

Dynamical $U(1)_R$ breaking in the metastable vacua

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ABSTRACT: In the Intriligator-Seiberg-Shih model, we parametrize spontaneous breaking of $U(1)_R$ symmetry with two gauge singlets with R-charges 1 and -1 . These singlets can play the role of the messengers. The messenger scale is dynamically generated, and hence there is no hierarchy problem between the supersymmetry breaking scale and the messenger scale. In the gauge mediation scenario, supersymmetry breaking scale turns out to be around $\mathcal{O}(10^6)$ GeV.

KEYWORDS: Supersymmetry Breaking, Supersymmetry Phenomenology.

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1. Introduction

Supersymmetry (SUSY) is still considered to be one of the best candidates for the solution of the gauge hierarchy problem. It explains the Higgs mass problem elegantly, but the question “Why is the electroweak scale so much smaller compared to the Planck scale?” needs a judicious setup for SUSY breaking. About thirty years ago Witten studied this question and suggested that the discrepancy between two scales might be understood if SUSY is broken dynamically [1]. Pre-2006 studies along this line have not shown any compelling model applicable to the real world phenomenology [2–7]. In some of these approaches there were SUSY preserving vacua by the interaction between the messenger sector and the SUSY breaking sector. Thus, these models already had the feature of metastable SUSY breaking vacua [3–5]. The idea of metastable SUSY breaking false vacua with SUSY preserving true minimum only in the SUSY breaking sector was introduced about ten years ago [8–10]. Recently, Intriligator, Seiberg, and Shih (ISS) gave an explanation that the metastability of the supersymmetry breaking vacua could be generic in SUSY gauge theories. Therefore, this argument could simplify and enrich the realm of the model building possibilities [11]. This work has especially attractive features in that it can be accommodated in string theory [13–15] and one can obtain low energy supersymmetry breaking.

Independently of these arguments on SUSY breaking, $U(1)_R$ symmetry breaking showed an important role on SUSY breaking itself. The origin of $U(1)_R$ symmetry is related to the rotation in the superspace which is defined by the fermionic coordinates,

$$U(1)_R : \theta \rightarrow e^{-i\alpha}\theta. \tag{1.1}$$

If we consider gravitational effect, $U(1)_R$ symmetry is explicitly broken by gravitational effect. We don’t, however, consider this case here. Early in the 1990s Nelson and Seiberg

showed that $U(1)_R$ symmetry is necessary for breaking SUSY spontaneously in models with generic and calculable superpotentials [16]. Therefore, the effect of $U(1)_R$ symmetry is protecting the SUSY breaking minimum from being invaded by the SUSY preserving minima. This pattern is also applicable to the case of the meta-stable vacua. Instead of the exact $U(1)_R$ symmetry, in the meta-stable vacuum case there exists an accidental and approximate one near the origin in the field configuration space. The effect of this symmetry is again to protect the meta-stable vacuum. In any case, it is necessary to have the $U(1)_R$ symmetry to obtain a stable SUSY breaking global or local minimum. Nevertheless, the $U(1)_R$ symmetry should be broken to obtain the soft gaugino masses and to make the R-axion heavy [3, 17, 9].

Related to the ISS model, many studies have been performed to break the accidental $U(1)_R$ symmetry. Some studies were done in direct mediation setup with an explicit R-symmetry breaking term [19, 21]. There have been studies on spontaneous R-symmetry breaking by introducing $U(1)$ gauge interaction in the direct gauge mediation scheme [21], but these studies have a problem with the Landau pole. The direct mediation scheme in the ISS scheme has problems because the ISS model has too many massive flavors. Some methods employed to avoid the Landau pole problem lead to too high SUSY breaking scale [19]. On the other hand, the Landau pole is inescapable if the SUSY is broken at the low scale [21]. Therefore, the ordinary gauge mediation is more persuasive in this sense. Some studies related to the R-symmetry breaking in the ISS setup introduce an explicit breaking term, which comes from the interaction between the messengers and the Goldstino super multiplet in the ordinary gauge mediation [20, 22], but the messenger masses are introduced by hand. Thus, the hierarchy between the SUSY breaking scale and the messenger scale remains unsolved.

