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Constraining the vacuum expectation values in naturally R-parity conserving supersymmetric models

K. Huitu [a], P.N. Pandita [a,b], K. Puolamäki [a]

[a] Helsinki Institute of Physics, FIN-00014 University of Helsinki, Finland
[b] Department of Physics, North Eastern Hill University, Shillong 793022, India [1]

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Abstract

We obtain a relation between right-handed gauge boson mass and soft squark mass in naturally R-parity conserving general supersymmetric left-right models. This relation implies that either W_R is lighter than twice the soft squark mass, or a ratio of vacuum expectation values (VEVs) in the model, denoted by $\tan\alpha$, is close to its value of unity in the limit of vanishing D-terms. Generally, we find that for heavy W_R $\tan\alpha$ is larger than one, and that the right-handed sneutrino VEV, responsible for spontaneous R-parity breaking, is at most of the order $M_{\rm SUSY}/h_{\Delta_R}$, where $M_{\rm SUSY}$ is supersymmetry breaking scale and h_{Δ_R} is the Yukawa coupling in Majorana mass term for right-handed neutrinos. These constraints follow from $SU(3)_c$ and $U(1)_{\rm em}$ gauge invariance of the ground state of the theory. © 1997 Elsevier Science B.V.

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Low energy supersymmetry is at present the only known extension of the Standard Model (SM) in which elementary Higgs scalar field, which is necessary to break the $SU(2)_L \times U(1)_Y$ symmetry, is natural [1]. The minimal supersymmetric standard model (MSSM) is constructed by simply doubling the number of degrees of freedom of the Standard Model, and adding an extra Higgs doublet (with opposite hypercharge) to cancel gauge anomalies and to generate masses for all quarks and leptons. However, one of the successes of SM, namely the automatic conservation of baryon number (B) and lepton number (L) by the renormalizable interactions, is not shared by the minimal supersymmetric standard model. In SM, the conservation of B and L follows from the particle content and the $SU(3)_c \times SU(2)_L \times U(1)_Y$ gauge invariance. In MSSM, baryon and lepton number violation can occur at tree level with catastrophic consequences unless the corresponding couplings are very small. The most common way to eliminate these tree level B and L violating terms is to impose a discrete Z_2 symmetry [2,3] known as a matter parity for superfields $(=(-1)^{3(B-L)})$ or equivalently R-parity on component fields ($R_p = (-1)^{3(B-L)+2S}$, S being the spin of the particle), with all the Standard Model particles having $R_p = +1$, while all the superpartners have $R_p = -1$. However, the assumption of R-parity conservation appears to be ad hoc, since it is not required for the internal consistency of the minimal supersymmetric standard model. Furthermore, all global symmetries,

[1] Permanent address.

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discrete or continuous, could be violated by the Planck scale physics effects [4]. The net effect would be the appearance of extremely tiny violations of otherwise exact global symmetries in the Lagrangians describing physics at the electroweak scale. The problem becomes acute for low energy supersymmetric models [5] because B and L are no longer automatic symmetries of the Lagrangian as they are in the Standard Model.

It would, therefore, be more appealing to have a supersymmetric theory where R-parity is related to a gauge symmetry, and its conservation is automatic because of the invariance of the underlying theory under an extended gauge symmetry. Indeed R_p conservation follows automatically in certain theories with gauged $(B-L)$, as is suggested by the fact that matter parity is simply a Z_2 subgroup of $(B-L)$. It has been noted by several authors [6,7] that if the gauge symmetry of MSSM is extended to $SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L}$ or $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, the theory becomes automatically R-parity conserving. Such a left-right supersymmetric theory (SUSY LR) solves the problems of explicit B and L violation of MSSM and has received much attention recently [8–12]. In particular, the dynamical issues connected with automatic R-parity conservation have been considered in case of left-right supersymmetric models [9]. It has been found in a wide class of such models [9] that R-parity must be spontaneously broken [13] because of the form of the scalar potential, although this could be avoided if non-renormalizable interactions are included. Thus, this model cures one of the major drawbacks of MSSM, although it leads to small L-violating terms, which are, however, suppressed and could be tested experimentally by searching for lepton number violation.

Furthermore, it has been shown [14] in the minimal SUSYLR model that the mass (m_{W_R}) of the right-handed gauge boson W_R has an upper limit related to the SUSY breaking scale, i.e., $m_{W_R} \leq gM_{\rm SUSY}/h_{\Delta_R}$, where g is the weak gauge coupling and h_{Δ_R} is the Yukawa coupling of the right-handed neutrinos with the triplet Higgs fields. This result is a consequence of electric charge conservation and low energy parity violation by the ground state of the theory.

