, and λ_8 and/or $\lambda_9 \neq 0$ for $N_B^{(1)}(\mathrm{D}) = \frac{1}{3}$, we can take

$$\Delta_{\mathrm{u}^c} = k_1 - k_5 , \quad \Delta_{\mathrm{d}^c} = k_2 , \quad \Delta_{\mathrm{e}^c} = k_3 ,$$

$$\Delta_{\mathrm{H}^c} = k_5 , \quad \Delta_{\mathrm{D}} = k_4 - \bar{k} , \quad \Delta_{\mathrm{D}^c} = \bar{k} ,$$

$$\Delta_{\mathrm{N}^c} = k_7 - k_2 - k_4 + \bar{k} , \tag{15}$$

and the other $\Delta_i = 0$ with

$$\begin{aligned}
\bar{k} &= k_8 \quad \text{for } \lambda_8 \neq 0 , \\
&= k_1 + k_3 + k_4 - k_5 - k_9 \quad \text{for } \lambda_9 \neq 0 .
\end{aligned} \tag{16}$$

These two expressions of $\bar{k}$ coincide when both λ_8 and λ_9 are nonzero, since the preservation of Z_n requires $k_8 = -k_9 + k_1 + k_3 + k_4 - k_5$ from an identity $\Gamma_8 = -\Gamma_9 + \sum_c B_c \Gamma_c$ with $B_c = (1,\ 0,\ 1,\ 1,\ -1)$ ($c = 1$–5). (If $\lambda_6 \neq 0$ as well as λ_7–$\lambda_9 \neq 0$, the identities (9) further require $k_6 - k_7 - k_8 = -k_2 - k_4 + k_5$. Then, Δ_{H^c} and Δ_{N_c} given in (15) also result in $A_{6i} X'_i = 0$ for λ_6.) We find similar Δ_i for the models with $N_B^{(2)} = -\frac{2}{3}$.

In this way, we have found that any Z_n symmetry which acts in a *generation-independent* way can be embedded in a suitable U(1) symmetry (as long as the baryon number conservation is ensured in accordance with (10) or (11)). In other words, the discrete symmetry has been "promoted" to a continuous symmetry in the low-energy effective theory [9].

Any global symmetry has a potential danger to cause phenomenological problems if it is spontaneously broken at a low mass scale; the appearance of Nambus–Goldstone bosons (or axions) for a U(1) [15], or the domain wall formation for a Z_n [16]. The gauge symmetry K of (1) may be spontaneously broken at $O(M_W \simeq 1\ \text{TeV})$ by the VEV's of H, H^c, N and N^c induced by the effects of the soft supersymmetry breaking terms. Then, one may worry if the global symmetries found so far are also broken together with the gauge group K, giving rise to phenomenological difficulties. The domain wall problem, however, does not arise if the discrete symmetries act in *generation-independent* ways, since they are embedded in the U(1) symmetries, as shown above.

As for the possible U(1) symmetries in (12) in connection with the problem of Nambu–Goldstone bosons, we will show that they actually remain unbroken. Note that X_i charges either for U(1) or Z_n in eq. (4) (or X'_i in eq. (14)) can be further redefined by taking a gauge transformation generated by Y, I and J:

$$X'' \equiv X + aY + bI + cJ\,. \tag{17}$$

If we choose

$$a = 2X_H + \tfrac{2}{5}(3X_N - 2X_{N^c})\,,$$

$$b = \tfrac{3}{5}(X_N + X_{N^c}),\quad c = -X_N + X_{N^c}\,, \tag{18}$$

we find

$$X''_i = 0 \quad (\text{modulo } n) \quad \text{for } \phi_i = H, H^c, N, N^c\,. \tag{19}$$

This is possible due to the fact that $X_H + X_{H^c} + X_N = 0$ (modulo n), derived from the term $\lambda_5 NH^cH$ in W. Therefore, the possible *generation-independent* global U(1) symmetries ($n = 0$) (and Z_n also) actually remain unbroken, and the Nambu–Goldstone bosons do not appear.

We now consider some variants of our model.

