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***R*-symmetric gauge mediation with Fayet-Iliopoulos term**

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ABSTRACT: We have studied *R*-symmetric gauge mediation models with Fayet-Iliopoulos terms. We give a concrete example of hidden sector with an $U(1)_H$ gauge theory and a Fayet-Iliopoulos term, which can induce distinctive soft terms in the visible sector, and help solving fine tuning problems in models of *R*-symmetric gauge mediation.

KEYWORDS: Supersymmetric Standard Model, Supersymmetry and Duality, Supersymmetry Breaking.

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1. Introduction

Of all the candidates to stabilize the hierarchy between the weak and Plank scales, supersymmetry seems to be the most plausible and predictive at LHC. Supersymmetry is broken in some hidden sector and mediated to the visible sector via gravity or gauge interactions [2, 3]. Such mechanisms induce the necessary soft terms in supersymmetric Standard Model (SSM) for which to be phenomenologically viable. In models with direct gauge mediation, a number of hidden models have been successfully constructed [21]. There exist now a general framework to calculate the soft masses [17].

Recently, a class of gauge mediation models with R -symmetry has been proposed [4, 5]. In usual gauge mediation models, R symmetry has to be spontaneously broken in order to generate Majorana gaugino masses, while R symmetry is retained in this class of models. By adding suitable C -parity chiral fields in adjoint representations, the gauginos can acquire Dirac masses. Phenomenologies are rather distinctive in these kind of models [5]. Unfortunately, some of the usual flavor problems persist in these models and fine tunings are needed.

In this paper, we will first have a close look at hidden sectors with R -symmetry and SUSY-breaking. It has been long understood that these two issues are closely connected to each other [16]. The R -symmetry can be broken at the tree level. If it is not broken at the tree level, it can still be spontaneously broken due to quantum corrections in general O'Raifeartaigh (OR) model, provided that the R -charges of superfields in the superpotential take values different from 0 and 2 [20]. This implies that the general form of superpotential is determined as $W = X^n f_n(\Phi)$, where X, Φ are R -charges 2 and 0 respectively, $f_n(\Phi)$ are polynomial of Φ constrained only by the renormalization of the theory. In section II, we will deform these models by including a Fayet-Iliopoulos (FI) term. And one sees that the R -symmetry can still be preserved.

In section III, we will construct a hidden sector with a FI term to realize the scenario of R -symmetric gauge mediation. Embedded into SUSY theories via an $U(1)$ group, the

phenomenological implications of the FI term have been extensively discussed in gauge and anomaly mediation.¹ In gauge mediation scheme, an U(1) FI term can be used to spontaneously break supersymmetry, while R -symmetry is usually unbroken as proposed in [16]. In $N = 2$ models, this can naturally generate Dirac gaugino masses [10]. In [11], an FI term is employed for $N = 1$ models with Majorana gaugino masses, in which R -symmetry is spontaneously broken by gaugino condensation in a strong-coupled Yang-Mills hidden sector. FI terms can also be applied in the construction of SUSY grand unification theories (GUT). If the role of messenger sector are replaced by an U(1) with FI term, there will be less modifications to RG running of SM gauge couplings. A GUT can be relatively easier realized [12]. Moreover, doublet-triplet splitting problem in high rank $SU(n)$ SUSY theories that include SU(5) GUT can be solved by introducing a new mass scale carried by FI term [13]. In anomaly mediation scheme, introducing a FI term can help to solve the tachyonic slepton problem, or even accommodate neutrino masses when suitable U(1) charges are chosen [14, 15].

In section IV, we calculate the soft masses, and $\mu/B\mu$ terms. The model preserves most distinctive features of R -symmetric gauge mediation. However, we can allievate the fine tunings and the $\mu/B\mu$ problem. To summarize, the model has the following features:

- R -symmetry is conserved in SUSY-breaking and visible sectors.
- There are no A terms, and adjoint fields Φ_r ($r=1,2,3$) that combined with gauginos have to be added to construct Dirac gauginos.
- There is an unbroken $U(1)_H$ gauge theory in the hidden sector. The SUSY-breaking effects contain the contributions of F - and D -terms at meantime.
- The negative sfermions masses squared coming from the D -term may solve the fine tuning problems in R -symmetric gauge mediation with only F -term induced visible effects [4]. The Dirac gaugino masses can be heavier than the sfermions masses.
- The little hierarchy between μ and $B\mu$ terms can also be obtained by adjusting the D -term.

