A Closer Look at Gaugino Masses
in
Pure Gravity Mediation Model/Minimal Split SUSY Model

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Abstract
We take a closer look at the gaugino masses in the context of pure gravity mediation models/minimal split SUSY models. We see that the gaugino mass spectrum has a richer structure in the presence of vector-like matter fields even when they couple to the supersymmetry breaking sector only through Planck suppressed operators. For example, the gluino mass can be much lighter than in anomaly mediation, enhancing the detectability of the gluino at the LHC experiments. The rich gaugino spectrum also allows new possibilities for dark matter scenarios such as the bino-wino co-annihilation, bino-gluino co-annihilation, or even wino-gluino co-annihilation scenarios, which affects future collider experiments as well as dark matter search experiments.
1 Introduction

For decades, supersymmetry (SUSY) has been intensively studied as a new fundamental ingredient of nature since it allows for a vast separation of low energy scales from high energy scales such as the Planck scale or the scale of the Grand Unified Theory (GUT) [1–4]. The unification of the gauge coupling constants at a very high energy scale also strongly motivates the supersymmetric standard model (SSM). Weak-scale supersymmetry with the superparticle masses in the hundreds GeV to a TeV range, in particular, has been placed as top prospect as it naturally solves the hierarchy problem while providing a good candidate for WIMP dark matter.

On the other hand, the supersymmetric standard models with heavier sfermions in the hundreds to thousands TeV range and the gauginos in the hundreds GeV to a TeV range have also gained attention as attractive alternatives to the weak-scale supersymmetry [5,6]. There, the neutral gauginos can be a good candidate of WIMP dark matter. One of the biggest advantages of this class of models is that they do not require any singlet SUSY breaking fields (i.e. the Polonyi fields), and hence, are free from the cosmological Polonyi problem [7,8]. These models are also free of the gravitino problem even for a very high reheating temperature [9,10], which is crucial for the successful thermal leptogenesis [11]. The suppressions of flavor-changing neutral currents and CP violating processes are another important advantages of this class of models.

One apparent drawback of the heavy sfermion models is that the fine-tuning required for the weak-scale is one part of $10^6–10^8$. In view of the above advantages, however, this drawback might be tolerable. Alternatively, the anthropic principle [12], which could be a manifestation of a large number of meta-stable vacua in string theory [13–16], may explain the weak-scale to Planck scale hierarchy [17,18], which further supports the heavy sfermion models discussed in the context of split supersymmetry [19,21].

We note in passing that the origin of the $\mu$-term has been a problem from the early stages of model building due to the absence of a singlet SUSY breaking fields [5,6]. This problem has been solved with a very simple model in Ref. [22] by utilizing a mechanism that generates a $\mu$-term similar in size to the gravitino mass from Planck suppressed
interactions to the $R$-symmetry breaking sector. The model in Ref. is now dubbed pure gravity mediation. It should be noted that pure gravity mediation has an identical structure to minimal split SUSY models. In this paper, we use the name “pure gravity mediation model”.

In pure gravity mediation models, the scalar bosons obtain SUSY breaking masses from the SUSY breaking sector via tree-level interactions of supergravity. The $\mu$-term is, as mentioned above, generated from the $R$-symmetry breaking sector via tree-level interactions of supergravity. The gaugino masses are, on the other hand, suppressed at the tree-level due to the absence of a singlet SUSY breaking field and are generated at the one-loop level mainly from the anomaly mediated SUSY breaking (AMSB) contributions. The observed 126GeV Higgs boson mass is explained with the gravitino mass in the hundreds to thousands TeV range since $\tan\beta$, the ratio of the two Higgs vacuum expectation values, is predicted to be $\mathcal{O}(1)$ in this class of model.

To date, the continuing absence of supersymmetric particles at the LHC is putting negative pressure on weak-scale supersymmetry and seems to favor heavier sfermion models. Furthermore, the observed Higgs mass at around 126 GeV points to a sfermion mass scales above the tens of TeV range as anticipated in Ref. In response to these results, the heavy sfermion models are gathering renewed attention from the view point of collider and dark matter phenomenology, low energy precision measurements, GUT or Planck scale theories, and cosmology.

In this paper, we give a closer look at the gaugino masses in the context of pure gravity mediation models. The gaugino masses are the most important phenomenological parameters in the foreseeable future. In particular, we pay attention to the threshold
corrections from vector-like matter fields. As pointed out in Refs. [57–60], these contributions can be comparable to the AMSB contributions even when the vector-like matter fields couple to the SUSY breaking sector only through Planck suppressed operators. As we will see, the gluino mass can be much lighter than the AMSB predictions for a given gravitino mass [29, 59, 60], enhancing the detectability of the gluino at the LHC experiments. We point out that the rich gaugino spectrum also allows for new possibilities for dark matter other than the standard thermal/non-thermal wino dark matter [61–64], including the bino-wino co-annihilation, the bino-gluino co-annihilation, or even the wino-gluino co-annihilation scenarios, which affects future dark matter search experiments. We emphasize the importance of the phases of the gaugino masses given by the threshold corrections, which are not fully discussed in the literatures [29, 59].

2 Gaugino Masses From Vector-Like Extra Matter

In this section, we calculate the gaugino masses when there exist vector-like matter fields $Q$ and $\bar{Q}$ which are charged under the SSM gauge symmetries. We assume that $R$ charges of $Q\bar{Q}$ add up to 0, so that the extra matter fields obtain a supersymmetric (Dirac) mass of order the gravitino mass from the $R$-breaking sector as is the case for the higgs multiplets in pure gravity mediation models [23, 24]. We also assume that the extra matter fields couple to the SUSY breaking sector only through Planck suppressed interactions similar to SSM matter fields in pure gravity mediation models.

2.1 SUSY breaking mass spectrum of extra matter

To illustrate how the SUSY breaking effects show up in the mass spectrum of extra matter, let us consider the simplest SUSY breaking sector with the effective superpotential,

$$ W = \Lambda^2 Z + \frac{m_{3/2}}{2} M_{PL}^2. \quad (1) $$

Here, $m_{3/2}$ denotes the gravitino mass representing the spontaneous (discrete) $R$-symmetry breaking, and $M_{PL}$ the reduced Planck scale. The SUSY breaking field $Z$ obtains an $F$-term vacuum expectation value (VEV) of $F_Z = -\Lambda^2$, and the flat universe condition gives
\[ \Lambda^4 = 3m_{3/2}^2M_{PL}^2 \]  
In the followings, we take \( m_{3/2} \) and \( \Lambda \) real and positive without loss of generality. As emphasized in the previous section, there is no singlet SUSY breaking field (i.e. the Polonyi field) in the pure gravity mediation model. The SUSY breaking field \( Z \) is assumed to be charged under some symmetries at the Planck scale or to be a composite field generated at some dynamical scale much lower than the Planck scale.

