Abstract

We study anomaly mediation models with gauge mediation effects from messengers which have a general renormalizable mass matrix with a supersymmetry-breaking spurion. Our models lead to a rich structure of supersymmetry breaking terms in the visible sector. We derive sum rules among the soft scalar masses for each generation. Our sum rules for the first and second generations are the same as those in general gauge mediation, but the sum rule for the third generation is different because of the top Yukawa coupling. We find the parameter space where the tachyonic slepton problem is solved. We also explore the case in which gauge mediation causes the anomalously small gaugino masses. Since anomaly mediation effects on the gaugino masses exist, we can obtain viable mass spectrum of the visible sector fields.
1 Introduction

Supersymmetry (SUSY) is one of the promising candidates for physics beyond the standard model. It can solve the hierarchy problem on the Higgs mass. In addition, the standard model gauge couplings are unified at a high-energy scale in the minimal supersymmetric extension of the standard model (MSSM). Moreover, the lightest supersymmetric particle (LSP) is a strong candidate for the dark matter in our universe.

If SUSY is realized in nature, it must be broken at an energy scale above the weak scale. We usually leave its dynamics to the hidden sector different from our visible sector. Anomaly-mediated supersymmetry breaking \cite{1,2} is one of the attractive mechanisms to transmit the SUSY breaking of the hidden sector into the visible sector fields without flavor problems. In anomaly mediation, the SUSY breaking in the hidden sector is encoded in the F-component of the superconformal compensator $\phi = 1 + \theta^2 \langle F_\phi \rangle$. The soft SUSY breaking parameters in the visible sector are given by the 1-loop suppressed form of the parameter $\langle F_\phi \rangle$. The gravitino naturally has $\mathcal{O}(10)$ TeV of a mass. Unfortunately, anomaly mediation suffers from the problem that the slepton masses become tachyonic. Then, we need to modify the original form by some additional effects in order to obtain a successful model \cite{3,4}.

On the other hand, in gauge-mediated supersymmetry breaking \cite{5,6}, the SUSY breaking of the hidden sector is transmitted into the visible sector by the standard model gauge interactions. As in anomaly mediation, the unwanted flavor-changing processes are suppressed due to the flavor blindness of the gauge interactions and hence gauge mediation is also considered to be a promising mediation mechanism of the SUSY breaking. However, in gauge mediation, we often encounter the anomalously small gaugino masses compared to the scalar masses. In this case, we cannot obtain $\mathcal{O}(1)$ TeV of the gaugino masses and the scalar masses at the same time. That may cause a hierarchy problem, again.

Both anomaly mediation and gauge mediation are quite interesting, but both may have problems. In particular, the pure anomaly mediation has the problem in the slepton sector, while gauge mediation may have the problem in the gaugino masses. Then, a natural approach to these problems would be to mix anomaly mediation and gauge mediation \cite{3,11,12,13,14}. The contributions to the soft masses from anomaly mediation and gauge mediation can be naturally comparable when we derive the mass scale of the messengers of gauge mediation from the parameter $\langle F_\phi \rangle$. The tachyonic slepton problem in the pure anomaly mediation can be cured by the contribution from gauge mediation while the anomalously small gaugino masses can be enhanced by the anomaly mediation contribution.

\footnote{The anomalously small gaugino mass problem can be seen in direct gauge mediation models \cite{7,8} (see also \cite{9}) and semi-direct gauge mediation models \cite{10}.}
In this paper, we extend the model of \cite{12}, which mix anomaly mediation and gauge mediation, to models with the following generalized fermion mass matrix of the messenger fields \cite{15}:

$$L_{mess} = \int d^2 \theta \phi M_{ij}(X) \psi_i \tilde{\psi}_j + h.c. = \int d^2 \theta \phi (\lambda_{ij} X + m_{ij}) \psi_i \tilde{\psi}_j + h.c.,$$ \hspace{1cm} (1.1)

where $X = \langle X \rangle + \theta^2 \langle F_X \rangle$ is a SUSY breaking spurion and $\lambda_{ij}, m_{ij}$ are constant matrices, which are in general independent of each other. The fields $\psi_i, \tilde{\psi}_j (i = 1, \ldots, N)$ denote the messengers which belong to the (anti-)fundamental representations under the $SU(5)$ group into which the standard model gauge symmetry is embedded. We can further generalize the messenger sector such that the superpotential has the doublet/triplet splitting as follows:

$$L'_{mess} = \int d^2 \theta \phi \left[ M^2_{ij}(X) \ell_i \tilde{\ell}_j + M^3_{ij}(X) q_i \tilde{q}_j \right] + h.c. = \int d^2 \theta \phi \left[ (\lambda_{2ij} X + m_{2ij}) \ell_i \tilde{\ell}_j + (\lambda_{3ij} X + m_{3ij}) q_i \tilde{q}_j \right] + h.c.,$$ \hspace{1cm} (1.2)

where $\ell_i, \tilde{\ell}_j$ and $q_i, \tilde{q}_j$ are $SU(2)$ doublets and $SU(3)$ triplets of the messengers respectively. We here take this general case with some conditions for simplicity. In this model, we study the soft mass spectrum of the model and identify the LSP.

The rest of the paper is organized as follows. In section 2, we will present our model which gives the messenger mass matrices \cite{11, 12} and analyze the soft mass spectrum of the visible sector fields. In section 3, we will show the numerical analyses of the soft masses. In addition, we will derive sum rules among the soft scalar masses for each generation. In section 4, we will conclude the discussions.

2 Generalities

In this section, we first show a model which leads to the messenger mass matrix \cite{11, 12} and analyze its vacuum structure, following the discussion of \cite{12}. Then, we present the soft mass formulae of the visible sector fields derived from anomaly mediation and gauge mediation.