In this paper, we study dynamical spontaneous $U(1)_R$ symmetry breaking in a modified ISS set up without the help of any explicit $U(1)_R$ breaking terms. We introduce two singlets with R-charges 1 and -1 , respectively, to keep the superpotential invariant under the $U(1)_R$ symmetry. By doing this, we can avoid the two superficial dilemma: obtaining $U(1)_R$ symmetry and breaking it. We find the $U(1)_R$ and SUSY breaking meta-stable minima without fine-tuning. Since the original work of ISS is related to the SQCD, we do not want to harm the good property that SUSY is restored at high energy scale by the dynamical interaction [11]. We find that our vacua are near the origin in the field space. It means that we can keep the strong points of the ISS model. The vacua can survive the transition to the true SUSY preserving vacua as long as its lifetime which is much longer than the age of our universe. Through this study, we obtain the dynamically generated messenger scale. At the same time, this scale is the SUSY breaking scale, and hence we can build a realistic model without introducing the messenger scale by hand. In addition, the new singlet fields can be used as the messenger fields in the gauge mediation scenario.

In section 2, we briefly review the ISS model, and consider the $U(1)_R$ symmetry breaking in the O’Raifeartaigh type SUSY breaking. In section 3, we introduce the model. In section 4, phenomenological implications are commented.

2. The ISS model and $U(1)_R$ symmetry

The O’Raifeartaigh model is the simplest mechanism breaking SUSY spontaneously [25]. The basic setup of the ISS model is related to this mechanism. In the early 1990s, the holomorphic property of SUSY gauge theory together with the global symmetry properties showed the good vacuum properties, such as the moduli space structure and the duality between the electric theory and its magnetic dual theory, which ordinary gauge theories do not possess [26–28]. SUSY QCD (SQCD) was classified by the number of flavors N_f and the number of color N_c . ISS started with the magnetic dual gauge theory which is infrared free, for which the electric theory is asymptotically free. This happens for $N_c < N_f < \frac{3}{2}N_c$.

The field contents in the magnetic dual theory are,

	Φ	φ	$\tilde{\varphi}$
SU(N)	1	\square	$\overline{\square}$
SU(N_f)	adj + 1	\square	$\overline{\square}$
U(1) _R	2	0	0

(2.1)

where $N = N_f - N_c$. Then the superpotential consistent with the symmetry is given by

$$W = h\text{Tr}\tilde{\varphi}\Phi\varphi - h\mu^2\text{Tr}\Phi. \tag{2.2}$$

If we assume that there is a kind of strongly coupled interaction, then we can make μ small enough by retrofitting as

$$\mu = \frac{\Lambda_s^3}{M_P^2}, \tag{2.3}$$

where Λ_s is the confining scale of a strong dynamics [29]. Then SUSY is broken via the incompatibility among the F-flat conditions, i.e. by the rank condition, at the origin of the field space. The tree level potential is

$$V_{\text{cl}} = (N_f - N)|h^2\mu^4| \tag{2.4}$$

with the flat direction given by

$$\Phi = \begin{pmatrix} 0 & 0 \\ 0 & X \end{pmatrix}, \quad \varphi = \begin{pmatrix} \varphi_0 \\ 0 \end{pmatrix}, \quad \tilde{\varphi}_0^T = \begin{pmatrix} \tilde{\varphi}_0 \\ 0 \end{pmatrix} \tag{2.5}$$

with $\tilde{\varphi}_0^T\varphi_0 = \mu^2\mathbf{1}_N$.