In this letter we report on further consequences of SUSYLR models. For phenomenological reasons, it is desirable to constrain the magnitudes of the $SU(2)_R$ breaking scale, represented by m_{W_R}, and R-parity breaking scale, corresponding to the right-handed sneutrino VEV. Here we will find a relation between the W_R mass and the soft squark mass $\tilde{m}$, independently of any Yukawa couplings. In addition, we constrain the order of the right-handed sneutrino VEV to be generally at most $M_{\rm SUSY}/h_{\Delta_R}$ for large W_R mass. These results follow from the conservation of electric charge and color by the ground state of the theory.

We begin by recalling the basic features of the SUSYLR models. The matter fields of the minimal model consist of the quark and lepton doublets, $Q(2,1,1/3)$; $Q^c(1,2,-1/3)$; $L(2,1,-1)$; $L^c(1,2,1)$, and the Higgs multiplets consist of $\Delta_L(3,1,-2)$; $\Delta_R(1,3,-2)$; $\delta_L(3,1,2)$; $\delta_R(1,3,2)$; $\Phi(2,2,0)$; $\chi(2,2,0)$. The numbers in the parentheses denote the representation content of the fields under the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. The most general superpotential of the model is given by

$$\begin{aligned} W = {} & h_{\phi Q} Q^T i\tau_2 \Phi Q^c + h_{\chi Q} Q^T i\tau_2 \chi Q^c \\ & + h_{\phi L} L^T i\tau_2 \Phi L^c + h_{\chi L} L^T i\tau_2 \chi L^c \\ & + h_{\delta_L} L^T i\tau_2 \delta_L L \\ & + h_{\Delta_R} L^{cT} i\tau_2 \Delta_R L^c + \mu_1 \mathrm{Tr}\left(i\tau_2 \Phi^T i\tau_2 \chi\right) \\ & + \mu_1' \mathrm{Tr}\left(i\tau_2 \Phi^T i\tau_2 \Phi\right) + \mu_1'' \mathrm{Tr}\left(i\tau_2 \chi^T i\tau_2 \chi\right) \\ & + \mathrm{Tr}\left(\mu_{2L} \Delta_L \delta_L + \mu_{2R} \Delta_R \delta_R\right). \end{aligned} \qquad (1)$$

The general form of the Higgs potential is given by

$$V = V_F + V_D + V_{\rm soft} \qquad (2)$$

and can be calculated in a straightforward manner. In the following we shall represent the scalar components of the superfields by the same symbols as the superfields themselves. The most general form of the vacuum expectation values of various scalar fields which preserves $Q_{\rm em}$ can be written as

$$\langle \Phi \rangle = \begin{pmatrix} \kappa_1 & 0 \\ 0 & e^{i\varphi_1}\kappa_1' \end{pmatrix}, \quad \langle \chi \rangle = \begin{pmatrix} e^{i\varphi_2}\kappa_2' & 0 \\ 0 & \kappa_2 \end{pmatrix},$$

$$\langle \Delta_L \rangle = \begin{pmatrix} 0 & v_{\Delta_L} \\ 0 & 0 \end{pmatrix}, \quad \langle \delta_L \rangle = \begin{pmatrix} 0 & 0 \\ v_{\delta_L} & 0 \end{pmatrix},$$

$$\langle \Delta_R \rangle = \begin{pmatrix} 0 & v_{\Delta_R} \\ 0 & 0 \end{pmatrix}, \quad \langle \delta_R \rangle = \begin{pmatrix} 0 & 0 \\ v_{\delta_R} & 0 \end{pmatrix},$$
$$\langle L \rangle = \begin{pmatrix} \sigma_L \\ 0 \end{pmatrix}, \quad \langle L^c \rangle = \begin{pmatrix} 0 \\ \sigma_R \end{pmatrix}. \tag{3}$$

For simplicity, we shall ignore the phases φ_1 and φ_2 in the following, although this does not affect the final conclusion. Due to the tiny mixing between the charged gauge bosons, κ_1' and κ_2' are taken to be much smaller than κ_1 and κ_2. Furthermore, since the electroweak ρ-parameter is close to unity, $\rho = 1.0002 \pm 0.0013 \pm 0.0018$ [15], the triplet vacuum expectation values $\langle \Delta_L \rangle$ and $\langle \delta_L \rangle$ must be small. For definiteness, we shall take $v_{\Delta_R} \sim v_{\delta_R} \sim v_R$, the generic scale of the right-handed symmetry breaking. In this class of models the spontaneous breakdown of R-parity is inevitable [9], and therefore we shall assume that σ_L and σ_R are non-zero. Since σ_L contributes to W_L mass, it is less than the weak scale. On the other hand, in electric charge preserving ground state σ_R is necessarily at least of the order of the typical SUSY breaking scale $M_{\rm SUSY}$ or the right-handed breaking scale v_R, whichever is lower [14].

