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H. Itoh – Theory Division, KEK
N. Okada – Theory Division, KEK
T. Yamashita – SISSA

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Low scale gravity mediation with warped extra dimension and collider phenomenology on the hidden sector

Hideo Itoh,[1,*] Nobuchika Okada,[1,2,†] and Toshifumi Yamashita[3,‡]

[1]*Theory Division, KEK, Oho 1-1, Tsukuba 305-0801, Japan*
[2]*Department of Particle and Nuclear Physics, The Graduate University for Advanced Studies (Sokendai), Oho 1-1, Tsukuba 305-0801, Japan*
[3]*SISSA, Via Beirut 4, I-34014 Trieste, Italy*
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We propose a scenario of gravity mediated supersymmetry breaking (gravity mediation) in a supersymmetric Randall-Sundrum model. In our setup, both the visible sector and the hidden sector coexist on the infrared (IR) brane. We introduce the Polonyi model as a simple hidden sector. Because of the warped metric, the effective cutoff scale on the IR brane is ''warped down,'' so that the gravity mediation occurs at a low scale. As a result, the gravitino is naturally the lightest superpartner (LSP) and contact interactions between the hidden and the visible sector fields become stronger. We address phenomenologies for various IR cutoff scales. In particular, we investigate collider phenomenology involving a scalar field (Polonyi field) in the hidden sector for the case with the IR cutoff around 10 TeV. We find a possibility that the hidden sector scalar can be produced at the LHC and the international linear collider (ILC). Interestingly, the scalar behaves like the Higgs boson of the standard model in the production process, while its decay process is quite different and, once produced, it will provide us with a very clean signature. The hidden sector may be no longer hidden.

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I. INTRODUCTION

Supersymmetric (SUSY) extension of the standard model is one of the most promising ways to solve the gauge hierarchy problem of the standard model. The minimal supersymmetric standard model (MSSM) is the simplest supersymmetric extension of the standard model, and its various phenomenological aspects have been investigated for many years. However, since no superpartner has been observed in the current experiments, SUSY should be broken at low energies. The origin of SUSY breaking and its mediation mechanism to the visible (MSSM) sector is one of the most important issues in any supersymmetric phenomenological models.

To be consistent with our observations that the nature is almost flavor blind and *CP* invariant, the way to transmit the SUSY breaking to the visible sector is severely constrained. For a few decades, various mechanisms for the SUSY breaking mediation have been proposed in the context of four-dimensional models and also brane world scenarios [1]. Each proposed model provides typical soft SUSY breaking mass spectra. Once superpartners are observed at future colliders and their mass spectra are precisely measured, the origin of the SUSY breaking mediation mechanism could be revealed.

The simplest model of SUSY breaking is the Polonyi model [2], where a chiral superfield singlet under the standard model gauge group and its tadpole term in superpotential are introduced. Then, the nonzero F-term is developed, and SUSY is broken. After SUSY is broken, the SUSY breaking is transmitted to the visible sector through some interactions such as gravity interactions or gauge interactions. Operators relevant to the SUSY breaking mediation are effectively described as higher dimensional contact operators between the hidden sector and the visible sector superfields. The scale of the SUSY breaking mediation is characterized by the mass scale of the contact operators. There are two well-known examples of SUSY breaking mediation. One is the gravity mediation in the minimal supergravity scenario [3], where the scale of the SUSY breaking mediation is the Planck scale which is nothing but the cutoff scale of supergravity. The other is the gauge mediated SUSY breaking (GMSB) [4], where SUSY breaking is transmitted through the standard model gauge interactions with the so-called ''messenger'' fields. The scale of the gauge mediation is characterized by the mass scale of the messenger fields, which is far below the Planck scale.

In this paper, we propose a new scenario of the gravity mediation in a supersymmetric Randall-Sundrum model. We introduce both of the visible and the hidden sectors on the infrared (IR) brane. As a simple hidden sector we take the Polonyi model, and consider the gravity mediation through contact operators between the hidden sector and the visible sector superfields. As first proposed by Randall and Sundrum [5], in four-dimensional effective theory, an original mass scale on the IR brane is ''warped down'' to a low scale by the warp factor. Therefore, in our model, the gravity mediation occurs at the low scale due to the warp-

*Electronic address: hideo@post.kek.jp
†Electronic address: okadan@post.kek.jp
‡Electronic address: yamasita@sissa.it

ing down of the original cutoff scale of the model.[1] We call this scenario "low scale gravity mediation."

As a result of the SUSY breaking mediation at the low scale, the gravitino is naturally the lightest superpartner (LSP), so that it can be a candidate of the dark matter in the present universe. Recently, this LSP gravitino scenario has been intensively studied in cosmology [6] and also in collider physics [7]. Our model can naturally provide this scenario.

Our model has further interesting features. The contact operators relevant to the gravity mediation also provide contact interactions between a scalar field (Polonyi field) in the hidden sector and the standard model fields. In the context of the warped extra dimension, the effective cutoff scale can be as low as 1 TeV without any serious fine-tuning for parameters in the model. We will find a possibility that the hidden sector scalar can be produced at the LHC and the ILC with a very clean signature, if the effective cutoff scale is low enough.

This paper is organized as follows. In the next section, we propose a SUSY model with a warped extra dimension which realizes the low scale gravity mediation. We also present a concrete model of the hidden sector as an example, which is nothing but the Polonyi model on the IR brane. In Sec. III, we address various phenomenological aspects of our model. In particular, we focus on collider phenomenologies involving the hidden sector scalar, and find a possibility that the hidden sector scalar can be discovered at the LHC and the ILC. The last section is devoted to summary and discussions.

II. LOW SCALE GRAVITY MEDIATION

We consider a SUSY model in the warped five-dimensional brane world scenario [5]. The fifth dimension is compactified on the orbifold S^1/Z_2 with two branes, ultraviolet (UV) and infrared (IR) branes, sitting on each orbifold fixed point. With an appropriate tuning for cosmological constants in the bulk and on the branes, we obtain the warped metric [5],

$$ds^2 = e^{-2kr|y|}\eta_{\mu\nu}dx^\mu dx^\nu - r^2 dy^2, \qquad (2.1)$$

for $-\pi \leq y \leq \pi$, where k is the AdS curvature, and r and y are the radius and the angle of S^1, respectively.

