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Moduli/Inflaton Mixing with Supersymmetry Breaking Field

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Abstract

A heavy scalar field such as moduli or an inflaton generally mixes with a field responsible for the supersymmetry breaking. We study the scalar decay into the standard model particles and their superpartners, gravitinos, and the supersymmetry breaking sector, particularly paying attention to decay modes that proceed via the mixing between the scalar and the supersymmetry breaking field. The impacts of the new decay processes on cosmological scenarios are also discussed; the modulus field generically produces too many gravitinos, and most of the inflation models tend to result in too high reheating temperature and/or gravitino overproduction.

I. INTRODUCTION

Scalar fields play an important role in the thermal history of the universe. Once a scalar field dominates the energy density of the universe, the subsequent evolution of the universe strongly depends on the reheating processes characterized by the decay temperature and the decay products.

Such scalar fields, symbolically denoted by ϕ, may be identified with an inflaton or moduli fields. A modulus field generally acquires nonzero vacuum expectation value (VEV) in the vacuum. Inflaton fields as well have non-vanishing VEVs in many inflation models. Once a scalar field obtains a nonzero VEV, $\phi_0 \equiv \langle\phi\rangle$, there is no remnant symmetry to forbid mixings of ϕ with the other fields, since the symmetries under which ϕ is charged, if any, are spontaneously broken in the vacuum. There is another important scalar field, z, which is responsible for the supersymmetry (SUSY) breaking. The presence of such SUSY breaking field is inevitable in the SUSY theories, because of an absence of the light superparticles. The SUSY breaking field, z, must be singlet under any unbroken symmetries at the vacuum in order for the auxiliary field, G_z, to obtain a finite VEV. Therefore the scalar field z as well generally obtains a VEV, $z_0 \equiv \langle z\rangle$.

We would like to stress that a scalar field ϕ with nonzero VEV, such as the inflaton and moduli, generically *mixes* with the SUSY breaking field z in the vacuum. In particular such mixing has impacts on the decay processes of ϕ. It has been recently argued that the modulus and inflaton decays may produce too many gravitinos and/or the lightest SUSY particle (LSP) [1, 2, 3, 4]. In Ref. [5], however, it has been demonstrated that the gravitino production rate can be suppressed by taking account of the mixing of ϕ with z in some explicit models. In this paper, we develop general analyses on the mixture of ϕ and z, and discuss its cosmological consequences, paying particular attention to the decay of ϕ via the mixing with z.

In the next section, we develop a formalism to obtain the mass-eigenstate basis and clarify the relation between the mass-eigenstate basis and the model basis. In Sec. III, we consider several decay processes in the mass eigenstates, especially those induced via the mixing with the SUSY breaking sector, in the gravity-mediated SUSY breaking scenario. We also discuss how the modulus and inflaton cosmology is affected by the mixing. In Sec. IV we take up the low energy SUSY breaking models such as the gauge-mediated SUSY breaking

(GMSB) models [6], clarifying the difference from the case of gravity mediation. Sec. V is devoted to discussions on miscellaneous topics. We give a summary in the last section. In Appendix. A, we show the goldstino interpretation of the scalar decay into gravitinos and see the equivalence between the two pictures.

II. MASS-EIGENSTATE BASIS

A scalar decay must be considered in its mass-eigenstate basis, while a model is often given in such a way that particles in the model are not mass eigenstates especially if some symmetries are spontaneously broken in the vacuum. In particular, it is quite probable that a scalar ϕ with nonzero VEV mixes with the SUSY breaking field z in the vacuum, since there is no remnant symmetry that forbids the mixing. The kinetic term and non-analytic (NA) and analytic (A) mass terms of ϕ and z in the model frame are given as

$$\mathcal{L}_{\rm kin.} = \partial_\mu \phi^\dagger \partial^\mu \phi + \partial_\mu z^\dagger \partial^\mu z + g_{z\bar{\phi}} \partial_\mu \phi^\dagger \partial^\mu z + g_{\phi\bar{z}} \partial_\mu z^\dagger \partial^\mu \phi, \tag{1}$$

$$-\mathcal{L}_{\rm mass}^{\rm (NA)} = M^2_{\phi\bar{\phi}} \phi^\dagger \phi + M^2_{z\bar{z}} z^\dagger z + M^2_{z\bar{\phi}} \phi^\dagger z + M^2_{\phi\bar{z}} z^\dagger \phi, \tag{2}$$

$$-\mathcal{L}_{\rm mass}^{\rm (A)} = \frac{1}{2} M^2_{\phi\phi} \phi\phi + \frac{1}{2} M^2_{zz} zz + M^2_{\phi z} \phi z + {\rm h.c.}, \tag{3}$$

where the fields are expanded around the VEV, $\phi \to \phi - \phi_0$ and $z \to z - z_0$. The mixings in the kinetic term, $g_{z\bar{\phi}}$ and $g_{\phi\bar{z}}$, are given by

$$g_{z\bar{\phi}} = \left\langle \frac{\partial^2 K}{\partial z \partial \phi^\dagger} \right\rangle, \quad g_{\phi\bar{z}} = \left\langle \frac{\partial^2 K}{\partial \phi \partial z^\dagger} \right\rangle, \tag{4}$$

where K is the Kähler potential, while $g_{\phi\bar{\phi}}$ and $g_{z\bar{z}}$ are normalized to be unity. Note that the cross term $g_{z\bar{\phi}} \partial_\mu \phi^\dagger \partial^\mu z$ naturally appears if there are higher order terms in the Kähler potential before the fields are expanded around the VEV. The purpose of this section is to clarify the relation between the model basis (ϕ, z) and the mass-eigenstate basis.

In the Einstein frame, the 4D $N = 1$ supergravity (SUGRA) Lagrangian contains the scalar potential, $V = e^G(G^i G_i - 3)$ [a]. The (non-)analytic mass terms can be written in terms of the total Kähler potential, $G = K + \ln |W|^2$, as

$$M^2_{ij^*} = \frac{\partial^2 V}{\partial \varphi^i \partial \varphi^{\dagger j}} = e^G \left(\nabla_i G_k \nabla_{j^*} G^k - R_{ij^* k \ell^*} G^k G^{\ell *} + g_{ij^*} \right), \tag{5}$$

$$M^2_{ij} = M^2_{ji} = \frac{\partial^2 V}{\partial \varphi^i \partial \varphi^j} = e^G \left(\nabla_i G_j + \nabla_j G_i + G^k \nabla_i \nabla_j G_k \right), \tag{6}$$

[a] Throughout this paper we assume that the D-term potential is negligible.

where we have assumed the vanishing cosmological constant, $G^i G_i = 3$, and used the potential minimization condition, $G^i \nabla_k G_i + G_k = 0$ in the vacuum. The gravitino mass is given by $m_{3/2} = \left\langle e^{G/2} \right\rangle$. Here and in what follows, the subscript i denotes a derivative with respect to the field φ^i, and the superscript is defined by $G^i = g^{ij^*} G_{j*}$. Here g_{ij^*} is the Kähler metric, $g_{ij^*} = G_{ij^*}$, and $R_{ij^*k\ell^*}$ is the curvature of the Kähler manifold, defined by $R_{ij^*k\ell^*} = g_{ij^*k\ell^*} - g^{mn^*} g_{mj^*\ell^*} g_{n^*ik}$. Also the covariant derivative of G_i is defined by $\nabla_i G_j = G_{ij} - \Gamma^k_{ij} G_k$, where the connection, $\Gamma^k_{ij} = g^{k\ell^*} g_{ij\ell^*}$, and $\nabla_k g_{ij^*} = 0$ is satisfied.

Throughout this paper, the scalar field, ϕ, is assumed to be much heavier than the gravitino due to a large supersymmetric mass, $m_\phi / m_{3/2} \equiv |\nabla_\phi G_\phi| \gg 1$. The SUSY breaking field z is such that it sets the cosmological constant to be zero, i.e., $G^z G_z \simeq 3$, while ϕ is assumed to give only subdominant contribution to the SUSY breaking, i.e., $|G_\phi| \ll 1$. Then, as long as $|\nabla_\phi G_z| \lesssim O(1)$ [b], the potential minimization condition for z,

$$G^z \nabla_z G_z + G^\phi \nabla_z G_\phi + G_z \ = \ 1 \,, \tag{7}$$

requires that the supersymmetric mass of z is equal to the gravitino mass, i.e., $|\nabla_z G_z| \simeq 1$ [c]. We assume that this is the case. It should be noted, however, that the scalar mass of z can be larger than $m_{3/2}$ due to the non-supersymmetric mass term, $e^G R_{z\bar{z}k\ell^*} G^k G^{\ell^*}$, if one adds, e.g. $\delta K = -|z|^4/\Lambda^2$ with a low cut-off scale $\Lambda \ll M_P$, to the Kähler potential, which leads to $m_z^2 \simeq 12 m_{3/2}^2 (M_P/\Lambda)^2$.

In the following we assume $M^2_{\phi\bar{\phi}}$ dominates over the other components of the mass terms. The results in the case of $M^2_{z\bar{z}} \gg M^2_{\phi\bar{\phi}} \gg$ (the other elements) will be given in the last of this section. The rest discussion of this section is however rather generic, and can be applied not only to the situation we stated above.

The kinetic term can be canonically normalized by a shift of z and a rescaling of ϕ;

$$\phi' \ = \ (1 - |g_{\phi\bar{z}}|^2)^{-1/2} \phi, \tag{8}$$

$$z' \ = \ z + g_{\phi\bar{z}} \phi, \tag{9}$$

$$\mathcal{L}_{\text{kin.}} \ = \ \partial_\mu \phi'^\dagger \partial^\mu \phi' + \partial_\mu z'^\dagger \partial^\mu z'. \tag{10}$$

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[b] In fact, according to the discussion of Ref. [1], $|\nabla_\phi G_z| \sim O(1)$ holds for modulus field with its VEV of the Planck scale, if the Kähler potential does not have any enhancement factor. In principle, $|\nabla_\phi G_z|$ could be larger than O(1) if the higher order term $g_{\phi z\bar{z}}$ in the Kähler potential is larger than unity. However, such a large mixing in the supersymmetric mass obscures the definitions (or roles) of the different two fields in the model basis. Also it makes the gravitino problem even worse.

[c] This was also noted in Ref. [7] in a different context.