We next construct the up- and down-squark mass matrices. To this end we write down explicitly the different components of the scalar potential (2) (g_L, g_R, g_{B-L} are the gauge couplings)

$$\begin{aligned} V_F = & |h_{\phi Q} Q^c Q^T (i\tau_2) + h_{\phi L} L^c L^T (i\tau_2) \\ & + \mu_1 (i\tau_2) \chi^T (i\tau_2) + 2\mu_1' (i\tau_2) \Phi^T (i\tau_2)|^2 \\ & + |h_{\chi Q} Q^c Q^T (i\tau_2) + h_{\chi L} L^c L^T (i\tau_2) \\ & + \mu_1 (i\tau_2) \Phi^T (i\tau_2) + 2\mu_1'' (i\tau_2) \chi^T (i\tau_2)|^2 \\ & + |(i\tau_2)(h_{\phi Q}\Phi + h_{\chi Q}\chi) Q^c|^2 \\ & + |Q^T (i\tau_2)(h_{\phi Q}\Phi + h_{\chi Q}\chi)|^2, \end{aligned} \tag{4}$$

$$\begin{aligned} V_D = & \tfrac{1}{8} g_L^2 \sum_a \big(\mathrm{Tr}(\Phi^\dagger \tau_a \Phi) + \mathrm{Tr}(\chi^\dagger \tau_a \chi) \\ & + 2\mathrm{Tr}(\Delta_L^\dagger \tau_a \Delta_L) + 2\mathrm{Tr}(\delta_L^\dagger \tau_a \delta_L) + L^\dagger \tau_a L \\ & + Q^\dagger \tau_a Q\big)^2 + \tfrac{1}{8} g_R^2 \sum_a \big(-\mathrm{Tr}(\Phi \tau_a \Phi^\dagger) \\ & - \mathrm{Tr}(\chi \tau_a \chi^\dagger) + 2\mathrm{Tr}(\Delta_R^\dagger \tau_a \Delta_R) \\ & + 2\mathrm{Tr}(\delta_R^\dagger \tau_a \delta_R) + L^{c\dagger} \tau_a L^c + Q^{c\dagger} \tau_a Q^c\big)^2 \\ & + \tfrac{1}{8} g_{B-L}^2 \big(-2\mathrm{Tr}(\Delta_R^\dagger \Delta_R) \\ & + 2\mathrm{Tr}(\delta_R^\dagger \delta_R) - 2\mathrm{Tr}(\Delta_L^\dagger \Delta_L) + 2\mathrm{Tr}(\delta_L^\dagger \delta_L) \\ & - L^\dagger L + L^{c\dagger} L^c + \tfrac{1}{3} Q^\dagger Q - \tfrac{1}{3} Q^{c\dagger} Q^c \big)^2, \end{aligned} \tag{5}$$

$$\begin{aligned} V_{\rm soft} = & \tilde{m}_Q^2 |Q|^2 + \tilde{m}_{Q^c}^2 |Q^c|^2 \\ & + \big(Q^T i\tau_2 (B_\phi \Phi + B_\chi \chi) Q^c + \text{h.c.}\big), \end{aligned} \tag{6}$$

where we have retained only those terms which are relevant for our discussion.

From (4), (5), and (6) it is straightforward to construct the various squark mass matrices. For definiteness we consider the up- and down-squark mass matrices for the lightest generation (ignoring the intergenerational mixing) which gives the tightest constraint in our case. The part of the potential containing the squark mass terms can be written as

$$\begin{aligned} V_{\rm squark} = & \begin{pmatrix} U_L^* & U_R^* \end{pmatrix} \tilde{M}_U \begin{pmatrix} U_L \\ U_R \end{pmatrix} \\ & + \begin{pmatrix} D_L^* & D_R^* \end{pmatrix} \tilde{M}_D \begin{pmatrix} D_L \\ D_R \end{pmatrix}. \end{aligned} \tag{7}$$

