Supersymmetric left-right models and light doubly charged Higgs bosons and Higgsinos

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(Received 12 December 1997; published 19 May 1998)

We point out that in a large class of supersymmetric left-right models with automatic R-parity conservation there is a pair of light doubly charged Higgs bosons and Higgsinos. Requiring the mass of these particles to satisfy the CERN LEP Z-width bound implies that the W_R mass must be above 10^9 GeV. [S0556-2821(98)02513-2]

PACS number(s): 12.60.Jv, 11.30.Er, 12.60.Cn

I. INTRODUCTION

Supersymmetric left-right models (SUSYLR) are attractive for several reasons: (1) they imply automatic conservation of baryon and lepton number [1], a property which makes the standard model so attractive, but is not shared by the minimal supersymmetric standard model (MSSM); (2) they provide a natural solution to the strong and weak CP problems of the MSSM [2]; (3) they yield a natural embedding of the seesaw mechanism for small neutrino masses [3]; (4) they arise as an intermediate scale theory in several SUSY grand unified theories (GUT's).

An essential feature of these models is that the $SU(2)_R \times U(1)_{B-L}$ symmetry is broken down to $U(1)_Y$ by the vacuum expectation values of a pair of Higgs multiplets which transform as the adjoint of the $SU(2)_R$ group with $B-L=\pm 2$. One of these is the same Higgs multiplet (denoted by Δ^c below) which is used to implement the seesaw mechanism for small neutrino masses. Both of them contain doubly charged Higgs bosons and Higgsinos. It has recently been shown [4] that in simple versions of this theory where the hidden sector SUSY breaking scale is above the M_{W_R}, the ground state of the theory breaks R-parity unless higher dimensional terms [2,5] or additional Higgs fields which break parity [4,6] are included.

It is the goal of this paper to show that the constraints of supersymmetry imply that the above mentioned pair of doubly charged particles are very light in a large class of interesting versions of the SUSYLR model. Since these masses depend on the scale $v_R \equiv v$ of $SU(2)_R$ breaking, one can use the CERN e^+e^- collider LEP Z-width constraints to fix a lower bound on v_R. Such a bound has already been noted in the minimal SUSYLR model [5] in the limit of exact supersymmetry.

The existence of this pair of light doubly charged fields is independent of the scale at which supersymmetry is broken provided that the effect of the breaking on the Higgs sector is soft. However the mass splitting between the Higgs bosons and Higgsinos is crucially dependent on whether the scale of supersymmetry breaking is higher than the W_R scale (as in a supergravity mediated scenario) or lower than the W_R scale (as in a gauge mediated scenario). This causes the bounds on the W_R scale to arise differently in the two cases. In the former, the bound arises from considering the Higgs boson masses and in the latter case from the Higgsino masses.

It also turns out that parity invariance does not play any role in our proof. Therefore the bound on v_R applies to models based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ without parity as long as the $SU(2)_R$ symmetry is broken by the same type of fields.

In the model of Ref. [5], it was demonstrated that it is possible to break both the gauge symmetry and parity using nonrenormalizable operators alone. This model, the minimal SUSYLR model predicts the existence of a total of nine light complex fields of which three are neutral, two singly charged and four doubly charged. These light fields have masses of order $\sim M_R^2/M_{Pl}$, which was used to set a lower bound on the W_R scale of 10^{10} GeV in the supersymmetric limit. However, since SUSY breaking terms are of order (100 GeV) [2], it is important that the soft SUSY breaking terms be considered in the full analysis. Since this model falls into the class we investigate, two of these four doubly charged fields correspond exactly to the ones predicted by us.

The underlying reason for the existence of this pair of light doubly charged particles can be understood by considering the result of Ref. [4]. The particle content of this model consists of bidoublets ϕ (2,2,0), Higgs fields Δ (3,1,2), $\bar{\Delta}$(3,1,−2), Δ^c(1,3,−2) and $\bar{\Delta}^c$(1,3,+2) and a singlet in addition to the usual quarks and leptons. [The numbers in parenthesis refer to their transformation properties under $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$.] If only renormalizable interactions of these fields are considered and R-parity is unbroken (i.e. $\langle \tilde{\nu}^c \rangle = 0$) then in the absence of supersymmetry breaking there exist a continuously connected set of vacua parametrized by a single angle θ such that

$$\Delta^c = \begin{bmatrix} 0 & v\cos\theta \\ v\sin\theta & 0 \end{bmatrix} \quad (1)$$

$$\bar{\Delta}^c = \begin{bmatrix} 0 & v\sin\theta \\ v\cos\theta & 0 \end{bmatrix}. \quad (2)$$

