

R mediation of dynamical supersymmetry breaking

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We propose a simple scenario of 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 communication by gravity and $U(1)_R$ gauge interactions, and supersymmetry breaking is mediated by the superconformal anomaly and $U(1)_R$ gauge interaction. Supersymmetry is dynamically broken by the interplay between the nonperturbative effect of the gauge interaction and the 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 a phenomenologically acceptable mass spectrum of superpartners with a 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 particle physics. If this is the case, supersymmetry must be spontaneously broken, and all superpartners must have appropriate masses, since their effect has not been observed yet. Therefore, finding a simple mechanism of 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 supersymmetry breaking and its mediation in supergravity theories is gravity mediation with a Polonyi potential in the hidden sector [1], but supersymmetry is not dynamically broken in this scenario. Moreover, it is well known that gravity mediation has a phenomenological problem: the degeneracy of squark masses at the Planck scale is distorted by the quantum effect at low energies, which causes the supersymmetric flavor problem. There is another conceptual problem with gravity mediation as pointed out in Ref. [2]: it is not the mediation by gravity, but the mediation by higher dimensional contact interactions introduced by hand. Although it is possible that the superspace density, which defines the supergravity Lagrangian, contains an 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 the sequestered sector [2]). In this case 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 are highly degenerate at low energies and there is no supersymmetric flavor problem, but sleptons have negative masses ($m_{\text{slepton}}^2 < 0$). There have been many attempts to solve this problem [2,5–7], and we usually need some additional fields which bring contact between 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 brings contact between two sectors.

It is also interesting to note that the Fayet-Iliopoulos term for $U(1)_R$ must exist due to the symmetry of supergravity, and this term can play an important role in supersymmetry breaking. In fact it has been shown that supersymmetry can be dynamically broken by the interplay between this Fayet-Iliopoulos term and the nonperturbative effect of a gauge interaction [11]. Since the auxiliary field of the $U(1)_R$ 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 Sec. III, and the analysis of the dynamics and mass spectrum is given in Sec. IV. Section V contains our conclusions.

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II. SUPERGRAVITY WITH $U(1)_R$ GAUGE SYMMETRY

In the superconformal framework [13–15] the general supergravity Lagrangian with $U(1)_R$ gauge symmetry is given by

$$\begin{aligned} \mathcal{L} = & -\frac{1}{2}[\bar{\Sigma}_c e^{-2g_R V_R} \Sigma_c \Phi(S_I, \bar{S}^I e^{2Q_{I^R} V_R} e^{2g_G V_G})]_D \\ & + [W(S_I) \Sigma_c^3]_F - \frac{1}{4}[f_R(S_I) W_R W_R]_F \\ & - \frac{1}{4}[f_{ab}(S_I) W_G^a W_G^b]_F, \end{aligned} \quad (1)$$

where we use the notation in Ref. [14]. Here, S_I are matter chiral multiplets with flavor index I and $U(1)_R$ charge Q_I , and V_R and V_G (W_R and W_G) are vector (chiral) multiplets corresponding to the gauge group of $U(1)_R$ and G , respectively. The multiplet Σ_c is the compensating multiplet, whose component should be appropriately fixed to obtain Poincaré supergravity. The functions Φ 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_α and V_{Gh} in these superspace densities, namely,

$$\begin{aligned} \Phi(S_I, \bar{S}^I e^{2Q_{I^R} V_R} e^{2g_G V_G}) = & \Phi_v(S_i, \bar{S}^i e^{2Q_{i^R} V_R} e^{2g_{Gv} V_{Gv}}) \\ & + \Phi_h(S_\alpha, \bar{S}^\alpha e^{2Q_{\alpha^R} V_R} e^{2g_{Gh} V_{Gh}}), \end{aligned} \quad (2)$$

$$W(S_I) = W_v(S_i) + W_h(S_\alpha), \quad (3)$$

where the indices i and α 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:

$$\begin{aligned} [f_{ab}(S_I) W_G^a W_G^b]_F \rightarrow & [f_{ab}^{Gv}(S_i) W_{Gv}^a W_{Gv}^b]_F \\ & + [f_{ab}^{Gh}(S_\alpha) W_{Gh}^a W_{Gh}^b]_F. \end{aligned} \quad (4)$$

