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Mediated Supersymmetry Breaking

T. Fukuyama – Ritsumeikan University

T. Kikuchi – KEK

N. Okada – KEK

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# D-term contributions to the mixed modulus–anomaly mediated supersymmetry breaking

Takeshi Fukuyama<sup>a</sup>, Tatsuru Kikuchi<sup>a,b,\*</sup>, Nobuchika Okada<sup>b,c</sup>

<sup>a</sup> *Department of Physics, Ritsumeikan University, Kusatsu, Shiga, 525-8577, Japan*

<sup>b</sup> *Theory Division, KEK, Oho 1-1, Tsukuba, Ibaraki, 305-0801, Japan*

<sup>c</sup> *Department of Particle and Nuclear Physics, The Graduate University for Advanced Studies, Oho 1-1, Tsukuba, Ibaraki, 305-0801, Japan*

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## Abstract

We investigate effects of  $D$ -term contributions to the mixed modulus–anomaly mediated supersymmetry breaking scenario. In the original scenario, the tachyonic slepton problem in the pure anomaly mediated supersymmetry breaking is cured by modulus contributions. We generalize the scenario so as to include contributions from the  $D$ -terms of  $U(1)_Y$  and the gauged  $U(1)_{B-L}$  which is motivated in a grand unified theory based on a higher rank gauge group such as  $SO(10)$ . As a consequence of additional  $D$ -term contributions to scalar masses, we obtain various soft supersymmetry breaking mass spectra, which are different from those obtained in the conventional mixed modulus–anomaly mediated supersymmetry breaking. Especially, we find that the lightest superpartner (LSP) neutralinos can be various types, such as higgsino-like, wino-like and bino-like degenerating with the next LSP sfermions.

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## 1. Introduction

Supersymmetry (SUSY) extension is one of the most promising way to solve the gauge hierarchy problem in the standard model [1]. However, since any superpartners have not been observed in current experiments, SUSY should be broken at low energies. Furthermore, soft SUSY breaking terms are severely constrained to be almost flavor blind and CP invariant. Thus, the SUSY breaking has to be mediated to the visible sector in some clever way not to induce too large CP and flavor violation effects. Some mechanisms to achieve such SUSY breaking mediations have been proposed [2].

The anomaly mediated supersymmetry breaking (AMSB) [3,4] is one of the most attractive scenario due to its flavor-blindness and ultraviolet (UV) insensitivity for the resultant soft SUSY breaking terms. Unfortunately, the pure AMSB scenario is obviously excluded, since it predicts slepton squared masses

being negative. There have been many attempts to solve this problem by taking into account additional positive contributions to slepton squared masses at tree level [3,5,6] or at quantum level [7,8]. Among them, adding  $D$ -terms of the  $U(1)_Y$  and the (gauged)  $U(1)_{B-L}$  may be the most interesting possibility, because these  $U(1)$  symmetries are anomaly-free with respect to the standard model gauge group so that the UV insensitivity is preserved [5]. We can expect such new contributions from the  $D$ -terms, if some grand unified theory (GUT) based on a higher rank gauge group such as  $SO(10)$  takes place at high energies, which includes new Higgs fields and the gauged  $U(1)_{B-L}$  as its subgroup. However, it has been found that this scenario requires a very small  $\tan\beta$  [9] to obtain the correct electroweak symmetry breaking. As a result, the top Yukawa coupling blows up far below the GUT scale, and the minimal supersymmetric standard model (MSSM) cannot be simply connecting into SUSY GUTs in the way usually expected.

Recently, Kachru–Kallosh–Linde–Trivedi (KKLT) [10] have proposed a way to stabilize the modulus in string theories with flux compactification. Interestingly, a stabilized modulus can induce additional SUSY breaking contributions comparable to

\* Corresponding author.

*E-mail addresses:* [fukuyama@se.ritsumei.ac.jp](mailto:fukuyama@se.ritsumei.ac.jp) (T. Fukuyama), [tatsuru@post.kek.jp](mailto:tatsuru@post.kek.jp) (T. Kikuchi), [okadan@post.kek.jp](mailto:okadan@post.kek.jp) (N. Okada).

the pure AMSB contributions, so as to solve tachyonic slepton problem. There have already been several studies on the SUSY breaking mediation in this KKL $T$  type setup [11–13], the so-called mixed modulus–anomaly mediation, and the characteristic sparticle mass spectrum have been obtained.

In this Letter, we generalize the mixed modulus–anomaly mediation scenario so as to include the effects of the  $D$ -term contributions. Contributions from the mixed modulus–anomaly mediation play a role to widen the allowed region of  $\tan\beta$ , so that the top Yukawa coupling remains perturbative until the GUT scale and the MSSM can simply connect into GUTs. On the other hand, the  $D$ -term contributions change sfermion mass spectrum from the one in the conventional mixed modulus–anomaly mediation scenario. As a result, the sparticle mass spectrum in our scenario can be quite different from the one obtained in the mixed modulus–anomaly mediation scenario or in the AMSB scenario with the  $D$ -term contributions.

This Letter is organized as follows. In Section 2, we briefly review on the KKL $T$  setup. In Section 3, we give formulas of the mixed modulus–anomaly mediation including  $D$ -term contributions from  $U(1)_Y$  and  $U(1)_{B-L}$ , which are necessary for our numerical analysis. The results of numerical analysis are presented in Section 4. The last section is devoted to summary and discussions.

