Supersymmetry and large scale left-right symmetry

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We show that the low energy limit of the minimal supersymmetric left-right models is the supersymmetric standard model with an exact R parity. The theory predicts a number of light Higgs scalars and fermions with masses much below the $B-L$ and $SU(2)_R$ breaking scales. The nonrenormalizable version of the theory has a striking prediction of light doubly charged supermultiplets which may be accessible to experiment. Whereas in the renormalizable case the scale of parity breaking is undetermined, in the nonrenormalizable one it must be bigger than about $10^{10}-10^{12}$ GeV. The precise nature of the seesaw mechanism differs in the two versions, and has important implications for neutrino masses. [S0556-2821(98)08119-3]

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I. INTRODUCTION

One of the central issues, if not the main one, in the minimal supersymmetric standard model (MSSM) is what controls the strength of R-parity breaking. The suppression of (some or all) R-parity violating couplings in the MSSM is essential to avoid catastrophic proton decay rates, and determines the fate of the lightest supersymmetric particle (LSP). The most appealing rationale for an otherwise *ad hoc* discrete symmetry would be to have it as an automatic consequence of a gauge principle [1].

This is more than an aesthetic issue, for only gauge symmetries are protected against possible high scale violations such as, for example, those arising from quantum gravitational effects. Since in the MSSM the action of R parity on the superfields may be written as $R=(-1)^{3(B-L)}$ [2], theories with gauged $B-L$ may be regarded as the minimal framework to implement this idea [1,3,4]. $B-L$ symmetry is naturally, indeed ineluctably, incorporated in left-right symmetric theories, which provide an understanding of parity violation in nature [5–8]. A construction of a consistent supersymmetric left-right theory for generic values of the parity breaking scale (M_R) thus becomes essential. A considerable amount of work has been done on theories with low M_R [that is, $M_R \sim (1-10) M_W$] regarding the construction of the theory [9,10]. On the other hand, only recently have there been attempts to study the more realistic case of $M_R \gg M_W$ [11,12]. In this paper we provide a systematic study of minimal supersymmetric left-right theories (MSLRM) with an arbitrarily large scale of parity breaking and controllable R-parity violation.

This forces us to focus on the version of the theory with the conventional implementation of the seesaw mechanism [13–15]. By this we mean that the right handed neutrino Majorana mass arises at the renormalizable level. However, the following problem arises here: such a renormalizable theory with minimal Higgs content simply does not allow for any spontaneous symmetry breaking whatsoever [16]. There are two possible ways out of this impasse: one can either extend the Higgs sector [16,10,11] or allow for nonrenormalizable terms in the superpotential [17,12].

We first concentrate on the renormalizable version of the theory, which is both more conventional and simpler to analyze from the point of view of vacuum structure. We then apply the same techniques to the nonrenormalizable version, and compare the physical implications of both models. This should not imply that we take the nonrenormalizable version less seriously; this is the minimal theory in terms of the particle spectrum, and it provides the supersymmetric version of the minimal left-right theory.

Although in [11,12] the vacuum structure of these theories was studied, in this paper we present for the first time a complete and correct analysis of the lifting of the dangerous D-flat directions. Among other things, we learned that unless the sign of various soft mass terms is positive many of the flat directions would not be lifted. This is discussed at length.

Our main conclusion is that unless electromagnetic charge invariance is violated, R parity remains unbroken. More precisely, the effective low-energy theory becomes the MSSM with exact R parity. This is true in both versions of the theory. On the other hand, the precise nature of the seesaw mechanism does depend on whether the symmetry breaking is achieved through a renormalizable or nonrenormalizable superpotential.

Besides R-parity conservation, another important experimental signature of these theories is the presence of a number of charged Higgs supermultiplets whose masses are much below M_R. For a reasonable choice of parameters, they are expected to lie near the electroweak scale. In the nonrenormalizable version, these light supermultiplets include doubly charged ones [12], which makes this model especially interesting from the point of view of experiment. We present in this paper the complete particle spectrum for both models.

Another important consequence of our work lies in the possible grand unified or superstring extension of left-right symmetry. Namely, in the literature, one often assumes the extended survival principle for Higgs supermultiplets. By this one means that the particles which by symmetries are allowed to be heavy do indeed become so. The essential lesson of our paper is that this is completely wrong, since we find a plethora of light Higgs states which evade the above principle. This conclusion is not new; it was noticed already in the early papers on the supersymmetric SO(10) grand unified theory [18,19]. Unfortunately this fact is usually overlooked in the literature.

In the next section, we introduce the left-right supersymmetric model and discuss possible minimal choices for the Higgs sector. We also summarize the standard method for studying the structure of supersymmetric vacua, namely, the one based on the characterization of the flat directions of the supersymmetric potential by holomorphic gauge invariants of the chiral superfields. In Sec. III we apply this method to analyze the structure of the vacuum of the renormalizable model. In Sec. IV we use these results to prove that R parity (and therefore both baryon and lepton number) remains an exact symmetry of the low energy effective theory. We devote Sec. V to the study of the spectrum of the theory, paying special attention to the light sector. Section VI is where the nonrenormalizable model is taken up, and compared to the renormalizable version. Finally, we present our conclusions and outlook in Sec. VII.

II. SUPERSYMMETRIC LEFT-RIGHT THEORIES

The left-right symmetric model of gauge interactions treats fermions of opposite chiralities in a symmetric way by extending the standard model gauge group to $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. Thus the anomaly-free global $B-L$ symmetry of the standard model is inescapably promoted to a gauge symmetry in this picture. To obtain left-right symmetric Yukawa interactions that can give rise to the fermion masses it is necessary to promote the standard model Higgs boson to a bidoublet, and realistic fermion mass matrices require at least two bidoublets in the supersymmetric case.

In the supersymmetric version of this theory we thus supersymmetrize the gauge sector in the standard way and introduce three generations of quark and leptonic chiral superfields with the following transformation properties:

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

$$\Phi_i=(1,2,2^*,0), \quad (i=1,2), \tag{2}$$

where the numbers in the brackets denote the quantum numbers under $SU(3)_c$, $SU(2)_L$, $SU(2)_R$, and $U(1)_{B-L}$, respectively (generation indices are understood). In our convention,

$$L=\begin{pmatrix}\nu\\ e\end{pmatrix}, \quad L_c=\begin{pmatrix}\nu_c\\ e_c\end{pmatrix}, \tag{3}$$

so that $L\rightarrow U_L L$ under $SU(2)_L$, but $L_c\rightarrow U_R^* L_c$ under $SU(2)_R$, and similarly for quarks. Also, $\Phi\rightarrow U_L \Phi U_R^\dagger$.

The nontrivial question that now arises concerns the mechanism for the spontaneous violation of left-right (LR) symmetry, namely, the selection of a suitable minimal set of Higgs fields to break this symmetry. Furthermore, since the model necessarily includes a right handed neutrino, a mechanism to explain the observed suppression of neutrino mass (if any) is also necessary. Indeed one of the most appealing features of these models is precisely that they provide a natural (''seesaw'') mechanism to explain this suppression. In the nonsupersymmetric case, of the two simplest choices for the Higgs fields, namely, doublets or triplets with respect to $SU(2)_{L,R}$, only the latter allows one to realize the above scenario. However, the inclusion of nonrenormalizable operators in the action can be used to introduce small masses for the neutrino even in the doublet case [20].