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

Note that the compensating multiplet Σ_c must have R charge, since the superpotential W has R charge. Therefore, the usual gauge choice to give Poincaré supergravity,

$$z_c = \sqrt{3}, \quad \chi_{Rc} = 0, \quad b_\mu = 0, \quad (5)$$

does not preserve $U(1)_R$ symmetry, where z_c and χ_{Rc} are scalar and spinor components of the compensating multiplet Σ_c and b_μ 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:

$$S_0 \equiv \Sigma_c [W(S_I)]^{1/3}. \quad (6)$$

The Lagrangian becomes

$$\begin{aligned} \mathcal{L} = & -\frac{1}{2}[\bar{S}_0 S_0 \tilde{\Phi}(S_I, \bar{S}^I e^{2Q_{I^R} V_R} e^{2g_G V_G})]_D + [S_0^3]_F \\ & - \frac{1}{4}[W_R W_R]_F - \frac{1}{4}[W_{Gv} W_{Gv}]_F - \frac{1}{4}[W_{Gh} W_{Gh}]_F, \end{aligned} \quad (7)$$

where

$$\begin{aligned} \tilde{\Phi}(S_I, \bar{S}^I e^{2Q_{I^R} V_R} e^{2g_G V_G}) \\ \equiv \frac{\Phi(S_I, \bar{S}^I e^{2Q_{I^R} V_R} e^{2g_G V_G})}{[\bar{W}(\bar{S}^I e^{2Q_{I^R} V_R} e^{2g_G V_G}) W(S_I)]^{1/3}}. \end{aligned} \quad (8)$$

The compensating multiplet S_0 is $U(1)_R$ singlet now. It was shown in Ref. [14] that the gauge fixing conditions of

$$z_0 = \sqrt{3} \tilde{\Phi}^{-1/2}(z_I, z_I^*), \quad \chi_{R0} = -z_0 \tilde{\Phi}^{-1} \tilde{\Phi}^J \chi_{RJ}, \quad b_\mu = 0 \quad (9)$$

directly give the standard form of the supergravity Lagrangian given in Ref. [16], where z_0 and χ_{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. [16] including covariant derivatives for $U(1)_R$ gauge symmetry. The Lagrangian is determined by a function

$$\begin{aligned} G(z_I, z_I^*) \equiv & -3 \ln \tilde{\Phi}(z_I, z_I^*) = -3 \ln \Phi(z_I, z_I^*) \\ & + \ln |W(z_I)|^2, \end{aligned} \quad (10)$$

where Φ 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 nontrivial R charge (see Ref. [8]).

The potential for scalar fields is given as follows:

$$V = V_F + V_D, \quad (11)$$

where the F -term contribution is

$$V_F = e^G [G_I (G^{-1})^I_J G^J - 3] \quad (12)$$

and the $U(1)_R$ D -term contribution is

$$V_D = \frac{g_R^2}{2} (G^I Q_I z_I)^2. \quad (13)$$

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

$$\langle \Phi_{vi} \rangle, \langle W_{vi} \rangle, \langle z_i \rangle = 0 \text{ or } \ll 1. \quad (14)$$

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

$$\langle V_{Fj}^i \rangle = m^{ik} \langle (G^{-1})_k^j \rangle m_{lj} + \frac{2}{3} \langle V_F \rangle \langle G_j^i \rangle, \quad (15)$$

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

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

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

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 the $SU(2)_H$ 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 nonperturbative effect of the $SU(2)_H$ 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 a 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 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 Yukawa couplings of

$$W_v^Y = g_u H Q \bar{U} + g_d \bar{H} Q \bar{D} + g_\nu H L \bar{N} + g_e \bar{H} L \bar{E}, \quad (19)$$

where we suppress generation indices, for simplicity.

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 Ω_i and Σ_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 $U(1)_R(\text{gravity})^2$ anomalies (see text).

	$SU(3)_c$	$SU(2)_L$	$U(1)_Y$	$SU(2)_H$	$U(1)_R$
Q	3	2	1/6	1	1
\bar{U}	3*	1	-2/3	1	1
\bar{D}	3*	1	1/3	1	1
L	1	2	-1/2	1	1
\bar{N}	1	1	0	1	1
\bar{E}	1	1	1	1	1
H	1	2	1/2	1	0
\bar{H}	1	2	-1/2	1	0
H'	1	2	1/2	1	2
\bar{H}'	1	2	-1/2	1	2
Ω_i	8	1	0	1	3/4
Σ_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

