Supersymmetry Breaking by Type II Seesaw Assisted Anomaly Mediation

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Supersymmetry breaking by type II seesaw assisted anomaly mediation

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Anomaly mediated supersymmetry breaking, when implemented in the minimal supersymmetric standard model, is known to suffer from the problem of negative slepton mass squared leading to the breakdown of electric charge conservation. We show, however, that when the minimal supersymmetric standard model is extended to explain small neutrino masses by including a pair of superheavy Higgs triplet superfields (the type II seesaw mechanism), the slepton masses can be deflected from the pure anomaly mediated supersymmetry breaking trajectory and become positive. In the simple model we present in this paper, the seesaw scale is about 10^{13}–10^{14} GeV. Gauge coupling unification can be maintained by embedding the triplet to $SU(5)$ **15**-multiplet. In this scenario, the b-ino is the lightest supersymmetric particle and its mass is nearly degenerate with the next-to-lightest supersymmetric particle slepton when the triplet mass is right around the seesaw scale.

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I. INTRODUCTION

Supersymmetry (SUSY) is considered to be a prime candidate for TeV scale physics since it resolves several conceptual issues of the standard model (SM), such as (i) radiative stability of the large hierarchy between Planck and weak scales and (ii) electroweak symmetry breaking. With additional assumptions, it develops other appealing features: for instance, if R-parity symmetry is assumed, it can provide a candidate for the dark matter of the Universe, and if no new physics or specific new physics is assumed, it can lead to the unification of gauge couplings at a very high scale.

Since there is no trace of supersymmetry in current observations, it must be a broken symmetry, and the question arises as to the origin of this breaking. While at the phenomenological level it is sufficient to assume soft breaking terms to implement this, low energy observations in the domain of flavor changing neutral currents (FCNC) imply strong constraints on it; i.e. the sparticle masses must be flavor degenerate. It is therefore reasonable to require that any mechanism for SUSY breaking must lead to such flavor degeneracy for slepton and squark masses. Indeed, there exist at least two well-known scenarios where this happens: gauge mediated SUSY breaking (GMSB) [1,2] and anomaly mediated SUSY breaking (AMSB) [3,4]. In the simplest examples of both these cases, the FCNC effects are dynamically suppressed. Both involve unknown physics in the hidden sector which breaks supersymmetry, and this SUSY breaking information is transmitted to the visible sector via certain messengers. In the GMSB scenario, the messenger sector generically involves new particles and forces, whereas in the AMSB scenario, SUSY breaking is transmitted via the conformal breaking induced by radiative corrections in supersymmetric field theories. However, they differ in the way the SUSY breaking manifests in the low energy sector: in GMSB (as in the minimal gravity-mediated supersymmetry breaking models), the detailed pattern of sparticle masses depends on ultraviolet physics, i.e. physics at mass scales much higher than the SUSY breaking scale, whereas AMSB models have the advantage that this pattern depends only on the low scale physics. They are therefore easier to test experimentally given a particular low scale theory.

However, it turns out that AMSB models, despite their elegance and predictive power, suffer from a fatal problem when the low scale theory is assumed to be the minimal supersymmetric standard model (MSSM); i.e. they predict the slepton mass squared to be negative and hence lead to a vacuum state that breaks electric charge conservation (called the tachyonic slepton problem henceforth). This is, of course, unacceptable, and this problem needs to be solved if AMSB models have to be viable. There are many attempts to solve this problem by taking into account additional positive contributions to the slepton mass squared [5–8].

An important thing to realize at this point is that the MSSM is not a complete theory of low energy particle

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physics and needs an extension to explain the small neutrino masses observed in experiments. The relevant question then is whether the MSSM extended to include new physics that explains small neutrino masses will cure the tachyonic slepton mass pathology of AMSB.

There are two simple extensions of the MSSM which provide a natural explanation of small neutrino masses: the two types of seesaw mechanisms, i.e. type I [9] and type II [10]. In the first case, a reasonable procedure is to extend the gauge symmetry of the MSSM to $SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times SU(3)_c$, which automatically introduces three right-handed neutrinos into the theory as well as new couplings involving the leptons. One could image that these new couplings can affect the slepton masses. In most discussions of the seesaw mechanism, it is commonly assumed that the seesaw scale is very high ($\geq 10^{13}$ GeV or so); so one would expect the associated new physics interactions to decouple. Such a generic scenario will not solve the tachyonic slepton problem. However, it has recently been pointed out [11] that there exists a class of minimal SUSY left-right symmetric models with high scale seesaw mechanisms, where left-handed weak isotriplets with $\mathrm{B-L} = +2$ and doubly charged Higgs fields with $\mathrm{B-L} = +2$ coupling to right-handed leptons have a naturally weak scale mass because of higher symmetries of the superpotential. Their couplings to leptons contribute to the slepton mass squared and can solve the tachyonic slepton mass problem [11].

