LEFT RIGHT SUSY AND THE FATE OF R-PARITY

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A fresh analysis of left right symmetric supersymmetric models in the generic case where the scale of right handed symmetry breaking $M_R >> M_{SUSY} \sim M_W$ is presented. We conclude that the low energy effective theory for such models is essentially the MSSM with R parity (and therefore B,L symmetry) but the spectrum includes heavy conjugate neutrino supermultiplets that permit a seesaw mechanism and several characteristic charged supermultiplets over and above those of the MSSM.

1 Introduction

The Standard Model (SM) has the appealing feature that (perturbative) baryon (B) and lepton (L) number conservation follows automatically from the requirements of renormalizability and gauge invariance of the lagrangian for the fields observed in nature (plus the Higgs required for spontaneous symmetry breaking). Supersymmetric extensions of the standard model, where each of the observed fields of the SM has a partner of opposite statistics, do not retain this property since the presence of scalar fields carrying B and L implies that one can write renormalizable gauge invariant interactions that violate B and L. These terms are described by the superpotential

$$W_2 = \mu' \bar{L} \hat{H} + \lambda_1 \bar{u} e \bar{d} e^c + \lambda_2 Q L d e + \lambda_3 L L e e^c$$

(1)

These interactions imply ultrarapid proton decay and appreciable amplitudes for a host of exotic processes unless the couplings $\lambda_i, \mu'$ are extremely small. The couplings in $W_2$ can be ensured to vanish if one invokes an ad-hoc $\mathbb{Z}_2$ symmetry, the so-called R parity under whose action the superfields pick up a phase $(-)^{3(B-L)}$ and the supercoordinates $\theta$ change sign. Equivalently one may describe it as a symmetry under which all superpartners of SM fields are odd while SM fields are even. An attractive rationale for this symmetry is for it to follow automatically from the presence of a gauged $B - L$ symmetry at the supersymmetric level. This would also protect it from the violation of global symmetries thought to occur due to Planck scale effects.

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violation of $R$ parity becomes tied to the spontaneous breaking of the $B - L$
gauge symmetry and can occur only if a scalar field with an odd value of $B - L$
aquires a vev. Whether and to what extent this happens has a crucial role in determining the low energy phenomenology of the SUSY theory. In particular it determines whether the LSP is stable and affects neutrino masses in interesting and predictable ways.

The Left Right symmetric model\( ^7 \) is one of the most appealing “Next to Minimal” models since it provides a “spontaneous” rationale for the maximal parity violation observed in the standard model along with a framework to understand the suppression of neutrino masses. Moreover its $U(1)$ gauge group turns out to be identical with $U(1)_{B-L}$ so that while incorporating the standard model it “takes care” to gauge its maximal anomaly free global symmetry! Thus a formulation of a minimal left right supersymmetric extension of the standard model in which $R$ parity is a part of the gauge symmetry offers the prospect of furnishing a model with rationales for several problematic features of the standard model.

In this talk I shall present recent progress in these directions made in collaboration with Karim Benakli, Alejandra Melfo, Andrija Rasin and Goran Senjanovic\( ^8, ^9, ^10 \) and embodied in two workable models: one with a renormalizable superpotential but an extended field spectrum that allows decoupling of the scales ($M_R, M_{B-L}$) at which the $SU(2)_R$ and $U(1)_{B-L}$ symmetries are broken and another in which these scales coincide but nonrenormalizable interactions (upto quartic) are retained in the superpotential. An important aspect of our work is that we employ a powerful method for the characterizing the flat directions of supersymmetric vacua by the holomorphic gauge invariants of chiral supermultiplets left independent after imposition of the holomorphic constraints required by the minimization of the potential in supersymmetric gauge theories. This method is useful only for analysis of symmetry breaking at scales where supersymmetry is a good symmetry. Hence our analysis is not applicable to the special case when the right handed and $B - L$ breaking scales are of the same order of magnitude as the supersymmetry breaking scale ($M_S$) but only to the generic case when these scales are much greater than $M_S$.

