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Revisiting R -invariant direct gauge mediation

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ABSTRACT: We revisit a special model of gauge mediated supersymmetry breaking, the “ R -invariant direct gauge mediation.” We pay particular attention to whether the model is consistent with the minimal model of the μ -term, i.e., a simple mass term of the Higgs doublets in the superpotential. Although the incompatibility is highlighted in view of the current experimental constraints on the superparticle masses and the observed Higgs boson mass, the minimal μ -term can be consistent with the R -invariant gauge mediation model via a careful choice of model parameters. We derive an upper limit on the gluino mass from the observed Higgs boson mass. We also discuss whether the model can explain the 3σ excess of the Z +jets+ E_T^{miss} events reported by the ATLAS collaboration.

KEYWORDS: Supersymmetry Phenomenology

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Contents

1	Introduction	1
2	$B\mu$-problem in R-invariant direct gauge mediation model	2
2.1	R -invariant direct gauge mediation model	2
2.2	Gauge mediated mass spectrum	3
2.3	$B\mu$ -problem	6
2.4	Gravitino dark matter	10
3	Upper bound on the gluino mass	11
4	Z+jets+missing E_T	13
5	Summary	15
A	R-charge assignments	17

1 Introduction

The model of gauge mediated supersymmetry (SUSY) breaking [1–6] is the most attractive candidate for phenomenologically successful minimal supersymmetric standard model (MSSM). In this case, soft SUSY breaking is mediated via the MSSM gauge interactions and, thus, the model is free from the infamous SUSY flavor changing neutral current problem.

One of the drawbacks of gauge mediation models is their somewhat cumbersome structure. In particular, careful model building is required to connect messenger fields to a SUSY breaking sector without destabilizing the SUSY breaking vacuum in the SUSY breaking sector. In fact, naive couplings between the SUSY breaking sector and messenger fields often lead to meta-stability of the SUSY breaking vacuum. In those models, the thermal history of the Universe and/or the masses of messenger fields are severely constrained [7].

Among various safe scheme to connect messengers to the SUSY breaking sector, the model developed in refs. [8, 9] is highly successful. In particular, the SUSY breaking vacuum is not destabilized, and hence the model is durable even when the reheating temperature of the Universe is very high. The stability of the SUSY breaking vacuum is achieved through R -symmetry, which is the origin of the name of the model, “ R -invariant direct gauge mediation.”¹

¹For simple embedding of the model into a dynamical SUSY breaking model with a radiative R -symmetry breaking, see ref. [10].

In this paper, we revisit the R -invariant direct gauge mediation model by paying particular attention to the consistency of the model with the minimal model that addresses the origin of the μ -term. Here the minimal model of the μ -term means one with a simple mass term for the Higgs doublets in the superpotential, which leads to a vanishing B -term at the messenger scale. As pointed out in ref. [11], it is difficult for the minimal model of the μ -term to be compatible with the R -invariant direct gauge mediation, for the model predicts rather suppressed gaugino masses compared with scalar masses. As we will see in this paper, the minimal μ -term can be consistent with the R -invariant gauge mediation model through a careful choice of model parameters, although the incompatibility is highlighted in view of the current experimental constraints on superparticle masses and the observed Higgs boson mass.

After discussing the compatibility of the R -invariant direct gauge mediation model with the minimal μ -term, we derive an upper bound on the gluino mass from the observed Higgs boson mass by exploiting the predicted ratio between the gluino and the stop masses in the R -invariant gauge mediation model. As a result of the upper bound, we find that a large portion of parameter space can be tested by the LHC Run-II with an integrated luminosity of 300 fb^{-1} , unless model parameters are highly optimized to obtain a large gluino mass.

We also discuss whether the R -invariant direct gauge mediation model can explain the 3σ excess of $Z+\text{jets}+E_T^{\text{miss}}$ events reported by the ATLAS collaboration [12]. We seek a spectrum similar to the one in ref. [13] where the gluino mainly decays into a gluon and a Higgsino via one-loop corrections. With such a spectrum, the excess can be explained while evading all the other constraints from SUSY searches at the LHC.

The paper is organized as follows. In section 2, we discuss the consistency between the R -invariant direct gauge mediation model and the minimal model of the μ -term. In section 3, we derive an upper bound on the gluino mass from the observed Higgs boson mass. In section 4, we discuss whether the R -invariant direct gauge mediation can explain the signal reported by the ATLAS collaboration. The last section is devoted to the summary of our discussions.

2 $B\mu$ -problem in R -invariant direct gauge mediation model

2.1 R -invariant direct gauge mediation model

We first review the minimal R -invariant direct gauge mediation model constructed in ref. [8, 9] (see also ref. [14]). This model introduces N_M sets of messenger fields, Ψ_i , $\bar{\Psi}_i$, Ψ'_i and $\bar{\Psi}'_i$, which are respectively $\mathbf{5}$, $\bar{\mathbf{5}}$, $\mathbf{5}$ and $\bar{\mathbf{5}}$ representations of the SU(5) gauge group of the grand unified theory (GUT). The index $i = 1, \dots, N_M$ labels each set of messengers.

Messengers of each set directly couple to a supersymmetry breaking gauge singlet field S in the superpotential,

$$W = W_{\text{SUSY}} + kS\Psi_i\bar{\Psi}_i + M_\Psi\Psi_i\bar{\Psi}'_i + M_{\bar{\Psi}}\bar{\Psi}_i\Psi'_i, \quad (2.1)$$

where k denotes a coupling constant and $M_{\Psi, \bar{\Psi}}$ are mass parameters. W_{SUSY} encapsulates a dynamical SUSY breaking sector such as those in ref. [15–17] whose effective theory is

simply given by

$$W_{\text{SUSY}} \simeq \Lambda^2 S, \quad (2.2)$$

where Λ denotes the associated dynamical scale. Due to the linear term of S in the superpotential, the SUSY breaking field obtains a non-vanishing F -term expectation value. It should be noted that the form of the superpotential in eq. (2.1) is protected by an R -symmetry with the charge assignments $S(2)$, $\Psi_i(0)$, $\bar{\Psi}_i(0)$, $\Psi'_i(2)$ and $\bar{\Psi}'_i(2)$,² which gives the origin of the name of the R -invariant direct gauge mediation.³ Due to this peculiar form of the superpotential, the SUSY breaking vacuum is not destabilized by the couplings to the messenger fields.

It should be emphasized that the R -symmetry needs to be broken spontaneously to generate non-vanishing MSSM gaugino masses. Such spontaneous R -symmetry breaking can be achieved, for example, through a simple extension of the dynamical SUSY breaking model [15, 16] with an extra $U(1)$ gauge interaction [10]. See also refs. [18–23] for radiative R -symmetry breaking in more generic models.⁴ Altogether, we postulate that the SUSY breaking field S obtains its expectation value,

$$\langle S(x, \theta) \rangle = S_0 + F \theta^2, \quad (2.3)$$

where S_0 denotes the vacuum expectation value of the A -term of S and θ is the fermionic coordinate of the superspace.