If $\theta=0$, electric charge is conserved; otherwise it is broken. Thus the only phenomenologically viable vacuum, the charge conserving one, is degenerate with a continuously connected set of other vacua. The excitation that corresponds to the flat direction connecting all these vacua must be a massless particle; it is straightforward to verify that in the charge conserving vacuum this particle is doubly charged. The other doubly charged particle is also massless but does not correspond to a flat direction. Once supersymmetry is broken the flat direction is lifted and the theory will choose to live either in the ''good'' vaccum or the ''bad'' vacuum. The main result of Ref. [4] is that if the scale at which supersymmetry is broken is higher than the W_R scale and electroweak symmetry breaking is ignored, then the theory necessarily lives in the charge violating vacuum if R-parity is unbroken. However if non-renormalizable operators suppressed by powers of M_{Planck} are added to the theory [2], then for sufficiently high right handed scale the theory can live in the charge preserving vacuum. However, this suggests that the previously massless Higgs boson corresponding to a flat direction is still light, since the flat direction has only been given a positive slope by these higher dimensional operators. We calculate the mass of this Higgs boson and verify that it is indeed light. We use the experimental lower bound on the mass of such a doubly charged particle to put a lower bound on the W_R scale when electroweak symmetry breaking is ignored. We then show that the inclusion of electroweak symmetry breaking does not alter this result.

If however the W_R scale is above the supersymmetry breaking scale, as in gauge mediated supersymmetry breaking scenarios [7], we show that even the renormalizable theory may live in the charge preserving vacuum. In this case the light Higgs bosons pick up a mass from the breaking of supersymmetry which is of the same order of magnitude as the masses of the superpartners of the standard model particles. Now however the corresponding Higgsinos are very light since the breaking of supersymmetry is assumed to be soft. Thus the non-renormalizable operators are once again needed, this time to give mass to the Higgsinos. The experimental lower bound on the mass of such particles can once again be used to put a lower bound on the right handed scale.

Finally, we also point out that a light doubly charged Higgs and Higgsino fields are also present in the version of the model where the vacuum state breaks R-parity since it was shown in Ref. [4] that in these models there is an upper limit on the W_R scale of order of a TeV and we expect the masses of all particles in the theory to be at most of the order of the W_R mass.

Although our detailed analysis is limited to a specific class of models, we consider whether our result holds for models with further matter content. We find that SUSYLR theories necessarily imply these light doubly charged Higgs superfields unless the model contains exotic light doubly charged $SU(2)_R$ singlets or certain additional Higgs multiplets which break $SU(2)_R$ while preserving hypercharge. Our results also have important implications for coupling constant unification in SUSY GUTs.

II. ANALYSIS FOR W_R SCALE BELOW SUSY BREAKING SCALE

In this section we calculate the masses of the doubly charged Higgs bosons and Higgsinos for a specific class of models when supersymmetry is broken above the W_R scale as in a supergravity mediated scenario. We first study the theory classically and subsequently justify our conclusions where required using an effective field theory analysis.

The matter content of the model we examine consists of the quarks $Q(2,1,1/3)$ and $Q^c(1,2,-1/3)$, the leptons $L(2,1,-1)$ and $L^c(1,2,1)$, the electroweak Higgs bidoublet $\phi(2,2,0)$, the triplets $\Delta(3,1,2), \Delta^c(1,3,-2), \overline{\Delta}(3,1,-2)$ and $\overline{\Delta}^c(1,3,2)$ and an arbitrary number of singlets $S_i(0,0,0)$ where the numbers in parentheses refer to their transformation properties under $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ respectively. In a later section we will consider models with additional matter content. Our procedure will be to first write down the most general potential involving the above fields consistent with the symmetries and from there to obtain the mass matrices of the doubly charged fields and show that one of the eigenvalues of the Higgs boson mass matrix is light.

We consider the most general superpotential consisting of the above fields. In order to account for the possibility that the right handed scale is large, we include, in addition to the renormalizable interactions, all possible nonrenormalizable interactions of the Δ's and Δ^c's among themselves to lowest order in $1/M_{Planck}$. Then, the relevant part of the superpotential is

$$\begin{aligned} W = & ifL^{c^T}\tau_2\Delta^c L^c + (M_0 + \lambda S_1)\mathrm{Tr}(\Delta^c\overline{\Delta}^c) \\ & + G(S_i, X_i) + A[\mathrm{Tr}(\Delta^c\overline{\Delta}^c)]^2/2 \\ & + B\ \mathrm{Tr}(\Delta^c\Delta^c)\mathrm{Tr}(\overline{\Delta}^c\overline{\Delta}^c)/2. \end{aligned} \tag{3}$$

In the above equation, X_i is a generic label for any field apart from the Δ^c's that the S_i's couple to; A, B, f, λ and M are parameters of the theory and $G(S_i, X_i)$ is the most general superpotential in the S_i's and X_i's alone. A and B are of order $1/M_{Planck}$.

Using Eq. (3), one can give a group theoretical argument for the existence of light doubly charged particles in the supersymmetric limit as follows. For this purpose let us first ignore the higher dimensional terms A and B as well as the leptonic couplings f. It is then clear that the superpotential has a complexified U(3) symmetry [i.e. a U(3) symmetry whose parameters are taken to be complex] that operates on the Δ^c and $\overline{\Delta}^c$ fields. This is due to the holomorphy of the superpotential. After one component of each of the above fields acquires a vacuum expectation value (VEV) as in the charge conserving case with $\theta = 0$ (and supersymmetry guarantees that both VEV's are parallel), the resulting symmetry is the complexified U(2). This leaves 10 massless fields. Once we bring in the D-terms and switch on the gauge fields, six of these fields become massive as a consequence of the Higgs mechanism of supersymmetric theories. That leaves four massless fields in the absence of higher dimensional terms. These are the two complex doubly charged fields. Of the two non-renormalizable terms A and B, only the A-term has the complexified U(3) symmetry. Hence the supersymmetric contribution to the doubly charged particles will come only from the B-term. Although the leptonic couplings do not respect this symmetry, they are unimportant in determining the vacuum structure as long as R-parity is conserved and hence they do not alter our result. Let us now proceed to prove this via explicit calculations.