By the compactification on the orbifold, $N = 1$ SUSY of the five-dimensional theory, which corresponds to $N = 2$ SUSY in the four-dimensional point of view, is broken down to four-dimensional $N = 1$ SUSY. The supergravity Lagrangian of this system can be described in terms of the superfield formalism of four-dimensional $N = 1$ SUSY theories [8–10]. For simplicity, here we consider only the gravity multiplet in the bulk whose Lagrangian is given by

$$\mathcal{L}_{\text{bulk}} = -3\int d^4\theta \frac{M_5^3}{k}(\phi^\dagger\phi - \omega^\dagger\omega), \qquad (2.2)$$

where M_5 is the five-dimensional Planck mass, $\phi = 1 + \theta^2 F_\phi$ is the compensating multiplet in the superconformal framework of supergravity [11], and $\omega = \phi e^{-\pi kT}$ with a radion chiral multiplet T whose real part of the scalar component is the fifth dimensional radius. Lagrangian for some chiral and gauge multiplets on the UV brane are generally described as

$$\mathcal{L}_{\text{UV}}^{\text{chiral}} = \int d^4\theta\phi^\dagger\phi\mathcal{K}_{\text{UV}} + \left(\int d^2\theta\phi^3 W_{\text{UV}} + \text{H.c.}\right),$$
$$\mathcal{L}_{\text{UV}}^{\text{gauge}} = \frac{1}{4}\int d^2\theta f_a \mathcal{W}^{a\alpha}\mathcal{W}^a_\alpha + \text{H.c.}, \qquad (2.3)$$

where $\mathcal{K}_{\text{UV}}$ and W_{UV} are Kahler potential and superpotential, respectively, and f_a is the gauge kinetic function. Replacing ϕ by ω due to the warped metric, we obtain the general Lagrangian for some chiral and gauge multiplets on the IR brane,

$$\mathcal{L}_{\text{IR}}^{\text{chiral}} = \int d^4\theta\omega^\dagger\omega\mathcal{K}_{\text{IR}} + \left(\int d^2\theta\omega^3 W_{\text{IR}} + \text{H.c.}\right),$$
$$\mathcal{L}_{\text{IR}}^{\text{gauge}} = \frac{1}{4}\int d^2\theta f_a \mathcal{W}^{a\alpha}\mathcal{W}^a_\alpha + \text{H.c.}. \qquad (2.4)$$

The setup of our model is that both the hidden and visible sectors reside on the IR brane. Except for gravity multiplet residing in the bulk, this is the same setup as in usual four-dimensional models. We introduce a simple hidden sector with a chiral superfield (X) singlet under the standard model gauge group, by whose F component SUSY is broken. We set free parts in the Kahler potential and the gauge kinetic functions for each superfield of the canonical form such as $\mathcal{K}_{\text{IR}}^{\text{free}} = \sum_i Q_i^\dagger Q_i + X^\dagger X$ and $f_a^{\text{free}} = 1$, where Q_i denotes matter and Higgs multiplets in the MSSM with flavor index i.

Now let us consider the gravity mediation on the IR brane, namely, SUSY breaking is transmitted through contact operators between the visible and the hidden sector superfields. For the gravity mediation in four-dimensional models, the contact operators are suppressed by the four-dimensional Planck mass, which is nothing but the cutoff of four-dimensional supergravity. In our case, the original cutoff should be the five-dimensional Planck mass. In addition to the free parts of the Kahler potential and the gauge kinetic functions, we introduce the following contact operators relevant to the gravity mediation,

$$\mathcal{L}_{\text{contact}} = -\int d^4\theta\omega^\dagger\omega\left(c_A^{ij}\frac{X + X^\dagger}{M_5} + c_0^{ij}\frac{X^\dagger X}{M_5^2}\right)Q_i^\dagger Q_j$$
$$- \frac{1}{4}\int d^2\theta c_a \frac{X}{M_5}\mathcal{W}^{a\alpha}\mathcal{W}^a_\alpha + \text{H.c.}, \qquad (2.5)$$

where c_A^{ij}, c_0^{ij}, and c_a are dimensionless parameters naturally of order one, and $a = 1, 2, 3$ corresponds to $U(1)_Y$,

[1] Here, the warp factor is not necessarily so strong to solve the hierarchy problem completely. The remaining hierarchy is solved by SUSY.

$SU(2)_L$, and $SU(3)_c$ gauge groups of the standard model. Although the coefficients, c_A^{ij} and c_0^{ij}, are generally flavor dependent, we assume the universal coefficients (minimal ansatz), $c_A^{ij} = c_A \delta^{ij}$ and $c_A^{ij} = c_0 \delta^{ij}$, as usual in the minimal supergravity scenario, otherwise flavor-changing-neutral-current (FCNC) processes through superpartners exceed the current experimental bounds.

Note that, in the present form, superfields have not yet been suitably normalized, because of the warped metric. The correct description in effective four-dimensional theory is given by replacing each chiral superfield as $Q_i, X \to Q_i/\omega, X/\omega$ so as to eliminate ω from their free kinetic terms. Now we arrive at the contact operators in effective four-dimensional theory of the form

$$\mathcal{L}_{\rm contact}^{\rm eff} = -\int d^4\theta \left(c_A \frac{X + X^\dagger}{\Lambda_{\rm IR}} + c_0 \frac{X^\dagger X}{\Lambda_{\rm IR}^2} \right) Q_i^\dagger Q_i - \frac{1}{4} \int d^2\theta c_a \frac{X}{\Lambda_{\rm IR}} \mathcal{W}^{a\alpha} \mathcal{W}_\alpha^a + \text{H.c.}. \tag{2.6}$$

Here, the effective cutoff $\Lambda_{\rm IR} = \omega M_5$ appears. This is the most important feature of a model with the warped extra dimension, that is, any dimensional parameters on the IR brane are inevitably warped down according to their mass dimensions in effective four-dimensional theory. As discussed in the original paper [5], $\Lambda_{\rm IR} \ll M_P$ can be achieved with a mild hierarchy among the original parameters. For example, $\Lambda_{\rm IR} \sim 1$ TeV can be realized by $M_5 \sim k \sim 11.3/r$. Here, $M_P = 2.4 \times 10^{18}$ GeV is the reduced Planck mass in four dimensions which is defined as $M_P^2 = M_5^3/k$ in the strongly warped case $\omega \ll 1$. Thus, we can take any value of the IR cutoff without theoretical difficulty.

Once the nonzero F-term of the hidden sector field, F_X, is developed, the contact operators introduced above lead to soft SUSY breaking terms in the visible sector. Assuming $\langle X \rangle \ll \Lambda_{\rm IR}$, for simplicity, scalar squared masses, the A-parameter and gaugino masses are extracted as[2]

$$\tilde{m}^2 = (c_A^2 + c_0) \frac{|F_X|^2}{\Lambda_{\rm IR}^2}, \tag{2.7}$$

$$A = 3 c_A \frac{F_X}{\Lambda_{\rm IR}}, \tag{2.8}$$

$$M_a = \frac{1}{2} c_a \frac{F_X}{\Lambda_{\rm IR}}. \tag{2.9}$$

For Higgs superfields, we can generally introduce contact terms between X and the gauge invariant product of up- and down-type Higgs superfields, $(H_u H_d)$. Such terms induce the μ-term and B-parameter of the order of $F_X/\Lambda_{\rm IR}$ through the Giudice-Masiero mechanism [12]. Now, since the scale of the gravity mediation is warped down to the low scale, the "low scale gravity mediation" has been realized.

