R-mediation of Dynamical Supersymmetry Breaking

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Abstract

We propose a simple scenario of the dynamical supersymmetry breaking in four dimensional supergravity theories. The supersymmetry breaking sector is assumed to be completely separated as a sequestered sector from the visible sector, except for the communication by the gravity and $U(1)_R$ gauge interactions, and the supersymmetry breaking is mediated by the superconformal anomaly and $U(1)_R$ gauge interaction. Supersymmetry is dynamically broken by the interplay between the non-perturbative effect of the gauge interaction and Fayet-Iliopoulos D-term of $U(1)_R$ which necessarily exists in supergravity theories with gauged $U(1)_R$ symmetry. We construct an explicit model which gives phenomenologically acceptable mass spectrum of superpartners with vanishing (or very small) cosmological constant.

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I. INTRODUCTION

Low energy supersymmetry may play an important role in solving many problems of the particle physics. If it is the case, supersymmetry must be spontaneously broken, and all superpartners must have appropriate masses, since their effect is not observed yet. Therefore, finding a simple mechanism of the supersymmetry breaking and its mediation without any phenomenological problems is an important task. If we believe low energy supersymmetry, it is natural to consider the supergravity framework.

The simplest scenario of the supersymmetry breaking and its mediation in supergravity theories is the gravity mediation with Polonyi potential in the hidden sector [1], but supersymmetry is not dynamically broken in this scenario. Moreover, it is well known that the gravity mediation has a phenomenological problem: the degeneracy of squark masses at Planck scale is distorted by the quantum effect at low energies, which causes supersymmetric flavor problem. There is another conceptual problem about the gravity mediation as pointed out in Ref. [2]: it is not the mediation by the gravity, but the mediation by the higher dimensional contact interactions introduced by hand. Although it is possible that the superspace density, which defines the supergravity Lagrangian, contains infinite number of higher dimensional contact interactions so that the Kähler potential has a simple canonical form, the origin of these interactions is mysterious.

There is another possibility that the visible sector and hidden sector are completely separated, namely, no contact interaction among them in the superspace density. This situation would be naturally realized, if two sectors are confined in the different branes separated in the direction of extra dimensions (now the hidden sector should be called as the sequestered sector [3]). In this case the supersymmetry breaking at the sequestered sector is transmitted to the visible sector only through the superconformal anomaly [2–4]. In this anomaly mediation the masses of squarks highly degenerate at low energies and there is no supersymmetric flavor problem, but sleptons have negative masses ($m_{\text{slepton}}^2 < 0$). There are many attempts to solve this problem [2–6], and we usually need some additional fields which communicate two sectors. In this paper we introduce this additional communication by gauging $U(1)_R$ symmetry in four dimensional supergravity theories [8–10]. Since the charge of $U(1)_R$ symmetry does not commute with supercharges, it is natural to consider that the $U(1)_R$ gauge boson propagates in whole space-time including extra dimensions, and communicates two sectors.

It is also interesting to note that Fayet-Iliopoulos term for $U(1)_R$ must exist due to the symmetry of supergravity, and this term can act an important role in the supersymmetry breaking. In fact it has been shown that the supersymmetry can be dynamically broken by the interplay between this Fayet-Iliopoulos term and the non-perturbative effect of a gauge interaction [11]. Since the auxiliary field of $U(1)_R$ gauge multiplet has vacuum expectation value, both squarks and sleptons can have positive masses of the order of the gravitino mass in an appropriate R-charge assignment, and the problem of the anomaly mediation can be avoided.

This paper is organized as follows. In the next section we give a general argument on the supergravity Lagrangian with $U(1)_R$ gauge symmetry. We give a general formula for the chirality-conserving scalar mass in the presence of $U(1)_R$ gauge symmetry, which is an extension of the formula given in Ref. [12]. An explicit model is constructed in section 3,
and the analysis of the dynamics and mass spectrum is given in section 4. Section 5 contains our conclusions.