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

$$W'_v = \mu_u H \bar{H}' + \mu_d \bar{H} H'. \quad (20)$$

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 canceled out, except for $[SU(3)_c]^2 U(1)_R$, $[SU(2)_L]^2 U(1)_R$, $[U(1)_R]^3$, and $U(1)_R(\text{gravity})^2$ anomalies. To cancel $[SU(3)_c]^2 U(1)_R$ and $[SU(2)_L]^2 U(1)_R$ anomalies, we further introduce additional fields Ω_i and Σ_i and Yukawa interactions with X :

$$W_v^X = g X (\Omega_i \Omega_i + \Sigma_i \Sigma_i) + m X X'. \quad (21)$$

The field X' and the mass term with X are required to have positive mass for X and to ensure dynamical supersymmetry breaking, since X has R charge less than $2/3$. The fields Ω_i and Σ_i become heavy by the vacuum expectation value of X which is generated by the one-loop effect of the Yukawa coupling in Eq. (21). The remaining anomalies $[U(1)_R]^3$ and $U(1)_R(\text{gravity})^2$ can be canceled out by introducing, for example, many fields of R charge 2 with appropriate values of q_1 and q_2 . There may be much more sophisticated and convincing ways to cancel these anomalies, but we leave this to 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, Φ . We simply assume as

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

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

$$V_F = \frac{1}{\Phi^2} \left\{ W_I^* W^I - \frac{1}{3} |z_I W^I|^2 + (W^* W^I z_I + W W_I^* z^{*I}) - 3|W|^2 \right\}, \quad (24)$$

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

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

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

$$W_h^{\text{eff}} = \lambda S [Q_1 Q_2] + \frac{\Lambda^5}{[Q_1 Q_2]}, \quad (26)$$

where Λ 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

$$V_F = \frac{1}{\Phi^2} \left\{ (\lambda v^2)^2 + 2v^2 \left(\lambda s - \frac{\Lambda^5}{v^4} \right)^2 - \frac{25}{3} \left(\frac{\Lambda^5}{v^2} \right)^2 + (\text{visible sector}) \right\}, \quad (27)$$

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

where

$$\Phi = 1 - \frac{1}{3} s^2 - \frac{2}{3} v^2 - \frac{1}{3} [z^{*I} z_I]_{\text{visible}}, \quad (29)$$

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

where v describes the flat direction of $SU(2)_H$, and 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 visible sector fields do not have vacuum expectation values at the tree level. It is rather trivial for fields with R charge larger than $2/3$, but it is nontrivial 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 μ_u , μ_d , and m are appropriately large, as will be explained at the end of this section. Therefore, in the following we consider stationary conditions for v and s , neglecting all contributions from the visible sector.

The analysis is almost the same as in Ref. [11]. In the 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 in 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 a vanishing (or very 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

$$v \simeq \sqrt{\frac{3}{5}} + \sqrt{\frac{15}{6}} s^2 - \frac{1}{g_R^2} \frac{243\lambda^2 + 6250\Lambda^{10}}{900\sqrt{15}}, \quad (32)$$

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

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

$$m_{3/2} \equiv \langle e^{G/2} \rangle \simeq 5.0 \frac{\Lambda^5}{M_P^4}. \quad (35)$$

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

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

where Q is the R charge of the scalar field. We see that these supersymmetry-breaking masses are the same order of magnitude. Phenomenologically acceptable values of these masses can be obtained by changing the value of Λ 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 anomaly mediation, since there should be no hidden (sequestered) sector field in the gauge kinetic function. Therefore,

$$m_{\lambda_i} \simeq \frac{\beta(g_i^2)}{2g_i^2} m_{3/2}, \quad (37)$$

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 a more precise formula). There are two contributions to the scalar mass:

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

where μ is the renormalization scale, and γ and Q are the anomalous dimension and R charge of the scalar field, respectively. The first term is the contribution by 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 $m_{3/2} \sim 10$ TeV to have gaugino masses heavier than about 100 GeV, the scalar mass becomes of the order of $(10 \text{ TeV})^2$.

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

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

where the Yukawa coupling of $W_{\text{Yukawa}} = \lambda \Phi_1 \Phi_2 \Phi_3$ is considered, and γ 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 superconformal symmetry:

$$B \simeq -m_{3/2}. \quad (40)$$

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:

$$[\mu_u^2 + \tilde{m}_R^2(q_H)][\mu_u^2 + \tilde{m}_R^2(q_{H'})] - \mu_u^2 B^2 > 0, \quad (41)$$

$$[\mu_d^2 + \tilde{m}_R^2(q_H)][\mu_d^2 + \tilde{m}_R^2(q_{H'})] - \mu_d^2 B^2 > 0, \quad (42)$$

where q_H is the R charge of H and \bar{H} and $q_{H'}$ is the R charge of H' and \bar{H}' . This condition is satisfied if μ_u and μ_d are slightly larger than $m_{3/2}$. On the other hand, μ_u^2 and μ_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 μ_u and μ_d are slightly larger than $m_{3/2}$ and electroweak symmetry breaking occurs radiatively at the one-loop level through the large Yukawa coupling of the top quark and the supersymmetry-breaking mass of the scalar top quark. We can analytically show that radiative electroweak symmetry breaking at the weak scale is possible in this model.