The mass matrix elements for the up-type squarks are

$$\begin{aligned} (\tilde{M}_U)_{U_L^* U_L} = & \tilde{m}_Q^2 + m_u^2 + \tfrac{1}{4} g_L^2 (\omega_\kappa^2 - 2\omega_L^2) \\ & + \tfrac{1}{6} g_{B-L}^2 (\omega_L^2 - \omega_R^2), \\ (\tilde{M}_U)_{U_R^* U_L} = & B_\phi \kappa_1' + B_\chi \kappa_2 - \mu_1 (h_{\phi Q} \kappa_2' + h_{\chi Q} \kappa_1) \\ & - 2 h_{\phi Q} \mu_1' \kappa_1 - 2 h_{\chi Q} \mu_1'' \kappa_2' \\ & + (h_{\phi L} h_{\phi Q} + h_{\chi L} h_{\chi Q}) \sigma_L \sigma_R \\ = & \big[(\tilde{M}_U)_{U_L^* U_R}\big]^*, \\ (\tilde{M}_U)_{U_R^* U_R} = & \tilde{m}_{Q^c}^2 + m_u^2 + \tfrac{1}{4} g_R^2 (\omega_\kappa^2 - 2\omega_R^2) \\ & + \tfrac{1}{6} g_{B-L}^2 (\omega_R^2 - \omega_L^2), \end{aligned} \tag{8}$$

and for down-type squarks

$$\begin{aligned} (\tilde{M}_D)_{D_L^* D_L} = & \tilde{m}_Q^2 + m_d^2 - \tfrac{1}{4} g_L^2 (\omega_\kappa^2 - 2\omega_L^2) \\ & + \tfrac{1}{6} g_{B-L}^2 (\omega_L^2 - \omega_R^2), \\ (\tilde{M}_D)_{D_R^* D_L} = & -B_\phi \kappa_1 - B_\chi \kappa_2' + \mu_1 (h_{\phi Q} \kappa_2 + h_{\chi Q} \kappa_1') \\ & + 2 h_{\phi Q} \mu_1' \kappa_1' + 2 h_{\chi Q} \mu_1'' \kappa_2 = \big[(\tilde{M}_D)_{D_L^* D_R}\big]^*, \end{aligned}$$

$$\left(\tilde{M}_D\right)_{D_R^* D_R} = \tilde{m}_{Q^c}^2 + m_d^2 - \tfrac{1}{4} g_R^2 \left(\omega_\kappa^2 - 2\omega_R^2\right) + \tfrac{1}{6} g_{B-L}^2 \left(\omega_R^2 - \omega_L^2\right), \tag{9}$$

where

$$m_u = h_{\phi Q} \kappa_1' + h_{\chi Q} \kappa_2, \quad m_d = h_{\phi Q} \kappa_1 + h_{\chi Q} \kappa_2', \tag{10}$$

and

$$\omega_L^2 = v_{\delta_L}^2 - v_{\Delta_L}^2 - \tfrac{1}{2}\sigma_L^2, \quad \omega_R^2 = v_{\Delta_R}^2 - v_{\delta_R}^2 - \tfrac{1}{2}\sigma_R^2,$$

$$\omega_\kappa^2 = \kappa_1^2 + \kappa_2'^2 - \kappa_2^2 - \kappa_1'^2. \tag{11}$$

In order not to break electromagnetism or color, none of the physical squared masses of squarks can be negative. Necessarily then all the diagonal elements of the squark mass matrices should be non-negative. Since the SUSY breaking scale is expected to be $\mathcal{O}(1\ \text{TeV})$, the weak scale is smaller than the other scales involved, v_R, σ_R or M_{SUSY}. We'll ignore the terms of the order of the weak scale or smaller in the following. Combining the diagonal elements of the mass matrices $\tilde{M}_U$ and $\tilde{M}_D$, it follows that

$$\tilde{m}_Q^2 + \tilde{m}_{Q^c}^2 \geq |\tfrac{1}{2} g_R^2 \omega_R^2| = \tfrac{1}{2} g_R^2 |v_{\Delta_R}^2 - v_{\delta_R}^2 - \tfrac{1}{2}\sigma_R^2|. \tag{12}$$

In arriving at (12), we have made no assumptions other than that of ignoring those quantities which are of the order of weak scale.

We define next an angle α with $\tan^2\alpha = (v_{\delta_R}^2 + \frac{1}{2}\sigma_R^2)/v_{\Delta_R}^2$ and write $\tilde{m}_Q^2 = \tilde{m}_{Q^c}^2 \equiv \tilde{m}^2$. Then Eq. (12) can be written as

$$\tilde{m}^2 \geq \tfrac{1}{4} g_R^2 v_{\Delta_R}^2 |1 - \tan^2\alpha|. \tag{13}$$

In order to understand the result (12) or (13), we recall that the right-handed gauge boson mass is given by (ignoring weak scale effects) [12]

$$m_{W_R}^2 = g_R^2\left(v_{\Delta_R}^2 + v_{\delta_R}^2 + \tfrac{1}{2}\sigma_R^2\right) = g_R^2 v_{\Delta_R}^2 (1 + \tan^2\alpha). \tag{14}$$

