A Viable Supersymmetric Model
with UV Insensitive Anomaly Mediation

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Abstract

We propose an electroweak model which is compatible with the UV insensitive anomaly mediated supersymmetry breaking. The model is an extension of the NMSSM by adding vector-like matter fields which can drive the soft scalar masses of the singlet Higgs field negative and the successful electroweak symmetry breaking is achieved. Viable parameter regions are found to preserve perturbativity of all the coupling constants up to the Planck scale. With this success, the model becomes a perfect candidate of physics beyond the standard model without the FCNC and CP problem. The cosmology is also quite interesting. The lightest neutralino is the wino which is a perfect cold dark matter candidate assuming the non-thermal production from the gravitino decay. There is no gravitino problem because it decays before the BBN era, and thus the thermal leptogenesis works. The cosmological domain wall problem inherent in the NMSSM is absent since the Z_3 symmetry is broken by the QCD instanton effect in the presence of the vector-like quarks. We also briefly comment on a possible solution to the strong CP problem *à la* the Nelson-Barr mechanism.

1 Introduction

Supersymmetric (SUSY) standard model (SM) is an attractive framework which realizes the standard model at low energy and suggests unification of gauge interactions at high energy. The SUSY breaking triggers the electroweak symmetry breaking while stabilizing the Higgs potential against the radiative corrections.

Despite the great success in the gauge and Higgs sectors, the matter sector is problematic. With generic SUSY breaking terms in the Lagrangian, the prediction of small CP violation and Flavor Changing Neutral Current (FCNC) in the SM may be destroyed by the new interactions among the fermions and their scalar partners. The problems suggest the special features of the SUSY breaking, which are flavor blind and CP conserving.

The gravitino, the SUSY partner of the graviton, also causes a problem in cosmology. In the gravity mediated SUSY breaking scenario, the gravitino mass is of the order of TeV and it decays during big-bang nucleosynthesis (BBN), destroying the successful prediction of the abundance of the light elements. In order to avoid the problem, the reheating temperature of the universe has to be lower than 10^6 GeV [1], which is too low for many baryogenesis scenarios with out-of-equilibrium decays of heavy particles.

An interesting possibility which solves all those problems automatically is the anomaly mediated SUSY breaking scenario [2, 3]. The anomaly mediation is the purely gravitational mediation mechanism of the SUSY breaking and is realized once the direct couplings between the hidden and visible sector fields are suppressed. The SUSY breaking only appears with the conformal anomaly of the theory and therefore it is insensitive to the ultraviolet (UV) physics and respect all the accidental symmetries in low energy such as CP and flavor conservation in the SM. Moreover, the soft SUSY breaking terms are suppressed by the loop factor $1/(4\pi)^2$ compared to the gravitino mass because of its quantum origin, and therefore the gravitino mass is naturally of $O(100 \text{ TeV})$. Such a heavy gravitino decays before the BBN era, opening a window of baryogenesis at high temperatures [4].

The UV insensitivity enables us to calculate all the SUSY breaking terms with known coupling constants. Unfortunately, the high predictability immediately excludes the pure anomaly mediation since the scalar leptons turn out to be tachyonic. There have been several proposals to cure the problem. For example, to go off the trajectory of the anomaly mediation, one considers non-decoupling effects by using flat directions [5, 6, 7, 8] or low energy thresholds [9]. The introduction of SM non-singlet particles which feel SUSY breaking directly from the hidden sector is shown to modify the spectrum through two-loop diagrams [10]. Also, attempts to modify the trajectory of the anomaly mediation itself by adding new interactions have been made [11, 12, 13, 14]. As modification of the initial conditions, several kinds of scenario have been proposed such as inclusion of the Kähler anomaly [15], adding D-terms [16, 17, 18], and adding the boundary interactions in the extra-dimensional setup [19]. The admixture of the gauge mediation and the anomaly mediation is recently considered in the context of the conformal sequestering scenario [20]. Among those various modifications, adding D-terms is clearly the safest since it has no danger of reintroductions of the CP or FCNC problems, and moreover it is shown to preserve the UV insensitivity which ensures the flavor blind and CP conserving soft terms in low energy even in the presence of the flavor changing or CP violating interactions in

high energy [16, 17].

Once we realize the UV insensitive anomaly mediation by the D-term modification, the next step is to consider the electroweak symmetry breaking to see if it is possible to have the desired vacuum [18, 21]. When all the SUSY breaking terms are calculable, it is highly non-trivial to have the correct vacuum expectation values (VEV) of the Higgs fields. In Ref.[21], it has been examined and found that the minimal SUSY standard model (MSSM) does not have a stable vacuum unless the $\tan\beta$ parameter, the ratio of two VEVs of the Higgs fields ($\langle H_2^0\rangle/\langle H_1^0\rangle$), is less than unity. In this case, the top-Yukawa coupling constant is very large and blows up right above the stop mass scale. Also, in the next to MSSM (NMSSM) the prediction to the Higgsino mass is too small in all the region of the parameter space. It is caused by the fact that the singlet Higgs field S only has Yukawa interactions which are asymptotically non-free. The soft mass squared is likely to be positive in that case, resulting in the small VEV of S. A successful model is found with linear term of the singlet field in the superpotential, which happens to be the low energy effective theory of the minimal SUSY fat Higgs model [22].

In this paper, we reexamine the NMSSM with D-term modified anomaly mediation by introducing additional vector-like quarks [23, 24, 25, 11]. We find that the coupling between the singlet Higgs field S and the vector-like quarks can modify the anomalous dimension of S significantly, and make the soft mass squared of S small enough to acquire a large VEV and thus there is no Higgsino mass problem.

This paper is organized as follows: In the next section, we review the UV insensitive anomaly mediation. We discuss the problems in the MSSM and NMSSM with the UV insensitive anomaly mediation in Section 3. We extend the NMSSM with additional vector-like matter fields and examine the electroweak symmetry breaking in Section 4, and discuss phenomenological and cosmological issues. Section 5 is devoted to conclusions.

2   UV insensitive anomaly mediation

Anomaly mediation is realized once we obtain a sequestered Kähler potential in the supergravity action. Since the fields in the visible sector can feel the SUSY breaking only through the gravitational interaction, the only source of the soft terms is the F-component of the gravity multiplet. In the superconformal formulation of the supergravity Lagrangian [26, 27, 28], the F-component is the auxiliary component of the chiral compensator multiplet Φ with which Lagrangian possesses superconformal symmetry. The supergravity Lagrangian is obtained by fixing the value of the components of Φ which breaks the superconformal symmetry explicitly down to the super-Poincaré symmetry. It is clear in this construction that the SUSY breaking appears only with the violation of the conformal symmetry. At the classical level, therefore, the soft scalar masses, the scalar cubic couplings (A-terms), and the gaugino masses vanish, and all those terms appear at quantum level with conformal anomaly. Explicitly, with the anomalous dimension γ_i of the chiral superfield Q_i and the beta function β_A of the gauge interaction labeled by A, soft terms at the scale μ are given by

$$A_{ijk} = -\lambda_{ijk}(\gamma_i + \gamma_j + \gamma_k)m_{3/2} \ , \quad \tilde{m}_i^2 = \frac{1}{2}\frac{d\gamma_i}{d\ln\mu}m_{3/2}^2 \ , \quad m_\lambda^A = \frac{\beta_A}{g_A}m_{3/2} \ , \tag{1}$$

where λ_{ijk} and g_A are the Yukawa and gauge coupling constants, respectively. The mass parameter $m_{3/2}$ is the gravitino mass which represents the F-component of the compensator multiplet. The above soft terms are defined by

$$\mathcal{L}_{\text{soft}} = -(A_{ijk}q_iq_jq_k + h.c.) - \tilde{m}_i^2|q_i|^2 - \frac{1}{2}m_\lambda^A\bar{\lambda}\lambda \;, \tag{2}$$

with q_i and λ being the scalar component of Q_i and gauginos, respectively. Of interest is that the soft terms at the scale μ are described by γ and β at that scale and hence do not depend on the high energy physics. The UV insensitivity is a phenomenologically desirable feature since it solves the SUSY FCNC and CP problem automatically.