The most general soft supersymmetry breaking potential compatible with the symmetries and relevant for our analysis has the form

$$\begin{aligned}V_S=&M_L^2L^{c\dagger}L^c+M_1^2\ \mathrm{Tr}(\Delta^{c\dagger}\Delta^c)+M_2^2\ \mathrm{Tr}(\overline{\Delta^c}^{\dagger}\overline{\Delta^c})\\&+\Sigma_i(\lambda_i' M_S\ \mathrm{Tr}(\Delta^c\overline{\Delta^c})S_i)\\&+M'^2\ \mathrm{Tr}(\Delta^c\overline{\Delta^c})+if'M''L^{cT}\tau_2\Delta^cL^c\\&+G'(S_i,X_i,S_i^\dagger,X_i^\dagger)+\text{H.c.}\end{aligned}\tag{4}$$

where M_L, M_1, M_2, M', λ_i', M_S, f' and M'' are parameters and $G'[S_i,X_i,S_i^\dagger,X_i^\dagger]$ is an arbitrary function in the S_i's, X_i's and their hermitian conjugates consistent with the soft breaking of SUSY. The relevant part of the D terms have the form

$$\begin{aligned}V_D=&\frac{g_2^2}{8}\Sigma_i(2\ \mathrm{Tr}\ \Delta^{c\dagger}\tau_i\Delta^c+2\ \mathrm{Tr}\ \overline{\Delta}^{c\dagger}\tau_i\overline{\Delta}^c\\&+\mathrm{Tr}\ \phi^\dagger\tau_i\phi+L^{c\dagger}\tau_iL^c+\overline{L}^{c\dagger}\tau_i\overline{L}^c)^2\\&+\frac{g_1^2}{2}\bigg(\mathrm{Tr}\ \Delta^{c\dagger}\Delta^c-\mathrm{Tr}\ \overline{\Delta}^{c\dagger}\overline{\Delta}^c\\&-\frac{1}{4}L^{c\dagger}L^c+\frac{1}{4}\overline{L}^{c\dagger}\overline{L}^c\bigg)^2.\end{aligned}\tag{5}$$

Since the masses of the doubly charged fields will be dependent on the neutral Higgs vacuum expectation values (VEVs), we first set out to determine these. The potential to be minimized in order to determine the Higgs VEVs consists of a sum of F,D and soft supersymmetry breaking terms. We assume that R-parity is unbroken so that $\langle\tilde{L}\rangle=\langle\tilde{L}^c\rangle=0$ and the terms involving these fields are unimportant in determining the classical vacuum. Then the potential to determine the Higgs VEVs can be written as

$$\begin{aligned}V=&[K\overline{\Delta}_j^{ci}+A\ \mathrm{Tr}(\Delta^c\overline{\Delta}^c)\overline{\Delta}_j^{ci}+B\ \mathrm{Tr}(\overline{\Delta}^c\overline{\Delta}^c)\Delta_j^{ci}][\text{H.c.}]\\&+[K\Delta_j^{ci}+A\ \mathrm{Tr}(\Delta^c\overline{\Delta}^c)\Delta_j^{ci}+B\ \mathrm{Tr}(\overline{\Delta}^c\overline{\Delta}^c)\overline{\Delta}_j^{ci}][\text{H.c.}]\\&+\beta[\mathrm{Tr}(\Delta^c\overline{\Delta}^c)][\text{H.c.}]+[\alpha_0\ \mathrm{Tr}(\Delta^c\overline{\Delta}^c)+\text{H.c.}]\\&+M_1^2\ \mathrm{Tr}(\Delta^{c\dagger}\Delta^c)+M_2^2\ \mathrm{Tr}(\overline{\Delta}^{c\dagger}\overline{\Delta}^c)+V_D.\end{aligned}\tag{6}$$

Here β, α_0, and K are parameters; the latter two in general depend upon the expectation value of the singlet and its F-component. While $K=\lambda\langle S_1\rangle+M$, α_0 picks up contributions from the F component of S_1 in the form $\langle\partial G/\partial S_1\rangle^\dagger[\mathrm{Tr}(\Delta^c\overline{\Delta}^c)]$ and also from M'^2 and λ_iM_S in the soft supersymmetry breaking terms. β is simply λ^2.