The main conclusion of our analysis is that the structure of the “minimal” supersymmetric left right vacua together with low energy data imply that unless electromagnetic charge invariance is violated $R$ parity remains unbroken so that the effective low energy theory is the MSSM with $R$ parity with certain characteristic additional supermultiplets in the spectrum. The most striking experimental signature of these theories is the presence of a number of charged Higgs supermultiplets with masses much less than $M_R$ including
doubly charged particles in the nonrenormalizable case.

2 Minimal Supersymmetric Left Right Models

As mentioned above the gauge group of these models is $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. The gauge bosons are accompanied by their corresponding gauginos while the matter fields of the standard model plus the right handed neutrino field required by left-right symmetry are grouped into three generations of quark and leptonic chiral superfields with the following transformation properties:

$$Q = (3, 2, 1, 1/3) \quad Q_c = (3^*, 1, 2, -1/3)$$
$$L = (1, 2, 1, -1) \quad L_c = (1, 1, 2, 1) \quad (2)$$

The doublet Higgs fields of the standard model are doubled in number by the requirement of anomaly cancellation in the MSSM and once again by left-right symmetry. Thus they fit into “bidoublets” w.r.t the left right gauge group. Realistic fermion mass matrices require that we introduce at least a pair of such bidoublets :

$$\Phi_i = (1, 2, 2^*, 0) \quad (i = 1, 2) \quad (3)$$

Where the numbers in brackets give the transformation properties of the chiral multiplet with respect to the gauge group.

In addition to these fields we must choose Higgs representations to accomplish the breaking of the $SU(2)_R \times U(1)_{B-L}$ symmetry down to the $U(1)_Y$. The two simplest choices are pairs of doublets and triplets respectively. The latter choice allows a see-saw mechanism even at the renormalizable level . Moreover in the supersymmetric case the doublets carry odd $B-L$ so that the $M_{B-L}$ and the scale of R parity breaking necessarily coincide . This implies that additional ad-hoc features must be introduced to control the R parity breaking couplings to lie within the stringent experimental limits. The triplet case is free of these problems so we shall consider only this case here. The fields we use are :

$$\Delta = (1, 3, 1, 2), \quad \Delta = (1, 3, 1, -2)$$
$$\Delta_c = (1, 1, 3, -2), \quad \Delta_c = (1, 1, 3, 2) \quad (4)$$

The doubling of Higgs triplets here relative to the non supersymmetric case again follows from anomaly cancellation. Left right symmetry is implemented
as a charge conjugation (convenient when one wishes to embed in a GUT) rather than as an (equally feasible and equivalent) parity transformation:

\[
Q \leftrightarrow Q_c, \quad L \leftrightarrow L_c, \quad \Phi_i \leftrightarrow \Phi_i^T, \\
\Delta \leftrightarrow \Delta_c, \quad \Sigma \leftrightarrow \Sigma_c. 
\]

(5)

With this minimal set of Higgs fields, however, the superpotential for these triplets which are to accomplish the breaking to SU(2) is merely

\[
W_{LR} = i\mathcal{F}(L^T \tau_2 \Delta L + L'^T \tau_2 \Delta_c L_c) + m_\Delta (\text{Tr} \, \Delta \Delta^c + \text{Tr} \, \Delta_c \Delta_c^c) 
\]

(6)

It is easy to see that the vanishing of F terms required by the minimization of the superpotential now requires that the vevs of \(\Delta, \Delta_c\) vanish while the vevs of \(\Delta, \Delta_c\) are proportional to those of \(fLL\) and \(f^* L_c L_c\) respectively. Moreover since squark vevs break color and charge we assume that they are forbidden by soft mass terms. The vanishing of \(D_{B-L}\) then implies that \(SU(2)_R\) and \(SU(2)_L\) are broken at the same scale. Although one can evade this problem by introducing a "parity odd" singlet it leads to additional difficulties in achieving a realistic breaking pattern and is aesthetically displeasing in a gauge theory to boot. Instead we employ either an additional pair of \(B-L\) neutral triplets

\[
\Omega = (1, 3, 1, 0), \quad \Omega_c = (1, 1, 3, 0) 
\]

(7)

where under Left-Right symmetry \(\Omega \leftrightarrow \Omega_c\), while maintaining renormalizability, or allow non-renormalizable terms.