It had been thought that sequestered Kähler potential is difficult to achieve in realistic model of quantum gravity. For instance, string theory tends to give rise to many moduli fields who can mediate additional supersymmetry breaking effects at the tree-level or one-loop level, which dominate over the anomaly-mediated contributions [29, 30]. In Ref. [3], the absence of such fields was explicitly assumed.* The physical separation of hidden and observable sectors along an extra dimension was used in Ref. [2] to justify the sequestered form of the Kähler potential. This, however, is not immune to the problem because there may be light bulk scalars such as radion. Only recently, a concrete mechanism to fix all of the moduli fields was proposed [33]. Moreover, the physical separation was shown not necessary to achieve sequestered Kähler potential if the hidden sector is nearly conformal [34, 35], even though the hidden sector needs to be of a special type [36]. Therefore achieving sequestered form of the Kähler potential does not appear to be an insurmountable problem any more.

An obvious problem of the framework is the tachyonic sleptons. The contribution from the gauge interaction to the scalar masses $\tilde{m}_i^2$ is positive (negative) for asymptotically free (non-free) gauge interaction. Since the sleptons only have $SU(2)_L$ and $U(1)_Y$ gauge interactions which are both asymptotically non-free, the scalar masses squared of the sleptons turn out to be negative with neglecting Yukawa coupling constants of the leptons.

However, we can easily solve the problem when we gauge the $B - L$ symmetry in the MSSM [17]. For example, if a $U(1)_A$ gauge symmetry in the hidden sector acquire the D-term, the kinetic mixing term between $U(1)_A$ and $U(1)_{B-L}$ induces the D-term of $U(1)_{B-L}$. The D-term of $U(1)_Y$ may also be generated in the same way. Remarkably those two D-terms can provide sufficiently large positive contributions for the scalar masses squared of both the left- and right-handed sleptons. Moreover, the most significant feature is that the modification with D-terms preserves the UV insensitivity of the soft terms [17].

The introduction of D-terms modifies the soft terms in the following way.

$$\tilde{m}_i^2 = \frac{1}{2}\frac{d\gamma_i}{d\ln\mu}m_{3/2}^2 - Q_Y^i D_Y - Q_{B-L}^i D_{B-L} \;, \tag{3}$$

where Q_Y^i and Q_{B-L}^i are the hypercharge and the $B - L$ charge of the corresponding super-field. The gaugino masses and the A-terms are not modified. Neglecting the Yukawa coupling

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*The paper discussed another possibility that the gaugino mass is generated at one-loop level by the anomaly-mediation effect whereas the scalar masses are not suppressed, realizing the split SUSY scenario [31, 32].

constants, we obtain the slepton masses as follows:

$$m_{\tilde{l}}^2 = \left(-\frac{11}{2} g_Y^4 - \frac{3}{2} g_2^4 \right) M^2 + \frac{1}{2} D_Y + D_{B-L} \ , \tag{4}$$

$$m_{\tilde{e}^c}^2 = -22 g_Y^4 M^2 - D_Y - D_{B-L} \ , \tag{5}$$

where g_Y and g_2 are the U(1)$_Y$ and SU(2)$_L$ gauge coupling constants, respectively, and $M = m_{3/2}/(4\pi)^2$. The additional D-term contributions are positive for both of the sleptons when $D_Y < -D_{B-L} < D_Y/2 < 0$.

The D-terms can be obtained in a consistent way with grand unified theories (GUT), even though U(1)$_Y$ D-term becomes gauge non-singlet. For example, in SO(10) grand unified theories, U(1)$_Y$ and U(1)$_{B-L}$ are both subgroups of SO(10). The kinetic mixing terms between the U(1)$_A$ gauge field and those U(1)'s in this case are generated after the GUT breaking assuming the presence of the following term:

$$\mathcal{L} \ni \int d^2\theta \frac{\Sigma^K}{M_{\mathrm{Pl}}} (W_{U(1)_A})^\alpha (W_{\mathrm{SO}(10)}^K)_\alpha + h.c. \ . \tag{6}$$

Here, Σ^K is a chiral superfield of 45 dimensional representation, which breaks SO(10) into SU(3)$_C$ × SU(2)$_L$ × U(1)$_Y$ × U(1)$_{B-L}$. Provided the D-term of U(1)$_A$ is generated at one-loop level, the above mixing induces the D-terms of U(1)$_Y$ and U(1)$_{B-L}$ of the order of $(4\pi)^2 (M_{\mathrm{GUT}}/M_{\mathrm{Pl}})M^2$ which is numerically similar size to the sfermion soft masses squared of $O(M^2)$. Alternatively, in the models with orbifold unification [37, 38, 39], we can simply write down a kinetic mixing term between U(1)$_Y$ and U(1)$_{B-L}$ gauge fields at the boundary.

3 Electroweak symmetry breaking in the MSSM and NMSSM

We now proceed to consider the electroweak symmetry breaking with the D-term modified anomaly mediation. Whether we obtain correct vacuum is quite non-trivial since all the soft terms are calculable with known coupling constants.

3.1 MSSM

In Ref.[21] it has been shown that the correct vacuum (correct Z boson mass) is realized only when $\tan\beta \lesssim 0.3$ which causes the Landau pole of the top-Yukawa coupling constant just above the SUSY breaking scale. The problem is caused by the explicit violation of the conformal symmetry by the μ-term, i.e., $W \ni \mu H_1 H_2$. Because it is a tree-level violation, the B-term associated with the μ-term is not suppressed by the loop factor and hence is large by a factor of $(4\pi)^2$ compared to other soft terms. Such a huge B-term requires large Yukawa couplings to fine-tune the VEV of the Higgs fields by enhancing soft masses of H_1 and H_2 through Eq.(1).
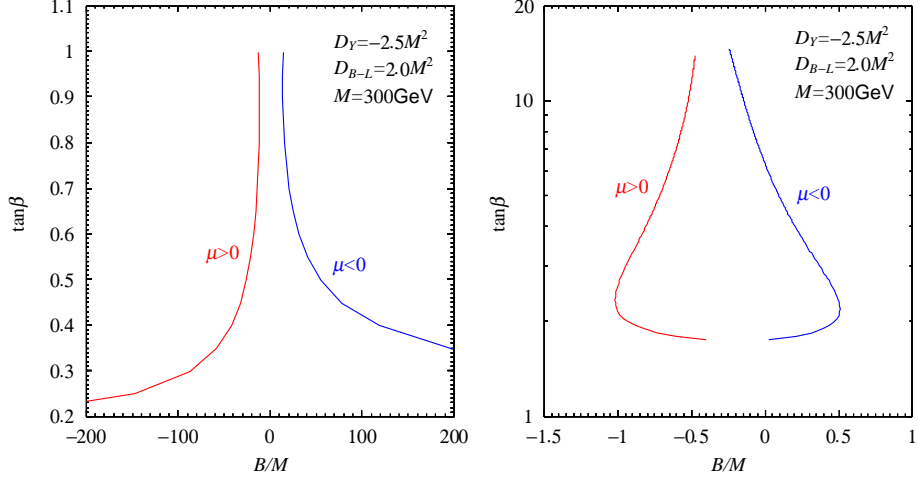
Figure 1: The relation between B-parameter and $\tan\beta$ is shown. We fixed the gravitino mass parameter ($M = m_{3/2}/(4\pi)^2$) to be 300GeV and D-terms ($D_Y = -2.5M^2$ and $D_{B-L} = 2.0M^2$).

