Simple Discrete Symmetries in Phenomenologically Viable E_6 Superstring Models

Manuel Drees and Xerxes Tata
Physics Department, University of Wisconsin, Madison, Wisconsin 53706
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We analyze the simplest discrete symmetries of the superpotential of superstring-inspired E_6 models necessary to suppress simultaneously (i) proton decay, (ii) neutrino masses, and (iii) tree-level flavor-changing neutral currents in E_6 superstring models without any intermediate scale. We show that there are two and only two different Z_2 symmetries that incorporate (i)–(iii). These symmetries put restrictions on the decays of exotic particles; in particular, the exotic quarks always decay semileptonically. A Z_3-symmetric model satisfying (i)–(iii) and with $\mu \rightarrow e\gamma$ forbidden is also constructed.

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Ever since the seminal papers of Candelas *et al.*[1] and Witten,[2] there has been a considerable amount of effort devoted to the study of the phenomenology[3–6] of the effective low-energy theory arising from the anomaly-free[7] superstring theory after compactification from ten to four dimensions. This is generally believed to be a supergravity theory based on the gauge group E_6 which is broken by the Hosotani mechanism[8] to a low-energy gauge group G with rank ≥ 5. The matter and Higgs fields are all contained in N_g **27** (and possibly additional $\mathbf{27}+\mathbf{27}^*$) representations of E_6. It has been shown[2] that the resulting low-energy superpotential is not the most general cubic polynomial invariant under G; rather, only those cubic terms present in the (originally) E_6-invariant $\mathbf{27}\times\mathbf{27}\times\mathbf{27}$ interactions are present. The relative couplings of these terms are arbitrary and not related by E_6 so that the most general superpotential can be written as[5,6]

$$f = h_L H L e^c + h_d H Q d^c + h_u \bar{H} Q u^c + h_\nu \bar{H} L \nu^c + h'_\nu D d^c \nu^c + \lambda_L L Q \bar{D} + \lambda'_L e^c u^c D + \lambda_B QQD + \lambda'_B u^c d^c \bar{D} + \alpha N H \bar{H} + \beta N D \bar{D}. \quad (1)$$

In Eq. (1), a sum over all the (N_g) generations is implicit. L and Q denote the usual lepton and quark doublets and e^c, d^c, and u^c the respective $SU(2)_L$ singlets. H and $\bar{H}$ denote the additional SU(2) doublets present in E_6, with negative and positive weak hypercharge, respectively. D and $\bar{D}$ denote the additional weak isosinglet quarks. Finally, ν^c and N are both $SU(3)\otimes SU(2)_L \otimes U(1)_Y$ singlets, with N being the singlet under the decomposition of the **27** under SO(10).

In addition to the **27** N_g superfields present in (1), it is possible that additional survivor superfields transforming as $\mathbf{27}\oplus\mathbf{27}^*$ of E_6 may also be present. This is not a generic feature, since, in general, the Hosotani mechanism renders all these fields superheavy. It is nevertheless possible for particular orientations of the adjoint vacuum expectation value (VEV) to get some of the survivor fields to be the Higgs fields that could break electroweak symmetry while ensuring that their color-triplet counterparts are heavy.[2,4] However, it has been shown by Kalyniak and Sunderesan[9] that (for rank-6 models) it is impossible to break the group G down to SU(3) $\otimes U(1)_{em}$ by use of just these fields. In other words, at least part of the electroweak breaking occurs via the vacuum expectation values of the fields present in the superpotential (1). A similar situation occurs for rank-5 models.[5] Since the masses can all be generated by the fields already present in the generational **27**'s of E_6, and since it requires a special choice of orientation of the adjoint E_6 VEV to get light survivors, we will ignore them from this point on.

It is clear that not all terms in (1) can be simultaneously present without phenomenological disasters. For example, if both λ_L- and λ_B-type terms are present in the superpotential, the D and $\bar{D}$ fields mediate unacceptably rapid proton decay. In order to forbid the appearance of such terms, several authors[2,3,5,6,10] have considered the imposition of discrete symmetries on the low-energy superpotential. Of course, in principle, this cannot be imposed from outside but should result from the structure of the compact manifold. Nevertheless, a study of the nature of discrete symmetries that would lead to phenomenologically acceptable models is of interest. The derivation of these from the underlying theory is an important but unrelated question.