## 2. KKL $T$ setup

In this section, we work out in the superconformal framework of supergravity [14]. A modulus superfield  $T$  plays a crucial role in the KKL $T$  setup, whose basic Lagrangian is given by

$$\mathcal{L} = -3 \int d^4\theta \phi^\dagger \phi e^{-K/3} + \int d^2\theta \phi^3 W + \text{h.c.}, \quad (1)$$

where  $\phi = 1 + \theta^2 F_\phi$  is the compensating multiplet. Here, the Kähler potential is taken to be the no-scale type,

$$K = -3 \ln(T + T^\dagger), \quad (2)$$

and the following superpotential is derived in the context of the type IIB string theory [10],

$$W = W_0 - C e^{-aT}, \quad (3)$$

where the first term is a constant, and the second term is generated through the  $SU(N_c)$  gaugino condensation with coefficients  $C$  and  $a = 8\pi^2/N_c$  being real and positive.<sup>1</sup>

With these Kähler potential and superpotential, the scalar potential is given by

$$V = \frac{T + T^\dagger}{3} |W_T|^2 - W W_T^\dagger - W^\dagger W_T, \quad (4)$$

where  $W_T = \partial W / \partial T$ . This scalar potential has a supersymmetric anti-de Sitter minimum,

$$V = -3 |F_\phi|^2 \left( T + T^\dagger + \frac{2}{a} \right) < 0, \quad (5)$$

with  $F_\phi = W_T^\dagger / 3$ . At the potential minimum, the  $F$ -term of the modulus is given by

$$F_T = \frac{2}{a} F_\phi = \frac{N_c}{4\pi^2} F_\phi. \quad (6)$$

In order to obtain a de Sitter (or Minkowski) vacuum, the lifting potential due to the presence of the anti-D3 brane is introduced [10],

$$\Delta V = \frac{D}{(T + T^\dagger)^n}, \quad (7)$$

where  $n$  is an integer ( $n = 2$  in the original KKL $T$  paper), and  $D$  is a constant whose value is tuned so as to realize the de Sitter (or Minkowski) vacuum. At the de Sitter (or Minkowski) vacuum, only  $\Re[T]$  of  $T$  has non-zero vacuum expectation value and the relation  $F_T \simeq F_\phi / a = N_c F_T / (4\pi^2)$  in Eq. (6) still holds.

Next, let us introduce the MSSM sector into the modulus Lagrangian. The Kähler potential is replaced to the one including the MSSM matter and Higgs superfields,  $K(T, T^\dagger) \rightarrow K(T, T^\dagger) + K_{\text{MSSM}}$ . For simplicity, we take the minimal Kähler potential for the MSSM superfields,  $K_{\text{MSSM}} = Q_i^\dagger e^{2g_a V_a} Q_i$ , where  $Q_i$  stands for the MSSM matter and Higgs superfields. Expanding  $e^{K/3}$ , the Kähler potential for the MSSM superfields is described as

$$\int d^4\theta \phi^\dagger \phi (T + T^\dagger) Q_i^\dagger e^{2g_a V_a} Q_i + \dots \quad (8)$$

For the gauge sector in the MSSM, the kinetic term is of the form,

$$\mathcal{L}_{\text{gauge}} = \frac{1}{4} \int d^2\theta f_a \mathcal{W}^{a\alpha} \mathcal{W}_\alpha^a. \quad (9)$$

We take the gauge kinetic function  $f_a = T$  in the following.

In the above setup, there are two SUSY breaking sources, namely  $F_\phi$  and  $F_T$ . Non-zero  $F_\phi$  induces soft SUSY breaking terms through the AMSB, and the resultant SUSY breaking mass scale is characterized by  $m_{\text{AMSB}} \sim F_\phi / (16\pi^2)$ . On the other hand, as can be easily seen from Eqs. (8) and (9), non-zero  $F_T$  leads to soft SUSY breaking terms at tree level, the modulus mediation. The resultant SUSY breaking mass scale in the modulus mediation is characterized by

$$m_{\text{modulus}} \sim \frac{F_T}{T + T^\dagger}. \quad (10)$$

Noting  $\Re[T] = 1/g_{\text{GUT}}^2 = \mathcal{O}(1)$  ( $g_{\text{GUT}}$  denotes the standard model gauge coupling at the GUT scale) and  $F_T \simeq N_c F_\phi / (4\pi^2)$ , we see that this contribution by the moduli mediation is comparable to the one by the AMSB,  $m_{\text{AMSB}} \sim m_{\text{modulus}}$ . This fact is the key of the mixed modulus–anomaly mediation scenario. According to the method developed in Ref. [15] (see also Ref. [7]), soft SUSY breaking terms (each gaugino masses  $M_a$ , sfermion squared masses  $\tilde{m}_i^2$  and  $A$ -parameters) at the GUT scale ( $\mu = M_{\text{GUT}} \simeq 2 \times 10^{16}$  GeV)<sup>2</sup>

<sup>1</sup> In the original work by KKL $T$  [10],  $C = 1$  and  $a = 0.1$  were used in order to realize de Sitter (or Minkowski) vacua.