Due to the vanishing \( R \)-charge of \( Q \bar{Q} \), the extra matter couples to the above SUSY breaking sector via the super- and Kähler potentials;

\[
W = \Lambda^2 Z \left(1 + y \frac{Q \bar{Q}}{M_{PL}^2}\right) + m_{3/2}M_{PL}^2 \left(1 + y' \frac{Q \bar{Q}}{M_{PL}^2}\right),
\]

\[
K = \lambda Q \bar{Q} + \lambda' Z \bar{Z} \frac{Q \bar{Q}}{M_{PL}^2} + \text{h.c.} + \cdots,
\]

where \( y, y', \lambda \) and \( \lambda' \) are dimensionless coupling constants. It should be noted that we can eliminate one of \( y, y' \) and \( \lambda \) through the Kähler-Weyl transformation when we are only interested in the masses of the extra matter fields (see also Ref. [65] for a related discussion). In fact, by using the Kähler-Weyl transformation\(^4\)

\[
K \rightarrow K - \lambda Q \bar{Q} - \lambda' Q \bar{Q} \bar{Q}, \quad W \rightarrow W \exp \left(\frac{\lambda Q \bar{Q}}{M_{PL}^2}\right),
\]

the super- and Kähler potential can be rewritten as,

\[
W' = (y' + \lambda) m_{3/2} Q \bar{Q} + \sqrt{3} (y + \lambda) m_{3/2} Z \frac{Q \bar{Q}}{M_{PL}^2} + \sqrt{3} m_{3/2} M_{PL} Z + \cdots,
\]

\[
K' = \lambda' Z \bar{Z} \frac{Q \bar{Q}}{M_{PL}^2} + \text{h.c.}.
\]

Therefore, we obtain the supersymmetric Dirac mass, \( M \), and the supersymmetry breaking mixing mass parameter, \( b \),

\[
M = (y' + \lambda) m_{3/2},
\]

\[
b = (3y - y' + 2\lambda - 3\lambda')m_{3/2}^2.
\]

In deriving the expression of \( b \), we have added up the contributions from the couplings to the SUSY breaking field and from the constant term in the superpotential through the

\(^3\)It is assumed that \( |\langle Z \rangle| \ll M_{PL} \).

\(^4\)Since the Kähler-Weyl transformation involves chiral rotations of fermion fields in chiral multiplets, it induces gauge kinetic functions which are proportional to \( \lambda Q \bar{Q}/M_{PL}^2 \). However, these terms do not contribute to gaugino masses.
supergravity interactions. As we expected, the parameters \( y, y' \) and \( \lambda \) show up in two combinations, \( (y + \lambda) \) and \( (y' + \lambda) \), although we keep all the parameters for later purpose. As we will show, the phase of \( b/M \) is a very important parameter for the gaugino masses.

### 2.2 Gaugino masses from threshold corrections

In order to calculate the threshold corrections, let us take the mass diagonalized basis for the extra matters. Here, it should be noted that in addition to the above mentioned SUSY breaking \( b \)-term, the scalar components of the extra matter generically obtain soft squared masses of order the gravitino mass just as the SSM matter fields do. Thus, the mass terms of the scalar components \( A, \bar{A} \) and the fermion components, \( \psi, \bar{\psi} \) are given by,

\[
\mathcal{L}_{\text{mass-scalar}} = -(|M|^2 + \tilde{m}^2_A)|A|^2 - (|M|^2 + \tilde{m}^2_{\bar{A}})|\bar{A}|^2 - (bA\bar{A} + \text{h.c.})
\]

\[
\equiv -m^2_A|A|^2 - m^2_{\bar{A}}|\bar{A}|^2 - (bA\bar{A} + \text{h.c.}),
\]

\[
\mathcal{L}_{\text{mass-fermion}} = -M\psi\bar{\psi} + \text{h.c.},
\]

respectively. Here, \( \tilde{m}^2_A \) and \( \tilde{m}^2_{\bar{A}} \) denote the soft squared masses. The mass terms of the scalar components are diagonalized by rotating the fields,

\[
\begin{pmatrix}
A_+ \\
A_-
\end{pmatrix} = \begin{pmatrix}
\cos \beta_Q & -e^{-i(\delta+\delta')}\sin \beta_Q \\
e^{i(\delta+\delta')}\sin \beta_Q & \cos \beta_Q
\end{pmatrix} \begin{pmatrix}
A \\
\bar{A}^\dagger
\end{pmatrix},
\]

\[
\tan \beta_Q = \frac{m^2_A - m^2_{\bar{A}} + \sqrt{(m^2_A - m^2_{\bar{A}})^2 + 4|b|^2}}{2|b|} > 0,
\]

\[
\delta = \arg(b/M), \quad \delta' = \arg(M),
\]

which leads to the mass eigenvalues,

\[
m^2_{\pm} = \frac{1}{2} \left( m^2_A + m^2_{\bar{A}} \pm \sqrt{(m^2_A - m^2_{\bar{A}})^2 + 4|b|^2} \right).
\]

The one-loop threshold correction from the extra matter with the above mass spectrum yields the gaugino masses \[67,\]

\[
\Delta m_{\lambda,\text{threshold}} = \frac{g^2}{16\pi^2} C_Q 2e^{i\theta}\sin 2\beta_Q |M| \left( \frac{m^2_+}{m^2_+ - |M|^2} \ln \frac{m^2_+}{|M|^2} - \frac{m^2_-}{|M|^2 - m^2_-} \ln \frac{|M|^2}{m^2_-} \right),
\]