Due to the stability of the SUSY breaking vacuum, the model is viable even when the reheating temperature of the Universe is very high. Thus, the model is consistent with thermal leptogenesis with $T_R \gtrsim 10^9 \text{ GeV}$ [27–30]. This feature should be compared with other types of direct gauge mediation models where a SUSY breaking vacuum is destabilized by messenger couplings such as

$$W = W_{\text{SUSY}} + k S \Psi_i \Psi_i, \quad (2.4)$$

(see e.g., ref. [31]). In such cases, the thermal history of the Universe and/or the masses of the messenger fields are severely restricted [7].⁵

2.2 Gauge mediated mass spectrum

We now summarize the gauge mediated mass spectrum of MSSM particles. The most distinctive feature of the MSSM spectrum in the R -invariant direct gauge mediation is that gaugino masses vanish at the one-loop level to the leading order of the SUSY breaking parameter, $\mathcal{O}(kF/M_{\text{mess}})$ [8, 9] and are suppressed by a factor of $\mathcal{O}(k^2 F^2/M_{\text{mess}}^4)$ in comparison with those in the conventional gauge mediation. In the following, we collectively

²The charges are assigned up to $U(1)$ messenger symmetries which are eventually broken by mixing with MSSM fields.

³To avoid completely stable messenger particles, more careful R -charge assignment is required as given in appendix A.

⁴It is also possible to construct O’Raifeartaigh models where spontaneous SUSY and R -symmetry breakings are achieved at *tree level* [24–26].

⁵For phenomenological studies of this class of models after the LHC Run-I experiment, see e.g., ref. [32].

denote the mass scale of the messenger sector by M_{mess} . Scalar masses, on the other hand, appear at the leading order of the SUSY breaking parameter at the two-loop level. Therefore, gauge mediated MSSM gaugino masses, M_a ($a = 1, 2, 3$), and MSSM scalar masses, m_{scalar} , are roughly given by

$$M_a \sim \frac{g_a^2}{16\pi} \frac{kF}{M_{\text{mess}}} \times \mathcal{O}\left(\frac{k^2 F^2}{M_{\text{mess}}^4}\right) \times (0.1 - 0.3), \quad (2.5)$$

$$m_{\text{scalar}} \sim \frac{g_a^2}{16\pi} \frac{kF}{M_{\text{mess}}}, \quad (2.6)$$

where g_a ($a = 1, 2, 3$) denote the gauge coupling constants of the MSSM gauge interactions. A factor of $\mathcal{O}(0.1)$ at the end of eq. (2.5) for the gaugino masses results from numerical analyses (see figure 1 and the following discussions). As a result, the predicted spectrum is hierarchical between gaugino masses and sfermion masses.

To date, searches for gluino pair production at the ATLAS and the CMS experiments have put severe lower limits on the gluino mass at around 1.4 TeV at 95% CL. The limits are applicable for cases where the bino either is stable [33, 34] or decays into a photon and a gravitino inside the detectors as the next-to-the lightest superparticle (NLSP) [35, 36]. To satisfy this constraint, we infer that

$$\frac{kF}{M_{\text{mess}}} \times \mathcal{O}\left(\frac{k^2 F^2}{M_{\text{mess}}^4}\right) = 10^{6-7} \text{ GeV}, \quad (2.7)$$

so that the gluino is sufficiently heavy (see eq. (2.5)).

Due to the hierarchy between the gaugino masses and sfermion masses in eqs. (2.5) and (2.6), the squarks are beyond the reach of the LHC Run-I when the gluino is heavier than 1.4 TeV. On the other hand, it should be noted that the squark masses are bounded from “above” by the correlation between the squark masses and the predicted lightest Higgs boson mass in the MSSM. In fact, unless the ratio of the Higgs vacuum expectation values, $\tan\beta$, is very close to unity, the scalar mass (especially the stop mass) should be around 10 – 100 TeV so that the lightest Higgs boson mass is consistent with the observed value [37–42], $m_h = 125.09 \pm 0.21 \pm 0.11 \text{ GeV}$ [43] (see discussions in section 3 for details). This requirement roughly leads to

$$\frac{kF}{M_{\text{mess}}} = 10^{6-7} \text{ GeV}. \quad (2.8)$$

Putting together conditions in eqs. (2.7) and (2.8), we find that the R -invariant direct gauge mediation is successful only when

$$\frac{kF}{M_{\text{mess}}} = 10^{6-7} \text{ GeV}, \quad (2.9)$$

$$\frac{kF}{M_{\text{mess}}^2} \sim 1. \quad (2.10)$$

In figure 1, we show a sample gauge mediated mass spectrum in the R -invariant direct gauge mediation model for $N_M = 1$. In the left plot, we show the spectrum as a function

Figure 1. The gauge mediated mass spectrum. The curves give the masses of various particles. In the figure, we take $M_\Psi = M_{\bar\Psi} = 2 \times 10^6$ GeV, $N_M = 1$ and $\tan\beta = 10$. In the left plot, we take $kS_0/\sqrt{M_\Psi M_{\bar\Psi}} = 1$ and show the spectrum as a function of $kF/(M_\Psi M_{\bar\Psi})$. In the right plot, we take $kF/(M_\Psi M_{\bar\Psi}) = 0.9$ and show the spectrum as a function of $kS_0/\sqrt{M_\Psi M_{\bar\Psi}}$.

of $kF/(M_\Psi M_{\bar\Psi})$ while fixing $M_\Psi = M_{\bar\Psi} = 2 \times 10^6$ GeV and $kS_0/\sqrt{M_\Psi M_{\bar\Psi}} = 1$. Here we take $\tan\beta = 10$ although the SUSY spectrum barely depends on $\tan\beta$. In our analysis, we use SOFTSUSY 3.6.2 [44] to calculate renormalization group evolution of soft parameters as well as to analyze the electroweak symmetry breaking conditions. The formulas of the gauge mediated spectrum at the messenger scale are given in ref. [9, 45]. As expected, the figure shows that gaugino masses become larger for a larger value of $kF/(M_\Psi M_{\bar\Psi})$, while the scalar masses are insensitive to this parameter. It should be noted that the messenger scalars are tachyonic for $kF/(M_\Psi M_{\bar\Psi}) > 1$. Therefore, the maximal gaugino masses are achieved for $kF/(M_\Psi M_{\bar\Psi}) \rightarrow 1$.

