Pure gravity mediation and chaotic inflation in supergravity

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We investigate the compatibility of pure gravity mediation (or minimal split supersymmetry) with chaotic inflation models in supergravity. We find that an approximate Z_2 parity of the inflaton is useful to suppress gravitino production from the thermal bath and to obtain consistent inflation dynamics. We discuss the production of the lightest supersymmetric particle through the decay of the inflaton with approximate Z_2 symmetry, and we find that a large gravitino mass is favored in order to avoid overproduction of the lightest supersymmetric particle, while a lower gravitino mass requires the tuning of parameters. This may explain why a gravitino mass of $O(100)$ TeV rather than $O(100)$ GeV may be natural.

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I. INTRODUCTION

High-scale supersymmetry (SUSY) with the gravitino mass $m_{3/2} = O(100)$ TeV is one of the most interesting models beyond the standard model. It not only explains the observed Higgs boson mass $m_h \simeq 126$ GeV [1,2] by stop- and top-loop radiative corrections [3–5], but also it is free from serious phenomenological and gravitino problems thanks to large sfermion and gravitino masses, $m_{\rm sfermion} \simeq m_{3/2} = O(100)$ TeV. Among high-scale SUSY models, pure gravity mediation (PGM) [6–8] is a particularly attractive scenario, because we do not need to introduce the Polonyi field to generate gaugino masses [9,10] and the SUSY-invariant mass (so-called μ) term of Higgs multiplets [11,12] in the minimal SUSY standard model (MSSM). Thus, the model is completely free from the cosmological Polonyi problem [13,14] (see also the minimal split SUSY [15], whose basic structure is identical to PGM).[1]

On the other hand, chaotic inflation [18] is one of the most attractive cosmic inflation scenarios [19,20]. It is free from the initial condition problem [21]; that is, inflation takes place for generic initial conditions of the inflaton field and the space-time. Chaotic inflation has been successfully realized in supergravity (SUGRA) [22].

In this paper, we investigate the compatibility of PGM with chaotic inflation. We show that in PGM, the inflaton should have a Z_2-odd parity to suppress the reheating temperature, thus avoiding gravitino overproduction from the thermal bath [23–26]. We also show that the Z_2 symmetry is helpful for the inflaton to have consistent dynamics without tuning the parameters in the inflaton sector.

We argue that, in order for the inflaton to decay, the Z_2 symmetry is softly broken by a small amount. We discuss the reheating process assuming the small breaking of the Z_2 symmetry, paying attention to the gravitino overproduction problem. It is known that the inflaton, in general, decays into gravitinos, and this leads to the overproduction of the lightest SUSY particle (LSP) [27–32]. We consider the LSP to be stable and a candidate for dark matter (DM) in the universe. We discuss how the overproduction of the LSP can be avoided. Assuming that leptogenesis [33] (for a review, see Ref. [34]) is responsible for the origin of the baryon asymmetry in the universe, we show that our solution to the above problem suggests a gravitino mass far larger than the electroweak scale, $m_{3/2} \gtrsim O(100)$ TeV, while fine-tuning the parameters in the SUSY-breaking sector and the MSSM sector is required for a smaller gravitino mass. We note that we do not use any constraints from the successful big bang nucleosynthesis (BBN) to derive the natural lower bound on the gravitino mass.

This may answer a fundamental question for high scale SUSY: why high-scale SUSY with $m_{3/2} = O(100)$ TeV is natural, but not so-called "natural SUSY" with $m_{3/2} = O(100)$ GeV. The gravitino mass was in fact expected to be $O(100)$ GeV before the Large Hadron Collider, because the electroweak scale is naturally obtained, without tuning parameters in the MSSM, when $m_{3/2} = O(100)$ GeV. In the landscape point of view [35–38], it seems difficult to understand how $m_{3/2} = O(100)$ TeV can be natural. As we show in this paper, the gravitino mass of $O(100)$ GeV requires fine-tuning to avoid LSP overproduction; otherwise, the DM density of the present universe would be outside the anthropic window [39,40]. Thus, $m_{3/2} = O(100)$ TeV may be as plausible as $m_{3/2} = O(100)$ GeV (see Fig. 1 for the schematic picture).

This paper is organized as follows. In the next section, we review chaotic inflation models in supergravity and show that the inflaton should have a Z_2-odd parity in PGM.

[1] High-scale SUSY models are also discussed in Refs. [16,17]. In Ref. [16], a mediation scale other than the Planck scale is introduced to generate soft scalar masses, and hence soft masses have a broader range than in the case of PGM. In Ref. [17], the Polonyi field is introduced to generate the μ term, and, hence, it is essentially different from the PGM.

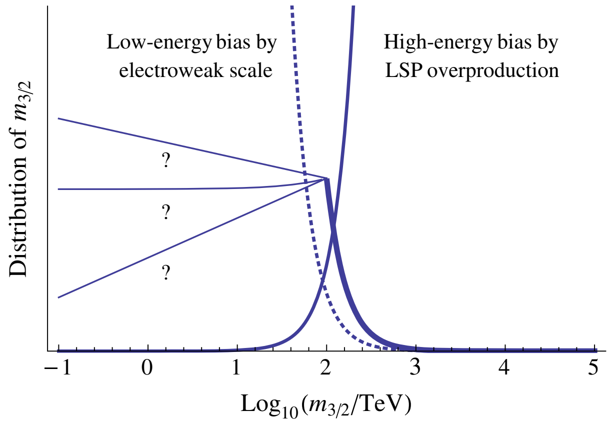

FIG. 1 (color online). A sketch of possible distributions of the gravitino mass.

In Sec. III, we discuss the decay of the inflaton into gravitinos and show how LSP overproduction can be avoided. We show that the solution to the LSP overproduction problem favors a gravitino mass far larger than the electroweak scale and that smaller gravitino masses require tuning of the parameters. The last section is devoted to discussion and conclusions.

II. CHAOTIC INFLATION MODEL IN SUPERGRAVITY

Chaotic inflation [18] is an attractive inflation model because it is free from the initial condition problem [21]: Inflation takes place for generic initial conditions of the inflaton field and the space-time. In this section, we review chaotic inflation in SUGRA, proposed in Ref. [22]. We first discuss the inflaton dynamics in the chaotic inflation model in SUGRA. Then, we show that the inflaton should have a Z_2-odd parity in PGM. In order for the inflaton to decay, we assume that the Z_2 symmetry is explicitly broken by a small value of a spurious field $\mathcal{E}$. Then, we discuss the decay of the inflaton into MSSM fields.