We expect that some of phenomenologies in the visible sector are model dependent. It would be interesting to develop a general framework to distinguish the general characters in R -symmetric gauge mediation, as done in [17]. On the other hand, the $\mu/B\mu$ problem should also be discussed in depth to find a more economical mechanism. Finally, we conclude in section V with discussions.

2. SUSY-breaking sectors with R -symmetry

According to [20], the R -symmetry is maintained only if that the R charges of superfields in the superpotential are either 0 or 2. Otherwise, radiative corrections will spontaneously

¹Applications in other fields, such as extra dimensions, cosmology and string theory, are beyond present discussions.

break the R -symmetry even if it is conserved at tree level. This implies that the general form of superpotential with R -symmetry reads,

$$W = X^n f_n(\Phi_m), \tag{2.1}$$

where $R(X^n) = 2$ and $R(\Phi_m) = 0$, n, m denote different superfields of X and Φ , respectively, repeated n implies summation. $f_n(\Phi)$ are n polynomial of Φ constrained only by the renormalization of the theory. The supersymmetry preserving vacua is given by vanishing F terms:

$$F_n = f_n(\Phi_m) = 0, \quad F_m = X^n \partial_m f_n(\Phi_m) = 0 \tag{2.2}$$

$F_m = 0$ can easily be satisfied by simply taking all $X^n = 0$. In cases with $n > m$, generally $F_n = 0$ can not be simultaneously satisfied, thus leading to supersymmetry spontaneously broken. The R -symmetry is preserved at the origin of the moduli space. In cases with $n \leq m$, there are two possibilities. In a generally renormalizable theory, $f_n(\Phi_m)$ always assumes the following form,

$$f_n(\Phi_m) = k_n + M_n^m \Phi_m + \lambda_n^{mm'} \Phi_m \Phi_{m'} \tag{2.3}$$

If all $k_n = 0$, we expect the supersymmetry and R symmetry are both unbroken at the origin of the moduli space. On the other hand, if some $k_n \neq 0$ as in the ordinary OR model, it is then possible to obtain SUSY-breaking models with R symmetry. For example, the ISS model constructed in [19] belongs to this type.

If there is an $U(1)_H$ gauge interaction in the hidden sector, one can also have a FI term in principle. This yields the mechanism of D -term scenario, in addition to the F -term scenario, to induce SUSY-breaking. Generically, a hidden sector with a $U(1)_H$ gauge can have the following potential,

$$V = F_m^2 + F_n^2 + \frac{D^2}{2g^2}, \quad D = g^2 (\xi + g^2 q_i (|\Phi_i|^2 - |\Phi'_i|^2)) \tag{2.4}$$

where g is the $U(1)_H$ gauge coupling, q_i are the $U(1)_H$ charges of Φ_i . Without k_n terms, one obviously cannot have $F = 0$ and $D = 0$ in the same time. The supersymmetry is thus broken spontaneously. In the specific case that the absolute minimum of the potential is at $\Phi = 0$, one would have $V = g^2 \xi^2 / 2$ and the SUSY-breaking comes only from the D -term.

In the case that there are non-zero k_n 's, the SUSY-breaking will come both from F - and D -term, if we can arrange the parameters in the model such that the absolute minimum of the potential is still at $\Phi = 0$. In this case, the R symmetry is unbroken with a minima $V = k_1^2 + \frac{1}{2} g^2 \xi^2$. As we will see, the R -symmetric gauge mediation with such a hidden sector has distinctive phenomenologies compared with only F term induced SUSY-breaking [4].

For simplicity, we have assumed that the Kahler potential is canonical, i.e, $K = X_n^\dagger X_n + \Phi_m^\dagger \Phi_m$, In principle, it can be modified by the radiative corrections which can in turn induce a non-trivial metric of moduli space.

	SU(5)	U(1) _H
Q_i	\square	+1
\bar{Q}_i	$\bar{\square}$	-1
Q_6	\square	0
\bar{Q}_6	$\bar{\square}$	0

Table 1: Q_i denote the first five chiral superfields. Note that the gauge groups in the hidden sector are the simplest extension of that in [4], which are well motivated and can be easily constructed in intersecting branes models.