The right plot shows the mass spectrum as a function of $kS_0/\sqrt{M_\Psi M_{\bar\Psi}}$. Here we take $M_\Psi = M_{\bar\Psi} = 2 \times 10^6$ GeV and fix $kF/(M_\Psi M_{\bar\Psi}) = 0.9$. In the region of $kS_0/\sqrt{M_\Psi M_{\bar\Psi}} < 1$, gaugino masses increase with $kS_0/\sqrt{M_\Psi M_{\bar\Psi}}$, while scalar masses are less sensitive to $kS_0/\sqrt{M_\Psi M_{\bar\Psi}}$. This behavior comes from the fact that the gaugino masses require R -symmetry breaking, while the scalar masses do not. In the region of $kS_0/\sqrt{M_\Psi M_{\bar\Psi}} > 1$, both the gaugino masses and the scalar masses are decreasing with $kS_0/\sqrt{M_\Psi M_{\bar\Psi}}$. This is because the messenger scale is dominated by kS_0 as in the conventional gauge mediation in that region.

For subsequent discussions, we split the messenger fields of SU(5) GUT multiplets into the SM gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y$ representations: $\mathbf{5} = (\mathbf{3}_{-1/3}, \mathbf{2}_{1/2})$ and $\bar{\mathbf{5}} = (\bar{\mathbf{3}}_{1/3}, \mathbf{2}_{-1/2})$, such as $\Psi^{(\prime)} = (D^{(\prime)}, L^{(\prime)})$ and $\bar{\Psi}^{(\prime)} = (\bar{D}^{(\prime)}, \bar{L}^{(\prime)})$, respectively. Accordingly, we also distinguish the parameters in eq. (2.1) for each messenger using subscripts: k_D and k_L , $M_{D,\bar{D}}$ and $M_{L,\bar{L}}$, respectively.

In the analysis of refs. [11, 45], it is assumed that k 's and M 's satisfy the so-called GUT conditions at the GUT scale:

$$k_D = k_L, \quad M_D = M_L, \quad M_{\bar{D}} = M_{\bar{L}}. \quad (2.11)$$

In this paper, we do not impose these conditions in view of the fact that the doublets and the

triplets of the Higgs multiplets in the GUT models are required to split. In fact, the doublet-triplet splitting in the Higgs sector is most naturally achieved in GUT models with product gauge groups [46]. In those models, the GUT conditions in eq. (2.11) are not expected to be satisfied generically (see also ref. [47] for a recent discussion). In the following, we simply take k 's and M 's of the D and the L -type messengers as independent parameters.

It should be emphasized that the R -invariant direct mediation model is free from the CP -problem from the messenger interactions. The phases of $k_{D,L}$ and $M_{D,\bar{D},L,\bar{L}}$ can be absorbed by appropriate phase rotations of $D^{(\prime)}$, $\bar{D}^{(\prime)}$, $L^{(\prime)}$, and $\bar{L}^{(\prime)}$. The phases of S_0 and F can also be absorbed by the phases of S and θ^2 , respectively. In the above arguments, we have tacitly made use of these phase rotations to make kF 's, kS_0 's and M 's positive.⁶

Before closing this subsection, let us comment on the upper limit on N_M from the requirement of perturbative unification of gauge couplings. In the R -invariant direct gauge mediation model, the $N_M = 1$ case includes two pairs of $(\mathbf{5}, \bar{\mathbf{5}})$. Besides, the messenger scale is at around 10^{6-7} GeV as discussed above. Therefore, the number of messengers in the messenger sector is severely constrained by the perturbative unification to $N_M \leq 2$. It should also be noted that the messenger fields in $\mathbf{10}$ and $\bar{\mathbf{10}}$ representations are also disfavored by the perturbative unification due to the doubled number of messengers in the R -invariant direct gauge mediation. In the following, we confine ourselves to the cases of $N_M = 1$ and $N_M = 2$ by taking the perturbative gauge coupling unification seriously.

2.3 $B\mu$ -problem

In the above analysis, we have not specified the origin of the μ -term. In fact, it is the long-sought problem about how to generate the μ and $B\mu$ -terms of a similar size to other soft parameters while not causing the SUSY CP -problem. The minimal possibility for the origin of the μ -term is to assume that it is given just as is:

$$W = \mu H_u H_d, \quad (2.12)$$

where the R charge of the two Higgs doublets is 2. As a notable feature of this type of μ -term, the B -term at the messenger scale vanishes at the one-loop level:⁷

$$B \simeq 0. \quad (2.13)$$

It should be also emphasized that this minimal model is favorable since it does not bring about the SUSY CP -problem.

One may consider more direct couplings between the Higgs doublets and the SUSY breaking sector to generate μ and B terms, in order to interrelate the sizes of those parameters to other soft parameters. Naïve couplings between the SUSY breaking sector and the Higgs doubles, however, lead to too large a B -term, which is nothing but the infamous

⁶For $N_M \geq 2$, one may allow in eq. (2.1) couplings among fields with different labels i . Such couplings, however, lead to non-trivial phases on the parameters that result in relative phases to the gaugino masses and may bring about the SUSY CP -problem. Those label-changing couplings can be suppressed by introducing a (approximate) $U(1)$ messenger symmetry for each label i , for example (see also ref. [10]).

⁷A non-vanishing B -term is obtained from the two-loop threshold corrections of the messenger fields [48, 49] and is expected to be at around 10 GeV for the wino/bino masses around one TeV.

$\mu/B\mu$ -problem. More intricate connections between the Higgs and the SUSY breaking sector might be elaborated. In those models, one should be very careful to avoid the SUSY CP -problem.

In view of the minimality and the safety from the SUSY CP -problem, the minimal model of the μ -term in eq. (2.12) seems to be the most favorable candidate. In fact, many phenomenological studies have been done based on this minimal model of the μ -term in the conventional gauge mediation models [48, 50, 51]. As pointed out in ref. [11], however, the almost vanishing B -term at the messenger scale has a tension in the case of the R -invariant direct gauge mediation model as we see shortly.

In models with the almost vanishing B -term at the messenger scale, the B parameter at the stop mass scale is dominated by renormalization group effect:

$$\frac{dB}{d\ln\mu_R} = \frac{1}{16\pi^2} \left[6a_t y_t + 6a_b y_b + 6g_2^2 M_2 + \frac{6}{5}g_1^2 M_1 \right], \quad (2.14)$$

where μ_R is the renormalization scale, $y_{t,b}$ the top and the bottom Yukawa coupling constants, and $a_{t,b}$ the corresponding trilinear soft parameters. In the gauge mediation models, $a_{t,b}$ are also small and dominated by renormalization group effects from the gluino mass. Roughly, the radiatively generated B -term at the stop mass scale is estimated to be

$$|B(m_{\text{stop}})| \sim \frac{1}{16\pi^2} \left[6a_t y_t + 6a_b y_b + 6g_2^2 M_2 + \frac{6}{5}g_1^2 M_1 \right] \log \frac{M_{\text{mess}}}{m_{\text{stop}}} < O(0.1) \times M_2, \quad (2.15)$$

where we have taken the messenger scale, $M_{\text{mess}} = O(10^{6-7})$ GeV, and $m_{\text{stop}} \sim 10$ TeV. In our analysis, we use the convention that gaugino masses, $B\mu$, and $\tan\beta$ are positive-valued. From eq. (2.15), the radiatively generated $B(m_{\text{stop}})$ is negative-valued at the low energy scale. Thus, the sign of μ is negative in our convention.