II. SUPERGRAVITY WITH U(1)<sub>R</sub> GAUGE SYMMETRY

In the superconformal framework [13–15] the general supergravity Lagrangian with U(1)<sub>R</sub> gauge symmetry is given by

\[
\mathcal{L} = -\frac{1}{2} \left[ \bar{\Sigma}_c e^{-2g_R V_R} \Sigma_c \Phi(S_I, \bar{S}_I e^{2Q_I g_R V_R} e^{2g_G V_G}) \right]_D + \left[ W(S_I) \Sigma_c^3 \right]_F - \frac{1}{4} \left[ f_R(S_I) W_R W_R \right]_F - \frac{1}{4} \left[ f_{ab}(S_I) W_a^a W_b^b \right]_F,
\]

where we use the notation in Ref. [14]. Here, \( S_I \) are matter chiral multiplets with flavor index \( I \) and U(1)<sub>R</sub> charge \( Q_I \), \( V_R \) and \( V_G \) (\( W_R \) and \( W_G \)) are vector (chiral) multiplets corresponding to the gauge group of U(1)<sub>R</sub> and \( G \), respectively. The multiplet \( \Sigma_c \) is the compensating multiplet, whose component should be appropriately fixed to obtain Poincaré supergravity. The functions \( \Phi \) and \( W \) are superspace densities in which interactions are described by the products of multiplets. Following the arguments in the previous section, we assume that there is no interaction between the visible sector fields \( S_i \) and \( V_{Gv} \) and the hidden (sequestered) sector fields \( S_\alpha \) and \( V_{Gh} \) in these superspace densities, namely,

\[
\Phi(S_I, \bar{S}_I e^{2Q_I g_R V_R} e^{2g_G V_G}) = \Phi_v(S_i, \bar{S}_i e^{2Q_i g_R V_R} e^{2g_{Gv} V_{Gv}}) + \Phi_h(S_\alpha, \bar{S}_\alpha e^{2Q_\alpha g_R V_R} e^{2g_{Gh} V_{Gh}}),
\]

\[
W(S_I) = W_v(S_i) + W_h(S_\alpha),
\]

where indices \( i \) and \( \alpha \) denote the flavors in the visible and hidden sectors, respectively, and \( Gv \) and \( Gh \) are gauge groups in each sector. The gauge kinetic function \( f_{ab}(S_I) \) should also be restricted as follows.

\[
\left[ f_{ab}(S_I) W_a^a W_b^b \right]_F \rightarrow \left[ f_{ab}^v(S_i) W_a^a W_b^b \right]_F + \left[ f_{ab}^h(S_\alpha) W_a^a W_b^b \right]_F.
\]

In the following we assume \( f_R(S_I) = 1 \) and \( f_{ab}^v = f_{ab}^h = \delta_{ab} \), for simplicity.

Note that the compensating multiplet \( \Sigma_c \) must have R-charge, since the superpotential \( W \) has R-charge. Therefore, the usual gauge choice to give Poincaré supergravity:

\[
z_c = \sqrt{3}, \quad \chi_{Rc} = 0, \quad b_\mu = 0
\]

does not preserve U(1)<sub>R</sub> symmetry, where \( z_c \) and \( \chi_{Rc} \) are scalar and spinor components of the compensating multiplet \( \Sigma_c \) and \( b_\mu \) is one of the gauge fields of the superconformal gauge group. We have to rescale the compensating multiplet to obtain the R-symmetric Poincaré supergravity:

\[
S_0 \equiv \Sigma_c (W(S_I))^{1/3}.
\]

The Lagrangian becomes
\[
\mathcal{L} = -\frac{1}{2} \left[ S_\sigma S_\sigma \tilde{\Phi}(S_I, \bar{S}_I e^{2Q_I g V_R e^{2g g V_R}}) \right]_D + \left[ S_0^3 \right]_F
- \frac{1}{4} [W_R W_R]_F - \frac{1}{4} [W_G W_G]_F - \frac{1}{4} [W_G h W_G h]_F,
\]
where
\[
\tilde{\Phi}(S_I, \bar{S}_I e^{2Q_I g V_R e^{2g g V_R}}) \equiv \frac{\Phi(S_I, \bar{S}_I e^{2Q_I g V_R e^{2g g V_R}})}{(\bar{W} S_I e^{2Q_I g V_R e^{2g g V_R}} W(S_I))^{1/3}}.
\]