The gauge group K may be of rank five. Consider, for example, a model of $K = SU(3)_C \times SU(2)_W \times U(1)_Y \times U(1)_I$ [8]. Then, the third $U(1)_J$ in E_6, which is no longer a local symmetry, still remains as a *global* symmetry in the superpotential. As long as the possibility of global symmetries is concerned, there is no essential difference between the rank-five model and the rank-six model except that in the rank-five model, $\langle N^c \rangle$ must be zero to avoid the Nambu–Goldstone boson associated with the *global* $U(1)_J$. (The $U(1)_J$ charges of H, H^c, N and N^c cannot be simultaneously absorbed into those of Y and I, as in eq. (17).)

The Higgs doublets, H and H^c, may come rather from an incomplete set in $(27 + \overline{27})$ [14]. Then, the multiplets contained in the complete 27's which have the same quantum numbers as the Higgs doublets must be regarded as heavy leptons, say L. They aquire masses through the coupling NL^cL. The couplings such as NH^cL, however, must be absent so that we can distinguish the heavy leptons from the Higgs doublets by the L-number conservation. The same arguments for the possible global symmetries as in the first model also apply to this case.

Therefore, our conclusions concerning the possibility of *generation-independent* global symmetries will be valid in a variety of superstring inspired models through the E_6 unification without intermediate gauge symmetry breaking scale:

(i) Phenomenologically acceptable models are singled out in accordance with the baryon number assignment to prevent rapid proton decay.

(ii) If such a model has a discrete symmetry, Z_n, (which would originate from the string theory), it is naturally embedded in a U(1) symmetry (except for tiny breaking terms suppressed by M_{Pl}).

(iii) Then, the global U(1) symmetries (through suitable redefinition in terms of the local U(1)'s) remain unbroken even after the spontaneous breakdown of the gauge symmetries at $O(M_W)$.

We finally present some comments on global symmetries which act in *generation-dependent* ways. Generally, the problems of Nambu–Goldston boson, visible axion and/or domain wall will arise. This is due to the fact that some Higgs multiplets with the same K-quantum numbers may transform differently under a generation-dependent symmetry. Then, Z_n symmetries may not be embedded in U(1)'s, or the transformations of the Higgs multiplets under a U(1) (or Z_n) symmetry may not be all rotated out (modulo n) in terms of the local gauge symmetries, as has been done in the case of generation-independent symmetries. (These difficulties might be evaded in some models [9].)

We examine several cases in the concrete. Let the chiral multiplets transform under a generation-dependent symmetry as

$$\phi_\alpha(i)\to\exp[i(2\pi/n)X_\alpha(i)]\cdot\phi_\alpha(i)\ , \qquad (20)$$

where $\phi_\alpha(i)=\mathrm{q}_\alpha$, u^c_α, d^c_α,, etc., and $\alpha=1$–3 is a generation index. Then, the global symmetry under consideration requires the following conditions, in particular, for the *nonvanishing* elements of the coupling matrices $(\lambda_1)_{\alpha\beta\gamma}$ and $(\lambda_2)_{\alpha\beta\gamma}$:

$$X_\alpha(\mathrm{q})+X_\beta(\mathrm{u}^c)+X_\gamma(\mathrm{H}^c)=0 \ \text{modulo}\ n\ , \qquad (21)$$

$$X_\alpha(\mathrm{q})+X_\beta(\mathrm{d}^c)+X_\gamma(\mathrm{H})=0 \ \text{modulo}\ n\ . \qquad (22)$$

(a) The Higgs multiplets, H, H^c, N and N^c, may still transform in a generation-independent way, while the quarks and leptons transform in a generation-dependent way:

$$X_\gamma(i)=X(i) \quad (i=\mathrm{H},\mathrm{H}^c,\mathrm{N},\mathrm{N}^c;\gamma=1\text{–}3)\ . \qquad (23)$$

Then, the phenomenological problems are avoided since the charges of the Higgs multiplets can be rotated out, as in (17)–(19).