The radiatively generated B -term is, generically, too small and renders too large a $\tan\beta$:

$$\tan\beta \simeq \frac{2}{\sin 2\beta} \simeq \frac{m_{H_u}^2 + m_{H_d}^2 + 2\mu^2}{B(m_{\text{stop}})\mu} \simeq \frac{m_{H_d}^2 + |m_{H_u}^2|}{|B(m_{\text{stop}})||m_{H_u}|} > \frac{2m_{H_d}}{|B(m_{\text{stop}})|} > O(100). \quad (2.16)$$

Such a large $\tan\beta$ leads to too large a bottom Yukawa coupling.⁸ In the third equality of eq. (2.16), we have used the electroweak symmetry breaking for a large $\tan\beta$,

$$\mu^2 = \frac{-m_{H_u}^2 \tan^2\beta + m_{H_d}^2}{\tan^2\beta - 1} + \frac{1}{2}m_Z^2 \simeq |m_{H_u}^2|. \quad (2.17)$$

We have also used $m_{H_u}^2 < 0$ and $m_{H_d}^2 > |m_{H_u}^2|$ which is valid for most parameter space. In the final inequality, we have used eqs. (2.5) and (2.15). It should be emphasized that this tension is due to the hierarchy between the gaugino mass and the scalar mass in the R -invariant direct gauge mediation.⁹

⁸More precisely, the Standard Model down-type Yukawa couplings are dominated by the radiatively generated non-holomorphic couplings to H_u [52] for such a large $\tan\beta$ (see also refs. [53, 54]). We do not preclude possibilities of the consistent parameter region with $\tan\beta = O(100)$. In this paper, however, we pursue the consistent parameter space with $\tan\beta \lesssim 50$ where the down-type Yukawa couplings come from the coupling to H_d in the superpotential. The possibilities with $\tan\beta = O(100)$ will be discussed elsewhere.

⁹The expected size of $\tan\beta$ is smaller for $N_M = 2$ compared with the case for $N_M = 1$, as the relative size of the gaugino mass (especially M_2) to the scalar mass becomes larger for $N_M = 2$.

Figure 2. The soft masses of H_u and H_d as functions of $\tan\beta$ for $N_M = 1$ (left) and $N_M = 2$ (right). The red and blue curves are m_{H_d} and $|m_{H_u}|$ ($m_{H_u}^2 < 0$) at the stop mass scale, respectively. The dashed curves are the corresponding soft masses at the messenger scale. The other parameters are explicitly indicated in each plot.

Figure 3. The B -term at the messenger scale (red curves) and the stop mass scale (blue curves) as functions of $\tan\beta$ for $N_M = 1$ (left) and $N_M = 2$ (right). The other parameters are indicated explicitly in each plot. The boundary condition $B(M) \simeq 0$ is satisfied for $\tan\beta \gtrsim 50$.

The above generic argument has a loophole. That is, we have assumed $m_{H_d}^2(m_{\text{stop}}) \simeq m_{H_d}^2(M_{\text{mess}})$. Although this relation is valid in most parameter space, it becomes invalid when $\tan\beta \gtrsim 50$ and the bottom Yukawa coupling y_b becomes comparable to the top Yukawa coupling. For such a large y_b , $m_{H_d}^2$ also receives a sizable negative contribution from the sbottom soft masses and gets smaller at the lower energy scale as $m_{H_u}^2$ does. When $m_{H_d}^2(m_{\text{stop}}) \ll m_{H_d}^2(M_{\text{mess}})$ is achieved, the resultant $\tan\beta$ can be much smaller than the one expected in eq. (2.16) and within a viable range. In this way, the R -invariant direct gauge mediation model can become consistent with the boundary condition with $B(M_{\text{mess}}) \simeq 0$.

In figure 2, we show m_{H_d} and $|m_{H_u}|$ as functions of $\tan\beta$. The plots show that $m_{H_d}^2(m_{\text{stop}}) \simeq m_{H_d}^2(M_{\text{mess}})$ for a moderate value of $\tan\beta$ as expected. For $\tan\beta \gtrsim 50$, on the other hand, the renormalization group effects on $m_{H_d}^2$ are sizable, and it becomes much smaller at the stop mass scale than at the messenger scale.

Armed with this observation, we have searched the parameter space for regions where the R -invariant direct gauge mediation is consistent with the boundary condition with

Figure 4. The required μ -term as functions of $\tan\beta$ for $N_M = 1$ (left) and $N_M = 2$ (right). The other parameters are indicated explicitly in each plot. The corresponding stop mass scales are $m_{\text{stop}} \simeq 12$ TeV ($N_M = 1$) and $m_{\text{stop}} \simeq 9$ TeV ($N_M = 2$), respectively.

an almost vanishing B -term. Figure 3 shows the B -term at the messenger scale (red curves) and the stop mass scale (blue curves) as functions of $\tan\beta$ in the R -invariant direct gauge mediation model. Here we impose the electroweak symmetry breaking condition with $m_Z \simeq 91.2$ GeV, instead of the boundary condition $B(M_{\text{mess}}) \simeq 0$. The plots show that the boundary condition, $B(M_{\text{mess}}) \simeq 0$, is compatible with the R -invariant model for $\tan\beta \gtrsim 50$, as expected.¹⁰ As a result, we find that the R -invariant direct gauge mediation model is consistent with the minimal μ -term.

It should be emphasized again that the consistency between the R -invariant direct gauge mediation model and the boundary condition $B(M_{\text{mess}}) \simeq 0$ is more difficult than in the case of the conventional gauge mediation. This difficulty stems from the hierarchy between the gaugino mass and the scalar mass as well as from the low messenger scale, $M \simeq 10^{6-7}$ GeV (see eq. (2.9)). In usual gauge mediation models, the messenger scale can be much larger while keeping the soft breaking mass scales in the TeV range, with which the radiatively generated B -term can be sizable due to a rather long interval of the renormalization group running. In the R -invariant direct gauge mediation model, on the other hand, one needs to take $k_L F / (M_L M_{\bar{L}})$ to be very close to 1, so that the gaugino mass, M_2 , takes a value as large as possible with which the radiatively generated B -term at the low energy becomes sizable.