The compensating multiplet \( S_0 \) is \( U(1)_R \) singlet now. It was shown in Ref. \cite{14} that the gauge fixing conditions of
\[
z_0 = \sqrt{3} \tilde{\Phi}^{-1/2}(z_I, z_I^*), \quad \chi_{R0} = -z_0 \tilde{\Phi} \tilde{\Phi}_J \chi_{RJ}, \quad b_\mu = 0
\]
directly give the standard form of the supergravity Lagrangian given in Ref. \cite{16}, where \( z_0 \) and \( \chi_{R0} \) are scalar and spinor components of the compensating multiplet \( S_0 \). \( \tilde{\Phi}_J \equiv \partial \tilde{\Phi}(z_I, z_I^*) / \partial z_J \) and \( z_J \) is the scalar components of \( S_J \). After all, the resultant Lagrangian in component fields has the standard form of Ref. \cite{16} including covariant derivatives for \( U(1)_R \) gauge symmetry. The Lagrangian is determined by a function
\[
G(z_I, z^{*I}) \equiv -3 \ln \tilde{\Phi}(z_I, z^{*I}) = -3 \ln \Phi(z_I, z^{*I}) + \ln |W(z_I)|^2,
\]
where \( \Phi \) and \( W \) satisfy the conditions of Eqs.\!(2) and \!(3). The difference of R-charges in covariant derivatives for each component field in a multiplet automatically appears due to the fact that \( W \) has non-trivial R-charge (see Ref. \cite{8}).

The potential for scalar fields is given as follows.
\[
V = V_F + V_D,
\]
where F-term contribution is
\[
V_F = e^G \left( G_I (G^{-1})_J G^J - 3 \right)
\]
and \( U(1)_R \) D-term contribution is
\[
V_D = \frac{g_R^2}{2} \left( G' Q_I z_I \right)^2.
\]
We take the reduced Planck scale as a unit of the mass scale. The chirality-conserving scalar mass can be obtained by differentiating this potential by \( z_I \) and \( z^{*J} \) and taking its vacuum expectation value. In addition to the conditions of Eqs.\!(2) and \!(3), we introduce the conditions of
\[
\langle \Phi_{vi} \rangle, \langle W_{vi} \rangle, \langle z_i \rangle = 0 \quad \text{or} \quad \ll 1.
\]
These conditions mean the assumption that the breaking scales of gauge symmetries in the visible sector should be much smaller than the reduced Planck scale. We obtain,
\[ \langle V_{F}^i \rangle = m^{ik} \langle (G^{-1})_k^j \rangle m_{lj} + \frac{2}{3} \langle V_F \rangle \langle G_j^i \rangle, \quad (15) \]

\[ \langle V_{D}^i \rangle = \frac{2}{3} \langle V_D \rangle \langle G_j^i \rangle + g_R^2 \left( Q_i - \frac{2}{3} \right) \langle D \rangle \langle G_j^i \rangle, \quad (16) \]

where \( m^{ik} \) is the supersymmetric mass and \( D \equiv G^I Q_I z_I \). The superpotential \( W \) has R-charge 2 in our convention. Therefore, the chirality-conserving supersymmetry-breaking scalar mass is obtained as

\[ \tilde{m}^2_{R_j} = \left( \frac{2}{3} \langle V \rangle + g_R^2 \left( Q_i - \frac{2}{3} \right) \langle D \rangle \right) \langle G_j^i \rangle. \quad (17) \]

The vacuum expectation value of the potential itself corresponds to the cosmological constant which should vanish in realistic models. We see that there is no gravity mediation, but there is “R-mediation” which is the tree-level contribution due to \( \langle D \rangle \neq 0 \).

**III. CONSTRUCTING A MODEL**

We construct an explicit model to show that the scenario which is described in the first section is possible. The particle contents of the model are summarized in Table I. In the following we simply introduce the role of each field without mentioning the dynamics in detail. The dynamics will be discussed in the next section.

As for the hidden sector, we take the same system which was introduced in Ref. [11]. It consists of two fields, \( Q_1 \) and \( Q_2 \), in the fundamental representation of SU(2)\(_H\) gauge group and a Yukawa interaction with a SU(2)\(_H\) singlet field \( S \):

\[ W_h = \lambda S [Q_1 Q_2], \quad (18) \]

where square brackets denote the contraction of SU(2)\(_H\) indices. Supersymmetry is dynamically broken by the interplay between the non-perturbative effect of SU(2)\(_H\) interaction and U(1)\(_R\) Fayet-Iliopoulos term, if there is no other field with R-charge less than 2/3.