There is a way of suppressing the B-term by using an accidental cancellation among different sources of the μ- and B-terms. For example, if we have the following terms in the Kähler and the superpotential,

$$K \ni cH_1H_2 + h.c. \ , \quad W \ni \tilde{\mu}H_1H_2 \ , \tag{7}$$

the μ- and B-terms from the SUSY breaking are given by

$$\mu = cm_{3/2} + \tilde{\mu} \ , \quad B\mu = m_{3/2}(-cm_{3/2} + \tilde{\mu}) \ . \tag{8}$$

Since the relative sign between the two contributions is different for the μ- and B-terms, there is a possibility to have small B-term by carefully tuning the c and $\tilde{\mu}$ parameters. In this situation, we can think of the B-parameter as a free parameter. We plot in Fig.1 the B-parameter dependence of $\tan\beta$ with fixing M, D_Y, and D_{B-L}. We have included the one-loop correction to the Higgs potential from the (s)top, (s)bottom, and (s)tau loop diagrams, and imposed the stability of all the scalar particles. For large $|B|/M$, such as $(4\pi)^2 \sim 158$, we need a small value of $\tan\beta$ so that the top-Yukawa coupling constant gives large contribution to the Higgs potential. The solutions with $\tan\beta \gtrsim 1$ are found with small values of B. Therefore, even in the MSSM, we could achieve the successful electroweak symmetry breaking with maintaining perturbativity up to the Planck/GUT scale at the expense of a fine-tuning of order 10^{-2} in Eq.(8).

Unfortunately, the model has a potential danger to introduce a new CP phase since c and $\tilde{\mu}$ are independent complex parameters. Although the reintroduction of the SUSY CP problem upsets the motivation of the UV insensitive anomaly mediation, there is an interesting observation. As we discuss later in Sec. 4.5, the anomaly mediation is a good framework for solving the strong CP problem by the scenario of the spontaneous CP violation [40, 41]. Since the scenario assumes the exact CP invariance in the Lagrangian, c and $\tilde{\mu}$ are real parameters

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if the fields which break the CP invariance do not couple with the Higgs fields. In this sense, we claim that the MSSM is a viable model with the UV insensitive anomaly mediation once we obtain small B-parameters.

3.2 NMSSM

Apart from the benefit of providing a solution to the μ-problem, the extension to the NMSSM is well-motivated in the context since the conformal symmetry is not violated at tree level, and hence there is no complication caused by the $(4\pi)^2$ enhanced soft parameters. The superpotential of the Higgs sector in the NMSSM is given by

$$W = \lambda S H_1 H_2 + \frac{h}{3} S^3 \ , \tag{9}$$

where S is the gauge singlet Higgs field whose VEV plays a role of the μ-parameter in the MSSM.

The electroweak symmetry breaking in the NMSSM is, however, not successful since the effective μ-parameter of $\lambda\langle S\rangle$ is too small (at most a few GeV), resulting in unacceptably light Higgsinos [21]. The problem is caused by the positivity of the soft mass squared for the singlet Higgs m_S^2. The m_S^2 parameter is given by the formula:

$$m_S^2 = \frac{1}{2}(4\pi)^2 \frac{d}{d\ln\mu}\left(2\lambda^2 + 2h^2\right) M^2 \ . \tag{10}$$

Ignoring the $SU(2)_L$ and $U(1)_Y$ gauge interactions, λ and h are asymptotically non-free, and thus m_S^2 is positive. The positivity indicates the stability of the origin of the potential, and therefore the VEV of S is only induced by the small shift of the origin through the linear term $\lambda A_\lambda\langle H_1^0\rangle\langle H_2^0\rangle S$ in the potential. Although the linear term can be enhanced by increasing the coupling λ and/or h, it is of no help since that also enhances m_S^2.

4 Modified NMSSM with Vector-like Matter Fields

4.1 The Model

Having understood the problems in the MSSM and the NMSSM, we now consider a model with vector-like matter fields. The problem in the NMSSM is caused by the high predictability of the anomaly mediation. The soft terms are calculable once we fix a model. Particularly the large positive m_S^2 parameter is problematic if we assume the minimal interactions. However, the interactions of S are phenomenologically unknown and we can easily modify the soft terms of the singlet. As an example, we consider a model with vector-like matter fields which couple to the S field.

A similar model has been considered in Ref.[11] in the context of the anomaly mediation without D-terms. It was pointed out that the calculability of the soft terms still makes the electroweak symmetry breaking difficult even in the presence of S and extra vector-like fields by the following reason. The $\tan\beta$ parameter is determined by the difference between two

soft masses of the Higgs doublets, $m_{H_1}^2$ and $m_{H_2}^2$, which only depends on the top Yukawa coupling constant f_t in that model. Therefore, the top-quark mass $m_t(= f_t v \sin \beta)$ is predicted rather than an input parameter. Unfortunately, successful electroweak symmetry breaking turns out to require unacceptably small m_t such as less than 145 GeV. They discussed a further extension of the model to overcome the situation by introducing three extra singlet superfields and appropriate interaction terms. In contrast, we show that the correct vacua are easily found in the model with D-terms. Since D_Y contributes to $m_{H_1}^2$ and $m_{H_2}^2$ with opposite signs, we do not have the unwanted correlation between $\tan \beta$ and f_t anymore. We can find phenomenologically viable parameter regions without introducing extra singlets except for S.

We introduce a pair of new chiral superfields D and $\bar{D}$ which have quantum numbers of $(\mathbf{3}, \mathbf{1})_{-1/3}$ and $(\bar{\mathbf{3}}, \mathbf{1})_{1/3}$ under the SM gauge group, and also L and $\bar{L}$ of $(\mathbf{1}, \mathbf{2})_{-1/2}$ and $(\mathbf{1}, \mathbf{2})_{1/2}$ so that the extra matter fields form complete SU(5) representations $\mathbf{5} + \bar{\mathbf{5}}$ to ensure the gauge coupling unification. We can write down the following superpotential:

$$W = \lambda S H_1 H_2 + \frac{h}{3} S^3 + k_D S \bar{D} D + k_L S \bar{L} L . \tag{11}$$

The introduction of the interactions between S and vector-like matter fields does not cause a new CP violation, since the coupling constants λ, h, k_D, and k_L can be made real without loss of generality by appropriate field redefinitions. We do not assume the presence of the direct couplings of D and L with the ordinary quarks and leptons, which may be forbidden by a U(1) symmetry or its discrete subgroup under which only the vector-like matter fields transform. The additional interactions in Eq.(11) modify the m_S^2 parameter from Eq.(10) to (see Appendix for the RGEs)

$$m_S^2 = \frac{1}{2} (4\pi)^2 \frac{d}{d \ln \mu} \left(2\lambda^2 + 2h^2 + 3k_D^2 + 2k_L^2 \right) M^2 . \tag{12}$$

Since k_D is asymptotically free in a wide range of parameter space due to the strong interaction of D and $\bar{D}$, the negative contribution may make m_S^2 negative and large in its absolute value such that we obtain large values of $\langle S \rangle$.

