5 September 1996

PHYSICS LETTERS B

ELSEVIER

Physics Letters B 383 (1996) 307–312

# R-parity violation and Peccei-Quinn symmetry in GUTs

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Received 4 March 1996; revised manuscript received 25 April 1996
Editor: L. Alvarez-Gaumé

**Abstract**

We address the question whether it is possible in GUTs to obtain R-parity violation with a large $\Delta L/\Delta B$ hierarchy of strengths so that the proton is stable while phenomenologically interesting $L$-violation is present. We consider versions of SU(5) with a built-in Peccei-Quinn symmetry spontaneously broken at an intermediate scale. The PQ symmetry and the field content guarantee a large suppression of the effective $B$-violating terms by a factor $\Lambda_{PQ}^3/M_P M_X^2$ while the effective $L$-violating terms stay large.

## 1. Introduction

A straightforward supersymmetrization of the Standard Model [1] allows the existence of low dimension operators ($D = 4,5$) that violate $B$- and $L$-number [2]. The $D = 4$ operators are usually avoided by imposing a discrete symmetry called R-parity [3] and defined as $R = (-1)^{3B+L+2S}$, $S$ being the spin. Similarly, the dangerous $D = 5$ operators are eliminated by imposing a suitable symmetry. If this symmetry is broken at some intermediate scale $\Lambda$, these operators will be suppressed by $\Lambda/M$, $M$ being a large mass scale. The Peccei-Quinn [4] symmetry proposed for the explanation of the vanishing vacuum angle theta is such a symmetry, suitable for the suppression of the $B$-violating $D = 5$ operators. Examples of GUTs incorporating a PQ symmetry have been constructed [5–7].

If R-parity is not a symmetry of the Standard Model, then the superpotential should include (directly or effectively) the terms

$$\lambda_{ijk} l_i l_j e_k^c + \lambda'_{ijk} d_i^c l_j q_k + \lambda''_{ijk} d_i^c d_j^c u_k^c + \epsilon_i l_i H . \qquad (1)$$

The indices are generation indices. The combination of the second and third term results in proton decay through squark exchange at an unacceptable rate unless $|\lambda'\lambda''| \leq 10^{-24}$. If one is restricted within the Standard Model it is possible, adopting a phenomenological attitude, to assume the existence of some of these couplings while forbidding the presence of others [8]. For example, setting $\lambda''_{ijk} = 0$ while keeping the rest leads to a number of $L$-violating phenomena. This is something that cannot be done in GUTs, at least in such a straightforward fashion. For instance, in SU(5) all terms in (1) can arise from

$$\lambda_{ijk}\phi_i(\bar{5})\phi_j(\bar{5})\psi_k(10) + \epsilon_i\phi_i(\bar{5})H(5) . \qquad (2)$$

In SU(5) all couplings in (1) are related by $\lambda''_{ijk} = \frac{1}{2}\lambda'_{ijk} = \lambda_{ijk}$ and should be present simultaneously. Then, if R-parity is not an exact symmetry, a large hierarchy in $B$-versus- $L$-violating strengths must be accounted for [9].

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*PII* S0370-2693(96)00758-7

Nevertheless, it is possible that these terms could be absent at the renormalizable level due to another symmetry, not directly related to R-parity, and show up as non-renormalizable effective interactions leading to small effective couplings suppressed by ratios of the breaking scale of this symmetry to some large mass scale. Note however that the required smallness of these couplings comes about almost exclusively from the need to suppress the $B$-violating interactions threatening the proton stability. $L$-violating couplings, if they were independent as in the Standard Model, would not be so severely constrained. A model of effective R-parity violation would be phenomenologically interesting if it were characterized by an effective large $B$-versus- $L$-violation disparity.