Before closing this subsection, let us comment on the required size of μ for successful electroweak symmetry breaking. As we have argued in eq. (2.17), the required size of μ is roughly given by

$$\mu^2 \simeq -m_{H_u}^2(m_{\text{stop}}) \quad (2.18)$$

for a large $\tan\beta$. Here the Higgs soft mass squared, $m_{H_u}^2$, is approximately given by

$$m_{H_u}^2(m_{\text{stop}}) \sim m_{H_u}^2(M_{\text{mess}}) - \frac{12y_t^2}{16\pi^2} m_{\text{stop}}^2 \log \frac{M_{\text{mess}}}{m_{\text{stop}}} \quad (2.19)$$

¹⁰The perturbative bound on the bottom Yukawa coupling constant (and hence on $\tan\beta$) is largely relaxed in comparison with the pure MSSM due to the presence of the messenger fields. The resultant perturbative limit is $\tan\beta \simeq 80$ for $N_m = 1$.

at the stop mass scale, which can be much smaller than $m_{\text{stop}} \simeq 10 \text{ TeV}$ for $M_{\text{mess}} \simeq 10^{6-7} \text{ GeV}$. As a result, the required size of μ -term is also much smaller than m_{stop} . This feature somewhat eases the electroweak fine-tuning problem while explaining the observed Higgs boson mass by a heavy stop mass of $\mathcal{O}(10) \text{ TeV}$. In figure 4, we show the required size of the μ -term as a function of $\tan \beta$ by taking the same parameter sets used in figure 3. The figure shows that the required μ -term is indeed smaller than the stop mass scale. It should be also noted that a smaller μ -term is also possible when $m_{H_u}^2(M_{\text{mess}})$ is slightly larger at the messenger scale. This property is important for the discussions in section 4.

2.4 Gravitino dark matter

In gauge mediation models, the gravitino is the lightest supersymmetric particle (LSP). By assuming R -parity conservation, it can serve as a candidate for dark matter. In the regime of much lighter than MeV, the gravitino is thermalized in the early Universe, and its relic abundance is estimated to be

$$\Omega_{3/2} h^2 \simeq 0.1 \left(\frac{100}{g_*(T_D)} \right) \left(\frac{m_{3/2}}{100 \text{ eV}} \right). \quad (2.20)$$

Here $m_{3/2}$ is the gravitino mass, and $g_*(T_D) \simeq 100$ denotes the effective massless degree of freedom in the thermal bath at the decoupling temperature [55]

$$T_D \sim \max \left[M_{\tilde{g}}, 160 \text{ GeV} \left(\frac{g_*(T_D)}{100} \right)^{1/2} \left(\frac{m_{3/2}}{10 \text{ keV}} \right)^2 \left(\frac{2 \text{ TeV}}{M_3} \right)^2 \right]. \quad (2.21)$$

As discussed above, a successful R -invariant direct gauge mediation requires

$$(kF)^{1/2} = 10^{6-7} \text{ GeV}. \quad (2.22)$$

By assuming that the SUSY breaking field S breaks supersymmetry dominantly, the gravitino mass is given by

$$m_{3/2} \simeq 10 \text{ keV} \times \left(\frac{0.1}{k} \right) \left(\frac{(kF)^{1/2}}{2 \times 10^6 \text{ GeV}} \right)^2. \quad (2.23)$$

In this case, the thermally produced gravitino abundance in eq. (2.20) is too large to be consistent with the observed dark matter density.¹¹

This tension is removed when the above relic density is diluted by entropy production by a factor of

$$\Delta \simeq 100 \times \left(\frac{100}{g_*(T_D)} \right) \left(\frac{m_{3/2}}{10 \text{ keV}} \right) \quad (2.24)$$

after the gravitino decouples from the thermal bath.¹² As shown in ref. [10], an appropriate amount of entropy can be provided by, for example, the decay of long-lived particles in the

¹¹The gravitino with a mass $m_{3/2} \simeq 100 \text{ eV}$ is not cold dark matter but hot dark matter. Hence, it is not a viable candidate for dark matter even if the thermal relic abundance is consistent with the observed dark matter density.

¹²In general, if the dilution factor is provided by a late-time decay of a massive particle which dominates the energy density of the Universe, it is given by $\Delta \simeq T_{\text{dom}}/T_{\text{decay}}$, where T_{dom} is the temperature at which the massive particle dominates the energy density of the Universe and T_{decay} is its decay temperature. In order not to affect the Big-Bang Nucleosynthesis, we require $T_{\text{decay}} \gtrsim \mathcal{O}(1-10) \text{ MeV}$, and hence the dilution factor is bounded from above by $\Delta < T_{\text{dom}}/\mathcal{O}(1-10) \text{ MeV}$.

dynamical SUSY breaking sector.¹³ Interestingly, the gravitino in this mass range is a good candidate for a slightly warm dark matter [11] enabled via an appropriate dilution factor.

Finally, let us comment on the decay length of the NLSP, which is the bino in most parameter space of the R -invariant direct gauge mediation model. The bino NLSP mainly decays into a gravitino and a photon/ Z -boson with the branching ratios

$$\Gamma(\tilde{b} \rightarrow \psi_{3/2} + \gamma) \simeq \frac{1}{48\pi} \frac{M_1^5}{M_{PL}^2 m_{3/2}^2} \cos^2 \theta_W, \quad (2.25)$$

$$\Gamma(\tilde{b} \rightarrow \psi_{3/2} + Z) \simeq \frac{1}{48\pi} \frac{M_1^5}{M_{PL}^2 m_{3/2}^2} \sin^2 \theta_W. \quad (2.26)$$

Here $M_{PL} \simeq 2.4 \times 10^{18}$ GeV denotes the reduced Planck scale, and θ_W is the weak mixing angle. Altogether, the decay length of the bino NLSP is given by

$$c\tau_{\tilde{B}} \simeq 0.6 \text{ m} \times \left(\frac{500 \text{ GeV}}{M_1} \right)^5 \left(\frac{m_{3/2}}{10 \text{ keV}} \right)^2. \quad (2.27)$$

Therefore, the bino may or may not decay inside the detectors, depending on the gravitino mass and the NLSP mass.

3 Upper bound on the gluino mass

As alluded to before, the R -invariant direct gauge mediation model predicts a hierarchy between gaugino masses and scalar masses. We have also argued that $\tan \beta$ is required to be large, $\tan \beta \gtrsim 50$, if we further assume that the μ -term is provided by the minimal μ -term in the superpotential, eq. (2.12). For such a large $\tan \beta$, the stop mass is restricted to be around 10 TeV to account for the observed Higgs boson mass, $m_h = 125.09 \pm 0.21 \pm 0.11$ GeV [36]. Thus, by remembering that the gluino mass is limited from above for a given squark mass, the observed Higgs boson mass leads to an upper bound on the gluino mass.