\footnote{If \( QQ \) couples to some flat directions, there also exist contributions to the \( b \)-term by \( F \)-terms of the flat directions \[66\]. We assume, however, that \( QQ \) do not couple to any flat directions.}
at the renormalization scale just below their threshold. Here, $C_Q$ is a Dynkin index of $Q$, which is normalized to be 1/2 for a fundamental representation, and $g$ the gauge coupling constant evaluated at around the scale of the extra matter. By adding the AMSB effects of the extra matter, $\Delta m_{\lambda,\text{AMSB}} = g^2/(16\pi^2)2C_Qm_3^2$, we obtain the final result\footnote{This formula can be applied to any cases, no matter the origin of the Dirac mass, $b$ term, and soft squared mass terms.}

\[
\Delta m_\lambda = \frac{g^2}{16\pi^2}2C_Q\left(e^{i\delta}\sin2\beta_Q|M|\left(\frac{m_+^2}{m_+^2 - |M|^2}\ln\frac{m_+^2}{|M|^2} - \frac{m_-^2}{|M|^2 - m_-^2}\ln\frac{|M|^2}{m_-^2}\right) + m_3^2\right).
\]

(11)

Before closing this subsection, several comments are in order. First, it can be proven that

\[
\frac{m_+^2}{m_+^2 - |M|^2}\ln\frac{m_+^2}{|M|^2} - \frac{m_-^2}{|M|^2 - m_-^2}\ln\frac{|M|^2}{m_-^2} > 0.
\]

(12)

Therefore, the phase of the gaugino mass contributed from the threshold correction is always determined by the phase of $b/M$.

Secondly, let us take the limit of small soft squared masses, i.e. $\tilde{m}_A^2, \tilde{m}_\tilde{A}^2 \ll |M|^2$. In this limit, the diagonalized scalar masses and mixing angle are reduced to

\[
m_\pm^2 = |M|^2 \pm |b|, \quad \tan\beta_Q = 1.
\]

(13)

With this mass spectrum, Eq. (11) is also reduced to

\[
\Delta m_\lambda = \frac{g^2}{16\pi^2}2C_Q\left[\frac{b}{M}F(|b/M|^2) + m_3^2\right],
\]

\[
F(x) = \frac{1 + x}{x^2}\ln(1 + x) + \frac{1 - x}{x^2}\ln(1 - x).
\]

(14)

In order for the scalar components of $Q\bar{Q}$ not to be tachyonic, the $b$ term should satisfy $|b| < |M|^2$, where the function $F$ takes values between 1 to $\ln(4) \simeq 1.4$.

Thirdly, let us consider the limit of $|y'| \gg 1$. In this case, the spectrum for the extra matter is similar to the case with a large Dirac mass term in the super-potential. Therefore, we expect that $Q\bar{Q}$ decouples and $\Delta m_\lambda = 0$ as expected from the ultraviolet insensitivity properties of the AMSB spectrum. Actually, since the Dirac mass term and
the $b$ term are given by $M = y'm_{3/2}$ and $b/M = -m_{3/2}$, and the soft squared mass terms are negligible, we obtain $\Delta m_\lambda = 0$ from Eq. (14).

Finally, let us take the limit of $|\lambda| \gg 1$, where the Dirac mass term and the $b$ term are given by $M = \lambda m_{3/2}$ and $b/M = 2m_{3/2}$. The soft squared mass terms are negligible and we obtain

$$\Delta m_\lambda = \frac{g^2}{16\pi^2}6C_Qm_{3/2}. \tag{15}$$

The AMSB effect and the threshold correction contribute to gaugino masses additively.

### 2.3 The SSM gaugino masses

The AMSB contributions from the minimal supersymmetric standard model (MSSM) fields to the gluino mass $M_3$, the wino mass $M_2$, and the bino mass $M_1$ are given by

$$M_3^{\text{AMSB}} = -\frac{g^2_3}{16\pi^2}3m_{3/2}, \quad M_2^{\text{AMSB}} = \frac{g^2_2}{16\pi^2}m_{3/2}, \quad M_1^{\text{AMSB}} = \frac{g^2_1}{16\pi^2} \frac{33}{5}m_{3/2}, \tag{16}$$

at around the renormalization scale $m_{3/2}$. Besides the contribution, the wino and the bino obtain threshold corrections from the higgsino as in Ref. [5],

$$\Delta M_2 = \frac{g^2_2}{16\pi^2}L, \quad \Delta M_1 = \frac{g^2_1}{16\pi^2} \frac{3}{5}L,$$

$L \equiv \mu \sin^2\beta \frac{m_A^2}{|\mu|^2 - m_A^2} \ln \frac{|\mu|^2}{m_A^2}$, \tag{17}

where $\beta$ is defined by the ratio of the vacuum expectation value of the up type higgs to that of the down type higgs, $\tan\beta = v_u/v_d$, and $m_A^2$ is the mass squared of heavy higgs bosons. As emphasized in Ref. [27], the higgsino threshold corrections can be comparable to AMSB contributions in large sections of parameter space in pure gravity mediation models since it predicts $\mu \sim m_A = \mathcal{O}(m_{3/2})$ and $\tan\beta = \mathcal{O}(1)$.

As we have found, the extra-matter fields give additional contributions to the gaugino masses as in Eq. (11). The physical gaugino masses are evaluated by solving renormalization equations,

$$\frac{d\ln M_i(\mu)}{d\ln \mu} = -\frac{g^2_i(\mu)}{8\pi^2} b_i, \quad (b_1, b_2, b_3) = (0, 6, 9),$$

$$M_i(M_{i,\text{phys}}) = M_{i,\text{phys}}, \tag{18}$$

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Boundary conditions are given by the sum of the contributions given in Eqs. (16), (17) and (11) evaluated at around the renormalization scale $m_{3/2}$.

In the following, let us suppose that the vector-like matter fields belong to $SU(5)$ GUT multiplets so that coupling unification is preserved. In this case, the contributions from the vector-like matter fields are given by,

$$\Delta M_i = \frac{g_i^2}{16\pi^2} e^{i\gamma} N_{\text{eff}} m_{3/2},$$

and hence, satisfy the so-called GUT relation. The definition of $N_{\text{eff}}$ can be understood by comparing Eqs. (11) and (19). It should be noted that $N_{\text{eff}}$ can be rather large either from small $m^2$ or from many extra matter fields. As we have discussed, the phase of $b/M$ is a free parameter, and hence, we take $\gamma$ as a free parameter. This is a crucial difference from axion models, where there is no phase freedom as reviewed in the appendix A.