In the supersymmetric case for simplicity we will not consider doublets. Since they carry one unit of $B-L$ charge, they are odd under R parity and thus the scales of R parity and LR symmetry breaking must coincide. Thus the smallness of R-parity violation in the doublet case can only be achieved *ad hoc* as in the MSSM. Therefore in this paper we choose to work with triplets. In the concluding section we will discuss the doublet alternative at greater length. The cancellation of $B-L$ anomalies requires the usual doubling of supermultiplets and thus the minimal choice of Higgs for LR breaking must include the multiplets below:

$$\Delta=(1,3,1,2), \quad \overline{\Delta}=(1,3,1,-2),$$
$$\Delta_c=(1,1,3,-2), \quad \overline{\Delta}_c=(1,1,3,2), \tag{4}$$

where $\Delta\rightarrow U_L \Delta U_L^\dagger$, but again $\Delta_c\rightarrow U_R^* \Delta_c U_R^T$. Left-right symmetry can be implemented in these theories either as a parity transformation

$$Q\leftrightarrow Q_c{}^*, \quad L\leftrightarrow L_c{}^*, \quad \Phi_i\leftrightarrow \Phi_i{}^\dagger,$$
$$\Delta\leftrightarrow \Delta_c{}^*, \quad \overline{\Delta}\leftrightarrow \overline{\Delta}_c^*, \tag{5}$$

or as a charge conjugation

$$Q\leftrightarrow Q_c, \quad L\leftrightarrow L_c, \quad \Phi_i\leftrightarrow \Phi_i{}^T,$$
$$\Delta\leftrightarrow \Delta_c, \quad \overline{\Delta}\leftrightarrow \overline{\Delta}_c. \tag{6}$$

The latter definition has an advantage from the point of view of grand unification, since it is an automatic gauge symmetry in SO(10). If one is not interested in the nature of CP violation, it makes no difference whatsoever which of the two definitions one uses. Strictly speaking, we do not even need this discrete symmetry in what follows, since in the supersymmetric limit all the minima are degenerate. However, the central challenge in left-right theories is the breaking of parity; so we include it in order to show that it can be done consistently and in accordance with experiment. For the sake of possible grand unification and transparency of our formula we choose the latter one.

With this set of multiplets, however, the most general renormalizable superpotential that one can write for the triplets which are to accomplish the breaking of parity is merely

$$W_{LR}=i\mathbf{f}(L^T\tau_2\Delta L+L_c^T\tau_2\Delta_c L_c)+m_\Delta(\text{Tr }\Delta\bar{\Delta}+\text{Tr }\Delta_c\bar{\Delta}_c). \tag{7}$$

Since we are considering the case where $M_R \gg M_S$, the supersymmetry (SUSY) breaking scale, the minimization of the potential is to be accomplished by setting the F terms for the chiral superfields and the D terms for the gauge fields to zero. Of course one must make sure that the soft supersymmetry breaking terms do not spoil the obtained hierarchy. This is not automatic for arbitrary soft terms but works in the case of gravity or gauge-mediated supersymmetry breaking. In these cases it is known that the soft terms imply a small perturbation of the SUSY preserving vacua. In what follows we assume the usual supergravity scenario of soft terms of order $m_{3/2}\simeq M_W$. Thus we have a small perturbation parameter M_W/M_R which in the leading approximation can be set to zero. Then it immediately follows that the vacuum expectation values (VEVs) of Δ,Δ_c must vanish identically while those of $\bar{\Delta},\bar{\Delta}_c$ are determined in terms of the VEVs of L,L_c respectively. Since the participation of squark VEVs in the symmetry breaking would lead to charge and color breaking (CCB) minima, we shall assume that their vanishing is ensured by suitable soft mass terms. Given that the squark VEVs vanish, the form of the the D term for the $B-L$ gauge field is then

$$D_{B-L}=-L^\dagger L-2\text{ Tr }\bar{\Delta}^\dagger\bar{\Delta}+L_c^\dagger L_c+2\text{ Tr }\Delta_c^\dagger\bar{\Delta}_c. \tag{8}$$

It is clear that vanishing of this D term requires that symmetry breaking in the left and right sectors must occur at the same scale.

To evade this difficulty we introduce an additional set of triplets,

$$\Omega=(1,3,1,0),\quad \Omega_c=(1,1,3,0), \tag{9}$$

where under left-right symmetry $\Omega\leftrightarrow\Omega_c$.

The inclusion of this set of multiplets has the additional attraction of allowing a separation of the scales where parity and $B-L$ symmetry are broken.

III. SYMMETRY BREAKING

We next turn to the minimization of the potential of the supersymmetric gauge theory introduced in Sec. II. The most general gauge-invariant superpotential that leads to a renormalizable action is

$$\begin{aligned}W_{LR}=&\mathbf{h}_l^{(i)}L^T\tau_2\Phi_i\tau_2L_c+\mathbf{h}_q^{(i)}Q^T\tau_2\Phi_i\tau_2Q_c+i\mathbf{f}(L^T\tau_2\Delta L\\&+L^{cT}\tau_2\Delta_cL_c)+m_\Delta(\text{Tr }\Delta\bar{\Delta}+\text{Tr }\Delta_c\bar{\Delta}_c)\\&+\frac{m_\Omega}{2}(\text{Tr }\Omega^2+\text{Tr }\Omega_c^2)+\frac{\mu_{ij}}{2}\text{ Tr }\tau_2\Phi_i^T\tau_2\Phi_j\\&+a(\text{Tr }\Delta\Omega\bar{\Delta}+\text{Tr }\Delta_c\Omega_c\bar{\Delta}_c)\\&+\alpha_{ij}(\text{Tr }\Omega\Phi_i\tau_2\Phi_j^T\tau_2+\text{Tr }\Omega_c\Phi_i^T\tau_2\Phi_j\tau_2),\end{aligned} \tag{10}$$

with $\mu_{ij}=\mu_{ji}$, $\alpha_{ij}=-\alpha_{ji}$, **f** and **h** are symmetric matrices, and generation and color indices are understood.

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. For example in a SU(5) SUSY gauge theory with a **24** of SU(5) as its only chiral multiplet one finds that the permissible vacua with a renormalizable potential are discrete [namely, the ones corresponding to the two maximal little groups of SU(5) besides the trivial one with the full SU(5) unbroken]. On the other hand, the introduction of additional matter multiplets such as a $\bar{\mathbf{5}}+\mathbf{10}$ anomaly-free pair quickly leads to a proliferation of flat directions emerging from these discrete minima.

In what follows we shall use an elegant and powerful method for characterizing the vacua of supersymmetric gauge theories [21–24]. The essence of this method is simply the following general result: (a) the space of field VEVs satisfying the D-flatness conditions $D_\alpha=0$ in a supersymmetric gauge theory is coordinatized by the independent holomorphic gauge invariants that may be formed from the chiral gauge multiplets in the theory. Further, (b) the space of field VEVs satisfying the D and F-flatness conditions is coordinatized by the holomorphic invariants left undetermined by the imposition of the conditions $F=0$ for each of the chiral multiplets in the theory.

The following simple example will serve to clarify the method. Consider a U(1) gauge theory with two chiral multiplets $\phi_\pm$ with gauge charges ±1. Then the condition $D=0$ requires only $|\phi_+|=|\phi_-|$. 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 $\phi_+\phi_-$. Now consider the effects of a superpotential $W=m\phi_+\phi_-$. The F-flatness condition now ensures that both VEVs vanish so that the D-flat manifold shrinks from the complex line parametrized by $c=\phi_+\phi_-$ to the single point $c=0$.

Thus, in principle, one should proceed by building all the holomorphic gauge invariants, establish which ones are left undetermined by the F-flatness conditions, and then discuss how the soft SUSY breaking terms may be used to lift those that are phenomenologically unacceptable—such as CCB directions. The analysis with the complete set of fields is, however, sufficiently complex to motivate a simplified approach to the problem.

On phenomenological grounds it is clear that the bidoublet and squark fields cannot obtain VEVs at the large scale. Therefore we omit them from our analysis of the symmetry breaking at the right handed scale. Even if they participate in flat directions running out of the parity breaking minimum, as long as their soft mass terms are taken (as usual) to be positive their VEVs at the high scale will vanish. On the other hand, since large VEVs for the sneutrinos in the right handed sector are *a priori* admissible, the L and L_c fields should be retained in the analysis.