For completeness, here we present a concrete model of the hidden sector as an example. When we discuss the SUSY breaking mechanism in extra dimension models, the radion field is generally involved and a mechanism to stabilize the extra dimensional radius is strongly related to the SUSY breaking mediation. In the supersymmetric warped extra dimension scenario, several ways to stabilize the radius have been proposed [8,13–15]. A model proposed in [15] is remarkable for our aim, because the radius is stabilized in supersymmetric way in the model, and the resultant supersymmetric radion mass is so heavy that the radion potential is little affected by the SUSY breaking on a brane. Here, we assume such a radius stabilization mechanism by which the vacuum expectation value (VEV) of the radion is completely fixed almost independently of the SUSY breaking mechanism on a brane. Then, ω is dealt with as a constant in the following.

We present a simple Lagrangian for the chiral superfield (X) in the hidden sector on the IR brane such that

$$\int d^4\theta \omega^\dagger \omega X^\dagger X + \left(\int d^2\theta \omega^3 m^2 X + \text{H.c.} \right), \tag{2.10}$$

where m is a mass parameter. This is nothing but the Polonyi model [2] on the IR brane. Rescaling X to give the canonical Kahler potential, $X \to X/\omega$, we obtain the SUSY breaking (nonzero F-term of X) in four-dimensional effective theory,

$$F_X = (\omega m)^2. \tag{2.11}$$

The SUSY breaking scale is controlled by the mass parameter m accompanied by the warp factor ω as expected. Depending on the value of $\Lambda_{\rm IR}$, we take a suitable value for the parameter m in the superpotential so as to provide the typical soft mass scale around the electroweak scale. Only with the canonical Kahler potential, there is a pseudoflat direction in the scalar potential and VEV of X is undetermined. A simple way to lift up the pseudoflat direction is to introduce higher order terms in the Kahler potential. When we simply add a term, $-c(X^\dagger X)^2/M_5^2$, with a dimensionless coefficient $c > 0$, the potential minimum is realized at $\langle X \rangle = 0$. In this simple case, the mass of the hidden sector scalar is given by $m_X = 2\sqrt{c} F_X/\Lambda_{\rm IR}$, which is the same order of the soft SUSY breaking mass scale in the visible sector.

Vacuum energy (cosmological constant) in supergravity has two contributions: One is positive from the SUSY breaking and the other is negative from VEV of the superpotential which couples to the compensating multiplet. This negative contribution is the result from the fact that

[2] In general, we can introduce higher dimensional terms among X and Yukawa couplings in superpotential, which induce additional A parameters. Throughout the paper, we do not consider higher dimensional operators in superpotential, for simplicity.

the Kahler potential of the compensating multiplet has a wrong sign in Eq. (2.2). To obtain the vanishing (almost zero) cosmological constant, we simply put a constant superpotential on the UV brane, $W_{\rm UV}$. From Eqs. (2.2) and (2.3), the total vacuum energy is described as

$$E_{\rm vac} \simeq |F_X|^2 - 3\frac{|W_{\rm UV}|^2}{M_P^2} \simeq 0, \tag{2.12}$$

where $M_P^2 = M_5^3/k$ as mentioned above, and the constant superpotential $W_{\rm UV}$ has been tuned so as to cancel out the positive contribution from the SUSY breaking.

Gravity multiplet resides in the bulk, whose zero mode represents the gravity multiplet in effective four-dimensional supergravity. Since the gravity sector in effective four-dimensional theory should be reproduced correctly, we obtain the usual formula for the gravitino mass in four-dimensional supergravity, $m_{3/2} \simeq W_{\rm UV}/M_P^2$. Considering the condition of the vanishing cosmological constant, the gravitino mass is usually expressed as

$$m_{3/2} \simeq \frac{W_{\rm UV}}{M_P^2} \simeq \frac{F_X}{M_P}. \tag{2.13}$$

In our scenario, the scale of the gravity mediation is warped down and the typical soft mass scale is given by $\tilde{m} \simeq F_X/\Lambda_{\rm IR}$, so that the gravitino mass is further rewritten as

$$m_{3/2} \simeq \frac{F_X}{M_P} \simeq \tilde{m} \times \left(\frac{\Lambda_{\rm IR}}{M_P}\right). \tag{2.14}$$

Therefore, in the above setup, the gravitino is naturally the LSP, because of the suppression factor $\Lambda_{\rm IR}/M_P$. A similar result has been discussed in the flat extra dimension model [16], where $\Lambda_{\rm IR}$ is replaced by M_5 smaller than M_P. We can reproduce this result by setting $M_5 < M_P$ and taking the flat space-time limit $k \to 0$.

III. PHENOMENOLOGY OF LOW SCALE GRAVITY MEDIATION

As discussed in the previous section, the IR cutoff, $\Lambda_{\rm IR}$, is the model parameter, and we can take any values for it. Accordingly, the SUSY breaking scale should be suitably chosen so as to provide the typical soft mass scale around the electroweak scale. In this section, we address phenomenologies of the low scale gravity mediation scenario for various IR cutoff scales.

A. Phenomenology with the LSP gravitino

As shown in the previous section, the gravitino is naturally the LSP due to the suppression factor $\Lambda_{\rm IR}/M_P$ in Eq. (2.14), so that it can be a candidate of the dark matter in the present universe. Since there is no such a suppression factor in the conventional minimal supergravity scenario, the gravitino mass is normally of the same order of the typical soft mass scale and the gravitino is not so likely to be the LSP. Again, note that, in the warped extra dimension scenario, we can take any values for $\Lambda_{\rm IR}$ without serious fine-tuning among the original parameters in the gravity sector, M_5, k, and r. Therefore, we can consider a wide range of the LSP gravitino mass according to values of the warp factor. This feature is similar to the GMSB scenario, where the gravitino mass varies with the messenger scale. The crucial difference is that our model, the same as the minimal supergravity scenario, has more flexibility for sparticle mass spectrum than the one in the GMSB scenario.