4.2   Electroweak symmetry breaking

Let us evaluate the VEV of the Higgs fields by minimizing the potential with soft terms calculated by Eqs.(1) and (3). The potential for the neutral components of the Higgs fields are given by

$$
\begin{aligned}
V &= (m_{H_1}^2 + |\lambda S|^2)|H_1^0|^2 + (m_{H_2}^2 + |\lambda S|^2)|H_2^0|^2 + m_S^2 |S|^2 \\
&\quad + |\lambda H_1^0 H_2^0 + h S^2|^2 + (A_\lambda S H_1^0 H_2^0 + \frac{A_h}{3} S^3 + h.c.) \\
&\quad + \frac{1}{8}(g_Y^2 + g_2^2)(|H_1^0|^2 - |H_2^0|^2)^2 .
\end{aligned} \tag{13}
$$

The soft terms are given as follows:

$$m_{H_1}^2 = \frac{1}{2}(4\pi)^2 \frac{d}{d\ln\mu} \left(f_\tau^2 + 3f_b^2 + \lambda^2 - \frac{1}{2}g_Y^2 - \frac{3}{2}g_2^2 \right) M^2 + \frac{1}{2}D_Y , \tag{14}$$

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$$m_{H_2}^2 = \frac{1}{2}(4\pi)^2 \frac{d}{d\ln\mu}\left(3f_t^2 + \lambda^2 - \frac{1}{2}g_Y^2 - \frac{3}{2}g_2^2\right)M^2 - \frac{1}{2}D_Y \ , \tag{15}$$

$$A_\lambda = -\lambda\left(3f_t^2 + f_\tau^2 + 3f_b^2 + 4\lambda^2 + 2h^2 + 3k_D^2 + 2k_L^2 - g_Y^2 - 3g_2^2\right)M \ , \tag{16}$$

$$A_h = -h\left(6\lambda^2 + 6h^2 + 9k_D^2 + 6k_L^2\right)M \ , \tag{17}$$

and m_S^2 is given in Eq.(12). Here, f_t, f_b, and f_τ are the Yukawa coupling constants of the top quark, the bottom quark, and the tau lepton, respectively, and we ignored those for the first and the second generations. The scale dependence of the Yukawa and gauge coupling constants are given in Appendix A. Notice that k_D enhances the A-terms while suppressing m_S^2 in Eq.(12), which makes it possible for S to acquire a large VEV. The minimization conditions with respect to the three Higgs fields S, H_1^0, and H_2^0 are given by

$$A_\lambda S + \lambda h S^2 = -\frac{\sin 2\beta}{2}(m_1^2 + m_2^2 + \lambda^2 v^2) \tag{18}$$

$$m_Z^2 = -\frac{m_1^2 - m_2^2}{\cos 2\beta} - (m_1^2 + m_2^2) \tag{19}$$

$$0 = \lambda^2 v^2 + 2h^2 S^2 + (\lambda h + \frac{A_\lambda}{2S})v^2 \sin 2\beta + A_h S + m_S^2 \ , \tag{20}$$

where $m_1 \equiv m_{H_1}^2 + \lambda^2 S^2$ and $m_2 \equiv m_{H_2}^2 + \lambda^2 S^2$, and $v^2 = \langle H_1^0\rangle^2 + \langle H_2^0\rangle^2$.

We numerically solve the Eqs.(18–20) while fixing $v = 174$ GeV. Once we postulate values of k_D, k_L, h, $\tan\beta$, and M, the three equations determine the sizes of λ, S, and D_Y. Stable minima are found with reasonable parameter sets as shown in the gray shaded region in Fig.2. The left and right figures correspond to solutions with negative and positive values of λ, respectively. The solutions are found for the range $1.5 \lesssim \tan\beta \lesssim 19$ and $0 < h \lesssim 0.3$. The parameters are fixed in the plots to be $k_D = 0.6$ and $k_L = 0.35$ motivated by the unification at the GUT scale, and $M = 600$ GeV. In the dark gray shaded region, the lightest SUSY particle (LSP) is a neutralino (mostly the wino). The stop and/or stau are lighter than the neutralino outside the region. The experimental bound on the lightest Higgs boson mass [42] is satisfied inside the dashed line. We can see the overlaps of the two regions for the $\lambda < 0$ case.

We plot those solutions λ, λS, and D_Y/M^2 as functions of h in Fig.3. Solutions are found in the gray shaded region and again the LSP is a neutralino in the dark gray shaded region. Lines correspond to various values of $\tan\beta$. We see in Fig.3 (a) and (b) that the values of $|\lambda|$ and h are restricted to be small such as less than 0.3. The large values are not preferred since those enhance m_S^2. Fig.3(c) and (d) show that we obtain large effective μ-parameter of the order of 1000 GeV such that the Higgsino is heavy enough. The negative values of D_Y in Fig.3(e) and (f) ensure the positive contributions to the slepton masses squared by taking D_{B-L} in an appropriate range.
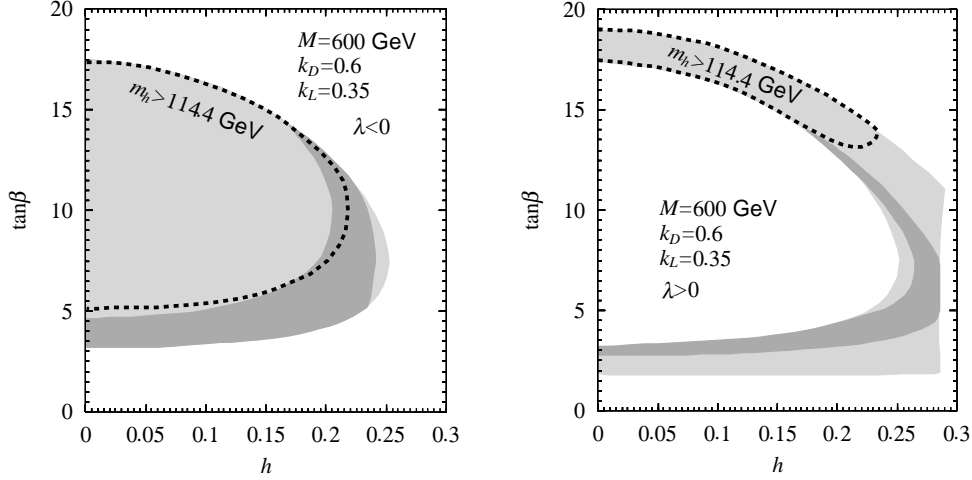
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Figure 2: The parameter regions where stable solutions are found in the modified NMSSM are plotted (the light gray shaded region). Left and right figures correspond to the solutions with $\lambda < 0$ and $\lambda > 0$, respectively. The LSP is a neutralino in the dark shaded region and the lightest Higgs boson is heavier than 114.4 GeV inside the dashed line. The other parameters are fixed to be $k_D = 0.6$, $k_L = 0.35$, and $M = 600$ GeV.