Then the mass terms become

$$
\begin{aligned}
-\mathcal{L}^{(NA)}_{\rm mass} &\equiv M^2_{\phi'\bar{\phi}'}\phi'^\dagger\phi' + M^2_{z'\bar{z}'}z'^\dagger z' + M^2_{z'\bar{\phi}'}\phi'^\dagger z' + M^2_{\phi'\bar{z}'}z'^\dagger\phi' \\
&\simeq (M^2_{\phi\bar{\phi}} - g_{\phi\bar{z}}M^2_{z\bar{\phi}} - g_{z\bar{\phi}}M^2_{\phi\bar{z}})\phi'^\dagger\phi' + M^2_{z\bar{z}}z'^\dagger z' \\
&\quad +(M^2_{z\bar{\phi}} - g_{z\bar{\phi}}M^2_{z\bar{z}})\phi'^\dagger z' + (M^2_{\phi\bar{z}} - g_{\phi\bar{z}}M^2_{z\bar{z}})z'^\dagger\phi', \qquad (11)\\
-\mathcal{L}^{(A)}_{\rm mass} &\simeq \frac{1}{2}(M^2_{\phi\phi} - 2g_{\phi\bar{z}}M^2_{\phi z})\phi'\phi' + \frac{1}{2}M^2_{zz}z'z' \\
&\quad +(M^2_{\phi z} - g_{\phi\bar{z}}M^2_{zz})\phi' z' + \text{h.c.}, \qquad (12)
\end{aligned}
$$

where $\mathcal{O}(g^2_{\phi\bar{z}})$ terms are omitted in each term.

Let us first diagonalize the non-analytic mass terms while keeping the kinetic term canonical by the following transformation:

$$
\begin{aligned}
\Phi &\equiv \phi' + \epsilon z', \\
Z &\equiv z' - \epsilon^*\phi', \qquad (13)
\end{aligned}
$$

where ϵ represents the mixing angle. Here we have assumed $|\epsilon| \ll 1$ and neglected those terms of $O(\epsilon^2)$. Since $M^2_{\phi'\bar{\phi}'}$ dominates over the other components in the mass matrix, the mixing angle is given by the ratio of $M^2_{\phi'\bar{\phi}'}$ to the off-diagonal component:

$$
\epsilon \simeq \frac{M^2_{z'\bar{\phi}'}}{M^2_{\phi'\bar{\phi}'}} \simeq \frac{M^2_{z\bar{\phi}} - g_{z\bar{\phi}}M^2_{z\bar{z}}}{M^2_{\phi\bar{\phi}}}. \qquad (14)
$$

Then the non-analytic mass matrix becomes diagonal in this basis:

$$
-\mathcal{L}^{(NA)}_{\rm mass} \simeq M^2_{\phi\bar{\phi}}\Phi^\dagger\Phi + \left(M^2_{z\bar{z}} - \frac{|M^2_{\phi\bar{z}}|^2}{M^2_{\phi\bar{\phi}}}\right)Z^\dagger Z, \qquad (15)
$$

We will call this basis (Φ, Z) as the NA mass-eigenstate basis in the following.

The physical processes become easy to be considered after the mass matrix is fully diagonalized. In particular, one should note that the analytic mass terms provide a further mixture between ϕ and $z^\dagger$ (z and $\phi^\dagger$). In the NA mass-eigenstate basis, the analytic mass terms become

$$
\begin{aligned}
-\mathcal{L}^{(A)}_{\rm mass} &\equiv \frac{1}{2}M^2_{\Phi\Phi}\Phi\Phi + \frac{1}{2}M^2_{ZZ}ZZ + M^2_{\Phi Z}\Phi Z + \text{h.c.} \\
&\simeq \frac{1}{2}\left(M^2_{\phi\phi} - 2(g_{\phi\bar{z}} - \epsilon^*)M^2_{\phi z}\right)\Phi\Phi + \frac{1}{2}\left(M^2_{zz} - 2\epsilon M^2_{\phi z}\right)ZZ \\
&\quad + \left(M^2_{\phi z} - M^2_{zz}g_{\phi\bar{z}} - M^2_{\phi\phi}\epsilon + M^2_{zz}\epsilon^*\right)\Phi Z + \text{h.c.} \qquad (16)
\end{aligned}
$$

up to $O(\epsilon)$. Let us concentrate on the mixings $\Phi - Z^\dagger$ and $Z - \Phi^\dagger$, since they become important in the following analyses. To diagonalize the analytic mass term, we take the following transformation,

$$\tilde{\Phi} \equiv \Phi + \tilde{\epsilon} Z^\dagger, \tag{17}$$

$$\tilde{Z} \equiv Z - \tilde{\epsilon} \Phi^\dagger, \tag{18}$$

where we have assumed that the mixing angle $\tilde{\epsilon}$ is much smaller than unity. Since the dominant contribution to the total mass matrix comes from the non-analytic mass component, $M^2_{\Phi\bar{\Phi}} \simeq m^2_\phi$, $\tilde{\epsilon}$ is given by the ratio of $M^2_{\Phi\bar{\Phi}}$ to $M^2_{\bar{\Phi}\bar{Z}}$:

$$\tilde{\epsilon} \simeq \frac{M^2_{\bar{\Phi}\bar{Z}}}{M^2_{\Phi\bar{\Phi}}} \simeq \frac{M^2_{\bar{\phi}\bar{z}} - g_{z\bar{\phi}} M^2_{\bar{z}\bar{z}}}{M^2_{\phi\bar{\phi}}}. \tag{19}$$

Thus obtained $(\tilde{\Phi}, \tilde{Z})$ is the desired mass-eigenstate basis. Although $M^2_{\Phi\Phi}$ (M^2_{ZZ}) can further mix the real and imaginary components of Φ (Z), it does not modify the following discussions qualitatively.

Here, we summarize the relation between the model basis (ϕ, z) and the mass-eigenstate basis $(\tilde{\Phi}, \tilde{Z})$.

$$\phi \simeq \tilde{\Phi} - \epsilon \tilde{Z} - \tilde{\epsilon} \tilde{Z}^\dagger, \tag{20}$$

$$z \simeq \tilde{Z} + (-g_{\phi\bar{z}} + \epsilon^*)\tilde{\Phi} + \tilde{\epsilon}\Phi^\dagger. \tag{21}$$

The explicit expressions for ϵ and $\tilde{\epsilon}$ are given by (14) and (19).

So far, we have assumed that $M^2_{\phi\bar{\phi}}$ dominates over the other components of the mass matrix. If z acquires a non-supersymmetric mass larger than the mass of ϕ, we can repeat a similar discussion to obtain the relation between the model basis and the mass-eigenstate basis:

$$\phi \simeq \tilde{\Phi} + (-g_{z\bar{\phi}} + \epsilon^*)\tilde{Z} + \tilde{\epsilon} Z^\dagger, \tag{22}$$

$$z \simeq \tilde{Z} - \epsilon\tilde{\Phi} - \tilde{\epsilon}\tilde{\Phi}^\dagger \tag{23}$$

with ϵ and $\tilde{\epsilon}$ given by

$$\epsilon \simeq \frac{M^2_{\phi'\bar{z}'}}{M^2_{z'\bar{z}'}} \simeq \frac{M^2_{\phi\bar{z}} - g_{\phi\bar{z}} M^2_{\phi\bar{\phi}}}{M^2_{z\bar{z}}}, \tag{24}$$

$$\tilde{\epsilon} \simeq \frac{M^2_{\bar{Z}\bar{\Phi}}}{M^2_{Z\bar{Z}}} \simeq \frac{M^2_{\bar{z}\bar{\phi}} - g_{\phi\bar{z}} M^2_{\bar{\phi}\bar{\phi}}}{M^2_{z\bar{z}}}. \tag{25}$$

In the following sections, we are particularly interested in the mixing of $\tilde{\Phi}(\tilde{\Phi}^\dagger)$ into z or $\tilde{Z}(\tilde{Z}^\dagger)$ into ϕ. Although three sets of the transformations are necessary to arrive at the mass eigenstate basis, the effective mixing angle is given by the largest mixing among them. To parametrize this effective mixing of $\tilde{\Phi}(\tilde{\Phi}^\dagger)$ into z, we define $\epsilon_{z\tilde{\Phi}}$ as

$$\epsilon_{z\tilde{\Phi}} \equiv \begin{cases} \mathrm{Max}\left\{ |g_{\phi\bar{z}}|, \left|\dfrac{M^2_{\phi\bar{z}}}{M^2_{\phi\bar{\phi}}}\right|, \left|\dfrac{M^2_{\phi z}}{M^2_{\phi\bar{\phi}}}\right| \right\} & \text{for} \quad M^2_{\phi\bar{\phi}} \gg M^2_{z\bar{z}}, \\ \\ \mathrm{Max}\left\{ \left|g_{\phi\bar{z}}\dfrac{M^2_{\phi\bar{\phi}}}{M^2_{z\bar{z}}}\right|, \left|\dfrac{M^2_{\phi\bar{z}}}{M^2_{z\bar{z}}}\right|, \left|\dfrac{M^2_{\phi z}}{M^2_{z\bar{z}}}\right| \right\} & \text{for} \quad M^2_{z\bar{z}} \gg M^2_{\phi\bar{\phi}}. \end{cases} \tag{26}$$

Similarly, we define the effective mixing of $\tilde{Z}(\tilde{Z}^\dagger)$ into ϕ as

$$\epsilon_{\phi\tilde{Z}} \equiv \begin{cases} \mathrm{Max}\left\{ \left|g_{\phi\bar{z}}\dfrac{M^2_{z\bar{z}}}{M^2_{\phi\bar{\phi}}}\right|, \left|\dfrac{M^2_{\phi\bar{z}}}{M^2_{\phi\bar{\phi}}}\right|, \left|\dfrac{M^2_{\phi z}}{M^2_{\phi\bar{\phi}}}\right| \right\} & \text{for} \quad M^2_{\phi\bar{\phi}} \gg M^2_{z\bar{z}}, \\ \\ \mathrm{Max}\left\{ |g_{\phi\bar{z}}|, \left|\dfrac{M^2_{\phi\bar{z}}}{M^2_{z\bar{z}}}\right|, \left|\dfrac{M^2_{\phi z}}{M^2_{z\bar{z}}}\right| \right\} & \text{for} \quad M^2_{z\bar{z}} \gg M^2_{\phi\bar{\phi}}. \end{cases} \tag{27}$$

Therefore, using these the effective mixing angles, the relations (20), (21), (22), and (23) can be roughly expressed as

$$\phi \sim \tilde{\Phi} - \epsilon_{\phi\tilde{Z}}\tilde{Z} \tag{28}$$

$$z \sim \tilde{Z} + \epsilon_{z\tilde{\Phi}}\tilde{\Phi}, \tag{29}$$

up to phase, where we have also dropped the distinction between $\tilde{\Phi}(\tilde{Z})$ and its conjugate.

