\textbf{B – L assisted Anomaly Mediation and the radiative B – L symmetry breaking}

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\textbf{Abstract}

Anomaly mediated supersymmetry breaking implemented in the minimal supersymmetric standard model (MSSM) is known to suffer from the tachyonic slepton problem leading to breakdown of electric charge conservation. We show however that when MSSM is extended to explain small neutrino masses by gauging the B-L symmetry, the slepton masses can be positive due to the $Z'$ mediation contributions. We obtain various soft supersymmetry breaking mass spectra, which are different from those obtained in the conventional anomaly mediation scenario. Then there would be a distinct signature of this scenario at the LHC.
1 Introduction

Supersymmetry (SUSY) extension is one of the most promising way to solve the gauge hierarchy problem in the standard model (SM) \[1\]. 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 not to induce too large CP and flavor violation effects. Some mechanisms to achieve such SUSY breaking mediation have been proposed \[2\].

The anomaly mediated supersymmetry breaking (AMSB) scenario \[3, 4, 5\] is one of the most attractive scenario due to its flavor-blindness and ultraviolet (UV) insensitivity for the resultant soft SUSY breaking terms. The pattern of SUSY breaking does not depend at all on physics at higher energy scales. On the eve of the Large Hadron Collider (LHC) operation at CERN, which start this year, there are several studies in the aspects of collider physics to discriminate the AMSB scenario from the other SUSY breaking mediation scenarios \[6, 7, 8\]. Despite the appeal of the AMSB, the original version of the AMSB is excluded because of its high predictivity. The slepton squared masses become negative at the weak scale, and hence the theory would break $U(1)_{em}$. There have been many attempts to solve this problem by incorporating additional positive contributions to slepton squared masses at tree level \[3, 9, 10, 11\] or at quantum level \[12, 13\].

An important thing to realize at this point is that MSSM is not a complete theory of low energy particle physics and needs extension to explain the small neutrino masses observed in experiments. The relevant question then is whether MSSM extended to include new physics that explains small neutrino masses will cure the tachyonic slepton mass pathology of AMSB. One of the simplest extensions of MSSM which provide natural explanation of small neutrino masses is to extend the gauge symmetry of MSSM to $SU(3)_c \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$ which naturally introduces three right-handed neutrinos into the theory in order for the anomaly cancellation. Once we incorporate the $U(1)_{B-L}$ gauge symmetry in SUSY models, the $U(1)_{B-L}$ gaugino $\tilde{Z}_{B-L}$ appears, and it can mediate the SUSY breaking \[14\] (the $Z'$ mediated SUSY breaking \[15, 16\]).

This paper is organized as follows. In Sec. 2 and Sec. 3, we give a brief review of the anomaly mediation and the $Z'$ mediated SUSY breaking, respectively. In Sec. 4, we combine these two scenarios and examine the numerical evaluations to give a sample mass spectra. Sec. 5 is devoted to summary and discussion.

\[3\] The similar idea has also been suggested in \[17, 18\].
2 Anomaly Mediation in the $B-L$ extended MSSM

In this section, we work out in the superconformal framework of supergravity [19], and we explain the anomaly mediation scenario in the $B-L$ extended MSSM.

In the superconformal framework of supergravity, the basic Lagrangian is given by

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

where $\phi = 1 + \theta^2 F_{\phi}$ is the compensating multiplet, $K$ is the Kähler potential in the conformal frame, and $W$ is the superpotential.

As for the gauge sector in the MSSM, the kinetic term is of the form,

$$ L_{\text{gauge}} = \frac{1}{4} \int d^2 \theta \tau_a \left( \frac{\mu_R}{\Lambda_{\phi}} \right) W^a_{\alpha} W^a_{\alpha}. \quad (2) $$

At the classical level, the compensator $\phi$ does not appear in the gauge kinetic term as the gauge chiral superfield $W^a_{\alpha}$ has a chiral weight $\frac{3}{2}$. It turns out that the dependence of $\phi$ comes out radiatively through the cutoff scale $\Lambda$ ($\mu_R$ is the renormalization scale). In the above setup, non-zero $F_{\phi}$ induces soft SUSY breaking terms through the AMSB, and the resultant SUSY breaking mass scale is characterized by $m_{\text{AMSB}} \sim F_{\phi}/(16\pi^2)$. Considering the anomaly mediation contribution to the soft scalar masses and A-terms, 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 is described as

$$ L_{\text{kin}} = \int d^4 \theta \phi^\dagger \phi Q_i^\dagger e^{2g_a V_a} Q_i + \cdots. \quad (3) $$

As discussed in Ref. [20], in softly broken supersymmetry, the soft terms associated to a chiral superfield $Q_i$ can be collected in a running superfield wave function $Z_i(\mu_R)$ such that

$$ \ln Z_i(\mu_R) = \ln Z_i(\mu_R) + [A_i(\mu_R) \theta^2 + h.c.] - \tilde{m}_i^2(\mu_R) \theta^4. \quad (4) $$

The running wave functions can be defined as $Z_i(\mu_R) = c_i(p^2 = -\mu_R^2)$, where $c_i$ is the coefficient of $Q_i^\dagger Q_i$ in the one point-irreducible (1PI) effective action. Therefore, turning on superconformal anomaly amounts to the shift $\mu_R \rightarrow \mu_R/(\phi^\dagger \phi)^{1/2}$.