We now assume the pattern of symmetry breaking to be such as to break $SU(2)_R$ while preserving electric charge and R-parity. As has been shown in Refs. [4, 5], for the class of models we are interested in, this is always true. In an arbitrary SUSYLR model of this class it is certainly not guaranteed that the pattern of left-right symmetry breaking will work out correctly. However what we intend to show is that for any model of this type which does have the right pattern of left-right symmetry breaking and lives in the charge preserving vacuum the doubly charged fields will be light. If our assumption that the theory lives in the "good" vacuum is true then the masses of the doubly charged Higgs bosons will turn out to be positive at the end of our calculation; if not at least one of these fields will turn out to have negative mass thereby giving us a consistency check on our assumption. In order to give our proof, let us expand the Higgs field in components as

$$\Delta^c=\begin{bmatrix}\Delta^{c-}/\sqrt{2} & v+\eta\\ \Delta^{c--} & -\Delta^{c-}/\sqrt{2}\end{bmatrix}\tag{7}$$

$$\overline{\Delta^c}=\begin{bmatrix}\overline{\Delta}^{c+}/\sqrt{2} & \overline{\Delta}^{c++}\\ \overline{v}+\overline{\eta} & -\overline{\Delta}^{c+}/\sqrt{2}\end{bmatrix}\tag{8}$$

while the electroweak Higgs bidoublet ϕ has the VEVs

$$\phi=\begin{bmatrix}\kappa_u & 0\\ 0 & \kappa_d\end{bmatrix}.\tag{9}$$

While M_1^2, M_2^2 and β are necessarily real, K, α_0, A and B are in general complex. By redefining the Δ^c's by phase factors we can make v and $\overline{v}$ real. Defining $\alpha=\mathrm{Re}[\alpha_0]$ and $\alpha'=\mathrm{Im}[\alpha_0]$ this implies that the imaginary part of α_0 satisfies the equation

$$\alpha' v\overline{v}(v^2+\overline{v}^2)+\mathrm{Im}[Av\overline{v}(K+Av\overline{v})^*]=0.\tag{10}$$

The potential to be minimized in order to determine $v,\overline{v}$ is

$$\begin{aligned}V=&\beta v^2\overline{v}^2+2\alpha v\overline{v}+(K+Av\overline{v})(K+Av\overline{v})^*(\overline{v}^2+v^2)\\&+M_1^2v^2+M_2^2\overline{v}^2+\frac{g_1^2}{2}[v^2-\overline{v}^2]^2\\&+\frac{g_2^2}{8}[2(v^2-\overline{v}^2)+\kappa_U{}^2-\kappa_D{}^2]^2.\end{aligned}\tag{11}$$

The equations which determine v and $\overline{v}$ then have the form

$$\begin{aligned}&\overline{v}(\alpha+\beta v\overline{v})+\mathrm{Re}[A^*(K+Av\overline{v})]\overline{v}(v^2+\overline{v}^2)\\&+v(K+Av\overline{v})(K+Av\overline{v})^*+v[M_1^2+g_1^2(v^2-\overline{v}^2)]\\&+v\frac{g_2^2}{2}[2(v^2-\overline{v}^2)+\kappa_u{}^2-\kappa_d{}^2]=0\end{aligned}\tag{12}$$

$$v(\alpha+\beta v\overline{v})+\mathrm{Re}[A^*(K+Av\overline{v})]v(v^2+\overline{v}^2)$$
$$+\overline{v}(K+Av\overline{v})(K+Av\overline{v})^*+\overline{v}[M_2^2-g_1^2(v^2-\overline{v}^2)]$$
$$-\overline{v}\frac{g_2^2}{2}[2(v^2-\overline{v}^2)+\kappa_u{}^2-\kappa_d{}^2]=0. \tag{13}$$

Defining

$$M^2=(M_1^2+M_2^2)/2 \tag{14}$$

$$\delta=(M_1^2-M_2^2)/2 \tag{15}$$

$$\chi=\delta+g_1^2[v^2-\overline{v}^2]+g_2^2[2(v^2-\overline{v}^2)+(\kappa_u)^2-(\kappa_d)^2]/2 \tag{16}$$

$$\Lambda=\delta+g_1^2[v^2-\overline{v}^2]-g_2^2[2(v^2-\overline{v}^2)+(\kappa_u)^2-(\kappa_d)^2]/2. \tag{17}$$

Let us keep in mind that the parameters M^2, δ, χ and Λ are all of order M_{SUSY}^2, the mass scale for the soft terms; A and B are of order $1/M_{Pl}$ and the α-term depends on the VEV's of the singlet fields and could therefore be arbitrary.

Let us now rewrite the extremization equations as

$$(\alpha+\beta v\overline{v})+\mathrm{Re}[A^*(v^2+\overline{v}^2)(K+Av\overline{v})]+(K+Av\overline{v})$$
$$\times(K+Av\overline{v})^*+M^2$$
$$=-\chi(v-\overline{v})/(v+\overline{v}) \tag{18}$$

$$-(\alpha+\beta v\overline{v})-\mathrm{Re}[A^*(K+Av\overline{v})](v^2+\overline{v}^2)$$
$$+(K+Av\overline{v})(K+Av\overline{v})^*+M^2$$
$$=-\chi(v+\overline{v})/(v-\overline{v}). \tag{19}$$

Let us note that $K+Av\overline{v}$ and $\alpha+\beta v\overline{v}$ vanish in the SUSY limit. If we assume that the only source of supersymmetry breaking is from the soft breaking terms and if none of the singlets has VEVs far exceeding the right handed scale M_R then careful analysis shows that these are generically at most of order M_{SUSY} and M_{SUSY}^2 respectively. This provides a qualitative way to see why the masses of the doubly charged fields are small compared to the $SU(2)_R$ scale since it is these combinations that appear in the mass matrix for the doubly charged particles.