The present paper focuses on an alternative approach which uses the type II seesaw mechanism for neutrino masses, to see how it affects the slepton masses. An advantage of this over the type I approach is that it does not involve extending the gauge symmetry, but it requires adding a pair of $Y = \pm 1$ $SU(2)_L$ triplet Higgs fields to the MSSM. The $SU(2)_L$ triplets have mass close to 10^{13} GeV, which is required to implement the type II seesaw mechanism for small neutrino masses. We further assume that the triplet masses arise from the vacuum expectation value (VEV) of a light singlet field with a high VEV. We then show that in the AMSB scenario, the F component of the singlet field acquires an induced VEV, leading to new set of SUSY breaking effects. These effects are gauge mediated contributions to sparticle masses in addition to the usual AMSB contributions. We find that these contributions solve the tachyonic slepton mass problem. Thus, the type II seesaw mechanism, in addition to solving the neutrino mass problem, also solves the problem of SUSY breaking by AMSB.[1] Of course, in this case one needs to assume R-parity symmetry to obtain stable dark matter.

[1]We note that pure gauge mediation in the presence of the type II seesaw mechanism has been considered recently [12]; our model is different since AMSB effects play a significant role in the final predictions.

This scenario makes predictions for the sparticles which are different from other scenarios. In particular, we find that the b-ino and sleptons are nearly degenerate with the messenger at the seesaw scale—a situation which is particularly advantageous for understanding the dark matter abundance in the Universe [13]. We also show that the model does preserve the unification of couplings.

The paper is organized as follows. In Sec. II, we explain the scenario of "deflected anomaly mediation" which plays a crucial role in our solution to the tachyonic slepton problem. In Sec. III, we present a simple superpotential for the singlet field and calculate the deflection parameter. Section IV contains the general formulas of sparticle masses in the deflected anomaly mediation. In Sec. V, we present the minimal model to solve the tachyonic slepton problem as well as generate light neutrino masses. Section VI contains the extended models which preserve the gauge coupling unification. We summarize our results in Sec. VII. In the Appendix, we present the calculation of the lifetime of the SUSY breaking local minimum.

II. DEFLECTED ANOMALY MEDIATION AND MESSENGER SECTOR

It is well known that, in the absence of additional supersymmetry breaking, the AMSB contribution to sparticle masses is ultraviolet insensitive. It has, however, been proposed that the presence of additional SUSY breaking effects could deflect the sparticle masses from the AMSB trajectory and lead to new predictions for the sparticle spectrum. This has been called the "deflected anomaly mediation" scenario [5,7]. A key ingredient of this scenario is the presence of gauge mediated contributions arising from new interactions in the theory. Typically, they involve the introduction of messengers Ψ and $\bar{\Psi}$ with the following coupling:

$$W = S\bar{\Psi}\Psi. \tag{1}$$

Clearly, $\bar{\Psi}$ and Ψ are the messenger chiral superfields in a vectorlike representation under the SM gauge group, and S is the singlet superfield. It is crucial for the messenger fields to be nonsinglets, at least under the $SU(2)_L \times U(1)_Y$ gauge group. In our model, the $SU(2)_L$ triplets which enforce the type II seesaw mechanism will play the role of these fields.[2] Once the scalar component (S) and the F component (F_S) in the singlet chiral superfield develop VEVs, the scalar lepton obtains new contributions to its mass squared through the same manner as in the gauge mediation scenario [1,2]. In our case, F_S is induced by the hidden sector SUSY breaking conformal compensator. The effect of nonzero F_S is to deflect the sparticle masses from

[2]In order to implement the type II seesaw mechanism in the MSSM, we only need one pair of triplets, and it turns out that one pair of triplets is sufficient to lift slepton masses and leave the b-ino as the LSP.

the pure AMSB trajectory of the renormalization group equations, thereby solving the tachyonic slepton problem.

As just noted, an important difference between the deflected AMSB from GMSB is that the SUSY breaking in the messenger sector is induced by the anomaly mediation, namely, F_ϕ, a nonzero F component of the compensator field, and F_ϕ therefore is the unique source of SUSY breaking in this scenario. Therefore, we can parametrize the SUSY breaking order parameter in the messenger sector such as

$$\frac{F_S}{S} = dF_\phi. \tag{2}$$

Here, d is the so-called "deflection parameter" which characterizes how much the sparticle masses are deflected from the pure AMSB results. Theoretical consistency constrains it to be $|d| < \mathcal{O}(1)$, because F_S/S is not the original SUSY breaking sector.