III. GRAVITY MEDIATION

Let us now consider the decay of $\tilde{\Phi}$ via the mixing with the SUSY breaking field z, and discuss its cosmological influence. To this end, we need to specify the way to mediate the SUSY breaking to the visible sector. In this section we consider the gravity mediation to exemplify how serious the problems caused by the mixing is.

A. Decay Modes

Let us study the scalar decay modes which proceed via the mixing with the SUSY breaking field. They are classified by the decay products: (i) the gravitinos; (ii) the SM particles (and their superpartners); (iii) the SUSY breaking fields. We discuss each case below.

1. Gravitino

The scalar field ϕ can decay into a pair of the gravitinos through the mixing with the SUSY breaking field [d]. The relevant couplings are [9, 10, 11]

$$
\begin{aligned}
e^{-1}\mathcal{L} \;=\; & -\frac{1}{8}\epsilon^{\mu\nu\rho\sigma}\left(G_\phi \partial_\rho \phi + G_z \partial_\rho z - \text{h.c.}\right)\bar{\psi}_\mu \gamma_\nu \psi_\sigma \\
& -\frac{1}{8}e^{G/2}\left(G_\phi \phi + G_z z + \text{h.c.}\right)\bar{\psi}_\mu \left[\gamma^\mu, \gamma^\nu\right]\psi_\nu,
\end{aligned}
\tag{30}
$$

where ψ_μ is the gravitino field, and we have chosen the unitary gauge in the Einstein frame with the Planck units, $M_P = 1$. One has to take account of the mixing between ϕ and $z(z^\dagger)$ discussed in the previous section, in order to evaluate the decay rate [5]. That is to say, we should rewrite the interactions in terms of the mass-eigenstate basis $(\tilde{\Phi}, \tilde{Z})$.

To this end, we first estimate the coupling to the gravitinos, G_Φ, in the NA mass-eigenstate basis $(\Phi,\, Z)$. In this basis, the off-diagonal components of the non-analytic mass term should vanish by definition:

$$
M^2_{\Phi\bar{Z}} = \nabla_\Phi G_\Phi \nabla_{\bar{Z}} G_{\bar{\Phi}} + \nabla_\Phi G_Z \nabla_{\bar{Z}} G_{\bar{Z}} - R_{\Phi\bar{Z}ij^*} G^i G^{j^*} = 0. \tag{31}
$$

Using $|\nabla_\Phi G_\Phi| \gg |\nabla_Z G_Z|$, we obtain

$$
\nabla_{\bar{Z}} G_{\bar{\Phi}} \;\simeq\; \frac{R_{\Phi\bar{Z}ij^*} G^i G^{j^*}}{\nabla_\Phi G_\Phi}. \tag{32}
$$

On the other hand, the potential minimization condition for Φ reads

$$
G_{\bar{\Phi}} \nabla_\Phi G_\Phi + G_{\bar{Z}} \nabla_\Phi G_Z + G_\Phi = 0, \tag{33}
$$

which can be solved for G_Φ:

$$
G_\Phi \;\simeq\; -\frac{\nabla_{\bar{\Phi}} G_{\bar{Z}}}{\nabla_{\bar{\Phi}} G_{\bar{\Phi}}} G_Z, \tag{34}
$$

where we have used $|\nabla_\Phi G_\Phi| \gg 1$ again. Substituting (32) into (34), we arrive at

$$
|G_\Phi| \;\simeq\; 3\sqrt{3}\frac{|R_{\Phi\bar{Z}Z\bar{Z}}|}{|\nabla_{\bar{\Phi}} G_{\bar{\Phi}}|^2}, \tag{35}
$$

where we have used $|G_Z| = |G^Z| \simeq \sqrt{3}$. Thus G_Φ is always proportional to $m_{3/2}^2/m_\phi^2 \ll 1$, while it can be enhanced if Z has a quite large SUSY breaking mass, $m_z^2 \simeq 3|R_{z\bar{z}z\bar{z}}|e^G \gg$

---

[d] If the large scalar mass originates from non-supersymmetric mass terms, the single-gravitino production rate (cf. [8]) dominates over the pair production rate, irrespective of the mixing with the SUSY breaking field.

e^G. It should be noted that (35) always holds in the NA mass-eigenstate basis as long as $m_\phi \gg m_{3/2}$, irrespective of the value of m_z. For the minimal Kähler potential, G_Φ is exactly zero in this basis. However one must keep in mind that Φ is generally not identical to the true mass eigenstate $\tilde{\Phi}$. In general, the true mass eigenstate $\tilde{\Phi}$ ($\tilde{Z}$) is no longer a scalar component of a chiral superfield [see Eqs. (17) and (18)], and hence the above consideration for G_Φ does not hold.

Now let us write down the interactions in the mass-eigenstate basis $(\tilde{\Phi}, \tilde{Z})$. In the NA mass-eigenstate basis, the couplings to the gravitinos are obtained by replacing (ϕ, z) with (Φ, Z) in (30). By performing the transformation (17) and (18), we can rewrite the interactions in terms of $\tilde{\Phi}$ and $\tilde{Z}$:

$$\begin{aligned} e^{-1}\mathcal{L} \;\simeq\; & -\frac{1}{8}\epsilon^{\mu\nu\rho\sigma}\left(\mathcal{G}_\Phi^{(-)}\,\partial_\rho\tilde{\Phi} - \mathcal{G}_\Phi^{(+)\dagger}\,\partial_\rho\tilde{\Phi}^\dagger + G_Z\partial_\rho\tilde{Z} - G_{\bar{Z}}\partial_\rho\tilde{Z}^\dagger\right)\bar{\psi}_\mu\gamma_\nu\psi_\sigma \\ & -\frac{1}{8}e^{G/2}\left(\mathcal{G}_\Phi^{(+)}\;\tilde{\Phi} + \mathcal{G}_\Phi^{(-)\dagger}\;\tilde{\Phi}^\dagger + G_Z\tilde{Z} + G_{\bar{Z}}\tilde{Z}^\dagger\right)\bar{\psi}_\mu\left[\gamma^\mu, \gamma^\nu\right]\psi_\nu, \end{aligned} \tag{36}$$

where we have defined

$$\mathcal{G}_\Phi^{(\pm)} \;\equiv\; G_\Phi \pm \tilde{\epsilon}^* G_{\bar{Z}}. \tag{37}$$

The decay rate of $\tilde{\Phi}$ is [1, 2]

$$\Gamma(\tilde{\Phi} \to 2\psi_{3/2}) \;\simeq\; \frac{|\mathcal{G}_\Phi^{(eff)}|^2}{288\pi}\frac{m_\phi^5}{m_{3/2}^2 M_P^2}, \tag{38}$$

for $m_\phi \gg m_{3/2}$, where we defined $|\mathcal{G}_\Phi^{(eff)}|^2 \equiv 1/2\,(|\mathcal{G}_\Phi^+|^2 + |\mathcal{G}_\Phi^-|^2) = |G_\Phi|^2 + |\tilde{\epsilon}^* G_{\bar{Z}}|^2$. The gravitino mass in the denominator arises from the longitudinal component of the gravitino. An interpretation in the goldstino limit is given in the Appendix. A.

Let us now evaluate the order-of-magnitude of $|\mathcal{G}_\Phi^{(eff)}|^2 = |G_\Phi|^2 + |\tilde{\epsilon}^* G_{\bar{Z}}|^2$. The first term can be related to m_z if z is heavier than the gravitino due to a non-supersymmetric mass, $m_z^2 \simeq 3|R_{z\bar{z}z\bar{z}}|e^G \gg m_{3/2}^2$:

$$|R_{\Phi\bar{Z}Z\bar{Z}}| \;\simeq\; \epsilon_{z\Phi}\frac{m_z^2}{3m_{3/2}^2}, \tag{39}$$

where $\epsilon_{z\Phi}$ represents the mixing of Φ into z, and it can be approximately given by

$$\epsilon_{z\Phi} \;\simeq\; \begin{cases} |g_{\phi\bar{z}}| + |\nabla_\phi G_z|\dfrac{m_{3/2}}{m_\phi} & \text{for} \quad m_\phi \gg m_z, \\ |g_{\phi\bar{z}}|\dfrac{m_\phi^2}{m_z^2} + |\nabla_\phi G_z|\dfrac{m_{3/2}m_\phi}{m_z^2} & \text{for} \quad m_z \gg m_\phi. \end{cases} \tag{40}$$

If $m_z \sim m_{3/2}$, however, $R_{\Phi\bar{Z}Z\bar{Z}}$ is not necessarily related to m_z. On the other hand, $|\tilde{\epsilon}|$ is [e]

$$
|\tilde{\epsilon}| \;\simeq\; \begin{cases} \sqrt{3}\,|g_{\bar{\Phi}ZZ}|\dfrac{m_{3/2}}{m_\phi} & \text{for} \quad m_\phi \gg m_z, \\ \sqrt{3}\,|g_{\bar{\Phi}ZZ}|\dfrac{m_{3/2}m_\phi}{m_z^2} & \text{for} \quad m_z \gg m_\phi. \end{cases} \tag{41}
$$

