Minimal seesaw as an ultraviolet insensitive cure for the problems of anomaly mediation

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(Received 11 September 2007; published 27 March 2008)

We show that an intermediate scale supersymmetric left-right seesaw scenario with automatic R-parity conservation can cure the problem of tachyonic slepton masses that arise when supersymmetry is broken by anomaly mediation, while preserving ultraviolet insensitivity. The reason for this is the existence of light $B - L = 2$ Higgses with Yukawa couplings to the charged leptons. We find these theories to have distinct predictions compared to the usual mSUGRA and gauge mediated models as well as the minimal anomaly mediated supersymmetry breaking models. Such predictions include a condensed gaugino mass spectrum and possibly a correspondingly condensed sfermion spectrum.

I. INTRODUCTION

The driving motivations for physics beyond the standard model are: (i) stabilizing the Higgs mass against radiative corrections, thereby providing a firmer understanding of the origin of mass; (ii) understanding the neutrino masses and mixings, especially the extreme smallness of its mass compared to those of charged fermions; (iii) finding a particle physics candidate for the dark matter of the universe; and (iv) solving the strong CP problem. Two prevalent ideas for resolving them are: supersymmetry (SUSY)—curing (i); and the seesaw mechanism [1–5]—curing (ii)—making SUSY seesaw very enticing. R-parity is assured as an automatic symmetry of the low energy Lagrangian [6–8] given $B - L$ is a gauged symmetry broken by $B - L = 2$ Higgs fields (these fields have independent motivations [9]). Conservation of R-parity would guarantee a stable lightest SUSY particle (LSP), providing a good dark matter candidate [10] as well as preventing catastrophic proton decay [caused by R-parity breaking terms of the minimal SUSY standard model (MSSM)]. Finally, gauged $B - L$ models embedded into the SUSY left-right framework provide a cure to the strong CP problem without the need for an axion [11,12]. This leads us to focus on the minimal SUSY left-right model and look at further constraints when the method of SUSY breaking is considered.

The nature and origin of SUSY breaking has been the focus of a great deal of attention. The detailed nature of new physics that breaks SUSY is unknown—although there are several interesting suggestions [10]. Here we focus on SUSY breaking via anomaly mediation which is related to the radiative breaking of conformal SUSY [13,14]. Anomaly mediated SUSY breaking (AMSB) predicts all the soft SUSY breaking parameters in terms of one mass parameter (the mass of the gravitino) and the β and γ functions of the low energy theory. As such, it is a highly predictive scenario which avoids the SUSY flavor problem (no new flavor physics is introduced) and solved the gravitino mass problem.

There is, however, a serious problem that afflicts any AMSB model whose low energy theory is the MSSM: the sleptons have negative mass-squared thereby leading to a ground state that breaks electric charge. Finding a cure to this problem is a difficult task given the predictability of AMSB and the fact that it generally decouples higher-scale physics. This forces solutions to include new couplings in the low energy theory or deflecting the AMSB trajectories. There are many proposed solutions along these lines [15–24], but the model we present here has the advantage of being well motivated by the seesaw mechanism.

In this paper we propose a new way to resolve this problem of AMSB using the minimal R-parity conserving SUSYLR seesaw model mentioned above. We present an instance of this class of bottom-up seesaw models that has an intermediate seesaw scale (of order 10^{11} GeV or so) and show that the slepton mass-square problem of AMSB is cured. Furthermore, ultraviolet (UV) insensitivity is preserved; a featured that is shared with only a few of the proposed AMSB solutions. The key to this is the existence of light doubly charged Higgses [25] and light left-handed triplets and their Yukawa couplings to the lepton superfields. The effects of these doubly charged fields can be discovered in low energy experiments as they lead to characteristic mass predictions which are different from those of other SUSY breaking scenarios. We will demonstrate these differences between our model, minimal supergravity (mSUGRA), minimal gauge mediated SUSY breaking (mGMSB), and AMSB with a universal scalar mass addition m_0^2 (mAMSB). Apart from experimental testability, a novel feature of our suggestion is that the cure is motivated from independent considerations. These are the defining phenomenological conditions:

(i) SUSY
(ii) local $B - L$ symmetry as part of the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ so that one can implement the seesaw mechanism
(iii) $B - L$ symmetry breaking is such that it leaves R-parity unbroken and assuring that there is a naturally stable dark matter candidate

(iv) SUSY is broken radiatively by conformal anomalies, hence keeping the soft terms (renormalization group equation) RGE invariant down to the TeV scale (UV insensitivity).