We consider a simple model which provides a sizable deflection parameter $|d| = \mathcal{O}(1)$. Let us begin with the supergravity Lagrangian for S in the superconformal framework [14,15] (supposing SUSY breaking in the hidden sector and fine-tuning of the vanishing cosmological constant),

$$\mathcal{L} = \int d^4\theta \phi^\dagger \phi S^\dagger S + \left\{ \int d^2\theta \phi^3 W(S) + \text{H.c.} \right\}, \tag{3}$$

where we have assumed the canonical Kahler potential (in the superconformal framework), W is the superpotential [except for Eq. (1)], and $\phi = 1 + \theta^2 F_\phi$ is the compensating multiplet with the unique SUSY breaking source F_ϕ, taken to be real and positive through $U(1)_R$ phase rotation.

The scalar potential can be read off as

$$V = |F_S|^2 - S^\dagger S |F_\phi|^2 - 3F_\phi W - 3F_\phi^\dagger W^\dagger \tag{4}$$

with the auxiliary field given by

$$F_S = -(SF_\phi + W_S^\dagger), \tag{5}$$

where W_S stands for $\partial W/\partial S$.

Using the stationary condition $\partial V/\partial S = 0$ and Eq. (5), we can describe the deflection parameter in the simple form

$$\frac{F_S}{S} = dF_\phi = -2\frac{W_S}{SW_{SS}} F_\phi, \tag{6}$$

where W_{SS} stands for $\partial^2 W/\partial S^2$. This is a useful formula, from which we can understand that S should be light in the SUSY limit in order to obtain a sizable deflection parameter $|d| = \mathcal{O}(1)$, because the SUSY mass term (W_{SS}) appears in the denominator.

III. SINGLET SUPERPOTENTIAL AND DEFLECTION PARAMETER

As a simple model, let us consider a superpotential

$$W = -mS^2 + \frac{S^4}{M}, \tag{7}$$

where m and M are mass parameters, and we assume them to be real and positive, and $m \ll M$.[3] The scalar potential is given by

$$V = |S|^2 \left| -2m + 4\frac{S^2}{M} \right|^2 + F_\phi \left(mS^2 + \frac{S^4}{M} \right) + \text{H.c.} \tag{8}$$

Changing a variable as $S^2 = xe^{i\varphi}$ with real parameters, $x \geq 0$ and $0 \leq \varphi \leq 2\pi$, the scalar potential is rewritten as

$$V(x, \varphi) = 4x\left(m^2 - 4\frac{m}{M} x\cos(\varphi) + 4\frac{x^2}{M^2} \right) + 2F_\phi \left(mx\cos(\varphi) + \frac{x^2}{M}\cos(2\varphi) \right). \tag{9}$$

It is easy to check that $\varphi = 0$ satisfies the stationary condition $\partial V/\partial\varphi = 0$, and we take $\varphi = 0$. Solving the stationary condition $\partial V(x, \varphi = 0)/\partial x = 0$, we find

$$x_\pm = \frac{M}{24}(8m - F_\phi \pm \sqrt{D}), \tag{10}$$

where $D = 16m^2 - 40F_\phi m + F_\phi^2$. It is easy to show that x_+ and x_- correspond to the local minimum and maximum of the potential, respectively. For a fixed F_ϕ, the potential minimum exists if $D > 0$; in other words,

$$m > \frac{5 + 2\sqrt{6}}{4} F_\phi. \tag{11}$$

From Eq. (6), the deflection parameter is given by

$$d = \frac{-2m + 4x_+/M}{m - 6x_+/M} = \frac{2(4m + F_\phi - \sqrt{D})}{3(4m - F_\phi + \sqrt{D})}. \tag{12}$$

The deflection parameter reaches its maximum value ($d_{\max}$) in the limit $m \to \frac{5+2\sqrt{6}}{4}F_\phi$, and

$$d_{\max} = \frac{2(3 + \sqrt{6})}{3(2 + \sqrt{6})} \simeq 0.816. \tag{13}$$

Squared masses of two real scalar fields in $S = (x + iy)/\sqrt{2}$ are found to be

$$m_x^2 = 8\frac{\sqrt{D}x_+}{M}, \qquad m_y^2 = \frac{2}{3}(24mF_\phi + (2m - F_\phi)\sqrt{D} + D^2), \tag{14}$$

which are roughly of order m^2. Through a numerical

[3] We have checked that there are no large scalar S mass terms induced by loop corrections in the theory.

calculation, we find $m_x \simeq 0.24F_\phi$ and $m_y \simeq 6.3F_\phi$ for m very close to its minimum value leading to $d = 0.81$.

The scalar potential of Eq. (8), in fact, has a SUSY minimum at $S = 0$, where the potential energy is zero, and the minimum at x_+ we have discussed is a local minimum. In the Appendix, we estimate the decay rate of the local minimum to the true SUSY minimum and find it is sufficiently small for $F_\phi \ll M$.