Combining (13) and (14), we find

$$m_{W_R}^2 |\cos 2\alpha| \leq 4\tilde{m}^2. \tag{15}$$

If the W_R boson is lighter than twice the soft squark mass $\tilde{m}$, Eq. (15) is fulfilled for any $\tan\alpha$. If $\tilde{m} \sim M_{\text{SUSY}} \sim 1$ TeV as is commonly assumed, m_{W_R} cannot be much less, since experimentally $m_{W_R} > 420$ GeV [16]. On the other hand, if m_{W_R} is much larger than $2\tilde{m}$, $\tan\alpha$ has to be close to one, e.g. for $m_{W_R} = 10$ TeV and $\tilde{m} = 1$ TeV, one would need $0.96 \leq \tan\alpha \leq 1.04$. It is interesting to note in this context that the vanishing of D-terms implies $\tan\alpha = 1$. To translate the limit for $\tan\alpha$ to an upper bound for the VEV $\langle \Delta_R^0 \rangle$, one needs a lower limit for g_R. This was found in [17] from $\sin^2\theta_W = e^2/g_L^2 = 0.23$, namely $g_R \geq 0.55\, g_L$. Consequently $v_{\Delta_R} \leq m_{W_R} |\cos\alpha|/(0.55\, g_L)$, e.g. in our example $v_{\Delta_R} \lesssim 20$ TeV.

To further analyze the situation with large m_{W_R}, we note that, if the right-handed scale and R-parity breaking scale differ from each other, one has $v_{\Delta_R}, v_{\delta_R} > \sigma_R$, since $\tan\alpha \sim 1$. We recall then the doubly charged Higgs mass matrix [12] given by (ignoring terms suppressed by σ_R/v_{Δ_R} or σ_R/v_{δ_R})

$$M_{\Delta^{++}\delta^{++}}^2 = \begin{pmatrix} m_{\Delta\delta}^2 \frac{v_\delta}{v_\Delta} - 4h_\Delta^2\sigma_R^2 - 2g_R^2\omega_R^2 & -m_{\Delta\delta}^2 \\ -m_{\Delta\delta}^2 & m_{\Delta\delta}^2 \frac{v_\Delta}{v_\delta} + 2g_R^2\omega_R^2 \end{pmatrix}, \tag{16}$$

where $m_{\Delta\delta}$ is the soft parameter mixing right-handed Higgs triplets. The two eigenvalues of the mass matrix need to be real and non-negative in order not to break $U(1)_{\text{em}}$. This leads to two conditions:

$$h_{\Delta_R}^2 \sigma_R^2 \leq \tfrac{1}{4} m_{\Delta\delta}^2 \left(\frac{v_{\delta_R}}{v_{\Delta_R}} + \frac{v_{\Delta_R}}{v_{\delta_R}} \right) \tag{17}$$

and

$$-2m_{\Delta\delta}^2 g_R^2 \omega_R^2 \left(\frac{v_{\delta_R}}{v_{\Delta_R}} - \frac{v_{\Delta_R}}{v_{\delta_R}} \right) + 4g_R^4\omega_R^4 + 4m_{\Delta\delta}^2 h_{\Delta_R}^2 \sigma_R^2 \frac{v_{\Delta_R}}{v_{\delta_R}} + 8h_{\Delta_R}^2 \sigma_R^2 g_R^2 \omega_R^2 \leq 0 \tag{18}$$

From (17) we see that $h_{\Delta_R}\sigma_R$ can be at most of the order of $m_{\Delta\delta} \sim M_{\text{SUSY}}$. To fulfill the inequality (18) we can consider two cases: $v_{\Delta_R} < v_{\delta_R}$ and $v_{\delta_R} < v_{\Delta_R}$. In both cases one must have $\omega_R^2 \leq 0$ or equivalently $\tan\alpha \geq 1$. The equality can hold for $\sigma_R = 0$ and $v_{\delta_R} = v_{\Delta_R}$.

In conclusion, we have studied the implications of SUSY left-right models, which naturally incorporate R-parity conservation. We have found that the W_R mass and the soft squark mass are related by (15), which implies that either the scale of the right-handed gauge symmetry breaking must be close to the SUSY breaking scale, or $\tan\alpha \sim 1$ corresponding to vanishing D-terms. In general, we have also found that for large m_{W_R}, the right-handed sneutrino VEV is constrained to be at most of the order $M_{\mathrm{SUSY}}/h_{\Delta_R}$, and that $\tan\alpha$ is larger than one.

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