R-parity violation in
superstring derived models

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Abstract

The ALEPH collaboration has recently reported a significant excess of four-jet events at the LEP 1.5 and LEP 2 experiments. While not yet confirmed by the other collaborations, it was recently proposed that this excess may be explained by certain R-parity violating operators in supersymmetric models. R-parity violating operators introduce the danger of inducing rapid proton decay. I discuss how the operators required to explain the ALEPH four jet events may arise from superstring derived models without inducing rapid proton decay.

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Particle physics has for long awaited the experimental deviation from the Standard Model predictions that will guide the way to the physics beyond the Standard Model. Such a discovery may be just around the corner. Recently, the ALEPH collaboration has reported an excess in the four–jet event cross section which is several sigmas above the Standard Model prediction [1, 2]. Perhaps even more intriguing is the fact that both at LEP 1.5 and LEP 2 runs ALEPH has observed a sharp peak at 106.1 ± 0.8 GeV, corresponding to 18 events with 3.1 expected from QCD background. The di–jet mass difference distribution of the selected 18 events is consistent with a value around 10 GeV. If interpreted as a particle pair production this together with the information on the di–jet mass sum suggests that the two particle produced have masses of about 58 and 48 GeV and production of same mass particles is disfavored. By extracting information on the primary parton [3], it is concluded that the pair produced particles have a sizable charge and that neutral particle production is disfavored [1, 2]. Absence of b–quarks in the final states disfavors the hypothesis of Higgs–boson production.

Recently, it was proposed [4] that the ALEPH excess of four jet events can be explained in supersymmetric models with R–parity violation [5]. According to this proposal left–handed and right–handed selectrons $\tilde{e}_L \tilde{e}_R$ are pair produced at LEP. The selectron pair then decay further by the R–parity violating operator

$$\lambda_{ijk} L_i Q^j d^k$$

where the standard notation for lepton and quark superfields has been used and $i, j, k$ denote the generation indices, and absence of top or bottom quarks in the final states restricts only the $\lambda_{1jk}$ with $j, k = 1, 2$ to be nonzero [4].

It is well known however that R–parity violation may induce rapid proton decay. If in addition to the operator in Eq. (1) one also has the operator

$$\eta_{ijk} u^i d^j d^k$$

with unsuppressed $\eta_{ijk}$ couplings then the proton decays rapidly. A quick estimate of the constraint from proton lifetime gives

$$g^2 \left( \frac{M_{\text{quark}}}{M_{\text{GUT}}} \right)^2$$

where

$$\lambda \eta < g^2 \left( \frac{M_{\text{quark}}}{M_{\text{GUT}}} \right)^2$$
Therefore, if $\lambda_{1jk} \sim 10^{-4}$ as proposed in ref. [4], and taking $M_{\text{squark}} \sim 1$ TeV $M_{\text{GUT}} \sim 10^{16}$ GeV, naively requires $\eta < 10^{-22}$. Thus, proton constraints essentially requires $\eta \equiv 0$ if the R–parity violation is to explain the excess of ALEPH four jet events.

Therefore, the problem is to understand why the couplings in Eq. (1) are allowed while the couplings in Eq. (2) are forbidden. In this paper I discuss this problem in the context of realistic superstring derived models.

To study this problem I examine the superstring models which are constructed in the free fermionic formulation [3, 7, 8, 9, 10, 11, 12, 13]. This class of superstring models reproduce many of the properties of the Standard Model, like three chiral generations with the Standard Model gauge group and existence of Higgs doublets which can generate realistic fermion mass spectrum. Two of the important features in this class of models is the existence of a stringy doublet–triplet splitting mechanism which resolves the GUT hierarchy problem [14] and the fact that the chiral generations all fall into the 16 of $SO(10)$. This last property admits the standard embedding of the weak hypercharge in $SO(10)$ and is crucial for the agreement of these models with $\sin^2 \theta_W(M_Z)$ and $\alpha_{\text{strong}}(M_Z)$ [12, 16].