Since couplings among the gravitino, particles, and sparticles in the MSSM are suppressed by the Planck mass, the gravitino cannot be in thermal equilibrium in the early universe. There are two generic ways through which LSP gravitinos are produced in the early universe. One is thermal production through scattering and decay processes of the MSSM particles in thermal plasma. In this case, the relic density of the gravitino is evaluated as [17]

$$\Omega^{\rm TP}h^2 \sim 0.2\left(\frac{T_R}{10^{10}\ {\rm GeV}}\right)\left(\frac{100\ {\rm GeV}}{m_{3/2}}\right)\left(\frac{M_3}{1\ {\rm TeV}}\right)^2, \tag{3.1}$$

where T_R is the reheating temperature after inflation (which should be smaller than $\Lambda_{\rm IR}$ due to theoretical consistency), and M_3 is the running gluino mass. The other is nonthermal production through the late time decay of a quasistable next LSP after its decoupling from the thermal plasma [6]. In this case, the relic density of the LSP gravitino is related to the relic density of the next LSP,

$$\Omega^{\rm NTP}h^2 = \frac{m_{3/2}}{M_{\rm NLSP}}\Omega_{\rm NLSP}h^2, \tag{3.2}$$

where $\Omega_{\rm NLSP}h^2$ would be the relic density of the next LSP if it were stable, and $M_{\rm NLSP}$ denotes its mass. By appropriately fixing the gravitino mass, the reheating temperature, and sparticle mass spectrum, the relic density suitable for the dark matter can be obtained. However, some cosmological constraints should be considered as well [18]. Since the gravitino couples to the MSSM particles very weakly, the next LSP decays into the LSP gravitino and standard model particles at late time. If it decays after big bang nucleosynthesis (BBN), its energetic daughters would destroy light nuclei through photodissociation and hadrodissociation and, as a result, upset the successful prediction of BBN. Furthermore, late time injection of energetic photons produced by the next LSP decay would distort the spectrum of the observed cosmic microwave background. These considerations will constrain the model parameters to consistently realize the gravitino dark matter scenario.

The quasistable next LSP opens up an interesting possibility in collider physics. The decay rate of the next LSP ($\tilde{\Psi}$) into a standard model particle (Ψ) and the LSP gravitino ($\psi_{3/2}$) is given by

$$\Gamma(\tilde{\Psi} \to \Psi \psi_{3/2}) = \frac{\kappa m_{\tilde{\Psi}}^5}{48\pi M_P^2 m_{3/2}^2}\left(1 - \frac{m_\Psi^2}{m_{\tilde{\Psi}}^2}\right)^4, \tag{3.3}$$

where $\kappa \sim 1$ is a model-dependent mixing parameter among superpartners and standard model particles. Thus, the lifetime of the next LSP is estimated as

$$\tau_{\tilde{\Psi}} \sim 10^8 \ \sec \times \left(\frac{100\ \text{GeV}}{m_{\tilde{\Psi}}}\right)^5 \left(\frac{m_{3/2}}{100\ \text{GeV}}\right)^2$$
$$\sim 10^8 \ \sec \times \left(\frac{100\ \text{GeV}}{\tilde{m}}\right)^3 \left(\frac{\Lambda_{\text{IR}}}{M_P}\right)^2. \tag{3.4}$$

Here, in the last equality, we have replaced the mass of the next LSP ($\tilde{m}_\Psi$) into the typical sparticle mass and used Eq. (2.14). If $\Lambda_{\text{IR}} \gg 10^{10}$ GeV, the decay length well exceeds the detector size of the LHC and the ILC, and the next LSP decay takes place outside the detector. In this case, there have been interesting proposals [7] for the way to trap quasistable next LSPs outside the detector, when the next LSP is a charged particle. Detailed studies of the next LSP decay may provide precise measurements of the gravitino mass and the four-dimensional Planck mass. On the other hand, if $\Lambda_{\text{IR}} \ll 10^{10}$ GeV, the next LSP decays within the detector. In the GMSB scenario where the next LSP can be neutralino and right-handed slepton [19], it has been pointed out [19,20] that the next LSP decay provides very characteristic SUSY signatures with leptons and/or photons accompanied by missing E_T. For detailed studies on general types of the next LSP, see Ref. [21]. Our model can naturally provide such general cases, since it has more flexibility for sparticle mass spectra than those in the GMSB scenario.

B. Collider phenomenology involving the hidden sector field

As mentioned above, we can take $\Lambda_{\text{IR}} = \mathcal{O}(1\ \text{TeV})$ without any serious fine-tuning for the original model parameters. If the effective cutoff scale is low enough, higher dimensional interactions suppressed by Λ_{IR} have an impact on collider physics. Note that the contact operators relevant to the gravity mediation also provide contact interactions between the hidden sector scalar field and the standard model fields. From the operator giving masses to gauginos in Eq. (2.6), we can extract interactions among the hidden sector scalar and the standard model gauge bosons such that

$$\mathcal{L}_{\text{int}} = -\frac{1}{4}\int d^2\theta c_a \frac{X}{\Lambda_{\text{IR}}} \mathcal{W}^{a\alpha}\mathcal{W}^a_\alpha$$
$$\supset -\frac{c_a}{4\sqrt{2}}\frac{\chi}{\Lambda_{\text{IR}}}\mathcal{F}^{a\mu\nu}\mathcal{F}^a_{\mu\nu} - \frac{c_a}{8\sqrt{2}}\frac{a}{\Lambda_{\text{IR}}}\mathcal{F}^{a\mu\nu}\tilde{\mathcal{F}}^a_{\mu\nu}, \tag{3.5}$$

where we have decomposed the hidden scalar field X into two real scalar fields, $X = (\chi + ia)/\sqrt{2}$, and $\mathcal{F}^a$ and $\tilde{\mathcal{F}}^a$ are the field strength and its dual of corresponding standard model gauge fields, respectively. In the case of $\langle X\rangle \ll \Lambda_{\text{IR}}$, the operators in Eq. (2.6) also give interactions between the hidden scalar and standard model fermions,

$$\mathcal{L}_{\text{int}} = \int d^4\theta c_A \frac{X + X^\dagger}{\Lambda_{\text{IR}}} Q_i^\dagger Q_i \supset \sqrt{2} c_A \frac{\chi}{\Lambda_{\text{IR}}} \mathcal{L}_{\text{kin}}^{\text{fermion}}, \tag{3.6}$$

where $\mathcal{L}_{\text{kin}}^{\text{fermion}}$ is the kinetic term for each standard model fermion.

Now we investigate collider phenomenologies involving the hidden scalar based on the above interactions. In the following, to make our discussion clear, we do not specify a concrete potential of the hidden sector fields, χ and a, so that their masses are dealt with as free parameters. Furthermore, we concentrate on the phenomenology involving only χ, the real part of X, for simplicity. The general case involving both χ and a can be investigated in the same way, and we will arrive at almost the same conclusions.