The F-flatness conditions that follow from the the superpotential (10) are as follows:

$$F_{\bar{\Delta}} = m_\Delta \Delta + a(\Delta\Omega - \tfrac{1}{2}\ \mathrm{Tr}\ \Delta\Omega) = 0,$$

$$F_{\bar{\Delta}_c} = m_\Delta \Delta_c + a(\Delta_c\Omega_c - \tfrac{1}{2}\ \mathrm{Tr}\ \Delta_c\Omega_c) = 0,$$

$$F_{\Delta} = m_\Delta \bar{\Delta} + i\mathbf{f}LL^T\tau_2 + a(\Omega\bar{\Delta} - \tfrac{1}{2}\ \mathrm{Tr}\ \Omega\bar{\Delta}) = 0,$$

$$F_{\Delta_c} = m_\Delta \bar{\Delta}_c + i\mathbf{f}L_cL_c^T\tau_2 + a(\Omega_c\bar{\Delta}_c - \tfrac{1}{2}\ \mathrm{Tr}\ \Omega_c\bar{\Delta}_c) = 0,$$

$$F_{\Omega} = m_\Omega \Omega + a(\bar{\Delta}\Delta - \tfrac{1}{2}\ \mathrm{Tr}\ \bar{\Delta}\Delta) = 0,$$

$$F_{\Omega_c} = m_\Omega \Omega_c + a(\bar{\Delta}_c\Delta_c - \tfrac{1}{2}\ \mathrm{Tr}\ \bar{\Delta}_c\Delta_c) = 0,$$

$$F_L = 2i\mathbf{f}\tau_2\Delta L = 0,$$

$$F_{L_c} = 2i\mathbf{f}\tau_2\Delta_c L_c = 0. \tag{11}$$

In the above we have self-consistently set the bidoublet and squark fields, which must have zero VEVs at scales $\gg M_R$, to zero.

Multiplying the triplet equations by triplet fields and taking traces it immediately follows that

$$\mathrm{Tr}\ \Delta^2 = \mathrm{Tr}\ \Delta\Omega = \mathrm{Tr}\ \bar{\Delta}\Omega = 0,$$

$$m_\Delta\ \mathrm{Tr}\ \Delta\bar{\Delta} = m_\Omega\ \mathrm{Tr}\ \Omega^2 = a\ \mathrm{Tr}\ \Delta\bar{\Delta}\Omega,$$

$$\mathrm{Tr}\ \Delta\bar{\Delta}(a^2\ \mathrm{Tr}\ \Omega^2 - 2m_\Delta^2) = 0, \tag{12}$$

with corresponding equations *mutatis mutandis* in the right handed sector. Thus it is clear that in either sector all three triplets are zero or nonzero together. By choosing the branch where $\mathrm{Tr}\ \Omega_c^2 = 2m_\Delta^2/a^2$ but $\mathrm{Tr}\ \Omega^2 = 0$ we ensure that the triplet VEVs break $SU(2)_R$ but not $SU(2)_L$. The field content of the triplets is

$$\Delta_c = \begin{pmatrix} \delta_c^-/\sqrt{2} & \delta_c^{--} \\ \delta_c^0 & -\delta_c^-/\sqrt{2} \end{pmatrix}, \quad \bar{\Delta}_c = \begin{pmatrix} \bar{\delta}_c^+/\sqrt{2} & \bar{\delta}_c^0 \\ \bar{\delta}_c^{++} & -\bar{\delta}_c^+/\sqrt{2} \end{pmatrix},$$

$$\Omega_c = \begin{pmatrix} \omega_c^0/\sqrt{2} & \omega_{c1}^- \\ \omega_{c2}^+ & -\omega_c^0/\sqrt{2} \end{pmatrix}, \tag{13}$$

where superscripts denote electromagnetic charges

$$Q_{em} = T_{3L} + T_{3R} + \frac{B-L}{2}. \tag{14}$$

One can use the three parameters of the $SU(2)_R$ gauge freedom to set the diagonal elements of Δ_c to zero so that it takes the form

$$\langle \Delta_c \rangle = \begin{pmatrix} 0 & \langle \delta_c^{--} \rangle \\ \langle \delta_c^0 \rangle & 0 \end{pmatrix}. \tag{15}$$

Now Eq. (12) 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 nonsingular if the seesaw mechanism which keeps the sneutrino light is to operate. Then it immediately follows from the condition $F_{L_c} = 0$, namely,

$$2i\mathbf{f}_{ab}\begin{pmatrix} 0 & 0 \\ \langle \delta_c^0 \rangle & 0 \end{pmatrix}\begin{pmatrix} \nu_c \\ e_c \end{pmatrix}^b = 0, \tag{16}$$

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. In the Appendix we shall exhibit the flat directions out of the parity breaking triplet sector vacua associated with the slepton fields and show that charge is broken in both the left and right handed lepton sectors along these flat directions.

In the case with triplets alone we first list the $SU(2)_{L,R}$ invariants with their $B-L$ charges. The gauge invariants can then be generated from these by multiplying invariants whose charges sum to zero. The invariants are

$B-L$ charge	Invariant
4	$x_1 = \mathrm{Tr}\ \Delta^2$
2	$x_2 = \mathrm{Tr}\ \Omega\Delta$
0	$x_3 = \mathrm{Tr}\ \Delta\bar{\Delta},\ x_4 = \mathrm{Tr}\ \Omega^2,\ x_5 = \mathrm{Tr}\ \Omega\Delta\bar{\Delta}$
-2	$x_6 = \mathrm{Tr}\ \Omega\bar{\Delta}$
-4	$x_7 = \mathrm{Tr}\ \bar{\Delta}^2$

(17)

plus the corresponding invariants x_i^c, with opposite charges, built from the right handed fields. Without the leptons, besides conditions (12) we have also

$$\mathrm{Tr}\ \bar{\Delta}^2 = \mathrm{Tr}\ \bar{\Delta}_c^2 = 0. \tag{18}$$

Notice that this fixes the values of all the $x_i^{(c)}$ and hence, in fact, all the values of all the gauge invariants that one can form from the triplet fields. Any flat directions running out of the vacua allowed by minimizing the potential of the triplets alone (i.e., the trivial and equal left-right VEVs vacua which preserve parity and the two asymmetric vacua that violate it) must involve the fields we have omitted from the analysis. If these fields have zero VEVs at the high scale due to positive soft mass terms, then the vacua at the high scale are isolated and, in particular, the parity breaking vacuum described above is phenomenologically viable.

It is easy to use Eqs. (12) and (15) to see that the VEVs of Ω_c, Δ_c are also fixed to have the charge preserving form

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

$$\langle \Delta_c \rangle = \begin{pmatrix} 0 & 0 \\ d & 0 \end{pmatrix}, \quad \langle \bar{\Delta}_c \rangle = \begin{pmatrix} 0 & \bar{d} \\ 0 & 0 \end{pmatrix}. \tag{19}$$

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

$$w=-\frac{m_\Delta}{a}\equiv -M_R\,,\quad d=\overline{d}=\left(\frac{2m_\Delta m_\Omega}{a^2}\right)^{1/2}\equiv M_{BL}\,. \tag{20}$$

Notice an interesting property of Eqs. (20). If we wish to have $M_R \gg M_{BL}$, we need $m_\Delta \gg m_\Omega$, i.e., a sort of inverse hierarchy of the mass scales. Furthermore, this hierarchy has a highly suggestive geometric form $m_\Omega \simeq M_{BL}^2/M_R$. One cannot help speculating that m_Ω could be originated by soft supersymmetry breaking terms. Namely, in the absence of m_Ω the superpotential (10) has a global $\mathrm{U}(1)_R$ symmetry with the following R charges:

$$\begin{array}{rl} \Delta,\overline{\Delta},\Delta_c,\overline{\Delta}_c & 1\\ \Omega & 0\\ L,L_c,Q,Q_c & 1/2. \end{array} \tag{21}$$

The idea is very simple. We assume the above $\mathrm{U}(1)_R$ symmetry which implies $m_\Omega=0$ and which will be broken by the soft breaking terms.[1] In the gravity-mediated scenario of supersymmetry breaking, it is easy to see that m_Ω can be substituted by soft supersymmetry breaking terms. In other words

$$m_\Omega \sim m_{3/2} \sim M_W\,. \tag{22}$$

This implies $M_{BL}^2 \simeq M_W M_R$ and thus we have only one new scale M_{BL} or M_R. Of course, for $M_R \gg M_I \equiv \sqrt{M_{Pl}M_W}$ (where M_{Pl} is the Planck scale), nonrenormalizable terms could in principle induce a bigger value for $m_\Omega \sim M_R^2/M_{Pl}$. We should stress, though, that neutrino physics strongly suggests M_R to be smaller or of order M_I.