$$ Z_i(\mu_R) = Z_i \left( \frac{\mu_R}{(\phi^\dagger \phi)^{1/2}} \right). \quad (5) $$

According to the method developed in Ref. [20] (see also Ref. [12]), soft SUSY breaking terms (each gaugino masses $M_a$, sfermion squared masses $\tilde{m}_i^2$ and $A$-parameters) at the scale
\( \mu_R \) can be extracted from renormalized gauge kinetic functions and SUSY wave function renormalization coefficients,

\[
M_a(\mu_R) = \frac{1}{16\pi^2} b_a g_a^2(\mu_R) F_\phi , \\
\tilde{m}_i^2(\mu_R) = \frac{1}{2} \frac{d \gamma_i(\mu_R)}{d \ln(\mu_R)} |F_\phi|^2 , \\
A_{ijk}(\mu_R) = -[\gamma_i(\mu_R) + \gamma_j(\mu_R) + \gamma_k(\mu_R)] F_\phi .
\]

(6)

Here, \( g_a \) are the gauge couplings, \( b_a \) are beta function coefficients, and \( \gamma_i \equiv -(1/2) d \ln Z/d \ln \mu \) are anomalous dimensions of the matter and Higgs superfields. All the soft mass parameters can be described by only one parameter, \( F_\phi \), so the anomaly mediation is highly predictive.

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. Although some fine-tuning among parameters is necessary to realize \( \mu \sim B \sim M_Z \), in the following analysis we treat them as free parameters so that the value of \( |\mu| \) and \( B \mu \) are determined by the stationary condition of the Higgs potential.

Let us consider the following superpotential,

\[
W = -(Y_U)_{ij} H_2 Q_i U^c_j + (Y_D)_{ij} H_1 Q_i D^c_j - (Y_L)_{ij} H_2 L_i N^c_j + (Y_E)_{ij} H_1 L_i E^c_j \\
- \mu H_1 H_2 - \mu' \Delta_1 \Delta_2 + \frac{1}{2} f_{ij} \Delta_1 N^c_i N^c_j ,
\]

(7)

where \( \Delta_1 \) and \( \Delta_2 \) have \( B - L \) charge \(-2\) and \(+2\) respectively. Neglecting Yukawa couplings for first two generations, anomalous dimensions are given by

\[
16\pi^2 \gamma_{Q_i} = -\frac{8}{3} g_3^2 - \frac{3}{2} g_2^2 - \frac{1}{18} g_Y^2 - \frac{2}{9} g_{B-L}^2 + (y_i^2 + y_b^2) \delta_{i3} , \\
16\pi^2 \gamma_{U^c_i} = -\frac{8}{3} g_3^2 - \frac{8}{9} g_Y^2 - \frac{2}{9} g_{B-L}^2 + 2y_i^2 \delta_{i3} , \\
16\pi^2 \gamma_{D^c_i} = -\frac{8}{3} g_3^2 - \frac{2}{9} g_Y^2 - \frac{2}{9} g_{B-L}^2 + 2y_b^2 \delta_{i3} , \\
16\pi^2 \gamma_{L_i} = -\frac{3}{2} g_2^2 - \frac{1}{2} g_Y^2 - 2g_{B-L}^2 + (y_i^2 + y_e^2) \delta_{i3} , \\
16\pi^2 \gamma_{N^c_i} = -2g_{B-L}^2 + f^2 + 2y_i^2 \delta_{i3} , \\
16\pi^2 \gamma_{E^c_i} = -2g_Y^2 - 2g_{B-L}^2 + 2y_e^2 \delta_{i3} , \\
16\pi^2 \gamma_{H_1} = -\frac{3}{2} g_2^2 - \frac{1}{2} g_Y^2 + 3y_b^2 + y_e^2 , \\
16\pi^2 \gamma_{H_2} = -\frac{3}{2} g_2^2 - \frac{1}{2} g_Y^2 + 3y_i^2 + y_e^2 , \\
16\pi^2 \gamma_{\Delta_1} = -8g_{B-L}^2 + f^2 , \\
16\pi^2 \gamma_{\Delta_2} = -8g_{B-L}^2 .
\]

(8)
The soft scalar masses are explicitly written as

\[
\begin{align*}
    m_{\tilde{q}}^2 &= \left( \frac{F_\phi}{16\pi^2} \right)^2 \left[ 8g_3^4 - \frac{3}{2}g_2^4 - \frac{11}{18}g_Y^4 - \frac{16}{3}g_{B-L}^4 + (y_t^2b_{yt} + y_b^2b_{yb})\delta_{i3} \right], \\
    m_{\tilde{u}}^2 &= \left( \frac{F_\phi}{16\pi^2} \right)^2 \left[ 8g_3^4 - \frac{88}{9}g_Y^4 - \frac{16}{3}g_{B-L}^4 + 2y_t^2b_{yt}\delta_{i3} \right], \\
    m_{\tilde{d}}^2 &= \left( \frac{F_\phi}{16\pi^2} \right)^2 \left[ 8g_3^4 - \frac{22}{9}g_Y^4 - \frac{16}{3}g_{B-L}^4 + 2y_b^2b_{yb}\delta_{i3} \right], \\
    m_{\tilde{e}}^2 &= \left( \frac{F_\phi}{16\pi^2} \right)^2 \left[ -\frac{3}{2}g_2^4 - \frac{11}{2}g_Y^4 - 48g_{B-L}^4 + (y_e^2b_{ye} + y^2b_{ye})\delta_{i3} \right], \\
    m_{\tilde{\nu}}^2 &= \left( \frac{F_\phi}{16\pi^2} \right)^2 \left[ -48g_{B-L}^4 + f^2b_f + 2y_e^2b_{ye}\delta_{i3} \right], \\
    m_{\tilde{\tau}}^2 &= \left( \frac{F_\phi}{16\pi^2} \right)^2 \left[ -22g_Y^4 - 48g_{B-L}^4 + 2y_{\tau}^2b_{\tau}\delta_{i3} \right], \\
    m_{\tilde{\Delta}_1}^2 &= \left( \frac{F_\phi}{16\pi^2} \right)^2 \left[ -192g_{B-L}^4 + f^2b_f \right], \\
    m_{\tilde{\Delta}_2}^2 &= \left( \frac{F_\phi}{16\pi^2} \right)^2 \left[ -192g_{B-L}^4 \right].
\end{align*}
\]