2.1 The models

In addition to the usual canonical Kähler potential, we consider the following terms in Lagrangian of the messenger fields $\psi_i, \tilde{\psi}_j$ and the singlet field $S$:

$$\Delta L = \int d^4 \theta \frac{\phi^4}{\phi} \left( \frac{1}{2} c_S S^2 + c_P \psi_i \tilde{\psi}_j \right) + \int d^2 \theta \left[ \frac{\lambda_S}{3!} S^3 + \lambda_P \psi_i \tilde{\psi}_j \right] + h.c.,$$ \hspace{1cm} (2.1)
where $c_S, \lambda_S, c_{Pij}$ and $\lambda_{Pij}$ are the real coupling constants. Here, we assume that the quadratic terms of $\psi_i, \tilde{\psi}_j$ and $S$ are absent. From the above Lagrangian, the scalar potential of this model is given by

$$
V = \left| c_S \langle F^\dagger \phi \rangle S + \frac{1}{2} \lambda_S S^2 + \lambda_{Pij} \psi_i \tilde{\psi}_j \right|^2 \\
+ \left| \left( c_{Pij} \langle F^\dagger \phi \rangle + \lambda_{Pij} S \right) \psi_i \right|^2 + \left| \left( c_{Pij} \langle F^\dagger \phi \rangle + \lambda_{Pij} S \right) \tilde{\psi}_j \right|^2 \\
+ \left| \langle F_\phi \rangle \right|^2 \left( \frac{1}{2} c_S S^2 + c_{Pij} \psi_i \tilde{\psi}_j \right) + h.c.
$$

(2.2)

We next consider the minimization of this potential. We assume that the messenger fields $\psi_i, \tilde{\psi}_j$ stabilize at the origin of their field space, $\langle \psi_i \rangle = \langle \tilde{\psi}_j \rangle = 0$ to preserve the standard model gauge symmetry. The expectation values of $S$ and $F_\phi$ on the minimum of the potential are then given by

$$
\langle S \rangle = -\frac{\langle F_\phi \rangle}{2 \lambda_S} \left( 3 c_S + \sqrt{c_S (c_S - 8)} \right),
$$

$$
\frac{\langle F_\phi \rangle}{\langle S \rangle} = -\frac{c_S + \sqrt{c_S (c_S - 8)}}{4 \lambda_P}.
$$

(2.3)

Note that this vacuum is the global minimum of the potential in a certain parameter range. Then, these vacuum expectation values lead to the following mass term for the messenger fields $\psi_i, \tilde{\psi}_j$:

$$
L'_{text{mess}} = \int d^2 \theta \phi \left( M_{ij} + F_{ij} \theta^2 \right) \psi_i \tilde{\psi}_j,
$$

(2.4)

where we define

$$
M_{ij} = \langle F_\phi \rangle c_{Pij} + \langle S \rangle \lambda_{Pij},
$$

$$
F_{ij} = -2 \langle F_\phi \rangle M_{ij} + \left( \langle F_\phi \rangle \langle S \rangle + \langle F_\phi \rangle \right) \lambda_{Pij}.
$$

(2.5)

Note that the matrix $M_{ij}$ is not proportional to the matrix $F_{ij}$ in general. Thus, we can obtain the form of the mass term of the messengers. The messenger scale is naturally the same order as the scale $\langle F_\phi \rangle$ when all the couplings in the model are of $O(1)$ and hence the anomaly-mediated and the gauge-mediated contributions are comparable. If we tune the parameters, we can realize the cases where the effect of anomaly mediation is dominant or that of gauge mediation is dominant. Since the vacuum we consider here is the global minimum

2 We can forbid these terms by a discrete $R$ symmetry such as $S(\theta) \rightarrow -S(i\theta)$, $\psi_i(\theta) \rightarrow -\psi_i(i\theta)$, $\tilde{\psi}_j(\theta) \rightarrow -\tilde{\psi}_j(i\theta)$ with the other fields even.

3 See [12, 13] for a detail of the potential analysis. In [13], it is pointed out that there is a UV divergent 1-loop linear term of $S$. However, such a term does not affect our results, because we parametrize $\langle S \rangle$ and $\langle F_\phi \rangle$ for our phenomenological purpose.
of the potential \(2.2\), we may worry about the consistency with the discussion of \([8]\) where the pseudomoduli space of the SUSY breaking vacuum cannot be locally stable everywhere in order to generate sizable gaugino masses in direct gauge mediation. However, there is no pseudomoduli space in our set-up and hence the discussion of \([8]\) cannot be applied as discussed in \([9]\). In this case, we can obtain sizable gaugino masses in the global minimum of the potential. The models discussed in \([9, 16, 17]\) with the minimal gauge mediation \([18]\) have the messenger sector separated from the SUSY breaking sector and they have the additional messenger gauge interaction (or the nonrenormalizable interaction in the Kähler potential) between these two sectors. Then, these models do not have a pseudomoduli space in the messenger sector and realize sizable gaugino masses in the global vacuum. In the present set-up, the messenger gauge interaction in the models of \([9, 16]\) is replaced to the interactions of the conformal compensator field \(\phi\) with the messenger sector fields. Then, the SUSY breaking in the messenger sector mediated by the compensator can generate nonzero leading order gaugino masses in the global vacuum.

We here comment on the difference between the model in \([15]\) and our set-up. In the model of \([15]\), the nontrivial R-charge assignment on the SUSY breaking field \(X\) and the messenger fields \(\psi_i, \tilde{\psi}_j\) restricts the determinant of the matrix \(\mathcal{M}(X)\) to the following form:

\[
\det \mathcal{M} = X^n G(m, \lambda), \quad n = \frac{1}{R(X)} \sum_{i=1}^{N} (2 - R(\psi_i) - R(\tilde{\psi}_i)),
\]

where \(G(m, \lambda)\) is some function of the coupling constants \(m, \lambda\) and \(R(X), R(\psi_i), R(\tilde{\psi}_j)\) are the R-charges of the fields \(X, \psi_i, \tilde{\psi}_j\). Then, when we introduce the doublet/triplet splitting into the messenger sector but take the same R-charge assignments of the \(SU(2)\) doublet and \(SU(3)\) triplet parts of the messengers, the following GUT relation among the gaugino masses is preserved:

\[
M_1 : M_2 : M_3 = \alpha_1 : \alpha_2 : \alpha_3,
\]

where \(M_1, M_2, M_3\) are the masses of the bino, wino and gluino fields respectively. On the other hand, in our model, any condition is not imposed on the coupling constants \(c_S, \lambda_S, c_{Pij}\) and \(\lambda_{Pij}\). Then, we have the general messenger mass matrix \(\mathcal{M}(X)\) and the determinant. Thus, in general, we do not have the above relation of the gaugino masses when we introduce the doublet/triplet splitting in the messenger sector. We can realize various ratios between the gaugino masses in our model. For simplicity, in order to analyze numerically spectra of our models in section 3, we require the same structures for the mass matrices of the \(SU(2)\) doublet and \(SU(3)\) triplet messengers \(\ell_i, \tilde{\ell}_j\) and \(q_i, \tilde{q}_j\) and preserve the GUT relation among the gaugino masses.
2.2 The soft mass formulae