In this paper, we attempt to obtain allowable discrete symmetries from the following requirements: (i) Proton decay is not too rapid; (ii) the neutrino is (essentially) massless; and (iii) flavor-changing neutral currents are *naturally* suppressed. [Since there are many neutral fields coupling to quarks and leptons in each generation, (iii) is a nontrivial requirement.] In addition, we require that (i)–(iii) be accomplished by one single symmetry. For simplicity, we will take this to be a two-valued symmetry so that the fields are all either odd or even under the symmetry operation. It is, of course, not *a priori* obvious that such a simple discrete symmetry satisfying the requirements (i)–(iii) even exists. Also, we will focus our attention on models without any intermediate scale. This is not to say that models with an intermediate scale $\simeq 10^{10}$ GeV are phenomenologically inconsistent. But it is fair to say that these models, because of their flexibility,[6] are certainly less predictive than their no-intermediate-scale counterparts. A critique of these models is given by Ellis *et al.*[11] Since, in this paper, we are interested in demonstrating the existence of simple phenomenologically consistent models with one discrete symmetry rather than analyzing all models, we will assume that there is no new physics between the weak and unification scales. In this case, the low-energy gauge group is uniquely[5] determined to be of the form

$$G = SU(3)_c \otimes SU(2)_L \otimes U(1) \otimes \tilde{U}(1).$$

We proceed by first recognizing that the superpotential (1), aside from being invariant under the group G, is also invariant under the additional global $U(1)_G$ contained in E_6 but not in G. This follows since (1) contains only those terms that were present in the E_6-invariant model. The charges of the various members of the **27** under $U(1)_G$ are $L = d^c = \frac{3}{4}$, $Q = e^c = u^c = -\frac{1}{12}$, $\nu^c = -\frac{11}{12}$, $H = \bar{D} = -\frac{2}{3}$, $\bar{H} = D = \frac{1}{6}$, and $N = \frac{1}{2}$. The invariance of the superpotential under global $G \otimes U(1)_G$ transformations then implies that three of the fields can be taken to be positive under our discrete symmetry without any loss of generality. For example, we can first fix u^c by rotating by a $U(1)_G$ transformation $u^c \to -u^c$, then d^c by rotating via a $\tilde{U}(1)$ transformation so that $d^c \to -d^c$ [note that this leaves u^c which has twice the charge of d^c under $\tilde{U}(1)$ unchanged] and finally we can fix the doublet Q by a pure $SU(2)_L$ transformation that takes all SU(2) doublets to their negatives.[12] Note that we have chosen to rotate the fields with the smallest nontrivial charges under the various groups. This ensures that the remaining fields go into either themselves or their negatives (rather than acquiring complex phases). For any two-valued discrete symmetry transformation S, we can thus without loss of generality choose u^c, d^c, and Q to transform into themselves.[13]

We now impose the requirements (i)–(iii) discussed earlier and examine the resulting consequences. The natural suppression of flavor-changing neutral currents requires that there exist a basis[14] in which at most one set of the fields $H, \bar{H}$ couples to quarks and leptons. Moreover, we know that in order to get quark masses, at least one set must couple to the quarks (these must acquire a VEV) which we can take to be H_3 and $\bar{H}_3$. Thus, H_3 and $\bar{H}_3$ must be even under S. Since the couplings of $H_{1,2}$ and $\bar{H}_{1,2}$ are forbidden, these must be odd under S. We now consider the fields N. We can always rotate to a basis where $\langle N_{1,2}\rangle = 0$, $\langle N_3\rangle \neq 0$. Since we require an $N_3 H_3 \bar{H}_3$ term in the superpotential to drive electroweak breaking,[5] N_3 must be even, whereas the requirement $\langle N_{1,2}\rangle = 0$ implies that there should be no $H_3 \bar{H}_3 N_{1,2}$ terms in the superpotential, since these lead to a VEV for $N_{1,2}$ unless there is an *unnatural* cancellation of terms. Thus $N_{1,2}$ should be odd under S. We remark also that if both N_1 and N_2 are even so that $\bar{H}_i N_{1,2}$ and $H_i N_{1,2}$ mass terms are absent, there would be a zero-mass state in the 12×12 neutral-gauge-Higgs-fermion sector.[15]