IV. SPARTICLE MASS SPECTRUM

We first give general formulas for sparticle masses in the deflected anomaly mediation with the nonzero deflection parameter d. Following the method developed in Ref. [16] (see also Ref. [5]), we can extract the sparticle mass formulas from the renormalized gauge couplings $[\alpha_i(\mu, S)]$ and the supersymmetric wave function renormalization coefficients $[Z_I(\mu, S)]$ at the renormalization scale (μ) and the messenger scale (S). With $F_S/S = dF_\phi$, the gaugino masses (M_i) and sfermion masses $(\tilde{m}_I)$ are given by

$$\frac{M_i}{\alpha_i(\mu)} = \frac{F_\phi}{2}\left(\frac{\partial}{\partial \ln\mu} - d\frac{\partial}{\partial \ln|S|}\right)\alpha_i^{-1}(\mu, S),$$
$$\tilde{m}_I^2(\mu) = -\frac{|F_\phi|^2}{4}\left(\frac{\partial}{\partial \ln\mu} - d\frac{\partial}{\partial \ln|S|}\right)^2 \ln Z_I(\mu, S). \tag{15}$$

For a simple gauge group, the gauge coupling and the wave function renormalizations are given by

$$\alpha_i^{-1}(\mu, S) = \alpha_i^{-1}(\Lambda_{\rm cut}) + \frac{b_i - N_i}{4\pi}\ln\left(\frac{S^\dagger S}{\Lambda_{\rm cut}^2}\right) + \frac{b_i}{4\pi}\ln\left(\frac{\mu^2}{S^\dagger S}\right), \tag{16}$$

$$Z_I(\mu, S) = \sum_i Z_I(\Lambda_{\rm cut})\left(\frac{\alpha_i(\Lambda_{\rm cut})}{\alpha_i(S)}\right)^{2c_i/(b_i-N_i)}\left(\frac{\alpha_i(S)}{\alpha_i(\mu)}\right)^{2c_i/b_i}, \tag{17}$$

where $\Lambda_{\rm cut}$ is the ultraviolet cutoff, b_i are the beta function coefficients for different groups, c_i are the quadratic Casimirs, N_i are the Dynkin indices of the corresponding messenger fields [for example, $N_i = 1$ for a vectorlike pair of messengers of a fundamental representation under the $SU(N)$ gauge group], and the sum is taken corresponding to the representation of the sparticles under the SM gauge groups. Substituting them into Eq. (15), we obtain

$$M_i(\mu) = \frac{\alpha_i(\mu)}{4\pi}F_\phi(b_i + dN_i), \tag{18}$$

$$\tilde{m}_I^2(\mu) = \sum_i 2c_i\left(\frac{\alpha_i(\mu)}{4\pi}\right)^2|F_\phi|^2 b_i G_i(\mu, S), \tag{19}$$

where

$$G_i(\mu, S) = \left(\frac{N_i}{b_i}\xi_i^2 + \frac{N_i^2}{b_i^2}(1 - \xi_i^2)\right)d^2 + 2\frac{N_i}{b_i}d + 1 \tag{20}$$

with

$$\xi_i \equiv \frac{\alpha_i(S)}{\alpha_i(\mu)} = \left[1 + \frac{b_i}{4\pi}\alpha_i(\mu)\ln\left(\frac{S^\dagger S}{\mu^2}\right)\right]^{-1}. \tag{21}$$

In the limit $d \to 0$, the pure AMSB results are recovered and Eq. (19) leads to the negative mass squared for an asymptotically nonfree gauge theory ($b_i < 0$). This result causes the tachyonic slepton problem in the pure AMSB scenario.

After integrating the messengers out, the scalar mass squared at the messenger scale is given by (taking $\xi_i = 1$)

$$\tilde{m}_I^2(S) = \sum_i 2c_i\left(\frac{\alpha_i(S)}{4\pi}\right)^2|F_\phi|^2[N_i d^2 + 2N_i d + b_i], \tag{22}$$

where the first, the second, and the third terms in the brackets correspond to pure GMSB, mixed GMSB and AMSB, and pure AMSB contributions, respectively. The sign of the second term is proportional to d, so that the sign of the deflection parameter results in a different sparticle mass spectrum. The case $d < 0$ has been investigated in Ref. [5]; the resultant sparticle mass spectrum at the electroweak scale is very unusual, and colored sparticles tend to be lighter than color-singlet sparticles. On the other hand, the case $d > 0$ examined in Ref. [7] leads to a mass spectrum similar to the GMSB scenario. In the following, we consider the case $d > 0$ based on the simple model discussed in Sec. III.

V. MINIMAL MODEL

From the above discussion, it is clear that to solve the tachyonic slepton problem, we need messenger fields which are nonsinglet under $SU(2)_L \times U(1)_Y$. If we now look at the way to implement the type II seesaw formula for the small neutrino mass [10], we find that we need a pair of $SU(2)_L$ triplet fields, $\bar{\Delta}$: $(\mathbf{3}, -1)$ and Δ: $(\mathbf{3}, +1)$, which can play the dual role of generators of neutrino masses as well as messenger fields.