In summary, $|\mathcal{G}_\Phi^{(eff)}|^2$ is given by

$$
|\mathcal{G}_\Phi^{(eff)}|^2 \;\simeq\; \left| 3\sqrt{3}\,R_{\Phi\bar{Z}Z\bar{Z}}\frac{m_{3/2}^2}{m_\phi^2}\right|^2 + \left|3\,g_{\bar{\Phi}ZZ}\,\frac{m_{3/2}}{m_\phi}\right|^2 \tag{42}
$$

for $m_\phi \gg m_z \sim m_{3/2}$,

$$
|\mathcal{G}_\Phi^{(eff)}|^2 \;\simeq\; \left|\sqrt{3}\,g_{\phi\bar{z}}\frac{m_z^2}{m_\phi^2}\right|^2 + \left|\sqrt{3}(\nabla_\phi G_z)\frac{m_{3/2}m_z^2}{m_\phi^3}\right|^2 + \left|3\,g_{\bar{\Phi}ZZ}\,\frac{m_{3/2}}{m_\phi}\right|^2 \tag{43}
$$

for $m_\phi \gg m_z \gg m_{3/2}$, and

$$
|\mathcal{G}_\Phi^{(eff)}|^2 \;\simeq\; \left|\sqrt{3}\,g_{\phi\bar{z}}\right|^2 + \left|\sqrt{3}(\nabla_\phi G_z)\frac{m_{3/2}}{m_\phi}\right|^2 + \left|3\,g_{\bar{\Phi}ZZ}\,\frac{m_{3/2}m_\phi}{m_z^2}\right|^2 \tag{44}
$$

for $m_z \gg m_\phi$ [f]. Note that, in the model basis, $|\nabla_\phi G_z| \simeq O(1)$ for a modulus field with its VEV of order M_P [1], while $|\nabla_\phi G_z| \sim \langle\phi\rangle$ for such scalar field ϕ with the Kähelr potential $K = |\phi|^2 + \cdots$ before expanding around the VEV [3, 4]. Therefore, the second term in Eq. (44) reproduces the partial decay rate of $\tilde{\Phi}$ into a pair of the gravitinos in Refs. [1, 2, 3, 4]. In addition, even in the case of $m_\phi \gg m_z$, the rate becomes sizable if $g_{\bar{\Phi}ZZ}$ is order unity.

2. SM (s)particles

In the gravity mediation, there are non-renormalizable interactions between the SUSY breaking field z and the SM sector to induce the soft SUSY breaking terms. For instance, the gaugino masses are obtained in the model frame by

$$
\mathcal{L} \;=\; \int d^2\theta\, c_z \frac{z}{M_P}\, W^{(a)}W^{(a)} + h.c. \tag{45}
$$

where $W^{(a)}$ is the supersymmetric field strength of the gauge field, and c_z is a coupling constant of order unity. The mixture between the heavy scalar field ϕ and z makes it

[e] Note that $\tilde{\epsilon}$ is at least $O(\langle\phi\rangle\, m_{3/2}^2/m_\phi^2)$ even in the case of the minimal Kähler potential.

[f] We are grateful to M. Ibe and Y. Shinbara for pointing out the 1st term in Eq.(44).

possible for $\tilde{\Phi}$ in the mass eigenstate to decay into the SM (s)particles through the above coupling. Using (29), the interaction between $\tilde{\Phi}$ and the SM (s)particles is given by

$$\mathcal{L}^{(\mathrm{mix})}_{\tilde{\Phi}WW} \sim \frac{c_z}{M_P}\epsilon_{z\tilde{\Phi}}\tilde{\Phi}\int d^2\theta\, W^{(a)}W^{(a)} + h.c. \tag{46}$$

which leads to

$$\Gamma^{(\mathrm{mix})}(\tilde{\Phi}\to gauge\ boson) \simeq \Gamma^{(\mathrm{mix})}(\tilde{\Phi}\to gaugino) \simeq \frac{3}{2\pi}\left(\frac{N_g}{12}\right)\epsilon^2_{z\tilde{\Phi}}|c_z|^2\frac{m_\phi^3}{M_P^2} \tag{47}$$

for $m_\phi \gg m_{3/2}$, where N_g is the number of final states, and $N_g = 12$ for $\mathrm{SU}(3)_c\times\mathrm{SU}(2)_L\times\mathrm{U}(1)_Y$. We notice that the decay rate of the gaugino production is comparable to that of the gauge boson [1, 2]. Note that this decay is always present as far as there is a mixing between the ϕ and z. As we will see, it will become important especially for the inflaton decay.

In the case of modulus decay, it also has a direct coupling to the SM sector, such as the dilatonic coupling with the gauge supermultiplet, $\mathcal{L} = \frac{\lambda_G}{M_P}\int d^2\theta\,\phi W^{(a)}W^{(a)} + \mathrm{h.c.}$. The decay rate through this coupling is given by

$$\Gamma^{(\mathrm{direct})}(\tilde{\Phi}\to gauge\ boson) \simeq \Gamma^{(\mathrm{direct})}(\tilde{\Phi}\to gaugino) \simeq \frac{3}{2\pi}\left(\frac{N_g}{12}\right)|\lambda_G|^2\frac{m_\phi^3}{M_P^2} \tag{48}$$

for $m_\phi \gg m_{3/2}$. In the gravity mediation, therefore, the direct decay of the modulus into the SM (s)particles becomes dominant over that through mixings, as long as $|\lambda_G| \sim 1$ and $\epsilon_{z\tilde{\Phi}} < 1$. Note that in the case of inflaton, it does not necessarily have the above direct coupling.

3. SUSY breaking sector

The heavy scalar can also decay into the hidden sector, which includes the SUSY breaking fields. Due to the mixing between the fields ϕ and $z(z^\dagger)$ in the model frame, the mass eigenstate $\tilde{\Phi}$ has a branch of the production of the hidden sector field $\tilde{Z}$, if kinematically allowed. In this subsection we discuss the decay $\tilde{\Phi}\to\tilde{Z}$ assuming $m_\phi \gg m_z$.

A possible interaction between ϕ and z comes from the Kähler potential, $K = g_{\phi\bar{z}z}\phi z^\dagger z + \mathrm{h.c.}$. Actually such an interaction is plausible once we consider an operator, $K = |\phi|^2|z|^2/M_P^2$, with taking account of the VEV of ϕ. The decay rate via this operator is however suppressed. If the decay is induced by the $D = 5$ operator in the Kähler potential, and if the final states have opposite chirality, the relevant coupling constant becomes proportional to the mass squared of the final state, m_Z^2, or that of the gravitino,$m_{3/2}^2$. Therefore the resultant decay becomes suppressed: $\Gamma \sim \mathrm{Max}[m_{3/2}^4, m_z^4]/m_\phi M_P^2$.

In the Kähler potential, there is another $D = 5$ operator, $K = g_{\bar{\phi}zz}\phi^\dagger zz + \mathrm{h.c.}$, which is different from the above one: the final states have the same chirality. This operator induces a larger decay rate:

$$\Gamma(\tilde{\Phi} \to \tilde{Z}\tilde{Z}) \;\simeq\; \frac{|g_{\bar{\phi}zz}|^2}{8\pi}\frac{m_\phi^3}{M_P^2}. \tag{49}$$

As discussed in the previous subsection, the Kähler mixing $g_{\bar{\phi}zz}$ also induces the decay into the gravitino [cf. (38), (42), and (43)]. It should be noted that these two decay rates can be correlated. As long as $|\mathcal{G}_\Phi^{(eff)}|$ is dominated by the last term in (42) or (43) that contains $|g_{\bar{\Phi}ZZ}|$, they become comparable: $\Gamma(\tilde{\Phi} \to 2\psi_{3/2})/\Gamma(\tilde{\Phi} \to 2\tilde{Z}) \simeq 1/4$.

It is stressed that production of the fermionic component of the z field is very different. This is because a combination of the fermionic components of ϕ and the SUSY breaking fields is absorbed into the gravitino as a goldstino. In the minimum setup, namely where there is a single SUSY breaking field, the fermionic component of z almost behaves as the goldstino, and that of ϕ provides the remnant massive degrees of freedom. Therefore, when the ϕ mass is given by the supersymmetric term, $\nabla_\phi G_\phi$, the mass of this massive fermion becomes close to m_ϕ, and hence the decay is kinematically suppressed or forbidden.

The produced $\tilde{Z}$ subsequently decays into the visible sector, and into the gravitino if kinematically allowed. The decay of the $\tilde{Z}$ field and its implications will be discussed in the following sections.

B. Modulus

In this and the next sections, we discuss how the decay via mixings with the SUSY breaking field affect the cosmological scenarios. We concentrate especially on the modulus and the inflaton, and see how disastrous the cosmological scenarios become due to such mixings.

Let us start from the modulus decay. We discuss two distinct cases $m_z > m_\phi$ and $m_\phi < m_z$ in turn. In both cases, the dominant decay channel is that into the SM (s)particles, whose rate is given by Eq. (48). A successful big bang nucleosynthesis (BBN) requires a temperature higher than ~ 5 MeV [12, 13], which leads to a lower bound on the modulus mass, $m_\phi \gtrsim 100$ TeV, and we assume this is the case.

Here, we should mention that there may be another cosmological moduli problem associated with the SUSY breaking field z. In this section we assume that it is the heavy ϕ field which dominates the energy density of the universe and causes the final reheating (before the BBN), and we will briefly discuss the z modulus problem in Sec. V.

1. $m_z > m_\phi$

In this case, the modulus field ϕ decays into the SM (s)particles and the gravitino, with the partial decay rates in Eq. (48) and Eq. (38), respectively. The branching ratio of the gravitino production then becomes

$$B_{3/2} = Br(\tilde{\Phi} \to 2\psi_{3/2}) \;=\; \frac{1}{72 N_g |\lambda_G|^2} \frac{m_\phi^2}{m_{3/2}^2} |\mathcal{G}_\Phi^{(eff)}|^2 \tag{50}$$

As can be seen in Eqs. (42) - (44), the coupling $|\mathcal{G}_\Phi^{(eff)}|$ depends on the model. If $m_z > m_\phi$, however, $|\mathcal{G}_\Phi^{(eff)}|$ is suppressed only by a single power of the gravitino mass and the branching ratio becomes

$$B_{3/2} \;\simeq\; \frac{1}{24 N_g |\lambda_G|^2} \left(|\nabla_\phi G_z|^2 + |g_{\phi\bar{z}}|^2 \frac{m_\phi^2}{m_{3/2}^2} \right). \tag{51}$$

Using $|\nabla_\phi G_z| \sim O(1)$ in the model basis [1], the first term is the same order as the one estimated in Refs. [1, 2]. As was shown there, such a large branching fraction of the gravitino production causes serious cosmological problems. The second term makes the problem even worse if $|g_{\phi\bar{z}}| \gg m_{3/2}/m_\phi$.