<sup>2</sup> In this Letter, we set the compactification scale of the string theory to be the GUT scale.

can be extracted from renormalized gauge kinetic functions and SUSY wave function renormalization coefficients [11],

$$\begin{aligned} M_a &= M(\alpha + b_a g_a^2), \\ \tilde{m}_i^2 &= M^2(\alpha^2 + 2\alpha(T + T^\dagger)\partial_T \gamma_i - 8\pi^2 \mu \partial_\mu \gamma_i), \\ A_{ijk} &= M(3\alpha - \gamma_i - \gamma_j - \gamma_k). \end{aligned} \quad (11)$$

Here,  $g_a$  ( $g_1 = g_2 = g_3 = g_{\text{GUT}}$ ) are the gauge couplings,  $b_a$  are the beta function coefficients, and  $\gamma_i$  are the anomalous dimensions which depend on  $T$  through the  $T$ -dependence of the gauge couplings and Yukawa couplings. Parameters  $M$  (typical soft SUSY breaking mass scale) and  $\alpha$  are defined as

$$M = \frac{F_\phi}{16\pi^2} \sim \frac{m_{3/2}}{16\pi^2}, \quad (12)$$

$$\alpha M = \frac{F_T}{T + T^\dagger} \quad (13)$$

with the gravitino mass  $m_{3/2}$ . The results of the pure AMSB are reproduced in the limit  $\alpha \rightarrow 0$ , while the limit  $\alpha \gg 1$  corresponds to the pure modulus mediation whose contribution to sfermion masses is positive. As discussed above,  $\alpha = \mathcal{O}(1)$  is expected in the KKLT setup, so that both the AMSB and the moduli mediation give important contributions to resultant soft SUSY breaking parameters.

There are remaining two parameters in the Higgs sector, namely  $\mu$  and  $B\mu$  terms, that are responsible for electroweak symmetry breaking and should be of the order of the electroweak scale. As in the AMSB scenario, the natural value of the  $B$ -parameter would be  $B \sim m_{3/2} \gg M$ , and the Higgs sector should be extended in order to achieve the  $B$ -parameter being at the electroweak scale. Although some fine-tuning among parameters is necessary, the way to realize  $\mu \sim B \sim M$  have been discussed in Ref. [11]. In our analysis, we treat them as free parameters as usual, that is,  $\mu$  and  $B\mu$  are replaced into two free parameters  $\tan\beta$  and  $\text{sgn}(\mu)$ , while the value of  $|\mu|$  is determined by the stationary condition of the Higgs potential. In the next section, we consider to add two  $D$ -terms of  $U(1)_Y$  and  $U(1)_{B-L}$ , and hence total set of free parameters in our analysis is

$$\{M, \alpha, D_Y, D_{B-L}, \tan\beta, \text{sgn}(\mu)\}. \quad (14)$$

### 3. Mixed modulus–anomaly mediation including $D$ -terms

Now let us introduce the  $D$ -terms to the mixed modulus–anomaly mediation. If there exists a  $U(1)$  gauge multiplet having a non-zero  $D$ -term, the kinetic term of a matter superfield gives

$$\mathcal{L} = \int d^4\theta Q_i^\dagger e^{q_i V} Q_i \supset q_i D \tilde{Q}_i^\dagger \tilde{Q}_i, \quad (15)$$

where  $q_i$  is the  $U(1)$  charge of the chiral multiplet  $Q_i$ . This leads to a shift for the scalar squared mass,

$$\tilde{m}_i^2 \rightarrow \tilde{m}_i^2 - q_i D. \quad (16)$$

The  $U(1)$  symmetry providing the  $D$ -term should be anomaly-free in order not to induce quadratic divergence in a theory.

As such a  $U(1)$  symmetry, there exist two candidates in the MSSM, namely  $U(1)_Y$  and gauged  $U(1)_{B-L}$ . Introduction of this  $U(1)_{B-L}$  gauge symmetry is well-motivated, if we assume that the MSSM is embedded into a GUT based on a higher rank gauge group such as  $SO(10)$  which includes the gauged  $U(1)_{B-L}$  as a subgroup. This possibility is our motivation to consider the  $D$ -term in addition to the mixed modulus–anomaly mediation. Normally, many extra Higgs fields are involved in such models, and some of them have non-zero vacuum expectation values to break the GUT symmetry at the supersymmetric level. Once soft SUSY breaking terms for these Higgs fields are taken into account, the vacuum would be realized at the point slightly away from the D-flat directions, so that non-zero  $D$ -terms are developed. Although it depends on the detailed structure of the Higgs sector, we may naturally expect the scale of the  $D$ -term to be  $D \sim M^2$ .

Calculating the anomalous dimensions [11], all the soft SUSY breaking terms can be obtained from Eq. (11). Taking  $U(1)_Y$  and  $U(1)_{B-L}$   $D$ -term contributions into account, the soft scalar masses for the first two generations at the GUT scale are explicitly written as

$$\begin{aligned} m_{\tilde{q}_{1,2}}^2 &= M^2 \left[ \frac{157}{25} g_{\text{GUT}}^4 - \frac{42}{5} g_{\text{GUT}}^2 \alpha + \alpha^2 - \frac{1}{6} \alpha_Y - \frac{1}{3} \alpha_{B-L} \right], \\ m_{\tilde{u}_{1,2}}^2 &= M^2 \left[ \frac{112}{25} g_{\text{GUT}}^4 - \frac{32}{5} g_{\text{GUT}}^2 \alpha + \alpha^2 + \frac{2}{3} \alpha_Y + \frac{1}{3} \alpha_{B-L} \right], \\ m_{\tilde{d}_{1,2}}^2 &= M^2 \left[ \frac{178}{25} g_{\text{GUT}}^4 - \frac{28}{5} g_{\text{GUT}}^2 \alpha + \alpha^2 - \frac{1}{3} \alpha_Y + \frac{1}{3} \alpha_{B-L} \right], \\ m_{\tilde{\ell}_{1,2}}^2 &= M^2 \left[ -\frac{87}{25} g_{\text{GUT}}^4 - \frac{18}{5} g_{\text{GUT}}^2 \alpha + \alpha^2 + \frac{1}{2} \alpha_Y + \alpha_{B-L} \right], \\ m_{\tilde{e}_{1,2}}^2 &= M^2 \left[ -\frac{198}{25} g_{\text{GUT}}^4 - \frac{12}{5} g_{\text{GUT}}^2 \alpha + \alpha^2 - \alpha_Y - \alpha_{B-L} \right]. \end{aligned} \quad (17)$$