We will show in Sec. III how these considerations produce slepton positive mass-squares, as well as introduce the model and give its sparticle spectrum. We will also show that low energy experimental constraints combined with AMSB decoupling dictate that the doubly charged Higgs masses must be greater than several TeV and less than or around F_ϕ (the gravitino mass). Section II will give a brief overview of AMSB and introduce its notation and terminology.

II. SUMMARY OF AMSB

AMSB has many attractive features: a large number of predictions, few parameters, an insensitivity to the UV, and a mathematical framework that elegantly describes its effects. The latter property allows one to express the SUSY breaking effects by analytically continuing parameters into superspace. AMSB then gives a method or set of rules on how to "promote" these parameters to superfields [26,27]: starting out with a conformally invariant Lagrangian (i.e. one with no dimensionful couplings), quantum corrections introduce a mass parameter, μ, which leads the renormalized parameters of the theory to be promoted to superfields by the rule

$$\mu \rightarrow \frac{\mu}{\sqrt{\phi^\dagger \phi}}. \tag{2.1}$$

The promotion of $Z(\mu)$ to a superfield $\mathcal{Z}(\mu)$ and $1/g^2(\mu)$ to the superfield $\mathcal{R}(\mu)$ gives rise to soft SUSY breaking terms. To obtain an expression for those terms it is convenient to choose a gauge where

$$\phi = 1 + F_\phi \theta^2. \tag{2.2}$$

This leads to the following form for the soft SUSY breaking parameters:

$$m_Q^2 = \frac{1}{2}|F_\phi|^2 \frac{d}{d\ln\mu} \gamma_Q(Y(\ln\mu), g_b(\ln\mu))$$

$$= \frac{1}{2}|F_\phi|^2 \left[\frac{1}{2} \beta_{g_a} \frac{\partial \gamma_Q}{\partial g_a} + \beta_Y \frac{\partial \gamma_Q}{\partial Y} + \text{c.c.} \right] \tag{2.3}$$

$$A_Q = \beta_{Y_Q} F_\phi \tag{2.4}$$

$$M_{\lambda_a} = \frac{\beta_{g_a}}{g_a} F_\phi, \tag{2.5}$$

where we are using the convention of [10]; that is, $\gamma_Q = -\frac{1}{2} \frac{d\ln Z}{d\ln\mu}$, $\beta_{g_a} = \frac{dg_a}{d\ln\mu}$, etc.

For the MSSM the contribution of the second term in Eq. (2.3) is negligible for at least the first and second generation sleptons, so the first term dominates. As both

$SU(2)_L \times U(1)_Y$ are infrared free, their contribution is negative and hence the sleptons get negative mass-squares. This is a fundamental problem because it implies the breakdown of electric charge in the ground state. Before AMSB models can be phenomenologically viable this problem must be solved, but it is worth pursuing a solution because AMSB provides decoupling of UV physics in an elegant manner (we discuss this below), naturally suppressed flavor-changing neutral-current interactions (the SUSY breaking parameters depend on the Yukawa couplings and are diagonalized with them), and high predictive power with a minimal number of arbitrary parameters (essentially all soft SUSY breaking terms depend on only F_ϕ). It is therefore of great interest to seek reasonable models where the slepton mass-squares are made positive without destroying those good features. We will present such a model in Sec. III, where we demonstrate that extending the MSSM to include neutrino mass—generated by an R-parity conserving seesaw mechanism—will simply and effectively achieve this goal. Yet for the moment we will continue our review of AMSB and address the decoupling of higher-scale physics.