To prove our result in more detail, let us multiply the above two equations to get a result which will be useful in the subsequent discussion:

$$[(K+Av\overline{v})(K+Av\overline{v})^*+M^2]^2-[(\alpha+\beta v\overline{v})+\mathrm{Re}[A^*(K+Av\overline{v})](v^2+\overline{v}^2)]^2=\chi^2. \tag{20}$$

We now calculate the mass matrix for the doubly charged Higgs bosons and obtain it to be

$$\begin{array}{c} \\ \Delta^{c++} \\ \overline{\Delta}^{c++} \end{array} \begin{array}{c} \begin{array}{cc} \Delta^{c--} & \qquad\qquad\qquad \overline{\Delta}^{c--} \end{array} \\ \begin{pmatrix} [K+(A+2B)v\overline{v}][\mathrm{H.c.}]+M^2+\Lambda & (\alpha_0^*+\beta v\overline{v})+(v^2+\overline{v}^2)(A+2B)^*(K+Av\overline{v}) \\ (\alpha_0+\beta v\overline{v})+(v^2+\overline{v}^2)(A+2B)(K+Av\overline{v})^* & [K+(A+2B)v\overline{v}[\mathrm{H.c.}]+M^2-\Lambda \end{pmatrix} \end{array}. \tag{21}$$

If either of the eigenvalues of this matrix is negative then the square of one of the scalar masses is negative and our assumption that the theory preserves electric charge is invalid. Rather than calculate the eigenvalues directly we choose to infer information by examining the trace, T and determinant, D of the above matrix. If either of these turns out to be negative the theory breaks electric charge. We first determine the trace which is the sum of the eigenvalues as

$$T=2[K+(A+2B)v\overline{v}][\mathrm{H.c.}]+2M^2. \tag{22}$$

This is typically of order $O(M_{SUSY}^2)$ or $O(M_R^4/M_{Planck}^2)$ where M_{SUSY} is the scale of the soft SUSY breaking mass terms and M_R the right handed scale. Since the product of the eigenvalues is merely the determinant we proceed to evaluate this:

$$D=[(K+(A+2B)v\overline{v})(\mathrm{H.c.})+M^2]^2-[(\alpha+\beta v\overline{v})+(v^2+\overline{v}^2)\mathrm{Re}[(A+2B)(K+Av\overline{v})^*]]^2$$
$$-[\alpha'+\mathrm{Im}[(A+2B)(K+Av\overline{v})^*](v^2+\overline{v}^2)]^2-\Lambda^2. \tag{23}$$

Using Eqs. (10) and (20), this simplifies to

$$D=[\chi^2-\Lambda^2]+(4\,\mathrm{Re}[Bv\overline{v}(K+Av\overline{v})^*]+4BB^*v^2\overline{v}^2)^2-(v^2+\overline{v}^2)^2[(2B)(K+Av\overline{v})][\mathrm{H.c.}]+8(\mathrm{Re}[(Bv\overline{v})(K+Av\overline{v})^*]$$
$$+BB^*v^2\overline{v}^2)[(K+Av\overline{v})(\mathrm{H.c.})+M^2]-4[\alpha+\beta v\overline{v}+\mathrm{Re}[A^*(K+Av\overline{v})](v^2+\overline{v}^2)][\mathrm{Re}[B^*(K+Av\overline{v})](v^2+\overline{v}^2)]. \tag{24}$$

We now have enough information to estimate the masses of the Higgs bosons. For simplicity we subdivide our analysis into two cases:

(1) low W_R scale, so that all terms suppressed by powers of M_{Planck} can be neglected;

(2) high W_R scale.

We first consider the low W_R scale case. This then corresponds to the renormalizable theory, i.e. $A=B=0$. The only term present is $\chi^2-\Lambda^2$ which we now examine in more detail. Using

$$\Lambda=\chi-g_2^2[2(v^2-\overline{v^2})+\kappa_u{}^2-\kappa_d{}^2] \tag{25}$$

and

$$\chi=-(v^2-\overline{v^2})(K^2+M^2)/(v^2+\overline{v^2}) \tag{26}$$

we find

$$\begin{aligned}\chi^2-\Lambda^2=&-g_2^4[2(v^2-\overline{v^2})+\kappa_u{}^2-\kappa_d{}^2]^2-4g_2^2(v^2-\overline{v^2})^2\\&\times[K^2+M^2]/(v^2+\overline{v^2})-2g_2^2(v^2-\overline{v^2})[K^2+M^2]\\&\times(\kappa_u{}^2-\kappa_d{}^2)/(v^2+\overline{v^2}).\end{aligned} \tag{27}$$

In the limit that electroweak effects are ignored (i.e. $\kappa_u=\kappa_d=0$) this is less than zero reproducing the known result [4] that the renormalizable theory has no charge conserving vacuum. We now see however that the last term in Eq. (27) may in fact alter this result if $|-\kappa_u{}^2+\kappa_d{}^2|>2|v^2-\overline{v^2}|$ and the W_R scale is low. We can estimate the mass of the lightest doubly charged boson to be $M_{++}\leqslant\frac{1}{4}(|\kappa_u^2-\kappa_d^2|)/\sqrt{v^2+\overline{v^2}}$. This implies that the scale of right-handed symmetry breaking (i.e. $\sqrt{v^2+\overline{v^2}}$) is less than about 400 GeV. For the case of manifest left-right symmetry, such a low value for M_{W_R} is inconsistent with observations. Thus for the R-parity conserving vacuum, low M_{W_R} scenario is inconsistent.