where \(b_{yt}, b_{yb}, b_{ye}, b_{yr},\) and \(b_f\) are given by

\[
\begin{align*}
    b_{yt} &= 6y_t^2 + y_b^2 + y_Y^2 - \frac{16}{3}g_3^2 - 3g_2^2 - \frac{13}{9}g_Y^2 - \frac{4}{9}g_{B-L}^2, \\
    b_{yb} &= y_t^2 + 6y_b^2 + y_Y^2 - \frac{16}{3}g_3^2 - 3g_2^2 - \frac{7}{9}g_Y^2 - \frac{4}{9}g_{B-L}^2, \\
    b_{ye} &= 3y_Y^2 + 4y_e^2 + y_{\tau}^2 + f^2 - 3g_2^2 - g_Y^2 - 4g_{B-L}^2, \\
    b_{yr} &= 3y_b^2 + 4y_{\tau}^2 + y_Y^2 - 3g_2^2 - 3g_Y^2 - 4g_{B-L}^2, \\
    b_f &= 4y_Y^2 + 3f^2 - 12g_{B-L}^2.
\end{align*}
\]

Also, the Higgs soft masses are given by

\[
\begin{align*}
    m_{H_1}^2 &= \left( \frac{F_\phi}{16\pi^2} \right)^2 \left[ -\frac{3}{2}g_2^4 - \frac{11}{2}g_Y^4 + 3y_b^2b_{yb} + y_{\tau}^2b_{yr} \right], \\
    m_{H_2}^2 &= \left( \frac{F_\phi}{16\pi^2} \right)^2 \left[ -\frac{3}{2}g_2^4 - \frac{11}{2}g_Y^4 + 3y_t^2b_{yt} \right].
\end{align*}
\]

The Higgs mass parameters, \(\mu\)-term and \(B\mu\)-term, are determined by the electroweak sym-
metry breaking conditions,

\[ |\mu|^2 = \frac{m_{H_1}^2 - m_{H_2}^2 \tan^2 \beta}{\tan^2 \beta - 1} - \frac{1}{2} M_Z^2, \]

\[ B_\mu = \frac{1}{2} [m_{H_1}^2 + m_{H_2}^2 + 2|\mu|^2] \sin 2\beta. \tag{12} \]

The $A$-parameters in the AMSB scenario are given by

\[ A_{ijk} = - (\gamma_i + \gamma_j + \gamma_k) F_\phi \tag{13} \]

with the above anomalous dimensions. Finally, the gaugino masses are given by

\[ M_{B-L} = 24 g_{B-L}^2 \left( \frac{F_\phi}{16\pi^2} \right), \]

\[ M_1 = 11 g_Y^2 \left( \frac{F_\phi}{16\pi^2} \right), \]

\[ M_2 = g_2^2 \left( \frac{F_\phi}{16\pi^2} \right), \]

\[ M_3 = -3 g_3^2 \left( \frac{F_\phi}{16\pi^2} \right). \tag{14} \]

The mass ratios are approximately $M_{B-L} : M_1 : M_2 : M_3 = 57 g_{B-L}^2 : 3 : 1 : 10$. So the Wino (rather than the more conventional Bino) is the lightest SUSY particle (LSP), and the gluino is an order of magnitude heavier than the LSP. Those predictions for the gaugino masses in the AMSB, that is, a Wino-like LSP, has interesting phenomenological consequences. The remarkable fact is that the lightest chargino mass is nearly degenerated with the lightest neutralino mass.

3 Contributions from the Z-prime mediation

Here we give a brief review of the Z-prime mediation of SUSY breaking \[15,16\] by discussing the pattern of the soft SUSY breaking parameters, the masses of the $Z'$-ino and of the MSSM squarks and gauginos, which are the most robust predictions of this scenario. At the SUSY breaking scale, $\Lambda_S$, SUSY breaking in the hidden sector is assumed to generate a SUSY breaking mass for the fermionic component of the $U(1)_{B-L}$ vector superfield. Given details of the hidden sector, its value could be evaluated via the standard technique of analytical continuation into superspace \[21\]. In particular, the gauge kinetic function of the field strength superfield $W_{B-L}^{\alpha}$ at the SUSY breaking scale is