A. Supergravity chaotic inflation model

In SUGRA, the scalar potential is given by the Kahler potential $K(\phi^i, \phi^{*\bar{i}})$ and the superpotential $W(\phi^i)$, where ϕ^i and $\phi^{*\bar{i}}$ are chiral multiplets and their conjugate antichiral multiplets, respectively. The scalar potential is given by

$$V = e^K[K^{\bar{i}i} D_i W D_{\bar{i}} W^* - 3|W|^2],$$
$$D_i W \equiv W_i + K_i W, \tag{1}$$

where subscripts i and $\bar{i}$ denote derivatives with respect to ϕ^i and $\phi^{*\bar{i}}$, respectively. $K^{\bar{i}i}$ is the inverse of the matrix $K_{i\bar{i}}$. Here and hereafter, we use a unit of the reduced Planck mass $M_{\mathrm{pl}} \simeq 2.4 \times 10^{18}$ GeV as unity.

Chaotic inflation requires a large field value of the inflaton during inflation. With the large field value, it seems to be difficult for the slow-roll inflation to take place in SUGRA, because of the exponential factor in the scalar potential, e^K. This problem was naturally solved in Ref. [22] by assuming a shift symmetry of the inflaton chiral multiplet Φ,

$$\Phi \to \Phi + iC, \tag{2}$$

where C is a real number. The inflaton is identified with the imaginary scalar component of Φ, $\phi \equiv \sqrt{2}\mathrm{Im}\Phi$. The exponential factor vanishes for ϕ, and hence the slow-roll inflation is naturally realized for a large field value of ϕ. Note that the chiral multiplet Φ must have a vanishing R charge to be consistent with the shift symmetry.

The inflaton potential is obtained by softly breaking the shift symmetry in the superpotential,[2]

$$W = mX\Phi, \tag{3}$$

where X is a chiral multiplet with an R charge of 2. The explicit breaking of the shift symmetry is expressed by the parameter m. Here, we have eliminated the term allowed by the R symmetry, $W \supset M^2 X$, where M is a constant, by the redefinition $\Phi \to \Phi - M^2/m$.

Let us discuss the inflaton dynamics. The Kahler potential consistent with the shift symmetry is given by

$$K = c(\Phi + \Phi^\dagger) + \frac{1}{2}(\Phi + \Phi^\dagger)^2 + XX^\dagger + \cdots, \tag{4}$$

where $\cdots$ denotes higher-dimensional terms, which we neglect for simplicity. The scalar potential is given by

$$V(\phi, \sigma) = \exp(\sigma^2 + \sqrt{2}c\sigma)\frac{1}{2}m^2(\phi^2 + \sigma^2), \tag{5}$$

where σ is the real scalar component of Φ, $\sigma \equiv \sqrt{2}\mathrm{Re}\Phi$. Since X is stabilized near the origin during inflation by a Hubble-induced mass term, we have set $X = 0$ [22]. Given $\phi \gg 1$, the scalar potential is minimized for $\sigma = -c/\sqrt{2}$. Thus, the scalar potential of ϕ during inflation is given by

$$V_{\mathrm{inf}}(\phi) \simeq \frac{1}{2}m_{\mathrm{eff}}^2\phi^2, \qquad m_{\mathrm{eff}} \equiv m \times e^{-c^2/4}. \tag{6}$$

The observed magnitude of the curvature perturbation, $\mathcal{P}_\zeta \simeq 2.2 \times 10^{-9}$ [45], determines m_{eff} as

$$m_{\mathrm{eff}} \simeq 6.0 \times 10^{-6} = 1.5 \times 10^{13}\ \mathrm{GeV}, \tag{7}$$

where we have assumed that the number of e-foldings corresponding to the pivot scale of 0.002 Mpc^{-1} is as large as 50–60.

[2]For discussion on the shift symmetry breaking in the Kahler potential, see Refs. [41–44].

After inflation, extrema of the potential are given by

$$\frac{\partial V}{\partial \phi} \propto \phi = 0, \quad \frac{\partial V}{\partial \sigma} \propto \sigma\left(\sigma^2 + \frac{c}{\sqrt{2}}\sigma + 1\right) + \phi^2\left(\sigma + \frac{c}{\sqrt{2}}\right) = 0. \tag{8}$$

For $c^2 < 8$, Eq. (8) has a unique solution at the origin. For $c^2 > 8$, Eq. (8) has three solutions for σ. One of the solutions, $\sigma = 0$, is the minimum with a vanishing potential, and another solution, $\sigma = -c/(2\sqrt{2}) - \mathrm{sgn}(c)\sqrt{c^2/8 - 1}$, is a local minimum with a nonvanishing potential. The other is a local maximum. Because σ is trapped at $\sigma = -c/\sqrt{2}$ for large ϕ values, σ moves to the local minimum with a nonvanishing potential as ϕ becomes small, which prevents the inflation from ending. Thus, it is required that $c^2 < 8$.

Around the origin, the mass of the inflaton is larger than m_{eff}. Since $c^2 < 8$, the mass of the inflaton at the origin, m, is within the range of[3]

$$1.5 \times 10^{13}\ \mathrm{GeV} = m_{\mathrm{eff}} \leq m < e^2 m_{\mathrm{eff}} = 1.1 \times 10^{14}\ \mathrm{GeV}. \tag{9}$$

B. Motivation of a Z_2 symmetry

Let us consider possible couplings of the inflaton to the MSSM particles. We first note that the field X has an R charge of 2. This is mandatory because the inflaton multiplet Φ must possess a shift symmetry, so its R charge must vanish. On the other hand, in PGM, the Higgsino Dirac mass term, called the μ term, is generated by the tree-level coupling of the Higgs multiplets to the R symmetry breaking [11,12]. This ensures that the μ term is of the same order as the soft scalar mass term, i.e., the gravitino mass. This mechanism requires the combination $H_u H_d$, where H_u and H_d are the up- and down- type Higgs multiplets, to have vanishing charges under any symmetry. Therefore, the following superpotential term is not forbidden by the R symmetry,:

$$W \supset g X H_u H_d, \tag{10}$$

where g is a constant.

The inflaton decays into Higgs pairs through the coupling in Eq. (10). The resultant reheating temperature is

$$T_{\mathrm{RH}} = 1.5 \times 10^9\ \mathrm{GeV}\,\frac{g}{m}\left(\frac{m}{1.5 \times 10^{13}\ \mathrm{GeV}}\right)^{1/2}. \tag{11}$$

For $g = O(1)$, the reheating temperature is so high that too many gravitinos are produced through thermal scatterings [23–26]. The coupling g must be strongly suppressed [22].[4] The suppression is easily achieved if X and Φ are odd under a Z_2 symmetry.[5] We note that the Z_2 symmetry is also helpful to ensure successful inflaton dynamics. As we have mentioned in the previous subsection, the superpotential term of $W \supset M^2 X$ is allowed by the R symmetry. The constant M is expected to be of order 1 without the Z_2 symmetry. As we shift Φ, $\Phi \to \Phi - M^2/m$, to eliminate the superpotential term, a large linear term in the Kahler potential, $c(\Phi + \Phi^\dagger)$ in Eq. (4), is induced. However, for inflation to end, the constant c in Eq. (4) must be smaller than $\sqrt{8}$, which requires tuning of the parameters in the Kahler potential. We can easily avoid the tuning if we have the Z_2 symmetry.