The visible sector is based on the system of the minimal supersymmetric standard model. At least the R-charges of leptons must be larger than 2/3 so that sleptons obtain positive masses through R-mediation of Eq.(17) (assuming \( \langle D \rangle > 0 \) in this section). We simply assume that all the quarks and leptons have unit R-charge. This means that the R-charges of Higgs fields must be zero (note that this is less than 2/3), since we need the Yukawa couplings of

\[ W^Y_v = g_u H Q \bar{U} + g_d H Q \bar{D} + g_e H L \bar{N} + g_e H L \bar{E}, \quad (19) \]

where we suppress the generation indices, for simplicity.

To ensure the dynamical supersymmetry breaking we introduce other two Higgs fields, \( H' \) and \( \bar{H} \), with the mass terms of

\[ W'_v = \mu_u H H' + \mu_d \bar{H} H'. \quad (20) \]
Although there are negative contributions to the masses of $H$ and $\bar{H}$ from Eq.(17), these mass terms can make all masses of Higgs fields positive at the tree level. Therefore, the electroweak symmetry must be broken radiatively [17].

At this stage, all the gauge anomalies are cancelled out, except for $(SU(3)_c)^2U(1)_R$, $(SU(2)_L)^2U(1)_R$, $(U(1)_R)^3$ and $U(1)_R$(gravity)$^2$ anomalies. To cancel $(SU(3)_c)^2U(1)_R$ and $(SU(2)_L)^2U(1)_R$ anomalies, we further introduce additional fields, $\Omega_i$ and $\Sigma_i$, and Yukawa interactions with $X$.

$$W_{\nu}^X = gX(\Omega_i\Omega_i + \Sigma_i\Sigma_i) + mXX'. \quad (21)$$

The field $X'$ and the mass term with $X$ are required to have positive mass for $X$ and to ensure the dynamical supersymmetry breaking, since $X$ has R-charge less than 2/3. The fields $\Omega_i$ and $\Sigma_i$ become heavy by the vacuum expectation value of $X$ which is generated by the 1-loop effect of the Yukawa coupling in Eq.(21). The remaining anomalies, $(U(1)_R)^3$ and $U(1)_R$(gravity)$^2$, can be cancelled out by introducing, for example, many fields of R-charge two with appropriate values of $q_1$ and $q_2$. There may be much more sophisticated and convincing way to cancel these anomalies, but we leave it for further studies.

IV. DYNAMICS OF THE MODEL

Before discussing the dynamics of the model in detail, we have to make an assumption about the superspace density, $\Phi$. We simply assume as

$$\Phi(z_I, z^*_I) = 1 - \frac{1}{3} \sum_I z^*_Iz_I \quad (22)$$

respecting the condition of Eq.(2), where $z_I$ are the scalar components of all the chiral multiplets in the model. This gives canonical kinetic terms in the first order of the $1/M_P$ expansion, where $M_P = M_{\text{planck}}/\sqrt{8\pi}$ is the reduced Planck scale. In this case the scalar potential can be written as

$$V = V_F + V_D \quad (23)$$

with

$$V_F = \frac{1}{\Phi^2} \left\{ W^*_IW^I - \frac{1}{3}|z_IW^I|^2 + \left(W^*W^Iz_I + WW^*_Iz^*_I\right) - 3|W|^2 \right\}, \quad (24)$$

$$V_D = \frac{1}{\Phi^2} \frac{g_R^2}{2} \left\{ \left(Q_I - \frac{2}{3}\right) z^*_Iz_I + 2 \right\}^2, \quad (25)$$

where we neglect the D-term contributions form other gauge interactions.

First, we discuss the dynamics of the supersymmetry breaking. The instanton effect of $SU(2)_H$ gauge interaction can be described as a dynamically generated superpotential [18]. The effective superpotential for the hidden sector is

$$W_h^{\text{eff}} = \lambda S[Q_1Q_2] + \frac{\Lambda^5}{[Q_1Q_2]}, \quad (26)$$
where $\Lambda$ is the scale of the dynamics of SU(2)$_H$. If we assume that the vacuum expectation values of $Q_1$ and $Q_2$ lie on the flat direction of SU(2)$_H$, we have

$$V_F = \frac{1}{\phi^2} \left\{ (\lambda v^2)^2 + 2v^2 \left( \lambda s - \frac{\Lambda^5}{v^4} \right)^2 - \frac{25}{3} \left( \frac{\Lambda^5}{v^2} \right)^2 + \text{(visible sector)} \right\},$$

$$V_D = \frac{g_R^2}{2} D^2, \quad (28)$$

where

$$\Phi = 1 - \frac{1}{3} s^2 - \frac{2}{3} v^2 - \frac{1}{3} \left[ z^I z_I \right]_{\text{visible}},$$

$$D = \frac{1}{\Phi} \left\{ \left( q_S - \frac{2}{3} \right) s^2 + \left( q_1 + q_2 - \frac{4}{3} \right) v^2 + 2 + \left[ \left( Q_I - \frac{2}{3} \right) z^I z_I \right]_{\text{visible}} \right\}, \quad (30)$$