IV. *R*-PARITY BREAKING

As discussed in the Introduction, at scales where supersymmetry is valid, the invariance of the action in minimal left-right symmetric theories under R parity is enforced by $B-L$ gauge symmetry and supersymmetry. Thus the only possible source of R-parity violation at scales $\gg m_S \sim 1$ TeV is spontaneous: when a field with odd $3(B-L)$ develops a VEV. In the natural and minimal versions of the left-right supersymmetric theory that we have considered here the only electrically neutral fields that can violate R parity spontaneously without breaking charge are the sneutrinos in the two sectors. On the other hand, we have seen that the right handed sneutrino VEV is strictly zero at the parity breaking minimum when working to leading order in the ratio $m_S/M_{BL} \sim M_W/M_{BL}$. What happens when the soft terms are switched on? As we said before, we assume the supergravity scenario of spontaneous supersymmetry breaking and the resulting soft terms. These terms are known to induce only small shifts of the order M_W/M_{BL}. Since the right handed sneutrino mass is large, i.e., of the order of $M_{BL} \gg M_W$, small corrections clearly cannot shift its VEV at all, and it will remain at the origin. This, however, is not necessarily true of the left handed sneutrino, whose mass is at most of the order of M_W. We discuss its case carefully below.

Any R-parity breaking in the right handed sector thus necessarily involves breaking of charge, at least at high scales $O(M_R)$. In fact, in the Appendix we show that the leptonic flat directions running out from the parity breaking vacuum necessarily violate charge in both left and right sectors. Thus we regard it as physically well motivated to assume that soft SUSY breaking mass terms must be such as to forbid excursions along the R-parity and charge violating flat directions, just as they must protect color.

Thus the effective theory below the scales M_R, M_{BL} will simply be the MSSM (with the addition of some new particle states: see next section) with R parity operative. It will therefore also possess an effective global lepton and baryon number symmetry.

Nevertheless, one may worry that the effects of the running of the coupling constants and masses may be such as to induce VEVs for the sneutrino fields. Such a thing happens in the MSSM when the Higgs bidoublet mass squared, although positive at high scales, suffers large negative corrections due to its strong Yukawa coupling to the top quark, developing a VEV at scales $O(M_W)$ [4,26]. The situation with sneutrino fields is, however, quite different since none of the leptonic Yukawa couplings are large. Moreover, even if the leptonic soft masses were small enough to be overcome by renormalization effects in going from M_{BL} down to the electroweak (EW) scale, a sneutrino VEV could develop only in the left handed sector since the right handed sneutrinos have superheavy masses. The global lepton number symmetry of the effective theory below M_{BL} then implies that a left handed sneutrino VEV would in fact lead to a Goldstone boson: the (doublet) Majoron [27]. Such a massless doublet Majoron is coupled strongly to the Z boson and leads to a large contribution to its width which is ruled out by experiment.

The one remaining possibility is that the violation of lepton number by the left sneutrino VEV, in combination with the electroweak VEVs, may trigger a VEV for the right handed sneutrino. Such explicit violation of the effective global Lepton number symmetry would in turn give a mass to the Majoron. If this mass were sufficiently large ($>M_Z/2$), then the contribution of this state to the Z width would be suppressed. To see how ν^c might get a VEV consider the allowed trilinear soft term in the potential

$$\Delta V_{\mathrm{soft}}=\cdots+m_S L^T\tau_2\Phi_i\tau_2 L_c+\cdots\,. \tag{23}$$

Once ν and Φ develop VEVs, this implies a linear term in the potential for the right handed sneutrino which will thus get an expectation value

[1]We thank Gia Dvali for bringing up this point. For an original application of this idea, see [25].

TABLE I. Mass spectrum for the Higgs supermultiplets in the renormalizable model.

State	Mass
δ_c^{++}, $\overline{\delta}_c^{++}$	$2aM_R$
$\delta_c^{+}-\frac{M_{BL}}{\sqrt{2}M_R}\omega_{c1}^{+}$	$aM_R\left[1+\frac{1}{2}\left(\frac{M_{BL}}{M_R}\right)^2\right]$
$\overline{\delta}_c^{+}-\frac{M_{BL}}{\sqrt{2}M_R}\omega_{c2}^{+}$	$aM_R\left[1+\frac{1}{2}\left(\frac{M_{BL}}{M_R}\right)^2\right]$
$\left(\omega_{c1}^{+}+\frac{M_{BL}}{\sqrt{2}M_R}\delta_c^{+}\right)-\left(\omega_{c2}^{+}+\frac{M_{BL}}{\sqrt{2}M_R}\overline{\delta}_c^{+}\right)$	$2gM_R\left[1+\frac{1}{4}\left(\frac{M_{BL}}{M_R}\right)^2\right]$
$\left(\omega_c^0+\frac{(\delta_c^0+\overline{\delta}_c^0)}{\sqrt{2}}\right)$	$aM_{BL}\left(1+\frac{M_{BL}}{4M_R}\right)$
$\left(\omega_c^0-\frac{(\delta_c^0+\overline{\delta}_c^0)}{\sqrt{2}}\right)$	$aM_{BL}\left(1-\frac{M_{BL}}{4M_R}\right)$
$\mathrm{Re}(\delta_c^0-\overline{\delta}_c^0)$	$2\sqrt{g^2+g'^2}M_{BL}$
Δ, $\overline{\Delta}$	aM_R
Ω	$aM_{BL}^2/2M_R$
H, $\overline{H}$	~ 0
H', $\overline{H}'$	$\sim M_R$

$$\langle \nu^c\rangle=\frac{m_S M_W\langle \nu\rangle}{M_{BL}^2}. \tag{24}$$

This would lead to effective R-parity and global lepton number violating terms of the form $m_\epsilon^2 LH$ where

$$m_\epsilon^2=\frac{m_S^2 M_W\langle \nu\rangle}{M_{BL}^2}. \tag{25}$$

Then the ''Majoron'' would get a mass squared of order

$$m_J^2\simeq m_\epsilon^2\frac{m_S}{\langle \nu\rangle}\simeq\frac{m_S^3 M_W}{M_{BL}^2}. \tag{26}$$

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 is a corner of parameter space we do not consider in this paper. We conclude that within the present scenario the bounds on the Z width rule out the possibility of R-parity violation due to sneutrino VEVs. In sum, *the low energy effective theory of the minimal supersymmetric left-right model is the MSSM (with some additional particle states) with strictly unbroken R parity, and the LSP is stable.*

V. MASS SPECTRUM

As we have seen, the symmetry breaking takes place in two stages. At a large scale $M_R=m_\Delta/a$, $SU(2)_R$ is broken down to $U(1)_R$ by the VEV of Ω_c. Later the VEVs of $\Delta_c,\overline{\Delta}_c$ are turned on at $M_{BL}=\sqrt{2m_\Delta m_\Omega}/a$, breaking $U(1)_R\times U(1)_{B-L}$ to $U(1)_Y$. However, a third scale appears in the superpotential, $m_\Omega=M_{BL}^2/M_R$. Let us now examine the mass spectrum, and see which scales are involved.

A. Higgs sector

We begin with the masses of the triplets. The results (for $M_R\gg M_{BL}$) are summarized in Table I.

As could be expected, almost all the particles get a mass at the scale M_R. In the right handed sector, the doubly charged particles $\delta_c^{++},\overline{\delta}_c^{++}$ do so, simply through their explicit m_Δ terms in the superpotential. The VEV of Ω_c will contribute with mass terms for the rest of the charged particles, giving to all of them a large mass M_R. However, the neutral particle masses correspond to the (in principle) lower scale M_{BL}. The reader can check the manifestation of the super-Higgs mechanism: we have states with masses equal to the charged (W_R) and neutral (Z_R) gauge boson masses:

$$M^2(W_R)=4g^2M_R^2+2g^2M_{BL}^2,$$
$$M^2(Z_R)=4(g^2+g'^2)M_{BL}^2. \tag{27}$$

In the left handed sector, on the other hand, masses come directly through the explicit terms in the superpotential. Δ and $\overline{\Delta}$ have a large mass of order M_R. But the mass of Ω is related to the third mass scale we mentioned above, M_{BL}^2/M_R. This is the most interesting prediction of the model: a complete $SU(2)_L$ triplet of scalars and fermions, at a relatively low mass scale, which could be accessible to future experiments. Notice that for the analysis in the previous section to be valid, $m_\Omega\simeq M_{BL}^2/M_R$ should not be below the scale of the soft supersymmetry breaking terms. In fact, as we have argued in Sec. III, the natural scale for m_Ω is of order $m_{3/2}\sim M_W$, at least in the physically motivated picture with the ratio M_{BL}^2/M_R generated dynamically. The dependence of these new light states on M_R is noteworthy. Both M_{BL} and M_R are likely to be large enough to be out of a direct experimental search. However, M_{BL} can be indirectly probed through the usual seesaw-induced neutrino mass (see below), and thus improving the experimental limits on new non-MSSM states will actually set *upper* limits on M_R. Finally, this indirect probe and the direct search for Ω may provide a crucial test of the consistency of the theory.