To illustrate the UV insensitivity of AMSB, consider a threshold $\Lambda \gg M \gg F_\phi$—such a scale may be an explicit mass term in the superpotential or the vev of the scalar component of the superfield X. In either case we assume that below M there are no remnant singlets in the effective theory; this is the same as requiring that as $\Lambda \rightarrow \infty$, M remains finite. The previous condition ensures that the effective theory's Lagrangian has the schematic form[1]

$$\mathcal{L}_{\text{eff}} = \mathcal{L}_Q + M^{-n} f(\underline{Q}, \psi_Q + M^4$$

$$= \mathcal{L}_Q + M^4 + \int d^2\theta \left(\frac{Q^4}{M\phi} + \frac{Q^5}{M^2\phi^2} + \cdots \right) + \text{H.c.,} \tag{2.6}$$

where $n > 0$ and $\mathcal{L}_Q$ represents the part of the Lagrangian involving only the various components of the matter superfields Q. This form of the effective theory makes explicit that the additional SUSY breaking effects from the threshold M go as $F_\phi/M \ll 1$. Thus, the rule $\mu \rightarrow \mu/|\phi|$ completely parametrizes all the SUSY breaking in both the high-scale and low-scale theories resulting in the maintenance of the AMSB trajectory below M.

Another means to see this decoupling is to realize that the replacement rule due to the threshold is

$$M \rightarrow M\phi + F_\phi \left[c_1 \frac{F_\phi}{M} + c_2 \left(\frac{F_\phi}{M} \right)^2 + \cdots \right] \theta^2 \approx M\phi \tag{2.7}$$

in addition to which there is the requirement pairing Λ with ϕ. The quantum corrections of the Lagrangian of Eq. (2.6)

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[1]We denote the scalar component of the superfield X as $\underline{X}$ to avoid confusing the two.

force M to appear in the effective theory as $\ln|M/\Lambda|$ and $\ln\mu/|M|$ (which comes when M is used as a cutoff in loop calculations). Using the replacement rules on these quantities gives

$$\ln\left|\frac{M}{\Lambda}\right| \rightarrow \ln\left|\frac{M\phi}{\Lambda\phi}\right| = \ln\left|\frac{M}{\Lambda}\right| \qquad (2.8)$$

$$\ln\frac{\mu}{|M|} \rightarrow \ln\frac{\mu}{|M\phi|} \qquad (2.9)$$

and once again only $\mu \rightarrow \mu/|\phi|$ is required to capture all the SUSY breaking.

The above argument may disturb the reader since the β functions change when crossing the threshold; however, what is actually happening is that when the threshold is crossed, the removal of the heavy fields adds a term $\Delta\beta$ that results in a shift of the higher-scale β functions, β^+, to the lower-scale β functions, β^-. This property, namely

$$\Delta\beta + \beta^+ = \beta^- \qquad (2.10)$$

is the one that keeps the theory in the AMSB form.

The UV decoupling of AMSB presents a major obstacle for fixing the negative mass-squares of the MSSM since any high-scale tinkering will leave little to no evidence at the low scale.

III. SUSYLR AND AMSB

The new feature of models combining AMSB and SUSYLR is that the effective theory below the ν_R scale contains Yukawa couplings to both the left- and right-handed electrons in addition to those of the MSSM; hence, the slepton masses can be made positive. Thus, the marriage of SUSYLR with AMSB gives positive slepton mass-squares and the resulting theory combines the predictive power of AMSB, explains small neutrino masses (through the seesaw mechanism), and retains a natural dark matter candidate (the LSP is stable due to R-parity conservation).

TABLE I. Assignment of the fermion and Higgs fields' representations of the left-right symmetry group [except for $U(1)_{B-L}$ where the charge under that group is given].