Finally, we discuss radiative mass generation for Ω_i and Σ_i which are introduced for the gauge anomaly cancellation. These fields should become heavy so that running gauge couplings do not blow up before the Planck scale. The point is whether X has a vacuum expectation value or not, since $\langle X \rangle \neq 0$ makes 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:

$$[m^2 + \tilde{m}_R^2(q_X)][m^2 + \tilde{m}_R^2(q_{X'})] - m^2 B^2 > 0, \quad (43)$$

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 dynamical supersymmetry breaking works. Therefore, if we naturally take $m^2 > m_{3/2}^2$, X cannot have a vacuum expectation value at the tree level. But it can have a vacuum expectation value radiatively at the one-loop level by the large coupling g and the supersymmetry-breaking masses of the scalar components of Ω_i and Σ_i . We can analytically show that the value of $g\langle X \rangle$ can be large enough ($g\langle X \rangle \simeq 10^{-2} M_P$, 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 supersymmetry breaking is mediated by the superconformal anomaly and $U(1)_R$ 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 nonperturbative effect of the gauge interaction and the Fayet-Iliopoulos term of $U(1)_R$. It was found that the problem of the tachyonic slepton in 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 acts the same or a rather stronger role. It forbids in the superpotential not only renormalizable terms but also all higher dimensional terms which violate baryon number symmetry. This is a simple realization of the idea proposed in Ref. [19].

We briefly summarize the phenomenological consequences of this model. All gauginos have masses of the order of 100 GeV, and the lightest superparticle would be a neutralino (B -ino or Z -ino). Further understanding of the gaugino spectrum and the nature of the lightest superparticle requires a more detailed analysis of the radiative correction to the spectrum as described in Ref. [3]. All scalar fermions have masses of the order of 10 TeV. 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 bosons and seven neutral Higgs bosons and all of them would have masses of the order of 10 TeV, except for one CP -even neutral Higgs boson which would have 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 has to be investigated in future: that is, to derive the four-dimensional effective theory from the fundamental theory in higher dimensions. For example, if we consider a five-dimensional theory as the fundamental theory, we have to integrate out the degrees of freedom which can propagate in the fifth dimension. In our scenario such degrees of freedom are gravity and the R gauge interaction. Especially, it has to be investigated how a $U(1)$ component of 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 a comprehensive analysis of 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 could naturally result in a very small (dimensionless) value of the gauge coupling in the four-dimensional effective theory if the extra dimensions have a relatively large volume. Other gauge and Yukawa couplings in this model are not suppressed by this mechanism, since the visible and sequestered sectors are assumed to be confined in each three-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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- [1] J. Polonyi, Budapest Report No. 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, J. High Energy Phys. **12**, 027 (1998).
 [4] J.A. Bagger, T. Moroi, and E. Poppitz, J. High Energy Phys. **04**, 009 (2000).
 [5] A. Pomarol and R. Rattazzi, J. High Energy Phys. **05**, 013 (1999).
 [6] Z. Chacko, M.A. Luty, I. Maksymyk, and E. Ponton, J. High Energy Phys. **04**, 001 (2000).
 [7] E. Katz, Y. Shadmi, and Y. Shirman, J. High Energy Phys. **08**, 015 (1999).
 [8] A.H. Chamseddine and H. Dreiner, Nucl. Phys. **B458**, 65 (1996).
 [9] D.J. Castaño, D.Z. Freedman, and C. Manuel, Nucl. Phys. **B461**, 50 (1996).
 [10] I. Jack and D.R.T. Jones, Phys. Lett. B **482**, 167 (2000).
 [11] N. Kitazawa, N. Maru, and N. Okada, Phys. Rev. D **62**, 077701 (2000); Nucl. Phys. **B586**, 261 (2000).
 [12] K. Inoue, M. Kawasaki, M. Yamaguchi, and T. Yanagida, Phys. Rev. D **45**, 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); **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. Ibáñez 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. D **26**, 287 (1982); L.J. Hall and I. Hinchliffe, Phys. Lett. **112B**, 351 (1982).
 [20] E.A. Mirabelli and M.E. Peskin, Phys. Rev. D **58**, 065002 (1998).