We now consider the high W_R case. Now, however, the theory with the non-renormalizable operators can lie in the charge preserving vacuum for sufficiently high W_R scale. We estimate the mass of the lighter of the doubly charged particles is either $O(M_R\sqrt{M_{SUSY}/M_{Planck}})$ or $O(M_R^2/M_{Planck})$, whichever is larger, while the mass of the heavier is the larger of $O(M_R^2/M_{Planck})$ and $O(M_{SUSY})$. The experimental lower bound on the mass of the lighter particle implies that $M_R>10^9$ GeV.

In both the low and high W_R cases as a consequence of supersymmetry breaking the Higgsinos pick up a mass from the following term in the superpotential:

$$W=[K+Av\overline{v}+2Bv\overline{v}]\Delta^{c--}\overline{\Delta}^{c++}. \tag{28}$$

This is a Dirac mass of order $O(M_{SUSY})$ or M_R^2/M_{Pl}, whichever is larger.

However in view of the fact that the high W_R scenario necessarily envisages a large hierarchy between the right handed scale and the scale of the soft SUSY masses, an effective field theory calculation would be more convincing than our tree level result. For simplicity we now restrict ourselves to the case of a single singlet S, which we integrate out at tree level along with the other heavy fields Δ^{c0}, $\overline{\Delta}^{c0}$, Δ^{c-}, and $\overline{\Delta}^{c+}$. The remaining effective field theory, consisting of Δ^{c--}, $\overline{\Delta}^{c++}$ and some of the X_i's, is then run down to the scale of the light Higgs fields. The potential we start from is the same as before except that $G(S_i,X_i)$ and $G'[S_i,X_i,S_i^\dagger,X_i^\dagger]$ become $G(S,X_i)$ and $G'[S,X_i,S_i^\dagger,X_i^\dagger]$ respectively while λ_i and λ' are now simply λ and λ'.

We expand Δ^{c0}, $\overline{\Delta}^{c0}$, and S about their vacuum expectation values as

$$\Delta^{c0}=v+\eta \tag{29}$$

$$\hat{\Delta}^{c0}=\overline{v}+\overline{\eta} \tag{30}$$

$$S=\langle S_0\rangle+S'. \tag{31}$$

Defining $\sigma_1=v\,\overline{\eta}+\overline{v}\,\eta$ and $\sigma_2=v$ Re $\eta-\overline{v}$ Re $\overline{\eta}$, we find that to the extent that supersymmetry breaking terms and terms suppressed by powers of M_{Planck} are small, σ_2 and two linearly independent combinations of σ_1 and S are approximate mass eigenstates. We define these two linearly independent combinations as S_1 and S_2. Then it is straightforward to verify that as a consequence of the cancellations of the tree level graphs involving the exchange of σ_1, S_1 and S_2, to order $O(M_{SUSY}/M_R)$ or $O(M_R/M_{Planck})$ the only residual interactions among the light fields are those in the effective potential below. We can write the part of the effective potential relevant for the light doubly charged Higgs field as

$$V=\left(\left[\int d^2\theta W+\text{H.c.}\right]+V_S+V_D\right) \tag{32}$$

where

$$W=[K+Av\overline{v}+2Bv\overline{v}]\Delta^{c--}\overline{\Delta}^{c++}+fe^ce^c\Delta^{c--} \tag{33}$$

$$\begin{aligned}V_{soft}=&[(\alpha_0+\beta v\overline{v}+(v^2+\overline{v^2})(A+2B)(K+Av\overline{v})^*)\\&\times\Delta^{c--}\overline{\Delta}^{c++}+\text{H.c.}]+(M^2+\Lambda)\Delta^{c++}\Delta^{c--}+(M^2\\&-\Lambda)\overline{\Delta}^{c++}\overline{\Delta}^{c--}+f'M''\tilde{e}^c\tilde{e}^c\Delta^{c--}+\text{H.c.}\end{aligned} \tag{34}$$

$$V_D=\frac{2g_1^2g_2^2}{(g_1^2+g_2^2)}[\Delta^{c++}\Delta^{c--}-\overline{\Delta}^{c--}\overline{\Delta}^{c++}]^2. \tag{35}$$

This effective theory must now be run down to the mass scale of the doubly charged fields M_Δ using the renormalization group equations (RGE's). On writing down the relevant RGE's it is clear that the only potentially large contribution to the evolution of the boson mass terms is likely to arise from the coupling to the leptons f and f' and is of order $O(\sqrt{f^2M_{SUSY}^2\ln[M_R/M_\Delta]/8\pi^2})$. Hence our tree level result for the mass terms may be corrected by about this amount.