\[ \mathcal{L}_{B-L} = \int d^2 \theta \left[ \frac{1}{g_{B-L}^2} + \beta_{B-L}^{\text{hid}} \ln \left( \frac{\Lambda_S}{M} \right) + \beta_{B-L}^{\text{vis}} \ln \left( \frac{\Lambda_S}{M_{Z_{B-L}}} \right) \right] W_{B-L}^{\alpha} W_{B-L}^{\alpha}, \tag{15} \]
where $M$ is the messenger scale, which we have assumed to be around the SUSY breaking scale, $M \sim \Lambda_S$. $\beta^{\text{hid}}_{B-L}$ and $\beta^{\text{vis}}_{B-L}$ are $\beta$-functions induced by U(1)$_{B-L}$ couplings to hidden and visible sector fields, respectively. Using analytical continuation, we replace $M$ with $M + \theta^2 F$, where $F$ is the SUSY breaking order parameter. We obtain the $\tilde{Z}_{B-L}$ mass as $M_{\tilde{Z}_{B-L}} \sim g^2_{B-L} \beta^{\text{hid}}_{B-L} F/M$. We assume that the U(1)$_{B-L}$ gauge symmetry is not broken in the hidden sector. And we assume some sequestering mechanism so that only the $B-L$ gaugino obtains a leading order mass term while the threshold corrections to the squarks and sleptons are only arisen at the next leading order as similar to the case of the gaugino mediation, where the $B-L$ gaugino lives in the bulk in a five dimensional setup while squarks and sleptons are put on the brane. In such a case, only the $B-L$ gaugino obtains a mass while the scalar masses receive negligible threshold corrections at the lowest order since they receive volume suppression.

Since all the chiral superfields in the visible sector are charged under U(1)$_{B-L}$, so all the corresponding scalars receive soft mass terms at 1-loop of order

$$m^2_{\tilde{q}_i} = \frac{8}{9} \frac{\alpha_{B-L}}{4\pi} M^2_{\tilde{Z}_{B-L}} \ln \left( \frac{\Lambda_S}{M_{\tilde{Z}_{B-L}}} \right),$$

$$m^2_{\tilde{\ell}_i} = \frac{8}{4\pi} \frac{\alpha_{B-L}}{4\pi} M^2_{\tilde{Z}_{B-L}} \ln \left( \frac{\Lambda_S}{M_{\tilde{Z}_{B-L}}} \right),$$

where $\alpha_{B-L} = g^2_{B-L}/(4\pi)$ and $Q^f_{B-L}$ is the U(1)$_{B-L}$ charge of $f$.

The MSSM gaugino masses, however, can only be generated at 2-loop level since they do not directly couple to the U(1)$_{B-L}$,

$$M_a = 4c_a \frac{\alpha_{B-L}}{4\pi} \frac{\alpha_a}{4\pi} M_{\tilde{Z}_{B-L}} \ln \left( \frac{\Lambda_S}{M_{\tilde{Z}_{B-L}}} \right),$$

where $(c_1, c_2, c_3) = (\frac{92}{15}, 4, \frac{4}{3})$.

From the discussion above, we see that the gauginos are considerably lighter than the sfermions. Taking $m_f \simeq 100 - 1000$ GeV, we find

$$M_{\tilde{Z}_{B-L}} \simeq 10^4 \text{ GeV}$$

and then the $Z'$ mediated contribution is well-suppressed:

$$M_a \simeq 10^{-4} M_{\tilde{Z}_{B-L}} \simeq 1 \text{ GeV},$$

which can be negligible compared to the contributions from anomaly mediation.
4 RGEs and its numerical evaluations

Now we consider the RGEs and analyze the running of the scalar masses $m^2_{\Delta_1}$ and $m^2_{\Delta_2}$. The key point for implementing the radiative $B-L$ symmetry breaking is that the scalar potential $V(\Delta_1, \Delta_2)$ receives substantial radiative corrections [22, 14]. In particular, a negative (mass)$^2$ would trigger the $B-L$ symmetry breaking. We argue that the masses of Higgs fields $\Delta_1$ and $\Delta_2$ run differently in the way that $m^2_{\Delta_1}$ can be negative whereas $m^2_{\Delta_2}$ remains positive.

The RGE for the $B-L$ coupling and mass parameters can be derived from the general results for SUSY RGEs of Ref. [23].

For the RGEs of the Yukawa couplings, we consider to include the additional contribution from the the $U(1)_{B-L}$ gauge sector.

\[ 16\pi^2 \frac{dy_A}{d\ln \mu} = b_A y_A , \quad (20) \]

where $A = (t, b, \nu, \tau, f)$, and $b_A$ is shown in the section 1. The RGEs of the MSSM gauge couplings are the same as MSSM, while the RGE of the $U(1)_{B-L}$ gauge coupling is given by

\[ 16\pi^2 \frac{dg_{B-L}}{d\ln \mu} = b_{B-L} g^3_{B-L} , \quad (21) \]

where $b_{B-L} = 24$. For the RGEs of the gaugino masses, it can be written as follows.

\[ 16\pi^2 \frac{dM_{\tilde{Z}_{B-L}}}{d\ln \mu} = 2b_{B-L} g^3_{B-L} M_{\tilde{Z}_{B-L}} , \quad (22) \]

where $(c_a) = (92/15, 4, 4/3)$. For the RGEs of the A-terms, it can be written as follows.