It is interesting to see whether this kind of symmetry reproduces the quark mass matrices of Fritzsch-type [17]:

$$M_\mathrm{f}\sim \begin{array}{cc} & f_\beta \\ f^c_\alpha & \begin{bmatrix} 0 & * & 0\\ * & 0 & * \\ 0 & * & * \end{bmatrix} \end{array} \ [\mathrm{f}=\mathrm{u},\mathrm{d}], \qquad (24)$$

where "*" denotes a nonzero entry. It is, however, immediately found that the symmetry conditions (21) and (22) for the nonzero elements in (24) require that for *all* the α, $\beta=1$–3,

$$X_\alpha(i)=X_\beta(i) \quad \text{modulo}\ n \quad (i=\mathrm{q},\mathrm{u}^c,\mathrm{d}^c)\ , \qquad (25)$$

which cannot ensure the vanishing of the relevant elements in (24). Therefore, Fritzsch-type mass matrices are not obtained in this case.

What kind of mass matrices are obtainable due to this sort of symmetry? We here present an example of Z_5 ($n=5$):

$$X_\alpha(\mathrm{q})=(3,2,2),\quad X_\alpha(\mathrm{u}^c)=(2,3,3)\ ,$$

$$X_\alpha(\mathrm{d}^c)=(3,2,3),\quad X(\mathrm{H})=X(\mathrm{H}^c)=0\ . \qquad (26)$$

This results in the following quark mass matrices:

$$M_\mathrm{u}\sim\begin{pmatrix} * & 0 & 0\\ 0 & * & *\\ 0 & * & *\end{pmatrix},\quad M_\mathrm{d}\sim\begin{pmatrix} 0 & * & *\\ * & 0 & 0\\ 0 & * & *\end{pmatrix}. \qquad (27)$$

The Z_5 may or may not be embedded in U(1), depending on the other couplings $(\lambda_5)_{\alpha\beta\gamma}\mathrm{N}_\alpha\mathrm{H}^c{}_\beta\mathrm{H}_\gamma$, etc., such that $X'_\alpha(\mathrm{q})=(-2,2,2)$, $X'_\alpha(\mathrm{u}^c)=(2,-2,-2)$ and $X'_\alpha(\mathrm{d}^c)=(-2,2,-2)$.

(b) The Higgs multiplets may also transform in a generation-dependent way. Consider an example of U(1):

$$X_\alpha(i)=(1,2,-2) \quad (i=\mathrm{q},\mathrm{u}^c,\mathrm{d}^c)\ ,$$

$$X_\alpha(j)=(4,0,-3) \quad (j=\mathrm{H},\mathrm{H}^c)\ , \qquad (28)$$

which constrain the elements of the coupling matrices λ_1 and λ_2 so as to reproduce the Fritzsch-type mass matrices. We may assign the charges of N_α, for example, as

$$X_\alpha(\mathrm{N})=(-1,3,-4)\ , \qquad (29)$$

so that some couplings among $(\lambda_5)_{\alpha\beta\gamma}N_\alpha\mathrm{H}^c{}_\beta\mathrm{H}_\gamma$ are also invariant under U(1). Furthermore, if $(\lambda_5)_{333}$ is nonzero as well as $(\lambda_5)_{312}$ and $(\lambda_5)_{321}$, U(1) is reduced to Z_{10}. This case with Z_{10} could be acceptable if the domain wall problem is solved by inflation [18], while if $(\lambda_5)_{333}=0$, it is ruled out by particle experiments such as $\mu\to e+G$ where G is the Nambu–Goldstone boson (or an axion) associated with U(1) [15].

In conclusion, we have investigated the phenomenological roles of possible global symmetries in a class of low-energy superstring inspired supersymmetric models. Such global symmetries may be present in the full string theory or may incidentally appear in the effective theory of the low-energy sector (except for explicit breakings $\propto M_{\mathrm{Pl}}^{-n}$). Phenomenologically acceptable models are singled out in accordance with the baryon number assignement to prevent rapid proton decay. If a model has a discrete symmetry, Z_n, which acts in a generation-independent way, it is naturally embedded in a U(1) symmetry. Then, such global U(1) symmetries (through suitable redefinition in terms of the local U(1)'s) remain unbroken even after the spontaneous breakdown of the gauge symmetries at $O(M_W)$. On the other hand, generation-dependent symmetries may reproduce interesting quark and lepton mass matrices. They will, however, generally suffer from the phenomenological difficulties associated with their spontaneous breakdown.

We would like to thank Q. Shafi for informing us about their articles, ref. [9], which dealt with some topics related to ours.

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