If we turn on the SU(N) which is IR free, there exists a scale Λ_m , above which the theory turns to be strongly coupled theory. The holomorphic gauge coupling of SU(N) is given by

$$e^{-8\pi^2/g^2(E)+i\theta} = \left(\frac{E}{\Lambda_m}\right)^{N_f-3N_c}. \tag{2.6}$$

Now we consider quantum correction to the potential, then the theory turns out to maintain the supersymmetry breaking feature at the origin even after we introduce the gauge

interaction. However, the gauge interaction plays a crucial role in the other place of the field configuration space. By non-perturbative effect we get the effective potential as

$$W_{\text{low}} = N(h^{N_f} \Lambda_m^{-(N_f-3N)} \det \Phi)^{1/N} - h\mu^2 \text{Tr} \Phi. \quad (2.7)$$

Next, by investigating the F-flat directions, we obtain a result that there exist supersymmetry preserving vacua, i.e. all the F terms vanish, with the value of

$$\langle h\Phi \rangle = \Lambda_m \epsilon^{2N/(N_f-N)} \mathbf{1}_{N_f} = \mu \frac{1}{\epsilon^{(N_f-3N)/(N_f-N)}} \mathbf{1}_{N_f}, \quad \text{where } \epsilon \equiv \frac{\mu}{\Lambda_m}. \quad (2.8)$$

The longevity of the metastable vacua is guaranteed by this inequality for $\epsilon \ll 1$,

$$|\mu| \ll |\langle h\Phi \rangle| \ll |\Lambda_m|. \quad (2.9)$$

If we return to the origin where the gauge interaction can be ignored, we find that there exists an accidental R-symmetry induced by quantum corrections. The existence of supersymmetry preserving vacua means that there is no exact $U(1)_R$ symmetry.

Nelson and Seiberg showed that generic and calculable models need $U(1)_R$ symmetry to break SUSY spontaneously [16]. However, SUSY is not spontaneously broken in spite of $U(1)_R$ when R-charges of all matter fields are either 2 or 0. This implies that R-charges of field contents are closely related to the SUSY breaking as mentioned above. Recently, Shih gave an explanation for the R-symmetry breaking, which is just a necessary condition for $U(1)_R$ symmetry breaking minimum in O’Raifeartaigh type models. Including the radiative corrections as in the Coleman-Weinberg potential, the pseudo-moduli gain masses at the origin in generic cases. He found that introducing some fields with R-charges except 0 and 2 can make the masses of the moduli negative by quantum corrections, but this depends on the condition of the parameter space. The negative masses of moduli mean that $U(1)_R$ symmetry is spontaneously broken via quantum corrections. Therefore, he led the conclusion that *in O’Raifeartaigh type SUSY breaking models, it is necessary to have matter fields of which R-charges are given neither 2 nor 0 for spontaneous $U(1)_R$ breaking* [31].

3. Spontaneous breaking of $U(1)_R$ in a modified ISS model

3.1 Model

In the O’Raifeartaigh-type models where all fields have R-charges of 0 or 2 only, the $U(1)_R$ symmetry is not spontaneously broken. Therefore, in order for $U(1)_R$ to be spontaneously broken, there has to be at least one field in the model with R-charge different from 0 and 2. [31]

Thus, we introduce new fields, A and B , with R-charge different from 0 and 2 to the original ISS type model. The matter fields we introduce are for $N_f > N$:

	Φ	φ	$\tilde{\varphi}$	A	B
$SU(N)$	1	\square	$\overline{\square}$	1	1
$SU(N_f)$	adj + 1	\square	$\overline{\square}$	1	1
$U(1)_R$	2	0	0	1	-1

(3.1)

Taking the canonical Kähler potential, the generic tree-level superpotential is

$$W = h\text{Tr}\tilde{\varphi}\Phi\varphi - h\mu^2\text{Tr}\Phi + \lambda AB\text{Tr}\Phi + mA^2 \quad (3.2)$$

which respects the $U(1)_R$ symmetry. The classical moduli space is obtained from

$$\begin{aligned} \frac{\partial W}{\partial A} = \lambda B\text{Tr}\Phi + 2mA, \quad \frac{\partial W}{\partial B} = \lambda A\text{Tr}\Phi, \quad \frac{\partial W}{\partial \Phi_{ij}} = h\tilde{\varphi}^i\varphi^j + (\lambda AB - h\mu^2)\delta^{ij}, \\ \frac{\partial W}{\partial \varphi} = h\tilde{\varphi}\Phi, \quad \frac{\partial W}{\partial \tilde{\varphi}} = h\Phi\varphi. \end{aligned} \quad (3.3)$$