## 2. Peccei-Quinn symmetry in SU(5) and R-parity violation

A PQ-symmetric version of the minimal supersymmetric SU(5) model can be constructed in a straightforward fashion at the expense of introducing an extra pair of Higgs pentaplets and singlets [6]. The superpotential of the model is

$$W = h_{ij}\psi_i\psi_j H + f_{ij}\psi_i\phi_j\overline{H} + \overline{H}'(M' + \lambda'\Sigma)H + \overline{H}(M'' + \lambda''\Sigma)H' + f\overline{H}HP + f'\overline{H}'H'\overline{P} + (M/2)\mathrm{Tr}(\Sigma^2) + (\lambda/3)\mathrm{Tr}(\Sigma^3)\,. \quad (3)$$

The extra fields are the pentaplets $H'$, $\overline{H}'$and the SU(5)-singlets $P$, $\overline{P}$. The charges under $U(1)_{PQ}$ are $\psi(1)$, $\phi(1)$,$H(-2)$, $\overline{H}(-2)$,$H'(2)$, $\overline{H}'(2)$, $\Sigma(0)$, $P(4)$, $\overline{P}(-4)$. In order to generate the required PQ-breaking we need to add to (3) suitable additional interactions among the singlets. Couplings $hP\overline{P}X$ to another (neutral) singlet $X$ with a mass of $O(M_P)$,when $X$ is integrated out, lead to effective non-renormalizable terms $h^2(P\overline{P})^2/M$. Such a term would be sufficient to induce spontaneous breaking of the PQ symmetry [10], in conjunction with the standard soft supersymmetry breaking terms in the potential $m_0^2(|P|^2 + |\overline{P}|^2)$ and $m_0Ah^2(P\overline{P})^2/M+$ h.c. The scale of PQ breaking is $\langle P\rangle = \langle\overline{P}\rangle \equiv \mu = -m_0M(A/6h^2)(1 + \sqrt{1 - 12/A^2}) \simeq 10^{10}$–$10^{12}$ GeV. This range of values is compatible with astrophysical and cosmological bounds [11].

R-parity, although not explicitly imposed, is an exact symmetry of the model even after PQ spontaneous breaking. Although we cannot exclude that R-parity is indeed an exact symmetry it is more interesting to explore the possibility that there exist additional interactions that ultimately lead to effective R-violating couplings among standard fields such as $\phi_iH$, $\phi_iH'$ and $\psi_i\phi_j\phi_k$. For instance, singlets carrying odd PQ charge could couple to the above operators. A viable model however should predict also the necessary suppression of these effective couplings. It is possible to construct models in which the PQ charges of the fields guarantee that the R-violating operators will appear at the non-renormalizable level. As one of the possible classes of models that could be constructed, we shall consider a pair of singlets $S$, $\overline{S}$ carrying PQ charges $1/2$, $-1/2$. This choice of charges ensures the absence of renormalizable couplings to the other fields. Then, the R-violating term

$$\lambda_i\phi_iHS^2/M \quad (4)$$

is possible. A PQ-breaking v.e.v. for $S$, $\overline{S}$ breaks R-parity and generates an effective Higgs-matter mixing through this term. A v.e.v. $\overline{\mu} = \langle S\rangle = \langle\overline{S}\rangle \sim \frac{m_0M}{\overline{h}^2}$, of the same order as the $P$ and $Q$ v.e.v., can be generated through the presence of a term $\overline{h}(S\overline{S})^2/M$ in conjunction with soft supersymmetry breaking. All these terms can arise as effective interactions from couplings $\phi HY + SS\overline{Y} + \overline{S}Y\overline{S} + S\overline{S}Z$ to singlets $Z$,$Y$,$\overline{Y}$ having masses of the order of the Planck-mass, after they are integrated out. Higher-order R-violating terms

$$(\phi H')S^2\overline{P}/M^2 \quad (5)$$

and

$$(\psi\phi\phi)\overline{P}S^2/M^3 \quad (6)$$

are also present but, their suppression with extra powers of the Planck mass makes them irrelevant. Terms with $\Sigma$ insertions can also be written down but they are suppressed by powers of $M_X/M$.

## 3. Higgs-matter mixing

Taking into account the interactions in (3) and (4), the Higgs-pentaplet mass matrices are

$$M^{(2)} = \begin{bmatrix} (f\mu) & M_2 & \epsilon \\ \overline{M_2} & (f'\mu) & 0 \\ 0 & 0 & 0 \end{bmatrix} \quad (7)$$

and

$$M^{(3)} = \begin{bmatrix} (f\mu) & M_3 & \epsilon \\ \overline{M_3} & (f'\mu) & 0 \\ 0 & 0 & 0 \end{bmatrix} \quad (8)$$

in a $H_2$, $H_2'/\overline{H_2}$, $\overline{H_2'}$, $l_0$ and $H_3$, $H_3'/\overline{H_3}$, $\overline{H_3'}$, $d_0^c$ basis. The matter fields $l_0$ and $d_0^c$ are the combinations appearing in the coupling (4)

$$(\lambda_i \langle S\rangle^2/M)\phi_i H = \epsilon_i \phi_i H = \epsilon(l_0 H_2 + d_0^c H_3)\,. \quad (9)$$