Let us begin with phenomenology at the LHC. If χ is light enough and Λ_{IR} are low enough, it may be possible to produce the hidden scalar at the collider through the interactions in Eq. (3.5). For $c_1 \simeq c_2 \simeq c_3 \simeq 1$, the dominant χ production process at the LHC is the gluon fusion process. Note that the dominant production process of the Higgs boson of the standard model is the same gluon fusion process through the effective interaction among the Higgs boson (h) and gluons induced by top quark one-loop diagram [22],

$$\mathcal{L}_{\text{eff}} = -\frac{\alpha_s}{16\pi} F_{1/2}(\tau_t)\frac{h}{v} G^{a\mu\nu}G^a_{\mu\nu}. \tag{3.7}$$

Here, $F_{1/2}$ is the form factor given as

$$F_{1/2} = -2\tau[1 + (1-\tau)f(\tau)], \tag{3.8}$$

where

$$f(\tau) = \begin{cases} [\sin^{-1}(1/\sqrt{\tau})]^2 & (\text{for } \tau \geq 1), \\ -\frac{1}{4}\left[\ln\left(\frac{1+\sqrt{1-\tau}}{1-\sqrt{1-\tau}}\right) - i\pi\right]^2 & (\text{for } \tau < 1), \end{cases} \tag{3.9}$$

$\tau_t = 4m_t^2/q^2$ with momentum transfer q^2 in the direction to the Higgs boson, and $v = 246$ GeV is VEV of Higgs field. Interestingly, the effective interaction is of the same form as the one in Eq. (3.5). Therefore, in the production process, the hidden scalar χ behaves like the Higgs boson in the standard model. Comparing coefficients of their interactions, we find that their production cross sections become comparable for $\Lambda_{\text{IR}} \simeq 10$ TeV, assuming the same masses for them.

When we consider the decay process of χ, we find a big difference between χ and the Higgs boson. From Eqs. (3.5) and (3.6), the partial decay width of χ into the standard model gauge bosons and fermions is easily calculated. The decay width into a pair of gauge bosons is found to be

$$\Gamma(\chi \to gg) = \frac{c_3^2}{16\pi} \frac{m_\chi^3}{\Lambda_{\rm IR}^2},$$

$$\Gamma(\chi \to \gamma\gamma) = \frac{(c_1 \cos^2\theta_w + c_2 \sin^2\theta_w)^2}{128\pi} \frac{m_\chi^3}{\Lambda_{\rm IR}^2},$$

$$\Gamma(\chi \to ZZ) = \frac{(c_1 \sin^2\theta_w + c_2 \cos^2\theta_w)^2}{1024\pi} \frac{m_\chi^3}{\Lambda_{\rm IR}^2} \times \beta_Z(3 + 2\beta_Z^2 + 3\beta_Z^4),$$

$$\Gamma(\chi \to WW) = \frac{c_2^2}{512\pi} \frac{m_\chi^3}{\Lambda_{\rm IR}^2} \beta_W(3 + 2\beta_W^2 + 3\beta_W^4),$$

$$\Gamma(\chi \to \gamma Z) = \frac{(c_1 - c_2)^2 \sin^2\theta_w \cos^2\theta_w}{64\pi} \frac{m_\chi^3}{\Lambda_{\rm IR}^2}\left(1 - \frac{m_Z^2}{m_\chi^2}\right)^3, \tag{3.10}$$

where m_χ is the mass of the hidden scalar χ, θ_w is the weak mixing angle, $\beta_Z = \sqrt{1 - 4(m_Z/m_\chi)^2}$, and $\beta_W = \sqrt{1 - 4(m_W/m_\chi)^2}$. The interaction of Eq. (3.6) gives the partial decay width into a fermion pair,

$$\Gamma(\chi \to f\bar{f}) = \frac{c_A^2}{4\pi} \frac{m_f^2 m_\chi}{\Lambda_{\rm IR}^2} \beta_f^3 \times N_c, \tag{3.11}$$

where m_f is the mass of the final state fermions, $\beta_f = \sqrt{1 - 4(m_f/m_\chi)^2}$, and N_c is the color factor for the final state fermions. Since fermions couple with χ through their kinetic terms, the decay width is proportional to m_f^2, so that decay channels into light fermions are very much suppressed compared to those into gauge boson pairs. This result contrasts with the fact that the dominant decay channel of the light Higgs boson with mass $m_h < 2m_W$ is into bottom and antibottom quarks, since the Higgs boson decay into gauge bosons occurs through one-loop radiative corrections.

We show the branching ratio of the χ decay in Fig. 1. Here, we have considered only the decay channels into the standard model particles, assuming that all the sparticles and Higgs bosons in the MSSM are heavier than χ. If χ is heavy enough, it can decay into sparticle pairs and Higgs boson pairs. Their interactions are found to be similar to Eq. (3.6), and the partial decay width into sparticle and Higgs boson pairs is proportional to the mass of the final states. We see that the branching ratio of the χ decay is quite different from that of the Higgs boson. In particular, the branching ratio of $\chi \to \gamma\gamma$ is large, $\text{Br}(\chi \to \gamma\gamma) \simeq 0.1$. On the other hand, the branching ratio of the Higgs boson into two photons is at most 10^{-3}, even when the Higgs mass is light $m_h < 2m_W$. This fact implies that once χ is produced at the LHC, the signature of χ is distinguishable from the Higgs boson one.

In the MSSM, the lightest Higgs boson is like the standard model Higgs boson and its mass is too light to

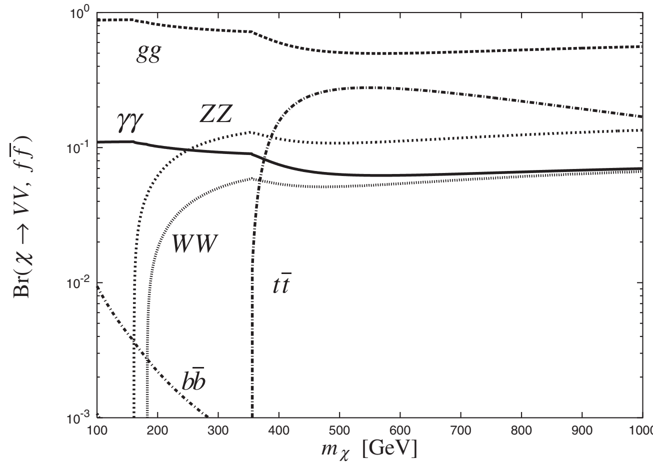

FIG. 1. The branching ratio of the hidden scalar (χ) as a function of its mass m_χ for $c_1 = c_2 = c_3 = c_A = 1$. The plot on the lower-left corner corresponds to $\text{Br}(\chi \to \tau\bar{\tau})$.