Fields	$SU(3)^c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$
Q	$(3, 2, 1, +\frac{1}{3})$
Q^c	$(\bar{3}, 1, 2, -\frac{1}{3})$
L	$(1, 2, 1, -1)$
L^c	$(1, 1, 2, +1)$
Φ	$(1, 2, 2, 0)$
Δ	$(1, 3, 1, +2)$
$\bar{\Delta}$	$(1, 3, 1, -2)$
Δ^c	$(1, 1, 3, -2)$
$\bar{\Delta}^c$	$(1, 1, 3, +2)$

A. The model

1. Definition and properties

The particle content of a SUSYLR model is shown in Table I. As the model is left-right symmetric, it contains both left- and right-handed Higgs bosons—in this case $B - L = \pm 2$ triplets so that R-parity may be preserved (a task for which the $B - L = 1$ doublets are not suitable). The presence of $SU(2)_L$ and $SU(2)_R$ triplets means that parity is a good symmetry until $SU(2)_R$ breaks. While the seesaw mechanism may be achieved with only $SU(2)_R$ Higgs fields, demanding parity forces the left-handed triplets to be present and these together then yield positive slepton masses.

The parity-respecting SUSYLR superpotential[2] is then

$$W_{\text{SUSYLR}} = W_Y + W_H + W_{\text{NR}} \qquad (3.1)$$

with

$$W_Y = iy_Q^a Q^T \tau_2 \Phi_a Q^c + iy_L^a L^T \tau_2 \Phi_a L^c + if_c L^{cT} \tau_2 \Delta^c L^c$$
$$+ ifL^T \tau_2 \Delta L \qquad (3.2)$$

$$W_H = (M_\Delta \phi - \lambda_S S)[\text{Tr}(\Delta^c \bar{\Delta}^c) + \text{Tr}(\Delta \bar{\Delta})] + \tfrac{1}{2}\mu_S \phi S^2$$
$$+ \tfrac{1}{3}\kappa_S S^3 + \lambda_N^{ab} N \,\text{Tr}(\Phi_a^T \tau_2 \Phi_b \tau_2) + \tfrac{1}{3}\kappa_N N^3 \qquad (3.3)$$

$$W_{\text{NR}} = \frac{\lambda_A}{M_{\text{Pl}}\phi}\,\text{Tr}^2(\Delta\bar{\Delta}) + \frac{\lambda_A^c}{M_{\text{Pl}}\phi}\,\text{Tr}^2(\Delta^c\bar{\Delta}^c) + \frac{\lambda_B}{M_{\text{Pl}}\phi}\,\text{Tr}(\Delta\Delta)\,\text{Tr}(\bar{\Delta}\,\bar{\Delta}) + \frac{\lambda_B^c}{M_{\text{Pl}}\phi}\,\text{Tr}(\Delta^c\Delta^c)\,\text{Tr}(\bar{\Delta}^c\bar{\Delta}^c)$$
$$+ \frac{\lambda_C}{M_{\text{Pl}}\phi}\,\text{Tr}(\Delta\bar{\Delta})\,\text{Tr}(\Delta^c\bar{\Delta}^c) + \frac{\lambda_D}{M_{\text{Pl}}\phi}\,\text{Tr}(\Delta\Delta)\,\text{Tr}(\Delta^c\Delta^c) + \frac{\bar{\lambda}_D}{M_{\text{Pl}}\phi}\,\text{Tr}(\bar{\Delta}\,\bar{\Delta})\,\text{Tr}(\bar{\Delta}^c\bar{\Delta}^c) + \cdots \qquad (3.4)$$

and the fields defined as

$$L = \begin{pmatrix} \nu \\ e \end{pmatrix}, \qquad L^c = \begin{pmatrix} e^c \\ -\nu^c \end{pmatrix}, \qquad \text{etc.} \qquad (3.5)$$

We have assumed that the singlet couplings absent from Eq. (3.3) are zero or small enough that they can be neglected. This condition is necessary to keep one singlet light (N) so that below the right-handed scale ν_R the theory is the next-to minimal supersymmetric standard model (NMSSM) with some additional particles. Although this may seem rather *ad hoc*, we do it out of convenience rather

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[2]In the following superpotential we suppress the generational and color indices for brevity.

than necessity: the low-scale theory must be such that it avoids an MSSM Higgs bilinear b term that is too large [28]; the superpotential given above happens to be one. However, it is not the *only* one and several alternative methods exist [13,17] to avoid this problem. As any of these alternatives are equally valid, and because the exact form of the electroweak scale theory is irrelevant to the conclusions, we merely select to use the superpotential above.