The potential of a supersymmetric gauge theory has the form

\[
V = \sum_{\text{Chiral}} |F_i|^2 + \sum_{\text{Gauge}} D^2 
\]

(8)

and its minimization thus requires that the F and D terms corresponding to each chiral multiplet and gauge generator respectively vanish i.e that the vacuum manifold is "D and F flat". The following powerful theorem is extremely useful in characterizing solutions of these equations:

**Theorem**:

a) The D-flat field space is coordinatized by the independent Holomorphic gauge invariants that one can form out of the chiral multiplets.

b) The D and F flat field space is coordinatized by the Holomorphic invariants left independent after imposition of the \(F = 0\) conditions.

As an example: consider a \(U(1)\) gauge theory with two chiral multiplets \(\phi_\pm\) with gauge charges \(\pm 1\). Then the condition \(D = 0\) requires only \(|\phi_+| =


φ−. Since gauge invariance can be used to rotate away one field phase we are left with a magnitude and a phase i.e one complex degree of freedom left undetermined. Result a) above predicts this since the only independent holomorphic gauge invariant in this case is simply φ+φ−. Now consider the effects of a superpotential \( W = mφ+φ− \). The F flatness condition now ensures that both vevs vanish so that the D flat manifold shrinks from the complex line parametrized by \( c = φ+φ− \) to the single point \( c = 0 \). Typically in the minimization of the potential of a SUSY gauge theory one finds that the space of vacua (the “moduli space”) may consist of several sectors corresponding to “flat directions” running out of various minima that would be isolated if a suitably smaller set of chiral multiplets had been used.

3 Renormalizable Model with Triplets

Let us now consider the pattern of symmetry breaking at \( M_R >> M_S \sim M_W \). We assume that the soft terms are such as to keep squark and bi-doublet vevs zero since these would break colour, charge or the \( SU(2)_L \) symmetry at the high scale. On the other hand since the neutral component of the \( L_c \) field could \textit{a priori} get a v.e.v without violating any low energy symmetry and we would like to consider the case where L-R symmetry is spontaneously violated, we keep both the \( L_c \) and the \( L \) fields besides the triplets. The most general renormalizable superpotential is then:

\[
W_{LR} = h^{(i)} L^T \tau_2 \Phi_i \tau_2 L_c + i f (L^T \tau_2 \Delta L + L'^T \tau_2 \Delta_c L_c) \\
\quad + m_\Delta (\text{Tr} \Delta \Xi + \text{Tr} \Delta_c \Xi_c) + \frac{m_\Omega}{2} (\text{Tr} \Omega^2 + \text{Tr} \Omega_c^2) \\
\quad + \mu_{ij} \text{Tr} \tau_2 \Phi_i^T \tau_2 \Phi_j + a (\text{Tr} \Omega \Xi + \text{Tr} \Delta_c \Omega_c \Xi_c) 
\]

(9)

Multiplying the the F flatness conditions for the triplet equations by triplet fields and taking traces it immediately follows

\[
\text{Tr} \Delta^2 = \text{Tr} \Delta \Omega = \text{Tr} \Delta_c \Omega_c = 0 \\
m_\Delta \text{Tr} \Delta \Xi = m_\Omega \text{Tr} \Omega^2 = a \text{Tr} \Delta \Xi \Omega \\
\text{Tr} \Delta_c \Xi (a^2 \text{Tr} \Omega^2 - 2 m_\Delta^2) = 0 
\]