We next show the formulae of the soft mass parameters of the visible sector derived from the anomaly mediation effects and the gauge mediation effects of the messengers whose fermion mass matrix is given by \((1.1)\). First, we calculate the gaugino masses which can be read off from the holomorphic gauge coupling \(\tau\) dependent on the spurion chiral superfield \(X\) and the compensator chiral superfield \(\phi\). We can write the gaugino mass as the following form:

\[
M_\lambda = \frac{i}{2\tau} \left( \frac{\partial \tau}{\partial \phi} \bigg|_{\phi=1} F_\phi + \frac{\partial \tau}{\partial X} \bigg|_{X=\langle X \rangle} F_X \right).
\]

(2.8)

The holomorphic gauge coupling at a scale \(\mu\) below the messenger scale is given by

\[
\tau(\mu) = \tau_0 + \frac{b'}{2\pi} \log \frac{1}{\Lambda} - \frac{i}{2\pi} \log \det \mathcal{M} + \frac{b}{2\pi} \log \frac{\mu}{\phi},
\]

(2.9)

where \(b'\) is the \(\beta\) function coefficient of the theory including the messenger fields while \(b\) is the \(\beta\) function coefficient in the effective theory below the mass scale of the messengers. The constant \(\Lambda\) is the cutoff scale of the model and \(\tau_0\) is the value of the coupling at that scale. Inserting this expression of the holomorphic coupling into (2.8), the gaugino mass is given by

\[
M_\lambda = \frac{\alpha}{4\pi} \left( bF_\phi + \frac{\partial}{\partial X} \log \det \mathcal{M} \bigg|_{X=\langle X \rangle} F_X \right),
\]

(2.10)

where \(\alpha \equiv g^2/4\pi\). The first term is considered to be the anomaly-mediated contribution and the second term is the gauge mediation contribution.

We next derive the soft scalar mass of the matter field in the visible sector, which can be read off from the wavefunction renormalization factor,

\[
Z = Z \left( \frac{\mu}{\Lambda|\phi|}, \frac{|X|}{\Lambda} \right),
\]

(2.11)

which is the function of the combinations of \(\mu/\Lambda|\phi|\) and \(|X|/\Lambda\). Then, the soft scalar mass of the matter field can be expressed as follows:

\[
m_Q^2 = -\frac{1}{4} \frac{\partial^2 \log Z}{\partial (\log \mu)^2} |F_\phi|^2 - \frac{1}{4} \frac{\partial^2 \log Z}{\partial (\log |X|)^2} \left| \frac{F_X}{X} \right|^2
\]

\[
+ \frac{1}{4} \frac{\partial^2 \log Z}{\partial \log \mu \partial \log |X|} F_\phi \frac{F_X^\dagger}{X^\dagger} + h.c.
\]

(2.12)

The first term is considered to be the anomaly-mediated contribution and here we can replace the derivative of the compensator field \(\phi\) to the derivative of the scale \(\mu\) because of the factor dependence of the wavefunction renormalization factor \((2.11)\). The second term is the gauge-mediated contribution. The rest terms are the mixing terms of both contributions and hence
the final result of the soft scalar mass is not the simple sum of the anomaly-mediated effect and the gauge-mediated effect. We deal with three contributions to the soft scalar mass one by one. The first term of the anomaly-mediated contribution can be rewritten in terms of the anomalous dimension $\gamma$ and the $\beta$ function at the scale $\mu$ such as

$$-\frac{1}{4} \frac{\partial^2 \log Z}{\partial (\log \mu)^2} |F_\phi|^2 = -\frac{1}{4} \left( \frac{\partial \gamma}{\partial g} \beta_g + \frac{\partial \gamma}{\partial y} \beta_y \right) |F_\phi|^2,$$

(2.13)

where $y$ is the Yukawa coupling. At the 1-loop order, the anomalous dimension and the $\beta$ functions of the gauge coupling and the Yukawa coupling are given by

$$\gamma = \frac{1}{16 \pi^2} \left( 4 C_2 g^2 - a y^2 \right),$$

$$\beta_g = -\frac{b g^3}{16 \pi^2},$$

$$\beta_y = \frac{y}{16 \pi^2} \left( e y^2 - f g^2 \right),$$

(2.14)

where $C_2$ is a quadratic Casimir and other coefficients in the above expressions are summarized in the Appendix. Next, we consider the rest of the terms in (2.12). Here, we can ignore the $X$ dependence of the Yukawa coupling in the above expression of the anomalous dimension $\gamma$. We assume that the mass eigenvalues of the messenger fields are the same order and take a common messenger scale $M_{mess}$. Then, we finally obtain the following soft scalar mass at the messenger scale $M_{mess}$:

$$m^2_{Q} = \left[ 2 b C_2 \left( \frac{\alpha}{4 \pi} \right)^2 + \frac{1}{2} a \frac{y^2}{(4 \pi)^2} \left( e \frac{y^2}{(4 \pi)^2} - f \frac{\alpha}{4 \pi} \right) \right] |F_\phi|^2$$

$$+ 2 C_2 \left( \frac{\alpha}{4 \pi} \right)^2 \sum_i \left( \frac{\partial \log |a_i|}{\partial \log |X|} \right)^2 \left| \frac{F_X}{X} \right|^2$$

$$+ 2 C_2 \left( \frac{\alpha}{4 \pi} \right)^2 \frac{\partial}{\partial \log |X|} \log |\det \mathcal{M}| |F_\phi F_X^{|X|} + h.c.,$$

(2.15)

where $a_i$ is the eigenvalue of the messenger mass matrix $\mathcal{M}$. In appendix, we write down the explicit soft masses of the MSSM fields. In the next section and Appendix, we set only the top Yukawa coupling $y_t$ non-vanishing, but the other Yukawa couplings vanishing.