Let us comment on $CP$ violations from the phase of the gaugino masses. First, we assume that some flavor symmetry controls the soft squared mass terms so that they are nearly diagonal, since otherwise constraints from the $K^0 - \bar{K}^0$ mixing suggest that the soft squared mass terms are larger than $\mathcal{O}(1000) \text{ TeV}$ [41, 68], even if $\gamma = 0$. Under this assumption, a one-loop contribution to the neutron electric dipole moment is much smaller than the experimental upper bound [48]. A two loop Barr-Zee type contribution, which dominates over the one loop contribution for large soft squared mass terms, is also far smaller than the experimental upper bound for $\mu = \mathcal{O}(100) \text{ TeV}$ [69].

In Figure 1, we show the physical gaugino masses in the presence of the extra matter fields for $m_{3/2} = 100 \text{ TeV}$ as a function of $N_{\text{eff}}$ for given values of $\gamma$. Here, we have neglected the higgsino threshold correction for simplicity, i.e. $L = 0$. It can be seen that the gluino mass can be much lighter than that predicted in pure AMSB. The figures also show that the gaugino mass spectrum strongly depends on the phase $\gamma$. It should be noted that it is even possible for all three gauginos to be degenerate for $\gamma \simeq 0$ and $N_{\text{eff}} \simeq 4 - 5$. This is caused by the fact that the MSSM contributions to the gluino mass is negative while those to the wino and the bino masses are positive. Thus, the addition of the extra matter contributions satisfying the GUT relation can reduce the gluino mass while increasing the wino and bino masses.

$N_{\text{eff}}$ is not identical to the number of flavors, $\sum_Q 2C_Q$. 




For a comparison, in Figure 2 we show the physical gaugino masses when there are no extra vector-like matter fields, i.e. $N_{\text{eff}} = 0$, but the higgsino threshold corrections are included, i.e. $L \neq 0$. Since the gluino mass does not receive any corrections from the higgsino threshold, it is difficult for the gauginos to have a rather degenerated spectrum.

In Figures 3, 4 and 5 we also show some parameter regions which are phenomenologically distinctive including:

- Regions in which the bino/gluino is the lightest supersymmetric particle (LSP)
- Regions in which the thermal abundance of the LSP is larger than the measured abundance of the cold dark matter, $\Omega_c h^2 = 0.1199 \pm 0.0027$ [70]
- Regions which are excluded by charged wino searches [71]
- Regions which are excluded by the gluino searches [72]
- Contour plots of the mass of the wino, $M_2$
- Contour plots of the ratio $M_3/M_{\text{LSP}}$

for various $m_{3/2}$, $L$ and $N_{\text{eff}}$. For the calculation of the thermal abundance of the LSP, we have utilized micrOMEGAs 3.2 [73]. For simplicity, we have neglected the Sommerfeld enhancement of the annihilation cross sections of the wino and gluino [74].

Let us first examine Figure 3 where $m_{3/2} = 50$ TeV. Usually, when the bino is the LSP, the thermal abundance of the LSP exceeds the observed value (the left most and the right most regions). However, there exist bands in which the bino is the LSP but the thermal abundance of the LSP does not exceed the observed value due to the co-annihilations with the wino [76] or gluino [77]. In these regions, it is possible that the bino is the dark matter and is difficult to be observed through direct/indirect detections, because the bino interactions are suppressed.

As the figures show, the ratio between the gluino mass and the LSP mass, $M_3/M_{\text{LSP}}$, strongly depends on the phase $\gamma$. If $M_3/M_{\text{LSP}}$ is close to one, for example, the decay products of gluinos are soft. Thus, in such cases, even if gluinos are copiously produced

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8 If one includes the effect of the Sommerfeld enhancement, the widths of the bino LSP bands discussed later are modified due to the change in the wino or gluino annihilation cross section. For a qualitative discussion, see Ref. [75]. Furthermore, the regions where the thermal abundance of the wino is consistent with the observed dark matter abundance are shifted to higher mass scale regions, $M_2 \simeq 3$ TeV [74].
at the LHC, the search for SUSY events requires initial state radiation \[82,83\]. Therefore, the phase parameter $\gamma$ is not only important for the dark matter properties but also important for collider searches.

In much of the right portion of Figure 5, the wino is the LSP and too heavy and the thermal abundance of the wino is larger than the observed value. However, due to co-annihilation with the gluino, there exist a region in which the thermal abundance is still smaller than the observed value (the lower right panel, a white band between the colored regions). This extends the thermal wino dark matter regions allowing heavier winos.

Let us comment about constraints on the dark matter from indirect dark matter searches. As discussed in Refs. \[78,79\], the wino LSP is constrained by indirect dark matter searches, in particular, by the gamma-ray searches in the Fermi and H.E.S.S. experiments from the Galactic center and the dwarf Spheroidal galaxies. The constraints from the gamma-ray searches from the Galactic center, however, suffer from large ambiguities of a dark matter profile (see e.g. Ref. \[80\]) and background estimations. Thus, by taking into account those ambiguities, the least uncertain constraints are set by the diffused gamma-ray search from the dwarf Spheroidal galaxies by the Fermi-LAT \[81\], which excludes the wino LSP mass below about 400GeV and in 2.2-2.5TeV. The bino LSP is, on the other hand, free from the constraints by the indirect dark matter searches.

In Table 1, we show some phenomenologically interesting benchmark points. Points A and C represent the bino-wino co-annihilation regions. Since the gluino is not degenerate with the LSP, a gluino search with hard jets and large missing energies will be effective at the LHC. Points B and D represent the bino-gluino co-annihilation regions. Since the gluino is degenerate with the LSP, a search utilizing initial state radiation is necessary at the LHC \[82,83\].

| \(m_{3/2} \) (TeV) | \(L/m_{3/2} \) | \(N_{\text{eff}} \) | \(\gamma \) | \(M_3 \) (GeV) | \(M_2 \) (GeV) | \(M_1 \) (GeV) |
|----------------|------------|--------|--------|-------------|-------------|-------------|
| A 50          | 0          | 4.5    | 0.85   | 1515        | 746         | 729         |
| B 50          | 1          | 5      | 0      | 955         | 993         | 877         |
| C 100         | 0          | 4.5    | 0.85   | 2878        | 1482        | 1470        |
| D 100         | 1          | 5      | 0      | 1808        | 1972        | 1768        |

Table 1: Some phenomenologically interesting benchmark points.
3 Discussion and conclusion

In this paper, we have investigated the gaugino masses in pure gravity mediation models. We have especially focused on the threshold correction from an additional vector-like matter, which gives a rich structure to the gaugino masses.