(10)

with corresponding equations \textit{mutatis mutandis} in the right handed sector. Thus it is clear that in either sector all three triplets are zero or non zero together. By choosing the branch where \( \text{Tr} \Omega_c^2 = \frac{2m_\Delta^2}{a^2} \) but \( \text{Tr} \Omega^2 = 0 \) we ensure that the triplet vevs break \( SU(2)_R \) but not \( SU(2)_L \).
One can use the 3 parameters of the $SU(2)_R$ gauge freedom to set the diagonal elements of $\Delta_c$ to zero so that it takes the form

$$
\langle \Delta_c \rangle = \begin{pmatrix} 0 & \langle \delta_c^{--} \rangle \\ \langle \delta_c^{--} \rangle & 0 \end{pmatrix}
$$

(11)

Now (10) gives $\langle \delta_c^{--} \rangle \langle \delta_c^{0} \rangle = 0$, which implies the electromagnetic charge-preserving form for $\langle \Delta_c \rangle$. Next it is clear that the Majorana coupling matrix $f_{ab}$ must be non-singular if the see saw mechanism which keeps the neutrino light is to operate. Then it immediately follows from the condition $F_{Lc} = 0$, namely,

$$
2i f_{ab} \begin{pmatrix} 0 & 0 \\ \langle \delta_c^{0} \rangle & 0 \end{pmatrix} \begin{pmatrix} \nu_c \\ e_c \end{pmatrix}^b = 0
$$

(12)

that the sneutrino vevs in the right-handed sector must vanish. Thus any vev of $L_c$ that appears at the high scale must necessarily break charge. We ensure that it (together with $L$) vanishes by suitably positive soft masses, just as for the squarks. When the sleptons have zero vevs one can show that all holomorphic invariants that one can form from the triplets are fixed (at each of the allowed minima) so the parity violating vacuum is isolated. and in fact the vevs can be shown to necessarily take the form

$$
\langle \Omega_c \rangle = \begin{pmatrix} w & 0 \\ 0 & -w \end{pmatrix}, \quad \langle L_c \rangle = 0
$$

$$
\langle \Delta_c \rangle = \begin{pmatrix} 0 & 0 \\ d & 0 \end{pmatrix}, \quad \langle \Xi_c \rangle = \begin{pmatrix} 0 & \bar{d} \\ 0 & 0 \end{pmatrix}
$$

(13)

In fact using the $B - L$ gauge invariance to fix the relative phase of $d$ and $\bar{d}$ one obtains

$$w = -\frac{m_\Delta}{a} \equiv -M_R, \quad d = \bar{d} = \left( \frac{2m_\Delta m_0}{a^2} \right)^{1/2} \equiv M_{BL}
$$

(14)

Furthermore one can show that the inclusion of slepton fields leads to 9 (complex) flat directions out of the discrete minima of the superpotential with triplets alone and these are parametrized in a gauge invariant way by $(a,b,. . .)$ are family indices)

$$ z_{[ab]d} = \sigma_{[ab]}(Tr \tilde{\Delta}^{(c)}_{d[ab]} \tilde{\Delta}^{(c)}) \tilde{\nu}_{a[c]} \tilde{e}_{b[c]} \tilde{e}_{d[c]}
$$

(15)

where the composite multiplets $\sigma_{ab}$ (singlet) and $\tilde{\Delta}_{ab}$ (triplet) are defined by
\[ L_a \bar{L}_b = \sigma_{ab} + \bar{\Sigma}_{ab} \]  

and similarly in the right handed sector.

4 R Parity

We have shown that R parity breaking at the high scale \( M_R \) necessarily implies breaking of electromagnetism so that one must assume that the soft terms are such as to forbid such breaking. Then the effective theory below \( M_{R,B-L} \) is the MSSM with R parity (and a heavy \( \nu^{(c)} \) superfield) hence a global B, L invariance! However, one may worry that, due to the running of couplings, the sneutrinos may obtain vevs; analogously to the Higgs in the MSSM which develops a vev inspite of a positive soft mass at high scales due to the effect of its large yukawa coupling to the top quark on its effective mass. In the present case however this is unlikely. Firstly \( \tilde{\nu}^{(c)} \) has a large mass and can shift away from zero only due to a linear term in the potential. Such a linear term can develop via the trilinear soft terms once \( \tilde{\nu} \) acquires a vev but is not present otherwise. On the other hand a vev for the left handed sneutrino implies the presence of a goldstone boson (the “doublet majoron”) with an appreciable coupling to the the low energy gauge group currents and is forbidden by the precise measurements of the Z width at LEP. One final possibility that one must consider is that the vev for \( \tilde{\nu} \) leads to a vev for \( \tilde{\nu}^{(c)} \) which in turn provides enough explicit L violation to give the putative majoron a mass greater than \( M_Z/2 \). However it is easy to see that