$v$ describes the flat direction of SU(2)$_H$, $s$ and $q_S$ are the vacuum expectation value and R-charge of $S$ ($q_S = 4$ and $q_1 + q_2 = -2$). It can be shown that all the visible sector fields do not have vacuum expectation values at the tree level. It is rather trivial for the fields with R-charge larger than $2/3$, but it is non-trivial for the fields $H$, $\bar{H}$ and $X$, since the vacuum expectation values of these fields negatively contribute to the vacuum energy in $V_D$. These fields do not have vacuum expectation values at the tree level, if the mass parameters $\mu_u$, $\mu_d$ and $m$ are appropriately large, as it will be explained at the end of this section. Therefore, in the following we consider the stationary conditions for $v$ and $s$ neglecting all the contributions from the visible sector.

The analysis is almost the same as in Ref. [11]. In case of $g_R^2 \gg \lambda \sim \Lambda^5$ there should be a solution of stationary conditions so that $v \simeq \sqrt{3/5}$ and $s \ll 1$, which results almost vanishing $D$. In this case the scalar potential approximately becomes

$$V \simeq \frac{1}{\phi^2} \left( \frac{3}{5} \right)^2 \left\{ \lambda^2 - 3 \left( \frac{5}{3} \right)^5 \Lambda^{10} \right\}. \quad (31)$$

Therefore, we can expect that there is a solution of vanishing (or vary small) cosmological constant with $\lambda \sim \sqrt{5}(5/3)^2\Lambda^5 \sim 6.2\Lambda^5$. Indeed, we can approximately obtain such a solution as

$$v \simeq \sqrt{\frac{3}{5}} + \frac{\sqrt{15}}{6}s^2 - \frac{1}{g_R^2} \frac{1243\lambda^2 + 6250\Lambda^{10}}{900\sqrt{15}}, \quad (32)$$

$$s \simeq \frac{675\Lambda^5}{486\lambda^2 + 6250\Lambda^{10}} \quad (33)$$

with vanishing cosmological constant by tuning $\lambda \simeq 6.9\Lambda^5$. A complete numerical analysis gives a solution

$$v \simeq \sqrt{3/5} + 0.012, \quad s \simeq 0.14 \quad (34)$$

with vanishing cosmological constant, where $g_R = 10^{-12}$, $\Lambda = 10^{-3}$ and $\lambda \simeq 6.9\Lambda^5$. At this vacuum the gravitino mass $m_{3/2}$ becomes
\[
m_{3/2} \equiv \langle e^{G/2} \rangle \simeq 5.0 \times \frac{\Lambda^5}{M_P^4}. \tag{35}
\]

The contribution to the mass of the scalar field due to \( \langle D \rangle \neq 0 \) can also be obtained from Eq. (17) as

\[
\tilde{m}_R^2(Q) = g_R^2 \langle D \rangle \left( Q - \frac{2}{3} \right) \simeq \left( 7.2 \times \frac{\Lambda^5}{M_P^4} \right)^2 \left( Q - \frac{2}{3} \right), \tag{36}
\]

where \( Q \) is the R-charge of the scalar field. We see that these supersymmetry-breaking masses are the same order of magnitude. The phenomenologically acceptable values of these masses can be obtained by changing the value of \( \Lambda \) within the same order of magnitude.

We summarize the spectrum of the supersymmetry-breaking masses and other supersymmetry breaking terms in the visible sector.