Taking those problems seriously, we assume, throughout this paper, a Z_2 symmetry under which X and Φ are odd. In order for the inflaton to decay into the MSSM particles, we assume that the Z_2 symmetry is broken by a small amount, which we express by a spurious field $\mathcal{E}$.[6] Here, the spurion $\mathcal{E}$ is odd under the Z_2 symmetry and a nonvanishing value of $\mathcal{E}$ represents the Z_2 symmetry breaking. In Table I, we summarize Z_2 and R charges of Φ, X, $\mathcal{E}$, and the combination $H_u H_d$.

C. Decay of the inflaton into MSSM fields

Based on the assumption of the broken Z_2 symmetry, we consider the following superpotential and Kahler potential for the inflaton and the MSSM sectors:

$$W = X(m\Phi - \mathcal{E}) + a_1 \mathcal{E} X H_u H_d + W_{\mathrm{MSSM}} + W_0,$$
$$K = XX^\dagger + \frac{1}{2}(\Phi + \Phi^\dagger)^2 + QQ^\dagger, \tag{12}$$

where W_{MSSM} is the superpotential of the MSSM, $W_0 = m_{3/2}$ is the constant term, Q denotes MSSM fields

[3]This range is slightly widened by taking higher-dimensional terms into account in the Kahler potential. Even if the mass of the inflaton is as large as 10^{14} GeV and, hence, the decay of the inflaton after inflation produces particles with extremely large momenta, the decay products thermalize soon after their production [46]. Thus, the standard estimation of the reheating temperature in the following discussion is valid.

[4]If g is not suppressed, the F term of X strongly depends on $H_u H_d$. The $H_u H_d$ direction works as a waterfall field in the hybrid inflation [47], and thus inflation ends for $|\phi| \gg 1$. This changes the prediction on the spectral index and the tensor fraction. We note that during the waterfall phase, the instability of H_u and H_d grows and the reheating temperature becomes extremely high.

[5]The Z_2 symmetry is consistent with the shift symmetry given in Eq. (2). g can be also suppressed if $H_u H_d$ carries a Peccei-Quinn charge. For the PGM model with Peccei-Quinn symmetry, see Refs. [48,49].

[6]Alternatively, the inflaton can decay into MSSM fields if MSSM fields are also charged under the Z_2 symmetry [50,51]. We do not consider this possibility in this paper.

TABLE I. Z_2 and R charges of Φ, X, $\mathcal{E}$, and the combination H_uH_d.

	Φ	X	$\mathcal{E}$	H_uH_d
Z_2	$-$	$-$	$-$	$+$
R	0	2	0	0

collectively, and a_1 is an order 1 coefficient. We take m to be real without loss of generality. To be concrete, we have assumed the minimal form of the Kahler potential. For clarity, we shift Φ as $\Phi \to \Phi + \mathcal{E}/m$. Then the superpotential and the Kahler potential is given by

$$W = mX\Phi + a_1\mathcal{E}XH_uH_d + W_{\mathrm{MSSM}} + W_0,$$
$$K = XX^\dagger + \frac{1}{2}(\Phi + \Phi^\dagger)^2 + c(\Phi + \Phi^\dagger) + QQ^\dagger, \quad (13)$$

where $c \equiv (\mathcal{E} + \mathcal{E}^\dagger)/m$ is a real constant. For a successful inflation, c must be smaller than $\sqrt{8}$, which indicates that $|\mathcal{E}| < O(m)$.

Let us discuss the decay of the inflaton into MSSM fields. First, the inflaton decays into Higgs pairs through the coupling in the superpotential in Eq. (12) with the width

$$\Gamma(\phi \to H_uH_d) = \frac{1}{4\pi}|a_1\mathcal{E}|^2 m. \quad (14)$$

Second, if a nonvanishing superpotential of MSSM fields exists, the inflaton automatically decays through the linear term of the inflaton field in the Kahler potential [31,32]. Assuming the presence of right-handed neutrinos with Majorana masses to explain the neutrino mass [52], dominant decay modes are provided by the following superpotential:

$$W = y_t Q_3\bar{u}_3H_u + \frac{1}{2}M_N NN, \quad (15)$$

where Q_3, $\bar{u}_3$, and N are the third-generation quark doublet, the third-generation up-type quark, and a right-handed neutrino, respectively. y_t and M_N are the top Yukawa coupling and right-handed neutrino mass, respectively. For simplicity, we assume that only one right-handed neutrino is lighter than the inflaton. Decay widths of the inflaton by these interactions are

$$\Gamma(\phi \to Q_3\bar{u}_3H_u) = \frac{3}{128\pi^3}c^2y_t^2m^3,$$
$$\Gamma(\phi \to NN) = \frac{1}{16\pi}c^2mM_N^2. \quad (16)$$

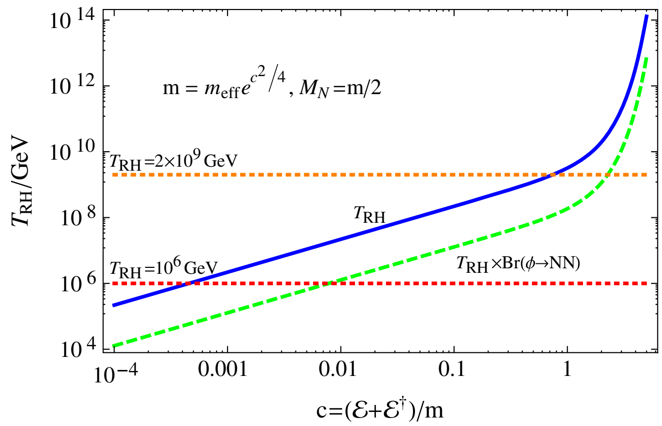

FIG. 2 (color online). The solid line shows the reheating temperature for a given parameter c. The dashed line shows the reheating temperature multiplied by the branching ratio of the inflaton into the right-handed neutrino.