We have set $\epsilon_i = \lambda_i \langle S\rangle^2/M$ and $\epsilon = (\sum \epsilon_i^2)^{1/2}$. Notice that $\epsilon$ is of the order of $\lambda\mu^2/M$, $\mu$ being the PQ breaking scale set by the $\langle S\rangle$, $\langle P\rangle$ v.c.v.'s.

The isodoublet mass eigenvalues can be read-off from $M^{(2)}(M^{(2)})^\dagger$. At this point we should impose the, inevitable, fine-tuning that will guarantee a light mass eigenvalue. It is convenient to put it in the form of the condition

$$(M_2\overline{M_2} - ff'\mu^2)^2 = \epsilon^2(M_2^2 + (f\mu)^2) \quad (10)$$

implying the appearance of a mass eigenvalue of the order of $\epsilon$. The resulting eigenvalues are

$$(m_2)_+ = (M_2^2 + \overline{M}_2^2 + (f\mu)^2 + (f'\mu)^2)^{1/2} \quad (11)$$

and

$$(m_2)_- = \epsilon\,. \quad (12)$$

Note that $(m_2)_+$ is of the order of $\mu$ since the condition (10) amounts to requiring that $M_2$, $\overline{M}_2$ are of that order. The combination

$$l = [(f'\mu)\overline{H}_2 - (\overline{M}_2)\overline{H}_2' + (M_2^2 + (f\mu)^2)^{1/2} l_0]/(m_2)_+ \quad (13)$$

is massless. The intermediate mass isodoublets $\overline{H}_+$, $H_+$ will have an appreciable influence on the running of gauge couplings. This is however within the limits allowed by existing data in correlation with proton decay [6]. The colour-triplet eigenvalues are both of order $M_3$, $\overline{M}_3$. The combination

$$d^c = N[d_0^c + \epsilon(M_3\overline{M}_3 - ff'\mu^2)^{-1} \times ((f'\mu)\overline{H}_3 - \overline{M}_3\overline{H}_3')] \quad (14)$$

is massless.

The standard down-quark Yukawa interactions written in terms of "mass-eigenstates"are

$$Y_i^{(d)}[(l_{0i}e_i^c + q_i'd_{0i}^c)\overline{H}_2 + (l_{0i}q_i' + u_i^c d_{0i}^c)\overline{H}_3] \quad (15)$$

with $q_i' = (u_i, V_{ij}d_j)$ in terms of the Kobayashi-Maskawa matrix $V_{ij}$. The combinations that mix with Higgses are $\epsilon_i l_{0i}$ and $\epsilon_i d_{0i}^c$. In general, all $\epsilon_i$'s are nonzero. We could always go to a new basis in which the combination $\epsilon_i\phi_i$ will define one family. For example, $l_1 = l_{01}$, $l_2 = l_{02}$ and $l_3 = \epsilon_i l_{0i}/\epsilon$.The new Yukawa's, according to (15) will be $Y_i' = Y_i - Y_3\epsilon_i/\epsilon_3$ for $i = 1,2$ and $Y_3' = Y_3\epsilon/\epsilon_3$. Nevertheless, it might be plausible [9], and certainly simplifying, to assume a family hierarchy in $\epsilon_i$ proportional to the hierarchical structure of the $Y_i$'s. In that case we could consider in $\epsilon_i\phi_i$ only the contribution of the (dominant) third family. Therefore, we proceed by assuming that only the third family has an appreciable R-violating coupling.