Here, we have defined  $\alpha_Y$  and  $\alpha_{B-L}$  as

$$\alpha_Y \equiv \frac{D_Y}{M^2}, \quad \alpha_{B-L} \equiv \frac{D_{B-L}}{M^2}, \quad (18)$$

and Yukawa couplings of the first two generations have been neglected as a good approximation. For the third generation sfermion masses, Yukawa couplings are involved,

$$\begin{aligned} m_{\tilde{q}_3}^2 &= M^2 \left[ \frac{157}{25} g_{\text{GUT}}^4 + y_t^2 b_{y_t} + y_b^2 b_{y_b} \right. \\ &\quad \left. - \left\{ \frac{42}{5} g_{\text{GUT}}^2 - 6(y_t^2 + y_b^2) \right\} \alpha + \alpha^2 - \frac{1}{6} \alpha_Y - \frac{1}{3} \alpha_{B-L} \right], \\ m_{\tilde{u}_3}^2 &= M^2 \left[ \frac{112}{25} g_{\text{GUT}}^4 + 2y_t^2 b_{y_t} - \left( \frac{32}{5} g_{\text{GUT}}^2 - 12y_t^2 \right) \alpha \right. \\ &\quad \left. + \alpha^2 + \frac{2}{3} \alpha_Y + \frac{1}{3} \alpha_{B-L} \right], \\ m_{\tilde{d}_3}^2 &= M^2 \left[ \frac{178}{25} g_{\text{GUT}}^4 + 2y_b^2 b_{y_b} - \left( \frac{28}{5} g_{\text{GUT}}^2 - 12y_b^2 \right) \alpha \right. \\ &\quad \left. + \alpha^2 - \frac{1}{3} \alpha_Y + \frac{1}{3} \alpha_{B-L} \right], \end{aligned}$$

$$\begin{aligned}
m_{\tilde{\ell}_3}^2 &= M^2 \left[ -\frac{87}{25} g_{\text{GUT}}^4 + y_\tau^2 b_{y_\tau} - \left( \frac{18}{5} g_{\text{GUT}}^2 - 6y_\tau^2 \right) \alpha \right. \\
&\quad \left. + \alpha^2 + \frac{1}{2} \alpha_Y + \alpha_{\text{B-L}} \right], \\
m_{\tilde{e}_3}^2 &= M^2 \left[ -\frac{198}{25} g_{\text{GUT}}^4 + 2y_\tau^2 b_{y_\tau} - \left( \frac{12}{5} g_{\text{GUT}}^2 - 12y_\tau^2 \right) \alpha \right. \\
&\quad \left. + \alpha^2 - \alpha_Y - \alpha_{\text{B-L}} \right], \tag{19}
\end{aligned}$$

where  $b_{y_t}$ ,  $b_{y_b}$  and  $b_{y_\tau}$  are given by

$$\begin{aligned}
b_{y_t} &= 6y_t^2 + y_b^2 - \frac{46}{5} g_{\text{GUT}}^2, \\
b_{y_b} &= y_t^2 + 6y_b^2 + y_\tau^2 - \frac{44}{5} g_{\text{GUT}}^2, \\
b_{y_\tau} &= 3y_b^2 + 4y_\tau^2 - \frac{24}{5} g_{\text{GUT}}^2. \tag{20}
\end{aligned}$$

When the condition,  $\alpha_Y < -\alpha_{\text{B-L}} < 1/2\alpha_Y$ , is satisfied with  $\alpha_Y < 0$  and  $\alpha_{\text{B-L}} > 0$ , slepton squared masses obtain positive contributions from  $D$ -terms. On the other hand, the  $D$ -terms in this region give negative contributions to  $m_{\tilde{q}_i}^2$  and  $m_{\tilde{u}_i}^2$ , while positive to  $m_{\tilde{d}_i}^2$ .

A-parameters are given by

$$A_{ijk} = M(3\alpha - \gamma_i - \gamma_j - \gamma_k) \tag{21}$$

with explicit formulas of the anomalous dimensions,

$$\begin{aligned}
\gamma_{q_i} &= \frac{21}{5} g_{\text{GUT}}^2 - (y_t^2 + y_b^2) \delta_{i3}, \\
\gamma_{u_i^c} &= \frac{16}{5} g_{\text{GUT}}^2 - 2y_t^2 \delta_{i3}, \\
\gamma_{d_i^c} &= \frac{14}{5} g_{\text{GUT}}^2 - 2y_b^2 \delta_{i3}, \\
\gamma_{\ell_i} &= \frac{9}{5} g_{\text{GUT}}^2 - y_\tau^2 \delta_{i3}, \\
\gamma_{e_i^c} &= \frac{6}{5} g_{\text{GUT}}^2 - 2y_\tau^2 \delta_{i3}, \\
\gamma_{H_1} &= \frac{9}{5} g_{\text{GUT}}^2 - 3y_b^2 - y_\tau^2, \\
\gamma_{H_2} &= \frac{9}{5} g_{\text{GUT}}^2 - 3y_t^2. \tag{22}
\end{aligned}$$