We now turn to the properties of the leptons and exotic quarks under the discrete symmetry transformation. We first note that, irrespective of the transformation properties of L and ν^c under S, one of the couplings $\bar{H}_3 L \nu^3$ or $\bar{H}_{1,2} L \nu^c$ must exist. (Recall that we assume that there is one and only one discrete symmetry to forbid couplings.) The former would lead to a Dirac mass $\simeq M_W$ for the neutrino and hence must be forbidden. Thus $\langle \bar{H}_{1,2}\rangle$ must vanish to forbid neutrino masses. Since $N_3 \bar{H}_i H_j$ $(i,j=1,2)$ terms are allowed, $\langle H_{1,2}\rangle = 0$ since, without fine tuning, a VEV for any two fields leads to a VEV for the third one. We thus conclude that $\langle H_3\rangle$ must also be responsible for lepton masses so that the product Le^c is even under S. Our analysis naturally divides into two cases, (A) and (B).

In case (A), L and e^c are both even which implies that ν^c is odd (to forbid Dirac neutrino masses). Also, we require an $N_3 D\bar{D}$ term to be present to give D and $\bar{D}$ quark masses so that D and $\bar{D}$ are either both even or both odd under S. The former case is ruled out because both λ_L- and λ_B-type terms in (1) are allowed which leads to rapid proton decay. If D and $\bar{D}$ are both odd, none of the terms of the λ_L or λ_B type is allowed. The decay of the exotic quarks takes place only via the $h'_\nu D d^c \nu^c$ term (note that D and $\bar{D}$ are mixed so that the mass eigenstates decay via this interaction), and so leads to a specific signature for their decay. We will return to this later.

We now consider case (B). L and e^c are both odd so that ν_L^c is even under S. Once again, we can choose D and $\bar{D}$ both even or both odd. The former allows λ_B-type terms and the h'_ν term in the superpotential and thus leads to a very short-lived proton, since ν^c is massless. The latter choice allows both λ_L-type terms but forbids h'_ν terms. We see that the exotic quarks decay only via λ_L couplings which again leads to a definite signature.

Our search for a discrete symmetry satisfying (i)–(iii) in a natural fashion has therefore yielded just two solutions. In the first cases, (A), all the usual quarks and leptons which are in the **10** and **5*** of SU(5) and the Higgs particles are even while the remaining fields are odd under S. We may, therefore, regard S an "exoticness parity." Note that S does not treat all the generations the same way, since it singles out the Higgs particles. Solution (B) differs from (A) only in the reversing of the S parity to the leptons. We remark that our starting point, viz., the assignment of $S=+$ to Q, d^c, and u^c, was just a convention. Starting with a different choice such as H_3 odd, and u^c and d^c even under S, leads to the same two solutions for the S parities up to a global

$G \otimes U(1)_G$ transformation.

We now turn to a brief discussion of the phenomenology of these models. Gauge invariance requires that ν_L^c is massless. Also, because of the S symmetry, terms of the type $N_3 H_3 \bar{H}_i$ ($i=1,2$) with two fields acquiring a VEV are forbidden. In this case the charged as well as the neutral gauge-Higgs-fermion and -boson matrices split into two decoupled sectors,[15] one containing the gauge and Higgs particles $(H_3, \bar{H}_3, N_3)$ and the other the particles $H_i \bar{H}_i$ and N_i, $i=1,2$. In this case, it has recently been shown[16] that the latter sector always contains a neutral fermion, $\tilde{n}$, with mass below about 115 GeV, and which in many cases may even be lighter. In fact, $\tilde{n}$ may well be the lightest supersymmetric particle. We note also that the ordinary and exotic charged $-\frac{1}{3}$ quarks do not mix and that neutrinos do not acquire masses radiatively as a consequence of the discrete symmetry.