\[ \langle \nu^{(c)} \rangle \approx \frac{m_S M_W \langle \nu \rangle}{M_{BL}^2} \]  

(17)

This would lead to effective R-parity and global lepton number violating terms of the form \( m_{\tilde{\nu}}^2 LH \) where

\[ m_{\tilde{\nu}}^2 \approx \frac{m_S^2 M_W \langle \nu \rangle}{M_{BL}^2} \]  

(18)

Then the “Majoron” would get a mass squared of order

\[ m_J^2 \approx m_{\tilde{\nu}}^2 \langle \nu \rangle \approx \frac{m_S^2 M_W}{M_{BL}^2} \]  

(19)

Thus in order that \( m_J \) be large enough to evade the width bound the scale \( M_{BL} \) would have to be \( O(m_S) \) : which possibility we do not consider here as discussed earlier.

Thus the effects of running cannot change our earlier conclusions regarding the effective theory at low energies.
5 Mass spectrum and Seesaw Mechanism

In the renormalizable model $SU(2)_R$ is broken down to $U(1)_R$ at a large scale $M_R = m_\Delta/a$, by the vev of $\Omega_c$. Later the vevs of $\Delta_c, \bar{\Delta}_c$ are turned on at $M_{BL} = \sqrt{2m_\Delta m_\Omega/a}$, breaking $U(1)_R \times U(1)_{BL}$ to $U(1)_Y$. However, a third scale appears in the superpotential, $m_\Omega = M_{BL}^2/M_R$. Although most of the extra fields of the renormalizable LR SUSY model get masses at $M_R$ or $M_{B-L}$ there are two notable exceptions.

Firstly a fine tuning is necessary in order to keep one pair of Electroweak doublets (out of the 2 pairs contained in the 2 bidoublets) light to serve to break the Electroweak group at $M_W$:

$$\mu_{11} \mu_{22} - (\mu_{12}^2 - \alpha_{12}^2 M_R^2) \simeq M_W^2$$

What is striking however is that the spectrum contains a complete $SU(2)_L$ triplet ($\Omega$) of scalars and fermions with masses $\sim M_{B-L}/M_R$. If the two scales are well separated this mass could near the lower limit $\sim M_W$ imposed by the consistency of the analysis.

Finally the seesaw mechanism takes its canonical form and is free of “contamination” namely the Majorana mass for the left handed neutrino that arises in non-SUSY theories since $\Delta$ acquires a small vev due to terms of form $\Delta^2 \Phi^2 \Delta_c$ in the potential which lead to a linear term for $\Delta$ after symmetry breaking.

6 Non-Renormalizable model

I now briefly describe the other alternative namely allowing non renormalizable terms in the superpotential while retaining the minimal set of fields. As before one ignores the squarks and bidoublets while analyzing the potential at $M_{B-L} = M_R$ so that the relevant superpotential to next to renormalizable order is

$$W_{nr} = m(\text{Tr} \Delta \bar{\Delta} + \text{Tr} \Delta_c \bar{\Delta}_c) + if(L^T \tau_2 \Delta L + L_c^T \tau_2 \Delta_c L_c)$$

$$+ \frac{a}{2M} [\text{(Tr} \Delta \bar{\Delta})^2 + (\text{Tr} \Delta_c \bar{\Delta}_c)^2] + \frac{c}{M} \text{Tr} \Delta \bar{\Delta} \text{Tr} \Delta_c \bar{\Delta}_c$$

$$+ \frac{b}{2M} [\text{Tr} \Delta^2 \tau_2 \Delta^2 + \text{Tr} \Delta_c^2 \tau_2 \Delta_c^2] + \frac{k}{M} L^T \tau_2 L L_c^T \tau_2 L_c +$$

$$+ \frac{1}{M} [d_1 \text{Tr} \Delta^2 \text{Tr} \bar{\Delta}^2 + d_2 \text{Tr} \Delta_c^2 \text{Tr} \bar{\Delta}_c^2] + \ldots$$