Substituting the expressions of $l_{03}$, $d_{03}^c$, $\overline{H}_2$ and $\overline{H}_3$ in terms of the light eigenstates, we obtain the leading order Yukawa coupling of the third generation

$$Y_3^{(d)}\Big[\frac{M_2}{\sqrt{M_2^2 + (f'\mu)^2}}(l_3\tau^c\overline{H}_-) + \frac{M_2^2 + 2(f\mu)^2}{\sqrt{(M_2^2 + (f'\mu)^2)(M_2^2 + (f\mu)^2)}}(q_3'b^c\overline{H}_-) - \frac{(f'\mu)M_2}{\sqrt{(M_2^2 + (f'\mu)^2)(M_2^2 + (f\mu)^2)}}(q_3'b^c l_3) + \ldots\Big] \quad (16)$$

No $B$-violating coupling appears due to colour antisymmetry. In contrast, the $L$-violating coupling $q_3'b^c l_3$ appears with an $O(1)$ coefficient. The Yukawa's of the other two generations are

$$\sum_{i=1,2} Y_i^{(d)}\Big[\frac{M_2^2 + 2(f'\mu)^2}{\sqrt{(M_2^2 + (f'\mu)^2)(M_2^2 + (f\mu)^2)}} \times (l_i e_i^c + q_i' d_i^c)\overline{H}_- - \frac{(f'\mu)M_2}{\sqrt{(M_2^2 + (f'\mu)^2)(M_2^2 + (f\mu)^2)}}(l_i e_i^c + q_i' d_i^c) l_3 + (\epsilon(f'\mu)/M_3\overline{M}_3)(q_i' l_i + u_i^c d_i^c) b^c\Big] \quad (17)$$

Note the presence of the $L$-violating interactions $\mu^c l_\mu l_\tau$, $e^c l_e l_\tau$, $s^c q_2' l_\tau$, $d^c q_1' l_\tau$ with $O(1)$ couplings while the $B$-violating operators $c^c s^c b^c$, $u^c d^c b^c$ carry a drastic suppression factor $\epsilon(f'\mu)/M_3\overline{M}_3$. This is a rather small number of the order of $10^{-20}$. This should be compared with the "direct" $B$-violating term $(\phi_i\phi_j\psi_k)\overline{P}S^2/M^3$, which carries an even smaller coefficient of the order of $(\mu/M)^3$.

The above hierarchy of $L$- versus $B$-non-conservation is sufficient to guarantee a stable proton since

$$\lambda'\lambda'' \sim (m_\mu/v_1)^2\epsilon(f'\mu)/M_3\overline{M}_3 \leq 10^{-24}. \quad (18)$$

Nevertheless a number of processes not respecting lepton-number result from (17). The interaction $\nu_\tau b' b^c$ generates at one loop a mass for the $\tau$-neutrino, roughly

$$\frac{Y_b^2}{16\pi^2}(m_b/\tilde{m}_b)^2 A m_{3/2}$$

which could be in agreement with existing cosmological bounds [11,13].

## 4. Other models

In the PQ-SU(5) model that has been analysed,the scale of PQ breaking has been "naturally" determined by the other scales present ($m_{3/2}$, $M_P$) and by the particular form of the superpotential couplings of the fields dictated by the symmetries. The suppression of R-violating terms, as in the analogous suppression of $D = 5$ operators that break Peccei-Quinn, is entirely independent of the fine-tuning required for the triplet-doublet splitting. This is much clearer in the so-called missing-doublet SU(5) model [12] endowed with a PQ symmetry [7]. This model has been constructed in order to avoid the fine numerical adjustment in the triplet-doublet mass splitting required in the minimal model. The superpotential is

$$W = \psi\psi H + \psi\phi\overline{H} + \overline{\lambda}H\Sigma\overline{\Theta} + \lambda\overline{H}\Sigma\Theta + \frac{M}{2}\text{Tr}(\Sigma^2) + \frac{h}{3}\text{Tr}(\Sigma^3) + \overline{\lambda}'H'\Sigma\overline{\Theta}' + \lambda'\overline{H}'\Sigma\Theta' + M_1\Theta\overline{\Theta}' + M_2\Theta'\overline{\Theta}. \quad (19)$$