4.3 Spectrum and phenomenology

We show a mass spectrum at a sample point in the parameter space and discuss phenomenological implications. As a point with the neutralino LSP, we choose parameters such as $h = 0.21$ and $\tan\beta = 10$ in Fig.2. The mass spectra are listed in Table 1. We take the $B - L$ charge of the extra-fields to be $D : 1/3$ and $L : -1$, although those are, in principle, arbitrary.

The lightest Higgs boson h_1^0 is likely to be mainly composed by the doublet Higgs fields H_1^0 and H_2^0. There is no significant mixing between the singlet and the doublet Higgs boson because of the small λ parameter. Therefore properties of the lightest Higgs boson are similar to those of the MSSM. In particular, h_1^0 is lighter than the Z-boson at tree level, and hence the radiative corrections from the (s)top loop diagrams are important to satisfy the experimental bound [43, 44, 45]. The large stop masses are required in order to obtain the large radiative corrections. As in the MSSM, this requires a relatively high SUSY scale such as $M \gtrsim 600$ GeV. The charged Higgs boson mass for this parameter is well above the limit from the $b \to s\gamma$, $m_{H^\pm} \gtrsim 350$ GeV [46].[†]

The LSP is mostly the wino, the $SU(2)_L$ gaugino. The large Higgsino mass parameter λS indicates the small gaugino-Higgsino mixing and thus the charged and neutral winos are highly degenerate. The dominant contribution to the mass splitting is the one-loop radiative correction from diagrams with the gauge boson loops, and is estimated to be 165 MeV for large M [47, 48].

[†]Because the solutions require $\lambda < 0$ and hence the chargino diagram constructively interferes with the charged Higgs boson diagram, the constraint is probably somewhat stronger than this in our case, but it is clear that a slightly higher M can satisfy the limit if needed. Detailed quantitative discussions on this constraint is beyond the scope of this paper.
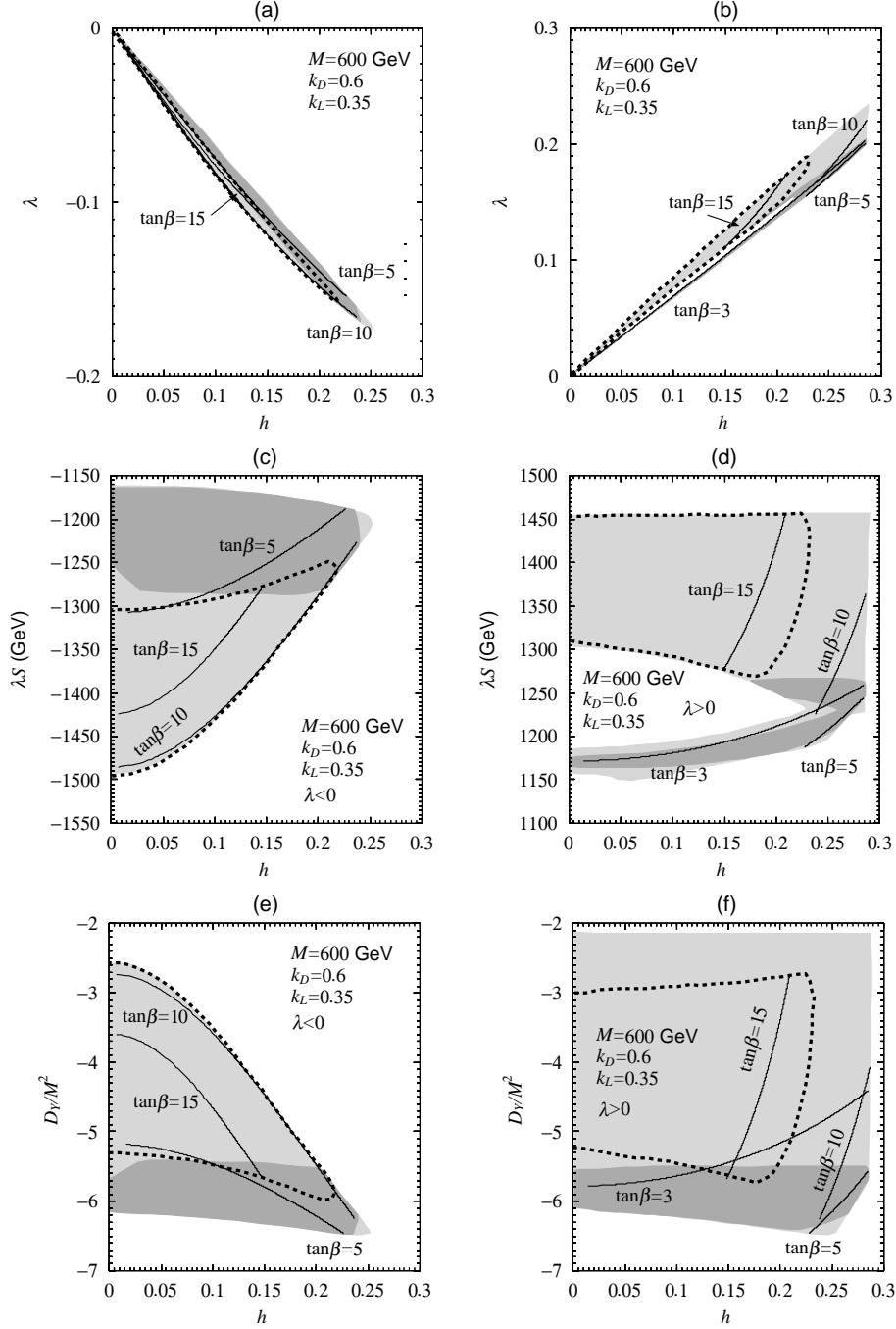
Figure 3: The regions with stable solutions are shown. Lines show the h dependence of λ (a,b), λS (c,d), and D_Y (e,f) with fixed values of $\tan\beta$. The neutralino LSP is realized in the dark shaded region. Left (a,c,e) and right (b,d,f) figures correspond to the solutions with $\lambda < 0$ and $\lambda > 0$, respectively. The other parameters are fixed to be $M = 600$ GeV, $k_D = 0.6$, and $k_L = 0.35$. The overlaps of the Higgs mass constraint and the dark shaded region in (f) is fictitious as the parameter space is covered multiple times.