The most copiously produced exotic particles, at least at hadron colliders, are likely to be the isosinglet quarks and scalar quarks. The mass eigenstate of the D-$\bar{D}$ quark system D_M (which has odd R parity) could decay via several two-body modes depending on their masses. If one of the supersymmetric partners $\tilde{D}_1$ of D_M is lighter than D_M, the decays $D_M \to \tilde{D}_1 +$gaugino which proceed by gauge interactions are likely to dominate the decays $D_M \to \tilde{D}_1 + \tilde{n}$. If this is not the case, for model (B) the decays $D_M \to d\tilde{\nu}_L, u\tilde{e}$ or $\tilde{d}\nu, \tilde{u}e$ (unlikely) or [for model (A)] the decay $D_M \to d\tilde{\nu}^c$ or $\tilde{d}\nu^c$ (unlikely) may be allowed. In the case where $\tilde{D}_1$ and all the usual scalar quarks and scalar leptons are heavier than D_M, the three-body decays $D_M \to$ neutralino (chargino) + quark + lepton, mediated by virtual $\tilde{D}$, scalar quarks, and scalar leptons, dominate. If $\tilde{D}_1$ is lighter than D_M, it will decay via $\tilde{D}_1 \to d\nu^c$ [model (A)] or $d\nu$ or du [model (B)]. Otherwise, the dominant decays are likely to be $\tilde{D}_1 \to D_M +$gaugino or $D_M + \tilde{n}$. We see that the isosinglet quarks always decay into lepton + quark. Their decays into charged leptons [for model (B)] are particularly characteristic since these hard leptons would provide clean signatures for these decays. Even in model (A) whenever the decays take place via gauge interactions in a substantial fraction of the events the next-to-lightest neutralino would be produced. This, in turn, would cascade down to the lightest neutralino; since the scalar-quark mass substantially exceeds the scalar-lepton mass in many models, the cascade could result in an enhanced branching[17] into leptons for this decay, thereby leading to a cleaner signature for the isosinglet quarks and scalar quarks. We also remark that, for model (B), there is the interesting possibility of the resonance production of $\tilde{D}_1$ at DESY HERA. The possibility of identifying isoscalar quark (scalar-quark) signals has been studied in detail by Angelopoulos *et al.*[18] In both cases (A) and (B), the S parity considered here rules out the possibility of the hard-to-identify diquark decay of D.

Finally we discuss the decays of the exotic particles contained in $H_{1,2}$, $\bar{H}_{1,2}$, and $N_{1,2}$. The neutral scalars can always decay via the $\nu\nu^c$ mode since the coupling h_ν to $H_{1,2}$ and $\bar{H}_{1,2}$ always exists. In this case, of course, they would escape detection. If the masses are in the right range, these may also decay into gauge bosons (fermions) and charged exotic bosons (fermions). The charged exotic scalars can always decay into $e^\pm + \nu^c$ (via h_ν terms) and possibly into $W^\pm +$neutral exotic scalars. The former decay would be very distinctive since it leads to very hard charged leptons in the final state. We should emphasize that in a generic model these exotic bosons would also decay into quarks and other lepton modes in much the same way as usual Higgs boson (except that their coupling is not proportional to the fermion mass). The restriction to the decay modes discussed above is a consequence of S-parity conservation. For the same reason, these can only be pair-produced by the Drell-Yan mechanism.

The exotic fermions can also be produced via the Drell-Yan process. The charged fermions can decay into the neutral exotic fermion $+W^\pm$ if the decay is kinematically allowed which it may well be since $\tilde{n}$ is light.[15] An alternative possibility for the case of relatively light scalar leptons is the $\tilde{\nu}^c e$ or $\nu^c \tilde{e}$ decay for the charged case and $\tilde{\nu}_L \nu^c$ or $\nu \tilde{\nu}_L^c$ for the neutral case. If none of these two-body modes are accessible, the fermion dominantly decays into three-body modes mediated by gauge or scalar leptons. Once again, it is a consequence of S-parity conservation that virtual scalar-quark decays are absent.

At this point, we note that both the models allow for exotic processes that are strongly constrained by experiment. For example, in model (B) the process $\mu + N \to e + N$ is generally mediated by $\tilde{D}$ and so leads to a bound $\simeq 10^{-3}$ on the λ_L and λ_L' couplings.[19] These are forbidden in model (A) which may thus be preferred. The decay $\mu \to e\gamma$, which proceeds via $\bar{H}_{1,2}$ loops, occurs in both models (h_ν terms), which leads to the bound[19] $h_\nu \lesssim 10^{-3}$. The problem of $\mu \to e$ transitions is not unique to this paper; this is present in most models in the literature. Since our discrete symmetry forbids h_L, h_d, and h_u couplings except to the Higgs multiplet, rare processes from kaon physics which would lead to the even stronger bounds[19] $\simeq 10^{-4}$–10^{-5} on these couplings are naturally forbidden.