\[ 16\pi^2 \frac{dA}{d\ln \mu} = [\text{MSSM + see-saw}] - 2a_A g^2_{B-L} (\tilde{A}_A - 2 M_{\tilde{Z}_{B-L}} Y_A) , \quad (23) \]

where $\tilde{A}_A = A_A Y_A$ with $A = (t, b, \nu, \tau)$ and $(a_t, a_b, a_\nu, a_\tau) = (2, 2, 2, 2)$. The RGE of the $A_f$-term can be written as

\[ 16\pi^2 \frac{d\tilde{A}_f}{d\ln \mu} = (9 \text{Tr}[f^\dagger f] + 2 \text{Tr}[Y_\nu^\dagger Y_\nu]) \tilde{A}_f + 8 f Y_\nu^\dagger \tilde{A}_f . \quad (24) \]

The RGEs of the soft scalar masses are given by

\[ 16\pi^2 \frac{dm^2_{\Delta_1}}{d\mu} = 2 \text{Tr}[f^\dagger f] m^2_{\Delta_1} + 4 \text{Tr}[f^\dagger m^2_N f] - 32 g^2_{B-L} |M_{\tilde{Z}_{B-L}}|^2 . \]
\[ 16\pi^2 \frac{dm^2_{\Delta_2}}{d\mu} = -32 g^2_{B-L} |M_{\tilde{Z}_{B-L}}|^2 . \]
\[ 16\pi^2 \frac{dm^2_{f}}{d\mu} = [\text{MSSM + see-saw}] - 8 g^2_{B-L} (Q_B^{f})^2 |M_{\tilde{Z}_{B-L}}|^2 . \quad (25) \]
where $Q_{B-L}^f$ is the $B-L$ charge of each chiral multiplet $f = Q, U^c, D^c, L, N^c$. For the RGEs of the $\mu$-term, it can be written as follows.

$$16\pi^2 \mu \frac{d}{d\mu} \mu' = (\text{Tr}[f^\dagger f] - 16g_{B-L}^2)\mu' .$$

(26)

In the numerical analysis, we fix $F_\phi$ to $10^5$ GeV for simplicity. So we have only three free parameters,

$$g_{B-L} , f , M_{Z_{B-L}} .$$

(27)

Once we fix $g_{B-L}, f$ and $M_{Z_{B-L}}$ at the SUSY breaking scale $\Lambda = \sqrt{F_\phi M_{pl}} \approx 10^{11}$ GeV, all the soft SUSY breaking parameters due to AMSB and $Z'$ mediation at $\Lambda$ are also fixed, and RGE evolutions provide us with informations at low scale.

Fig. 1 shows the evolutions of the soft mass for the field $\Delta_1$. In Fig. 1 from top to the bottom curves, we varied the value of $f$ as $f = 1.5, 2.5, 3.5$ with $F_\phi = 100$ TeV, $g_{B-L} = 0.1$ and $M_{Z_{B-L}} = 5$ TeV.

For example, for the case of $f = 2.5$, the soft mass squared for the fields $\Delta_1$ goes across the zeros toward negative value, that is nothing but the realization of the radiative symmetry breaking of $U(1)_{B-L}$ gauge symmetry. The seesaw scale is found to be at $v_{B-L} = 10^4$ GeV. Hence the right-handed neutrinos obtain their masses of $M_N = f v_{B-L} \approx 10^4$ GeV. The running behavior in Fig. 1 can be understood in the following way. Starting from the high energy scale, the soft mass squared increases because of the gauge coupling contributions, and decrease of the mass squared is caused by the Yukawa coupling $f$ that dominate over the gauge coupling contribution.

Fig. 2 show the evolutions of the soft mass for sleptons, where the Yukawa coupling $f$ is fixed to 2.5, since their spectra are almost independent of the value of $f$. As seen in Fig. 2, the larger $M_{Z_{B-L}}$ gives the more positive slepton mass. This behavior is easily understood from Eq. (14), the RGE of the slepton. On the other hand, the larger $g_{B-L}$ gives the degenerate mass spectra. This is because, $m_{\tilde{\ell}_1}^2$ and $m_{\tilde{\ell}_2}^2$ at $\Lambda$ depend only on $g_{B-L}$ in the case of the large $g_{B-L}$. These degenerate mass spectra are one of the outstanding feature of this scenario.

In Table 1, we show some example data of the resultant sparticle mass spectrum and Higgs boson masses, where we took $\tan \beta = 10, F_\phi = 50$ TeV and $f = 2.5$. Here, the standard model-like Higgs boson mass is evaluated by including one-loop corrections through top and scalar top quarks,

$$\Delta m_h^2 = \frac{3}{4\pi^2} y_t^4 v^2 \sin^4 \beta \ln \left( \frac{m_{\ell_1} m_{\ell_2}}{m_t^2} \right) ,$$

(28)

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 presented in the previous section, the resultant soft SUSY breaking
parameters are proportional to $F_{\phi}$. Thus, as we take $F_{\phi}$ larger, sparticles become heavier
and, accordingly, Higgs boson masses become larger.

5 Dark matter relic density

In this section we discuss the cosmological features of the lightest neutralino. The recent
Wilkinson Microwave Anisotropy Probe (WMAP) satellite data [24] provide estimations of
various cosmological parameters with greater accuracy. The current density of the universe
is composed of about 73% of dark energy and 27% of matter. Most of the matter density is
in the form of the CDM, and its density is estimated to be [24]

$$\Omega_{\text{CDM}} h^2 = 0.1143 \pm 0.0034 .$$  \hspace{1cm} (29)

If the R-parity is conserved in SUSY models, the LSP is stable. The lightest neutralino, if
it is the LSP, is the plausible candidate for the CDM.