2. $m_\phi > m_z$

Now one has to consider a new decay mode, $\tilde{\Phi} \to 2\tilde{Z}$, in addition to the channels discussed above. As discussed in Sec.III A 3, if $g_{\bar{\phi}zz}$ is sizable, the $\tilde{\Phi}$ produces roughly as many $\tilde{Z}$ as the gravitino. Here we discuss the fate of the produced $\tilde{Z}$ and its implications.

If $m_z \lesssim 2\, m_{3/2}$, the $\tilde{Z}$ dominantly decays into the visible sector via the interaction Eq. (45), which leads to $\Gamma(\tilde{Z} \to visible) \sim m_z^3/M_P^2$. Note that this rate is comparable to the decay rate of the gravitino for $m_z \sim m_{3/2}$. Therefore it causes qualitatively similar problems as the gravitino, such as changing light element abundances and producing too many LSPs [1, 2]. The details of the constraint on the model depends on the mass and couplings of the z field.

If on the other hand $m_z \gg m_{3/2}$, due to SUSY breaking mass term, $\tilde{Z}$ dominantly decays into the gravitino. (Note that $|G_Z| \simeq \sqrt{3}$ and the rate is enhanced by $(m_z/m_{3/2})^2$. cf. Eq. (38).) Recall that there are gravitinos directly produced by the Φ decay. The net effect is therefore just an enhancement of the gravitino abundance by an order one factor. The subsequent gravitino decay is subject to the cosmological constraints [1, 2].

To summarize, $\tilde{\Phi}$ produces roughly as many $\tilde{Z}$ as the gravitinos, and the produced $\tilde{Z}$ will cause a similar problem as the gravitino does.

C. Inflaton

We now turn to discuss the inflaton decay. We assume that the SUSY breaking field z is lighter than the inflaton ϕ [g]. Therefore, the inflaton can decay into the SM (s)particles, gravitinos, and the SUSY breaking sector fields. The importance of the inflaton decay into the gravitino has been recently pointed out in Ref. [3].

Let us first consider the inflaton decay into the SM (s)particles through the interaction (46). The mixing with the SUSY breaking field may enhance the decay rate of the inflaton, which leads to a higher reheating temperature, T_R. Since T_R is bounded from above due to the abundance of the gravitinos produced by thermal scatterings, such mixing must be small enough.

The presence of the interaction (46) sets a lower bound on the reheating temperature:

$$T_R \gtrsim 3 \times 10^8 \text{ GeV} \, \epsilon_{z\tilde{\Phi}} |c_z| \left(\frac{m_\phi}{10^{12}\text{GeV}} \right)^{\frac{3}{2}} \tag{52}$$

where we have used $N_g = 12$ for the SM gauge groups and the relativistic degrees of freedom $g_* \simeq 200$. For $m_{3/2} \simeq O(0.1 - 1\,\text{TeV})$, the bound from the gravitino problem reads $T_R < O(10^6 - 10^8)\,\text{GeV}$ [14, 15], where the upper bound depends on the gravitino mass and the hadronic branching ratio B_h. Combining this with (52), we obtain

$$\epsilon_{z\tilde{\Phi}} \lesssim (3 \times 10^{-3} - 0.3)\, c_z^{-1} \left(\frac{m_\phi}{10^{12}\text{GeV}} \right)^{-\frac{3}{2}}. \tag{53}$$

[g] Note, however, that this may not be the case in the new inflation models [3].

The heavier the inflaton mass is, the severer this bound becomes. For the new inflation model [16, 17], the inflaton mass is relatively small, $m_\phi \sim 10^{10}$ GeV, and therefore the bound (53) does not give any sensible constraint on the mixing. For the hybrid inflation models [18, 19, 20], however, the inflaton mass can be very large, $m_\phi \sim O(10^{11} - 10^{15})$ GeV [h] Then we obtain $\epsilon_{z\tilde{\Phi}} \lesssim O(10^{-7} - 10)$. To translate this bound into that on parameters in the model basis, let us estimate $\epsilon_{z\tilde{\Phi}}$,

$$\epsilon_{z\tilde{\Phi}} \simeq \mathrm{Max}\left[|g_{\phi\bar{z}}|, \sqrt{3}\phi_0 \frac{m_{3/2}}{m_\phi}, \sqrt{3}|g_{\bar{\phi}zz}|\frac{m_{3/2}}{m_\phi}\right], \tag{54}$$

where we assumed $m_\phi \gg m_z$. Since the second and third terms are highly suppressed due to the ratio of the gravitino mass to the inflaton mass, the bound on $\epsilon_{z\tilde{\Phi}}$ is effectively that on $|g_{\phi\bar{z}}|$. In the case of the hybrid inflation model, therefore, we obtain a nontrivial bound, $|g_{\phi\bar{z}}| \lesssim O(10^{-7} - 10)$.

To see how severe the bound on the mixing is, it is necessary to consider explicit interactions in the Kähler potential [i]. Let us consider the following interactions in the model basis before expanding the fields around their VEVs,

$$\delta K = k_1|\phi|^2(z + z^\dagger) + k_2|\phi|^2|z|^2 + \frac{k_3}{2}|\phi|^2(zz + z^\dagger z^\dagger)\cdots, \tag{55}$$

where k_i ($i = 1, 2, 3$) are numerical coefficients, and we have dropped several terms like $\phi^2(z + z^\dagger)$, assuming that ϕ is charged under some symmetry. As long as z is a singlet, all the coefficients are expected to be order unity. Then $g_{\phi\bar{z}}$ is non-vanishing if ϕ and z take non-zero VEVs,

$$g_{\phi\bar{z}} = k_1\phi_0^* + k_2\phi_0^* z_0 + k_3\phi_0^* z_0^*. \tag{56}$$

Therefore the constraint on $g_{\phi\bar{z}}$ can be interpreted as that on the numerical coefficients k_i, which are otherwise unconstrained from any symmetries of ϕ. If k_i is severely constrained from cosmological considerations, it indicates either that there is still unknown symmetry or mechanism to suppress the couplings, or that such inflation model with vanishing ϕ_0 is favored. As an example, let us take the hybrid inflation model with $\phi_0 \sim 10^{-3}$. Then $|g_{\phi\bar{z}}| \lesssim O(10^{-7} - 10)$ can be rephrased as $|k_1 + k_2 z_0 + k_3 z_0^*| \lesssim O(10^{-4} - 10^4)$. Therefore, a

[h] The hybrid inflation models contain two types of the fields: the inflaton field and the waterfall fields. Although the bound on $|g_{\phi\bar{z}}|$ applies to both fields, we identify ϕ with the waterfall field when we substitute the VEV of ϕ into (56).

[i] We assume here that ϕ and z are not coupled in the superpotential for simplicity.

considerable part of the parameter spaces are disfavored if $B_h \simeq 1$. (Note, however, that the hybrid inflation model is already disfavored only from the direct gravitino production [3].)

Let us now consider the inflaton decay into the gravitinos. Recently, it was pointed out in Ref. [3] that the gravitino production from the inflaton decay can put a severe constraint on the inflation models (in particular, the hybrid inflation model is excluded unless the higher order terms in the Kähler potential is extremely suppressed). The decay rate into a pair of the gravitinos is given by (38). The constraint on the inflation models can be read from Fig. 1 in Ref. [3] by replacing G_ϕ with $|\mathcal{G}_\Phi^{(eff)}| \simeq |3 g_{\bar{\Phi}ZZ}\, m_{3/2}/m_\phi|$. Thus, we can rephrase the results of Ref. [3] that the hybrid inflation model is excluded unless $|g_{\bar{\Phi}ZZ}|$ is extremely small in the gravity mediation.

Lastly, let us consider the inflaton decay into the SUSY breaking sector. As in the case of the moduli, the inflaton decay into $\tilde{Z}$ is always concomitant with almost same amount of the gravitino production, since the both production rates are proportional to $|g_{\bar{\Phi}ZZ}|^2$. Therefore the produced $\tilde{Z}$ only causes a problem which is at most as severe as the gravitino overproduction problem.

IV. LOW ENERGY SUSY BREAKING MODELS

In this section we consider low energy SUSY breaking models, as represented by the GMSB models. Compared to the gravity mediation, there are two major differences. One is that the SUSY breaking field couples to the visible sector more strongly, which is a general feature of the low energy SUSY breaking models. This enhances the decay rate of $\tilde{\Phi}$ due to the mixing. The other is the existence of the messenger sector fields, which is characteristic to the GMSB models. Since the messenger sector contains another scalar field, s, we need to consider the scalar mixings of both $\phi - z$ and $\phi - s$.

In the messenger sector, there is a chiral superfield, s, with nonzero VEVs of the scalar and auxiliary components, which couples to the messengers Ψ_M and $\bar{\Psi}_M$ by

$$W \;=\; y_M\, s\, \Psi_M \bar{\Psi}_M, \tag{57}$$

where y_M is a coupling constant. The scalar VEV, $M_s \equiv \langle s \rangle$, sets the messenger mass scale, $M_{\rm mess} \equiv y_M M_s$, while the F-term, F_s, provides the mass splitting between the messenger fermions and bosons, $\sim y_M F_s$. The SUSY breaking is transmitted radiatively to the visible sector by the SM gauge interactions, under which Ψ_M and $\bar{\Psi}_M$ are charged. The SUSY breaking scale in the visible sector is therefore determined by F_s/M_s. For example, the gaugino mass is induced by [j]

$$\mathcal{L} = \int d^2\theta \, \frac{\alpha}{4\pi} \frac{s}{M_s} W^{(a)} W^{(a)} + h.c., \tag{58}$$

where α denotes the gauge coupling, and we have assumed the messenger index $N = 1$. The gaugino mass M_λ is therefore given by

$$M_\lambda = \frac{\alpha}{4\pi} \frac{F_s}{M_s}. \tag{59}$$

By using $G_s = F_s/(m_{3/2} M_P)$, we can relate M_s to $m_{3/2}$:

$$m_{3/2} \simeq 9 \times 10^{-4}\,\mathrm{GeV}\, G_s^{-1} \left(\frac{m_{\tilde{g}}}{1\,\mathrm{TeV}}\right) \left(\frac{M_s}{10^{10}\,\mathrm{GeV}}\right), \tag{60}$$

where the gluino mass $m_{\tilde{g}}$ is determined at the messenger scale. In contrast to the gravity mediation, it is nontrivial (and therefore model-dependent) how large G_s is. In fact, in the direct gauge-mediation scenario, $|G_s| \sim 1$ if s is identified with z, while $|G_s| \ll 1$ in such models that the SUSY breaking effects is radiatively transmitted from a secluded sector (that contains z) to the messenger sector. If $|G_s| \sim 1$, there is no significant difference between s and z. If $|G_s| \ll 1$, we need to consider the mixings $\phi - z$ and $\phi - s$, separately (for simplicity we neglect the mixing between z and s). The formalism developed in Sec. II can also be applied to the $\phi - s$ mixing. Since $|G_z| \sim \sqrt{3}$, it is the mixing with z that determines the decay of ϕ into the gravitinos. On the other hand, it is s that determines the decay into the SM (s)particles, since s (not z) couples to the SM (s)particles via the messengers Ψ_M and $\bar{\Psi}_M$. Lastly ϕ may decay into both s and z via the mixings. Thus, although there are two SUSY breaking fields s and z in the GMSB models, we can similarly discuss the decay of ϕ as we did in Sec. III.