It is possible that the gluino is much lighter than in the anomaly mediated gaugino spectrum, which affects the gluino searches at the LHC. Lighter gluino masses enhances the detectability of the gluino at the LHC experiments. If the gluino is degenerate with the LSP, a search utilizing initial state radiation is necessary \[82,83\].

The rich structure also allows for new dark matter possibilities such as bino-wino co-annihilation, bino-gluino co-annihilation, or even wino-gluino co-annihilation. It is possible that the bino is dark matter and is difficult to be observed through direct/indirect detections, because the bino interactions are suppressed. If the bino decays through an \( R \) parity violating interaction, however, it can provide signals in cosmic rays. For example, a decaying bino yields positron signals consistent with the anomalous results \[84\] reported.
by the PAMELA [85], Fermi-LAT [86] and AMS-02 [87] collaborations.

We have also noted that the corrections to the gaugino masses in general have different phases than the contributions from AMSB i.e. when the corrections are from vector-like matter as heavy as the gravitino. We have pointed out that the phases are important phenomenologically. For example, in the bino-wino co-annihilation region, the gluino masses are strongly dependent on the phases, as can be seen from Figures 3 and 4.

If the mentioned above rich structures are confirmed, especially for a gluino much lighter than in the anomaly mediated gaugino spectrum, it suggests the existences of additional vector-like matter. It is remarkable that one can probe higher energy physics by measuring the gaugino masses. In particular, when the vector-like matter obtains a Dirac mass from $R$ symmetry breaking and hence are as heavy as the gravitino, they may be found in future collider experiments of $\mathcal{O}(100)$ TeV.

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9 In Ref. [84], the argument is made based on the wino dark matter scenario. As for a decay through an $R$ parity violation by $LLE^c$ operator, the phenomenology is essentially the same as in the bino dark matter case.
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A Review on gaugino masses from axion model

In this section, we review the contribution of axion models to the gaugino masses, following Ref. [60]. In general supersymmetric axion models, there are axion chiral multiplets, which couple to vector-like matter fields. Since the axion multiplets are flat directions and hence not fixed, they generally obtain non-zero $F$ terms. Thus, the gaugino masses receive threshold corrections from the vector-like matters.

A.1 KSVZ type models

Let us consider the so-called KSVZ [88,89] type axion model in which the anomaly of the Peccei-Quinn (PQ) symmetry [90–92] of the QCD is mediated by additional SSM charged matters $Q$ and $\bar{Q}$. Here, we assume a super-potential,

$$W = \lambda X(\psi \bar{\psi} - v^2) + y \frac{\psi^n}{M_{PL}^{n-1}} Q\bar{Q},$$

where $X, \psi, \bar{\psi}$ are fields which have $(PQ, R)$ charges $(0, 2), (1, r_\phi)$ and $(-1, -r_\phi)$, respectively. Without loss of generality, we take $\lambda, y$ and $v$ to be positive and real by field redefinitions. We assume that the axion multiplet is the only flat direction, $\lambda v \gg m_{3/2}$. We also assume that $y \langle \psi \rangle^n / M_{PL}^{n-1} \gg m_{3/2}$.

The scalar potential of the scalar components of $X, \psi,$ and $\bar{\psi}$ is given by

$$V = \lambda^2 |\psi \bar{\psi} - v^2|^2 + \lambda^2 |X|^2 (|\psi|^2 + |\bar{\psi}|^2)
+ m_{3/2}^2 (a_X |X|^2 + a_\psi |\psi|^2 + a_{\bar{\psi}} |\bar{\psi}|^2)
+ \left(2\lambda v^2 m_{3/2} X + \tilde{b} m_{3/2}^2 \psi \bar{\psi} + \text{h.c.} \right).$$
Here, we assume that $X$, $\psi$ and $\bar{\psi}$ couple to the SUSY breaking sector only through Planck suppressed interactions, and hence, $a_X$, $a_\psi$ and $a_{\bar{\psi}}$ are at largest $\mathcal{O}(1)$. It should be noted that the $\tilde{b}$ term, $\tilde{b}m_{3/2}^2\psi\bar{\psi}$ with $\tilde{b} = \mathcal{O}(1)$, can arise from the $R$ symmetry breaking effect \cite{23,24} because the combination $\psi\bar{\psi}$ is neutral under the PQ and $R$ symmetry. As we will see, however, the $\tilde{b}$ term does not affect gaugino masses.

The minimum of the potential is given at

$$
\langle X \rangle = -\frac{2m_{3/2}v^2}{\lambda (|\langle \psi \rangle|^2 + |\langle \bar{\psi} \rangle|^2)} \left(1 + \mathcal{O}\left(\frac{m_{3/2}^2}{\lambda^2 v^2}\right)\right),
$$

$$
\langle \psi \rangle = \left(\frac{a_{\bar{\psi}}m_{3/2}^2 + \lambda^2 |\langle X \rangle|^2}{a_{\psi}m_{3/2}^2 + \lambda^2 |\langle X \rangle|^2}\right)^{1/4} v \left(1 + \mathcal{O}\left(\frac{m_{3/2}^2}{\lambda^2 v^2}\right)\right),
$$

$$
\langle \bar{\psi} \rangle = \left(\frac{a_{\psi}m_{3/2}^2 + \lambda^2 |\langle X \rangle|^2}{a_{\bar{\psi}}m_{3/2}^2 + \lambda^2 |\langle X \rangle|^2}\right)^{1/4} v \left(1 + \mathcal{O}\left(\frac{m_{3/2}^2}{\lambda^2 v^2}\right)\right). \tag{22}
$$

Here, we take $\langle \psi \rangle$ to be positive and real by field redefinitions. Note that at the leading order in $m_{3/2}/(\lambda v)$, the vacuum expectation values do not depend on the $\tilde{b}$ term. This is because the direction $\psi\bar{\psi}$ is fixed by the super-potential.

In order to calculate the gaugino masses, let us calculate the $b$ term of $QQ$. It is given by

$$
\mathcal{L}_{b-term} = \frac{y \langle \psi \rangle^n}{M_{PL}^n} m_{3/2} A\bar{A} + n y \frac{\langle \psi \rangle^{n-1}}{M_{PL}^{n-1}} \langle F_\psi \rangle A\bar{A} + h.c., \tag{23}
$$

$$
F_\psi = - \left(W_\psi^\dagger + m_{3/2}\psi\right) = -\lambda X^\dagger \bar{\psi} - m_{3/2}\psi, \tag{24}
$$

where $A$ and $\bar{A}$ are the scalar components of $Q$ and $\bar{Q}$, respectively, as in the main text.

