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R-PARITY VIOLATION AND PECCEI-QUINN SYMMETRY IN GUTS

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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 P-Q 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.

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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 $R = (-1)^{3B+L+2S}$ with 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 Λ, these operators will be supressed by Λ/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 P-Q symmetry have been constructed[5, 6, 7].

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

$$\lambda_{ijk} l_i l_j e^c_k + \lambda'_{ijk} d^c_i l_j q_k + \lambda''_{ijk} d^c_i d^c_j u^c_k + \epsilon_i l_i H \tag{1}$$

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, addopting a phenomenological attitude, to assume the existence of some of these couplings while forbidding the presense 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

$$\lambda_{ijk} \phi_i(\overline{5}) \phi_j(\overline{5}) \psi_k(10) + \epsilon_i \phi_i(\overline{5}) H(5) \tag{2}$$

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].

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 exclusivelly 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, they 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 P-Q -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

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

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 P-Q 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 P-Q 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 explicitely imposed, is an exact symmetry of the model even after P-Q spontaneous breaking. Although we cannot exclude that R-parity is indeed an exact symmetry it is more interesting to explore the possibility that additional interactions exist which ultimately lead to effective R-violating couplings among standard fields such as ϕ_iH, ϕ_iH' and $\psi_i\phi_j\phi_k$. For instance, singlets carrying odd P-Q charge could couple to the above operators. A viable model however should predict also the neccessary suppression of these effective couplings. It is possible to costruct models in which the P-Q 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 P-Q charge 1/2. This choice of charge ensures the absence of renormalizable couplings to the other fields. Then, the R-violating term

$$\lambda_i\phi_iHS^2/M \tag{4}$$

is possible. A P-Q-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 of 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

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

and

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

are also present but their suppression with extra powers of the Planck mass makes them not relevant.Terms with Σ 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

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

and

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

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)

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

We have set $\epsilon_i = \lambda_i\langle S\rangle^2/M$ and $\epsilon = (\sum \epsilon_i^2)^{1/2}$. Notice that ϵ is of the order of $\lambda\mu^2/M$, μ being the P-Q breaking scale set by the $\langle S\rangle$, $\langle P\rangle$ v.e.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-tunning that will guarantee a light mass-eigenvalue. It is convenient to put it in the form of the condition

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

implying the appearence of a mass-eigenvalue of the order of ϵ. The resulting eigenvalues are

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

and

$$(m_2)_- = \epsilon \tag{12}$$

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

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

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

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

is massless.

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

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

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 ϵ_i's are non-zero. 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 ϵ_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 procceed by assuming that only the third family has an appreciable R-violating coupling.

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

$$Y_3^{(d)}[\frac{M_2}{\sqrt{M_2^2+(f'\mu)^2}}(l_3\tau^c\overline{H}_-)+\frac{M_2^2+2(f\mu)^2}{\sqrt{(M_2^2+(f'\mu)^2)(M_2^2+(f\mu)^2)}}(q'_3b^c\overline{H}_-)$$

$$-\frac{(f'\mu)M_2}{\sqrt{(M_2^2+(f'\mu)^2)(M_2^2+(f\mu)^2)}}(q'_3b^cl_3)+\ldots] \qquad (16)$$

No B-violating coupling appears due to colour antisymmetry. In contrast, the L-violating coupling $q'_3b^cl_3$ appears with an O(1) coefficient. The Yukawa's of the other two generations are

$$\sum_{i=1,2}Y_i^{(d)}[\frac{M_2^2+2(f'\mu)^2}{\sqrt{(M_2^2+(f'\mu)^2)(M_2^2+(f\mu)^2)}}(l_ie_i^c+q'_id_i^c)\overline{H}_- \qquad (17)$$

$$-\frac{(f'\mu)M_2}{\sqrt{(M_2^2+(f'\mu)^2)(M_2^2+(f\mu)^2)}}(l_ie_i^c+q'_id_i^c)l_3+(\epsilon(f'\mu)/M_3\overline{M}_3)(q'_il_i+u_i^cd_i^c)b^c]$$

Note the presence of the L-violating interactions $\mu^cl_\mu l_\tau$, $e^cl_el_\tau$, $s^cq'_2l_\tau$, $d^cq'_1b^c$ with O(1) couplings while the B-violating operators $c^cs^cb^c$, $u^cd^cb^c$ carry a drastic suppression factor $\epsilon(f'\mu)/M_3\overline{M}_3$. This is a rather small number of the orderd of 10^{-20} . This should be compared to 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

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

Nevertheless a number of procceces 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 $\frac{Y_b^2}{16\pi^2}(m_b/\tilde{m}_b)^2Am_{3/2}$ which, being of the order of Mev, is easily in agreement with existing cosmological bounds[11, 13].

4.Other models. In the P-Q SU(5) model that has been analyzed,the scale of P-Q 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-tunning required for the triplet-doublet splitting. This is much more clear in the so-called missing-doublet SU(5) model[12] endowed with a P-Q 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

$$W=\psi\psi H+\psi\phi\overline{H}+\overline{\lambda}\overline{H}\Sigma\overline{\Theta}+\lambda H\Sigma\Theta+\frac{M}{2}Tr(\Sigma^2)+\frac{h}{3}Tr(\Sigma^3) \qquad (19)$$

$$+\overline{\lambda}'\overline{H}'\Sigma\overline{\Theta}' + \lambda' H'\Sigma\Theta' + M_1\Theta\overline{\Theta}' + M_2\Theta'\overline{\Theta}$$

The SU(5) and $U(1)_{PQ}$ quantum numbers of the fields are $\psi(10,\alpha/2)$, $\phi(\overline{5},\beta/2)$, $H(5,-\alpha)$, $\overline{H}'(\overline{5},\alpha)$, $\overline{H}(\overline{5},-(\alpha+\beta)/2)$, $H'(5,(\alpha+\beta)/2)$, $\Theta(50,\alpha)$, $\Theta'(50,-(\alpha+\beta))$, $\overline{\Theta}(\overline{50},(\alpha+\beta)/2)$, $\overline{\Theta}'(\overline{50},-\alpha)$, $\Sigma(75,0)$ .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

$$\psi\psi H + \psi\phi\overline{H} + H_3'\overline{H}_3 M_3 + H_3\overline{H}_3'\overline{M}_3 \tag{20}$$

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 bellow the unification scale, namely 10^{14}–10^{15}Gev. There is 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 P-Q are introduced. Being a little different than the case of the minimal P-Q 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

$$fP\overline{H}'H' \tag{21}$$

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

$$P^3Q/M + (\overline{H}H)P^2Q/M^2 + \tilde{\lambda}_i S^3(\phi_i H)/M^2 \tag{22}$$

All these terms can be written down for charges defined for independent α's and β'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 $U(1)_{PQ}$ procceeds 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

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

and

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

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 P-Q 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

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

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

$$Y_3^{(d)}[l_\tau\tau^c\overline{H}_- + \frac{\hat{\epsilon}}{\sqrt{\epsilon^2+\hat{\epsilon}^2}}(q_3'b^c\overline{H}_-) + \frac{\epsilon}{\sqrt{\epsilon^2+\hat{\epsilon}^2}}(q_3'b^cl_\tau)] \tag{26}$$

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

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

5.Brief summary. The L-violating couplings of (26) aswell 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 analyzed in the literature[8] and will not be considered here. In the present article we addressed the question of whether it is posssible 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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