To obtain the limit on the gluino mass from the observed Higgs boson mass, let us first consider the ratio between the gluino mass and the stop mass, which is shown as a function $k_DF/(M_DM_{\bar{D}})$ for $N_M = 1$ and $N_M = 2$ in figure 5. Here we take $k_DS_0 = (M_DM_{\bar{D}})^{1/2}$ to maximize the gluino mass, as given in figure 1. We have also checked that a heavier gluino mass cannot be achieved even for $M_D \neq M_{\bar{D}}$. The figure indicates that the ratio is a monotonically increasing function of $k_DF/(M_DM_{\bar{D}})$ and that $m_{\tilde{g}}/m_{\text{stop}} \lesssim 0.2$ for $N_M = 1$ and $m_{\tilde{g}}/m_{\text{stop}} \lesssim 0.3$ for $N_M = 2$.

From the observed Higgs boson mass, the stop mass is restricted to be in the $\mathcal{O}(10)$ TeV regime. The left plot of figure 6 shows the Higgs boson mass as a function of the stop mass for $\tan \beta \gtrsim 50$. In our analysis, we use **SusyHD** [59] to calculate the Higgs boson mass. The band represents uncertainties of our prediction, originating from (i) higher-order corrections of the Higgs mass calculation, (ii) uncertainties of Standard Model input parameters, and (iii) choices of model parameters that affect the SUSY spectrum other than stop masses. In

¹³A mass of 10^{6-7} GeV for the messenger is too light to provide a sufficient dilution factor [56] (see also appendix A). For other mechanisms of entropy production after the decoupling of gravitinos, see refs., [57, 58].

Figure 5. The ratio between the gluino mass and the stop mass, $m_{\text{stop}} \equiv \sqrt{m_{\tilde{t}_1} m_{\tilde{t}_2}}$, as a function of $k_D F / (M_D M_{\bar{D}})$ for $N_M = 1$ and $N_M = 2$. In the figure, we fix $M_D = M_{\bar{D}}$ and $k_D S_0 = M_D$ in order to optimize the gluino mass.

our analysis, we take the first uncertainty to be 1 GeV to make our discussions conservative (see also ref. [59]). As for the uncertainties from Standard Model parameters, the top mass uncertainty $m_t = 173.21 \pm 0.51 \pm 0.71$ GeV is the most important one for the Higgs mass prediction, and amounts to an error of about 0.3%. The effect of model parameters to which the stop mass is negligible. In the right plot, we show the stop mass dependence of $\Delta\chi^2$ defined by

$$\Delta\chi^2 = \frac{(125.09 - m_h)^2}{\sqrt{\sigma_{\text{ex}}^2 + \sigma_{\text{th}}^2}}, \quad (3.1)$$

where σ_{ex} denotes the experimental error and σ_{th} denotes the theoretical error as listed above. The 1σ (2σ) upper limit corresponds to $m_h \simeq 126.2$ GeV (127.2 GeV). From the plot, we find that $m_{\text{stop}} \gtrsim 21$ TeV is excluded at 2σ level (as indicated by the horizontal dashed line).¹⁴ Thus, by remembering that $m_{\tilde{g}}/m_{\tilde{t}} \lesssim 0.2$ for $N_M = 1$ and $m_{\tilde{g}}/m_{\tilde{t}} \lesssim 0.3$ for $N_M = 2$, we immediately find respectively $m_{\tilde{g}} \lesssim 4$ TeV and $m_{\tilde{g}} \lesssim 6$ TeV.

For a closer look, we show in figure 7 the gluino mass as a function of $k_D F / (M_D M_{\bar{D}})$ for $k_D F / (M_D M_{\bar{D}})^{1/2} \simeq 2.5 \times 10^6$ GeV and $N_M = 1$ and for $k_D F / (M_D M_{\bar{D}})^{1/2} \simeq 1.8 \times 10^6$ GeV and $N_M = 2$. Such parameter choices correspond to the upper limit of the stop mass, $m_{\text{stop}} \simeq 20$ TeV for each N_M . In the figure, we also show the expected 95% CL lower limits on the gluino mass at the 14-TeV LHC with the integrated luminosity 300 fb^{-1} and $m_{\tilde{g}} < 2.3$ GeV, at the HL-LHC with the integrated luminosity 3000 fb^{-1} and $m_{\tilde{g}} < 2.7$ GeV, and at a 33-TeV hadron collider with the integrated luminosity 3000 fb^{-1} and $m_{\tilde{g}} < 5.8$ GeV [64]. The figure shows that the LHC experiment will cover a large portion of the parameter space for $N_M = 1$ unless the gluino mass is highly optimized. The figure also shows that a 33-TeV hadron collider will cover almost the entire gluino mass range.

¹⁴If we use FeynHiggs 2.11.2 [60–63], we obtain a Higgs boson mass larger by about 4 GeV for m_{stop} in the $\mathcal{O}(10)$ TeV regime. This leads to a more stringent bound on the gluino mass.

Figure 6. The predicted Higgs boson mass as a function of the stop mass (left). $\Delta\chi^2$ as a function of the stop mass (right). The band in the left plot shows the theory uncertainties.

Figure 7. The upper limit on the gluino mass as a function of $k_D F / (M_D M_{\bar{D}})$ for $N_M = 1$ (left) and $N_M = 2$ (right). Each plot corresponds to the stop mass around 20 TeV, the upper limit obtained from the Higgs boson mass. The blue curves show the upper limits for an optimal R -symmetry breaking, $k_D S_0 = (M_D M_{\bar{D}})^{1/2}$, and the green curves show those for a less optimal R -symmetry breaking, $k_D S_0 = (M_D M_{\bar{D}})^{1/2}/2$. The dashed lines are the expected 95% CL limits on the gluino mass [64].

4 Z +jets+missing E_T

Recently, the ATLAS collaboration reported a 3σ excess in the search for events with a Z boson (decaying into a lepton pair) accompanied by jets and a large missing transverse energy (E_T^{miss}) [12].¹⁵ They observed 29 events in a combined signal region with di-electrons and di-muons at the Z -pole in comparison with an expected background of 10.6 ± 3.2 events. Although the significance of the signal is not sufficiently high at this point, many attempts have been made to explain the excess using the MSSM [13, 66–69]. The signal requires colored SUSY particles lighter than about 1.2 TeV [66]. In this section, we briefly discuss whether the R -invariant direct gauge mediation model can explain the reported signal.

As discussed in the previous section, the R -invariant direct gauge mediation model predicts a mass hierarchy between the gauginos and the scalars. Hence, the candidate col-

¹⁵Using the early data of LHC Run-II with an integrated luminosity of 3.2 fb^{-1} , the ATLAS experiment reported an excess in the same channel with a significance of 2.2σ [65]. With further data taking, the signal will be refuted/confirmed in the near future.