The superpotential of Eq. (3.3) dictates that

$$\langle S \rangle = \frac{M_\Delta}{\lambda_S} \phi \qquad (3.6)$$

$$\langle \Delta^c \rangle \langle \bar{\Delta}^c \rangle = \langle S \rangle \Big(\frac{M_\Delta \kappa_S}{\lambda_S^2} + \frac{\mu_S}{\lambda_S} \Big) \phi. \qquad (3.7)$$

Equation (3.6) should be evident from the form of the superpotential; Eq. (3.7) requires Eq. (3.3) to be recast as

$$W_{\mathrm{H}} \supset [-\lambda_S \mathrm{Tr}(\Delta^c \bar{\Delta}^c) + \tfrac{1}{2}\mu_S \phi S + \tfrac{1}{3}\kappa_S S^2]S. \qquad (3.8)$$

The inclusion of the nonrenormalizable terms of Eq. (3.4) (which are necessary if R-parity is conserved [25,29]) will shift the vevs of $\underline{\Delta}^c$, $\bar{\Delta}^c$, and $\underline{S}$ by $\sim M_\Delta^2/M_{\mathrm{Pl}} \ll M_\Delta$ so they may be safely ignored. It is worth noting that the nonrenormalizable terms are only irrelevant because as $M_{\mathrm{Pl}} \rightarrow \infty$ the vevs all remain finite; that is, they depend at most on $1/M_{\mathrm{Pl}}$.

2. Neutrino masses

When the $SU(2)_R$ triplets acquire a vev they induce a large mass term for the right-handed neutrinos from the term

$$\mathrm{i} f_c L^{cT} \tau_2 \Delta^c L^c \supset -(f_c)_{ij} \Delta^{c0} \nu_i^c \nu_j^c \rightarrow -\langle \Delta^{c0} \rangle (f_c)_{ij} \nu_i^c \nu_j^c. \qquad (3.9)$$

As $v_R \gtrsim 10^{10}$ GeV, this term is the well-known large Majorana mass for right-handed neutrinos. With the inclusion of electroweak symmetry breaking, the neutrinos have a mass matrix, in the basis (ν_i, ν_j^c):

$$M_\nu = \begin{pmatrix} 0 & y_L^a \langle \Phi_u^0 \rangle_a \\ (y_L^a)^T \langle \Phi_u^0 \rangle_a & f_c \langle \Delta^{c0} \rangle \end{pmatrix} \qquad (3.10)$$

which may be block-diagonalized to obtain the lighter states masses, m_ν, using the type I seesaw mechanism formula:

$$m_\nu = -\frac{\langle \Phi_u^0 \rangle_a \langle \Phi_u^0 \rangle_b}{\langle \Delta^{c0} \rangle} y_L^a f_c^{-1} (y_L^b)^T. \qquad (3.11)$$

Experimental constraints on lepton flavor violation [30] will compel f_c to be diagonal in flavor space, while neutrino mixings may be obtained through flavor violation in the y_L^a's. Since one linear combination of the y_L^a's forms the conventional low-scale y_e while the orthogonal set forms y_ν, this will not reintroduce a flavor violation problem for

the electron and its siblings. Additionally, because ν^c has such a large mass it is integrated out at the right-handed scale, taking y_ν (and its flavor violation) out of the renormalization group equations. These considerations ensure that choosing f_c and f diagonal at the right-handed scale will not introduce significant radiative corrections when running below that scale.