The vacua along the classical moduli space are

$$A = 0, \quad B = 0, \quad \Phi = \begin{pmatrix} 0 & 0 \\ 0 & X \end{pmatrix}, \quad \varphi = \begin{pmatrix} \varphi_0 \\ 0 \end{pmatrix}, \quad \tilde{\varphi}_0^T = \begin{pmatrix} \tilde{\varphi}_0 \\ 0 \end{pmatrix} \quad (3.4)$$

or

$$A = 0, \quad B = \text{arbitrary}, \quad \Phi = \begin{pmatrix} 0 & 0 \\ 0 & X \end{pmatrix}, \quad \varphi = \begin{pmatrix} \varphi_0 \\ 0 \end{pmatrix}, \quad \tilde{\varphi}_0^T = \begin{pmatrix} \tilde{\varphi}_0 \\ 0 \end{pmatrix} \quad \text{where } \text{Tr}X = 0 \quad (3.5)$$

with $\tilde{\varphi}_0^T\varphi_0 = \mu^2\mathbf{1}_N$. Then, the classical scalar potential is given by

$$V = (N_f - N) |h\mu^2|^2. \quad (3.6)$$

This model does not have a SUSY ground state because it is impossible for F_A , F_B , and F_Φ terms to vanish simultaneously. However, at

$$A = \frac{h\mu^2}{\lambda} \frac{1}{B}, \quad \text{Tr}\Phi = 0, \quad \varphi = 0, \quad \tilde{\varphi} = 0, \quad (3.7)$$

the classical potential has a runaway direction toward $B \rightarrow \infty$,

$$V = 4|mA|^2 = 4 \left| \frac{h\mu^2 m}{\lambda B} \right|^2 \rightarrow 0. \quad (3.8)$$

No static vacuum exists along this direction, and SUSY is asymptotically restored as $B \rightarrow \infty$.

3.2 One-loop lifting of pseudo-moduli

The minima of the tree-level scalar potential considered here occur along the pseudo-moduli space (3.4). Expanding around the above classical moduli space (3.4),¹

$$\begin{aligned} \Phi = \begin{pmatrix} \delta\Phi_{11} & \delta\Phi_{12} \\ \delta\Phi_{21} & X + \delta\Phi_{22} \end{pmatrix}, \quad \varphi = \begin{pmatrix} \mu Y + \delta\varphi_1 \\ \delta\varphi_2 \end{pmatrix}, \quad \tilde{\varphi}_0^T = \begin{pmatrix} \mu/Y + \delta\tilde{\varphi}_1 \\ \delta\tilde{\varphi}_2 \end{pmatrix} \\ A = \delta A, \quad B = \delta B, \end{aligned} \quad (3.9)$$

¹We choose the case of $B = 0$ and X is arbitrary even though the case of $B = \text{arbitrary}$ and $\text{Tr} X = 0$ is also possible. If we take A and B as the messengers, we don't have to consider the latter.