$\tilde{g}$	810	GeV	$\tilde{u}_L$	14	TeV
$\chi_1^0(\tilde{b})$	521	GeV	$\tilde{u}_R$	13	TeV
$\chi_2^0(\tilde{H})$	623	GeV	$\tilde{d}_L$	14	TeV
$\chi_3^0(\tilde{H})$	624	GeV	$\tilde{d}_R$	13	TeV
$\chi_4^0(\tilde{w})$	1.5	TeV	$\tilde{t}_1$	12	TeV
$\chi_1^\pm$	619	GeV	$\tilde{t}_2$	13	TeV
$\chi_2^\pm$	1.5	TeV	$\tilde{b}_1$	12	TeV
$\tilde{\ell}_L$	5.63	TeV	$\tilde{b}_2$	13	TeV
$\tilde{e}_R$	2.8	TeV	h	126.4	GeV
$\tilde{\tau}_1$	2.3	TeV	H	2.2	TeV
$\tilde{\tau}_2$	5.5	TeV	A	2.2	TeV
$\tilde{\nu}$	5.5	TeV	$H^\pm$	2.3	TeV

$\tan \beta$	55
$\text{sign}(\mu)$	-1
N_M	2
$M_D = M_{\bar{D}}$	1.965×10^6 GeV
$M_L = M_{\bar{L}}$	1.225×10^6 GeV
$k_D F / (M_D M_{\bar{D}})$	0.610
$k_L F / (M_L M_{\bar{L}})$	0.999
$m_t(\text{pole})$	173.21 GeV
$\alpha_s(m_Z)$	0.1185

Table 1. A sample mass spectrum which explains the Z -boson signal following ref. [13]. At this sample point, we take $k_{D,L}S_0 = M_{D,L}$, respectively. We also require $B(M_{\text{mess}}) \simeq 0$ by assuming minimal μ -term, as detailed in section 2. Due to the choice of $k_L F / (M_L M_{\bar{L}}) \simeq 0.999$, the lightest L -type messenger scalar particle is light, with a mass of around 27 TeV. We calculate the Higgs boson mass using *SusyHD* [59] whose theoretical uncertainties from higher-order corrections are estimated to be about 1 GeV (see also discussions in section 3).

ored SUSY particle required for the signal is inevitably the gluino. For such a light gluino, constraints from searches for SUSY particles with jets+ E_T^{miss} are usually severe, and most parameter region has been excluded [33, 34]. As shown in ref. [13], however, a careful study shows that the reported signal can be explained by the gluino production while evading all the other constraints. In this case, the gluino decays mainly into a gluon and a neutral Higgsino which subsequently decays into a Z -boson and a stable bino. Such a two-body decay of the gluino is induced by top-stop-loop diagrams, which can be dominant when the mass difference between the gluino and the Higgsino is sufficiently small and squark masses are much heavier than gaugino masses [13]. The dominance of the two-body decay mode is important to evade constraints from SUSY searches of multi-jets+ E_T^{miss} . It is also important to suppress the decay of the neutral Higgsino into a Higgs boson and a bino by requiring $M_{\tilde{H}} - M_{\tilde{b}} \lesssim 100$ GeV. In addition, preventing the wino from appearing in the gluino decay chain is also important to avoid constraints from other SUSY searches.

To attain a light Higgsino, we remind readers that the μ -term is related to $m_{H_u}^2$ via the electroweak symmetry breaking condition

$$\mu^2 \simeq -m_{H_u}^2(m_{\text{stop}}) \simeq -m_{H_u}^2(M_{\text{mess}}) + \frac{12y_t^2}{16\pi^2} m_{\text{stop}}^2 \log \frac{M_{\text{mess}}}{m_{\text{stop}}}. \quad (4.1)$$

Here we assume $\tan \beta \gtrsim 50$, as discussed in the previous section. This relation shows that a slightly larger $m_{H_u}^2(M_{\text{mess}})$ for a given m_{stop} at the messenger scale results in a smaller μ -term and hence lighter Higgsinos. A larger $m_{H_u}^2$ also corresponds to larger bino and

wino masses, which are also favorable to explain the signal. Through a careful parameter choice, we find that the desired spectrum can be achieved, as given in table 1), where the Higgsino masses are placed between the gluino mass and the bino mass.¹⁶

At the model point in table 1, the gluino decay is dominated by the radiatively induced two-body modes with the branching ratios

$$Br(\tilde{g} \rightarrow \tilde{H} + g) \simeq 0.47, \quad (4.2)$$

$$Br(\tilde{g} \rightarrow \tilde{B} + g) \simeq 0.06. \quad (4.3)$$

Here we use SDECAY v1.3 [70] to calculate the decay widths of MSSM particles. It should be noted that the MSSM parameters in table 1 is not optimal for the dominance of the two-body decay modes. Thus, a rather light gluino is required to account for the observed signals [13]. The production cross section for a pair of gluinos is given by

$$\sigma = 132 \pm 13 \text{ fb}, \quad (4.4)$$

as calculated at the next-to-leading-logarithmic accuracy by NLL-Fast v1.2 [71–76].

For a simplified estimate, we rely on the analyses given in CheckMATE v1.2.1 [77, 78] which incorporates DELPHES 3 [79] and FastJet [80, 81]. We generate signal events using MadGraph5 v2.2.3 [82, 83] connected to Pythia 6.4 [84]. The MLM matching scheme is used with a matching scale at 150 GeV [85]. We choose CTEQ6L1 [86] for the parton distribution functions. As a result, we obtain about 10 events in the signal region [12], while evading the constraints from the multi-jets+ E_T^{miss} search [87], the mono-jet search [88], as well as the CMS on- Z search [89] at 95%CL.¹⁷ Therefore, the model parameters in table 1 can successfully provide signal events consistent with the excess at 1.4σ level.¹⁸

5 Summary

In this paper, we revisited a spacial model of gauge mediated supersymmetry breaking, the “ R -invariant direct gauge mediation.” The model is favorable as it is durable even when the reheating temperature of the Universe is very high. We paid particular attention to the consistency of the model with the minimal model addressing the origin of the μ -term. As a result, we found that the minimal model can be consistent with the R -invariant gauge mediation model with a careful choice of model parameters, although incompatibility was highlighted in view of the current experimental constraints on superparticle masses and the observed Higgs boson mass. We also found that the μ -term was generically smaller than the stop mass, which might ease the electroweak fine-tuning problem while explaining the observed Higgs boson mass with a heavy stop mass of $\mathcal{O}(10)$ TeV.

¹⁶We assume a somewhat heavy gravitino, $m_{3/2} \gtrsim 100 \text{ keV}$, so that the bino NLSP is stable inside the detector (see eq. (2.27)).

¹⁷Here we do not take into account the constraints on the ZZ mode from the four lepton + E_T^{miss} searches [90, 91]. Due to smaller branching ratios of the gluino into a Higgsino and a gluon in our model, the constraints from those searches are weaker than the one discussed in ref. [13]. We have also confirmed that the constraint from the Z +dijet+ E_T^{miss} searches [92, 93] is less important.