3. Curing AMSB

With the model defined and neutrino masses reviewed, the opportunity is ripe to explore how this all comes together to fix AMSB. As stated previously, no new SUSY breaking effects are introduced at the right-handed scale; therefore, Eqs. (3.6) and (3.7) are still valid below this regime and the theory respects the AMSB trajectory. Yet even though the particles remain on their AMSB trajectory, the negative slepton mass-squares problem is still solved. This comes about because of the additional Yukawa couplings f and f_c which survive to the lower-scale theory.

The existence of the f coupling at the lower scale can be seen from the superpotential Eq. (3.3): when S gets the vev of Eq. (3.6), the mass term for the $SU(2)_L$ triplets vanishes while the $SU(2)_R$ triplets also get a vev, so their mass term remains. This would leave Δ and $\bar{\Delta}$ massless below the right-handed breaking scale except that the nonrenormalizable terms contribute a mass through

$$\frac{\lambda_C}{M_{\mathrm{Pl}} \phi} \mathrm{Tr}(\Delta \bar{\Delta}) \mathrm{Tr}(\Delta^c \bar{\Delta}^c) \rightarrow \frac{\lambda_C}{M_{\mathrm{Pl}} \phi} \langle \Delta^c \rangle \langle \bar{\Delta}^c \rangle \mathrm{Tr}(\Delta \bar{\Delta})$$

$$\simeq \frac{\lambda_C v_R^2 \phi}{M_{\mathrm{Pl}}} \mathrm{Tr}(\Delta \bar{\Delta}). \qquad (3.12)$$

The same mass value of v_R^2/M_{Pl} is also responsible for f_c surviving to the low scale, but this time in the context of light doubly charged particles. It is well known that the class of SUSYLR models considered here have light doubly charged particles [25] with a mass as mentioned above. The question that needs to be addressed here is "how light?". If their mass is large, $F_\phi \ll m_{\mathrm{DC}} \ll v_R$, then these particles merely introduce another trajectory preserving threshold which decouples from the lower-scale theory. For the right-handed selectron, this would be disastrous as it would have a purely negative AMSB contribution to its mass. Thus, it makes sense to demand that the doubly charged particles have a mass $m_{\mathrm{DC}} \sim F_\phi$.

The existence of the $SU(2)_L$ triplets and the doubly charged particles below or around $m_{3/2}$ means that their couplings remain in the low-scale superpotential and are therefore important. For the sleptons, the relevant terms are

$$W \supset f_c \Delta^{c--} e^c e^c + \mathrm{i} f L^T \tau_2 \Delta L. \qquad (3.13)$$

The survival of these Yukawa couplings f_c and f allows the scalar e^c and e mass-squares to be positive. Assuming that f, f_c are diagonal in flavor space (an assumption

validated by lepton flavor violating experiments [30]), we need only $f_1 \simeq f_2 \simeq f_{c1} \simeq f_{c2} \simeq \mathcal{O}(1)$ to make the sleptons positive. The only constraint here is from muonium-antimuonium oscillations [31] which demands that $f_{c1}f_{c2}/4\sqrt{2}m_{DC}^2 \approx f_1f_2/4\sqrt{2}m_{DC}^2 < 3 \times 10^{-3}G_F$; however, with both the doubly charged fields and $SU(2)_L$ triplets having a mass $m_{DC} \simeq F_\phi \sim 10$ TeV, this is easily satisfied. Furthermore, this constraint limits the range for v_R as $m_{DC} \simeq v_R^2/M_{Pl} \simeq F_\phi$ implies that $v_R \simeq 10^{11}$–10^{12} GeV.

The net result is that AMSB and SUSYLR yield a sfermion sector that depends on very few parameters: F_ϕ, f_{c1}, f_{c3}, in addition to the usual $\tan\beta$ and $\operatorname{sgn}\mu$ (because of parity, $f_1 = f_{c1}$ and $f_3 = f_{c3}$). Interestingly, two of the new parameters—the f_c Yukawa couplings—also have implications for neutrino oscillations.