The superstring models under consideration are constructed in two steps. In the first step the observable gauge symmetry is broken to $SO(10) \times SO(6)^3$. There are 48 generations in the chiral 16 representation of $SO(10)$ with $N = 1$ space–time supersymmetry. In the second step the $SO(10)$ symmetry is broken to one of its subgroups, $SU(5) \times U(1)$, $SO(6) \times SO(4)$ or $SU(3) \times SU(2) \times U(1)^2$. The flavor $SO(6)^3$ symmetries are broken to $U(1)^n$, where $n$ may vary between 3–9, and the number of generations is reduced to three. The symmetry is then broken further in the effective field theory and the weak hypercharge is some linear combination of the Cartan subgenerators. For example, in the standard–like models the weak hypercharge is given by,

$$U(1)_Y = \frac{1}{3} U(1)_C + \frac{1}{2} U(1)_L$$  \hspace{1cm} (4)$$

The chiral generations in these superstring models are obtained from the 16 multiplets of $SO(10)$ and carry charges under the flavor symmetries. These models
typically contain an “anomalous” $U(1)$ symmetry which requires that some fields in the massless string spectrum obtain non–vanishing VEVs \cite{17}. Further details on the construction of the realistic free fermionic models are given in ref. \cite{10}.

In general in string models one expects the appearance of R–parity violating terms of the form of Eq. (1) and (2). If both are not suppressed then the proton decays much too fast. If the $B – L$ generator is gauged like in $SO(10)$ then these terms are forbidden at the cubic level by gauge invariance. However, they may still be generated from nonrenormalizable terms that contain the right–handed neutrino.

$$\eta_1(uddN)\Phi + \eta_2(QLdN)\Phi.$$  \hspace{1cm} (5)

where $\Phi$ is a combination of fields that fixes the string selection rules and gets a VEV of $O(M_{Pl})$ and $N$ is the Standard Model singlet in the 16 of $SO(10)$. Thus, the ratio $\langle N \rangle/M_{Pl}$ controls the rate of proton decay. In general, terms of the form of Eq. (5) are expected to appear in string models at different orders of nonrenormalizable terms. For example, in the model of ref. \cite{11} such terms appear at order $N = 6$

$$
(u_3d_3 + Q_3L_3)d_2N_2\Phi_{45}\bar{\Phi}_2^- \\
+ (u_3d_3 + Q_3L_3)d_1N_1\Phi_{45}\bar{\Phi}_1^+ \\
+ u_3d_2d_2N_3\Phi_{45}\Phi_2^- + u_3d_1d_1N_3\Phi_{45}\Phi_1^+ \\
+ Q_3L_1d_3N_1\Phi_{45}\Phi_3^+ + Q_3L_1d_1N_3\Phi_{45}\Phi_3^+ \\
+ Q_3L_2d_3N_2\Phi_{45}\bar{\Phi}_3^- + Q_3L_2d_2N_3\Phi_{45}\bar{\Phi}_3^-.
$$

(6)

In this model the states from the sector $b_3$ are identified with the lightest generation. It is therefore seen that if any of $N_1$, $N_2$ or $N_3$ gets a Planck scale VEV, dimension four operators may be induced which would result in rapid proton decay. It is interesting to note that all the terms in Eq. (6) contain the field $\Phi_{45}$. If the VEV of $\Phi_{45}$ vanishes then all the higher order terms are identically zero. In this specific model due to the anomalous $U(1)$ symmetry $\Phi_{45}$ must get a VEV and, in general, dimension four operators may be induced. Nevertheless, this observation suggests the possibility that slight variation of the model will result in a field appearing in these terms which is not required to get a VEV. However, even if such a possibility can work we see
that both the desired terms of the form $QLd$ and the undesired terms of the form $udd$ are induced, or forbidden, simultaneously.

In the flipped $SU(5)$ model similar terms may arise from the terms

$$FF\bar{f}H\Phi^n.$$ (7)

Here $F$ and $H$ are in the $(10,1/2)$ representation and $\bar{f}$ is in the $(\bar{5},-3/2)$ representation of $SU(5) \times U(1)$. The field $F$ contains the $Q$, $d$, $N$ fields and $\bar{f}$ contains the $u$ and $L$ fields. The Standard Model singlet, $N$ in the Higgs field $H$ obtains a VEV which breaks the $SU(5) \times U(1)$ symmetry to the Standard Model symmetry. Thus, terms of the form of Eq. (7) produce simultaneously the terms in Eq. (1) and Eq. (2). Terms of the form of Eq. (7) are in general found in the string models [18]. Therefore, to produce only the terms of the form of Eq. (1) while preventing the terms in Eq. (2) requires a different mechanism.