3. A concrete example of hidden sector

Consider a hidden sector with the gauge groups of $SU(5) \times U(1)$, two types of chiral superfields Q in the fundamental representation, and two types of chiral superfields \bar{Q} in the anti-fundamental representation, as shown in the table. The global flavor symmetry is $SU(6)$. The $U(1)_H$ charges of Q 's and \bar{Q} 's are either +1 or -1, so the model is anomaly free.

In the infrared, the strong coupling $SU(5)$ theory can be described by the dual magnetic theory of ISS-type [19], with a superpotential

$$W_{\text{mag}} = \lambda \bar{q} \mathcal{M} q - f^2 \text{Tr} \mathcal{M} \tag{3.1}$$

where $q = (\varphi, \psi)$ are the dual quarks and \mathcal{M} the mesons,

$$\mathcal{M} = \begin{pmatrix} \omega M & kN \\ k\bar{N} & k'Y \end{pmatrix} \tag{3.2}$$

In detail, the superpotential of our theory can be written as,²

$$\begin{aligned} W_{\text{mag}} &= X^n f_n(\Phi_m) \\ f_1 &= -f^2 \omega + \lambda \varphi \bar{\varphi}, \\ f_2 &= \lambda k \bar{\varphi} \psi, \\ f_3 &= \lambda k \varphi \bar{\psi}, \\ f_4 &= -f^2 + \lambda k' \bar{\psi} \psi, \end{aligned} \tag{3.3}$$

where X_n denote the mesons singlets (M, N, \bar{N}, Y) of R -charge 2, the rest of R -charge 0. When the $U(1)_H$ coupling constant $g \rightarrow 0$, the theory returns to the usual Seiberg duality [18] and corresponding ISS model.

Here is the rational for such an assumption. In the standard model the strong $SU(3)_c$ quark theory can be well described by the dual effective theory of composite mesons and baryons at low energy, which are not invalidated by extra electroweak interactions. To

²We can also begin with a hidden sector with superpotential W_{mag} and D -term, and consider the $SU(5) \times U(1)_H$ theory as one ultraviolet completion. For instance, abelian SUSY theory with a set of chiral superfields including singlets is another candidate.

reach the macroscopic superpotential (3.3), we have assumed that the $U(1)_H$ gauge theory does not spoil the validity of Seiberg duality with small enough coupling constant g . There are spontaneously broken $N = 2$ dual theories with FI terms [9], and $N = 1$ dualities which are not spoiled by deformations of IR irrelevant couplings [7]. There are also Seiberg dualities with non-simple Lie groups. For example, there are $N = 1$ theories with two gauge groups [8]. Similar assumption has been applied to study other topics in earlier works [6].

We now read out the $U(1)_H$ charges of the dual mesons and quarks. The singlet mesons M, Y are also $U(1)_H$ singlets. The N, \bar{N} mesons are not $U(1)_H$ singlets, but of $U(1)_H$ charges $+1, -1$, respectively. The dual quarks q carry the same $U(1)_H$ charges as those of Q . In summary, all dual fields are $U(1)_H$ singlets except $\varphi, \bar{\varphi}, N, \bar{N}$.

It can be shown that the supersymmetry is broken in our model. The absolute minimum of V is located at the origin of moduli space with

$$F_{Tr M} = \omega f^2, \quad \bar{\psi}\psi = v^2 = f^2/(\lambda k') \quad (3.4)$$

Since $\psi, \bar{\psi}$ have no $U(1)_H$ charges, their nonzero vacuum expectation values do not break the $U(1)_H$ gauge symmetry spontaneously.³ However, the global $SU(6)$ symmetry is spontaneously broken to $SU(5)$. The corresponding Nambu-Goldstone (NG) bosons acquire significant masses due to interactions with the messenger sector [4].