¹⁸Here, we consider the p -value corresponding to the probability that the signal+background can explain the observed event numbers consistently, and 1.4σ corresponds to $p \simeq 0.16$ (see e.g., ref. [94]).

We found that there existed an upper limit on the gluino mass from the observed Higgs boson mass when the μ -term was given by the minimal model. Due to a hierarchy between gaugino masses and sfermion masses as well as the requirement for a large $\tan\beta$, the observed Higgs boson mass led to an upper limit on the stop mass of about 20 TeV and a corresponding upper limit on the gluino mass of about 4 TeV. This result is encouraging because the LHC experiment will be able to cover a large portion of the parameter space unless the model parameters are highly optimized to achieve a large gluino mass. This situation is parallel to, for example, high-scale supersymmetry breaking models with anomaly mediated gaugino mass [95–98] such as pure gravity mediation model/minimal split SUSY [99–102] (see also e.g. ref. [103]), which also predicts that the gluino is within the reach of the future collider experiments [104–106], while explaining the observed Higgs boson mass with a large m_{stop} .

We also discussed whether the R -invariant direct gauge mediation model could explain the 3σ excess of the $Z+\text{jets}+E_T^{\text{miss}}$ events reported by the ATLAS collaboration [12]. With carefully chosen parameters, we found it possible to explain the excess and the masses of Higgsinos were placed in between those of gluino and bino.¹⁹

Finally, we comment on some ideas that provide the appropriate size of the μ -term. In our analysis, we have only discussed that the required size of the μ -term for a successful electroweak symmetry breaking is in the TeV range or smaller. Since we assume that the μ -term is consistent with the R -symmetry, the smallness of the μ -term requires some additional symmetry. One popular idea is to generate the μ -term from the breaking of a Peccei-Quinn symmetry [108] via a dimension-5 operator [109, 110].²⁰ As another possibility, we propose to make use of a Z_2 symmetry (we name here 10-parity) under which only H_d and the MSSM matter fields incorporated in the **10** representation of SU(5) GUT group change their signs:

$$\begin{aligned} (Q_L, \bar{U}_R, \bar{E}_R) &\rightarrow -(Q_L, \bar{U}_R, \bar{E}_R), & (\bar{D}_R, L_L) &\rightarrow (\bar{D}_R, L_L), & \bar{N}_R &\rightarrow \bar{N}_R, \\ H_u &\rightarrow H_u, & H_d &\rightarrow -H_d. \end{aligned} \quad (5.1)$$

In this case, the small μ -term can be explained by a tiny breaking of the 10-parity.

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¹⁹As another interesting feature of the R -invariant direct gauge mediation model, it is often accompanied by a pseudo Nambu-Goldstone boson associated with R -symmetry breaking, the R -axion. With the gluino mass range suggested by the $Z+\text{jets}+E_T^{\text{miss}}$, it is also possible to search for the R -axion which can be produced via gluon-fusion at the LHC experiment [107].

²⁰The scalar partner of the axion could cause some cosmological problem in the gauge mediation scenario, although we do not go into details in this paper.

	S	Ψ	$\bar{\Psi}$	Ψ'	$\bar{\Psi}'$	H_u	H_d	$\mathbf{10}_{\text{MSSM}}$	$\bar{\mathbf{5}}_{\text{MSSM}}$	$\bar{N}_R$
R	2	11/5	-11/5	-1/5	21/5	4/5	6/5	3/5	1/5	1

Table 2. The R -charge assignments that allow the messenger-SUSY breaking interactions in eq. (2.1), the minimal μ -term, the MSSM Yukawa interactions, the mass term of the right-handed neutrinos, and the messenger-matter mixing in eq. (A.1).

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A R -charge assignments

For the models to be cosmologically viable, the messenger fields are required to decay first enough in the thermal history of the Universe. In this appendix, we summarize the R -charge assignments that allow the appropriate mixing between the messengers and the MSSM fields.

Explicitly, we seek R -charge assignments which allow a tiny mixing via

$$W \sim \langle W \rangle \Psi' \bar{\mathbf{5}}_{\text{MSSM}}. \quad (\text{A.1})$$

Hereafter, we use SU(5) GUT representations for the MSSM matter fields: $\bar{\mathbf{5}}_{\text{MSSM}} = (\bar{D}_R, L_L)$ and $\mathbf{10} = (Q_L, \bar{U}_R, \bar{E}_R)$. By assuming that the messenger-SUSY breaking interactions given in eq. (2.1), the minimal μ -term, the MSSM Yukawa interactions and the mass term of the right-handed neutrinos are consistent with the R -symmetry, leading to the R -charge assignments given in table 2. Here the charge assignments for the messenger fields are different from the one discussed in section 2.1, which can be obtained by appropriately mixing the R -symmetry and messenger rotation. It should be noted that the resultant mixing masses are of $\mathcal{O}(10)$ keV, and hence, they do not affect the messenger spectrum or the gauge mediated soft masses.

Through the tiny mixing term, the lightest messenger decays into MSSM particles with a decay width

$$\Gamma_{\text{mess}} \sim \frac{g_a^2}{16\pi} \frac{m_{3/2}^2}{M_{\text{mess}}}, \quad (\text{A.2})$$

which corresponds to the decay temperature

$$T_{\text{decay}} \sim \mathcal{O}(1) \text{ GeV} \left(\frac{m_{3/2}}{10 \text{ keV}} \right) \left(\frac{10^6 \text{ GeV}}{M_{\text{mess}}} \right)^{1/2}. \quad (\text{A.3})$$

This decay temperature is much higher than the temperature at which the messenger field would dominate over the energy density of the Universe,

$$T_{\text{dom}} \sim M_{\text{mess}} Y_{\text{mess}}, \quad (\text{A.4})$$

	S	Ψ	$\bar{\Psi}$	Ψ'	$\bar{\Psi}'$	H_u	H_d	$\mathbf{10}_{\text{MSSM}}$	$\bar{\mathbf{5}}_{\text{MSSM}}$	$\bar{N}_R$
R	2	$-1/5$	$1/5$	$9/5$	$11/5$	$4/5$	$6/5$	$3/5$	$1/5$	1

Table 3. The R -charge assignments when Ψ' in eq. (A.1) is replaced by Ψ .

where the thermal yield of the lightest messenger (the doublet messenger) [56]

$$Y_{\text{mess}} \sim 10^{-10} \left(\frac{M_{\text{mess}}}{10^6 \text{ GeV}} \right). \quad (\text{A.5})$$

Finally, let us note that the above choice in table 2 is not the unique choice for cosmologically viable models. For example, we may consider the charge assignment which allows the mixing term with Ψ' in eq. (A.1) replaced by Ψ , leading to different R -charge assignments (see table 3).

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