3 Spectrum and phenomenology

In this section, we show the soft mass spectrum in the MSSM numerically by using the formula shown in the previous section. We investigate two cases where the leading contribution to the gaugino mass from gauge mediation is zero or nonzero. We also give a comment on the $\mu$-term and $B$-term.
3.1 Numerical analyses

As discussed in the introduction, the messenger fields generally do not form complete SU(5) multiplets as far as they preserve the unification of the standard model gauge couplings. The SU(2) doublet and SU(3) triplet parts of the messengers can have different supersymmetric masses and SUSY breaking mass splittings. We here consider this general situation and take the different couplings $c_2^{P_{ij}}, c_3^{P_{ij}}$ and $\lambda_2^{P_{ij}}, \lambda_3^{P_{ij}}$ for the doublet and triplet parts of the messenger fields to obtain the general spectra of the visible sector fields. Then, we have five continuous parameters $F_\phi, \Lambda_2^g, \Lambda_3^g, \Lambda_2^X, \Lambda_3^X$ defined as follows:

$$\Lambda_{2,3}^{g,x} = \frac{\partial}{\partial X} \log \det \mathcal{M}_{2,3} \bigg|_{X=(X)} F_X,$$

$$(\Lambda_{X}^{2,3})^2 = \sum_i \left( \frac{\partial \log |a_i^{2,3}|}{\partial \log |X|} \right)^2 \frac{F_X}{X}^2 + \frac{\partial}{\partial \log |X|} \log |\det \mathcal{M}_{2,3}| F_\phi \frac{F_X^*}{X} + h.c. \quad (3.1)$$

where the upper indices 2, 3 represent the doublet and triplet contributions and $a_i^{2,3}$ denote eigenvalues of the messenger mass matrix of doublet and triplet messenger fields $\mathcal{M}_{2,3}$. For simplicity, here we assume $\Lambda^2_g = \Lambda^3_g$ and denote them as $\Lambda_g (= \Lambda^2_g = \Lambda^3_g)$ to parametrize the gaugino masses. It is straightforward to extend the following numerical analysis to the case with $\Lambda^2_g \neq \Lambda^3_g$. We can set the overall scale by the scale $F_\phi$ and express all the soft mass parameters as functions of dimensionless parameters $r_1 \equiv \Lambda_g/\Lambda_X^2$, $r_2 \equiv \Lambda_X^2/F_\phi$ and $r_3 \equiv \Lambda_X^3/F_\phi$.

Before numerical analysis, we comment on sum rules of the soft scalar masses. For each of three generations, the sfermion masses at the messenger scale satisfy the following sum rules,

$$\text{Tr } (B-L) m^2 = 2m_Q^2 - m_U^2 - m_D^2 - 2m_L^2 + m_E^2 = 0. \quad (3.2)$$

In addition, the following sum rule:

$$\text{Tr } Y m^2 = m_Q^2 - 2m_U^2 + m_D^2 - m_L^2 + m_E^2 = 0,$$

is also satisfied for each of the first and second generations. These sum rules are the same as those derived in general gauge mediation [19]. The latter sum rule (3.3) is violated in the third generation because of effects from the top Yukawa coupling. Thus, our parameter space is different from one of general gauge mediation. Instead of the above sum rule, the sfermion masses in the third generation satisfy the following sum rule,

$$m_Q^2 - 2m_U^2 + m_D^2 - m_L^2 + m_E^2 + m_{H_2}^2 - m_{H_1}^2 = 0. \quad (3.4)$$

4 These sum rules have also been derived in the context of various SUSY breaking models [20].
These sum rules have corrections due to renormalization group (RG) effects between the messenger scale and the weak scale \[19, 21\]. However, our natural messenger scale is low such as \(O(10)\) TeV. Thus, such RG corrections on the sum rules are small.

Also, we give a comment on the gauge coupling unification. The doublet/triplet splitting leads to corrections such as \(\log(M^2/M^3)\) in the gauge coupling unification. We assume that \(c_P^2/c_P^3 = O(1)\) and \(\lambda^2_P/\lambda^3_P = O(1)\). That would lead to \(M^2/M^3 = O(1)\). Then, we assume that the doublet/triplet splitting would have a sufficiently small effect on the gauge coupling unification with leading to \(O(1)\) of splitting for \(r_2/r_3\).

In numerical analysis, we take the messenger scale as 10 TeV. The soft SUSY breaking masses are evaluated using the expressions of (2.10) and (2.15), and evolved down to the weak scale with RG effects. Here we take \(\tan\beta = 10\).

One of the important constraints is the condition to avoid the tachyonic slepton masses, which always appear in the pure anomaly mediation. The slepton masses squared become positive for the following parameter region,

\[
1.5r_2^2 + r_3^2 \gtrsim 19,
\]

at the messenger scale. There are RG corrections due to the bino mass between the messenger scale and the weak scale, but such corrections are small and the allowed region does not change drastically.

In Figure 1, we show the spectrum at the weak scale with \(r_2 = r_3\). The upper panel corresponds to the case with \(r_1 = 0\), where the gaugino masses are induced only by the pure anomaly mediation. The lower panel corresponds to the case with \(r_1 = 1\), that is, both anomaly mediation and gauge mediation contribute to the gaugino masses. The red, green and blue solid lines represent the bino, wino and gluino, respectively. The pink dot, yellow dash, violet dashdot, brown longdash, gold spacedash and black spacedot lines represent the soft masses of the left-handed stop \(\tilde{Q}_3\), the right-handed stop \(\tilde{U}_3\), the right-handed down-sector squarks \(\tilde{D}\), the left-handed sleptons \(\tilde{L}\), the right-handed sleptons \(\tilde{E}\) and the up-sector Higgs \(H_2\), where the soft mass of the down-sector Higgs \(H_1\) is the same as one of the left-handed sleptons \(\tilde{L}\). In Figure 1, we have omitted masses of the first two generations of the left-handed squarks \(\tilde{Q}_{1,2}\) and the right-handed up-sector squarks \(\tilde{U}_{1,2}\). Those are heavier than \(\tilde{Q}_3\) and \(\tilde{U}_3\). In Figures 2, 3 and 4, we take \(r_2 = 2r_3, 5r_3, 1/2r_3\), the others are the same as Figure 1.