To see their role in the neutrino sector, we add to the MSSM superpotential the following couplings of the triplets to the lepton doublets (L_i) and the up-type Higgs doublet (H_u),

$$W_{\rm seesaw} = Y_{ij}L_i\Delta L_j + \lambda H_u\bar{\Delta}H_u, \tag{23}$$

where i, j denotes the generation index, and Y_{ij} is the Yukawa coupling. If they couple to the singlet field S discussed above as

$$W_{\rm mess} = S\,{\rm tr}[\bar{\Delta}\Delta], \tag{24}$$

then once $\langle S\rangle \neq 0$, it will give heavy mass to the triplets. Integrating out the heavy messengers with mass $M_{\rm mess} = \langle S\rangle$, this superpotential leads to the light neutrino mass

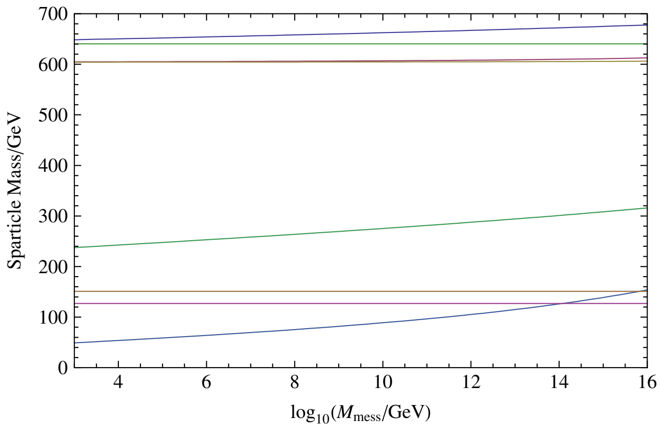

FIG. 1 (color online). Sparticle masses at $\mu = 500$ GeV as a function of the messenger scale in the type II seesaw model with one pair of $SU(2)_L$ triplet messengers. Here $d = 0.81$ and $F_\phi = 25$ TeV have been taken. Each line corresponds to the left-handed squark ($m_{\tilde{Q}}$), the gluino (M_3), the right-handed up squark ($m_{\tilde{u}^c}$), the right-handed down squark ($m_{\tilde{d}^c}$), the left-handed slepton ($m_{\tilde{L}}$), the W-ino (M_2), the b-ino ($|M_1|$), and the right-handed slepton ($m_{\tilde{e}^c}$) from above at $M_{\rm mess} = 10^3$ GeV. Two lines of $m_{\tilde{u}^c}$ and $m_{\tilde{d}^c}$ are overlapping and not distinguishable. For the messenger scale $M_{\rm mess} \gtrsim 10^{14}$ GeV, the b-ino becomes the lightest superparticle.

matrix $M_\nu \sim Y_{ij}\lambda\langle H_u\rangle^2/M_{\rm mess}$. This is the type II seesaw mechanism. If the messenger scale lies around the intermediate scale $M_{\rm mess} = 10^{13-14}$ GeV, the seesaw mechanism provides the correct scale for light neutrino masses with $Y_{ij}\lambda$ of order 1.

Note that since $F_S \neq 0$, the triplets can also serve as messenger superfields as in the usual GMSB models, and make additional contributions to slepton masses. In this minimal case, with a given d and the formulas in Eqs. (18)–(21), we now calculate the sparticle mass spectrum including the effects of AMSB and anomaly deflection. The beta function parameters needed for this purpose are $(b_1, b_2, b_3) = (-33/5, -1, +3)$, $(N_1, N_2, N_3) = (18/5, 4, 0)$. Neglecting the effects of Yukawa couplings,[4] the sparticle masses (in GeV) evaluated at $\mu = 500$ GeV are depicted in Fig. 1 as a function of the messenger scale $\log_{10}(M_{\rm mess}/{\rm GeV})$. Here, we have taken $d = 0.81$, $F_\phi = 25$ TeV, and the standard model gauge coupling constants at the Z pole as $\alpha_1(m_Z) = 0.0168$, $\alpha_2(m_Z) = 0.0335$, and $\alpha_3(m_Z) = 0.118$. Since the Higgs triplet pair does not carry color quantum number, the gluino mass stays on the AMSB trajectory and does not depend on the messenger scale as shown in Fig. 1. Note that for the messenger scale $M_{\rm mess} \gtrsim 10^{14}$ GeV, the b-ino becomes the lightest superparticle (LSP) and the b-ino-like neutralino would be a candidate for dark matter in our scenario [17]. For a small $\tan\beta$, annihilation processes of b-ino-like neutralinos are dominated by a p-wave, and since this annihilation process is not so efficient, the resultant relic density tends to exceed the upper bound on the observed dark matter density. This problem can be avoided, if the neutralino is quasidegenerate with the next LSP slepton, and the coannihilation process between the LSP neutralino and the next LSP slepton can lead to the right dark matter density. It is very interesting that our results show this degeneracy happening at $M_{\rm mess} \simeq 10^{14}$ GeV, which is, in fact, the correct seesaw scale.