Gauginos in the visible sector can have masses only through the anomaly mediation, since there should be no hidden (sequestered) sector field in the gauge kinetic function. Therefore,

\[
m_{\lambda_i} \simeq \frac{\beta(g_i^2)}{2g_i^2} m_{3/2}, \tag{37}
\]

where \( g_i \) and \( \beta(g_i^2) \) are the gauge coupling and its beta function of the gauge group \( i \) in the visible sector, respectively (see Ref. [4] for more precise formula). There are two contributions to the scalar mass:

\[
\tilde{m}^2 \simeq -\frac{1}{4} \frac{d\gamma}{d\ln \mu} m_{3/2}^2 + \tilde{m}_R^2(Q), \tag{38}
\]

where \( \mu \) is the renormalization scale, and \( \gamma \) and \( Q \) are the anomalous dimension and R-charge of the scalar field, respectively. The first term is the contribution by the anomaly mediation and the second term is the contribution by R-mediation given by Eq. (36). The second contribution always dominates the first contribution, since the second contribution is the tree-level one. Therefore, the scalar field with \( Q > 2/3 \) naturally has positive mass. If we take \( m_{3/2} \sim 10 \text{TeV} \) to have gaugino masses heavier than about 100GeV, the scalar mass becomes of the order of \( (10 \text{TeV})^2 \).

Other supersymmetry breaking terms also emerge through the anomaly mediation. The A-term emerges corresponding to each Yukawa coupling through the anomaly mediation:

\[
A_{\Phi_1, \Phi_2, \Phi_3} = -\frac{1}{2} (\gamma_{\Phi_1} + \gamma_{\Phi_2} + \gamma_{\Phi_3}) m_{3/2}, \tag{39}
\]

where the Yukawa coupling of \( W_{\text{Yukawa}} = \lambda \Phi_1 \Phi_2 \Phi_3 \) is considered, and \( \gamma \) denotes the anomalous dimension of each field. The B-term emerges corresponding to each mass term at the tree level, since the mass term explicitly breaks the superconformal symmetry:

\[
B \simeq -m_{3/2}. \tag{40}
\]

Note that the order of the magnitude of \( A \) is always smaller than that of \( B \), since the B-term emerges at the tree level.
Next, we discuss the radiative electroweak symmetry breaking in our model. When we neglect the hidden sector, Higgs fields do not have vacuum expectation values at the tree level, if the following conditions are satisfied.

\[
\left( \mu_u^2 + \tilde{m}_R^2(q_H) \right) \left( \mu_d^2 + \tilde{m}_R^2(q_H') \right) - \mu_u^2 B^2 > 0, \\
\left( \mu_d^2 + \tilde{m}_R^2(q_H) \right) \left( \mu_u^2 + \tilde{m}_R^2(q_H') \right) - \mu_d^2 B^2 > 0,
\]

where \( q_H \) is the R-charge of \( H \) and \( \tilde{H} \) and \( q_H' \) is the R-charge of \( H' \) and \( \tilde{H}' \). This condition is satisfied, if \( \mu_u \) and \( \mu_d \) are slightly larger than \( m_{3/2} \). On the other hand, \( \mu_u^2 \) and \( \mu_d^2 \) must be larger than \( |\tilde{m}_R^2(q_H)| \simeq \frac{2}{3} m_{3/2}^2 \) so that our mechanism of the dynamical supersymmetry breaking works. Therefore, it is natural to consider that the masses \( \mu_u \) and \( \mu_d \) are slightly larger than \( m_{3/2} \) and the electroweak symmetry breaking occurs radiatively at the 1-loop level through the large Yukawa coupling of the top quark and the supersymmetry-breaking mass of the scalar top. We can analytically show that the radiative electroweak symmetry breaking at the weak scale is possible in this model.

Finally, we discuss the radiative mass generation for \( \Omega_i \) and \( \Sigma_i \) which are introduced for the gauge anomaly cancellation. These fields should become heavy so that running gauge couplings do not blow up before Planck scale. The point is whether \( X \) has vacuum expectation value or not, since \( \langle X \rangle \neq 0 \) make these fields massive through the first term of Eq. (21). When we neglect the hidden sector, the vacuum expectation value of \( X \) vanishes at the tree level, if the following condition is satisfied.

\[
\left( m^2 + \tilde{m}_R^2(q_X) \right) \left( m^2 + \tilde{m}_R^2(q_{X'}) \right) - m^2 B^2 > 0,
\]

where \( q_X \) and \( q_{X'} \) are the R-charges of \( X \) and \( X' \), respectively. This condition is satisfied, if \( m \) is larger than \( m_{3/2} \). On the other hand, \( m^2 \) must be larger than \( |\tilde{m}_R^2(q_X)| \simeq \frac{1}{6} m_{3/2}^2 \) so that our mechanism of the dynamical supersymmetry breaking works. Therefore, if we naturally take \( m^2 > m_{3/2}^2 \), \( X \) can not have vacuum expectation value at the tree level. But, it can have vacuum expectation value radiatively at the 1-loop level by the large coupling \( g \) and the supersymmetry-breaking masses of the scalar components of \( \Omega_i \) and \( \Sigma_i \). We can analytically show that the value of \( g \langle X \rangle \) can be large enough \( (g \langle X \rangle \simeq 10^{-2} M_P, \text{for example}) \) with \( m \sim m_{3/2}, g \sim 1 \) and very small \( g_R \sim 10^{-12} \).