decay into weak gauge boson pairs. The most important channel for the lightest Higgs boson search at the LHC is its decay process into two photons. Therefore, the χ production and its decay process into two photons have a great impact on the (lightest SUSY) Higgs boson search at the LHC. To see this, let us evaluate a ratio between two photon events from the χ decay and the Higgs boson decay. The ratio of χ and the Higgs boson production rates can be estimated from the ratio between the coefficients of the $\mathcal{F}^{a\mu\nu}\mathcal{F}^a_{\mu\nu}$ terms in Eqs. (3.5) and (3.7). For $m_\chi = m_h = 120$ GeV and $c_1 = c_2 = c_3 = c_A = 1$, the event number ratio as a function of the effective cutoff scale is depicted in Fig. 2. We can see that, for $\Lambda_{\rm IR} = 10$ TeV, the number of events from χ production is 2 orders of magnitude larger than that from the Higgs boson production. Even for $\Lambda_{\rm IR} = \mathcal{O}(100$ TeV), the ratio is still of order one. Therefore, if

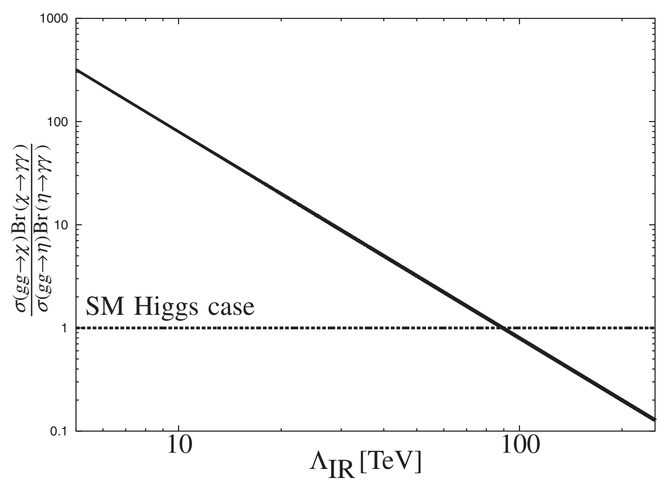

FIG. 2. The ratio between two photon events from the χ production and the Higgs boson production at the LHC, as a function of $\Lambda_{\rm IR}$, for $c_1 = c_2 = c_3 = c_A = 1$ and $m_\chi = m_h = 120$ GeV.

$\Lambda_{\rm IR}$ is around 10 TeV, the hidden sector scalar χ can be discovered at the LHC with a very clean signature.

We discuss more on interesting features of the low scale gravity mediation scenario. Note that there is one-to-one correspondence between gaugino masses and the partial decay width of the hidden sector scalar into gauge boson pairs, because they are originated from the same contact operators. Considering that the quantity M_a/α_a is invariant under renormalization group equations [23], the ratio between gaugino masses at the typical soft mass scale $\tilde{m}$ is given by

$$M_1{:}M_2{:}M_3 = c_1\frac{\alpha_1(\tilde{m})}{\alpha_1(\Lambda_{\rm IR})}{:}c_2\frac{\alpha_2(\tilde{m})}{\alpha_2(\Lambda_{\rm IR})}{:}c_3\frac{\alpha_3(\tilde{m})}{\alpha_3(\Lambda_{\rm IR})}, \tag{3.12}$$

which is determined by the ratio between c_a. As shown in Eq. (3.10), the ratio between the partial decay width into pairs of gauge bosons is also fixed by the ratio between c_a. Therefore, once gauginos and the hidden sector scalar are discovered at future colliders and their masses and the partial decay width of the hidden scalar are precisely measured, we can check the origin of SUSY breaking mediation by examining this one-to-one correspondence.

Finally, let us investigate phenomenology at the ILC. The ILC, the liner e^+e^- collider, is the so-called Higgs boson factory where a large number of Higgs bosons will be produced. The most clean channel of the Higgs boson production at the ILC is the associated Higgs production (Higgsstrahlung production), $e^+e^- \to Zh$, through the standard model interaction $\mathcal{L}_{\rm int} = (m_Z^2/v)hZ^\mu Z_\mu$. Since the hidden sector scalar χ has the vertex among Z-boson in Eq. (3.5), we can consider the same associated process for the χ production at the ILC.[3] In the case of the universal couplings, $c_1 = c_2 = c_3$, for simplicity, the cross section of the process $e^+e^- \to Z\chi$ is found to be[4]

$$\begin{aligned}\frac{d\sigma}{d\cos\theta}(e^+e^- \to Z\chi) &= \frac{1}{64\pi s}\sqrt{\frac{E_Z^2 - m_Z^2}{s}}\frac{c_2^2}{2}\left(\frac{e}{\sin\theta_w\cos\theta_w}\right)^2 \\ &\quad\times (g_L^2 + g_R^2)\left(\frac{s}{s - m_Z^2}\right)^2 \\ &\quad\times \frac{E_Z^2}{\Lambda_{\rm IR}^2}\left(1 + \cos^2\theta + \frac{m_Z^2}{E_Z^2}\sin^2\theta\right),\end{aligned} \tag{3.13}$$

where $\cos\theta$ is the scattering angle of the final state Z-boson, $g_L = -1/2 + \sin^2\theta_w$, $g_R = \sin^2\theta_w$, and $E_Z = \frac{\sqrt{s}}{2}[1 + (m_Z^2 - m_\chi^2)/s]$. Since Higgs boson couples to a pair of Z-bosons at tree level, its production cross section is mostly larger than the one of χ production. In Fig. 3, we show the ratio of the total cross sections between χ and Higgs boson productions as a function of $\Lambda_{\rm IR}$ at the ILC with the collider energy $\sqrt{s} = 1$ TeV. The ratio, $\sigma(e^+e^- \to Z\chi)/\sigma(e^+e^- \to Zh)$, becomes one for $\Lambda_{\rm IR} \simeq 1.3$ TeV, and it decreases proportionally to $1/\Lambda_{\rm IR}^2$.