To mediate the SUSY-breaking to the visible sector, we gauge the remaining global $SU(5)$ flavor symmetry. The ϕ and N fields will serve as the messengers. The scalar mass matrix for messengers is given by,

$$\begin{pmatrix} \varphi^* & \bar{\varphi} & N^* & \bar{N} \end{pmatrix} \begin{pmatrix} M^2 + D & -zM^2 & 0 & 0 \\ -zM^2 & M^2 - D & 0 & 0 \\ 0 & 0 & M^2 + D & 0 \\ 0 & 0 & 0 & M^2 - D \end{pmatrix} \begin{pmatrix} \varphi \\ \bar{\varphi}^* \\ N \\ \bar{N}^* \end{pmatrix} \quad (3.5)$$

where $M = \sqrt{\lambda\omega/z}f$ and $z = \omega k'/k^2$. From eq. (3.5) we see that the eigenstates of the upper 2×2 block of scalar mass matrix are $\phi_{\pm} = (\bar{\varphi}^* \pm \varphi)/\sqrt{2}$ of eigenvalues,

$$\phi : m_{\pm}^2 = M^2(1 \pm \tilde{z}), \quad \tilde{z} = \sqrt{z^2 + x^2} \quad (3.6)$$

where $x = D/M^2$. In order to avoid tachyons in the spectrum, we need to impose $x^2 < (1 - z^2)$. The eigenvalues of N, \bar{N} are given by,

$$N : m_{\pm}^2 = M^2(1 \pm x) \quad (3.7)$$

The fermion mass matrix for messengers φ, N is off-diagonal,

$$\begin{pmatrix} \varphi & N \end{pmatrix} \begin{pmatrix} 0 & Me^{i\theta} \\ Me^{-i\theta} & 0 \end{pmatrix} \begin{pmatrix} \bar{\varphi} \\ \bar{N} \end{pmatrix} \quad (3.8)$$

³Recently, it is proposed that unbroken weak coupling $U(1)_H$ theory can work as a model of dark matter [23].

They are all degenerate at M . The spectra for NG particles and the remaining messengers $X, \psi, \bar{\psi}$ are the same as those given in [4]. Compared with the messenger spectra in [4], the masses of scalar messengers ϕ and N are modified by the D -term. The sfermions masses squared and Dirac gaugino masses in the visible sector will be induced by the mass splitting of ϕ and N in the loop(s).

4. Soft terms and fine tunings

We now analyze the soft terms in the visible sector. Before going to the details, we outline the main characteristics of R -symmetric gauge mediation with D -term,

- The sfermion masses are decreased by negative D -term induced contribution. When $\sqrt{D} \sim M \sim 10^3 \text{TeV}$, the sfermions are still of order $\mathcal{O}(1) \text{TeV}$. On the other hand, the Dirac gaugino masses increase by positive D -term induced contribution. Without D -term effects the sfermions masses are usually heavier than gauginos masses. However, this can easily be reversed with the D -term. This reversion of the gaugino and scalar mass ratio helps to evade constraints on flavor structures.
- As shown in [4], fine tunings are needed to have viable diagonal and off-diagonal coefficients c_D and c_{OD} for the sfermion mass matrix. An adjustable D -term helps to obtain reasonable relations $c_D \sim c_{OD} \sim 1$.
- $B\mu$ and μ terms receive the nonzero and zero contributions from D -term respectively, which can be used to adjust the small hierarchy between μ and $B\mu$.
- Gravitino is the lightest superparticle (LSP) of a mass in the order of eV , which is determined by $m_{3/2} \sim (D + f^2)/M_{Pl}$ in supergravity.

The sfermion masses squared receive contributions from the ultraviolet (UV) and infrared (IR) physics. The IR contribution is due to the mass splitting of the messengers induced by the SUSY breaking in the hidden sector, starting from two-loops [22]. The UV contribution can be written generically,

$$\int d^4\theta \frac{c_{ij}}{\Lambda^2} (\Xi \Xi^\dagger) Q_i^\dagger Q_j \tag{4.1}$$

where $\Xi = \langle Tr M \rangle = \theta^2 \omega f^2$. In total, we have

$$\begin{aligned} (\tilde{m}^2)_{ij} &= c_{ij} (\tilde{m}_{UV}^2)_{ij} - (\tilde{m}_{IR}^2)_{ij}, \\ (\tilde{m}_{IR}^2)_{ij} &= \delta_{ij} \frac{g_s^4 M^2}{(16\pi^2)^2} J(x, \tilde{z}), \\ \tilde{m}_{UV} &= \begin{pmatrix} z \\ \lambda \end{pmatrix} \begin{pmatrix} M \\ \Lambda \end{pmatrix} M \end{aligned} \tag{4.2}$$