Third, the inflaton couples with gauge multiplets through radiative corrections [32]. Radiative corrections induce couplings of the inflaton in kinetic functions,[7]

$$\left[\frac{1}{g^2} + i\frac{\theta_{\mathrm{YM}}}{8\pi^2} + \frac{1}{16\pi^2}c\Phi(T_{\mathrm{G}} - T_M)\right]W^\alpha W_\alpha, \quad (17)$$

where g, θ_{YM}, and W^α are the gauge-coupling constant, the theta angle, and the field-strength superfield, respectively. T_G is the Dynkin index of the adjoint representation and T_M is the total Dynkin index of matter fields. The decay width of the inflaton into the gauge multiplet V by the gauge kinetic function is given by

$$\Gamma(\phi \to VV) = \frac{\alpha^2}{256\pi^3}N_G(T_G - T_M)^2c^2m^3, \quad (18)$$

where $\alpha = g^2/4\pi$ and N_G is the number of the generator of the gauge symmetry. Because of the suppression by a one-loop factor, this decay mode is subdominant in the MSSM. As we will see, however, this decay mode plays an important role in considering the decay of the inflaton into the SUSY-breaking sector in Sec. III.

In Fig. 2, we show the relation between the reheating temperature $T_{\mathrm{RH}} \equiv 0.2\sqrt{\Gamma_{\mathrm{tot}}}$ and the parameter c, where Γ_{tot} is the total decay width of the inflaton. Here, we assume $a_1 = 1$ and that $\mathcal{E}$ is real.

Let us put a restriction on the reheating temperature, which is crucial for the discussion on the gravitino problem

[7] When one moves on to the Einstein frame and canonicalizes fields, one encounters inflaton-dependent chiral rotations of fermion fields. Thus, in the Einstein frame with canonical normalization for matter and gauge fields, the shift symmetry also involves chiral rotations of fermion fields, which is anomalous. The coupling in Eq. (17) can be understood as the counterterm for the anomaly.

in the next section. Throughout this paper, we assume that leptogenesis [33] is responsible for the origin of the baryon asymmetry of the universe. Thermal leptogenesis requires $T_{\rm RH} \gtrsim 2\times 10^9$ GeV [53,54], therefore $c \gtrsim 0.7$ is required. Since the inflaton decays into the right-handed neutrino, nonthermal leptogenesis [55–57] is also possible.[8] In Fig. 2, we also show $T_{\rm RH}\times {\rm Br}(\phi\to NN)$ by a dashed line. Here, it is assumed that $M_N = m/2$, so that the decay width of the inflaton into the right-handed neutrino is at a maximum. Nonthermal leptogenesis requires $T_{\rm RH}\times {\rm Br}(\phi\to NN)\times (2M_N/m) \gtrsim 10^6$ GeV [56,59], and hence $c \gtrsim 0.008$.

In the following, we require at least $T_{\rm RH}\times {\rm Br}(\phi\to NN)\times(2M_N/m) > 10^6$ GeV, that is, $c > 0.008$, so that nonthermal leptogenesis is possible. We also consider the more severe constraint from the successful thermal leptogenesis, $T_{\rm RH} > 2\times 10^9$ GeV; that is, $c > 0.7$. This constraint should be satisfied when $M_N \ll m$ and, hence, ${\rm Br}(\phi\to NN)\times(2M_N/m)$ is suppressed.

III. THE GRAVITINO PROBLEM AND GRAVITINO MASS

It is known that gravitinos are, in general, produced through the decay of the inflaton, which results in overproduction of the LSP [27–32]. In this section, we first discuss how gravitinos are produced from the decay of the inflaton. Then we discuss how a large gravitino mass is required in order to avoid LSP overproduction.

A. Review of the decay of the inflaton into gravitinos

Let us consider the simplest SUSY-breaking model with the following (effective) superpotential:

$$W = \mu^2 Z, \tag{19}$$

where $\mu^2 = \sqrt{3}m_{3/2}$ is the SUSY-breaking scale and Z is the SUSY-breaking field. Since the SUSY-breaking field Z does not obtain its mass from the superpotential, it should obtain its mass from the Kahler potential; otherwise, the SUSY-breaking field would obtain a large amplitude in the early universe, causing the cosmological Polonyi problem [13,14]. The Kahler potential term that yields the mass term is

$$K = -\frac{1}{\Lambda^2}ZZ^\dagger ZZ^\dagger = -\frac{m_Z^2}{12m_{3/2}^2}ZZ^\dagger ZZ^\dagger, \tag{20}$$

where $\Lambda \ll 1$ is an energy scale and m_Z is the mass of the scalar component of Z. This term is provided by the interaction of the SUSY-breaking field with other fields in the SUSY-breaking sector. The inflaton, in general, decays into those fields in the SUSY-breaking sector, as is the case with MSSM fields. Since the SUSY-breaking sector fields couple to the SUSY-breaking field Z, they eventually decay into the gravitino. We examine this issue for concrete examples later.

The inflaton also decays into a pair of gravitinos through the mass mixing between the inflaton and the scalar component of the SUSY-breaking field Z [27]. For the Kahler potential and superpotential in Eq. (13), the mass mixing is given by

$$V_{\rm mix} = \sqrt{3}cm_{3/2}mZX^\dagger + {\rm H.c.} \tag{21}$$

at around $Z=\Phi=X=0$. The mixing angle between the inflaton and the SUSY-breaking field is given by

$$\theta = \sqrt{\frac{3}{2}}c\frac{m_{3/2}m}{m_Z^2 - m^2}. \tag{22}$$

The coupling between the scalar component of Z and the gravitino, that is, the Goldstino ψ, is given by the Kahler potential in Eq. (20) as

$$\mathcal{L} = -\frac{\sqrt{3}}{6}\frac{m_Z^2}{m_{3/2}}Z^\dagger\psi\psi + {\rm H.c.} \tag{23}$$

From Eqs. (22) and (23), we obtain the decay width of the inflaton decaying into a pair of gravitinos,

$$\Gamma(\phi\to 2\psi_{3/2}) = \frac{c^2}{64\pi}m^3\left(\frac{m_Z^2}{m_Z^2-m^2}\right)^2 \simeq \begin{cases} \frac{c^2}{64\pi}m^3 & (m_Z \gg m) \\ \frac{c^2}{64\pi}\frac{m_Z^4}{m} & (m_Z \ll m). \end{cases} \tag{24}$$

The decay width is of the same order as that of the inflaton decaying into MSSM fields if $m_Z \gg m$.

Now it is clear that the inflaton, in general, decays into gravitinos. The gravitino eventually decays into the LSP. The density parameter of the LSP is given by

$$\Omega_{\rm LSP}h^2 \simeq \sum_f n_f {\rm Br}(\phi\to f)\frac{3T_{\rm RH}}{4m}\frac{m_{\rm LSP}}{3.6\times 10^{-9}\ {\rm GeV}}. \tag{25}$$

Here, f denotes decay modes and n_f is the number of gravitinos produced per decay mode. For example, $n_f = 2$ for $f = 2\psi_{3/2}$.