Once $SU(2)_R$ is broken, the bidoublets Φ_1,Φ_2 get split into four $SU(2)_L$ doublets, and as usual one fine-tuning is necessary to keep one pair of them light. Namely, when Ω_c gets a VEV the mass terms for the bidoublets in the superpotential become

$$W(m_\Phi)=\frac{\mu_{ij}}{2}\,\mathrm{Tr}\;\tau_2\Phi_i^T\tau_2\Phi_j+M_R\alpha_{ij}\;\mathrm{Tr}\;\tau_3\Phi_i^T\tau_2\Phi_j\tau_2. \tag{28}$$

Now, writing the bidoublets in terms of $SU(2)_L$ doublets $H_i,\overline{H}_i$ as

$$\Phi_i=(H_i,\overline{H}_i)\equiv\begin{pmatrix}\phi_i^0 & \overline{\phi}_i^+\\ -\phi_i^- & \overline{\phi}_i^0\end{pmatrix}, \tag{29}$$

the mass terms are seen to correspond to

$$W(m_{H,\overline{H}})=\mu_{ij}H_i\overline{H}_j+\alpha_{ij}M_R(H_i\overline{H}_j-\overline{H}_iH_j). \tag{30}$$

With the fine-tuning condition

$$\mu_{11}\mu_{22}-(\mu_{12}^2-\alpha_{12}^2M_R)\simeq 0, \tag{31}$$

the bidoublets get split into two heavy left handed doublets $H',\overline{H}'$ with masses $\sim M_R$, and two MSSM Higgs doublets $H,\overline{H}$. A comment is in order. The fine-tuning of Higgsino masses does not automatically guarantee the same for the Higgs scalar masses. However, in the context of the supergravity scenario (and gauge mediation) it is well known that corrections are at most of order $m_{3/2}\simeq M_W$.

B. Neutrino mass

Another distinct prediction of this model is that the seesaw mechanism takes its canonical form. By canonical form we mean (in the single-generation case)

$$\begin{pmatrix} 0 & m_D \\ m_D & M \end{pmatrix}, \tag{32}$$

where m_D is the usual Dirac neutrino mass and M is the large Majorana mass of the right handed neutrino. The mass M is induced through the VEV of Δ_c, and thus $M\sim M_R$. Interestingly enough, the form (32) is hard to achieve in nonsupersymmetric theories, for Δ in general acquires a small VEV after electroweak breaking. Its origin is a term in the potential linear in Δ, i.e., of the form $\Delta\Phi^2\Delta_c$ [28]. In other words, the (1,1) mass element in general is not zero, and this has important implications for light neutrino mass spectrum. Namely, the light neutrino mass is given by

$$m_\nu=\frac{m_D^2}{M_R}. \tag{33}$$

In the supersymmetric version we are considering, though, the form (32) is exact up to order $1/M_{Pl}$. Simply, at the renormalizable level there are no terms linear in Δ in the potential. If one admits nonrenormalizable terms cut off by the Planck scale, along the lines of [28], one finds that at electroweak breaking Δ gets a VEV of order $(M_W^2M_{BL})/(M_RM_{Pl})$. The (1,1) element in Eq. (32) is thus suppressed with respect to the usual seesaw mechanism by $M_{BL}^2/(M_RM_{Pl})$, and is completely negligible for physics much below the Planck scale.

VI. NONRENORMALIZABLE MODEL

As has been pointed out in [12], it is possible to break parity even with just the minimal field content given in Eqs. (1)–(4), if one allows for nonrenormalizable interactions, suppressed by inverse powers of a large scale M. Including dimension-4 operators, the most general superpotential becomes

$$\begin{aligned}
W_{nr}=&\,m(\mathrm{Tr}\,\Delta\overline{\Delta}+\mathrm{Tr}\,\Delta_c\overline{\Delta}_c)+i\mathbf{f}(L^T\tau_2\Delta L+L_c^T\tau_2\Delta_cL_c)+\frac{a}{2M}[(\mathrm{Tr}\,\Delta\overline{\Delta})^2+(\mathrm{Tr}\,\Delta_c\overline{\Delta}_c)^2]+\frac{c}{M}\,\mathrm{Tr}\,\Delta\overline{\Delta}\,\mathrm{Tr}\,\Delta_c\overline{\Delta}_c\\
&+\frac{b}{2M}[\mathrm{Tr}\,\Delta^2\,\mathrm{Tr}\,\overline{\Delta}^2+\mathrm{Tr}\,\Delta_c^2\,\mathrm{Tr}\,\overline{\Delta}_c^2]+\frac{1}{M}[d_1\,\mathrm{Tr}\,\Delta^2\,\mathrm{Tr}\,\Delta_c^2+d_2\,\mathrm{Tr}\,\overline{\Delta}^2\,\mathrm{Tr}\,\overline{\Delta}_c^2]\\
&+\mathbf{h}_l^iL^T\tau_2\Phi_i\tau_2L_c+\mathbf{h}_q^iQ^T\tau_2\Phi_i\tau_2Q_c+\mu_{ij}\,\mathrm{Tr}\,\tau_2\Phi_i^T\tau_2\Phi_j+\frac{\lambda_{ijkl}}{M}\,\mathrm{Tr}\,\tau_2\Phi_i^T\tau_2\Phi_j\,\mathrm{Tr}\,\tau_2\Phi_k^T\tau_2\Phi_l\\
&+\frac{\alpha_{ij}}{M}(\mathrm{Tr}\,\overline{\Delta}\Delta\Phi_i\tau_2\Phi_j^T\tau_2+\mathrm{Tr}\,\overline{\Delta}_c\Delta_c\Phi_i^T\tau_2\Phi_j\tau_2)+\frac{\beta_{ij}}{M}\,\mathrm{Tr}\,\tau_2\Phi_i^T\tau_2\Phi_j[\mathrm{Tr}\,\Delta\overline{\Delta}+\mathrm{Tr}\,\Delta_c\overline{\Delta}_c]\\
&+\frac{\eta_{ij}}{M}\,\mathrm{Tr}\,\Phi_i\tau_2\Delta_c\Phi_j^T\tau_2\Delta+\frac{\overline{\eta}_{ij}}{M}\,\mathrm{Tr}\,\Phi_i\tau_2\overline{\Delta}_c\Phi_j^T\tau_2\overline{\Delta}+\frac{\mathbf{k}_{ql}}{M}Q^T\tau_2LQ_c^T\tau_2L_c+\frac{\mathbf{k}_{qq}}{M}Q^T\tau_2QQ_c^T\tau_2Q_c\\
&+\frac{\mathbf{k}_{ll}}{M}L^T\tau_2LL_c^T\tau_2L_c+\frac{\mathbf{j}}{M}[Q^T\tau_2QQ^T\tau_2L+Q_c^T\tau_2Q_cQ_c^T\tau_2L_c].
\end{aligned} \tag{34}$$

Of course, not all of the terms above play an equally important role. If a certain renormalizable interaction is already present in the potential [as is the case, for example, with the term $(L_c^\dagger L_c)^2$], one can safely neglect small corrections of order $1/M$. It is only when there are no cubic or quartic couplings that we *must* keep the nonrenormalizable ones.

As in the renormalizable case, we assume that the soft terms are such as to drive the VEVs of the squark and bidoublet fields to zero. With this assumption it follows, exactly as in the renormalizable case, that for a parity breaking but charge preserving VEV of Δ_c the VEV of the right handed sneutrinos is necessarily zero. Thus any R-parity breaking due to L_c getting a VEV at scales $\gg M_S$ would necessarily break charge. Therefore we assume that the soft terms forbid any VEVs for the L_c fields. We are therefore left with the problem of analyzing the F-flatness conditions for the four triplets $\Delta,\overline{\Delta},\Delta_c,\overline{\Delta}_c$ to determine whether or not they admit a charge preserving but parity breaking isolated minimum.