In the simple superpotential of the singlet discussed in Sec. III, the messenger scale is given by $M_{\rm mess} = \langle S\rangle \sim \sqrt{F_\phi M}$. To obtain $M_{\rm mess} \sim 10^{13}$–$10^{14}$ GeV with $F_\phi = \mathcal{O}(10)$ TeV, we can specify the superpotential in Eq. (7) as

$$W \sim -mS^2 + \eta\frac{S^4}{M_{\rm Pl}} \qquad (25)$$

with $\eta \sim 10^{-3}$–10^{-5}, where $M_{\rm Pl}$ is the Planck scale.

VI. MINIMAL MODEL WITH GRAND UNIFICATION

The messengers we have introduced in the minimal model are $SU(3)_c$ singlets, and the existence of such particles below the grand unification scale $M_{\rm GUT} \sim 10^{16}$ GeV spoils the successful gauge coupling unification in the MSSM. As is well known, the gauge coupling unification can be kept if the messenger fields introduced are in the $SU(5)$ grand unified theory multiplets. There are two possibilities for such messengers that play two different roles in the neutrino sector by the seesaw mechanism. One is to introduce the messengers of $\mathbf{15} + \overline{\mathbf{15}}$ multiplets under $SU(5)$, which include Δ and $\bar{\Delta}$ as submultiplets. The other possibility is to introduce **24** multiplets [18].

Let us first consider the **15** and $\overline{\mathbf{15}}$ case in the $SU(5)$ grand unified theory model. We introduce the superpotentials

$$W_{\rm mess} = S\bar{T}T, \qquad W_{\rm seesaw} = Y_{ij}\bar{\mathbf{5}}_i\bar{\mathbf{5}}_jT + \lambda\mathbf{5}_H\mathbf{5}_H\bar{T}, \quad (26)$$

where T and $\bar{T}$ are **15** and $\overline{\mathbf{15}}$ multiplets. After integrating the heavy messengers out, we obtain the light neutrino mass matrix as $M_\nu \sim \langle\mathbf{5}_H\rangle^2/\langle S\rangle$ through the type II seesaw mechanism.

Sparticle masses can be evaluated in the same manner as before, but in this case, $N_1 = N_2 = N_3 = 7$. The resultant sparticle masses at $\mu = 500$ GeV are depicted in Fig. 2 as a function of the messenger scale $\log_{10}[S/{\rm GeV}]$. Here, we have taken $d = 0.48$ and $F_\phi = 25$ TeV. The b-ino becomes the LSP, degenerating with right-handed sleptons for the messenger scale $M_{\rm mess} \sim 10^{13}$ GeV.

[4]In general, there are Yukawa mediation contributions to the $SU(2)_L$ doublet slepton mass due to the coupling $Y_{ij}L_iL_j\Delta$. In this paper, we consider the case in which $Y_{ij} \leq 0.1$ by adjusting the seesaw scale and also the parameter λ, so that the Yukawa mediation contributions are negligible.

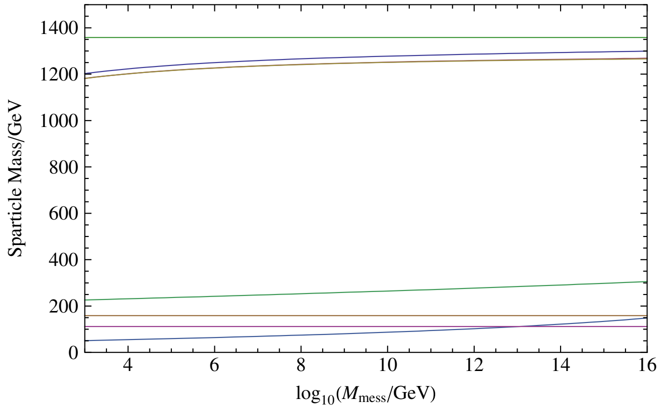

FIG. 2 (color online). Sparticle masses at $\mu = 500$ GeV as a function of the messenger scale in the type II seesaw model with one pair of $\overline{\mathbf{15}} + \mathbf{15}$ messengers. Here $d = 0.48$ and $F_\phi =$ 25 TeV have been taken. Each line corresponds to M_3, $m_{\tilde{Q}}$, $m_{\tilde{u}^c}$, $m_{\tilde{d}^c}$, $m_{\tilde{L}}$, M_2, $|M_1|$, and $m_{\tilde{e}^c}$ from above at $M_{\rm mess} =$ 10^3 GeV. Two lines of $m_{\tilde{u}^c}$ and $m_{\tilde{d}^c}$ are overlapping and not distinguishable. For the messenger scale $M_{\rm mess} \gtrsim 10^{13}$ GeV, the b-ino becomes the LSP.