B. The gravitino problem in a strongly coupled SUSY-breaking model

We first discuss a strongly coupled SUSY-breaking model. To be concrete, let us consider the $SU(5)$ SUSY-breaking model [60,61]. The model is composed of a $SU(5)$ gauge theory with **10** and $\bar{\mathbf{5}}$ representations. Because

[8] Leptogenesis from inflaton decay is also discussed in Ref. [58], where the mechanism of generating the lepton asymmetry depends on the grand unification scale spectrum.

there is no parameter except for the gauge coupling, Λ and μ are as large as the dynamical scale of the $SU(5)$ gauge theory, Λ_5. Assuming a naive dimensional analysis [62,63], the Kahler potentials and superpotentials are evaluated as

$$W = c_1 \frac{\Lambda_5^2}{4\pi} \mathcal{Z},$$
$$K = \mathcal{Z}\mathcal{Z}^\dagger - c_2 \frac{16\pi^2}{\Lambda_5^2} \mathcal{Z}\mathcal{Z}^\dagger \mathcal{Z}\mathcal{Z}^\dagger, \tag{26}$$

where c_1 and c_2 are order 1 coefficients and $\mathcal{Z}$ is a composite field responsible for the SUSY breaking. Here, we have assumed that only one composite field has a nonvanishing SUSY-breaking F term, for simplicity.

As shown in Eq. (18), the inflaton decays into the $SU(5)$ gauge multiplet V_5 through the kinetic function if the dynamical scale is small enough, $m \gtrsim 2\Lambda_5$. Here, we assume that masses of hadrons of the $SU(5)$ gauge theory are as large as Λ_5. The decay rate is given by

$$\Gamma(\phi \to V_5 V_5) = \frac{27\alpha_5^2}{32\pi^2} c^2 m^3, \tag{27}$$

where α_5 is the fine structure constant of the $SU(5)$ gauge theory. Note that this decay rate is of the same order as the decay rate into MSSM particles [see Eqs. (14) and (16)], and hidden hadrons eventually decay into gravitinos.

Even if the decay mode is kinematically closed, $m \lesssim 2\Lambda_5$, the direct decay into gravitinos is unsuppressed because $m_{\mathcal{Z}} \sim \Lambda_5 \gtrsim m$ [see Eq. (24)]. Thus, the decay of the inflaton inevitably produces gravitinos, and the resultant density parameter of the LSP is

$$\Omega_{\rm LSP} h^2 \simeq \frac{T_{\rm RH}}{m} \frac{m_{\rm LSP}}{3.6 \times 10^{-9}\ {\rm GeV}}. \tag{28}$$

The universe is overclosed by the LSP unless

$$m_{\rm LSP} \lesssim 10\ {\rm MeV} \frac{m}{1.5 \times 10^{13}\ {\rm GeV}} \frac{10^6\ {\rm GeV}}{T_{\rm RH}}. \tag{29}$$

When $m_{\rm LSP}$ is so small, however, thermally produced LSPs overclose the universe (recall the Lee-Weinberg bound [64]) unless the LSP is degenerated with a charged SUSY particle. Such a light charged SUSY particle is already excluded by various experiments.[9] We will not consider the strongly coupled SUSY-breaking model below.

C. The gravitino problem in a SUSY-breaking model with weak coupling

The origin of the failure of the strongly coupled SUSY-breaking model is that the decay of the inflaton either into gravitinos or into SUSY-breaking sector fields is unsuppressed. Note that simultaneous suppression of these two decay modes is achieved by realizing the following hierarchy:

$$m_Z \ll m \ll m_{\text{SUSY-breaking}}, \tag{30}$$

where $m_{\text{SUSY-breaking}}$ is the mass scale of SUSY-breaking sector fields. We show in this subsection that this hierarchy is easily achieved if the SUSY-breaking sector involves weak couplings [65].

To be concrete, let us consider the Izawa-Yanagida-Intriligator-Thomas SUSY-breaking model [66,67] with the $SU(2)$ gauge theory. We introduce four fundamental representations of the $SU(2)$, $Q_i (i = 1\text{–}4)$. Below the dynamical scale of the $SU(2)$, Λ_2, the theory is described by meson fields with the deformed moduli constraint [68]

$$W_{\rm dyn} = 4\pi \Xi \left({\rm Pf} M_{ij} - \frac{\Lambda_2^2}{16\pi^2} \right), \tag{31}$$

where $M_{ij} = -M_{ji} \sim Q_i Q_j / \Lambda_2$ are meson fields and Ξ is a Lagrange multiplier field. Pf denotes the Pfaffian over indices i, j. Here, we again assume a naive dimensional analysis and put order 1 coefficients to unity. It can be seen that there are flat directions, in which ${\rm Pf} M_{ij} = \Lambda_2^2/16\pi^2$.

To fix the flat directions, let us introduce five singlet chiral multiplets, Z_a ($a = 1\text{–}5$), and assume the following superpotential:

$$W_{\rm tree} = \lambda c_{a,ij} Z_a Q_i Q_j, \tag{32}$$

where λ and $c_{a,ij}$ are constants. To simplify our discussion, we assume a global $SO(5)$ symmetry under which Z_a and Q_i are the vector and spinor representation of the $SO(5)$ symmetry. $c_{a,ij}$ should be the appropriate Clebsch-Gordan coefficients. Adding Eqs. (31) and (32), we obtain the effective superpotential

$$W = \frac{\lambda}{4\pi} \Lambda_2 Z_a M_a + 4\pi \Xi \left(M_a M_a + M^2 - \frac{\Lambda_2^2}{16\pi^2} \right), \tag{33}$$

where we take linear combinations of meson fields and form a vector representation of the $SO(5)$, M_a ($a = 1\text{–}5$). M is the remaining independent linear combination. Now, flat directions are fixed and the vacuum is given by $Z_a = M_a = 0$, $M = \Lambda_2/4\pi$.