Finally, we show that it is possible to construct a model with Z_3 as the discrete symmetry group in which the $\mu \to e$ decays discussed above are forbidden. A generic Z_3 transformation is $\psi \to \exp[(2\pi i/3)Q_\psi]\psi$, with $Q_\psi = 0$, 1,2. We choose $Q=0$ for all the usual quarks (this allows general flavor mixings), for the electron family and for the Higgs fields H_3, $\bar{H}_3$, and N_3. In order to forbid $\mu \to e$ transitions, we choose $Q=1$ for the μ and τ doublets and $Q=2$ for the respective singlets. Also, we choose $Q_{\nu^c}=1$ for all ν^c in order to forbid neutrino masses. In order to forbid tree-level flavor-changing neutral currents, we need $Q_{H_3} \neq Q_{H_i}$ and $Q_{\bar{H}_3} \neq Q_{\bar{H}_i}$. We take $Q_{H_1} = Q_{\bar{H}_2} = Q_{N_1} = 1$ and $Q_{\bar{H}_1} = Q_{H_2} = Q_{N_2} = 2$. This ensures that $H_3\bar{H}_iN_k$, $\bar{H}_3H_iN_k$, $N_3\bar{H}_iH_k$, and $H_3\bar{H}_3N_3$ terms all exist and that the unwanted terms are absent. We have checked that the 6×6 neutral exotic fermion matrix has no generic zero eigenvalues. Finally, we take $Q_D = 1$ and $Q_{\bar{D}} = 2$ which allows a mass for D fermions. λ_B- and λ_L-type terms are forbidden so that D decays only via the h'_ν term in (1). We see that our model incorporates the absence of $\mu \to e$ transitions naturally although $\tau \to \mu\gamma$ decays are allowed. The experimental bound on the $B(\tau \to \mu\gamma) \lesssim 5\times10^{-4}$ leads to a bound on the coupling of about a few times 10^{-2} and so requires no unnaturally small parameters.

To summarize, we have shown that for superstring models without any intermediate scale, if we require a single Z_2 symmetry, S, to suppress (i) proton decay, (ii) neutrino masses, and (iii) tree-level flavor-changing neutral currents, only two solutions are possible. Both these solutions require discrete symmetry operators that distinguish between generations. The first solution requires the introduction of a multiplicatively conserved "exoticness parity" under which the usual quark, lepton, and Higgs superfields are even with the remaining fields in the **27** of E_6 being odd. The second solution interchanges the S parity of the leptons including ν_L^c. In both models the exotic quarks or scalar quarks which will be copiously produced at the Fermilab Tevatron or the proposed Superconducting Super Collider are predicted to decay via relatively clean leptonic modes. Finally, a Z_3 model in which (i)–(iii) are satisfied and $\mu \to e\gamma$ is forbidden has been constructed.

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[12]The discrete symmetries distinguish between generations since the Higgs doublets (which are in one of the **27**'s) are distinguished from the other doublets with the same quantum numbers. The quarks of the different generations are, however, treated as copies even under this symmetry in order to allow for complete flavor mixing as required by the Kobayashi-Maskawa matrix. For the leptons, this is not the case. We have, however, checked that if we require all $\bar{H}_3L_i\nu_j^c$ terms to be absent to prevent neutrino masses, it is not possible to have a Z_2 symmetry that distinguishes lepton families. It is possible to have a Z_3 symmetry which does so as is evident from the example given in the text.

[13]Note that we cannot fix yet another field by a $U(1)_Y$ transformation. The field with the smallest weak hypercharge is the doublet Q. A hypercharge transformation that takes $Q \to -Q$ does so for all doublets and so we cannot fix two doublets by independent SU(2) and $U(1)_Y$ transformations. Nothing can be fixed by SU(3) since -1 is not an SU(3) group element.

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