At the upper panels of Figures, 1, 2, 3 and 4, the gaugino masses are obtained as the pure anomaly mediation, because of \(r_1 = 0\). That is, the wino is the lightest among the gaugino fields, and the gluino is much heavier. Thus, the LSP is the wino-like neutralino in the allowed region, where the slepton masses squared are positive \(3.5\). The next-to-LSP
is the chargino and the mass difference between the LSP and the next-to-LSP is small like anomaly mediation. Squarks are much heavier than other sparticles in the allowed regions. Obviously, a large value of $r_2$ like $r_2 = r_3/2$ in the upper panel of Figure 4 makes squarks heavier and a small value like $r_2 = 2r_3$ and $r_2 = 5r_3$ in the upper panels of Figures 2 and 3 make them lighter. The soft scalar mass of $H_2$ behaves non-trivially depending on $r_2$ and $r_3$. For $r_2 = r_3$ and $r_2 = r_3/2$ in the upper panels of Figures 1 and 4, the soft scalar mass squared $m^2_{H_2}$ decreases as $|r_2|$ increases. For $r_2 = 2r_3$ in the upper panel of Figures 2, the soft scalar mass $m_{H_2}$ is almost constant against $r_2$. On the other hand, for $r_2 = 5r_3$ in the upper panel of Figures 3, the soft scalar mass squared $m^2_{H_2}$ increases as $|r_2|$ increases. The soft scalar mass squared $m^2_{H_2}$ becomes positive for $|r_2| \gtrsim 7.4$, while it is always negative in the other figures. This non-trivial behavior is originated from the negative radiative corrections due to the stop masses between the messenger scale and the weak scale. The stop masses are quite heavy in the upper panels of Figures 1 and 4. In particular, they become much heavier as $|r_2|$ increases. They lead to largely negative radiative corrections in $m^2_{H_2}$. In the case with $r_2 = 2r_3$ and $r_2 = 5r_3$, the stop masses are not heavy compared with the above cases. In addition, the gauge mediation effect on $m^2_{H_2}$ is positive and it increases as $|r_2|$ increases. This leads to the behaviors shown in the upper panels of Figures 2 and 3. For $r_2 \gtrsim 2r_3$, there appears the parameter region of $r_2$, where $m^2_{H_2}$ is positive at the weak scale. In such a region, the successful electroweak symmetry breaking does not occur. Then, we have the excluded region for a large $|r_2|$.

In the lower panels of Figures 1, 2, 3 and 4, it seems that several masses vary variously depending on a value of $r_2$ as well as $r_3$. Such a behavior is originated from the fact that for $r_1 \neq 0$ the gaugino masses have contributions due to both anomaly mediation and gauge mediation, and they vary depend on $r_1r_2$ in our parametrization. Varying the gaugino masses also affect behaviors of the scalar masses through radiative corrections between the messenger scale and the weak scale. However, such radiative corrections on the slepton masses are very small and the behaviors of the sleptons in the lower panels of Figures 1, 2, 3 and 4 are almost the same as the corresponding upper panels, although the squark masses and the up-sector Higgs soft mass change significantly. That is, the region for the non-tachyonic slepton masses corresponds to Eq. (3.5). The three gaugino masses, $M_1, M_2$ and $M_3$, are proportional to $|r_1r_2 - 33/5|$, $|r_1r_2 - 1|$ and $|r_1r_2 + 3|$, respectively, up to the gauge couplings. As a result, the gaugino masses, $M_1$, $M_2$ and $M_3$, are very suppressed around $r_1r_2 \approx 33/5$, $r_1r_2 \approx 1$ and $r_1r_2 \approx -3$, respectively. On the other hand, far away from those points, the gaugino masses, in particular the gluino mass, become heavier. For example, for $r_1 = 1$, those points correspond to $r_2 = 33/5$, 1 and $-3$, respectively. Because of this behavior, the wino is heavier.
than the right-handed slepton in most of the parameter space, and the wino can not be the LSP. When we take a large value of $|r_1|$ like $|r_1| \gtrsim 2$, we would have the parameter region with the wino LSP. For $r_2 > 0$, the bino is the LSP except the parameter region, where the slepton has a tachyonic mass or the right-handed slepton is the LSP. On the other hand, for $r_2 < 0$, the gluino can be the LSP in a narrow region, but in the other region the right-handed slepton is the LSP. Such a parameter region would be unfavorable. In addition, far away from the point $r_2 = -3$, the gluino becomes heavier. One can see this behavior by comparing the upper and lower panels in Figures 1, 2, 3 and 4.

We comment on the squark masses in the lower panels of Figures 1, 2, 3 and 4. The radiative corrections due to the gluino mass between the messenger scale and the weak scale are important in the squark masses. Since the gluino mass is smaller for $r_2 < 0$ than one for $r_2 > 0$, the squark masses are also smaller for $r_2 < 0$ than those for $r_2 > 0$. Note that for $r_1 = 0$ the squark masses as well as the slepton masses are symmetric under the $Z_2$ reflection $r_2 \leftrightarrow -r_2$. Furthermore, a small value of $r_3$ like $r_2 = 2r_3$ and $r_2 = 5r_3$ also leads to smaller squark masses as the upper panels of Figures 2 and 3. In particular, the right-handed stop can have a tachyonic mass in a certain parameter region, as pointed out already in Ref. [12]. Such a parameter region with the tachyonic right-handed stop becomes wider as $r_3$ becomes smaller like $r_2 = 2r_3$ and $r_2 = 5r_3$ in the lower panels of Figures 2 and 3. Thus, in these cases, there is a wide region excluded by tachyonic masses of the right-handed slepton and the right-handed stop. On the other hand, since far away from the point $r_2 = -3$, the gluino mass becomes heavier in particular for positive $r_2$, the stop masses also become heavier in these parameter region. One can see this behavior by comparing the stop and gluino masses in the upper and lower panels of Figures 1, 2, 3 and 4.

Here, we comment on the up-sector Higgs soft mass in the lower panels of Figures 1, 2, 3 and 4. The behavior of the up-sector Higgs soft mass depends on the stop masses. When the stop masses of the lower panels of Figures 1, 2, 3 and 4 are similar to those in the corresponding upper panels, values of $m_{H_u}^2$ are similar between the upper and lower panels. When the stop masses of the lower panels become heavier than those in the upper panels, a value of $m_{H_u}^2$ is driven to a negative direction in the lower panels compared with those in the upper panels. For a large value of $|r_2|$ like $r_2 = 5r_3$, there is a parameter space with $m_{H_u}^2 > 0$ in the lower panel of Figure 3, similar to the corresponding upper panel. Thus, the parameter region is constrained by realization of the successful electroweak symmetry breaking as well as avoiding the tachyonic slepton and/or stop masses. Indeed, the allowed region corresponds to $4 \lesssim r_2 \lesssim 15$ in the lower panel of Figure 3. To summarize phenomenological aspects shown in the lower panels of Figures 1, 2, 3 and 4, certain parameter regions are excluded by the
masses of the right-handed slepton and/or the right-handed stop. In the half of the allowed region, i.e. $r_2 > 0$, the bino is the LSP, while in the other half $r_2 < 0$ the right-handed slepton would be the LSP except a narrow region with the gluino LSP.