The F terms for the triplets now read

$$F_\Delta=\left(m+\frac{a}{M}\,\mathrm{Tr}\,\Delta\overline{\Delta}+\frac{c}{M}\,\mathrm{Tr}\,\Delta_c\overline{\Delta}_c\right)\overline{\Delta}+\left(\frac{b}{M}\,\mathrm{Tr}\,\overline{\Delta}^2+\frac{d_1}{M}\,\mathrm{Tr}\,\Delta_c^2\right)\Delta=0, \tag{35}$$

$$F_{\overline{\Delta}}=\left(m+\frac{a}{M}\,\mathrm{Tr}\,\Delta\overline{\Delta}+\frac{c}{M}\,\mathrm{Tr}\,\Delta_c\overline{\Delta}_c\right)\Delta+\left(\frac{b}{M}\,\mathrm{Tr}\,\Delta^2+\frac{d_2}{M}\,\mathrm{Tr}\,\overline{\Delta}_c^2\right)\overline{\Delta}=0,$$

$$F_{\overline{\Delta}_c}=\left(m+\frac{a}{M}\,\mathrm{Tr}\,\Delta_c\overline{\Delta}_c+\frac{c}{M}\,\mathrm{Tr}\,\Delta\overline{\Delta}\right)\Delta_c+\left(\frac{b^*}{M}\,\mathrm{Tr}\,\Delta_c^2+\frac{d_2}{M}\,\mathrm{Tr}\,\overline{\Delta}^2\right)\overline{\Delta}_c=0,$$

$$F_{\Delta_c}=\left(m+\frac{a}{M}\,\mathrm{Tr}\,\Delta_c\overline{\Delta}_c+\frac{c}{M}\,\mathrm{Tr}\,\Delta\overline{\Delta}\right)\overline{\Delta}_c+\left(\frac{b}{M}\,\mathrm{Tr}\,\overline{\Delta}_c^2+\frac{d_1}{M}\,\mathrm{Tr}\,\Delta^2\right)\Delta_c=0.$$

In the renormalizable model with an extra triplet Ω, the left and right handed sectors are completely decoupled, and the potential admits two discrete minima in each sector. The trivial one is chosen by the left handed sector, while the right handed triplets reside in the nontrivial one, which is charge and color preserving. In the nonrenormalizable case, on the contrary, the two sectors are coupled in the F equations. As we now show, the physically relevant minimum, for which the VEV of the left handed triplets vanish and that of the right handed triplets respects charge, does not involve any flat direction. This guarantees the stability of the vacuum against the presence of soft terms.

The $SU(2)_{L,R}$ invariants are now just

$B-L$ charge	Invariant
4	$y_1=\mathrm{Tr}\,\Delta^2$
0	$y_2=\mathrm{Tr}\,\Delta\overline{\Delta}$
-4	$y_3=\mathrm{Tr}\,\overline{\Delta}^2$

(36)

together with their right handed counterparts with opposite $B-L$ charges. Now, the equations for F_Δ and $F_{\overline{\Delta}}$ give

$$\left[\left(m+\frac{a}{M}\,\mathrm{Tr}\,\Delta\overline{\Delta}+\frac{c}{M}\,\mathrm{Tr}\,\Delta_c\overline{\Delta}_c\right)^2-\left(\frac{b}{M}\,\mathrm{Tr}\,\overline{\Delta}^2+\frac{d_1}{M}\,\mathrm{Tr}\,\Delta_c^2\right)\times\left(\frac{b}{M}\,\mathrm{Tr}\,\Delta^2+\frac{d_2}{M}\,\mathrm{Tr}\,\overline{\Delta}_c^2\right)\right]\mathrm{Tr}\,\Delta\overline{\Delta}=0 \tag{37}$$

and the corresponding equation for the right handed sector.

Out of the two branches allowed by Eq. (37) we focus on the one specified by $y_2=0$. With this choice the conditions

$$(by_3+d_1y_1^c)y_1=0;\quad (by_1+d_2y_3^c)y_3=0 \tag{38}$$

follow from $F_\Delta=F_{\overline{\Delta}}=0$. These equations are both satisfied on the branch specified by $y_1=y_2=y_3=0$ and it then follows that the only gauge invariants which might remain undetermined, and therefore allow a flat direction out of this branch, are y_2^c and $y_1^cy_3^c$. We emphasize we have not chosen a point on a flat direction but a branch of solutions of the field equations specified by the conditions $y_i=0$. The equations in the right handed sector are now simply

$$(mM+(a+b)y_2^c)y_1^c=0, \tag{39}$$

$$(mM+ay_2^c)y_2^c+by_1^cy_3^c=0. \tag{40}$$

The two branches of solutions of Eq. (39) are (i) $y_1^c=0$ and (ii) $y_2^c=-(mM)/(a+b)$. It easily follows from Eq. (40) that apart from the trivial solution where all invariants vanish the other possibilities are

$$\text{(a)}\quad y_1^c=y_3^c=0,\quad y_2^c=-\frac{mM}{a}, \tag{41}$$

$$\text{(b)}\quad y_2^c=-(mM)/(a+b),\quad y_1^cy_3^c=\frac{M^2m^2}{(a+b)^2}. \tag{42}$$

Solution (a) is equivalent to the one found in the renormalizable case: it is charge preserving and breaks parity. On the other hand, using $SU(2)_R$ invariance to put the diagonal elements of Δ_c to zero it immediately follows that solution (b) implies breaking of charge.

Thus in this model the triplets get a VEV just as in the renormalizable one, Eq. (19), where

$$d=\overline{d}=\sqrt{-\frac{mM}{a}}. \tag{43}$$

A. Mass spectrum

It can be seen that the nonrenormalizable model has distinct features. Symmetry breaking occurs in one stage, at a scale

$$M_R\equiv\sqrt{-\frac{Mm}{a}}. \tag{44}$$

In the analysis, it has been fundamental to assume m not to be smaller than the supersymmetry breaking scale. With the large scale M of order M_{Pl}, and $m\gtrsim 1$ TeV, the parity

TABLE II. Mass spectrum for the Higgs supermultiplets in the nonrenormalizable model.

State	Mass
$\delta_c^+ - \overline{\delta}_c^+$	$\sqrt{2} g M_R$
$\mathrm{Re}(\delta_c^0 - \overline{\delta}_c^0)$	$2\sqrt{g^2+g'^2} M_R$
$\delta_c^0 + \overline{\delta}_c^0$	$2aM_R^2/M$
δ_c^{++}, $\overline{\delta}_c^{++}$	$2bM_R^2/M$
Δ, $\overline{\Delta}$	$(a-c)M_R^2/M$
H, $\overline{H}$	~ 0
H', $\overline{H}'$	$\sim M_R^2/M$

breaking scale becomes $M_R \gtrsim 10^{11}$ GeV. However, the mass spectrum of the theory in fact involves two scales. Some of the particles in the Higgs multiplets will remain ''light'' after $\mathrm{SU}(2)_R$ breaking, with masses of order $m \sim M_R^2/M$. Thus as in the renormalizable case we have a very interesting phenomenology involving light charged, doubly charged, and neutral supermultiplets as shown in Table II. It is appealing also here to invoke soft supersymmetry breaking terms to dynamically generate m, in which case it would be of the order of the electroweak scale.

Notice again the super-Higgs mechanism being operative, since supersymmetry is not broken. In this case, only the states that belong to super-Higgs multiplets get a mass of order M_R; the rest of the states can only get a Planck-suppressed mass. Of particular interest is the presence of two sets of light doubly charged scalars and fermions ($\delta^{++}, \overline{\delta}^{++}$), and two full $\mathrm{SU}(2)_L$ triplets (Δ and $\overline{\Delta}$). They could have masses as low as the supersymmetry breaking scale. In contrast to the renormalizable case, the search for these particles sets a *lower* limit on M_R.