In the AMSB scenario or its extension with $Z'$ mediation, the lightest neutralino is
mostly Wino-like, and it undergoes rapid annihilation though reaction: $\tilde{W}\tilde{W} \rightarrow W^+ W^-$. The resultant relic abundance is too small, which can roughly be estimated to be [4]

$$\Omega_{\tilde{W}} h^2 \simeq 5 \times 10^{-4} \left( \frac{M_{\tilde{W}}}{100 \text{ GeV}} \right)^2 .$$  \hspace{1cm} (30)

So the mass of the DM neutralino has to be very heavy to satisfy the WMAP data. If
the Wino-like neutralino with SU(2)$_L$ charge is much heavier than the weak gauge boson as
described above, the weak interaction is a long-distance force for non-relativistic two-bodies
states of such particles. If this non-perturbative effect (namely, Sommerfeld enhancement)
of the dark matter at the freeze-out temperature is taken into account, the abundance can
be reduced by about 50% [25, 26]. Therefore, the allowed region exists for large value of $F_{\phi}$.

Such a large value of soft mass is disfavored in view of the little hierarchy problem. In
order to keep the neutralino DM light, non-thermal production of the DM should be
considered as proposed in [27]. Once we accept the non-thermal production of the LSP
neutralino from the moduli decays, then it is possible to produce sufficient relic abundance
of the LSP neutralino even for the light Wino-like neutralino DM.

6 Summary and discussion

Anomaly mediation of supersymmetry breaking (AMSB) is very attractive because the re-
sultant soft supersymmetry breaking parameters at a given energy scale are determined only
by physics at that energy scale (UV insensitivity) and hence is highly predictive (only one
parameter, $F_\phi$). However, there is the so-called tachyonic slepton problem. In this paper, we have constructed a viable anomaly mediation scenario of SUSY breaking by adding a contribution from the $Z'$ mediated SUSY breaking contributions. In the $Z'$ mediated SUSY breaking scenario, while the scalar masses are generated at the 1-loop level, however, gaugino masses can only be generated at 2-loop level, so the gaugino masses are completely determined by the pure anomaly mediation itself. Therefore, the characteristic signature of the present model predictions appear in the scalar partners mass spectra. We have investigated the scalar partners mass spectra for several choices of parameters in this model, for instance, for different values of the $Z'$ gaugino mass. The resultant sparticle mass spectra was found to be interesting in scope of the LHC.

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A  RGEs in the MSSM with right-handed neutrinos

A.1  The 2-loop RGE for the gauge couplings

\[
16\pi^2 \mu \frac{d}{d\mu} g_1 = \frac{33}{5} g_1^3 + \frac{g_1^3}{16\pi^2} \left( \frac{199}{25} g_1^2 + \frac{27}{5} g_2^2 + \frac{88}{5} g_3^2 \right),
\]

\[
16\pi^2 \mu \frac{d}{d\mu} g_2 = g_2^3 + \frac{g_2^3}{16\pi^2} \left( \frac{9}{5} g_1^2 + 25 g_2^2 + 24 g_3^2 \right),
\]

\[
16\pi^2 \mu \frac{d}{d\mu} g_3 = -3 g_3^3 + \frac{g_3^3}{16\pi^2} \left( \frac{1}{5} g_1^2 + 9 g_2^2 + 14 g_3^2 \right).
\]

Here $g_2 \equiv g$ is the $SU(2)_L$ gauge coupling constant and $g_1 \equiv \sqrt{\frac{5}{3}} g'$ is the $U(1)$ gauge coupling constant with the GUT normalization ($g_1 = g_2 = g_3$ at $\mu = M_{GUT}$).
A.2 The 1-loop RGE for the Yukawa couplings

\[
16\pi^2 \mu^2 \frac{d}{d\mu} Y_u = Y_u \left[ \left\{ -\frac{13}{15} g_1^2 - 3 g_2^2 - \frac{16}{3} g_3^2 + 3 \text{Tr}(Y_u^\dagger Y_u) + \text{Tr}(Y_\nu^\dagger Y_\nu) \right\} 1_{3\times3} \right.
\]
\[
+ 3 (Y_u^\dagger Y_u) + (Y_d^\dagger Y_d) \right], \quad (34)
\]

\[
16\pi^2 \mu^2 \frac{d}{d\mu} Y_d = Y_d \left[ \left\{ -\frac{7}{15} g_1^2 - 3 g_2^2 - \frac{16}{3} g_3^2 + 3 \text{Tr}(Y_d^\dagger Y_d) + \text{Tr}(Y_e^\dagger Y_e) \right\} 1_{3\times3} \right.
\]
\[
+ 3 (Y_d^\dagger Y_d) + (Y_u^\dagger Y_u) \right], \quad (35)
\]

\[
16\pi^2 \mu^2 \frac{d}{d\mu} Y_\nu = Y_\nu \left[ \left\{ -\frac{3}{5} g_1^2 - 3 g_2^2 + 3 \text{Tr}(Y_u^\dagger Y_u) + \text{Tr}(Y_\nu^\dagger Y_\nu) \right\} 1_{3\times3} \right.
\]
\[
+ 3 (Y_\nu^\dagger Y_\nu) + (Y_e^\dagger Y_e) \right], \quad (36)
\]

\[
16\pi^2 \mu^2 \frac{d}{d\mu} Y_e = Y_e \left[ \left\{ -\frac{9}{5} g_1^2 - 3 g_2^2 + 3 \text{Tr}(Y_d^\dagger Y_d) + \text{Tr}(Y_e^\dagger Y_e) \right\} 1_{3\times3} \right.
\]
\[
+ 3 (Y_e^\dagger Y_e) + (Y_\nu^\dagger Y_\nu) \right]. \quad (37)
\]