In the case of the $SO(6) \times SO(4)$ superstring models the Standard Model fermions are embedded in the

$$F_L \equiv (4,2,1) = Q + L$$

$$\bar{F}_R \equiv (\bar{4},1,2) = u + d + e + N$$ (8)

representations of the $SU(4) \times SU(2)_L \times SU(2)_R$. Note that $F_L + \bar{F}_R$ make up the 16 spinorial representation of $SO(10)$. The dangerous dimension four operators are obtained in this case from the operator

$$F_L F_L \bar{F}_R \bar{H}_R \quad \text{and} \quad \bar{F}_R \bar{F}_R \bar{F}_R \bar{H}_R$$ (9)

where $\bar{H}_R$ is the Higgs representation which breaks the extended non–Abelian symmetry. We observe that in the $SO(6) \times SO(4)$ type models, like the $SU(3) \times SU(2) \times U(1)^2$ type models, the operator in Eqs. (1) and (2) arise from two distinct operators.

Next, I turn to the model of Ref. [13]. The detailed spectrum of this model and the quantum numbers are given in Ref. [13]. In this model the observable gauge group formed by the gauge bosons from the Neveu–Schwarz sector alone is

$$SU(3)_C \times SU(2)_L \times U(1)_C \times U(1)_L \times U(1)_{1,2,3,4,5,6}$$ (10)
However, in this model two additional gauge bosons appear from the twisted sector \( 1 + \alpha + 2\gamma \). These new gauge bosons are singlets of the non–Abelian gauge group but carry \( U(1) \) charges. Referring to this generators as \( T^\pm \), then together with the linear combination

\[
T^3 \equiv \frac{1}{4} \left[ U(1)_C + U(1)_4 + U(1)_5 + U(1)_6 + U(1)_7 - U(1)_9 \right]
\]

the three generators \( \{T^3, T^\pm\} \) together form an enhanced \( SU(2)_{\text{custodial}} \) symmetry group. Thus, the original observable symmetry group is enhanced to

\[
SU(3)_C \times SU(2)_L \times SU(2)_{\text{cust}} \times U(1)_C' \times U(1)_L \times U(1)_{1,2,3} \times U(1)_{4',5',7''}
\]

The different combinations of the \( U(1) \) generators are given in ref. [13, 16]. The weak hypercharge is still defined as a combination of \( U(1)_C \) and \( U(1)_L \). However in the present model \( U(1)_C \) is part of the extended \( SU(2)_{\text{custodial}} \) symmetry. We can express \( U(1)_C \) in terms of the new orthogonal \( U(1) \) combinations,

\[
\frac{1}{3} U(1)_C = \frac{2}{5} \left\{ U(1)_C' + \frac{5}{16} \left[ T^3 + \frac{3}{5} U_7'' \right] \right\}.
\]

and the weak hypercharge is given as before by the linear combination

\[
U(1)_Y = \frac{1}{3} U(1)_C + \frac{1}{2} U(1)_L
\]

The weak hypercharge depends on the diagonal generator of the custodial \( SU(2) \) gauge group. We can therefore instead define the new linear combination with this term removed,

\[
U(1)_{Y'} \equiv U(1)_Y - \frac{1}{8} T^3
\]

\[
= \frac{1}{2} U(1)_L + \frac{5}{24} U(1)_C
\]

\[
- \frac{1}{8} \left[ U(1)_4 + U(1)_5 + U(1)_6 + U(1)_7 - U(1)_9 \right],
\]

so that the weak hypercharge is expressed in terms of \( U(1)_{Y'} \) as

\[
U(1)_Y = U(1)_{Y'} + \frac{1}{2} T^3 \quad \implies \quad Q_{\text{e.m.}} = T^3_L + Y = T^3_L + Y' + \frac{1}{2} T^3_{\text{cust}}.
\]
The final observable gauge group then takes the form

$$SU(3)_C \times SU(2)_L \times SU(2)_{\text{cust}} \times U(1)_{Y'} \times \left\{ \text{seven other } U(1) \text{ factors} \right\}. \quad (17)$$

These remaining seven $U(1)$ factors must be chosen as linear combinations of the previous $U(1)$ factors so as to be orthogonal to the each of the other factors in (17).