A. Decay Modes

The decay channels of the heavy scalar $\tilde{\Phi}$ are quite similar to those in the gravity mediation. In particular, the gravitino production rate is independent of the couplings between

[j] Note that such an interaction as (58) always exists in the low energy SUSY breaking models, even if the messenger sector does not exist. In this case M_s simply parametrizes the strength of the interaction between s and the visible sector.

the SUSY breaking field and the visible sector. In addition, the decays of $\tilde{\Phi}$ into z and s (if kinematically allowed) are also similar to the gravity mediation case, i.e., they are dominated by the decays induced by the higher order couplings in the Kähler potential $g_{\bar{\phi}zz}$ and $g_{\bar{\phi}ss}$, respectively, if these couplings are sizable, and otherwise suppressed. Here, we focus on the new features of the low energy SUSY breaking scenario.

When the SUSY breaking effects are mediated to the visible sector with the interactions with a lower cutoff, M_s, the mixing-induced $\tilde{\Phi}$ decay into the visible sector depends on M_s rather than M_P. Since the field s couples with the visible sector, the decay rate is evaluated from the operator Eq. (58) as

$$\Gamma^{(\mathrm{mix})}(\tilde{\Phi} \to gauge\ boson) \ \simeq\ \Gamma^{(\mathrm{mix})}(\tilde{\Phi} \to gaugino) \ \simeq\ \frac{\alpha_s^2}{16\pi^3}\left(\frac{N_g}{8}\right)\epsilon_{s\tilde{\Phi}}^2 \frac{m_\phi^3}{M_s^2}, \tag{61}$$

where $\epsilon_{s\tilde{\Phi}}$ is defined as Eq. (29) with $z \to s$, and we assumed the decay is dominated by the gluon/gluino production. The decay rate thus depends on the mixing, $\epsilon_{s\tilde{\Phi}}$. If the mixing is dominated by the non-renormalizable term in the Kähler potential, $K = |\phi|^2|s|^2/M_P^2$, $\epsilon_{s\tilde{\Phi}}$ is given by $\sim M_s/M_P$. Then the mixing-induced decay rate becomes $\Gamma \sim m_\phi^3/M_P^2$.

In the GMSB setups, there exist the messenger fields, Ψ_M and $\bar{\Psi}_M$. Since they interacts with the s field by the renormalizable coupling, Eq. (57), the heavy scalar field, $\tilde{\Phi}$, can decay into the fermionic component of Ψ_M and $\bar{\Psi}_M$ rapidly as long as the channel is kinematically allowed. From Eq. (57), the decay rate is estimated as

$$\Gamma(\tilde{\Phi} \to \psi_\Psi \psi_{\bar{\Psi}}) \ \simeq\ N_{\mathrm{mess}} \frac{|y_M \epsilon_{s\tilde{\Phi}}|^2}{32\pi} m_\phi, \tag{62}$$

where N_{mess} is a number of the possible final states, for instance, $N_{\mathrm{mess}} = 5$ when $\Psi_M + \bar{\Psi}_M$ are charged as $\mathbf{5} + \bar{\mathbf{5}}$ under $SU(5)_{\mathrm{GUT}}$. Therefore unless y_M and/or $\epsilon_{s\tilde{\Phi}}$ is extremely suppressed, the dominant channel of $\tilde{\Phi}$ becomes the production of the messenger fermion.

B. Modulus

The modulus decay in the low energy SUSY breaking scenario is similar to that in the gravity mediation, as long as $\epsilon_{s\tilde{\Phi}} \lesssim M_s/M_P$. This is the case if the mixing mainly comes from e.g., $\delta K = \kappa\, |\phi|^2|s|^2/M_P^2$. The modulus decay into the SM (s)particles via the mixing with s then proceeds with the rate (61) that is at most comparable to (48). Therefore a successful BBN requires $m_\phi \gtrsim 100$ TeV as in the case of gravity mediation. In the general low energy SUSY breaking models, however, s can be a singlet field and therefore such an interaction as $\delta K = |\phi|^2(s+s^\dagger)/M_P$ may exist [k]. In this case, the modulus decay into the visible sector via the mixing can exceed the rate (48), and the modulus may decay before the BBN even if its mass m_ϕ is smaller than 100 TeV. Although this may relax the moduli problem in the low energy SUSY breaking models, it strongly depends on the nature of s whether such an interaction exists at all.

Actually, there may be a cosmological moduli problem associated with the z and/or s fields. Here, we discuss the case where the universe is dominated and reheated by the heavy field ϕ, and leave the other cases for discussion in Sec. V.

The gravitino production occurs via the mixing of ϕ with z (and s if $G_s \sim O(1)$ as in the direct gauge mediation), and the decay rate is given by (38). The cosmological constraints on the stable gravitinos from the modulus decay are as given in Ref. [1].

If kinematically allowed, and if $g_{\bar{\phi}zz}$ and $g_{\bar{\phi}ss}$ are non-vanishing, the modulus decays into z and s. Although the masses of z and s are considered to be comparable or larger than $m_{3/2}$, they are model-dependent. So here we take m_z and m_s as free parameters. The abundance of z is the same order of the gravitino abundance if $g_{\bar{\phi}zz}$ is sizable, just as the case in the gravity mediation. On the other hand, unless $G_s \sim O(1)$, the s abundance is not necessarily correlated with the gravitinos. For $m_z > 2m_{3/2}$, the produced z can decay into a pair of gravitinos, and the rate is enhanced for larger m_z. The z field may decay into the SM (s)particles, if it has a direct coupling to the visible sector. The strength of the coupling must be such that the F-term of z does not give dominant contributions to the soft masses. In addition, it is possible that z has relatively strong interactions with s and decays into s. Then we only have to consider the decay processes of s, which comes both directly from the modulus decay and through the decay of z. The interaction of s with the visible sector is given by (58). Assuming that s dominantly decays into two gluons, the decay temperature is given by

$$T_d^{(s)} \simeq 0.04\,\mathrm{GeV}\left(\frac{m_s}{10\,\mathrm{GeV}}\right)^{\frac{3}{2}}\left(\frac{M_s}{10^{10}\,\mathrm{GeV}}\right)^{-1}, \tag{63}$$

where we take $g_* = 10.75$. To be conservative, we require that s decays before the BBN

[k] Note that this coupling enhances the gravitino production rate due to the first term in Eq. (44) if G_s is sizable.

starts, i.e., $T_d^{(s)} \gtrsim 5$ MeV [12, 13]. Then the mass of s must satisfy

$$m_s \gtrsim 3\,\mathrm{GeV}\left(\frac{M_s}{10^{10}\,\mathrm{GeV}}\right)^{\frac{2}{3}}. \tag{64}$$

It should be noted however that, even if this inequality is satisfied, s may produce the too many LSPs and/or gravitinos, if kinematically allowed.

Lastly let us comment on the modulus decay into the messengers. Although there exist the messenger fields in the GMSB models, it is unlikely that the modulus decays into them, since the messenger scale $M_{\rm mess}$ is typically larger than the modulus mass.

C. Inflaton

The low energy SUSY breaking models may contain the messenger sector as in the GMSB models. If the messenger scale $M_{\rm mess} = y_M M_s$ is smaller than the inflaton mass m_ϕ, the inflaton can decay into the messenger sector as well via the $\phi - s$ mixing. In the following we discuss the cases with and without such a channel separately.

1. *Decay into visible sector, the gravitinos, s and z*

Let us first consider the case without the decay into the messenger sector. The inflaton then decays into the SM (s)particles, the gravitinos, s and z. In the following we assume $m_\phi \gg m_s,\, m_z$.

The decay into the SM (s)particles may proceed via the mixing with s. The interaction (58) sets a lower bound on the reheating temperature:

$$T_R \gtrsim 2\times 10^{11}\ \mathrm{GeV}\,\frac{\epsilon_{s\tilde{\Phi}}}{|G_s|}\left(\frac{m_{\tilde{g}}}{1\,\mathrm{TeV}}\right)\left(\frac{m_{3/2}}{1\,\mathrm{GeV}}\right)^{-1}\left(\frac{m_\phi}{10^{12}\,\mathrm{GeV}}\right)^{\frac{3}{2}}, \tag{65}$$

where we set $N_g = 8$ and $g_* \simeq 200$. In the low energy SUSY breaking models, the upper bound on T_R is given by [21, 22]

$$T_R \lesssim 5\times 10^7\ \mathrm{GeV}\left(\frac{m_{\tilde{g}}}{1\,\mathrm{TeV}}\right)^{-2}\left(\frac{m_{3/2}}{1\,\mathrm{GeV}}\right), \tag{66}$$

for $m_{3/2} = 10^{-4} - 10$ GeV, and

$$T_R \lesssim O(100)\mathrm{GeV}, \tag{67}$$

for 1 keV $\lesssim m_{3/2} \lesssim 10^{-4}$ GeV, in order for the gravitino abundance not to exceed the dark matter abundance. Here and in what follows we neglect the difference of the values of $m_{\tilde{g}}$ at the messenger scale and at the reheating temperature. In the GMSB model, the assumption $M_{\rm mess} = y_M M_s > m_\phi$ sets a lower limit on the gravitino mass. Although the inflaton mass m_ϕ strongly depends on the inflation models, it is typically larger than $O(10^9)$ GeV. Using (60), the gravitino mass should be larger than $O(10^{-4})$ GeV in this case. For the low energy SUSY breaking models without the messenger sector, such a lower limit is not necessarily applied. Combining (65) with (66) or (67), we obtain the severe bound on the mixing:

$$\epsilon_{s\tilde{\Phi}} \lesssim 3\times 10^{-10}|G_s|\left(\frac{m_{\tilde{g}}}{1\,{\rm TeV}}\right)^{-3}\left(\frac{m_{3/2}}{1\,{\rm MeV}}\right)^{2}\left(\frac{m_\phi}{10^{12}\,{\rm GeV}}\right)^{-\frac{3}{2}}, \tag{68}$$

for $m_{3/2} = 10^{-4} - 10$ GeV and

$$\epsilon_{s\tilde{\Phi}} \lesssim O(10^{-15})|G_s|\left(\frac{m_{\tilde{g}}}{1\,{\rm TeV}}\right)^{-1}\left(\frac{m_{3/2}}{10^{-5}\,{\rm GeV}}\right)\left(\frac{m_\phi}{10^{12}\,{\rm GeV}}\right)^{-\frac{3}{2}}, \tag{69}$$

for 1 keV $\lesssim m_{3/2} \lesssim 10^{-4}$ GeV. The bounds becomes severer for smaller G_s, larger m_ϕ, and smaller $m_{3/2}$.