The nonrenormalizable terms are suppressed by inverse powers of a large scale $M$ (for instance the Planck mass). If a renormalizable interaction (such as
\((L_c \bar{L}_c)^2\) is already present then we can safely neglect corrections \(\sim M^{-1}\). Now using arguments exactly similar to those we gave for the renormalizable case one can show that the conjugate sneutrino vev must vanish at the high scale where SUSY is good and hence R parity and electric charge can only be broken together. Thus one must again invoke soft terms to avoid this unpleasant possibility. With \(L, \bar{L}_c\) vevs zero one uses the holomorphic invariants technique to show that there is an isolated parity violating vacuum (besides the trivial one with unbroken symmetry and other charge violating ones). The detailed analysis given in [10] is a good illustration of the power of this technique. The scale of symmetry breaking is the geometric mean of the mass parameter \(m/a\) and the large scale \(M\). Thus if \(M \sim M_P\) then \(m \geq 1\) TeV gives \(M_R \geq 10^{11}\) GeV.

As before one can finetune to keep a pair of doublets light. However what is remarkable is that there is a plethora of supermultiplets with masses as low as \(m \sim 1\) TeV namely one neutral and two doubly charged fields from the triplets \(\Delta, \bar{\Delta}_c\) together with two complete \(SU(2)_L\) triplets \(\Delta, \bar{\Delta}\) and two \(SU(2)_L\) doublets (over and above the pair of \(SU(2)_L\) doublets of the MSSM). Thus such “minimal” Left-Right Supersymmetric model have the spectacular signature of 4 distinct light doubly charged supermultiplets. It appears that this property is quite robust since it is connected with the fact that without non-renormalizable terms the vacuum breaks charge so that the addition of these relatively small corrections can only lift the mass of the supermultiplet associated with the charge breaking flat directions by a small amount. Furthermore accurately measured quantities such as the Z width are sensitive to the presence of such particles in the low energy spectrum so that they actually constrain \(M_R\) to lie above \(10^9\) GeV in such theories. The presence of 4 massless \(SU(2)_L\) doublets down to scales \(\sim 1\) TeV has important implications for the solutions to the strong CP problem based on parity [17, 19].

The seesaw mechanism in this case is also quite distinct since it is “impure” i.e the potential contains terms of form

\[
\frac{m}{M} \Phi^\dagger \Delta_c \Delta + m^2 \Delta^2
\]

which gives a vev for \(\Delta\)

\[
\langle \Delta \rangle \approx \frac{\langle \Phi^2 \rangle}{\sqrt{mM}} \sim \frac{M^2_{R}}{M_{R}}
\]

exactly as in the non-supersymmetric case [9].
7 Conclusions

We have shown that Left Right Supersymmetric theories offer an attractive matrix in which the MSSM can be embedded with a gain of rationality as far as its problems with naturality and options vis vis neutrino masses are concerned. The theory makes definite testable predictions regarding the low energy spectrum in both the viable and consistent “minimal ” models analyzed. In particular the non-renormalizable minimal models predict low mass doubly charged supermultiplets that place lower bounds on the scale $M_R$. We have made extensive use of the powerful technique of holomorphic invariants while analyzing the potential and its possible flat directions. This permitted a clear and intelligible picture of the low energy theory to emerge inspite of the plethora of multiplets involved. Our work should provide a fresh impetus to analysis of GUTS with a left right symmetric intermediate symmetry and serve as reference point for future work on such theories.

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