Also, the Higgs soft masses at the GUT scale are given by

$$\begin{aligned}
m_{H_1}^2 &= M^2 \left[ -\frac{87}{25} g_{\text{GUT}}^4 + 3y_b^2 b_{y_b} + y_\tau^2 b_{y_\tau} \right. \\
&\quad \left. - \left( \frac{18}{5} g_{\text{GUT}}^2 - 18y_b^2 - 6y_\tau^2 \right) \alpha + \alpha^2 + \frac{1}{2} \alpha_Y \right], \\
m_{H_2}^2 &= M^2 \left[ -\frac{87}{25} g_{\text{GUT}}^4 + 3y_t^2 b_{y_t} - \left( \frac{18}{5} g_{\text{GUT}}^2 - 18y_t^2 \right) \alpha \right. \\
&\quad \left. + \alpha^2 - \frac{1}{2} \alpha_Y \right]. \tag{23}
\end{aligned}$$

The Higgs mass parameters,  $\mu$ -term and  $B\mu$ -term, are determined from the electroweak symmetry breaking conditions,

$$\begin{aligned}
\mu^2 &= \frac{m_{H_1}^2 - m_{H_2}^2 \tan^2 \beta}{\tan^2 \beta - 1} - \frac{1}{2} M_Z^2, \\
B\mu &= \frac{1}{2} [m_{H_1}^2 + m_{H_2}^2 + 2\mu^2] \sin 2\beta. \tag{24}
\end{aligned}$$

In numerical analysis presented in the next section, we found that the condition to provide the correct electroweak symmetry breaking gives the most severe constraint on the parameter space ( $\alpha$ ,  $\alpha_Y$ ,  $\alpha_{\text{B-L}}$ ), rather than that to provide non-tachyonic sfermion masses  $m_{\tilde{f}}^2 > 0$ .

Finally, we give the explicit formulas of the gaugino masses at the GUT scale,

$$\begin{aligned}
M_1 &= M \left( \alpha + \frac{33}{5} g_{\text{GUT}}^2 \right), \\
M_2 &= M(\alpha + g_{\text{GUT}}^2), \\
M_3 &= M(\alpha - 3g_{\text{GUT}}^2). \tag{25}
\end{aligned}$$

Inputting the soft SUSY breaking terms expressed above at the GUT scale as the boundary conditions, the soft SUSY breaking terms at the electroweak scale are obtained through the renormalization group equations (RGEs). In the next section, we show the resultant soft SUSY breaking mass spectrum for various inputs of  $M$ ,  $\alpha$ ,  $\alpha_Y$  and  $\alpha_{\text{B-L}}$  with a given  $\tan \beta$ .

#### 4. Numerical results

Now we are ready to perform a numerical evaluation by using the formulas presented in the previous sections. With given  $\tan \beta$  and the parameter set ( $\alpha$ ,  $\alpha_Y$ ,  $\alpha_{\text{B-L}}$ ), we input the formulas for soft SUSY breaking terms at the GUT scale, and then evolve them according to the one-loop RGEs [16]. In our analysis, we take an averaged soft SUSY breaking mass scale as  $M_S = 500$  GeV and evaluate all the soft SUSY breaking parameters at this scale. As examples, we investigate the cases of  $\tan \beta = 10$  and  $\tan \beta = 45$  with the unified gauge coupling constant  $\alpha_{\text{GUT}}^{-1} = 25.4$  at  $M_{\text{GUT}}$ . As a good approximation, we consider Yukawa couplings only for fermions in the third generation with input values at  $M_{\text{GUT}}$  as

$$y_t = 0.635, \quad y_b = 0.0616, \quad y_\tau = 0.0687, \tag{26}$$

for  $\tan \beta = 10$  and

$$y_t = 0.749, \quad y_b = 0.449, \quad y_\tau = 0.454, \tag{27}$$

for  $\tan \beta = 45$ , respectively.

First we examine the allowed region of the parameter space ( $\alpha$ ,  $\alpha_Y$ ,  $\alpha_{\text{B-L}}$ ) for given  $\tan \beta = 10$  and 45 and  $M = 500$  GeV. Sparticle mass spectrum for various inputs in the range of  $0 \leq \alpha \leq 6$  and  $-10 \leq \alpha_Y, \alpha_{\text{B-L}} \leq 10$  has been calculated in every 0.2 intervals for  $\alpha$  and in every 0.5 intervals for  $\alpha_Y$  and  $\alpha_{\text{B-L}}$ . The allowed parameter sets of ( $\alpha$ ,  $\alpha_Y$ ) and ( $\alpha$ ,  $\alpha_{\text{B-L}}$ ) are plotted in Figs. 1 and 2, for which resultant sfermion squared masses are all positive and the electroweak symmetry breaking is correctly achieved. In both  $\tan \beta = 10$  and 45 cases, the allowed region is severely constrained for  $\alpha \lesssim 2$  mainly due to the condition for the correct electroweak symmetry breaking. In particular, for a large  $\tan \beta$ , the soft mass squared of the down-type