B. Numerical analysis

We now present the resulting mass spectrum for this model. For this analysis we start by running the parameters of the standard model up to M_{SUSY}, match at that point to the NQNMSSM (not-quite NMSSM: the NMSSM with doubly charged particles, left-handed triplets, and two additional Higgs doublets), and use the appropriately modified RGEs of [32] to get to the right-handed scale. Without loss of generality, we assume that only one up-type Higgs and one down-type Higgs get a vev [33]. Additionally, we take the standard simplifying assumption that only the third generation Higgs Yukawa couplings are important.

Figure 1 shows the mass spectrum of the general SUSYLR model and the MSSM with other popular SUSY breaking scenarios (the figure is truly only comprehensible in color [34]). The comparison was obtained by matching the gluino mass between the models, and then running the masses down to the scale Q using ISAJET [35]. The spectra in Fig. 1 contain the generic features, though the figure was generated using the points listed in Table II. It is also interesting to note that the heavier sfermion mass eigenstates are mostly right-handed contrary to most mSUGRA and GMSB scenarios.

One of the more striking features of the SUSYLR + AMSB spectrum is that gaugino sector masses are all relatively close to each other. This is unique from the popular scenarios displayed in Fig. 1 and is due to the contributions of the $SU(2)_L$ and $U(1)_Y$ extended particle content at low energy. Such a massive wino consequently relaxes the naturalness arguments made in [36,37]. These arguments proceed along the lines that squark masses and the μ term must be below around 1 TeV to preserve the naturalness of SUSY. Therefore, a naturalness upper bound can be put on the wino mass. Such an upper bound suggests that run II of the tevatron should have explored most of the viable wino parameter space, which would not be the case here.
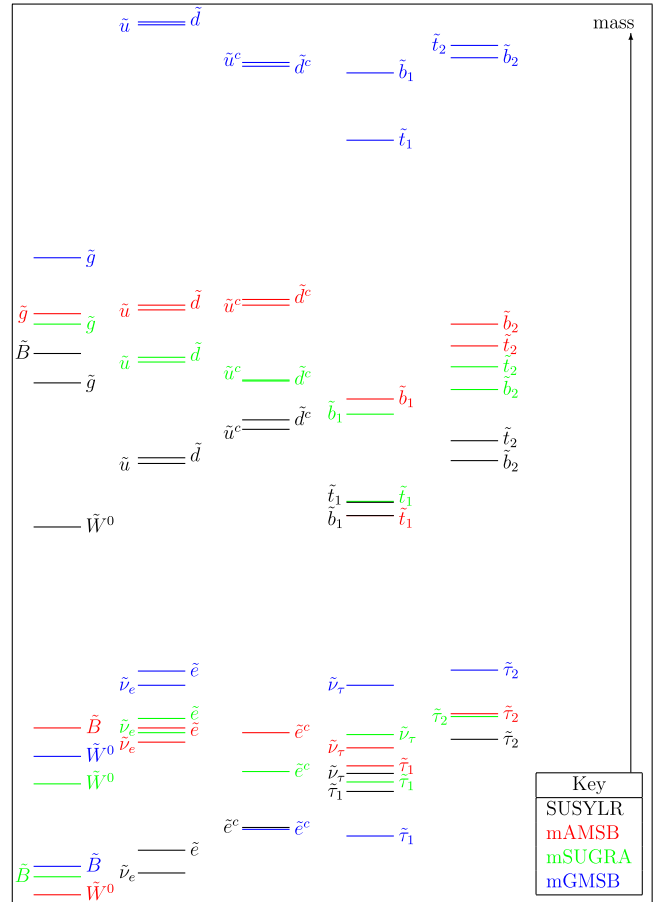

FIG. 1 (color online). The mass spectrum for the superpartners of the standard model for four different models (in four different colors [34]): SUSYLR + AMSB, mAMSB, mSUGRA, and mGMSB. Note that for the SUSYLR + AMSB, $\tilde{t}_2$ and $\tilde{b}_2$ are mostly right-handed; in contrast with the usual mSUGRA or mGMSB cases where they are typically mostly left-handed.