When we simply compare between the cases with $r_1 = 0$ and $r_1 = 1$, the allowed region for $r_1 = 0$ is wider than one for $r_1 = 1$. If we take different values of $r_1$, the situation would change.

We give several representative points in Tables 1 and 2. We take $F_\phi$, i.e. the gravitino mass to satisfy the wino mass bound $m_{\tilde{W}} > 100\text{GeV}$ and the bino mass bound $m_{\tilde{B}} > 50\text{GeV}$. The $A$-term is induced by the pure anomaly mediation. We also estimate the $A$-term corresponding to the top Yukawa coupling. It is found that $A_t$ is smaller than the stop mass in most parameter space, so the stop mixing is relatively-small. At all of the points shown in Tables 1 and 2, the stop masses are heavy, so that the lightest Higgs mass is heavy enough to satisfy the LEP bound $m_h > 114\text{GeV}$.

In the model, which we discussed in section 2, the natural messenger scale would be of $\mathcal{O}(10)$-$\mathcal{O}(100)\text{ TeV}$. However, other types of models would lead to a mixture between anomaly mediation and gauge mediation with different messenger scales [3]. Following such a rather phenomenological viewpoint, finally we show examples with higher messenger scales. Figure 5 shows the soft masses in the case with $r_1 = 0$ and $r_2 = r_3$, which are the same as the upper panel in Figure 1. The upper and lower panels in Figure 5 correspond to $10^{10}\text{ GeV}$ and $2 \times 10^{16}\text{ GeV}$ as the messenger scale. It seems that the qualitative results except the up-sector Higgs soft mass are roughly similar to the case that the messenger scale is 10 TeV. One of the important differences is that the spectrum of the soft scalar masses are rather compact in the cases with higher messenger scales compared with the case of the 10 TeV messenger scale, that is, the squarks are lighter in particular for a large value of $|r_2|$. That also affects on the soft scalar mass of the up-sector Higgs. Because of the lighter stop masses, the soft mass squared of the up-sector Higgs becomes positive for a large value of $|r_2|$. This situation is similar to one in the upper panel of Figure 3. Another important point is the long logarithmic RG running. In particular, the slepton masses receive such a long logarithmic RG running effect due to the bino mass, and they tend to become positive. Thus, the region excluded by the tachyonic slepton becomes narrow in Figure 5 compared with one in the upper panel of Figure 1. Similarly, we can study other values of $r_1$ and $r_3$ for higher messenger scales for a purely phenomenological purpose.

We have taken $\Lambda_2^2 = \Lambda_3^3$ just for simplicity in all of the above analyses. It would be interesting to study for other values of the ratio, $\Lambda_2^2/\Lambda_3^3 \neq 1$. 

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3.2 The $\mu - B\mu$ problem

Here, we comment on the $\mu$ term and the $B\mu$ term. In our model, it would be simple to generate the $\mu$ term and the $B\mu$ term by the following terms,

$$\int d^4\theta c_H \frac{\delta^i}{\phi} H_1 H_2 + \int d^2\theta \lambda_H S H_1 H_2 + \text{h.c.}$$

(3.6)

Then, we obtain

$$\mu = \langle F_\phi \rangle c_H + \langle S \rangle \lambda_H,$$

(3.7)

$$B\mu = -\langle F_\phi \rangle \mu + (\langle F_\phi \rangle \langle S \rangle + \langle F_S \rangle) \lambda_H.$$

These are independent of each other because there are two parameters, $c_H$ and $\lambda_H$. However, the natural scale of $B$ would be of $O(F_\phi)$, and such a scale is too large to realize successfully the electroweak symmetry breaking. For example, when $|m_{H_1}^2| \sim |m_{H_2}^2|$, it would be required that $\mu^2$ and $B\mu$ are of the same order. Such a value of $\mu$ can be obtained for $c_H, \lambda_H = O(0.1 - 0.01)$ since we assume that $\langle F_\phi \rangle, \langle S \rangle = O(10)$ TeV. However, since the natural scale of $B$ would be of $O(F_\phi)$, we would need a few percent of fine-tuning between $c_H$ and $\lambda_H$ to realize

$$\mu B \sim \mu^2, |m_{H_1}^2|, |m_{H_2}^2|,$$

(3.8)

that is, the $\mu - B\mu$ problem. Each of anomaly mediation and gauge mediation has the $\mu - B\mu$ problem. Obviously, their mixture studied here also has the same problem unless we have any definite mechanism to cancel out the contributions due to anomaly mediation and gauge mediation in the $B\mu$ term to lead to a suppressed value of the $B\mu$ term.

On the other hand, in Figure 3, there is a parameter region, where $|m_{H_2}^2| \ll |m_{H_1}^2|$, e.g. Point 4 in Table 1. When the relation

$$\mu^2 \sim m_{H_2}^2 \ll B\mu \ll m_{H_1}^2,$$

(3.9)

is satisfied, we can realize the successful electroweak symmetry breaking, that is, a large value of $B$ may not be problematic. For example, at Point 4, we have

$$m_{H_1}^2(M_Z) \simeq 1.61 \text{ TeV}^2, \quad m_{H_2}^2(M_Z) \simeq -7870 \text{ GeV}^2.$$

(3.10)

By using

$$|\mu|^2 = \frac{M_Z^2}{2} - \frac{m_{H_2}^2 \tan^2 \beta - m_{H_1}^2}{\tan^2 \beta - 1}, \quad \sin 2\beta = \frac{2B\mu}{2|\mu|^2 + m_{H_1}^2 + m_{H_2}^2},$$

(3.11)