The full massless spectrum of this model is given in Ref. [13]. In this model the charged and neutral leptons transform as doublets of the $SU(2)_{\text{custodial}}$ symmetry while the quarks are singlets. Therefore, because of the custodial $SU(2)$ symmetry the terms of the form

$$QLdN \quad (18)$$

are invariant under the custodial $SU(2)$ symmetry, while the terms of the form

$$uddN \quad (19)$$

are not invariant. We could contemplate tagging another $N$ field to Eq. (19) which will render it invariant under $SU(2)_{\text{custodial}}$. However, this will spoil the invariance under $U(1)_L$. We therefore find that the baryon number violating operators, Eq. (19) vanish to all orders in the model of ref. [13]. Therefore, this model admits the type of custodial symmetries which allow the R–parity lepton–number violating operators of Eq. (1) while they forbid the baryon–number violating operators of Eq. (2). This conclusion was verified by an explicit search of nonrenormalizable terms up to order $N = 10$. On the other hand we find already at order $N = 6$ the non–vanishing terms

$$Q_1d_3L_3N_1\Phi_{45}\Phi_1 + Q_1d_2L_2N_2\Phi_{45}\Phi_{45}$$
$$Q_1d_1L_3N_3\Phi_{45}\Phi_1 + Q_1d_1L_2N_2\Phi_{45}\Phi_{45}$$
$$Q_2d_3L_3N_2\Phi_{45}\Phi_2 + Q_2d_2L_1N_1\Phi_{45}\Phi_{45}$$
$$Q_2d_2L_3N_3\Phi_{45}\Phi_2 + Q_2d_1L_1N_2\Phi_{45}\Phi_{45} \quad (20)$$

At higher orders additional terms will appear. It is therefore seen that while the R–parity baryon number violating operators are forbidden to all orders of nonrenormalizable terms the lepton number violating operators are allowed. This is precisely
what is required if the R–parity violation interpretation of the excess of four jet events observed by the ALEPH collaboration is correct.

Let us note some further remarks with in regard to the model proposed in Ref. [4]. As claimed there the R–parity interpretation prefers low values of tan β and therefore to allow perturbative unification requires some intermediate thresholds. This is precisely the scenario suggested by the class of superstring standard–like models [19]. In this class of models the top–bottom quarks mass hierarchy arises due to the fact that only the top quark gets its mass from a cubic level term in the superpotential while the bottom quark gets its mass term from a higher order term. Thus, in this class of models the top–bottom mass splitting arises due to a hierarchy of the Yukawa couplings rather than a large value of tan β. It has similarly been proposed in the context of this class of superstring models that intermediate matter thresholds are required for resolution of the string scale gauge coupling unification problem [12, 16].

To conclude, it was shown in this paper that string models can give rise to dimension four R–parity lepton number violating operators while forbidding the baryon number violating operators. Thus, R–parity violation is allowed while proton decay is forbidden. It will be of further interest to examine whether similar mechanism can operate in other string models [20, 21]. For example, the $SO(6) \times SO(4)$ type models are of particular interest as they also can in principle differentiate between the lepton–number and baryon–number violating operators. It is of further interest to study whether the string models can actually give sizable R–parity violation which is not in conflict with any observation. Finally, we eagerly await the experimental resolution of the observed excess in the ALEPH four jet events.