V. CONCLUSION

We have proposed a sequestered sector scenario in which the supersymmetry breaking are mediated by the superconformal anomaly and \( \text{U}(1)_R \) gauge interaction without gravity mediation at the tree level. We constructed an explicit model in which supersymmetry is dynamically broken by the interplay between the non-perturbative effect of the gauge interaction and Fayet-Iliopoulos term of \( \text{U}(1)_R \). It was found that the problem of the tachyonic slepton in the anomaly mediation can be avoided in this scenario. The spectrum of the supersymmetry breaking masses is very simple and there is no supersymmetric flavor problem. We have also proposed a mechanism to radiatively generate the mass of the field which should not appear at low energies.

We mention a remarkable fact in this model: R-parity is not necessary. In the minimal supersymmetric standard model R-parity is usually assumed to forbid interactions which
violate baryon number symmetry. In our model $U(1)_R$ gauge symmetry naturally act the same or rather stronger role. It forbids in the superpotential not only renormalizable terms but also all the higher dimensional terms which violate baryon number symmetry. This is a simple realization of the idea proposed in Ref. [19].

We briefly summarize phenomenological consequences of this model. All the gauginos have the masses of the order of 100GeV, and the lightest superparticle would be a neutralino (bino or zino). Further understanding of the gaugino spectrum and the nature of the lightest superparticle requires more detailed analysis on the radiative correction to the spectrum as described in Ref. [3]. All the scalar fermions have the masses of the order of 10TeV. Therefore, in near future collider experiments we could not discover scalar fermions, but gauginos. The Higgs sector in this model is very different from the one in the minimal supersymmetric standard model, since it includes four Higgs doublets. There would be three charged Higgs and seven neutral Higgs, and all of them would have the masses of the order of 10TeV, except for one CP-even neutral Higgs which would have the mass of the order of the weak scale. Therefore, we could see one Higgs boson in near future collider experiments, but it would be impossible to see other Higgs particles.

There is an important point which have to be investigated in future. That is to derive the four dimensional effective theory from the fundamental theory in higher dimension. For example, if we consider a five dimensional theory as a fundamental theory, we have to integrate out the degrees of freedom which can propagate in fifth dimension. In our scenario such degrees of freedom are gravity and $R$ gauge interaction. Especially, it have to be investigated how an $U(1)$ component of the larger $R$ gauge symmetry in five dimensions is projected out to $U(1)_R$ gauge symmetry in four dimensional effective theories. It is also to be investigated how completely two sectors are separated in the superspace densities in four dimensional effective theories.

Although we still do not have the comprehensive analysis about deriving four dimensional effective theories from higher dimensional gauged supergravity theories, it is possible to expect that the very small value of the $U(1)_R$ gauge coupling in this model could be naturally obtained in the large extra dimension scenario. The result of Ref. [20] suggests that the natural (dimensionful) value of the gauge coupling in the higher dimensional theory would naturally result very small (dimensionless) value of the gauge coupling in the four dimensional effective theory, if the extra dimensions have relatively large volume. Other gauge and Yukawa couplings in this model are not suppressed by this mechanism, since the visible and sequestered sector are assumed to be confined in each 3-brane and only the supergravity multiplet and $U(1)_R$ gauge boson can propagate the bulk. If the volume of the extra dimensions is very large, the $U(1)_R$ gauge boson and graviton would be observed in near future collider experiments.

Finally, we want to emphasize that the proposed scenario is very simple, and we can rather easily construct calculable models which have concrete predictions. We believe that this direction is worth investigating further.