the superpotential can be expressed as

$$\begin{aligned}
W &= h\text{Tr}\tilde{\varphi}\Phi\varphi - h\mu^2\text{Tr}\Phi + \lambda AB\text{Tr}\Phi + mA^2 \\
&= -h\mu^2\text{Tr}(X + \delta\Phi_{22}) + m(\delta A)^2 \\
&\quad + h\text{Tr}\left[\delta\Phi_{11}\left(\mu Y\delta\tilde{\varphi}_1 + \delta\varphi_1\frac{\mu}{Y}\right) + \frac{\mu}{Y}\delta\Phi_{12}\delta\varphi_2 + \delta\tilde{\varphi}_2\delta\Phi_{21}\mu Y\right] \\
&\quad + h\text{Tr}[\delta\tilde{\varphi}_1\delta\Phi_{11}\delta\varphi_1 + \delta\tilde{\varphi}_1\delta\Phi_{12}\delta\varphi_2 + \delta\tilde{\varphi}_2\delta\Phi_{21}\delta\varphi_1 + (X + \delta\Phi_{22})\delta\varphi_2\delta\tilde{\varphi}_2] \\
&\quad + \lambda\delta A\delta B\text{Tr}[\delta\Phi_{11} + (X + \delta\Phi_{22})].
\end{aligned} \tag{3.10}$$

Now let us consider the Coleman-Weinberg potential for the pseudo-moduli(i.e. X) [32]

$$\begin{aligned}
V_{\text{eff}}^{(1)} &= \frac{1}{64\pi^2}\text{STr}\left(\mathcal{M}^4\log\frac{\mathcal{M}^2}{\Lambda^2}\right) \\
&\equiv \frac{1}{64\pi^2}\left[\text{Tr}\left(\mathcal{M}_B^4\log\frac{\mathcal{M}_B^2}{\Lambda^2}\right) - \text{Tr}\left(\mathcal{M}_F^4\log\frac{\mathcal{M}_F^2}{\Lambda^2}\right)\right],
\end{aligned} \tag{3.11}$$

where \mathcal{M}_B^2 and \mathcal{M}_F^2 are the tree-level boson and fermion mass matrices

$$\mathcal{M}_B^2 = \begin{pmatrix} W_{ik}^\dagger W^{kj} & W_{ijk}^\dagger W^k \\ W^{ijk} W_k^\dagger & W^{ik} W_{kj}^\dagger \end{pmatrix}, \quad \mathcal{M}_F^2 = \begin{pmatrix} W_{ik}^\dagger W^{kj} & 0 \\ 0 & W^{ik} W_{kj}^\dagger \end{pmatrix}. \tag{3.12}$$

As usual, W^i stands for $\partial W/\partial\delta\varphi_i$. From now on, we work at $Y = 1$.

Before going further, we have to check whether there exist tachyonic modes at the pseudo-moduli space of (3.4). Among the eigenvalues of \mathcal{M}_B^2 , there are two modes which could be potentially dangerous,

$$\begin{aligned}
&\lambda^2 X^2 + 2m^2 - \sqrt{4m^2\lambda^2 X^2 + 4hm\lambda^2\mu^2 X + h^2\lambda^2\mu^4 + 4m^4} \\
&\lambda^2 X^2 + 2m^2 - \sqrt{4m^2\lambda^2 X^2 - 4hm\lambda^2\mu^2 X + h^2\lambda^2\mu^4 + 4m^4}.
\end{aligned} \tag{3.13}$$

There are no tachyonic modes for the following range of the pseudo-modulus X :

$$|X| \gtrsim \max\{\mu, (4m\mu^2)^{1/3}\} \tag{3.14}$$

where we set the couplings $h \simeq \lambda \simeq 1$. The two modes become tachyonic outside of the range of (3.14) and the pseudo-moduli space (3.4) is locally unstable there. As a result, the fields can roll down to the supersymmetric runaway vacua along the tachyonic directions [33, 34].

In the range of (3.14) for the pseudo-modulus X , there does not exist a tachyonic mode, and hence we can compute the one-loop Coleman-Weinberg potential (3.11) in that pseudo-moduli space (3.4). In order for the $U(1)_R$ symmetry to be broken spontaneously, we need to find a field, which has a non-zero VEV but charged with $U(1)_R$. In fact, X is a flat direction in the tree level and has a R-charge 2, but it is lifted and has a locally stable minimum at $X \sim \mathcal{O}(1)\mu$ by the one-loop Coleman-Weinberg potential. The $U(1)_R$ symmetry is spontaneously broken there. Note that

$$|\langle X \rangle| \sim \mathcal{O}(1)|\mu| \quad \longrightarrow \quad |\langle X \rangle| \ll |\Lambda_m|. \tag{3.15}$$