On the other hand, the bidoublet splitting proceeds in an equivalent way as in the renormalizable model, with the higher order interactions effectively playing the role of the Ω field. The mass terms for Φ_i come from

$$W(m_\Phi) = \frac{\mu_{ij}}{2} \mathrm{Tr}\ \tau_2 \Phi_i^T \tau_2 \Phi_j + \frac{\alpha_{ij}}{M} \mathrm{Tr}\langle \overline{\Delta}_c \Delta_c \rangle \Phi_i^T \tau_2 \Phi_j \tau_2 + \frac{\beta_{ij}}{M} \mathrm{Tr}\ \tau_2 \Phi_i^T \tau_2 \Phi_j\ \mathrm{Tr}\langle \overline{\Delta}_c \Delta_c \rangle \tag{45}$$

or

$$W(m_\Phi) = \left(\frac{\mu_{ij}}{2} + \frac{M_R^2}{M} \beta_{ij} \right) \mathrm{Tr}\ \tau_2 \Phi_i^T \tau_2 \Phi_j + \frac{M_R^2}{2M} \alpha_{ij}\ \mathrm{Tr}(1+\tau_3) \Phi_i^T \tau_2 \Phi_j \tau_2. \tag{46}$$

In terms of the SU(2) doublets $H_i, \overline{H}_i$ we have now

$$W(m_{H,\overline{H}}) = \mu'_{ij} H_i \overline{H}_j + \frac{M_R^2}{2M} \alpha_{ij} (H_i \overline{H}_j - \overline{H}_i H_j), \tag{47}$$

where $\mu'_{ij} = \mu_{ij} + M_R^2 (2\beta_{ij} + \alpha_{ij})/M$. The crucial difference with the previous model can be seen immediately: one will now fine-tune

$$\mu'_{11} \mu'_{22} - \left({\mu'_{12}}^2 - \frac{M_R^2}{2M} \alpha_{12}^2 \right) \simeq 0, \tag{48}$$

so that the two heavy doublets will not have masses at the large scale M_R but at $m = M_R^2/M \gtrsim 1$ TeV. Just as in the renormalizable case, it is appealing to have a small scale m generated softly of order M_W. In that case, one could imagine the μ terms being large and being fine-tuned among each other. However, this does not work, since there would be no splitting of the doublets within the bidoublets, but rather just splitting of the two bidoublets in the L-R symmetric manner (one complete bidoublet would remain light, and the other one heavy). This in the usual manner would imply the vanishing of quark mixing angles. In other words $\mu_{ij} \simeq m$, and both generated by the soft supersymmetry breaking terms. The dynamical generation of mass terms of the theory is a rather appealing and long pursued scenario, and furthermore in this case there is then no need for the fine-tuning of large and independent mass scales. The clear phenomenological test of this idea is the necessary existence of two more weak doublet supermultiplets at the experimentally accessible energy scale.

Therefore, although parity is broken at M_R, the model does not reduce to the MSSM until much later, at the lower scale m. This has important consequences for the solution of the strong CP problem. With four Higgs doublets with masses much below M_R, the running of the Yukawa matrices quickly generates a sizable strong CP phase [29]. This forces any viable solution to the strong CP problem based on parity in supersymmetric models to have M_R of the order of the weak scale [30].

B. Neutrino mass

Another important difference has to do with neutrino mass. We have seen that in the renormalizable case the see-saw mechanism takes its canonical form (33). Now the situation is completely different, and resembles the nonsupersymmetric case. The nonrenormalizable terms now are essential, since they provide the interaction terms and determine the scale M_R. One finds in the potential the relevant terms (written schematically)

$$\frac{m}{M} \Phi^2 \overline{\Delta}_c \Delta + m^2 \Delta^2, \tag{49}$$

which gives a VEV for Δ,

$$\langle \Delta \rangle = \frac{\langle \Phi^2 \rangle}{\sqrt{mM}} \sim \frac{M_W^2}{M_R}, \tag{50}$$

exactly as in the nonsupersymmetric case [28]. Once again, the two models lead to different phenomenological implications.

VII. DISCUSSION AND OUTLOOK

In this paper we have offered a complete analysis of symmetry breaking in minimal supersymmetric left-right models with a scale M_R much above the electroweak scale. We were led by the requirement that R parity not be introduced *ad hoc* and that its breaking be controlled without fine-tuning. This rules out the possibility of using doublet Higgs fields to break parity, and leads naturally to triplets, further motivated by the seesaw mechanism. By minimal, then, we mean theories with the minimal Higgs sector needed to achieve the complete symmetry breaking down to $\mathrm{SU}(3)_C\times\mathrm{U}(1)_{em}$. This in practice means the following.

(i) At the renormalizable level, one needs to introduce a new physical scale associated with $B-L$ breaking. This is achieved through the introduction of an additional pair of $B-L$ neutral triplets. The alternative possibility of utilizing parity-odd singlets does not work [10].

(ii) If one accepts nonrenormalizable terms, the Higgs sector consists just of the supersymmetric extension of the usual states needed to generate consistent fermion and gauge boson masses.

In both cases, the requirement of minimality meant the exclusion of further singlet states. The central theoretical result of our analysis is the proof that the physically acceptable minimum does not lie on a flat direction. Being an isolated point, with a large barrier separating it from other (nonphysical) minima, there is no danger that it will not be stable on a cosmological scale.

The two versions of the theory have in common two extremely important characteristics: (a) R parity never gets broken, and (b) the low-energy effective theory, besides the usual MSSM states, necessarily contains light charged or doubly charged superfield multiplets.

What is different is the nature of the seesaw mechanism and of the precise spectrum of light states. Whereas in the renormalizable version the seesaw mechanism takes its canonical form, in the nonrenormalizable case, the situation parallels the one in the nonsupersymmetric left-right models. This in general leads to different neutrino mass spectra, and is experimentally distinguishable.

Other important differences arise in the Higgs boson mass spectrum, as displayed in Tables I and II. In both cases there are two relevant physical mass scales above M_W. In the renormalizable case they are M_{BL} and M_R (and we discuss the physically interesting case $M_{BL}\ll M_R$). The new light supermultiplet is the left handed Higgs triplet Ω, with a mass of order M_{BL}^2/M_R which could naturally lie near the weak scale [11]. In particular it does so in the case of a dynamically generated ratio of the two new physical scales M_{BL} and M_R. In other words, m_Ω is the result of soft terms, which break both supersymmetry and an otherwise automatic continuous R symmetry. The light particles comprise both neutral and single-charged scalars and fermions.

In the nonrenormalizable model, the scales M_{BL} and M_R coincide, but there is a new high scale M_{Pl}. In this case there is a plethora of new light states with a mass M_R^2/M_{Pl}, which among other fields include the experimentally very interesting doubly charged scalars and fermions [12]. This is the crucial difference between the two theories. The doubly charged states are of utmost interest due to their spectacular experimental signatures. In the past there have been a number of papers devoted to the phenomenological implications of supersymmetric left-right theories with low M_R [31–36], in which the doubly charged states are discussed. Most of this analysis carries on to our case.

Light doubly charged particles continue to exist even if one adds an arbitrary number of gauge singlets [37]. In this paper, however, we do not include singlets, for the whole point of our work has been minimality. The nonrenormalizable version is obviously the minimal supersymmetric left-right theory. However, since renormalizability provides the cornerstone of field theories, the version with an intermediate $B-L$ scale can also be considered minimal.

Another, equally important, implication of the existence of the new light supermultiplets is its impact on the running of the gauge couplings. The analysis done in the past often relied on the survival principle, assuming that all states which by symmetry are allowed to be heavy become so. This is manifestly wrong, for, as we have shown, there are a number of light scalars and fermions whose existence defies this principle. Clearly, a new analysis of unification is required. It is not enough to take the result of our paper, for there may be additional light states which survive the large scale breaking of the underlying grand unified theory. This has already been noted in the early works on the SO(10) grand unification of this theory [18,19], but no running has been performed in these papers. In view of this, it is not clear at all that these models can be successfully unified, but we reserve the final judgment for the future.