When we calculate the gaugino masses via the $QQ$ loop, the contributions from the first term in Eq. \eqref{23} cancel with the AMSB contribution.\footnote{This cancellation happens only when $y \langle \psi \rangle^n / M_{PL}^{n-1} \gg m_{3/2}$. For gaugino masses with $y \langle \psi \rangle^n / M_{PL}^{n-1} \sim m_{3/2}$, see Sec. \ref{sec:m3/2}.} This is the famous decoupling of heavy vector-like matter \cite{5}. The contributions from the second term, on the other hand, do not cancel, which lead to corrections to the gaugino masses given by

$$
\Delta m_\lambda = -\frac{ny \langle \psi \rangle^{n-1}}{y \langle \psi \rangle^n m_{3/2}/M_{PL}^{n-1}} \cdot \frac{g^2}{16\pi^2} 2C_Q m_{3/2} = \frac{g^2}{16\pi^2} 2C_Q \times \frac{-n \langle F_\psi \rangle}{\langle \psi \rangle}, \tag{25}
$$
where \( C_Q \) is a Dynkin index of \( Q \), which is normalized to be \( 1/2 \) for a fundamental representation.

From Eqs. (22) and (24), the \( F \) term of \( \psi \) is given by

\[
F_\psi = -m_{3/2} \langle \psi \rangle \frac{a_\bar{\psi} - a_\psi}{a_\bar{\psi} + a_\psi + 2\lambda^2 |\langle X \rangle|^2/m_{3/2}^2} \left( 1 + \mathcal{O} \left( \frac{m_{3/2}^2}{\lambda^2 v^2} \right) \right) \equiv -m_{3/2} \langle \psi \rangle \epsilon, \tag{26}
\]

where \( \epsilon \) is order one, unless the soft squared mass terms of \( \psi \) and \( \bar{\psi} \) accidentally coincide with each others.

By substituting Eq. (26) into Eq. (25), we obtain the contribution from the axion model to gaugino masses,

\[
\Delta m_\lambda = \frac{g^2}{16\pi^2} 2C_Q n \epsilon m_{3/2}. \tag{27}
\]

Note that the phase is aligned with the AMSB contribution. This is because the phase of \( \langle \psi \rangle \) and \( \langle F_\psi \rangle \) are aligned with each others.

Let us comment on the case with several flavors of vector-like matters, as is the case with axion models presented in Ref. [93]. Even if there are several flavors of vector-like matters, we can always diagonalize their mass matrix. Each mass eigenstates contribution to the gaugino masses is as given in Eq. (27). The gaugino masses are simply multiplied by the number of the flavors.

### A.2 The SSM gaugino masses

In the presence of the axion model described above, the gaugino masses receive threshold corrections at the scale of the mass of \( Q\bar{Q} \). However, \( m_\lambda/g^2 \) is a renormalization invariant in supersymmetric theory at an one-loop level. Hence, it is not necessary to solve the renormalization equations from the mass scale of \( Q\bar{Q} \) to the gravitino mass scale for an one-loop analysis. We can treat the corrections given by Eq. (27) as if it is generated at the gravitino mass scale, and solve the renormalization equations (18). Therefore, in this axion model, gaugino masses are given by the upper left panel of Fig. 1. Phenomenology can be read off from the line \( \gamma = 0 \) of Figures 3, 4 and 5. In axion models with a large number of additional matter [93], \( N_{\text{eff}} \) would be considerable.
References

[1] L. Maiani. in Proceedings: Summer School on Particle Physics, Paris, France (1979).

[2] M. J. G. Veltman, Acta Phys. Polon. B 12, 437 (1981).

[3] E. Witten, Nucl. Phys. B 188, 513 (1981).

[4] R. K. Kaul, Phys. Lett. B 109, 19 (1982), and references therein.

[5] G. F. Giudice, M. A. Luty, H. Murayama and R. Rattazzi, JHEP 9812, 027 (1998) [hep-ph/9810442].

[6] J. D. Wells, hep-ph/0306127; J. D. Wells, Phys. Rev. D 71, 015013 (2005) [hep-ph/0411041].

[7] G. D. Coughlan, W. Fischler, E. W. Kolb, S. Raby and G. G. Ross, Phys. Lett. B 131, 59 (1983).

[8] M. Ibe, Y. Shinbara and T. T. Yanagida, Phys. Lett. B 639, 534 (2006) [hep-ph/0605252].

[9] H. Pagels and J. R. Primack, Phys. Rev. Lett. 48, 223 (1982); S. Weinberg, Phys. Rev. Lett. 48, 1303 (1982); M. Y. Khlopov and A. D. Linde, Phys. Lett. B 138, 265 (1984).

[10] M. Kawasaki, K. Kohri and T. Moroi, Phys. Rev. D 71 (2005) 083502 [arXiv:astro-ph/0408426]; K. Jedamzik, Phys. Rev. D 74, 103509 (2006) [arXiv:hep-ph/0604251]; M. Kawasaki, K. Kohri, T. Moroi and A. Yotsuyanagi, Phys. Rev. D 78, 065011 (2008) [arXiv:0804.3745 [hep-ph]], and references therein.

[11] M. Fukugita and T. Yanagida, Phys. Lett. B174 (1986) 45; For reviews, W. Buchmuller, P. Di Bari and M. Plumacher, Annals Phys. 315, 305 (2005) [hep-ph/0401240]; W. Buchmuller, R. D. Peccei and T. Yanagida, Ann. Rev. Nucl. Part. Sci. 55, 311 (2005) [arXiv:hep-ph/0502169]; S. Davidson, E. Nardi and Y. Nir, Phys. Rept. 466, 105 (2008) [arXiv:0802.2962 [hep-ph]].

[12] S. Weinberg, Phys. Rev. Lett. 59, 2607 (1987).

[13] R. Bousso and J. Polchinski, JHEP 0006, 006 (2000) [hep-th/0004134].
[14] S. Kachru, R. Kallosh, A. D. Linde and S. P. Trivedi, Phys. Rev. D 68, 046005 (2003) [hep-th/0301240].