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h_1^0	115	$\tilde{l}_L$	630		
h_2^0	464	$\tilde{e}_R$	509	D	4969
h_3^0	2743	$\tilde{\tau}_1$	495	$\tilde{D}_1$	2294
A_1^0	462	$\tilde{\tau}_2$	628	$\tilde{D}_2$	6973
A_2^0	3703	$\tilde{\nu}$	625	L	2898
$H^\pm$	470	$\tilde{\nu}_\tau$	621	$\tilde{L}_1$	1927
$\chi_1^\pm$	489	$\tilde{u}_L$	1564	$\tilde{L}_2$	3560
$\chi_2^\pm$	1277	$\tilde{u}_R$	1397		
χ_1^0	489	$\tilde{t}_1$	552		
χ_2^0	979	$\tilde{t}_2$	1280		
χ_3^0	1273	$\tilde{d}_L$	1566		
χ_4^0	1280	$\tilde{d}_R$	2013		
χ_5^0	3478	$\tilde{b}_1$	1256		
$\tilde{g}$	1459	$\tilde{b}_2$	1984		

Table 1: A sample mass spectrum is shown in the unit of GeV. We take $M = 600$ GeV, $\lambda = -0.15$, $h = 0.21$, $D_Y = -5.7M^2$, $D_{B-L} = 4.6M^2$, and $\tan\beta = 10$. The k_D and k_L coupling constants are taken to be 0.6 and 0.35, respectively. The Higgs potential has a minimum with the correct size of the Higgs VEV and $S = 8.3$ TeV. We have included the one-loop correction to the lightest Higgs boson mass.

In this case, the charged wino mainly decays into a charged pion and a neutral wino with long lifetimes, so that we may see the highly-ionizing charged tracks at the hadron or e^+e^- linear colliders [48, 49]. We discuss a scenario with the wino dark matter in the next subsection.

The little fine-tuning problem in the MSSM is left in the model [50, 22]. The non-observation of the Higgs boson requires relatively heavy stops, indicating a high SUSY breaking scale. A certain degree of fine-tuning is necessary to obtain the correct size of the Higgs VEV. In the NMSSM, the λSH_1H_2 coupling may raise the Higgs boson mass without upsetting the perturbativity if $\lambda \lesssim 0.7$. However, in the model with anomaly mediation, it is of no help since the λ parameter is necessary to be as small as 0.3.

4.4 Wino cold dark matter

The wino is a perfect candidate for the dark matter of the universe despite its large annihilation cross section [48, 51]. The current abundance can be explained by the non-thermal production from the decay of the gravitinos which were produced when the universe had high temperatures. In the following, we discuss the relic mass density of the wino including the effects of the non-thermal production.

By assuming the instantaneous decay of the gravitino at the decay temperature T_d, the yield[‡] of the wino from the gravitino decay is equal to that of the gravitino at T_d [52, 53] and is given

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[‡]The yield is defined relative to the entropy density.

by

$$Y_\chi^{\rm NT}(T_d) \simeq Y_{3/2}(T_{RH}) \simeq 1.9 \times 10^{-12} \times \left(\frac{T_{RH}}{10^{10}\,{\rm GeV}} \right), \tag{21}$$

where T_{RH} is the reheating temperature of the universe. The decay temperature T_d is given by

$$T_d \simeq 9\,{\rm MeV} \left(\frac{10}{g_*(T_d)} \right)^{1/4} \left(\frac{m_{3/2}}{100\,{\rm TeV}} \right)^{3/2} \simeq 0.8\,{\rm MeV} \left(\frac{10}{g_*(T_d)} \right)^{1/4} \left(\frac{m_{\rm wino}}{100\,{\rm GeV}} \right)^{3/2}, \tag{22}$$

where $g_*(T)$ denotes the number of the effective massless degrees of freedom and we use $m_{\rm wino} \simeq 5.2 \times 10^{-3} m_{3/2}$ (see Eq. (1)). In the limit of the pure wino LSP, the dominant annihilation process is the pair annihilation into two W-bosons via t-channel charged wino exchange diagram. Because T_d is much lower than the wino mass, we can take the non-relativistic limit of the cross section as a good approximation and that is given by [51]

$$\langle \sigma v \rangle = \frac{g_2^4}{2\pi} \frac{1}{m_{\rm wino}^2} \frac{(1 - m_W^2/m_{\rm wino}^2)^{3/2}}{(2 - m_W^2/m_{\rm wino}^2)^2}. \tag{23}$$

With the annihilation cross section, we obtain the yield Y_χ at low temperature T by solving the Boltzmann equation as follows [54]:

$$Y_\chi(T) = \left[\frac{1}{Y_\chi^{\rm TH}(T_d) + Y_\chi^{\rm NT}(T_d)} + \frac{1}{Y_\chi^{\rm ann}(T_d, T)} \right]^{-1}. \tag{24}$$

Here, $Y_\chi^{\rm TH}(T_d)$ is the thermal relic of the wino, and $Y_\chi^{\rm ann}(T_d, T)$ represents the annihilation effects after the non-thermal production. Those quantities are given by

$$Y_\chi^{\rm TH}(T_d) \simeq 10^{-14} \times \left(\frac{m_{\rm wino}}{100\,{\rm GeV}} \right), \tag{25}$$

$$Y_\chi^{\rm ann}(T_d, T) \simeq \sqrt{\frac{45}{8\pi^2 g_*(T_d)}} \frac{1}{\langle \sigma v \rangle M_{\rm Pl}(T_d - T)} \tag{26}$$

$$\simeq 2 \times 10^{-10} \times \left(\frac{g_*(T_d)}{10} \right)^{-1/4} \left(\frac{m_{\rm wino}}{100\,{\rm GeV}} \right)^{1/2}. \tag{27}$$

We can neglect the second term in Eq. (24) for $T_{RH} \lesssim 10^{12}$ GeV, since $Y_\chi^{\rm NT}(T_d)$ is much smaller than $Y_\chi^{\rm ann}(T_d, T)$. In this range, the yield Y_χ and the mass density parameter of the wino Ω_χ are simply given by

$$Y_\chi(T) \simeq 10^{-14} \times \left(\frac{m_{\rm wino}}{100\,{\rm GeV}} \right) + 1.9 \times 10^{-12} \times \left(\frac{T_{RH}}{10^{10}\,{\rm GeV}} \right), \tag{28}$$

$$\begin{aligned} \Omega_\chi h^2 &= \frac{m_{\rm wino} Y_\chi}{3.6221 \times 10^{-9}\,{\rm GeV}} \\ &\simeq 2.8 \times 10^{-4} \times \left(\frac{m_{\rm wino}}{100\,{\rm GeV}} \right)^2 + 5.3 \times 10^{-2} \times \left(\frac{m_{\rm wino}}{100\,{\rm GeV}} \right) \left(\frac{T_{RH}}{10^{10}\,{\rm GeV}} \right). \end{aligned} \tag{29}$$
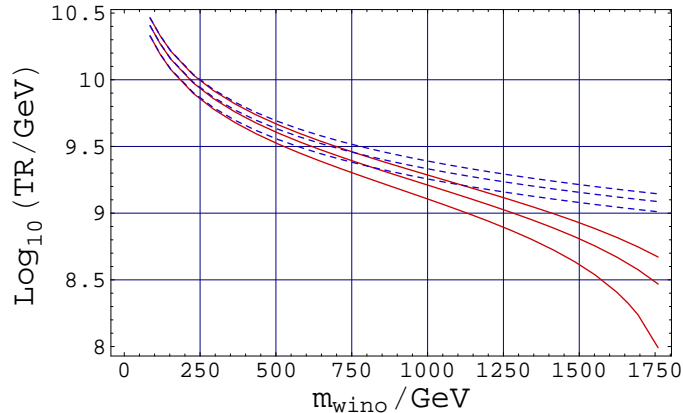
13

Figure 4: The required reheating temperature in order for the wino to be the dark matter. The solid lines corresponds to the required reheating temperature for $\Omega_\chi h^2 \simeq 0.095, 0.11, 0.13$ and the dashed lines for $\Omega_\chi^{\mathrm{NT}} h^2 \simeq 0.095, 0.11, 0.13$ from the bottom to the top, respectively.