A.3 The 2-loop RGE for the gaugino masses

\[
16\pi^2 \mu^2 \frac{d}{d\mu} M_1 = \frac{66}{5} g_1^2 M_1
\]
\[
+ \frac{2 g_1^2}{16\pi^2} \left\{ \frac{199}{5} g_1^2 (2 M_1) + \frac{27}{5} g_2^2 (M_1 + M_2) + \frac{88}{5} g_3^2 (M_1 + M_3) \right\}, \quad (38)
\]

\[
16\pi^2 \mu^2 \frac{d}{d\mu} M_2 = 2 g_2^2 M_2
\]
\[
+ \frac{2 g_2^2}{16\pi^2} \left\{ \frac{9}{5} g_2^2 (M_1 + M_2) + 25 g_2^2 (2 M_2) + 24 g_3^2 (M_2 + M_3) \right\}, \quad (39)
\]

\[
16\pi^2 \mu^2 \frac{d}{d\mu} M_3 = -6 g_3^2 M_3
\]
\[
+ \frac{2 g_3^2}{16\pi^2} \left\{ \frac{11}{5} g_1^2 (M_1 + M_3) + 9 g_2^2 (M_2 + M_3) + 14 g_3^2 (2 M_3) \right\}. \quad (40)
\]
\[ \begin{align*}
16\pi^2 \mu \frac{d}{d\mu} (m_q^2)_{ij} &= \left( -\frac{2}{15} g_t^2 |M_1|^2 + \frac{6}{3} g_2^2 |M_2|^2 + \frac{32}{3} g_3^2 |M_3|^2 \right) \delta_{ij} + \frac{1}{5} g_t^2 S \delta_{ij} \\
& \quad + \left( m_q^2 Y_u^\dagger Y_u + m_q^2 Y_d^\dagger Y_d + Y_u^\dagger Y_u m_q^2 + Y_d^\dagger Y_d m_q^2 \right)_{ij} \\
& \quad + 2 \left( Y_u^\dagger m_q^2 Y_u + m_{H_u}^2 Y_u^\dagger Y_u + A_u^\dagger A_u \right)_{ij} \\
& \quad + 2 \left( Y_d^\dagger m_q^2 Y_d + m_{H_d}^2 Y_d^\dagger Y_d + A_d^\dagger A_d \right)_{ij}, \\
\end{align*} \]

\[ \begin{align*}
16\pi^2 \mu \frac{d}{d\mu} (m_\nu^2)_{ij} &= \left( \frac{32}{15} g_t^2 |M_1|^2 + \frac{32}{3} g_2^2 |M_2|^2 \right) \delta_{ij} - \frac{4}{5} g_t^2 S \delta_{ij} \\
& \quad + \left( m_\nu^2 Y_\nu^\dagger Y_\nu + Y_\nu^\dagger Y_\nu m_\nu^2 \right)_{ij} \\
& \quad + 4 \left( Y_\nu^\dagger m_\nu^2 Y_\nu + m_{H_\nu}^2 Y_\nu^\dagger Y_\nu + A_\nu A_\nu^\dagger \right)_{ij}, \\
\end{align*} \]

\[ \begin{align*}
16\pi^2 \mu \frac{d}{d\mu} (m_\ell^2)_{ij} &= \left( \frac{6}{5} g_t^2 |M_1|^2 + \frac{6}{5} g_2^2 |M_2|^2 \right) \delta_{ij} - \frac{3}{5} g_t^2 S \delta_{ij} \\
& \quad + \left( m_\ell^2 Y_\ell^\dagger Y_\ell + m_\ell^2 Y_\ell^\dagger Y_\ell m_\ell^2 \right)_{ij} \\
& \quad + 2 \left( Y_\ell^\dagger m_\ell^2 Y_\ell + m_{H_\ell}^2 Y_\ell^\dagger Y_\ell + A_\ell A_\ell^\dagger \right)_{ij}, \\
\end{align*} \]

\[ \begin{align*}
16\pi^2 \mu \frac{d}{d\mu} (m_\gamma^2)_{ij} &= -\frac{24}{5} g_t^2 |M_1|^2 \delta_{ij} + \frac{2}{5} g_t^2 S \delta_{ij} + 2 \left( m_\gamma^2 Y_\gamma^\dagger Y_\gamma + Y_\gamma^\dagger Y_\gamma m_\gamma^2 \right)_{ij} \\
& \quad + 4 \left( Y_\gamma^\dagger m_\gamma^2 Y_\gamma + m_{H_\gamma}^2 Y_\gamma^\dagger Y_\gamma + A_\gamma A_\gamma^\dagger \right)_{ij}, \\
\end{align*} \]