To break the SUSY, we add an additional singlet chiral multiplet Z and add the superpotential

$$\Delta W = y Z c_{ij} Q_i Q_j, \tag{34}$$

where y is a constant and c_{ij} is an appropriate Clebsch-Gordan coefficient, to form a singlet of the $SO(5)$. Adding Eqs. (33) and (34), we obtain the superpotential

[9] In PGM, the photino LSP of a mass of $O(10)$ MeV is naturally obtained if $m_{3/2} = O(1)$ GeV. In this case, however, the electroweak symmetry-breaking scale is also $O(1)$ GeV.

$$W=\frac{y}{4\pi}\Lambda_2 ZM+\frac{\lambda}{4\pi}\Lambda_2 Z_a M_a+4\pi\Xi\left(M_a M_a+M^2-\frac{\Lambda_2^2}{16\pi^2}\right). \tag{35}$$

Assuming $y \ll \lambda$, the vacuum is given by $Z_a \simeq M_a \simeq 0$, $M \simeq \Lambda_2/4\pi$. The F term of Z is nonzero and, hence, the SUSY is spontaneously broken.

Let us discuss the decay of the inflaton into the SUSY-breaking sector. If the dynamical scale is small enough, $m \gtrsim 2\Lambda_2$, the inflaton decays into gauge multiplets of the $SU(2)$, as shown in Eq. (18). The decay rate and n_f are

$$\Gamma(\phi \to V_2 V_2) = \frac{3\alpha_2^2}{16\pi^3}c^2 m^3 (\text{for } m > 2\Lambda_2), \qquad n_{V_2V_2} \geq 4, \tag{36}$$

where α_2 is the fine structure constant of the $SU(2)$ gauge theory. As in the case of the $SU(5)$ model, for $m > 2\Lambda_2$, this decay mode is as dominant as decay modes into MSSM fields, and, hence, the gravitino is overproduced.

If $m < 2\Lambda_2$, on the other hand, the mass of the inflaton is not far above the dynamical scale, and, hence, we can treat the decay of the inflaton into SUSY-breaking sector fields as composite fields. The decay rate through mass terms in Eq. (35) and n_f are

$$\begin{aligned}\Gamma(\phi \to Z_a M_a) &= \frac{5}{16\pi}c^2 m\left(\frac{\lambda}{4\pi}\Lambda_2\right)^2\left(\text{for } m > 2\frac{\lambda}{4\pi}\Lambda_2\right),\\ n_{Z_aM_a} &= 4,\\ \Gamma(\phi \to ZM) &= \frac{1}{16\pi}c^2 m\left(\frac{y}{4\pi}\Lambda_2\right)^2 (\text{for } m > \Lambda_2),\\ n_{ZM} &= 4.\end{aligned} \tag{37}$$

To be conservative, we assume that $\lambda \simeq 4\pi$. In this case, the decay into $Z_a M_a$ is kinematically forbidden.

Now, we are at the point of showing that the desired hierarchy in Eq. (30) can be realized. After integrating out Z_a, M_a, and M, we are left with the effective superpotential

$$W_{\text{eff}} = \frac{y}{16\pi^2}c_3\Lambda_2^2 Z, \tag{38}$$

where $c_3 = 1$ is a constant, which we leave as a free parameter for later convenience. The dynamical scale Λ_2 is related to the gravitino mass as

$$\begin{aligned}\Lambda_2 &= 3^{1/4}m_{3/2}^{1/2}4\pi y^{-1/2}c_3^{-1/2}\\ &= 2.6\times 10^{12}\ \text{GeV}y^{-1/2}\left(\frac{m_{3/2}}{10\ \text{TeV}}\right)^{1/2}c_3^{-1/2}.\end{aligned} \tag{39}$$

The mass of the scalar component of Z is given by the Kahler potential

$$K = -\frac{y^4}{16\pi^2\Lambda_2^2}ZZ^\dagger ZZ^\dagger, \tag{40}$$

and is as large as

$$m_Z = \frac{2y^3}{(4\pi)^3}\Lambda_2 c_3 = 2.6\times 10^9\ \text{GeV}y^{5/2}\left(\frac{m_{3/2}}{10\ \text{TeV}}\right)^{1/2}c_3^{1/2}. \tag{41}$$

It can be seen that the hierarchy in Eq. (30) is achieved for small y, and, hence, the overproduction of the LSP is avoided.

For a small y, however, the scalar component of Z is light and the oscillation of the scalar Z is induced in the early universe [65]. The oscillation eventually decays into gravitinos, which may lead to overproduction of the LSP. Let us estimate the abundance of the LSP from this contribution. The potential of the scalar component of Z during inflation is given by

$$V(Z) = a_2 H_{\text{inf}}^2|Z|^2 + m_Z^2|Z|^2 - (2\sqrt{3}m_{3/2}^2 Z + \text{H.c.}), \tag{42}$$

where H_{inf} is the Hubble scale during inflation and a_2 is an order 1 constant, which we assume to be positive. Since $H_{\text{inf}} \simeq 10^{14}$ GeV is larger than m_Z for the parameter of interest, the Hubble-induced mass term traps Z to its origin during inflation. After inflation, as the Hubble scale drops below m_Z, Z begins its oscillation around the origin,

$$Z_0 = 2\sqrt{3}m_{3/2}^2/m_Z^2 = 1.2\times 10^8\ \text{GeV}y^{-5}\frac{m_{3/2}}{10\ \text{TeV}}c_3^{-1}, \tag{43}$$

with an initial amplitude $Z_i = Z_0$. As anticipated, the amplitude is larger for smaller y. The LSP abundance originating from the oscillation of Z is given by

$$\Omega_{\text{osc}}h^2 = \frac{T_{\text{RH}}}{4m_Z}\frac{Z_i^2}{M_{\text{pl}}^2}\frac{m_{\text{LSP}}}{3.6\times 10^{-9}\ \text{GeV}}. \tag{44}$$

Let us show how a large gravitino mass is required. In Fig. 3, we show the constraint on $m_{3/2}$ and y. Here, we assume that $c = 0.7$ (i.e., thermal leptogenesis is possible) and $m_{\text{LSP}} = 3\times 10^{-3}m_{3/2}$. In the red-shaded region ($\Omega_{\text{SUSY}}h^2 > 0.1$), the universe is overclosed by the LSP due to the decay of the inflaton into SUSY-breaking sector fields [see Eqs. (36) and (37)]. The right edge of this region is determined by the kinematical threshold, $m = \Lambda_2$. In the blue-shaded region ($\Omega_{3/2}h^2 > 0.1$), the decay of the inflaton into a pair of gravitinos causes the overclosure. In the yellow-shaded region ($\Omega_{\text{osc}}h^2 > 0.1$), the oscillation of the SUSY-breaking field leads to the overclosure. In Fig. 4, we show the same constraint for $c = 0.008$ (i.e., nonthermal