where

$$J(x, \tilde{z}) = \frac{7}{9}(x^4 + \tilde{z}^4) + \frac{38}{75}(x^5 + \tilde{z}^5) + \mathcal{O}(x^6, \tilde{z}^6) \tag{4.3}$$

c_{ij} are the coefficients appearing in the UV operator. The UV operators responsible for Dirac gaugino masses are,

$$\int d^2\theta \frac{\bar{D}^2 D_\alpha(\Xi\Xi^\dagger)}{\Lambda^3} Tr(W^\alpha\Phi) \quad (4.4)$$

where W' , W refer to the $U(1)_H$ and SSM spinor superfields respectively. The IR contribution to gaugino masses is again due to the mass splitting of the messengers, starting from one-loop. Since the masses of the scalars N, \bar{N} in the loop are shifted, gaugino masses are changed, in comparison with those in [4]. Explicitly, we have

$$\begin{aligned} m_{1/2} &= m_{\text{IR}} + m_{\text{UV}}, \\ m_{\text{IR}} &= \frac{g_s y}{16\pi^2} M \cos\left(\frac{\theta}{v}\right) Q(x, \tilde{z}) \end{aligned} \quad (4.5)$$

where

$$m_{\text{UV}} \simeq \left(\frac{\tilde{m}_{\text{UV}}}{\Lambda}\right) \tilde{m}_{\text{UV}} \quad (4.6)$$

and θ is the NG bosons, The function $Q(x, \tilde{z})$ is defined by ,

$$Q(x, \tilde{z}) = \frac{1}{\tilde{z}} ((1 + \tilde{z}) \log(1 + \tilde{z}) - (1 - \tilde{z}) \log(1 - \tilde{z}) - 2\tilde{z}) + (\tilde{z} \rightarrow x) \quad (4.7)$$

The coefficients c_{ij} can also be written as [4],

$$c_{\text{D}} = \frac{\tilde{m}^2 + \tilde{m}_{\text{IR}}^2}{\tilde{m}_{\text{UV}}^2}, \quad c_{\text{OD}} = \delta \left(\frac{\tilde{m}^2}{\tilde{m}_{\text{UV}}^2} \right) \quad (4.8)$$

which refer to the non-perturbative behavior of the hidden sector, which arises when the RG scale near the Landau pole Λ in the direct gauge mediation.⁴ In principle, c_{ij} are of the order $\mathcal{O}(1)$ and cannot be calculated in perturbative method. However, they are constrained by ratio of gaugino mass over sfermion mass to avoid flavor problems. In the ISS model, c_{D} should be smaller than 10^{-2} to be phenomenologically viable. This implies that the non-perturbative physics of the hidden sector is seriously constrained and this is the origin of fine tuning.

In models with D -term breaking, c_{D} can be in the neighborhood of unity. Take $M \sim 10^3 \text{TeV}$ and $\Lambda \sim 10^4 \text{TeV}$ for illustrations, in which $\alpha_3(M) \sim 0.15$. We start with the following parameter space,

$$z = 0.1 \quad \text{and} \quad \lambda = 1, \quad (4.9)$$

which can yield the typical mass relations,

$$\tilde{m} \sim \tilde{m}_{\text{UV}} \sim \tilde{m}_{\text{IR}} \sim 10^{-3} M, \quad m_{1/2} \sim m_{\text{IR}} \sim 10^{-2} M \quad (4.10)$$

⁴The messengers introduced to mediate the SUSY-breaking effects substantially modify the slopes of gauge coupling α_s running, lead to the divergence of α_s at Λ .