\[ \begin{align*}
16\pi^2 \mu \frac{d}{d\mu} (m_\delta^2)_{ij} &= 2 \left( m_\delta^2 Y_\delta^\dagger Y_\delta + Y_\delta^\dagger Y_\delta m_\delta^2 \right)_{ij} + 4 \left( Y_\delta^\dagger m_\delta^2 Y_\delta + m_{H_\delta}^2 Y_\delta^\dagger Y_\delta + A_\delta A_\delta^\dagger \right)_{ij}.
\end{align*} \]
\[
16\pi^2 \frac{d}{d\mu} (m_{H_u}^2) = - \left( \frac{6}{5} g_1^2 |M_1|^2 + 6g_2^2 |M_2|^2 \right) + \frac{3}{5} g_1^2 S \\
+ \ 6 \text{Tr} (m_q^2 Y_u^\dagger Y_u + Y_u^\dagger (m_u^2 + m_{H_u}^2) Y_u + A_u^\dagger A_u) \\
+ \ 2 \text{Tr} (m_q^2 Y_\nu^\dagger Y_\nu + Y_\nu^\dagger (m_\nu^2 + m_{H_u}^2) Y_\nu + A_\nu^\dagger A_\nu), \quad (47)
\]

\[
16\pi^2 \frac{d}{d\mu} (m_{H_d}^2) = - \left( \frac{6}{5} g_1^2 |M_1|^2 + 6g_2^2 |M_2|^2 \right) - \frac{3}{5} g_1^2 S \\
+ \ 6 \text{Tr} (m_q^2 Y_d^\dagger Y_d + Y_d^\dagger (m_d^2 + m_{H_d}^2) Y_d + A_d^\dagger A_d) \\
+ \ 2 \text{Tr} (m_q^2 Y_e^\dagger Y_e + Y_e^\dagger (m_e^2 + m_{H_d}^2) Y_e + A_e^\dagger A_e), \quad (48)
\]

where
\[
S \equiv \text{Tr}(m_q^2 + m_d^2 - 2m_u^2 - m_\nu^2 + m_{\nu}^2 - m_{H_d}^2 + m_{H_u}^2). \quad (49)
\]

**A.5 The 1-loop RGE for the soft SUSY breaking A-terms**

\[
16\pi^2 \frac{d}{d\mu} A_{u_{ij}} = \left\{ -\frac{13}{15} g_1^2 - 3g_2^2 - \frac{16}{3} g_3^2 + 3 \text{Tr}(Y_u^\dagger Y_u) + \text{Tr}(Y_\nu^\dagger Y_\nu) \right\} A_{u_{ij}} \\
+ \ 2 \left\{ \frac{13}{15} g_1^2 M_1 + 3g_2^2 M_2 + \frac{16}{3} g_3^2 M_3 + 3 \text{Tr}(Y_u^\dagger A_u) + \text{Tr}(Y_\nu^\dagger A_\nu) \right\} Y_{u_{ij}}, \quad (50)
\]

\[
16\pi^2 \frac{d}{d\mu} A_{d_{ij}} = \left\{ -\frac{7}{15} g_1^2 - 3g_2^2 - \frac{16}{3} g_3^2 + 3 \text{Tr}(Y_d^\dagger Y_d) + \text{Tr}(Y_e^\dagger Y_e) \right\} A_{d_{ij}} \\
+ \ 2 \left\{ \frac{7}{15} g_1^2 M_1 + 3g_2^2 M_2 + \frac{16}{3} g_3^2 M_3 + 3 \text{Tr}(Y_d^\dagger A_d) + \text{Tr}(Y_e^\dagger A_e) \right\} Y_{d_{ij}}, \quad (51)
\]

\[
16\pi^2 \frac{d}{d\mu} A_{e_{ij}} = \left\{ -\frac{9}{5} g_1^2 - 3g_2^2 + 3 \text{Tr}(Y_e^\dagger Y_e) + \text{Tr}(Y_\nu^\dagger Y_\nu) \right\} A_{e_{ij}} \\
+ \ 2 \left\{ \frac{9}{5} g_1^2 M_1 + 3g_2^2 M_2 + 3 \text{Tr}(Y_e^\dagger A_e) + \text{Tr}(Y_\nu^\dagger A_\nu) \right\} Y_{e_{ij}}, \quad (52)
\]

\[
16\pi^2 \frac{d}{d\mu} A_{e_{ij}} = \left\{ -\frac{3}{5} g_1^2 - 3g_2^2 + 3 \text{Tr}(Y_u^\dagger Y_u) + \text{Tr}(Y_\nu^\dagger Y_\nu) \right\} A_{e_{ij}} \\
+ \ 2 \left\{ \frac{3}{5} g_1^2 M_1 + 3g_2^2 M_2 + 3 \text{Tr}(Y_u^\dagger A_u) + \text{Tr}(Y_\nu^\dagger A_\nu) \right\} Y_{e_{ij}}, \quad (53)
\]
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Figure 1: The evolution of the soft mass for the field $\Delta_1$ from the SUSY breaking scale to the $B-L$ gaugino mass scale. The solid black, gray and dashed black lines are for $f = 1.5, 2.5, 3.5$, respectively. Here we have chosen $F_\phi = 100$ TeV, $g_{B-L} = 0.1$ and $M_{Z_{B-L}} = 5$ TeV.
Figure 2: The running behavior of the soft mass parameters $m_{\tilde{\ell}}$ (solid line) and $m_{\tilde{e}}$ (dotted line) are shown. Here we have chosen $F_\phi = 50$ TeV and $f = 2.5$. 
Table 1: Sparticle and Higgs boson mass spectra (in units of GeV) in the case of \( \tan \beta = 10 \), \( F_\phi = 50 \) TeV and \( f = 2.5 \).