From a cosmological point of view, there is a potential problem with the increase in $SU(2)_L$ and $U(1)_Y$ gauge coupling strengths at the right-handed scale: they cause tachyonic squark masses at that scale [remember these gauge couplings give a negative contribution in Eq. (2.3)]. Theories with tachyonic squark masses have been studied in the GUT framework and were found to be safe albeit unsavory [38]. Reheating temperatures around 10^9 GeV will cause charge violating vacua to disappear [39], while still remaining below 10^{11} GeV so that the breaking of parity is not an issue. Also, the tunneling rates to the bad vacua are too small in most of the parameter space [40,41] to cause a problem.

Continuing along cosmological lines, both mass spectrums shown above indicate that the sneutrino is the LSP in this model. Both the tau and electron sneutrinos are LSP candidates depending on the relative sizes of f_3 and f_1. Although sneutrino dark matter is highly constrained [42,43], there could be other dark matter candidates such

TABLE II. The benchmark points for creating the spectrum of Fig. 1. The parameters shown were chosen by matching the gluino mass for SUSYLR + AMSB to mGMSB; a Polonyi-like model was used for mSUGRA matching m_0 to $F_\phi/16\pi^2$. Q is the scale at which the masses are reported by ISAJET. Because it is not widely known, we remind the reader that the metric prefix E in the table means "exa" and is 10^{18}.

SUSYLR + AMSB	AMSB + m_0^2	mGMSB	mSUGRA
$\tan\beta = 15$	$\tan\beta = 15$	$\tan\beta = 15$	$\tan\beta = 15$
$\text{sgn}\,\mu = +1$	$\text{sgn}\,\mu = +1$	$\text{sgn}\,\mu = +1$	$\text{sgn}\,\mu = +1$
$Q = 550$ GeV	$Q = 558$ GeV	$Q = 899$ GeV	$Q = 537$ GeV
$F_\phi = 30$ TeV	$F_\phi = 30$ TeV	$\Lambda = 90$ TeV	$m_0 = 190$ GeV
	$m_0 = 290$ GeV	$M_{\text{mess}} = 180$ TeV	$m_{1/2} = 285$ GeV
$v_R = 135$ EeV			$A_0 = 241$ GeV
$f_1 = f_{c1} = 0.52$			
$f_3 = f_{c3} = 0.6$			

as light singlet fields mixed with Higgsinos. It could also be that the sneutrinos generated from late decay of the gravitino are dark matter. Several other possibilities for dark matter have also been discussed in [44], and we are currently investigating possible scenarios.

Finally, let us consider the sleptons masses—the main purpose of this paper. As advertised earlier, these are positive and depend on just a few parameters: F_ϕ, f_1, f_3 (since we have preserve parity at the high scale in this paper $f_{c1} = f_1$ and $f_{c3} = f_3$ at the right-handed scale) and to a lesser extent on $\tan\beta$ and the right-handed scale. The relative sizes of the masses are controlled by relative f coupling: the larger the coupling the larger the mass, e.g. increasing f_1 would raise the mass of the left-handed slepton. Such an affect contrasts strongly with other non-AMSB models with light doubly charged Higgses where the right-handed stau mass drops with increase in f_{c3} type coupling [45,46].

IV. CONCLUSION

We have presented a new way to solve the negative mass-squared slepton problem of AMSB using a minimal, bottom-up extension of the MSSM that incorporates neutrino masses (via the seesaw mechanism), solves the strong CP problem, and resolves the R-parity violation problem

of the MSSM. Slepton masses are rescued from the red by their couplings to both remnant doubly charged fields and left-handed triplets. Constraints from low energy physics and the nondecoupling of these additional fields require that their mass range be several TeV to F_ϕ. This in turn means the seesaw scale must be around 10^{11} GeV clearly distinguishing our model from GUT seesaw models.

The model we presented has soft terms which remain on their AMSB trajectory down to the SUSY scale. We have shown the sparticle spectrum for this model and compared it with typical predictions from other SUSY breaking scenarios finding significant deviations, especially in the gaugino sector.

ACKNOWLEDGMENTS

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