The SU(5) and $U(1)_{PQ}$ quantum numbers of the fields are $\psi(10,\alpha/2)$, $\phi(\overline{5},\beta/2)$, $H(5,-\alpha)$, $\overline{H}'(\overline{5},\alpha)$, $\overline{H}(\overline{5},-(\alpha+\beta)/2)$, $H'(5,(\alpha+\beta)/2)$, $\Theta(\overline{50},\alpha)$, $\overline{\Theta}'(\overline{50},-(\alpha+\beta)/2)$, $\Theta(50,(\alpha+\beta)/2)$, $\Theta'(50,-\alpha)$, $\Sigma(75,0)$. The masses $M_1$, $M_2$ are taken to be of the order of the Planck-mass in order to avoid an increase of the gauge coupling beyond the perturbativity limit due to the presence of too many light fields. Integrating out the superheavy 50's we obtain the effective superpotential

$$\psi\psi H + \psi\phi\overline{H} + H_3'\overline{H}_3M_3 + H_3\overline{H}_3'\overline{M}_3 \quad (20)$$

in which only the colour-triplets appear with masses $M_3 = \lambda\overline{\lambda}'\langle\Sigma\rangle^2/M_1$, $\overline{M}_3 = \lambda'\overline{\lambda}\langle\Sigma\rangle^2/M_2$. Both these masses are slightly below the unification scale, namely $10^{14}$–$10^{15}$ GeV. There are no mass terms for the doublets as a consequence of the absence of direct mass terms for the pentaplets.

In addition to the interactions appearing in (19), new interaction terms are possible if gauge-singlet fields, charged under PQ are introduced. Being a little different from the case of the minimal PQ-SU(5), we introduce $P(-(3\alpha+\beta)/2)$, $Q(3(3\alpha+\beta)/2)$ and $S((\alpha-\beta/2)/3)$. No other renormalizable terms are possible with these fields except

$$fP\overline{H}'H'. \quad (21)$$

Again, various non-renormalizable interactions are present. They are

$$P^3Q/M + (\overline{H}H)P^2Q/M^2 + \tilde{\lambda}_iS^3(\phi_iH)/M^2. \quad (22)$$

All these terms can be written down for charges defined for independent $\alpha$'s and $\beta$'s. This reflects the existence of two U(1)'s of which one can be broken by an extra interaction of leading non-renormalizable order $1/M$ among the fields $P$, $Q$, $S$ that forces a relation among the phases. For example the interaction $P^2QS/M$ enforces the, peculiar, phase relation $2\beta = -11\alpha$. In any case, the breaking of the $U(1)_{PQ}$ proceeds in a similar way as in the minimal model, coming out again in the range $10^{10}$–$10^{12}$ GeV.

The Higgs pentaplet mass matrices are

$$M^{(2)} = \begin{bmatrix} \hat{\epsilon} & 0 & \epsilon \\ 0 & (f\mu) & 0 \\ 0 & 0 & 0 \end{bmatrix} \quad (23)$$

and

$$M^{(3)} = \begin{bmatrix} \hat{\epsilon} & \overline{M}_3 & \epsilon \\ M_3 & (f\mu) & 0 \\ 0 & 0 & 0 \end{bmatrix} \quad (24)$$

in a $H_2, H_2'/\overline{H}_2, \overline{H}_2', l_{03}$ and $H_3, H_3'/\overline{H}_3, \overline{H}_3', d_{03}^c$ basis. Again for simplicity we have assumed that the R-non-conserving coupling is exclusively to the third generation. The appearing parameters are $\epsilon = \tilde{\lambda}\langle S\rangle^3/M^2 \sim 10^2$–$10^3$ GeV, for an intermediate PQ scale choice of $10^{11}$ GeV, and $\hat{\epsilon} = \langle P\rangle^2\langle Q\rangle/M^2$, roughly of the same order. The doublet mass matrix leads to eigenvalues $m_+^2 = (f\mu)^2$ and $m_-^2 = \epsilon^2 + \hat{\epsilon}^2$. The combination $l_\tau = (\epsilon\overline{H}_2 - \hat{\epsilon}l_{03})/m_-$ is massless. The triplet eigenvalues are both of order $M_3 \sim \overline{M}_3$. The combination

$$b^c = N[\epsilon(f\mu)\overline{H}_3 - \epsilon M_3\overline{H}_3' + (M_3\overline{M}_3 - \hat{\epsilon}f\mu)d_{03}^c] \quad (25)$$

is massless. Expressing the down-quark Yukawa's in terms of eigenstates we obtain

$$Y_3^{(d)}[l_\tau\tau^c\overline{H}_- + \frac{\hat{\epsilon}}{\sqrt{\epsilon^2+\hat{\epsilon}^2}}(q_3'b^c\overline{H}_-) + \frac{\epsilon}{\sqrt{\epsilon^2+\hat{\epsilon}^2}}(q_3'b^cl_\tau)]\sum_{i=1,2}Y_i^{(d)}[(q_i'l_i + d_i^cu_i^c)b^c \times (\frac{\epsilon f\mu}{M_3^2})(\frac{M_3\overline{M}_3}{M_3^2-\overline{M}_3^2} - (\frac{M_3}{\overline{M}_3})^3) + (q_i'd_i^c + l_ie_i^c)(\overline{H}_-\hat{\epsilon} + l_\tau\epsilon)/\sqrt{\epsilon^2+\hat{\epsilon}^2}] \quad (26)$$