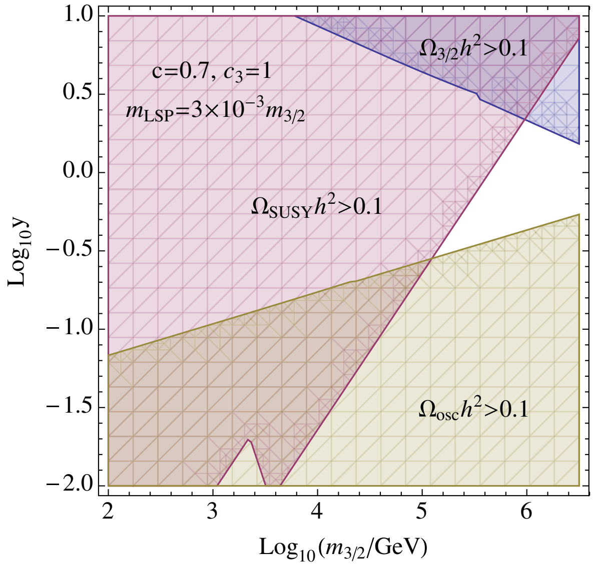

FIG. 3 (color online). Constraint on the gravitino mass $m_{3/2}$ and the coupling of the SUSY-breaking field y. In the red-shaded region ($\Omega_{\rm SUSY}h^2 > 0.1$), the blue-shaded region ($\Omega_{3/2}h^2 > 0.1$), and the yellow-shaded region ($\Omega_{\rm osc}h^2 > 0.1$), the universe is overclosed by the LSP due to the decay of the inflaton into SUSY-breaking sector fields, that of the inflaton into gravitino pairs, and that of the SUSY-breaking field into gravitinos, respectively. We assume $c = 0.7$, $c_3 = 1$, and $m_{\rm LSP} = 3 \times 10^{-3} m_{3/2}$.

leptogenesis is possible). From both figures, we see the constraint on the gravitino mass,[10]

$$m_{3/2} > O(100)\ \text{TeV}. \tag{45}$$

It is remarkable that the constraint in Eq. (45) coincides with what is expected in PGM [6–8].

Let us discuss how we can avoid the constraint on the gravitino mass. First, we have assumed that $m_{\rm LSP} = 3 \times 10^{-3} m_{3/2}$ to obtain the constraint, since it is determined by the anomaly mediation [9,10]. However, a lower mass for the LSP can be obtained by canceling the anomaly-mediated contribution, using the Higgsino threshold correction [9]. In Fig. 5, we show the constraint on $m_{3/2}$ and y for $(c, m_{\rm LSP}) = (0.7, 3 \times 10^{-6} m_{3/2})$. It can be seen that regions with $m_{3/2} = O(10)$ TeV are allowed.

Let us compare the plausibility of $m_{3/2} = O(10)$ TeV with that of $m_{3/2} = O(100)$ TeV in the landscape point of view. Since we have no knowledge about the distribution of parameters in the landscape, the following discussions are based on our naive expectations. We note that different assumptions about the distributions lead to different conclusions.

[10]A similar conclusion is derived in Ref. [69], where the BBN constraints are used. Notice that we have obtained Eq. (45) solely from constraints on the LSP DM density.

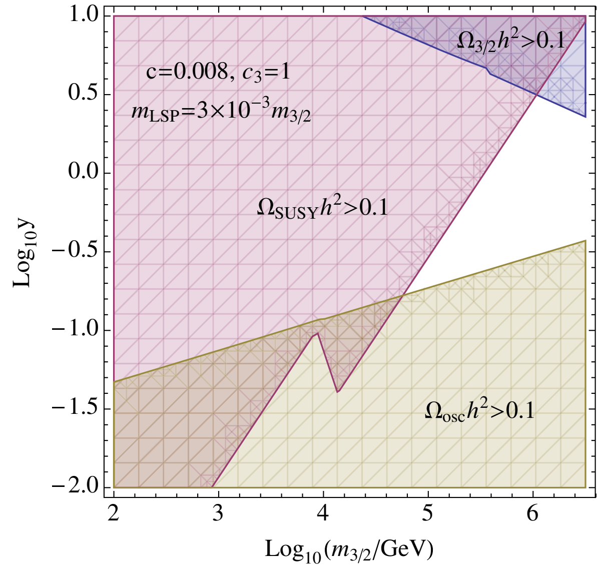

FIG. 4 (color online). Same as for Fig. 3 but with $c = 0.008$, $c_3 = 1$, and $m_{\rm LSP} = 3 \times 10^{-3} m_{3/2}$.

For the electroweak scale, $m_{3/2} = O(10)$ TeV would be more natural than $m_{3/2} = O(100)$ TeV by a factor of $(100\ \text{TeV})^2/(10\ \text{TeV})^2 = 100$. Because the LSP mass is a complex parameter, $m_{\rm LSP} = 3 \times 10^{-6} m_{3/2}$ would require tuning of $(3 \times 10^{-6}/3 \times 10^{-3})^2 \sim 10^{-6}$. Thus, we naively expect that the region with $m_{3/2} = O(100)$ TeV may be more natural than the region with $m_{3/2} = O(10)$ TeV.

Second, to simplify our discussion, we have assumed the $SO(5)$ symmetric Izawa-Yanagida-Intriligator-Thomas model. Without the $SO(5)$ symmetry, c_3 in Eq. (38) is a constant that is determined by coupling constants in the

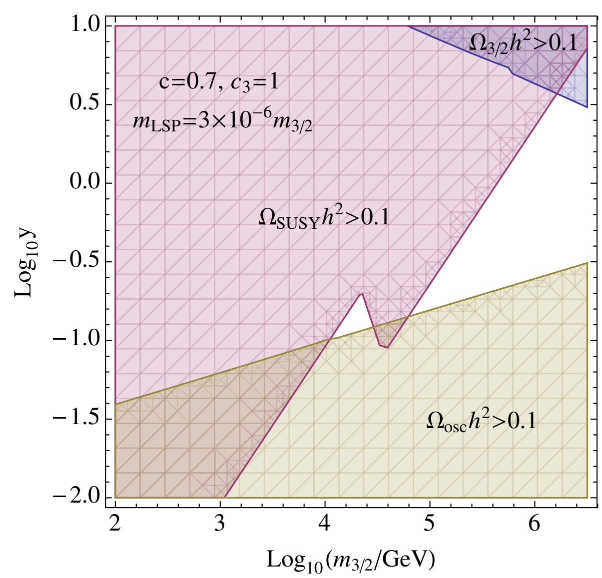

FIG. 5 (color online). Same as for Fig. 3 but with $c = 0.7$, $c_3 = 1$, and $m_{\rm LSP} = 3 \times 10^{-6} m_{3/2}$.