\footnote{For example, for the mixture of anomaly mediation and moduli mediation, there is a certain type of cancellation mechanisms in the $B\mu$ term [22]. We need such a cancellation mechanism in this scenario.}
we obtain
\[ |\mu| \simeq 120 \text{GeV}, \quad B \simeq 1.4 \text{TeV}, \quad (3.12) \]
e.g. for \( \tan \beta = 10 \). In this case, the fine-tuning to be required between \( c_H \) and \( \lambda_H \) is ameliorated such as \( \mathcal{O}(10) \% \). However, we need another fine-tuning for \( r_2 \) to realize \( |m_{H_2}^2| \ll |m_{H_1}^2| \), unless we have any definite mechanism to set a proper value of \( r_2 \). Thus, the simple way to generate the \( \mu \) term and the \( B\mu \) term requires a fine-tuning. We could consider another way to generate the \( \mu \) term and the \( B\mu \) term \([12, 13]\), where we do not need fine-tuning, but here we do not pursue further. Finally, we comment on the LSP. One of specific aspects at Point 4 is that the LSP is higgsino-like, while different points lead to another LSP such as the bino or wino.

|       | Point 1 | Point 2 | Point 3 | Point 4 | Point 5 |
|-------|---------|---------|---------|---------|---------|
| \( r_1 \) | 0       | 0       | 0       | 0       | 0       |
| \( r_2 \) | 3.5     | 4.0     | 4.0     | 7.37    | 2.5     |
| \( r_3 \) | \( r_2 \) | \( r_2/2 \) | \( r_2/5 \) | \( r_2/5 \) | 2\( r_2 \) |
| \( m_{3/2} \) | 50 TeV  | 50 TeV  | 50 TeV  | 50 TeV  | 50 TeV  |
| \( m_\tilde{B} \) | 462     | 462     | 462     | 462     | 462     |
| \( m_\tilde{t} \) | 135     | 135     | 135     | 135     | 135     |
| \( m_\tilde{\chi}_1^0 \) | 1430    | 1430    | 1430    | 1430    | 1430    |
| \( m_{\tilde{Q}_3} \) | 2450    | 1800    | 1450    | 1910    | 3200    |
| \( m_{\tilde{u}_3} \) | 1970    | 1130    | 465     | 753     | 2810    |
| \( m_{\tilde{Q}_{1,2}} \) | 2590    | 1940    | 1580    | 2040    | 3370    |
| \( m_{\tilde{t}_{1,2}} \) | 2302    | 1480    | 970     | 1260    | 3180    |
| \( m_{\tilde{D}} \) | 2530    | 1820    | 1430    | 1624    | 3340    |
| \( m_{\tilde{L}}(m_{H_1}) \) | 586     | 163     | 671     | 1268    | 412     |
| \( m_{\tilde{E}} \) | 209     | 192     | 166     | 431     | 232     |
| \( -m_{H_2} \) | 1375    | 1020    | 856     | 47.5    | 1790    |
| \( A_t \) | 1300    | 1300    | 1300    | 1300    | 1300    |

Table 1: We give various representative points at the weak scale in the parameter space with \( r_1 = 0 \). For the gaugino masses we take \(|M|\), and for the scalar masses we take \(|m^2|^{1/2} \times \text{sign}(m^2)\). The messenger scale is taken to be 10 TeV. All masses except \( m_{3/2} \) are in GeV.

4 Conclusion

We have studied the models that anomaly mediation and gauge mediation are competed. This mixture can avoid the tachyonic sleptons in a certain parameter space. Our messenger structure, (1.1) and (1.2), leads to a quite rich pattern of the SUSY breaking terms. Still,
Table 2: Same as Table 1, but for $r_1 = 1$.

|         | Point 6 | Point 7 | Point 8 | Point 9 | Point 10 |
|---------|---------|---------|---------|---------|---------|
| $r_1$   | 1       | 1       | 1       | 1       | 1       |
| $r_2$   | 4.0     | 4.5     | 4.5     | 2.5     | 3.5     |
| $r_3$   | $r_2$   | $r_2/2$ | $r_2/5$ | $2r_2$  | $2r_2$  |
| $m_{3/2}$ | 15 TeV  | 20 TeV  | 20 TeV  | 25 TeV  | 15 TeV  |
| $m_{B}$  | 54.6    | 58.8    | 58.8    | 144     | 65.1    |
| $m_{\tilde{W}}$ | 122     | 189     | 189     | 101     | 101     |
| $m_{\tilde{G}}$ | 998     | 1430    | 1425    | 1306    | 926     |
| $m_{\tilde{Q}_3}$ | 961     | 1070    | 954     | 1713    | 1376    |
| $m_{\tilde{U}_{1,2}}$ | 835     | 875     | 743     | 1527    | 1260    |
| $m_{\tilde{D}_{1,2}}$ | 1010    | 870     | 1000    | 1800    | 1440    |
| $m_{\tilde{E}}$ | 936     | 994     | 856     | 1710    | 1400    |
| $m_{\tilde{L}}(m_{H_1})$ | 987     | 1080    | 954     | 1788    | 1430    |
| $m_{\tilde{E}}$ | 207     | 310     | 671     | 207     | 185     |
| $m_{\tilde{H}_2}$ | 76.1    | 91.0    | 80.2    | 112     | 112     |
| $A_t$   | 487     | 507     | 441     | 932     | 716     |
|         | 590     | 819     | 819     | 858     | 564     |

there are the parameter regions excluded by the tachyonic sleptons, the tachyonic stops or the positive soft scalar mass of the up-sector Higgs field. It seems that the allowed parameter space for $r_1 = 0$ is wider than one for $r_1 = 1$. Thus, the models, where the gaugino masses are generated by the pure anomaly mediation, would be interesting. Obviously, our models naturally solve the SUSY flavor problem, because of the mixture between gauge mediation and anomaly mediation. The LSP can be bino-like, wino-like or higgsino-like depending on the parameters, while the LSP might be the stau or the stop in a certain region. Thus, we have a dark matter candidate as usual and it would be interesting to study dark matter physics in our models. The gravitino is heavy such as $\langle F_\phi \rangle \sim \mathcal{O}(10)$ TeV.

We have studied the model with the $5 + \overline{5}$ messenger fields. We can extend our models by adding $10 + \overline{10}$ messenger fields. When we split 10 into $(3, 2) + (3, 1) + (1, 1)$ [26], our model would cover the parameter space corresponding to the general gauge mediation [19] with anomaly mediation. Those models can lead to much richer structure of the SUSY breaking terms.

We have derived the sum rules among the soft scalar masses. The sum rules for the first and the second generations are the same as those in general gauge mediation, but the third generation leads to the different sum rule. Thus, our parameter space is different from one in general gauge mediation. It is important to study theoretical implications and
phenomenological aspects of our sum rules.