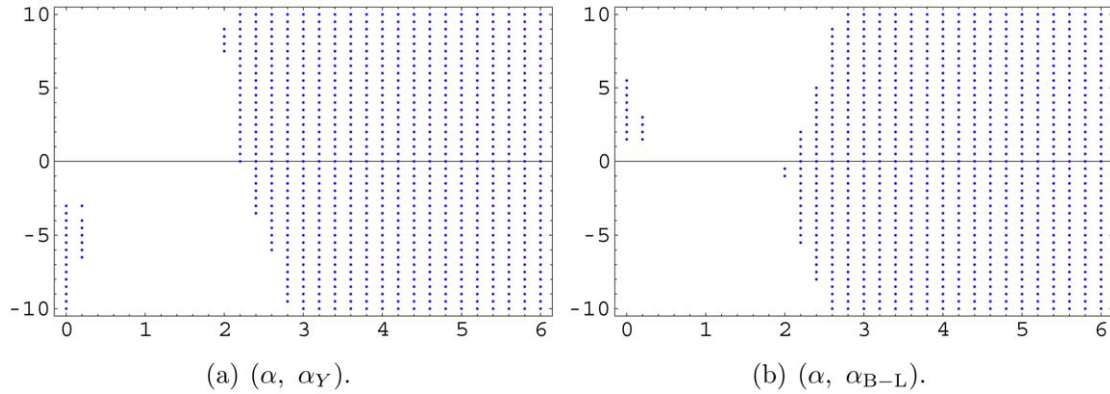


Fig. 1. The allowed parameter set which provides all the sfermion squared masses positive and the correct electroweak symmetry breaking in the case of  $\tan\beta = 10$  and  $M = 500$  GeV.

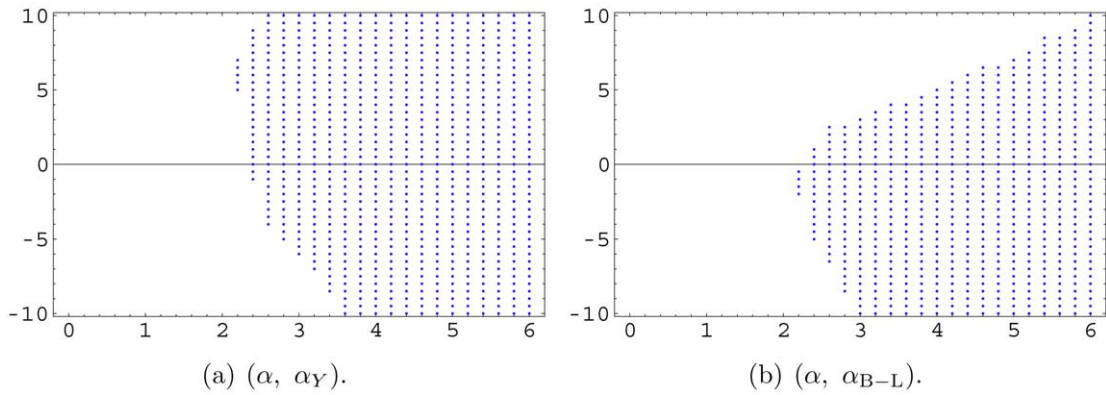


Fig. 2. The allowed parameter set in the case of  $\tan\beta = 45$  and  $M = 500$  GeV.

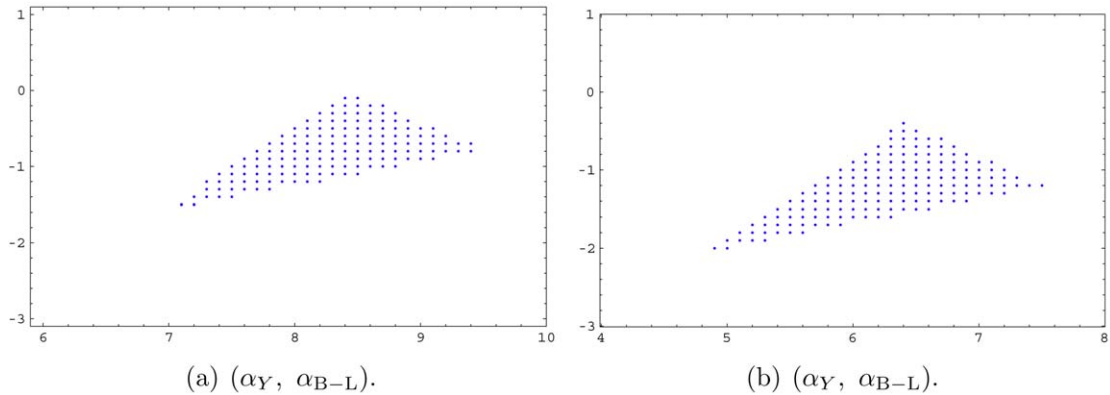


Fig. 3. The allowed parameter set which provides all the sfermion squared masses positive and the correct electroweak symmetry breaking in the case of  $\alpha = 2$ ,  $\tan\beta = 10$  and  $M = 500$  GeV, and (b)  $\alpha = 2.2$ ,  $\tan\beta = 45$  and  $M = 500$  GeV.

Higgs doublet is likely to be  $m_{H_1} \lesssim m_{H_2} < 0$ , so that it becomes difficult to achieve the correct electroweak symmetry breaking.

We have performed the same analysis in the case of the conventional mixed modulus–anomaly mediation ( $\alpha_Y = \alpha_{B-L} = 0$ ), and found that the allowed region is constrained to be  $\alpha \gtrsim 2.5$  for  $\tan\beta = 10$  and  $\alpha \gtrsim 3.2$  for  $\tan\beta = 45$ . Thus, Figs. 1 and 2 show that the allowed region of  $\alpha$  is widened in the presence of  $D$ -term contributions. In order to explicitly show this fact, we present the allowed parameter sets of  $(\alpha_Y, \alpha_{B-L})$  for fixed  $\alpha$  in Fig. 3. The point,  $\alpha_Y = \alpha_{B-L} = 0$ , corresponding to

the conventional mixed modulus–anomaly mediation is not allowed.