[15] L. Susskind, In *Carr, Bernard (ed.): Universe or multiverse?* 247-266 [hep-th/0302219].

[16] F. Denef and M. R. Douglas, JHEP 0405, 072 (2004) [hep-th/0404116].

[17] V. Agrawal, S. M. Barr, J. F. Donoghue and D. Seckel, Phys. Rev. Lett. 80, 1822 (1998) [hep-ph/9801253].

[18] T. E. Jeltema and M. Sher, Phys. Rev. D 61, 017301 (2000) [hep-ph/9905494].

[19] N. Arkani-Hamed and S. Dimopoulos, JHEP 0506, 073 (2005) [hep-th/0405159].

[20] G. F. Giudice and A. Romanino, Nucl. Phys. B 699, 65 (2004) [Erratum-ibid. B 706, 65 (2005)] [hep-ph/0406088].

[21] N. Arkani-Hamed, S. Dimopoulos, G. F. Giudice and A. Romanino, Nucl. Phys. B 709, 3 (2005) [hep-ph/0409232].

[22] M. Ibe, T. Moroi and T. T. Yanagida, Phys. Lett. B 644, 355 (2007) [hep-ph/0610277].

[23] K. Inoue, M. Kawasaki, M. Yamaguchi and T. Yanagida, Phys. Rev. D 45, 328 (1992).

[24] J. A. Casas and C. Munoz, Phys. Lett. B 306, 288 (1993) [hep-ph/9302227].

[25] G. F. Giudice and A. Masiero, Phys. Lett. B 206, 480 (1988).

[26] M. Ibe and T. T. Yanagida, Phys. Lett. B 709, 374 (2012) [arXiv:1112.2462 [hep-ph]].

[27] M. Ibe, S. Matsumoto and T. T. Yanagida, Phys. Rev. D 85, 095011 (2012) [arXiv:1202.2253 [hep-ph]].

[28] N. Arkani-Hamed, IFT Inaugural Conference (2011), http://www.ift.uam.es/workshops/Xmas11/?q=node/2.

[29] N. Arkani-Hamed, A. Gupta, D. E. Kaplan, N. Weiner and T. Zorawski, arXiv:1212.6971 [hep-ph].

[30] L. J. Hall and Y. Nomura, JHEP 1201, 082 (2012) [arXiv:1111.4519 [hep-ph]].
[31] A. Arvanitaki, N. Craig, S. Dimopoulos and G. Villadoro, JHEP 1302, 126 (2013) [arXiv:1210.0555 [hep-ph]].

[32] For a review, H. P. Nilles, Phys. Rept. 110 (1984) 1.

[33] M. Dine and D. MacIntire, Phys. Rev. D 46, 2594 (1992) [hep-ph/9205227].

[34] L. Randall and R. Sundrum, Nucl. Phys. B 557, 79 (1999) [hep-th/9810155].

[35] G. Aad et al. [ATLAS Collaboration], Phys. Lett. B 716, 1 (2012) [arXiv:1207.7214 [hep-ex]].

[36] S. Chatrchyan et al. [CMS Collaboration], Phys. Lett. B 716, 30 (2012) [arXiv:1207.7235 [hep-ex]].

[37] Y. Okada, M. Yamaguchi and T. Yanagida, Phys. Lett. B 262, 54 (1991); see also Y. Okada, M. Yamaguchi and T. Yanagida, Prog. Theor. Phys. 85, 1 (1991); J. R. Ellis, G. Ridolfi and F. Zwirner, Phys. Lett. B 257, 83 (1991); H. E. Haber and R. Hempfling, Phys. Rev. Lett. 66, 1815 (1991).

[38] D. S. M. Alves, E. Izaguirre and J. G. Wacker, arXiv:1108.3390 [hep-ph].

[39] G. F. Giudice and A. Strumia, Nucl. Phys. B 858, 63 (2012) [arXiv:1108.6077 [hep-ph]].

[40] R. Sato, S. Shirai and K. Tobioka, JHEP 1211, 041 (2012) [arXiv:1207.3608 [hep-ph]].

[41] B. Bhattacharjee, B. Feldstein, M. Ibe, S. Matsumoto and T. T. Yanagida, Phys. Rev. D 87, 015028 (2013) [arXiv:1207.5453 [hep-ph]].

[42] L. J. Hall, Y. Nomura and S. Shirai, JHEP 1301, 036 (2013) [arXiv:1210.2395 [hep-ph]].

[43] M. Ibe, S. Matsumoto and R. Sato, Phys. Lett. B 721, 252 (2013) [arXiv:1212.5989 [hep-ph]].

[44] R. Sato, S. Shirai and K. Tobioka, arXiv:1307.7144 [hep-ph].

[45] J. Hisano, T. Kuwahara and N. Nagata, Phys. Lett. B 723, 324 (2013) [arXiv:1304.0343 [hep-ph]].
[46] J. Hisano, D. Kobayashi, T. Kuwahara and N. Nagata, JHEP 1307, 038 (2013) [arXiv:1304.3651 [hep-ph]].

[47] W. Altmannshofer, R. Harnik and J. Zupan, arXiv:1308.3653 [hep-ph].

[48] K. Fuyuto, J. Hisano, N. Nagata and K. Tsumura, arXiv:1308.6493 [hep-ph].

[49] A. Linde, Y. Mambrini and K. A. Olive, Phys. Rev. D 85, 066005 (2012) [arXiv:1111.1465 [hep-th]].

[50] B. S. Acharya, G. Kane and P. Kumar, Int. J. Mod. Phys. A 27, 1230012 (2012) [arXiv:1204.2795 [hep-ph]].

[51] J. L. Evans, M. Ibe, K. A. Olive and T. T. Yanagida, Eur. Phys. J. C 73, 2468 (2013) [arXiv:1302.5346 [hep-ph]].

[52] J. L. Evans, K. A. Olive, M. Ibe and T. T. Yanagida, arXiv:1305.7461 [hep-ph].

[53] K. Nakayama and F. Takahashi, JCAP 1205, 035 (2012) [arXiv:1203.0323 [hep-ph]].

[54] B. Feldstein and T. T. Yanagida, Phys. Lett. B 720, 166 (2013) [arXiv:1210.7578 [hep-ph]].

[55] K. Harigaya, M. Kawasaki and T. T. Yanagida, Phys. Lett. B 719, 126 (2013) [arXiv:1211.1770 [hep-ph]].