Eq.(29) indicates that the wino can be the dominant component of the dark matter (i.e. $\Omega_\chi h^2 \simeq 0.11$) for an appropriate reheating temperature.

We show in Fig. 4 the required reheating temperature of the universe as a function of the wino mass in order for the wino to be the dark matter. Solid lines correspond to $\Omega_\chi h^2 \simeq 0.095, 0.11, 0.13$ from the bottom to the top, respectively. We also plot in the same figure the reheating temperature which satisfy $\Omega_\chi^{\mathrm{NT}} h^2 = m_{\mathrm{wino}} Y_\chi^{\mathrm{NT}}/(3.6221 \times 10^{-9}\, \mathrm{GeV}) \simeq 0.095, 0.11, 0.13$ as dashed lines. The mass density $\Omega_\chi^{\mathrm{NT}} h^2$ is the component of the non-thermally produced wino through the gravitino decay. As we see, $\Omega_\chi^{\mathrm{NT}}$ dominates the mass density of the dark matter for $m_{\mathrm{wino}} \lesssim 500$ GeV.

Remarkably, the required reheating temperature is consistent with the lower bound on T_{RH} for the thermal leptogenesis, $T_{RH} \gtrsim 4 \times 10^9$ GeV [55, 56]. The lower bound indicates the upper bound on the wino mass to be $m_{\tilde{W}} \lesssim 500$ GeV.

4.5 A solution to the strong CP problem

The UV insensitive anomaly mediation has a potential to solve the strong CP problem by the Nelson-Barr [40, 41]. The mechanism is claimed not to be a good solution in SUSY models because the SUSY breaking effect reintroduces the strong CP phase $\bar{\theta}$ at one-loop level [57]. However, with the UV insensitivity of the SUSY breaking terms, it is revived as a solution to the problem (see also [58, 59]).

The mechanism needs CP to be an exact symmetry of the Lagrangian such that the $\bar{\theta}$ parameter as well as the phase in the CKM matrix vanish. The observed non-vanishing CP phase in the CKM matrix can be induced at low energies by the spontaneous CP violation. Nelson and Barr proposed models in which we obtain only the CKM phase by the spontaneous CP violation while $\bar{\theta}$ remains vanishing at tree level. The simplest model is the following. We

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introduce a pair of vector-like quarks d_4^c and $\bar{d}_4^c$ which have the same and opposite quantum numbers as the right-handed down-type quarks. With the superpotential

$$W = aX d_i^c \bar{d}_4^c + f_d^{ij} q_i H_d d_j^c + m d_4^c \bar{d}_4^c , \tag{30}$$

where $i, j = 1, 2, 3$ and $\alpha = 1 - 4$, the phase in the VEV of X induce the CKM phase, but the determinant of the mass matrix remains to be real indicating no contribution to $\bar{\theta}$.

In general, such new matter fields with flavor-dependent couplings reintroduce the SUSY flavor problem through their loops. It is the non-trivial virtue of the UV insensitivity that decouples these interactions from the low-energy soft terms so that we can discuss such additional interactions without conflicting the phenomenological limits. In turn, the UV insensitivity turns out to be crucial for the mechanism as well as we discuss below.

Once we realize the situation with the non-zero CKM phase and vanishing $\bar{\theta}$ at tree level, the non-renormalization theorem in SUSY ensures vanishing $\bar{\theta}$ even at quantum level. Potentially problematic is the SUSY breaking effect as mentioned before. One-loop diagrams with SUSY breaking couplings induce the $\bar{\theta}$ parameter and acceptably small values of $\bar{\theta}$ require extraordinarily high degrees of degeneracy among SUSY breaking parameters [57]. However, the UV insensitive anomaly mediation as well as gauge mediation [60] does not suffer from the problem. The UV insensitivity ensures that all the SUSY breaking terms are described by the SUSY invariant quantities such as the Yukawa and gauge coupling constants, and therefore there is no new Jarlskog invariant other than that of the SM, i.e., $\Im(\det[m_u m_u^\dagger, m_d m_d^\dagger])$. With the small Jarlskog invariant and significant loop suppression factors, the $\bar{\theta}$ parameter remains to be very small such as $10^{-29} - 10^{-19}$ [59], which is much smaller than the experimental upper bound of 10^{-10}.

4.6 The cosmological domain wall and the tadpole problems

One may worry about the formation of the cosmological domain wall associated with the Z_3 symmetry breaking in the NMSSM [61]. However, with the presence of the vector-like quarks D and $\bar{D}$, there is no domain wall problem. Since the Z_3 symmetry, under which all the chiral superfields has a unit charge, is anomalous with respect to the $SU(3)_C$ gauge group, the instanton effect can give a sufficient energy shift among domains such that the domain wall is unstable [62].

Related to this issue, there is a problem of the instability of the Higgs potential caused by a loop correction associated with the SUSY breaking effect. The tadpole diagrams of S with gravitational interactions diverge quadratically, and reintroduce the hierarchy problem [61, 63]. The appearance of the linear term can be forbidden if the gravitational interaction respects the Z_3 symmetry. It is, however, argued that the quantum gravity violates all the global symmetry. Since the Z_3 symmetry is anomalous, we cannot think of the symmetry as a gauge symmetry, and thus it is expected to be broken at the Planck scale.

Although it may not be a problem once we understand quantum gravity, it is possible to control the tadpole divergence by embedding Z_3 to a higher anomaly free symmetry Z_{3N} such that the Z_3 symmetry is realized as an accidental approximate symmetry after breaking of Z_{3N}. For example, the Z_{3N} symmetry, under which all the chiral superfields have charge N, can be

made anomaly free when we add N pairs of vector-like quarks $\mathcal{D}$ and $\bar{\mathcal{D}}$ which have the same quantum numbers as D and $\bar{D}$ under the SM gauge group and charge -1 under Z_{3N}. When the Z_{3N} symmetry is broken spontaneously by the VEV of a field Σ which has Z_{3N} charge 2, the vector-like fields acquire the mass of the order of $\langle \Sigma \rangle$ and decouple from the low-energy physics if the VEV is much larger than the electroweak scale. The mass term of the vector-like quarks D and $\bar{D}$ or the μ-term given through the VEV of Σ is naturally small by the Z_{3N} symmetry, because it restricts the form of the lowest dimensional interaction to be $\Sigma^X/M_{\mathrm{Pl}}^{X-1} D\bar{D}$ where $X = N/2$ or $2N$ for even or odd N, respectively. The tadpole term in the Lagrangian is also suppressed as $m_{3/2}\langle \Sigma \rangle^N/M_{\mathrm{Pl}}^{N-2} S$, which can be small enough.

Also, the Z_{3N} symmetry may naturally forbids the large mixing between ordinary quarks d^c and $\bar{D}$. For example, if the Z_{3N} charges of D and $\bar{D}$ are $N+n$ and $N-n$, respectively, the mixing is naturally suppressed by the Z_{3N} breaking effect. The small but finite mixing makes the vector-like fields unstable, which ensures the absence of the problem with overclosure of the universe [64].