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| $F$ | SEC     | $SU(3)_C \times SU(2)_L \times SU(2)_c$ | $Q_C'$ | $Q_L$ | $Q_1$ | $Q_2$ | $Q_3$ | $Q_4'$ | $Q_5'$ | $SU(5)_H \times SU(3)_H$ | $Q_{v'}$ | $Q_v$ |
|-----|---------|-------------------------------------|--------|-------|-------|-------|-------|-------|-------|--------------------------|----------|-------|
| $L_1$ | $b_1^\oplus$ | (1, 2, 2) | $-\frac{1}{2}$ | 0 | $\frac{1}{2}$ | 0 | 0 | $-\frac{1}{2}$ | $-\frac{1}{2}$ | (1, 1) | $-\frac{2}{3}$ | |
| $Q_1$ | $1 + \alpha + 2\gamma$ | (3, 2, 1) | $\frac{1}{6}$ | 0 | $\frac{1}{2}$ | 0 | 0 | $-\frac{1}{2}$ | $-\frac{1}{2}$ | (1, 1) | $\frac{1}{3}$ | |
| $d_1$ | | (3, 1, 1) | $-\frac{1}{6}$ | $-1$ | $\frac{1}{2}$ | 0 | 0 | $\frac{1}{2}$ | $\frac{1}{2}$ | (1, 1) | $-\frac{4}{3}$ | |
| $N_1$ | | (1, 1, 2) | $\frac{1}{2}$ | $-1$ | $\frac{1}{2}$ | 0 | 0 | $\frac{1}{2}$ | $\frac{1}{2}$ | (1, 1) | $\frac{4}{3}$ | |
| $e_1$ | | (1, 1, 2) | $\frac{1}{2}$ | 1 | $\frac{1}{2}$ | 0 | 0 | $\frac{1}{2}$ | $\frac{1}{2}$ | (1, 1) | $\frac{1}{3}$ | |
| $u_1$ | | (3, 1, 1) | $-\frac{1}{6}$ | $-1$ | $\frac{1}{2}$ | 0 | 0 | $\frac{1}{2}$ | $\frac{1}{2}$ | (1, 1) | $-\frac{4}{3}$ | |
| $L_2$ | $b_2^\oplus$ | (1, 2, 2) | $-\frac{1}{2}$ | 0 | 0 | $\frac{1}{2}$ | 0 | $\frac{1}{2}$ | $-\frac{1}{2}$ | (1, 1) | $-\frac{2}{3}$ | |
| $Q_2$ | | (3, 2, 1) | $\frac{1}{6}$ | 0 | 0 | $\frac{1}{2}$ | 0 | $\frac{1}{2}$ | $-\frac{1}{2}$ | (1, 1) | $\frac{1}{3}$ | |
| $d_2$ | | (3, 1, 1) | $-\frac{1}{6}$ | 1 | 0 | $\frac{1}{2}$ | 0 | $-\frac{1}{2}$ | $\frac{1}{2}$ | (1, 1) | $-\frac{4}{3}$ | |
| $N_2$ | | (1, 1, 2) | $\frac{1}{2}$ | $-1$ | 0 | $\frac{1}{2}$ | 0 | $-\frac{1}{2}$ | $\frac{1}{2}$ | (1, 1) | $\frac{2}{3}$ | |
| $e_2$ | | (1, 1, 2) | $\frac{1}{2}$ | 1 | 0 | $\frac{1}{2}$ | 0 | $-\frac{1}{2}$ | $\frac{1}{2}$ | (1, 1) | $\frac{2}{3}$ | |
| $u_2$ | | (3, 1, 1) | $-\frac{1}{6}$ | $-1$ | 0 | $\frac{1}{2}$ | 0 | $-\frac{1}{2}$ | $\frac{1}{2}$ | (1, 1) | $-\frac{4}{3}$ | |
| $L_3$ | $b_3^\oplus$ | (1, 2, 2) | $-\frac{1}{2}$ | 0 | 0 | 0 | $\frac{1}{2}$ | 0 | 1 | (1, 1) | $-\frac{2}{3}$ | |
| $Q_3$ | | (3, 2, 1) | $\frac{1}{6}$ | 0 | 0 | 0 | $\frac{1}{2}$ | 0 | 1 | (1, 1) | $\frac{1}{3}$ | |
| $d_3$ | | (3, 1, 1) | $-\frac{1}{6}$ | 1 | 0 | 0 | $\frac{1}{2}$ | 0 | $-1$ | (1, 1) | $-\frac{4}{3}$ | |
| $N_3$ | | (1, 1, 2) | $\frac{1}{2}$ | $-1$ | 0 | 0 | $\frac{1}{2}$ | 0 | $-1$ | (1, 1) | $\frac{2}{3}$ | |
| $e_3$ | | (1, 1, 2) | $\frac{1}{2}$ | 1 | 0 | 0 | $\frac{1}{2}$ | 0 | $-1$ | (1, 1) | $\frac{2}{3}$ | |
| $u_3$ | | (3, 1, 1) | $-\frac{1}{6}$ | $-1$ | 0 | 0 | $\frac{1}{2}$ | 0 | $-1$ | (1, 1) | $-\frac{4}{3}$ | |

Table 1: Three generations of massless states and their quantum numbers in the model of Ref. [13].