To exemplify how severe the bound is, let us rewrite this bound to that on coefficients of the higher order interactions in the Kähler potential. We consider the following interactions in the model basis before expanding around the VEVs:

$$\delta K \;=\; k_1|\phi|^2(s+s^\dagger) + k_2|\phi|^2|s|^2, \tag{70}$$

where k_1 and k_2 are numerical coefficients. The first term can be forbidden if s has some symmetries (i.e., $k_1 = 0$), but we include it here to see how severely such an interaction is constrained. On the other hand, k_2 is unconstrained from any symmetries, so it is expected to be order unity. This Kähler potential leads to

$$g_{\phi\bar{s}} \;=\; k_1\,\phi_0^* + k_2\,\phi_0^* M_s, \tag{71}$$

where we have taken the VEV of s real, for simplicity. Barring cancellations, we obtain the constraints on k_1 and k_2 from (68) and (69), since $\epsilon_{s\tilde{\Phi}}$ is roughly equal to $|g_{\phi\bar{s}}|$ for $m_\phi \gg m_s$:

$$\begin{aligned}
|k_1| \;&\lesssim\; 8\times 10^{-7}\,|G_s|\left(\frac{m_{\tilde{g}}}{1\,{\rm TeV}}\right)^{-3}\left(\frac{m_{3/2}}{1\,{\rm MeV}}\right)^{2}\left(\frac{m_\phi}{10^{12}\,{\rm GeV}}\right)^{-\frac{3}{2}}\left(\frac{|\phi_0|}{10^{15}\,{\rm GeV}}\right)^{-1}, \\
|k_2| \;&\lesssim\; 2\times 10^{2}\left(\frac{m_{\tilde{g}}}{1\,{\rm TeV}}\right)^{-2}\left(\frac{m_{3/2}}{1\,{\rm MeV}}\right)\left(\frac{m_\phi}{10^{12}\,{\rm GeV}}\right)^{-\frac{3}{2}}\left(\frac{|\phi_0|}{10^{15}\,{\rm GeV}}\right)^{-1},
\end{aligned} \tag{72}$$

for $m_{3/2} = 10^{-4} - 10$ GeV and

$$\begin{aligned}
|k_1| &\lesssim O(10^{-9})|G_s| \left(\frac{m_{\tilde{g}}}{1\,\mathrm{TeV}}\right)^{-1} \left(\frac{m_{3/2}}{1\,\mathrm{MeV}}\right) \left(\frac{m_\phi}{10^{12}\,\mathrm{GeV}}\right)^{-\frac{3}{2}} \left(\frac{|\phi_0|}{10^{15}\,\mathrm{GeV}}\right)^{-1}, \\
|k_2| &\lesssim O(0.1) \left(\frac{m_\phi}{10^{12}\,\mathrm{GeV}}\right)^{-\frac{3}{2}} \left(\frac{|\phi_0|}{10^{15}\,\mathrm{GeV}}\right)^{-1},
\end{aligned} \tag{73}$$

for 1 keV $\lesssim m_{3/2} \lesssim 10^{-4}$ GeV, where we have used (60). The bounds become severer for larger inflaton VEV and mass. It should be noted that the coefficient k_1 is tightly constrained, which strongly disfavors the existence of a singlet field s in the low energy SUSY breaking model. In other words, the s field must be charged under some symmetry (e.g. $U(1)$ symmetry) to forbid such interaction as $\sim |\phi|^2(s + s^\dagger)$. Let us now focus on the constraint on k_2, which is generically unsuppressed. As an example, let us consider the hybrid inflation model. For $m_{3/2} > 10^{-4}$ GeV, only some fraction of the parameter space is disfavored, while fairly wider ranges of the model parameters require fine-tunings on the higher order interactions in the Kähler potential for $m_{3/2} < 10^{-4}$ GeV.

Next let us consider the $\tilde{\Phi}$ decay into the gravitinos. Since the decay is induced by the mixing with z with $G_z \sim 1$, the gravitino overproduction problem is similar to that considered in the previous section. The only difference is the upper bound on T_R from the abundance of the gravitinos produced by thermal scatterings [cf. (66) and (67)]. Although the upper bound on T_R depends on $m_{3/2}$ for $m_{3/2} > 10^{-4}$ GeV, that on the abundance of the gravitinos directly produced by the inflaton decay does not depend on $m_{3/2}$. Therefore the gravitino overproduction problem sets a more or less similar bounds on the inflation models given in Ref. [3].

Lastly let us consider the decay of $\tilde{\Phi}$ into the SUSY breaking sector, s and z. As discussed in Sec. III, the decay rate into z is comparable to the gravitino production. However, in contrast to the gravity mediation, z may have relatively strong coupling with the messenger sector or the visible sector. If this coupling is so strong that z decays mainly into the visible sector before the BBN, z may not be cosmologically problematic. Even if the coupling is weak, z decays into the gravitinos as far as $m_z > 2m_{3/2}$, and it only increases the gravitino abundance by $O(1)$ factor. Although s is also produced from the $\tilde{\Phi}$ decay if $g_{\bar{\phi}ss}$ is sizable, it does not cause any cosmological difficulties if the inequality (64) is satisfied.

2. Decay into the messenger sector

If there is a messenger sector as in the GMSB scenario, and if the messenger scale $M_{\rm mess} = y_M M_s$ is smaller than the inflaton mass, the inflaton can decay directly into the messenger sector. Indeed, such a decay may make the reheating temperature of the inflation even higher than that discussed in the previous subsection. Using the decay rate (62), the reheating temperature becomes

$$T_R \gtrsim 2\times 10^{14}\,{\rm GeV}\,|y_M\,\epsilon_{s\tilde\Phi}|\left(\frac{m_\phi}{10^{12}\,{\rm GeV}}\right)^{\frac{1}{2}}, \tag{74}$$

where we set $N_{\rm mess}=5$. Combined with (66) or (67), we obtain

$$\epsilon_{s\tilde\Phi} \;\lesssim\; 3\times 10^{-10}|y_M|^{-1}\left(\frac{m_{\tilde g}}{1\,{\rm TeV}}\right)^{-2}\left(\frac{m_{3/2}}{1\,{\rm MeV}}\right)\left(\frac{m_\phi}{10^{12}\,{\rm GeV}}\right)^{-\frac{1}{2}} \tag{75}$$

for $m_{3/2} = 10^{-4} - 10$ GeV and

$$\epsilon_{s\tilde\Phi} \;\lesssim\; O(10^{-13})|y_M|^{-1}\left(\frac{m_\phi}{10^{12}\,{\rm GeV}}\right)^{-\frac{1}{2}} \tag{76}$$

for 1 keV $\lesssim m_{3/2} \lesssim 10^{-4}$ GeV.

Assuming that the mixing is provided by the interaction (70), we can rewrite the above bounds as

$$\begin{aligned} |k_1| \;&\lesssim\; 8\times 10^{-7}\,|y_M|^{-1}\left(\frac{m_{\tilde g}}{1\,{\rm TeV}}\right)^{-2}\left(\frac{m_{3/2}}{1\,{\rm MeV}}\right)\left(\frac{m_\phi}{10^{12}\,{\rm GeV}}\right)^{-\frac{1}{2}}\left(\frac{|\phi_0|}{10^{15}\,{\rm GeV}}\right)^{-1}, \\ |k_2| \;&\lesssim\; 2\times 10^{2}\,\left(\frac{m_{\tilde g}}{1\,{\rm TeV}}\right)^{-2}\left(\frac{m_{3/2}}{1\,{\rm MeV}}\right)\left(\frac{M_{\rm mess}}{10^{10}\,{\rm GeV}}\right)^{-1}\left(\frac{m_\phi}{10^{12}\,{\rm GeV}}\right)^{-\frac{1}{2}}\left(\frac{|\phi_0|}{10^{15}\,{\rm GeV}}\right)^{-1}, \end{aligned} \tag{77}$$

for $m_{3/2} = 10^{-4} - 10$ GeV and

$$\begin{aligned} |k_1| \;&\lesssim\; O(10^{-9})\left(\frac{m_\phi}{10^{12}\,{\rm GeV}}\right)^{-\frac{1}{2}}\left(\frac{|\phi_0|}{10^{15}\,{\rm GeV}}\right)^{-1}, \\ |k_2| \;&\lesssim\; O(0.1)\left(\frac{M_{\rm mess}}{10^{10}\,{\rm GeV}}\right)^{-1}\left(\frac{m_\phi}{10^{12}\,{\rm GeV}}\right)^{-\frac{1}{2}}\left(\frac{|\phi_0|}{10^{15}\,{\rm GeV}}\right)^{-1}, \end{aligned} \tag{78}$$

for 1 keV $\lesssim m_{3/2} \lesssim 10^{-4}$ GeV. Thus k_1 is severely constrained as before, which indicates that k_1 must be vanishing due to a symmetry. The constraint on k_2 depends on the messenger scale $M_{\rm mess} = |y_M|M_s$; it becomes severer for larger $M_{\rm mess}$. Note that the constraint is more or less similar to that obtained from the $\tilde\Phi$ decay into the SM (s)particles [cf. (72) and (73)].