5   Conclusions

We have examined electroweak symmetry breaking in the NMSSM with UV insensitive anomaly mediation. We added a vector-like pair of matter fields to solve the light Higgsino problem in the NMSSM with the anomaly mediation. Viable parameter regions are found to preserve the perturbative coupling unification.

With the success of the electroweak symmetry breaking, the model is a phenomenologically and cosmologically perfect package for particle physics. The SUSY FCNC and CP problems are solved thanks to the UV insensitivity of the anomaly mediation. Also, the strong CP phase is not induced by the SUSY breaking effect once we set vanishing $\bar{\theta}$ at tree level, which is claimed to be natural in the string theory context [65]. The lightest neutralino is the wino which is a good candidate of the cold dark matter assuming the non-thermal production from the gravitino decay. The thermal leptogenesis works without contradicting the constraint on the gravitino abundance from the BBN theory, since the gravitino decays before the BBN era. There is no cosmological domain wall problem associated with the spontaneous Z_3 symmetry breaking in the NMSSM because the Z_3 symmetry is broken by the QCD instanton effect in the presence of the vector-like quarks. The tadpole problem in the NMSSM may be avoided by imposing anomaly free Z_{3N} symmetry at high energy.

Acknowledgments

A   Anomalous dimensions and beta functions

We list the anomalous dimensions and the beta functions in the model which are necessary to compute the soft SUSY breaking terms.

The anomalous dimensions are given by

$$(4\pi)^2 \gamma_{H_1} = f_\tau^2 + 3f_b^2 + \lambda^2 - \frac{1}{2}g_Y^2 - \frac{3}{2}g_2^2 \ , \tag{31}$$

$$(4\pi)^2 \gamma_{H_2} = 3f_t^2 + \lambda^2 - \frac{1}{2}g_Y^2 - \frac{3}{2}g_2^2 \ , \tag{32}$$

$$(4\pi)^2 \gamma_S = 2\lambda^2 + 2h^2 + 3k_D^2 + 2k_L^2 \ , \tag{33}$$

$$(4\pi)^2 \gamma_{l_i} = f_\tau^2 \delta_{i3} - \frac{1}{2}g_Y^2 - \frac{3}{2}g_2^2 \ , \tag{34}$$

$$(4\pi)^2 \gamma_{e_i^c} = 2f_\tau^2 \delta_{i3} - 2g_Y^2 \ , \tag{35}$$

$$(4\pi)^2 \gamma_{q_i} = (f_b^2 + f_t^2)\delta_{i3} - \frac{1}{18}g_Y^2 - \frac{3}{2}g_2^2 - \frac{8}{3}g_3^2 \ , \tag{36}$$

$$(4\pi)^2 \gamma_{d_i^c} = 2f_b^2 \delta_{i3} - \frac{2}{9}g_Y^2 - \frac{8}{3}g_3^2 \ , \tag{37}$$

$$(4\pi)^2 \gamma_{u_i^c} = 2f_t^2 \delta_{i3} - \frac{8}{9}g_Y^2 - \frac{8}{3}g_3^2 \ , \tag{38}$$

$$(4\pi)^2 \gamma_D = k_D^2 - \frac{2}{9}g_Y^2 - \frac{8}{3}g_3^2 \ , \tag{39}$$

$$(4\pi)^2 \gamma_{\bar{D}} = k_D^2 - \frac{2}{9}g_Y^2 - \frac{8}{3}g_3^2 \ , \tag{40}$$

$$(4\pi)^2 \gamma_L = k_L^2 - \frac{1}{2}g_Y^2 - \frac{3}{2}g_2^2 \ , \tag{41}$$

$$(4\pi)^2 \gamma_{\bar{L}} = k_L^2 - \frac{1}{2}g_Y^2 - \frac{3}{2}g_2^2 \ , \tag{42}$$

where we neglect the Yukawa coupling constants for the first and second generations.

The beta functions for the gauge coupling constants are the following:

$$(4\pi)^2 \frac{dg_Y}{d\ln\mu} = \frac{38}{3}g_Y^3 \ , \tag{43}$$

$$(4\pi)^2 \frac{dg_2}{d\ln\mu} = 2g_2^3 \ , \tag{44}$$

$$(4\pi)^2 \frac{dg_3}{d\ln\mu} = -2g_3^3 \ . \tag{45}$$

The beta functions for the Yukawa coupling constants are expressed by the anomalous dimensions as follows:

$$\begin{aligned}(4\pi)^2 \frac{df_\tau}{d\ln\mu} &= (4\pi)^2 f_\tau(\gamma_{l_3} + \gamma_{H_1} + \gamma_{e_3^c}) \\ &= f_\tau(4f_\tau^2 + 3f_b^2 + \lambda^2 - 3g_Y^2 - 3g_2^2) \ , \end{aligned} \tag{46}$$

$$\begin{aligned}(4\pi)^2 \frac{df_b}{d\ln\mu} &= (4\pi)^2 f_b(\gamma_{q_3} + \gamma_{H_1} + \gamma_{d_3^c}) \\ &= f_b(f_\tau^2 + 6f_b^2 + f_t^2 + \lambda^2 - \frac{7}{9}g_Y^2 - 3g_2^2 - \frac{16}{3}g_3^2) \ , \end{aligned} \tag{47}$$

$$\begin{aligned}(4\pi)^2 \frac{df_t}{d\ln\mu} &= (4\pi)^2 f_t(\gamma_{q_3} + \gamma_{H_2} + \gamma_{u_3^c}) \\ &= f_t(6f_t^2 + f_b^2 + \lambda^2 - \frac{13}{9}g_Y^2 - 3g_2^2 - \frac{16}{3}g_3^2) \ , \end{aligned} \tag{48}$$

$$\begin{aligned}(4\pi)^2 \frac{d\lambda}{d\ln\mu} &= (4\pi)^2 \lambda(\gamma_S + \gamma_{H_1} + \gamma_{H_2}) \\ &= \lambda(f_\tau^2 + 3f_b^2 + 3f_t^2 + 4\lambda^2 + 2h^2 + 3k_D^2 + 2k_L^2 - g_Y^2 - 3g_2^2) \ , \end{aligned} \tag{49}$$

$$\begin{aligned}(4\pi)^2 \frac{dh}{d\ln\mu} &= (4\pi)^2 3h\gamma_S \\ &= 3h(2\lambda^2 + 2h^2 + 3k_D^2 + 2k_L^2) \ , \end{aligned} \tag{50}$$

$$
\begin{aligned}
(4\pi)^2 \frac{dk_D}{d\ln\mu} &= (4\pi)^2 k_D(\gamma_S + \gamma_D + \gamma_{\bar{D}}) \\
&= k_D\left(2\lambda^2 + 2h^2 + 5k_D^2 + 2k_L^2 - \frac{4}{9}g_Y^2 - \frac{16}{3}g_3^2\right) ,
\end{aligned} \tag{51}
$$

$$
\begin{aligned}
(4\pi)^2 \frac{dk_L}{d\ln\mu} &= (4\pi)^2 k_L(\gamma_S + \gamma_L + \gamma_{\bar{L}}) \\
&= k_L\left(2\lambda^2 + 2h^2 + 3k_D^2 + 4k_L^2 - g_Y^2 - 3g_2^2\right) .
\end{aligned} \tag{52}
$$

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