Clearly this will not qualitatively alter our result that the doubly charged Higgs bosons will be light. The fermion masses suffer only wave function renormalization and also remain light.

III. ANALYSIS FOR W_R SCALE ABOVE SUSY BREAKING SCALE

In the previous section we assumed that the scale at which supersymmetry is broken is higher than the W_R scale. However this need not be the case and in particular there has recently been a lot of interest in theories where gauge interactions are the mediators of supersymmetry breaking at a relatively low scale [7]. This is the case we now study in detail.

Our analysis essentially will differ from that of the previous section only in that the soft SUSY breaking terms are now generated explicitly only at the scale at which the messenger fields are integrated out, and are not explicitly present at the W_R scale. Since they are generated by loop graphs involving the gauge bosons of the residual symmetries, their form will be such as to respect only the surviving gauge symmetries. We will show that this difference has consequences for phenomenology. Our procedure must therefore be to integrate out the heavy fields at the W_R scale, run the theory down to the messenger scale, integrate out the messengers thereby generating the soft SUSY breaking mass terms, and then run the theory down to the M_Δ scale. We make no assumption about the messenger fields except that they carry electroweak quantum numbers, and do not couple directly to the Higgs sector. However for simplicity we restrict ourselves once again to the one singlet case.

After integrating out the heavy fields at the right handed scale, the effective field theory has the form

$$V=\left[\int d^2\theta W+\text{H.c.}\right]+V_D \tag{36}$$

where V_D is the same as in Eq. (5) but W is now simply

$$W=2Bv\bar{v}\Delta^{c--}\bar{\Delta}^{c++}+fe^ce^c\Delta^{c--}. \tag{37}$$

This potential suffers only wave function renormalization down to the messenger scale. Then on integrating out the messenger fields soft SUSY breaking terms will be generated, the form of which are to some extent dependent on the nature of the messengers. However, if the messengers couple to the Higgs sector only through gauge interactions and not directly through the superpotential then the relevant part of these terms generically have the form,

$$V_S=M^2(\Delta^{c++}\Delta^{c--}+\bar{\Delta}^{c--}\bar{\Delta}^{c++}). \tag{38}$$

These terms arise from two loop diagrams involving the messenger fields coupling via hypercharge gauge interactions to the Higgs sector. The total potential must then be run down to the electroweak scale using the same RGE's as in the previous section, and the mass term will once again receive some modification. However, this cannot alter the basic result that the doubly charged Higgs bosons only acquire a mass of order O(M_{SUSY}) or O(M_R^2/M_{Planck}) and therefore remain light. Our analysis however brings up the following interesting question; in such a scenario can the theory live in the charge preserving vacuum without the need for higher dimensional operators? After all, since the light fields all have positive mass even in the absence of the higher dimensional operators the charge conserving vacuum is at least a local vacuum of the theory even without them. Notice however that the doubly charged Higgsinos then have no mass. Thus the nonrenormalizable terms are still required, this time to give mass to the Higgsinos. The experimental lower bound on the mass of such a doubly charged fermion puts a bound on the W_R scale of $M_{W_R}\geq 10^{10}$ GeV.

It is interesting that the theory can in fact live in the charge preserving vacuum without the higher dimensional operators. The possibility exists that by bringing the W_R and SUSY breaking scales close together or by reconsidering the assumption that the messenger and Higgs sectors do not directly couple in the superpotential it may be possible to do away with the higher dimensional operators altogether, thereby altering the bound. This is a possible direction for future research.

IV. MODELS WITH ADDITIONAL MATTER CONTENT

In this section we consider the effects of relaxing our earlier restrictions on the matter content of the model we studied on our result that the doubly charged Higgs bosons and Higgsinos are light. In order for the additional matter content to affect our result there must be a difference between the vacuum energies of the charge conserving and charge violating vacua or at least a barrier between them. This must occur in the limit of exact supersymmetry because any correction to the masses from SUSY breaking effects will be at most of order M_{SUSY} and too small to fundamentally alter our result. We also ignore nonrenormalizable terms because corrections to the masses arising from these can reasonably be expected to be small. Hence we will be studying renormalizable theories in the limit of exact supersymmetry and observing the effect of the additional matter on the masses of the doubly charged Higgs fields. We also assume that R-parity is unbroken.

We now attempt to systematically go over some possibilities for the additional matter content. Since the new fields transform as representations of $SU(2)_R$ we proceed in order of increasing dimensionality of the representation.

(1) Charged singlets $T(1,4)$ and $\bar{T}(1,-4)$

Here the numbers within the parenthesis denote the $SU(2)_R$ and $U(1)_{B-L}$ quantum numbers respectively. These allow couplings of the form $T\ \mathrm{Tr}(\Delta^c\Delta^c)$ and $\bar{T}\ \mathrm{Tr}(\bar{\Delta}^c\bar{\Delta}^c)$ in the superpotential. It is straightforward to verify that these do indeed lift the flat direction giving the components of the doubly charged Higgs superfield a mass at the W_R scale.