In Tables 1 and 2, we show some example data of the resultant sparticle mass spectrum and Higgs boson masses for  $\text{sgn}(\mu) > 0$ . Here, the standard model-like Higgs boson mass is evaluated by including one-loop corrections through top and scalar top quarks,

$$\Delta m_h^2 = \frac{3}{4\pi^2} y_t^4 v^2 \sin^4 \beta \ln \left( \frac{m_{\tilde{t}_1} m_{\tilde{t}_2}}{m_t^2} \right), \quad (28)$$

Table 1  
Sparticle and Higgs boson mass spectra (in units of GeV) in the case of  $\tan\beta = 10$

$M$ [GeV]	250	110	110	120
$\alpha$	0	5	5	5
$(\alpha_Y, \alpha_{B-L})$	(-6, 3)	(0, 0)	(12.5, 12.5)	(-12, 6)
$m_{\tilde{\chi}_{1,2,3,4}^0}$	99.0, 351, 568, 576	393, 491, 712, 730	393, 491, 714, 731	429, 535, 740, 759
$m_{\tilde{\chi}_{1,2}^\pm}$	99.0, 575	490, 729	490, 730	534, 758
$m_{\tilde{g}}$	902	940	940	1030
$m_{\tilde{e}, \tilde{\mu}_{1,2}}$	178, 219	575, 618	419, 732	646, 697
$m_{\tilde{\tau}_{1,2}}$	161, 225	556, 614	397, 726	628, 692
$m_{\tilde{u}, \tilde{c}_{1,2}}$	784, 826	911, 943	915, 940	971, 1020
$m_{\tilde{t}_{1,2}}$	569, 774	429, 834	464, 816	434, 891
$m_{\tilde{d}, \tilde{s}_{1,2}}$	829, 930	907, 946	918, 934	1020, 1030
$m_{\tilde{b}_{1,2}}$	733, 918	763, 891	729, 912	826, 1000
$m_h$	116	114	115	115
$m_H$	373	911	913	936
$m_A$	373	911	913	936
$m_{H^\pm}$	382	914	917	939

Table 2  
Sparticle and Higgs boson mass spectra (in units of GeV) in the case of  $\tan\beta = 45$

$M$ [GeV]	450	280	110	110
$\alpha$	2.5	3	5	5
$(\alpha_Y, \alpha_{B-L})$	(-3, -4.5)	(-3, -4.5)	(0, 0)	(6, -6)
$m_{\tilde{\chi}_{1,2,3,4}^0}$	174, 181, 1130, 1135	480, 499, 714, 775	392, 486, 633, 656	393, 489, 657, 677
$m_{\tilde{\chi}_{1,2}^\pm}$	178, 1130	486, 774	485, 655	488, 676
$m_{\tilde{g}}$	1110	958	940	940
$m_{\tilde{e}, \tilde{\mu}_{1,2}}$	637, 1440	614, 947	575, 618	583, 585
$m_{\tilde{\tau}_{1,2}}$	263, 118	420, 716	243, 516	254, 480
$m_{\tilde{u}, \tilde{c}_{1,2}}$	1000, 1230	902, 1010	911, 943	920, 950
$m_{\tilde{t}_{1,2}}$	343, 948	387, 787	473, 749	489, 759
$m_{\tilde{d}, \tilde{s}_{1,2}}$	948, 1230	879, 1020	907, 946	882, 953
$m_{\tilde{b}_{1,2}}$	504, 888	531, 723	570, 709	546, 706
$m_h$	115	114	115	116
$m_H$	203	154	112	272
$m_A$	203	154	112	272
$m_{H^\pm}$	219	173	138	284

which is important to push up the Higgs boson mass so as to satisfy the LEP II experimental bound,  $m_h \gtrsim 114$  GeV. As can be understood from the RGEs and the soft SUSY breaking parameters at the GUT scale presented in the previous section, the resultant soft SUSY breaking parameters are proportional to  $M$ . Thus, as we take  $M$  larger with fixed  $(\alpha, \alpha_Y, \alpha_{B-L})$ , sparticles become heavier and, accordingly, Higgs boson masses become larger. In the first column in Table 1, the LSP neutralino is wino-like as the same as in the pure AMSB scenario, while bino-like in the other columns. In the last two columns, the LSP neutralino well degenerates with the next LSP sfermion. Depending on values of  $\alpha_Y$  and  $\alpha_{B-L}$ , stau or stop can be the next LSP. This shows remarkable effects due to the  $D$ -term contributions.

The LSP neutralino is a good candidate for the dark matter in cosmology [17]. For small  $\tan\beta$ , if the LSP neutralino is

bino-like, its annihilation processes are dominated by p-wave, and are not so effective that the neutralino relic density tends to over-close the present universe. If the LSP neutralino well degenerates with the next LSP sfermions, its co-annihilation process with the next LSP plays an important role to make neutralino annihilation processes effective. Our results show that this case is possible due to the effects of the  $D$ -term contributions on the sfermion masses.

In the first two columns in Table 2, the lightest neutralino is higgsino-like, while bino-like in the last two columns. For  $\alpha \gtrsim 3$ , we found that the LSP is stau and this region is cosmologically disfavored. Light Higgs boson masses shown in the table indicate that, in the case of large  $\tan\beta$  and small  $\alpha$ , it is difficult to achieve the correct electroweak symmetry breaking.