In the case of **24** multiplets (Σ), the relevant superpotential is given by

$$W_{\rm mess} = S\,{\rm tr}[\Sigma^2], \qquad W_{\rm seesaw} = Y_i \bar{\mathbf{5}}_i \Sigma \mathbf{5}_H. \qquad (27)$$

After integrating out the heavy **24**, the light neutrino mass matrix is given by $M_\nu \sim Y_i Y_j \langle \mathbf{5}_H \rangle^2 / \langle S \rangle$. Note that the rank of this matrix is 1. We need to introduce at least two **24** messengers to incorporate the realistic neutrino mass matrix. As an example, we consider two **24** messengers with the same masses. We evaluate sparticle masses

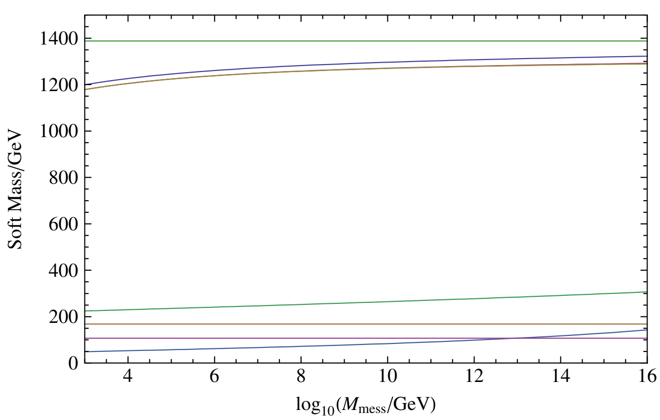

FIG. 3 (color online). Sparticle masses at $\mu = 500$ GeV as a function of the messenger scale in the model with two pairs of **24** messengers. Here $d = 0.35$ and $F_\phi = 25$ TeV have been taken. Each line corresponds to M_3, $m_{\tilde{Q}}$, $m_{\tilde{u}^c}$, $m_{\tilde{d}^c}$, $m_{\tilde{L}}$, M_2, $|M_1|$, and $m_{\tilde{e}^c}$ from above at $M_{\rm mess} = 10^3$ GeV. Two lines of $m_{\tilde{u}^c}$ and $m_{\tilde{d}^c}$ are overlapping and not distinguishable. For the messenger scale $M_{\rm mess} \gtrsim 10^{13}$ GeV, the b-ino becomes the LSP.

with $N_1 = N_2 = N_3 = 2 \times 5 = 10$ in this case. The resultant sparticle masses at $\mu = 500$ GeV are depicted in Fig. 3 as a function of the messenger scale $\log_{10}[S/{\rm GeV}]$. Here, we have taken $d = 0.35$ and $F_\phi = 25$ TeV. The b-ino becomes the LSP, degenerate with right-handed sleptons for the messenger scale $M_{\rm mess} \sim 10^{13}$ GeV.

VII. CONCLUSION

In conclusion, we have pointed out that a minimal extension of the MSSM needed to explain small neutrino masses via the seesaw mechanism can also cure the tachyonic slepton mass problem of anomaly mediated supersymmetry breaking. We have presented the sparticle spectrum for these models and shown that they can preserve the unification of gauge couplings. We find it interesting that the same mechanism that explains the smallness of neutrino masses also cures the tachyonic slepton problem of AMSB.

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APPENDIX: LIFETIME OF THE LOCAL MINIMUM

The scalar potential in Sec. III, $V(S)$, has the global SUSY minimum at the origin, and the minimum we have discussed is a local minimum. If our world is trapped in the local minimum, it will eventually decay into the SUSY minimum. The lifetime of the local minimum should be sufficiently long, at least longer than the age of the Universe, $\tau_U \sim 4.3 \times 10^{17}$ s for our model to be viable. Here we estimate the decay rate of the false vacuum within the parameters of our model.

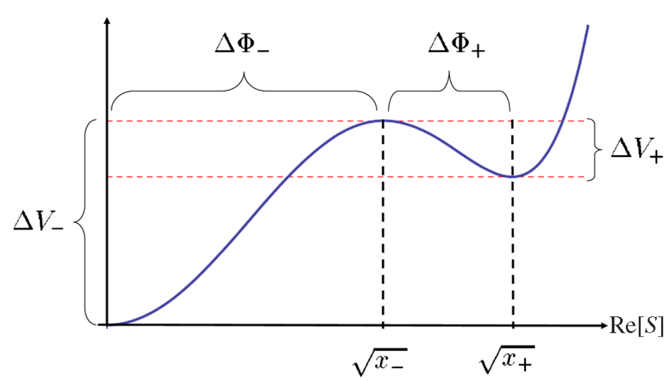

FIG. 4 (color online). Schematic picture of the scalar potential $V(S)$ as a function of the real part of S.