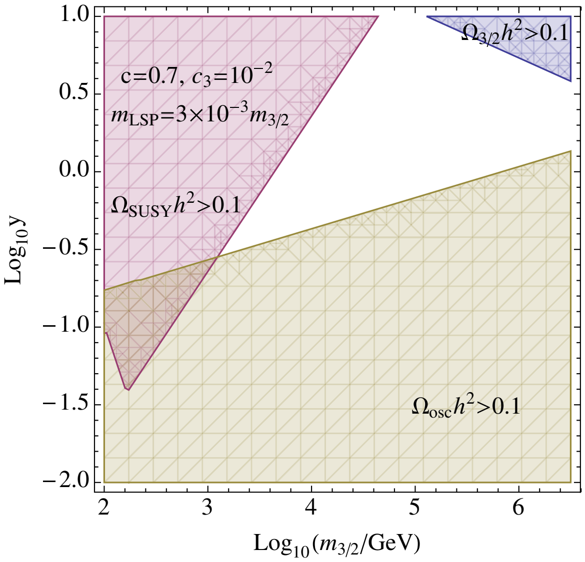

FIG. 6 (color online). Same as for Fig. 3 but with $c = 0.7$, $c_3 = 10^{-2}$, and $m_{\rm LSP} = 3 \times 10^{-3} m_{3/2}$.

SUSY-breaking model. If there is fine-tuned cancellation between condensation of hidden quarks that couple to the SUSY-breaking field, c_3 can be much smaller than $O(1)$. This cancellation further separates the SUSY-breaking scale from the dynamical scale. For given $m_{3/2}$ and y, the constraints shown in Figs. 3 and 4 are relaxed. In Fig. 6, we show the constraint for $(c, c_3) = (1.4, 10^{-2})$. It can be seen that the region with $m_{3/2} = O(1)$ TeV survives.

Let us again naively compare the plausibility of $m_{3/2} = O(1)$ TeV with that of $m_{3/2} = O(100)$ TeV. For the electroweak scale, $m_{3/2} = O(1)$ TeV would be more natural than $m_{3/2} = O(100)$ TeV by a factor of $(100\ \mathrm{TeV})^2/(1\ \mathrm{TeV})^2 = 10^4$. On the other hand, since c_3 is a complex parameter, $c_3 = 10^{-2}$ would require fine-tuning of 10^{-4}. These two regions, $m_{3/2} = O(1)$ TeV and $O(100)$ TeV, may be equally plausible.

Third, we have assumed the minimal form of the Kahler potential. By considering higher-dimensional terms in the Kahler potential and tuning their coefficients, the decay of the inflaton into the SUSY-breaking sector can be suppressed. In principle, the gravitino mass of $O(100)$ GeV survives because of the tuning. However, to suppress all the decay modes, all the coefficients of the higher-dimensional terms must be carefully chosen; this may require more fine-tuning.

We should stress, finally, that all of the above arguments are merely a sketch of the kinds of fine-tuning that are needed to ensure that the gravitino mass is below $O(100)$ TeV. Because we do not know the distributions of relevant parameters in landscape, it is impossible for us to draw any definite conclusions of the most plausible gravitino mass. However, the present analysis shows that it is not necessarily surprising if the gravitino mass is of $O(100)$ TeV, even if the SUSY-breaking scale is at a low energy in order to obtain the electroweak scale naturally.

IV. DISCUSSION AND CONCLUSION

In this paper, we have investigated the compatibility of PGM with chaotic inflation in supergravity. We have shown that the inflaton should have a Z_2-odd parity to suppress the reheating temperature, thus avoiding gravitino overproduction from the thermal bath in PGM. We have also shown that the Z_2 symmetry is helpful for the inflaton to have consistent dynamics without tuning the parameters in the inflaton sector.

We assume that, in order for the inflaton to decay, the Z_2 symmetry is broken by a small amount. We have discussed the reheating process and the gravitino problem under the assumption of a small breaking of the Z_2 symmetry. We have discussed how gravitino overproduction by the decay of the inflaton can be avoided, and we have shown that the solution to the overproduction problem favors a gravitino mass far larger than the electroweak scale, $m_{3/2} \gtrsim O(100)$ TeV.

This consideration gives a new insight into the fine-tuning problem in the high-scale SUSY. It is usually assumed that the gravitino mass of $O(100)$ GeV is natural, because the electroweak scale is obtained without tuning the parameters in the MSSM. It is hard to understand how a gravitino mass of $O(100)$ TeV can be natural. However, as we have shown in this paper, the gravitino mass of $O(100)$ GeV requires an amount of fine-tuning to avoid LSP overproduction. Therefore, it may be natural to have high-scale SUSY with a gravitino mass $O(100)$ TeV.

In this paper, in order to suppress the reheating temperature, we have assumed a Z_2 symmetry. Another option is to assume the spatial separation of the inflaton sector and the MSSM sector in a higher-dimensional theory. Our discussion of the LSP overproduction is also applicable to this case.

We should note that we can replace the inflaton mass m in Eq. (12) by a vacuum expectation value of some field. Consider a B-L gauge symmetry, for example, that is broken by a vacuum expectation value of a chiral multiplet S with a B-L charge of $+1$. We assume that X carries a B-L charge of -1, f so that the following superpotential is allowed [70]:

$$W = k\Phi S X, \qquad k\langle S\rangle = m. \tag{46}$$

The Yukawa coupling k represents a shift symmetry breaking. We may take $k = O(0.1)$ and $\langle S\rangle = O(10^{-4})$ as an example. The unwanted linear term $W = M^2 X$ is replaced by $W = \mathcal{M}\langle S\rangle X$, and the required condition $M^2 = \mathcal{M}\langle S\rangle \lesssim m$ may be explained by $\mathcal{M} \lesssim 0.1$ without the Z_2 symmetry.

In this paper, we have assumed that $m_{\rm LSP} = O(10^{-3}) m_{3/2}$, which is true in the case of PGM. In general gravity mediation models with a singlet SUSY-breaking field (i.e., a Polonyi field), the LSP mass is expected to be of the same order as the gravitino mass. If this is the case, thermally produced LSPs will easily overclose the universe unless the reheating temperature is far smaller than the LSP mass. For $T_{\rm RH} \gtrsim 10^6$ GeV, the gravitino mass smaller than 10^7 GeV is excluded.

Finally, let us comment on other inflation models. The lower bound on the gravitino mass in Eq. (45) is basically obtained from the condition that masses of SUSY-breaking sector fields are larger than the inflaton mass. Thus, for models with the inflation mass of $O(10^{13})$ GeV, a similar bound to that on the gravitino mass in Eq. (45) will be obtained. On the contrary, if models have the maximal reheating temperature, the lower bound on the gravitino mass may be obtained [71] so that enough LSPs are produced through the gravitino production in the thermal bath to explain the DM density.

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