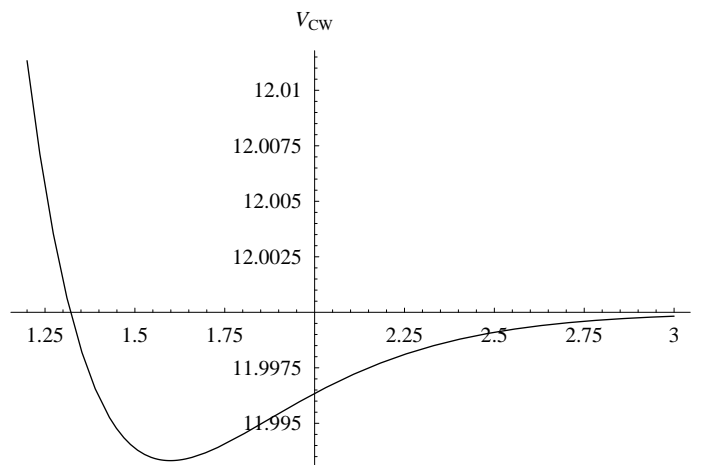


Figure 1: Plot of the Coleman-Weinberg potential without gauge interaction terms. We neglect the numerical factor of $\frac{1}{64\pi^2}$. We set $X/\mu = e^t$.

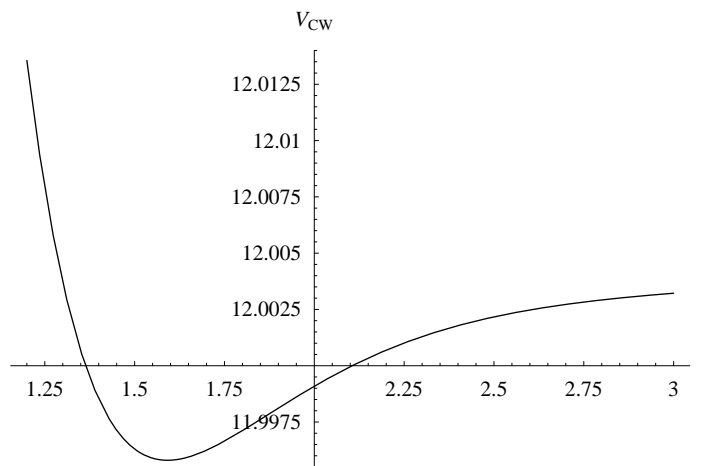


Figure 2: Same as figure 1 except for the gauge interaction $g = 0.1$.

Thus, our addition of new fields, A and B , does not ruin the merit of the original ISS model i.e. the longevity of the metastable vacua. It is important that there exist stable minima for a wide range of parameters of h , λ and m :

$$h \gtrsim \mathcal{O}(0.1), \quad \lambda \gtrsim \mathcal{O}(0.1), \quad m \lesssim 1.3\mu. \tag{3.16}$$

As an example, we show figure 1 for $h = 1$, $\lambda = 1$, and $m = 0.1\mu$.

So far, we considered the superpotential without the gauge interaction terms in the ISS model. Therefore, we also have to check whether the gauge interaction terms spoil the basic feature of our ISS model and whether the spontaneous $U(1)_R$ breaking feature is maintained. It is easy to obtain the one-loop Coleman-Weinberg potential in the presence of the gauge bosons and gauginos, in the same range of the pseudo-modulus X given above.

Obviously, if the gauge coupling g is small, our result is not modified significantly. As we can see in figure 2, there exist stable minima at the order of $\mathcal{O}(1)\mu$ even for the case of a bit large coupling. In addition, this effect increases the slope of the Coleman-Weinberg potential at large field values compared to the case without the gauge interaction terms. As a result, the addition of the gauge interaction terms rather improves the stability of the $U(1)_R$ breaking minima. A plot of the Coleman-Weinberg potential with the gauge interaction terms is given in figure 2 for $h = 1$, $\lambda = 1$, $m = 0.1\mu$, and $g = 0.1$.