[56] W. Buchmuller, V. Domcke, K. Kamada and K. Schmitz, arXiv:1309.7788 [hep-ph].

[57] A. E. Nelson and N. J. Weiner, hep-ph/0210288.

[58] K. Hsieh and M. A. Luty, JHEP 0706, 062 (2007) [hep-ph/0604256].

[59] A. Gupta, D. E. Kaplan and T. Zorawski, arXiv:1212.6969 [hep-ph].

[60] K. Nakayama and T. T. Yanagida, Phys. Lett. B 722, 107 (2013) [arXiv:1302.3332 [hep-ph]].

[61] T. Moroi and L. Randall, Nucl. Phys. B 570, 455 (2000) [hep-ph/9906527].

[62] T. Gherghetta, G. F. Giudice and J. D. Wells, Nucl. Phys. B 559, 27 (1999) [hep-ph/9904378].

[63] J. Hisano, S. Matsumoto and M. M. Nojiri, Phys. Rev. Lett. 92, 031303 (2004) [hep-ph/0307216].
[64] M. Ibe, R. Kitano, H. Murayama and T. Yanagida, Phys. Rev. D 70, 075012 (2004) [hep-ph/0403198].

[65] F. D’Eramo, J. Thaler and Z. Thomas, arXiv:1307.3251 [hep-ph].

[66] A. Pomarol and R. Rattazzi, JHEP 9905, 013 (1999) [hep-ph/9903448].

[67] E. Poppitz and S. P. Trivedi, Phys. Lett. B 401, 38 (1997) [hep-ph/9703246].

[68] F. Gabbiani, E. Gabrielli, A. Masiero and L. Silvestrini, Nucl. Phys. B 477, 321 (1996) [hep-ph/9604387].

[69] N. Arkani-Hamed, S. Dimopoulos, G. F. Giudice and A. Romanino, Nucl. Phys. B 709, 3 (2005) [hep-ph/0409232].

[70] P. A. R. Ade et al. [Planck Collaboration], arXiv:1303.5076 [astro-ph.CO].

[71] ATLAS-CONF-2013-069 [ATLAS Collaboration].

[72] ATLAS-CONF-2013-047 [ATLAS Collaboration].

[73] G. Belanger, F. Boudjema, A. Pukhov and A. Semenov, arXiv:1305.0237 [hep-ph].

[74] J. Hisano, S. Matsumoto, M. Nagai, O. Saito and M. Senami, Phys. Lett. B 646, 34 (2007) [hep-ph/0610249].

[75] K. Harigaya, K. Kaneta and S. Matsumoto, in preperation.

[76] A. Birkedal-Hansen and B. D. Nelson, Phys. Rev. D 64, 015008 (2001) [hep-ph/0102075].

[77] S. Profumo and C. E. Yaguna, Phys. Rev. D 69, 115009 (2004) [hep-ph/0402208].

[78] T. Cohen, M. Lisanti, A. Pierce and T. R. Slatyer, arXiv:1307.4082 [hep-ph].

[79] J. Fan and M. Reece, arXiv:1307.4400 [hep-ph].

[80] F. Nesti and P. Salucci, JCAP 1307, 016 (2013) arXiv:1304.5127 [astro-ph.GA].

[81] M. Ackermann et al. [Fermi-LAT Collaboration], arXiv:1310.0828 [astro-ph.HE].

[82] J. Alwall, K. Hiramatsu, M. M. Nojiri and Y. Shimizu, Phys. Rev. Lett. 103, 151802 (2009) arXiv:0905.1201 [hep-ph].

[83] B. Bhattacherjee, A. Choudhury, K. Ghosh and S. Poddar, arXiv:1308.1526 [hep-ph].
[84] M. Ibe, S. Matsumoto, S. Shirai and T. T. Yanagida, JHEP 1307, 063 (2013) [arXiv:1305.0084 [hep-ph]].

[85] O. Adriani et al. [PAMELA Collaboration], Nature 458, 607 (2009) [arXiv:0810.4995 [astro-ph]].

[86] M. Ackermann et al. [Fermi LAT Collaboration], Phys. Rev. Lett. 108, 011103 (2012) [arXiv:1109.0521 [astro-ph.HE]].

[87] M. Aguilar et al. [AMS Collaboration], Phys. Rev. Lett. 110, no. 14, 141102 (2013).

[88] J. E. Kim, Phys. Rev. Lett. 43, 103 (1979).

[89] M. A. Shifman, A. I. Vainshtein and V. I. Zakharov, Nucl. Phys. B 166, 493 (1980).

[90] R. D. Peccei and H. R. Quinn, Phys. Rev. Lett. 38, 1440 (1977); R. D. Peccei and H. R. Quinn, Phys. Rev. D 16, 1791 (1977).

[91] S. Weinberg, Phys. Rev. Lett. 40, 223 (1978).

[92] F. Wilczek, Phys. Rev. Lett. 40, 279 (1978).

[93] K. Harigaya, M. Ibe, K. Schmitz and T. T. Yanagida, arXiv:1308.1227 [hep-ph].
Figure 3: The gaugino masses as functions of $N_{\text{eff}}$ and $\gamma$ for a given $m_{3/2}$ and $L$. The solid lines show the ratio between $M_3/M_{\text{LSP}}$. The dashed lines show the wino mass. The LSP abundance is larger than the observed abundance in the light shaded regions (light-blue), where the LSP is the bino. In the pink shaded regions, the LSP is the bino but the relic abundance does not exceed the observation due to coannihilation effects. Here, we have neglected the effect of the Sommerfeld enhancement for simplicity (see footnote 8). The gluino is the LSP in the brown region. The wino LSP has been excluded in the dark shaded regions by disappearing track searches [71]. The green shaded regions are excluded by gluino searches [72].
Figure 4: The same as Figure 3 but for $m_{3/2} = 100$ TeV and $L = 0, m_{3/2}$. Here, we have neglected the effect of the Sommerfeld enhancement for simplicity.
Figure 5: The same as Figure 3 but for $m_{3/2} = 300$ TeV and $L = 0$. In the light shaded region (light-blue) in the right hand side, the LSP is mostly wino unlike the previous two figures. The closeup view of the gluino-wino coannihilation region is shown in the lower right panel. The lower left panel is the closeup view of the wino-bino coannihilation region. Here, we have neglected the effects of the Sommerfeld enhancement for simplicity.