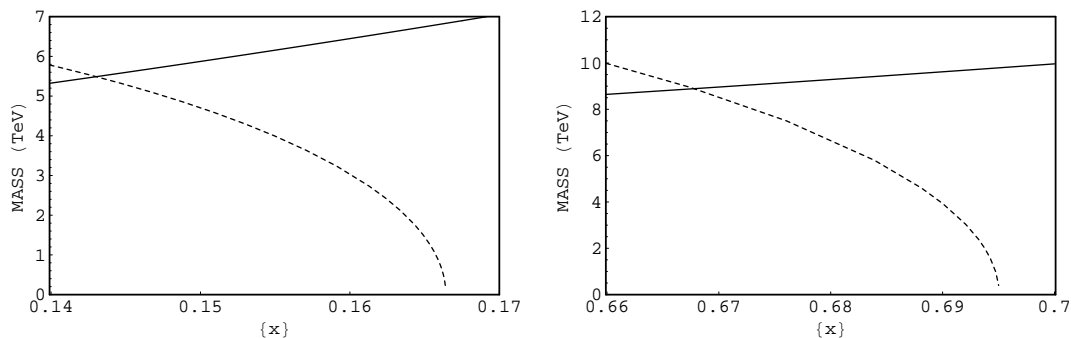


Figure 1: The masses of sfermion (dashed line) and gaugino (solid line) (TeV scale) as functions of x parameter. Left figure is for the region of small $c_D = 0.1$, where a large mass ratio can be obtained near $x \sim 0.16$ ($y = 4$ and $M = 3 \times 10^3$ TeV). Right figure is for the region of large $c_D = 8$, where a large mass ratio can be obtained near $x \sim 0.69$ ($y = 3$ and $M = 10^3$ TeV).

Explicit calculations show that $c_D \leq 10$ and $c_{OD} \leq 1$ are phenomenologically viable. Shown in figure. 1 are the masses of sfermions and gauginos for two typical regions of c_{ij} . Clearly, for the given parameter space (4.9), there are no fine tunings and no extra Landau pole coming from large Yukawa coupling λ .

When going up along the positive direction of z , one has to increase $x \sim 1$ and (or) λ in order to restore the typical mass relations (4.10). However, there is the upper bound on $x \sim (1 - z)$, as one wants to avoid tachyons. Roughly z cannot be greater than 0.5.

Finally, we address the $\mu/B\mu$ term in R -symmetric gauge mediation. Generally, there are three mechanisms to generate $\mu/B\mu$ term of weak scale in gauge mediation without R -symmetry. One is by introducing a gauge singlet X , usually dubbed as NMSSM [24]. Another is by introducing massive vector-like pairs. Thirdly, one takes the conformal sequestering into account [25]. In R -symmetric gauge mediated theories, the unbroken R -symmetry severely restricts the μ term, while $B\mu$ of weak scale can be generated from either the F - or D -terms. An economical scheme to the μ term is to introduce two extra SU(2) doublets $R_{u,d}$ of R -charge 2 [4]. The UV operators corresponding to μ and $B\mu$ are

$$\int d^4\theta c_F \frac{\Xi^\dagger}{\Lambda} (H_u R_d + H_d R_u) \quad (4.11)$$

and,

$$\int d^4\theta c'_F \frac{\Xi \Xi^\dagger}{\Lambda^2} H_u H_d + \int d^2\theta c'_D \frac{W'_\alpha W'^{\alpha}}{M^2} H_u H_d \quad (4.12)$$

respectively. Here the second term comes from the unbroken $U(1)_H$ gauge theory. They are both at the weak scale if $\sqrt{D} \sim M \sim 10^3$ TeV

$$\mu \sim c_F m_{UV}, \quad B\mu \sim c'_F m_{UV}^2 + c'_D \frac{D^2}{M^2} \quad (4.13)$$

Note that the simplest way to generate μ term is to include a R -charge zero spurion field of nonzero F -term in the messenger sector. According to discussions in section 2,

this can be realized. For instance, the Wess-Zumino model with only chiral superfields is a feasible candidate. It would be interesting to develop a scheme to discuss the general R -symmetric gauge mediation as done in ordinary gauge mediation [17], to see which of the features outlined above are general and which are model dependent.

5. Conclusions

In this paper, we have discussed the possibilities to obtain SUSY-breaking hidden sectors with R -symmetry, with an extra $U(1)_H$ sector and a FI term. A concrete example of hidden sector is constructed. In this particular model, we find that in the visible sector, the ratio of Dirac gaugino mass over sfermion mass substantially increases compared with those with only F -term [4]. The μ and $B\mu$ terms receive zero and nonzero contribution, respectively. These help evading the fine tunings in R -symmetric gauge mediation with interesting flavor phenomenologies.

As we point out in section 2, there are other candidates as hidden sectors in R -symmetric gauge mediation. It is worth developing a general framework to see which characters are model independent, especially the generations of $\mu/B\mu$ terms, which closely connect with some important topics such as electro-weak symmetry breaking and the dark matter model of supersymmetric neutralino.

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