(2) Charged vector triplets $T(3,2)$ and $\bar{T}(3,-2)$

These have the same quantum numbers as Δ^c and $\bar{\Delta}^c$ Hence superpotential couplings of the form $M\ \mathrm{Tr}(\Delta^cT)$, $M\ \mathrm{Tr}(\bar{T}\bar{\Delta}^c)$, $S_i\ \mathrm{Tr}(\Delta^cT)$ and $S_i\ \mathrm{Tr}(\bar{T}\bar{\Delta}^c)$ are now possible. However these will not lift the flat direction because the

VEVs of T and $\overline{T}$ can always have the same form as those of $\overline{\Delta}^c$ and Δ^c without any effect of the angle θ on the vacuum energy, i.e. if

$$\Delta^c = v\begin{bmatrix} 0 & \cos\theta \\ \sin\theta & 0 \end{bmatrix} \tag{39}$$

then

$$\overline{T} = t\begin{bmatrix} 0 & \cos\theta \\ \sin\theta & 0 \end{bmatrix}. \tag{40}$$

Thus we do not expect our result to change.

Apart from the charged singlet above, no field transforming as a higher dimensional representation of $SU(2)_R$ when added to the theory can preserve the charge conserving vacuum unless it itself also breaks $SU(2)_R$. This may be verified by explicit calculation for all the three cases below which exhaust the possibilities. This is because supersymmetry is explicitly broken by the superpotential in the charge preserving vacuum unless the additional matter fields pick up VEVs thereby breaking $SU(2)_R$ themselves.

(3) A neutral triplet [6] $\Omega^c(3,0)$ that couples to the Higgs sector as $\mathrm{Tr}(\Omega^c\overline{\Delta}^c\Delta^c)$

(4) A neutral quintuplet $T_{ij}^{kl}(5,0)$ that couples as $\Delta_k^{ci}\overline{\Delta}_l^{cj}T_{ij}^{kl}$.

(5) Charged quintuplets $T_{ij}^{kl}(5,4)$ and $\overline{T}_{ij}^{kl}(5,-4)$ that couple to the Higgs sector as $\Delta_k^{ci}\Delta_l^{cj}T_{ij}^{kl}$ and $\overline{\Delta}_k^{ci}\overline{\Delta}_l^{cj}\overline{T}_{ij}^{kl}$.

If the multiplets that have been added to the theory break $SU(2)_R$ then only a detailed analysis of the vacuum structure for each individual theory can determine whether the relevant flat direction is lifted or not. There are also far more possibilities than the three above. A careful analysis has been performed for the neutral triplet [6], which shows that the flat direction is successfully lifted, but not for the other cases. To do so for the other cases is beyond the scope of the present paper.

Thus the conclusion of this section is that the light doubly charged Higgs superfield can be avoided if a doubly charged $SU(2)_R$ singlet is present in the theory or if there are certain specific Higgs multiplets which break $SU(2)_R$ while preserving hypercharge.

V. CONCLUSION

In summary, we find that the combination of supersymmetry with left-right symmetry leads to nontrivial constraints on the mass spectrum of the theory if the starting theory is assumed to be automatically R-parity conserving. In particular, we find the interesting result that for the R-parity conserving scenario, the mass of the doubly charged Higgs bosons and/or Higgsinos which are part of the $SU(2)_R$ multiplets used to implement the seesaw mechanism will be unacceptably light unless the W_R mass is larger than 10^9 GeV. Thus in a large class of simple models where the low scale of parity restoration and supersymmetry are consistent with each other is when R-parity is dynamically broken by the vacuum.

Our result has the following interesting implications:

(i) This should give new impetus to the experimental searches for the W_R boson, since it implies that if experiments exclude a low mass W_R, then its mass can only be in the 10^9 GeV in a large class of simple models. This latter range is of course of great deal of interest in connection with solutions to the solar and atmospheric neutrino puzzles. On the other hand, if a low mass W_R is discovered, it would imply that in the context of simple models that R-parity must be dynamically broken.

(ii) The lightness of the doubly charged fields is now valid even if the $SU(2)_R$ scale is in the superheavy range (i.e. $>10^9-10^{10}$ GeV); discovering the phenomenological effects of these light particles [8] acquires a new urgency and importance. The experimental discovery of such particles would provide spectacular evidence for the supersymmetric left-right model and their masses could provide valuable information about the W_R scale.

(iii) Our result will have important implications for gauge and Yukawa coupling unification in left-right and SO(10) models with automatic R-parity conservation. In particular, the evolution equations will have to include the effects of the doubly charged particles at a much earlier scale than the $SU(2)_R$ breaking scale. Otherwise the exotic multiplets that help us to avoid our result will have to be included above the W_R scale. This is presently under investigation.

(iv) Finally, our results about the lightness of the doubly charged Higgs bosons hold even when R-parity is spontaneously broken, since as already emphasized, in this case the $SU(2)_R$ scale is bound to be in the TeV range [4].

ACKNOWLEDGMENTS

One of us (Z.C.) would like to thank Markus Luty for discussions. This work has been supported by the National Science Foundation grant number PHY-9421386.

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