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APPENDIX

In this appendix we analyze the flat directions of the superpotential when the slepton fields in both sectors are retained in the analysis. Since doublets must occur bilinearly in all invariants it is convenient to define composite singlet ($\sigma_{ab}=-\sigma_{ba}$) and triplet fields ($\overline{\Delta}'_{ab}=\overline{\Delta}'_{ba}$) by

$$L_a\tilde{L}_b=\frac{1}{2}\sigma_{ab}+\frac{1}{2}\,\overline{\Delta}'_{ab} \tag{A1}$$

and similarly in the right handed sector. They obey Fierzing constraints like

$$\sigma_{ab}\sigma_{cd}=\frac{1}{2}\sigma_{ad}\sigma_{bc}+\frac{1}{4}\mathrm{Tr}(\overline{\Delta}'_{ad}\overline{\Delta}'_{bc}). \tag{A2}$$

Then in addition to the $\mathrm{SU}(2)_{L,R}$ invariants listed in Eq. (17) one can form the additional invariants (not all independent)

$B-L$ charge	Invariant
-4	$x_8^{abcd}=\mathrm{Tr}\,\overline{\Delta}'_{ab}\overline{\Delta}'_{cd}$
-4	$x_9^{ab}=\mathrm{Tr}\,\overline{\Delta\Delta}'_{ab}$
0	$x_{10}^{ab}=\mathrm{Tr}\,\Delta\overline{\Delta}'_{ab}$
-2	$x_{11}^{ab}=\mathrm{Tr}\,\Omega\overline{\Delta}'_{ab}$
-2	$x_{12}^{ab}=\mathrm{Tr}\,\Delta\overline{\Delta\Delta}'_{ab}$
0	$x_{13}^{ab}=\mathrm{Tr}\,\Delta\Omega\overline{\Delta}'_{ab}$
-4	$x_{14}^{ab}=\mathrm{Tr}\,\overline{\Delta}\Omega\overline{\Delta}'_{ab}$
-2	$x_{15}^{abcd}=\mathrm{Tr}\,\overline{\Delta\Delta}'_{ab}\overline{\Delta}'_{cd}$
-6	$x_{16}^{abcdef}=\mathrm{Tr}\,\overline{\Delta}'_{ab}\overline{\Delta}'_{cd}\overline{\Delta}'_{ef}$
-2	$x_{17}^{ab}=\sigma_{ab}$

(A3)

and a similar set in the right handed sector. It is easy to show that Eq. (12) continue to hold in the present case. Thus one can again choose the parity breaking vacuum by selecting the solution where all (unprimed) left handed triplets except $\overline{\Delta}$ and x_1^c,x_2^c,x_6^c vanish while the VEVs of x_3^c,x_5^c are fixed by the solution $x_4=2m_\Delta^2/a^2$. This leaves us with the invariants $x_7^{(c)}-x_{17}^{(c)}$ to consider. In the right handed sector, as we saw in Sec. III, the equation for F_{L_c} ensures that the right handed sneutrino VEVs necessarily vanish. Since $x_{17}^{(c)}=-i\nu^c_{[a}e^c_{b]}$, it vanishes. Because of the Fierz relations, Eq. (A1) (and a similar one relating the $x_{16}^{(c)}$ to products of $x_{17}^{(c)}$ and $x_8^{(c)}$), $x_8^{(c)}$ and $x_{16}^{(c)}$ all vanish while

$$\overline{\Delta}'_{(c)ab}=e^c_a e^c_b\tau_+ . \tag{A4}$$

Then using the F equations in the right handed sector it is easy to convince oneself that all invariants are fixed in terms of $x_9^{(c)ab}$. However, because of Eq. (A4), only three are of these are independent and these can be taken to be $x_9^{(c)aa}$.

In the left handed sector one finds that $\overline{\Delta}$ is determined in terms of $\overline{\Delta}'_{ab}$ by the equation for F_Δ; thus only invariants involving σ_{ab} or $\overline{\Delta}'_{(c)ab}$ are left. But since all of these may be written in terms of products of $x_{17}^{[ab]}$, its is clear that these are the only independent $\mathrm{SU}(2)_L$ invariants left undetermined in the left sector and are 3 in number. Since x_{17}^{ab} carries $B-L$ charge (-2) and $x_9^{(c)aa}$ carries (-4), it is clear that one can form nine independent gauge invariants which are left undetermined after imposition of the F constraints:

$$z_{[ab]d}=x_{17}^{ab}[x_9^{(c)dd}]^{1/2}. \tag{A5}$$

Thus the manifold of flat directions running out of the parity breaking vacuum is parametrized by these nine complex coordinates. From Eq. (A4) and $\sigma_{ab}=-i\nu_{[a}e_{b]}$, it is thus clear that the coordinates $z_{[ab]d}$ all involve a product of selectron and antiselectron VEVs, and hence these flat directions violate charge.

[1] R. Mohapatra, Phys. Rev. D **34**, 3457 (1986).
[2] S. P. Martin, Phys. Rev. D **46**, 2769 (1992).
[3] A. Font, L. Ibáñez, and F. Quevedo, Phys. Lett. B **228**, 79 (1989).
[4] L. Ibáñez and G. G. Ross, Phys. Lett. B **260**, 291 (1991).
[5] J. Pati and A. Salam, Phys. Rev. D **10**, 275 (1974).
[6] R. Mohapatra and J. Pati, Phys. Rev. D **11**, 2558 (1975).
[7] G. Senjanović and R. Mohapatra, Phys. Rev. D **12**, 1502 (1975).
[8] G. Senjanović, Nucl. Phys. **B153**, 334 (1979).
[9] R. Kuchimanchi and R. N. Mohapatra, Phys. Rev. D **48**, 4352 (1993).
[10] R. Kuchimanchi and R. N. Mohapatra, Phys. Rev. Lett. **75**, 3989 (1995).
[11] C. S. Aulakh, K. Benakli, and G. Senjanović, Phys. Rev. Lett. **79**, 2188 (1997).
[12] C. S. Aulakh, A. Melfo, and G. Senjanović, Phys. Rev. D **57**, 4174 (1998).
[13] M. Gell-Mann, P. Ramond, and R. Slansky, in *Supergravity*, edited by P. van Niewenhuizen and D. Freedman (North-Holland, Amsterdam, 1979).
[14] T. Yanagida, in *Proceedings of the Workshop on Unified Theory and Baryon number in the Universe*, Tsukuba, Japan, 1979, edited by O. Sawada and A. Sugamoto (KEK, Tsukuba, 1979).
[15] R. N. Mohapatra and G. Senjanović, Phys. Rev. Lett. **44**, 912 (1980).
[16] M. Cvetič, Phys. Lett. **164B**, 55 (1985).
[17] R. N. Mohapatra and A. Rašin, Phys. Rev. Lett. **76**, 3490 (1996).
[18] C. Aulakh and R. N. Mohapatra, Phys. Rev. D **28**, 217 (1983).
[19] J. Maalampi and J. Pulido, Nucl. Phys. **B228**, 242 (1983).
[20] E. K. Akhmedov, Z. G. Berezhiani, and G. Senjanović, Phys. Rev. Lett. **69**, 3013 (1992).
[21] F. Buccella, J. P. Derendinger, S. Ferrara, and C. A. Savoy, Phys. Lett. **115B**, 375 (1982).
[22] I. Affleck, M. Dine, and N. Seiberg, Nucl. Phys. **B241**, 493 (1984).
[23] I. Affleck, M. Dine, and N. Seiberg, Nucl. Phys. **B256**, 557 (1985).
[24] M. A. Luty and W. Taylor, Phys. Rev. D **53**, 3399 (1996).
[25] G. Dvali, Phys. Lett. B **324**, 59 (1994).
[26] L. Álvarez-Gaume, M. Claudson, and M. B. Wise, Nucl. Phys. **B207**, 96 (1982).
[27] C. S. Aulakh and R. N. Mohapatra, Phys. Lett. **119B**, 136 (1982).
[28] R. N. Mohapatra and G. Senjanović, Phys. Rev. D **23**, 165 (1981).

[29] R. N. Mohapatra and A. Rašin, Phys. Rev. D **54**, 5835 (1996).
[30] R. N. Mohapatra, A. Rašin, and G. Senjanović, Phys. Rev. Lett. **79**, 4744 (1997).
[31] K. Huitu, J. Maalampi, and M. Raidal, Nucl. Phys. **B420**, 449 (1994).
[32] K. Huitu, J. Maalampi, and M. Raidal, Phys. Lett. B **328**, 60 (1994).
[33] K. Huitu, J. Maalampi, A. Pietila, and M. Raidal, Nucl. Phys. **B487**, 27 (1997).
[34] F. Cuypers and S. Davidson, Eur. Phys. J. **C2**, 503 (1998).
[35] J. Gunion, C. Loomis, and K. T. Pitts, hep-ph/9610237.
[36] F. Cuypers and M. Raidal, Nucl. Phys. **B501**, 3 (1997).
[37] Z. Chacko and R. N. Mohapatra, Phys. Rev. D **58**, 015001 (1998).