Lastly we comment on the lightest messenger particle (LMP). If the inflaton dominantly decays into the messengers, the LMP is also produced. Note that, since the messenger number is conserved in Eq. (57) and in the gauge interactions, all the produced ψ_Ψ and $\psi_{\bar{\Psi}}$ eventually decay into the lightest messenger particle, which is a combination of the bosonic components of the messenger fields. If T_R exceeds $M_{\rm mess}$, the LMPs are thermalized, while, if not, they are non-thermally produced. It has been known that thus generated LMP easily overcloses the universe if it is stable [23]. So it must be unstable due to a direct or indirect interaction with the visible sector. If the LMP decays fast enough, the constraints (77) and (78) are valid. However, if the LMP decay rate is small enough, they may dominate the universe and produce large entropy at late time, diluting the pre-existing gravitinos [24]. In this case the constraints (77) and (78) cannot be applied. The detailed discussion on the LMP abundance and its effect on the thermal history may be important for constraining the mixings, but it is beyond the scope of this paper.

V. OTHER ISSUES

So far, we have assumed that it is the heavy scalar field $\tilde{\Phi}$ which dominates the energy density of the universe and which causes the reheating. However, there is a potential cosmological problem of the z modulus field, which gives dominant contribution to the SUSY breaking, i.e., $G^z G_z \simeq 3$. In fact, as we have seen in Sec. II, its mass is comparable to the gravitino mass unless there is a significant SUSY breaking effect on the z mass. Since the z field, which corresponds to $\tilde{Z}$ in the mass eigenstate basis, couples with the visible sector only via the non-renormalizable interactions, it may cause the moduli problem of itself. The importance of the $\tilde{Z}$ field for cosmology was also mentioned in Ref. [7].

The evolution of the energy density of $\tilde{Z}$ field depends on the model and cosmological scenario. In fact $\tilde{Z}$ might be displaced far from the potential minimum during the inflation, which would lead to an universe dominated by the $\tilde{Z}$'s oscillation. To be concrete, let us consider the gravity mediation. Then, if $m_z \gtrsim m_{3/2}$, the $\tilde{Z}$ decay would produce too many gravitinos with $B_{3/2} \simeq O(1)$. In addition, $\tilde{Z}$ must decay before the BBN starts. Therefore the $\tilde{Z}$–dominated universe could be consistent only if $m_{3/2} \gtrsim m_z \gtrsim 100$ TeV, which is a very challenging constraint on the structure of the SUSY breaking sector. Note that, even if the initial displacement of the z field is set to be zero in the model basis by some mechanism, the mass eigenstate $\tilde{Z}$ can obtain a finite amplitude after ϕ starts oscillating, through the $\phi - z$ mixing. Since the thermal history associated with the decay of the SUSY breaking field is strongly dependent on the detailed structure of the SUSY breaking sector as well as the cosmological scenario, further studies are required for this issue.

In the low energy SUSY breaking models, $\tilde{Z}$ may have stronger couplings with the messenger and/or the visible sector, through which $\tilde{Z}$ decays fast enough. There is an additional field scalar s in the messenger sector, which may cause a similar problem. However, the potential s-modulus problem is not serious if (64) is satisfied.

So far we have discussed the SUSY breaking models that contain direct couplings between the visible sector and the SUSY breaking field. In the case of the anomaly mediation [25], the visible sector is sequestered from the SUSY breaking sector, for example, by the geometrical separation. Since the sequestered Kähler potential is not minimal, the models generally contain finite mixings. Then the gravitino and $\tilde{Z}$ productions can be one of the dominant channels of the $\tilde{\Phi}$ decay. The distinct difference from the gravity mediation lies in the interactions between the SUSY breaking field and the visible sector: they are generally quite suppressed because of the sequestering. Thus, in the anomaly mediation, we need to investigate the subsequent decay of the SUSY breaking field with a special care, and to this end, we must go into details of the hidden sector. For instance, the minimal setup of the anomaly mediation is known to suffer from the tachyonic sleptons. To cure the charge-breaking vacuum, one might introduce an extra field to mediate the SUSY breaking effects. Then the field may affect the cosmological scenario related to the $\tilde{\Phi}$ decay as well as that of $\tilde{Z}$. Thus, the analysis quite depends on the models.

VI. CONCLUSIONS AND DISCUSSION

In this paper, we studied the decay processes of the heavy scalar ϕ, especially paying attention to the effects of the mixture between ϕ and the SUSY breaking field, z. The scalar field generally mixes with the SUSY breaking field in the Kähler potential. Then the decay amplitudes of the heavy scalar field into the lighter particles can be modified significantly by the mixing-induced interactions. We explicitly estimated the production rates of the SM (s)particles, gravitino and the SUSY breaking field. In particular, we obtained the general form of the gravitino coupling in the mass-eigenstate basis.

The mixture of ϕ with the SUSY breaking field is particularly important for the thermal history, once the field ϕ dominates the energy density of the universe. Such a scalar field may be identified, for example, with a modulus and inflaton. In this paper, we also discussed the impacts on the cosmology due to the mixing-induced decay both in the modulus and inflaton cases. Particularly, it was found that the modulus decay generally suffers from the moduli-induced gravitino problem. In the inflaton decay, it is stressed that the mixture, if any, provides a lower bound on the reheating temperature because the inflaton can decay into the SM (s)particles via the interaction that mediates the SUSY breaking effects to the visible sector. In the GMSB setup, the inflaton may rapidly decay into the SM (s)particles or the messengers due to the $\phi - s$ mixing, resulting in too high reheating temperature. Such a feature becomes prominent for the models with a large inflaton mass and VEV, like the hybrid inflation model. As well as the gravitino overproduction problem due to the mixings in the Kähler potential, such a high temperature suffers from too much abundance of the gravitino produced by the thermal scatterings.

All these difficulties are originated from the mixture between the heavy scalar and the SUSY breaking sector fields. One of the simplest solutions, especially for the inflaton, is to postulate a symmetry of ϕ which is preserved at the vacuum. In many inflation models, the inflaton acquires a VEV in the vacuum, therefore the mixings are not protected by any symmetries. In a simple class of the chaotic inflation, however, the inflaton field is invariant under a Z_2 discrete symmetry, $\phi \to -\phi$. Then the scalar VEV as well as the auxiliary component of the inflaton will be vanishing. Thus the inflaton field does not mix with the SUSY breaking field in this case. Another solution is to introduce large entropy production at a late time. However, we always need to pay attention to whether the additional field that induces the entropy dilution is free from the mixing with the SUSY breaking sector field or not.

It is a symmetry that determines whether a field mixes with another, since the symmetry dictates structure of the interactions. Once the symmetry is broken spontaneously or explicitly, there is no reason that the mixings should not occur. To construct a successful cosmological scenario, one must always check whether the mixings might affect the dynamics concerned. Although this might involve the detailed structure of e.g. the SUSY breaking sector, we believe that thus obtained constraints on the mixings will shed light on the true structure of the high energy physics.

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APPENDIX A: DECAY INTO THE GRAVITINOS IN THE GOLDSTINO PICTURE

According to the goldstino–equivalence theorem [26], the scalar–gravitino–gravitino interaction discussed in Refs. [1, 2] should also be understood in the goldstino limit, i.e., in the context of spontaneously broken global SUSY. Here we show it explicitly. The generic form of the scalar–goldstino–goldstino interaction has been derived in Ref. [27] in the context of Higgs–goldstino–goldstino interaction:

$$\mathcal{L} = \frac{1}{2F_{\rm total}} \left(\frac{\langle F^i \rangle}{F_{\rm total}} \right) \left(M^2_{ij^*} \varphi^{*j} + M^2_{ij} \varphi^j \right) \overline{\tilde{\chi}} P_L \tilde{\chi} \ + \ \text{h.c.}, \tag{A1}$$

where $M^2_{ij^*} = \langle V_{ij^*} \rangle$, $M^2_{ij} = \langle V_{ij} \rangle$ and $\tilde{\chi}$ is the (4-component) goldstino field. Notice that, when interpreted in terms of SUGRA, the interaction *does* have an enhancement factor $1/F_{\rm total} \propto 1/m_{3/2}$, and there is no chirality suppression. Assuming $M^2_{ij^*} \gg M^2_{ij}$ and taking a basis where $M^2_{ij^*} = \delta_{ij} m_i^2$, we obtain

$$\mathcal{L}_{\phi\tilde{\chi}\tilde{\chi}} = \frac{m_\phi^2}{2F_{\rm total}} \left(\frac{\langle F_\phi \rangle}{F_{\rm total}} \right) \phi^* \overline{\tilde{\chi}} P_L \tilde{\chi} \ + \ \text{h.c.}, \tag{A2}$$

leading to

$$\Gamma(\phi_{\rm R,I} \to \tilde{\chi}\tilde{\chi}) = \frac{1}{32\pi} \frac{m_\phi^5}{F_{\rm total}^2} \left(\frac{|\langle F_\phi \rangle|}{F_{\rm total}} \right)^2 . \tag{A3}$$

This can be rewritten in terms of SUGRA by using $F_{\rm total} = \sqrt{3} m_{3/2} M_P$:

$$\Gamma(\phi_{\rm R,I} \to 2\psi_{3/2}) = \frac{1}{96\pi} \frac{m_\phi^5}{m_{3/2}^2 M_P^2} \left(\frac{|\langle F_\phi \rangle|}{F_{\rm total}} \right)^2 . \tag{A4}$$

which, by using $F_\phi / F_{\rm total} = G_\phi / \sqrt{3}$, reproduces the result obtained in Refs. [1, 2]. Whether it is suppressed or enhanced by the gravitino mass depends on the fractional contribution of the ϕ–multiplet to the total amount of the SUSY breaking, $\langle F_\phi \rangle / F_{\rm total}$. In the extreme case where ϕ itself is the dominant source of the SUSY breaking, $\langle F_\phi \rangle / F_{\rm total} \simeq 1$, the rate is indeed enhanced. For $\langle F_\phi \rangle / F_{\rm total} \simeq m_{3/2}/m_\phi$ the $m_{3/2}$ dependence cancels out, and for $\langle F_\phi \rangle / F_{\rm total} \lesssim (m_{3/2}/m_\phi)^2$, the rate is suppressed by the gravitino mass.

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