The coupling manner among χ and the Z-boson pair is different from that of the Higgs boson, and this fact reflects into the difference of the angular distribution of the final state Z-boson. In the high energy limit, we find $\frac{d\sigma}{d\cos\theta} \times (e^+e^- \to Z\chi) \propto 1 + \cos^2\theta$, while $\frac{d\sigma}{d\cos\theta}(e^+e^- \to Zh) \propto 1 - \cos^2\theta$. Figure 4 shows the angular distributions of

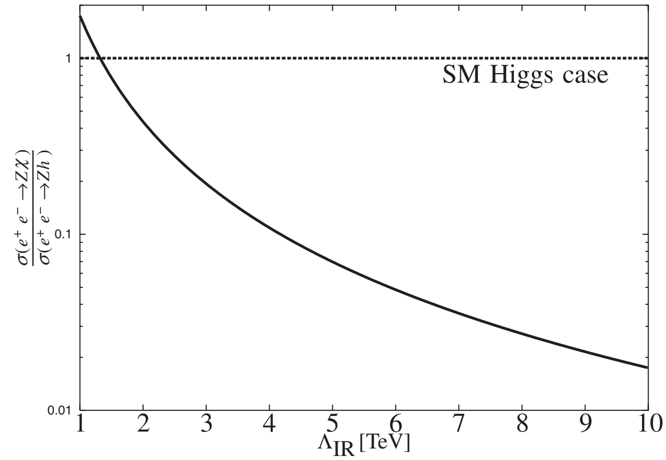

FIG. 3. The ratio of total cross sections between the associated χ and Higgs productions as a function of $\Lambda_{\rm IR}$, at the ILC with the collider energy $\sqrt{s} = 1$ TeV. Here, we have fixed the parameters such as $m_\chi = m_h = 120$ GeV and $c_1 = c_2 = c_3 = c_A = 1$. The ratio becomes one for $\Lambda_{\rm IR} \simeq 1.3$ TeV.

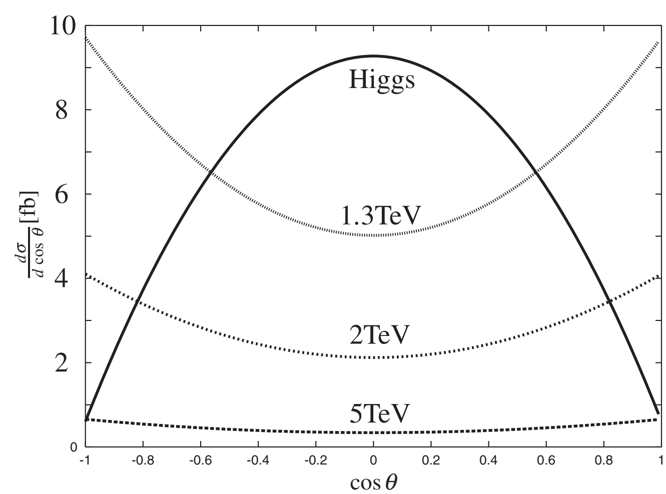

FIG. 4. The angular dependence of the cross sections for $m_\chi = m_h = 120$ GeV and $c_1 = c_2 = c_3 = c_A = 1$, at the ILC with the collider energy $\sqrt{s} = 1$ TeV. The standard model Higgs boson case is depicted as the solid line, while the others correspond to the χ productions with $\Lambda_{\rm IR} = 1.3$, 2, and 5 TeV, respectively, from above.

[3]For χ productions, we can also consider the process $e^+e^- \to \gamma\chi$ as one of main production processes, while such a process is negligible for the Higgs boson production. Studies on this process itself would be interesting.

[4]In the general case $c_1 \neq c_2$, the process $e^+e^- \to \gamma^* \to Z\chi$ should be included.

the associated χ and Higgs boson productions, respectively. Even if $m_\chi = m_h$ and the cross sections of χ and Higgs boson productions are comparable, the angular dependence of the cross section can distinguish the χ production from the Higgs boson one. Of course, detecting two photons from the χ decay with the sizable branching ratio $\mathrm{Br}(\chi \to \gamma\gamma) \sim 0.1$ would be an easy way to distinguish χ from Higgs boson, as discussed in the case of the LHC.

IV. CONCLUSIONS AND DISCUSSIONS

We have proposed the low scale gravity mediation scenario with the warped extra dimension. The setup of the scenario is that both of the hidden and visible sectors coexist on the IR brane. This setup is the same as in the four-dimensional minimal supergravity scenario except for the gravity multiplet residing in the bulk. We have considered the gravity mediated SUSY breaking through the contact operators between the hidden and the visible sector superfields. The crucial point is that the effective cutoff on the IR brane is warped down, so that the gravity mediation takes place at low energies. As a result, the gravitino is naturally the LSP, just as in the GMSB scenario. However, our gravity mediated scenario has more flexibility for sparticle mass spectra than those in the GMSB scenario. We have briefly discussed phenomenologies related to the LSP gravitino scenario.

If the effective cutoff scale is low enough, for example, $\Lambda_{\mathrm{IR}} = \mathcal{O}(10\ \mathrm{TeV})$, our scenario provides interesting phenomenologies at the future colliders. The contact operators relevant to the gravity mediation also provide contact interactions among the hidden sector scalars and the standard model particles. We have investigated collider physics involving the hidden sector scalar fields at the LHC and the ILC. Interestingly, the hidden sector scalars behave like the standard model Higgs boson in their production processes and, therefore, the existence of such scalars has a great impact on the Higgs boson search at the colliders. However, the decay process of the hidden scalars is quite different from the Higgs boson one, and, once produced, they will provide us with a very clean signature. The hidden sector may be no longer hidden.

Several discussions are in order.

In this paper, we have concentrated our discussion only on the contact operators relevant to the gravity mediation. In general, we may introduce contact operators among the visible sector fields themselves, which induce contact interactions among the standard model particles. For such contact interactions, the lower bound on Λ_{IR} by the current experiments should be taken into account. The electroweak precision measurements give the lower bound, $\Lambda_{\mathrm{IR}} \gtrsim 5\ \mathrm{TeV}$ [24]. If contact operators which cause FCNC processes are considered, rough estimation gives a more severe bound, $\Lambda_{\mathrm{IR}} \gtrsim 100\ \mathrm{TeV}$. We may expect that the severely constrained operators are forbidden by some underlying flavor symmetry which justifies the minimal ansatz.

Next, if $\Lambda_{\mathrm{IR}} = \mathcal{O}(10\ \mathrm{TeV})$, in other words, $\omega \sim 10^{-14}$, taken as in the previous section, the gravitino mass becomes too small, $m_{3/2} \sim 10^{-3}$–10^{-4} eV, to account for the dark matter density in the present universe. In this case, our model must be extended so as to implement a suitable candidate for the cold dark matter. Among various possibilities, we notice that, in extra dimensional models, there is more flexibility for the scale of the gravitino mass. In fact, as discussed in a series of papers [13,14,25,26], it is generally possible for the gravitino mass to be even the Planck scale. An important feature is that the gravity multiplet residing in the bulk couples to fields on both branes. Thus, when we introduce an additional hidden sector on the UV brane, the gravitino directly picks up the SUSY breaking on the UV brane and becomes massive. If the SUSY breaking scale is much larger than the effective SUSY breaking scale on the IR brane, total vacuum energy in Eq. (2.12) is replaced into

$$E_{\mathrm{vac}} \simeq |F_Y|^2 - 3\frac{|W_{\mathrm{UV}}|^2}{M_P^2} \simeq 0, \tag{4.1}$$

where F_Y is the large SUSY breaking on the UV brane, so that the gravitino mass is dominantly induced from this SUSY breaking, $m_{3/2} \sim F_Y/M_P \gg F_X/M_P$. On the other hand, the visible sector residing on the IR brane cannot directly feel the SUSY breaking on the UV brane, because two branes are spatially separated.