Again, the $\Delta B/\Delta L$ hierarchy is of order $\epsilon\mu/M_3^2$ and $\lambda'\lambda'' \sim (m_\mu/v_1)^2(\epsilon f\mu/M_3^2) \le 10^{-24}$.

## 5. Discussion and brief summary

Both models analyzed above contain an extra pair of intermediate mass isodoublets which could, in principle, jeopardize the agreement of low-energy data with unification. The analysis has been done in Ref. [7]. Using the experimental data quoted there, i.e. $\alpha^{-1}(M_Z) = 127.9 \pm 0.2$, $\sin^2\theta_W(M_Z) = 0.2326 \pm 0.0008$ and $\alpha_3(M_Z) = 0.118 \pm 0.007$, these authors obtain, in the case of the *minimal PQ-SU(5) model*, the inequality

$$2.2 \times 10^{13}\ \text{GeV} \le (M_{H_c}M_{\overline{H}_c}/M_{\overline{H}_+}) \le 2.3 \times 10^{17}\ \text{GeV} \quad (27)$$

and, in the case of the *missing doublet PQ-SU(5) model*,

$$3.7 \times 10^{17}\ \text{GeV} \le (M_{H_c}M_{\overline{H}_c}/M_{\overline{H}_+}) \le 3.8 \times 10^{21}\ \text{GeV}. \quad (28)$$

With $M_{\overline{H}_+}$ we denote the mass of the heavy isodoublets. The effective scale appearing in these inequalities is also the scale appearing in the proton decay rate through $D = 5$ operators. Note that (28) puts a much weaker constraint than (27) and this model is still consistent with the lower limit on the nucleon lifetime even for the case of large $\tan\beta$ [7]. The missing doublet PQ-SU(5) model is easily in agreement with low-energy data. Note that the difference in the two models arises due to extra terms in the renormalization group equations coming from the mass- splittings within the $\Sigma(75)$ Higgs field. Using the more recent data [7] $\alpha^{-1}(M_Z) = 127.9 \pm 0.2$, $\sin^2\theta_W(M_Z) = 0.2314 \pm 0.0004$ and $\alpha_3(M_Z) = 0.118 \pm 0.0007$, the inequality becomes

$$1.4 \times 10^{17}\ \text{GeV} \le (M_{H_c}M_{\overline{H}_c}/M_{\overline{H}_+}) \le 5.5 \times 10^{20}\ \text{GeV}. \quad (29)$$

For $M_{H_c} \sim M_{\overline{H}_c} \sim 10^{13-16}$ GeV, we can have $M_{\overline{H}_+} \sim 10^{10}$ GeV. We do not wish to go any further into the analysis of the predictions of these models here, since most of it has already been done [6,7,14], and since it is not the central issue of the present article. The reason for employing both versions of the PQ-SU(5) model is to illustrate that our proposed $R$-parity violation paradigm is, to a large extent, model- independent.

The $L$-violating couplings of (26) as well as of (17) lead to a number of phenomenological implications apart from neutrino masses, like new exotic decays or just new important contributions to various processes. Most of these have been analysed in the literature [8] and will not be considered here. In the present article we addressed the question of whether it is possible in GUTs to obtain R-parity violation with a large $\Delta L/\Delta B$ hierarchy of strengths so that the proton stability is ensured while interesting $L$-non-conserving processes exist at appreciable rates. We considered variants of the SU(5) GUT with a built-in Peccei-Quinn symmetry suitable for suppressing $D = 5$ $B$-violating operators. It turns out that a spontaneously broken Peccei-Quinn symmetry in conjunction with an appropriate field content can result in an effective R-parity breaking characterized by a large hierarchy.

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