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**A: Summary of the soft parameters**

Here, we show the explicit formula of the soft masses in our scenario. The group-theoretical factors $C_2$ and $a$ are summarized in Table 3. The other coefficients are obtained as

$$b_1 = -\frac{33}{5}, \quad b_2 = -1, \quad b_3 = 3,$$

$$f^1 = \frac{13}{15}, \quad f^2 = 3, \quad f^3 = \frac{16}{3}, \quad e = 6.$$  \hspace{1cm} (A.1)

In addition, we define

$$\tilde{\alpha}_i = \left(\frac{g_i}{4\pi}\right)^2, \quad Y_i = \left(\frac{y_i}{4\pi}\right)^2.$$  \hspace{1cm} (A.2)

At the messenger scale, the explicit formula of the gaugino masses and the soft scalar
masses are obtained as follows. The gaugino masses are written by

\[ M_1 = \tilde{\alpha}_1 \left[ -\frac{33}{5} F_\phi + \left( \frac{3}{5} \Lambda_2^2 + \frac{2}{5} \Lambda_3^2 \right) \right], \]

\[ M_2 = \tilde{\alpha}_2 \left[ -F_\phi + \Lambda_2^2 \right], \]

\[ M_3 = \tilde{\alpha}_3 \left[ 3F_\phi + \Lambda_3^2 \right]. \tag{A.3} \]

The stop masses are obtained as

\[ m_{\tilde{t}_3}^2 = \left[ -\frac{11}{50} \tilde{\alpha}_1^2 - \frac{3}{2} \tilde{\alpha}_2^2 + 8\tilde{\alpha}_3^2 + Y_t^2 \left( 6Y_t^2 - \frac{13}{15} \tilde{\alpha}_1 - 3\tilde{\alpha}_2 - \frac{16}{3} \tilde{\alpha}_3 \right) \right] |F_\phi|^2 \]

\[ + \left[ \frac{1}{30} \tilde{\alpha}_1^2 \left( \frac{2}{5} (\Lambda_3^3)^2 + \frac{3}{5} (\Lambda_3^3)^2 \right) + \frac{3}{2} \tilde{\alpha}_2^2 (\Lambda_3^2)^2 + \frac{8}{3} \tilde{\alpha}_3^2 (\Lambda_3^3)^2 \right], \tag{A.4} \]

\[ m_{\tilde{t}_3}^2 = \left[ -\frac{88}{25} \tilde{\alpha}_1^2 + 8\tilde{\alpha}_3^2 + 2Y_t^2 \left( 6Y_t^2 - \frac{13}{15} \tilde{\alpha}_1 - 3\tilde{\alpha}_2 - \frac{16}{3} \tilde{\alpha}_3 \right) \right] |F_\phi|^2 \]

\[ + \left[ \frac{8}{15} \tilde{\alpha}_1^2 \left( \frac{2}{5} (\Lambda_3^3)^2 + \frac{3}{5} (\Lambda_3^3)^2 \right) + \frac{8}{3} \tilde{\alpha}_3^2 (\Lambda_3^3)^2 \right]. \]

The masses for the first and second generations of the up-sector left-handed and right-handed squarks are obtained in the same form except taking \( Y_t = 0 \). The right-handed down-sector squark masses are obtained as

\[ m_{\tilde{D}}^2 = \left[ -\frac{22}{25} \tilde{\alpha}_1^2 + 8\tilde{\alpha}_3^2 \right] |F_\phi|^2 + \left[ \frac{2}{15} \tilde{\alpha}_1^2 \left( \frac{2}{5} (\Lambda_3^3)^2 + \frac{3}{5} (\Lambda_3^3)^2 \right) + \frac{8}{3} \tilde{\alpha}_3^2 (\Lambda_3^3)^2 \right]. \tag{A.5} \]

The slepton masses are obtained as

\[ m_{\tilde{L}}^2 = \left[ -\frac{99}{50} \tilde{\alpha}_1^2 - \frac{3}{2} \tilde{\alpha}_2^2 \right] |F_\phi|^2 + \left[ \frac{3}{10} \tilde{\alpha}_1^2 \left( \frac{2}{5} (\Lambda_3^3)^2 + \frac{3}{5} (\Lambda_3^3)^2 \right) + \frac{3}{2} \tilde{\alpha}_2^2 (\Lambda_3^2)^2 \right], \]

\[ m_{\tilde{E}}^2 = \left[ -\frac{198}{25} \tilde{\alpha}_1^2 \right] |F_\phi|^2 + \left[ \frac{6}{5} \tilde{\alpha}_1^2 \left( \frac{2}{5} (\Lambda_3^3)^2 + \frac{3}{5} (\Lambda_3^3)^2 \right) \right]. \tag{A.6} \]

The Higgs soft masses are obtained as

\[ m_{H_1}^2 = m_{\tilde{L}}^2 \]

\[ m_{H_2}^2 = \left[ -\frac{99}{50} \tilde{\alpha}_1^2 - \frac{3}{2} \tilde{\alpha}_2^2 + 3Y_t^2 \left( 6Y_t^2 - \frac{13}{15} \tilde{\alpha}_1 - 3\tilde{\alpha}_2 - \frac{16}{3} \tilde{\alpha}_3 \right) \right] |F_\phi|^2 \]

\[ + \left[ \frac{3}{10} \tilde{\alpha}_1^2 \left( \frac{2}{5} (\Lambda_3^3)^2 + \frac{3}{5} (\Lambda_3^3)^2 \right) + \frac{3}{2} \tilde{\alpha}_2^2 (\Lambda_3^2)^2 \right]. \tag{A.7} \]
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Figure 1: Spectrum of the superpartner masses as a function of $r_2$. We take $r_2 = r_3$ and the messenger scale as 10 TeV. For the gaugino masses we plot $|M|$, and for the scalar masses we plot $|m_2|^{1/2} \times \text{sign}(m_2)$ in unit of $\Lambda_\phi$. 
Figure 2: Same as Figure 1, but for $r_2 = 2r_3$. 
Figure 3: Same as Figure 1, but for \( r_2 = 5r_3 \).
Figure 4: Same as Figure 1, but for $r_2 = 1/2r_3$. 
Figure 5: Same as Figure 1, but we take the messenger scale as $10^{10} \text{ GeV}$ at the upper panel and $2 \times 10^{16} \text{ GeV}$ at the lower panel.