We must consider some possibilities on the SUSY breaking mediation from the UV brane to the IR brane over the bulk space. One is due to quantum corrections through the supergravity multiplet in the bulk, whose contribution is evaluated as [27]

$$\Delta\tilde{m} \sim m_{3/2}\omega^2. \tag{4.2}$$

Even if $m_{3/2} \sim M_P$, this is negligible compared to the low scale gravity mediation on the IR brane with the strong warp factor, $\omega \sim 10^{-14}$. In supergravity, the SUSY breaking mediation through the superconformal anomaly [28] (anomaly mediation) always exists. In warped extra dimensional models, the anomaly mediation contribution on the IR brane can be characterized by [25]

$$\Delta\tilde{m}_{\mathrm{AMSB}} \sim \frac{F_\omega}{\omega}. \tag{4.3}$$

This contribution highly depends on a mechanism to stabilize the fifth dimension. For example, in models of the radius stabilization proposed in [13,14], we obtain $\Delta\tilde{m}_{\mathrm{AMSB}} \sim m_{3/2}\omega^n$ with a model parameter n of order one. The setup of these models is that the visible sector resides on the IR brane while the hidden sector resides only on the UV brane.[5] Thus, it is easy to combine our model with these models by introducing the hidden sector also on the IR brane. When we fix the model parameter appropriately, $n > 1$, for example, we can realize the situation that the gravity mediation on the IR brane gives the dominant contribution to the soft SUSY breaking parameters. In this case, the gravitino is much heavier than sparticles in the MSSM, and the lightest neutralino can be the LSP and the candidate of the cold dark matter as usual.

Taking the flat space-time limit, $k \to 0$, in our model, we obtain the effective cutoff scale as $\Lambda_{\rm IR} = M_5$. If we take M_5 to be much smaller than the four-dimensional Planck scale, we can realize the low scale gravity mediation without the warp factor. However, in this case, the low scale cutoff, $M_5 \ll M_P$, implies $1/r \ll M_5$ in order to correctly reproduce the four-dimensional Planck scale through the relation, $M_P^2 \sim M_5^3 r$. Thus, one may claim a hierarchy problem between $1/r \ll M_5$. In the warped extra dimension scenario, there is no such a hierarchy problem, thanks to the warp factor. In the following, we will show that there exists a theoretical lower bound on M_5 in the flat extra dimension scenario, even if we admit the hierarchy between $1/r \ll M_5$.

Using the condition of the vanishing cosmological constant in Eq. (2.12) and the typical soft mass scale $\tilde{m} \sim F_X/M_5$, we obtain the relation[6]

$$W_{\rm UV} \sim F_X M_P \sim \tilde{m} M_5 M_P. \tag{4.4}$$

Since M_5 is the cutoff scale of the original theory, the theoretical consistency, $W_{\rm UV} \leq M_5^3$, implies the lower bound on the scale M_5 such as

$$M_5 \geq \sqrt{\tilde{m} M_P} \sim 10^{10}\ \text{GeV} \tag{4.5}$$

for $\tilde{m} \simeq 100$ GeV. Therefore, we cannot take M_5 as low as 1 TeV. In the warped extra dimension models, the above condition is replaced by $M_5 \geq \sqrt{\omega \tilde{m} M_P}$. For any M_5 satisfying this condition, we can realize the effective low scale cutoff by the warp factor, $\Lambda_{\rm IR} = \omega M_5$. There is no lower bound on the effective cutoff in the theoretical point of view.

Finally, let us consider an issue related to the scale of the SUSY breaking. We can express the SUSY breaking scale in terms of the typical soft mass scale and the effective cutoff, $\sqrt{F_X} \sim \sqrt{\tilde{m}\Lambda_{\rm IR}}$. When the scale of the SUSY breaking mediation, $\Lambda_{\rm IR}$, is very high, for example, $\Lambda_{\rm IR} = M_P$ in the usual minimal supergravity scenario, the hierarchy between $\sqrt{F_X} \ll \Lambda_{\rm IR}$ is necessary to provide soft SUSY breaking masses around the electroweak scale. How to generate such a hierarchy could be an important issue when one constructs a concrete SUSY breaking model. Dynamical SUSY breaking [29] is a remarkable possibility, in which the SUSY breaking scale is controlled by the dynamical scale of some strong interaction induced through the dimensional transmutation. Thus, there is no problem on the hierarchy.

In the warped extra dimension scenario, any original dimensional parameters on the IR brane are warped down according to their mass dimensions such as

$$M_5 \to \omega M_5 = \Lambda_{\rm IR}, \qquad \sqrt{F_X} \to \omega \sqrt{F_X}. \tag{4.6}$$

As discussed before, we can realize, for example, $\Lambda_{\rm IR} \sim 10$ TeV only with the mild hierarchy, $M_5 \sim k \sim M_P$ and $1/r \sim 0.1 M_P$. In the same way, if we introduce a mild hierarchy for the original SUSY breaking scale, $\sqrt{F_X} \sim 0.1 M_P$, we obtain the effective SUSY breaking scale such as $\sqrt{F_X} \sim 0.1 M_P \to \sqrt{F_X} \sim 1$ TeV. Then, the typical soft SUSY breaking mass scale appears around the electroweak scale, $\tilde{m} \sim 100$ GeV. This result implies that, in order to provide the correct electroweak scale, we do not need to introduce any additional mass scales except for the four-dimensional Planck scale. For the μ-parameter in the Higgs sector of the MSSM, we can follow the same manner. When a mildly hierarchical μ-parameter such as $\mu \sim 0.1 M_P$ is introduced in the original superpotential on the IR brane, it becomes a suitable scale, $\mu \sim 1$ TeV, in effective four-dimensional theory. Taking $\Lambda_{\rm IR} \leq 10$ TeV can make everything go well only with the mild hierarchy, and so it would be the most natural setting.

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[5]To be precise, SUSY on the UV brane is explicitly broken, nevertheless we obtain softly broken SUSY theory on the IR brane. This is a scenario proposed in [13,14], "emergent supersymmetry."

[6]In the flat extra dimensional scenario, there is no difference between superpotentials on the IR and UV brane, since the AdS curvature is zero and, thus, $\omega = \phi$.

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