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REFERENCES

[1] J. Polonyi, Budapest preprint KFKI-93 (1977).
[2] L. Randall and R. Sundrum, Nucl. Phys. B557, 79 (1999).
[3] G.F. Giudice, M.A. Luty, H. Murayama and R. Rattazzi, JHEP 9812, 027 (1998).
[4] J.A. Bagger, T. Moro and E. Poppitz, JHEP 0004, 009 (2000).
[5] A. Pomarol and R. Rattazzi, JHEP 9905, 013 (1999).
[6] Z. Chacko, M.A. Luty, I. Maksymyk and E Ponton, JHEP 0004, 001 (2000).
[7] E. Katz, Y. Shadmi and Y. Shirman, JHEP 9908, 015 (1999).
[8] A.H. Chamseddine and H. Dreiner, Nucl. Phys. B458, 65 (1996).
[9] D.J. Castro, D.Z. Freedman and C. Manuel, Nucl. Phys. B461, 50 (1996).
[10] I. Jack and D.R.T. Jones, Phys. Lett. B482, 167 (2000).
[11] N. Kitazawa, N. Maru and N. Okada, hep-ph/9911251 to be published in Phys. Rev. D; hep-ph/0003240 to be published in Nucl. Phys. B.
[12] K. Inoue, M. Kawasaki, M. Yamaguchi and T. Yanagida, Phys. Rev. D45, 328 (1992).
[13] M. Kaku and P.K. Townsend, Phys. Lett. 76B, 54 (1978).
[14] T. Kugo and S. Uehara, Nucl. Phys. B222, 125 (1983); ibid B226, 49 (1983).
[15] S. Ferrara, L. Girardello, T. Kugo and A. Van Proeyen, Nucl. Phys. B223, 191 (1983).
[16] E. Cremmer, S. Ferrara, L. Girardello and A. Van Proeyen, Nucl. Phys. B212, 413 (1983).
[17] L. Ibañez and G.G. Ross, Phys. Lett. 110B, 215 (1982); K. Inoue, A. Kakuto, H. Komatsu and S. Takeshita, Prog. Theor. Phys. 68, 927 (1982); L. Alvarez-Gaume, M. Claudson and M.B. Wise, Nucl. Phys. B207, 96 (1982).
[18] I. Affleck, M. Dine and N. Seiberg, Nucl. Phys. B241, 493 (1984).
[19] S. Weinberg, Phys. Rev. D26, 287 (1982); L.J. Hall and I. Hinchliffe, Phys. Lett. 112B, 351 (1982).
[20] E.A. Mirabelli and M.E. Peskin, Phys. Rev. D58, 065002 (1998).
|       | SU(3)$_c$ | SU(2)$_L$ | U(1)$_Y$ | SU(2)$_H$ | U(1)$_R$ |
|-------|-----------|-----------|-----------|-----------|-----------|
| $Q$   | 3         | 2         | 1/6       | 1         | 1         |
| $\tilde{U}$ | 3$^*$   | 1         | -2/3      | 1         | 1         |
| $\tilde{D}$ | 3$^*$    | 1         | 1/3       | 1         | 1         |
| $L$   | 1         | 2         | -1/2      | 1         | 1         |
| $\tilde{N}$ | 1       | 1         | 0         | 1         | 1         |
| $\tilde{E}$ | 1       | 1         | 1         | 1         | 1         |
| $H$   | 1         | 2         | 1/2       | 1         | 0         |
| $\tilde{H}$ | 1       | 2         | -1/2      | 1         | 0         |
| $H'$  | 1         | 2         | 1/2       | 1         | 2         |
| $H''$ | 1         | 2         | -1/2      | 1         | 2         |
| $\Omega_i$ | 8       | 1         | 0         | 1         | 3/4       |
| $\Sigma_i$ | 1       | 3         | 0         | 1         | 3/4       |
| $X$   | 1         | 1         | 1         | 1         | 1/2       |
| $X'$  | 1         | 1         | 1         | 1         | 3/2       |
| $Q_1$ | 1         | 1         | 1         | 2         | $q_1$     |
| $Q_2$ | 1         | 1         | 1         | 2         | $q_2$     |
| $S$   | 1         | 1         | 1         | 1         | 4         |

**TABLE I.** Particle contents of the model. The system of the fields $Q_1$, $Q_2$ and $S$ with SU(2)$_H$ gauge symmetry constitutes the hidden (sequestered) sector. Other fields are the member of the visible sector. The index $i$ runs from 1 to 4 for $\Omega_i$ and $\Sigma_i$. The charges $q_1$ and $q_2$ follow the constraint of $q_1 + q_2 = -2$. The concrete values of these charges are determined through the cancellations of $(U(1)_R)^3$ and $U(1)_R$ (gravity)$^2$ anomalies (see text).