In our calculation, the scalar potential is treated in the triangle approximation [19]. A schematic picture of the scalar potential is depicted in Fig. 4. Let us take the path in the direction of Re[S]: climbing up from the local minimum at Re[S] $= \sqrt{x_+}$ to the local maximum at Re[S] $= \sqrt{x_-}$, then rolling down to the SUSY minimum at $S = 0$. In the triangle approximation, parameters characterizing the potential are

$$\Delta V_\pm, \qquad \Delta\Phi_\pm, \tag{A1}$$

where $\Delta V_+ (\Delta V_-)$ is the height difference of the potential between the barrier and the local (global) minima, and $\Delta\Phi_+ (\Delta\Phi_-)$ is the width difference of the potential between the barrier and the local (global) minima. Following Ref. [19], we define

$$c \equiv \frac{\Delta V_- \Delta\Phi_+}{\Delta V_+ \Delta\Phi_-} \tag{A2}$$

and the decay rate per unit volume is estimated as $\Gamma/V \sim e^{-B}$ with

$$B = \frac{32\pi^2}{3} \frac{1+c}{(\sqrt{1+c}-1)^4} \frac{\Delta\Phi_+^4}{\Delta V_+}. \tag{A3}$$

The consistency condition to apply the triangle approximation is given by [19]

$$\left(\frac{\Delta V_-}{\Delta V_+}\right)^{1/2} \geq \frac{2\Delta\Phi_-}{\Delta\Phi_- - \Delta\Phi_+}. \tag{A4}$$

For the scalar potential analyzed in Sec. III,

$$\Delta\Phi_+ = \sqrt{x_+} - \sqrt{x_-}, \Delta\Phi_- = \sqrt{x_-},$$
$$\Delta V_+ = V(x_-, 0) - V(x_+, 0), \qquad \Delta V_- = V(x_-, 0). \tag{A5}$$

In order to get the deflection parameter as large as possible, let us consider the case where the local minimum and maximum points are very close, namely, $\Delta\Phi_+$ and ΔV_+ are very small. In this case, the condition Eq. (A4) is satisfied, and we can apply the triangle approximation. With a small parameter $0 < \epsilon \ll 1$, we parametrize

$$m = \frac{5+2\sqrt{6}}{4} F_\phi (1+\epsilon). \tag{A6}$$

In the limit $\epsilon \to 0$, the local minimum and maximum collide, and the local minimum disappears. The deflection parameter is approximately described as

$$d \simeq d_{\max} - \frac{\sqrt{12+5\sqrt{6}}}{3} \epsilon^{1/2} \simeq d_{\max} - 1.64\epsilon^{1/2}. \tag{A7}$$

The straightforward calculations give the following results:

$$\begin{aligned}
\Delta\Phi_+ &\simeq \sqrt{\frac{12+5\sqrt{6}}{54+24\sqrt{6}}} \sqrt{F_\phi M} \epsilon^{1/2}, \\
\Delta\Phi_- &\simeq \frac{1}{2}\sqrt{\frac{9+4\sqrt{6}}{6}} \sqrt{F_\phi M}, \\
\Delta V_+ &= \frac{(12+5\sqrt{6})^{3/2}}{27} F_\phi^3 M \epsilon^{3/2}, \\
\Delta V_- &= \frac{1107+452\sqrt{6}}{288} F_\phi^3 M.
\end{aligned} \tag{A8}$$

Also, we find

$$B \simeq \frac{\pi^2 128(12+5\sqrt{6})^{3/2}}{9(6937+2832\sqrt{6})} \frac{M}{F_\phi} \epsilon^{3/2} \simeq 1.21 \times \frac{M}{F_\phi} \epsilon^{3/2}. \tag{A9}$$

Recalling that the messenger scale is roughly given by $M_{\rm mess} \sim \sqrt{F_\phi M}$ and $F_\phi \simeq 10$ TeV to obtain sparticle masses around 100 GeV–1 TeV, we can rewrite B as

$$\begin{aligned}
B &\simeq 1.21 \left(\frac{M_{\rm mess}}{F_\phi}\right)^2 \epsilon^{3/2} \\
&= 1.21 \times 10^{20} \left(\frac{M_{\rm mess}/10^{14}\ {\rm GeV}}{F_\phi/10\ {\rm TeV}}\right)^2 \epsilon^{3/2}.
\end{aligned} \tag{A10}$$

For the parameters chosen as in Fig. 1, $M_{\rm mess} \simeq 10^{14}$ GeV, $F_\phi = 25$ TeV, $d = 0.81$ and correspondingly $\epsilon \simeq 1.57 \times 10^{-5}$, we find $B \simeq 1.20 \times 10^{12}$. The lifetime of the local minimum is extremely long.

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