4. Applications

The exact $U(1)_R$ symmetry is broken at high energy scale, but it seems that the approximate $U(1)_R$ can forbid the gaugino mass terms. In our case, however, it is broken even at the low energy scale. Therefore, we can generate the Majorana mass terms for the gauginos. Since we obtain $U(1)_R$ symmetry breaking meta-stable vacua dynamically and spontaneously, the gaugino masses are generated without introducing messenger mass by hand. In addition to this, the messengers are not tachyonic, for it is always guaranteed that $\langle F \rangle \leq \langle X \rangle^2$ in the allowed parameter range. If we accept the gauge mediation scheme, then the messenger mass scale will be determined by the scales of $\langle \text{Tr}\Phi \rangle = \langle X \rangle > \mu$ and $\langle F \rangle \sim h\mu^2$. In the gauge mediation scheme ², we get the gaugino soft mass terms as

$$\begin{aligned}
 m_{\frac{1}{2}}(t) &= k_r \frac{\alpha_r(t)}{4\pi} \Lambda_G \\
 \Lambda_G &= \sum_{i=1}^{N_f} n_i \frac{F_i}{M_i} g(F_i/M_i^2) \\
 g(x) &= \frac{1}{x^2} [(1+x) \ln(1+x) + (1-x) \ln(1-x)],
 \end{aligned}
 \tag{4.1}$$

where $\alpha_r(t)$ is related to the visible sector gauge coupling at the messenger scale. k_r and n_i are constants of $\mathcal{O}(1)$, which depend on the messenger structure [36]. A rough estimation of $m_{\frac{1}{2}} \sim \mathcal{O}(1)\text{TeV}$ gives a phenomenologically acceptable parameter range for h , λ and m . Namely, we get $\mu \sim \mathcal{O}(10^6)\text{GeV}$. Even though we introduced the messenger fields, it is found that the structure of the potential is not dangerous. As we have seen in Sec 3.2, the mass scales of the newly introduced singlets are determined by the requirement of $U(1)_R$ symmetry breaking stable minima. Moreover, the newly introduced singlets have an interaction term with the Goldstino supermultiplet fields (i.e. X), and so they can be used as the messenger fields. However, we do not consider any specific models here. Their phenomenological applications are left for future communication.

Now we will turn to the cosmological aspect. In the original ISS, $U(1)_R$ symmetry is broken in the microscopic view point. On the other hand, in the macroscopic view point R-symmetry is broken by the non-perturbative term in (2.7). Thus, this R-symmetry is not exact but approximate and accidental. A spontaneously broken approximate global symmetry gives arise to a pseudo goldstone boson. In our model, $U(1)_R$ is an approximate

²For the review, see [35].

symmetry because of the Λ_m suppressed term and what we have obtained are the vacua which break it spontaneously. Therefore, we do not have to worry about the R-axion.

5. Conclusion

We have studied the radiatively generated spontaneous $U(1)_R$ symmetry breaking in the ISS setup. We introduced two gauge singlet fields with R-charges 1 and -1 respectively to keep the superpotential generic up to $U(1)_R$ symmetry. Since the $U(1)_R$ is an approximate symmetry, we need not worry about the R-axion. For the study of the pseudo-moduli space, we have found the spontaneously broken radiatively generated meta-stable vacua with a large range of parameter space. We can easily satisfy the meta-stability of our vacua because it appears to be at the scale of $O(\mu)$ where μ can be made small by the retrofitting argument. We can obtain the dynamically generated messenger scale which turns out to be the SUSY breaking scale. As a result, we can build a realistic model without introducing the messenger scale by hand. In addition, new singlets can act as the messengers. Finally, if the gauge mediation scheme is used for the SUSY breaking, the SUSY breaking scale is $\mathcal{O}(10^6)\text{GeV}$.

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