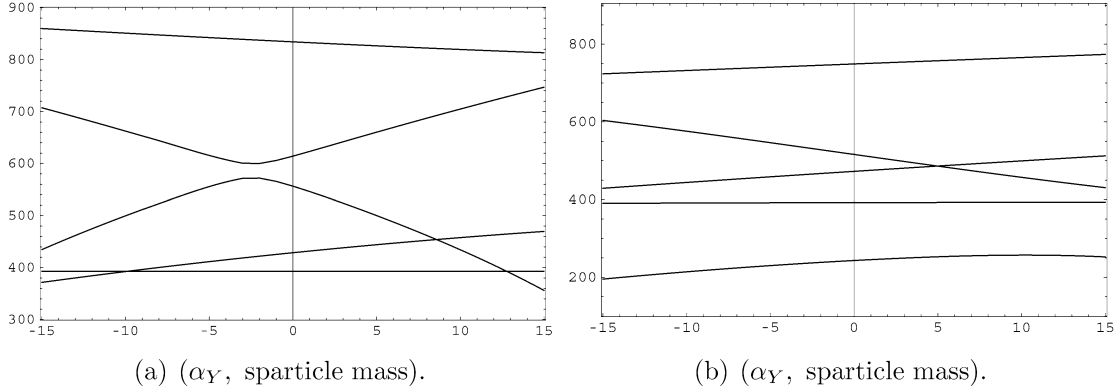


Fig. 4. Sparticle mass spectrum (in units of GeV) as a function of  $\alpha_Y$ . The figure (a) shows the result in the case of  $\alpha = 5$ ,  $\tan \beta = 10$ ,  $M = 110$  GeV and  $\alpha_{B-L}$  fixed to be  $\alpha_{B-L} = \alpha_Y$ . Each plot corresponds to  $m_{\tilde{\tau}_2}$ ,  $m_{\tilde{\tau}_1}$ ,  $m_{\tilde{\chi}_1^0}$  and  $m_{\tilde{\tau}_1}$ , respectively, from top to bottom at  $\alpha_Y = 0$ . In the figure (b),  $\alpha = 5$ ,  $\tan \beta = 45$ ,  $M = 110$  GeV and  $\alpha_{B-L}$  has been taken to be  $\alpha_{B-L} = -\alpha_Y$ . Each plot corresponds to  $m_{\tilde{\tau}_2}$ ,  $m_{\tilde{\tau}_1}$ ,  $m_{\tilde{\chi}_1^0}$  and  $m_{\tilde{\tau}_1}$ , respectively, from top to bottom at  $\alpha_Y = 0$ .

Finally we show sparticle mass spectrum as a function of the parameters of the  $D$ -terms. Fig. 4(a) shows several sparticle masses as a function of  $\alpha_Y$  in the case of  $\alpha = 5$ ,  $\tan \beta = 10$ ,  $M = 110$  GeV and  $\alpha_{B-L}$  fixed to be  $\alpha_{B-L} = \alpha_Y$ . This figure includes the sparticle masses presented in the second and third columns in Table 1. The point of  $\alpha_Y = 0$  corresponds to the sparticle mass spectrum in the conventional mixed modulus–anomaly mediation. We can see that the  $D$ -term contributions dramatically change the resultant sparticle masses, in particular slepton masses, from those in the conventional mixed modulus–anomaly mediation. For  $\alpha_Y < -2.5$ , lighter stau is mostly left-handed stau, while mostly right-handed stau for  $\alpha_Y > -2.5$ . As discussed above, there exists the parameter region where lighter stau or stop degenerates with the LSP neutralino. The case of  $\alpha = 5$ ,  $\tan \beta = 45$ ,  $M = 110$  GeV and  $\alpha_{B-L}$  fixed as  $\alpha_{B-L} = -\alpha_Y$  is depicted in Fig. 4(b). This includes the results in the last two columns in Table 2. Sparticle masses moderately depend on  $\alpha_Y$  in this case.

## 5. Summary and discussion

We have extended the mixed modulus–anomaly mediation so as to include  $D$ -term contributions from  $U(1)_Y$  and  $U(1)_{B-L}$ . Such  $D$ -term contributions can generically be expected when we consider some grand unified theory based on a higher rank gauge group such as  $SO(10)$ . We have evaluated soft SUSY breaking terms and obtained various sparticle mass spectra for various input values of  $(\alpha, \alpha_Y, \alpha_{B-L})$ , that are different from those obtained in the conventional mixed modulus–anomaly mediation. Especially, we have found that the LSP neutralino can be various types such as wino-like, higgsino-like and bino-like. In addition, stau or stop can be the next LSP with degenerate masses with bino-like neutralino due to the  $D$ -term contributions. This indicates that the co-annihilation channel can be opened up, when we consider the dark matter physics for the bino-like LSP neutralino. Evaluating the dark matter relic density in our scenario is an interesting subject. We leave this for future works.

Non-zero  $D$ -term of  $U(1)_{B-L}$  has further phenomenological importance, that is, new flavor violating effects can be generated through it. In the presence of the  $D$ -term of  $U(1)_{B-L}$ , off-diagonal elements of the slepton mass squared matrix can be generated [18],

$$(\Delta m_{\tilde{\ell}}^2)_{ij} = \frac{1}{8\pi^2} (Y_v^\dagger Y_v)_{ij} D_{B-L}, \quad (29)$$

where  $Y_v$  is the neutrino Dirac Yukawa coupling matrix. Depending on the Yukawa coupling matrix and the value of  $D_{B-L}$ , the lepton flavor violating (LFV) processes may be sizable. The off-diagonal elements can also be induced by RGEs through the neutrino Yukawa coupling as usually discussed [19]. Taking these contributions all together, analyzing the LFV processes in our scenario is worth investigating. In this analysis, concrete information about neutrino Yukawa coupling is necessary [20].

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