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Supergravity Inflation Free from Harmful Relics

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We present a realistic supergravity inflation model which is free from the overproduction of potentially dangerous relics in cosmology, namely moduli and gravitinos which can lead to the inconsistencies with the predictions of baryon asymmetry and nucleosynthesis. The radiative correction turns out to play a crucial role in our analysis which raises the mass of supersymmetry breaking field to intermediate scale. We pay a particular attention to the non-thermal production of gravitinos using the non-minimal Kähler potential we obtained from loop correction. This non-thermal gravitino production however is diminished because of the relatively small scale of inflaton mass and small amplitudes of hidden sector fields.

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I. INTRODUCTION

There exist generic and well known problems in constructing supergravity inflation models with broken local supersymmetry in vacuum. First, one must carefully choose the superpotential and Kähler potential such that non-renormalizable terms do not spoil the flatness of the inflaton potential, which is also essential to obtain the observed CMB spectrum. Second, most supergravity inflation models run into cosmological problems at late times due to an over-abundance of harmful relics such as moduli and gravitinos, and we need to check that a model is free from these problems to be consistent with the data of baryon asymmetry and nucleosynthesis predictions.

Several supergravity inflation models free from abundant moduli have been proposed [1–5], and we give a brief review for this moduli problem with emphasis on the importance of the radiative corrections. Especially when the supersymmetry breaking field has flat direction at tree level, radiative correction has a significant effect on its potential to lead to the necessity for the modification of minimal Kähler potential.

More recently, the nonthermal production of gravitinos has been drawn attention [6–10]. These problems on gravitino production during preheating era, however, have been analyzed only using the minimal Kähler potential so far. The gravitino interactions depend on the form of Kähler potential and consequently the nonthermal production of gravitinos can depend on its form as well. Our treatment is the first analysis of gravitino non-thermal production taking account of the non-minimal Kähler potential obtained from the loop correction of supersymmetry breaking field.

The detailed analysis for nonthermal production of gravitinos in a system of coupled fields has been done only for models with non-renormalizable hidden sector supersymmetry breaking, and only the Polonyi model in particular [7,11]. It was shown, in this model, that dominant fermion fields which are created efficiently through preheating mechanism are inflatinos rather than gravitinos (thus free from gravitino problem). This type of supergravity inflation model is a good toy model to investigate the nonthermal particle production in preheating era, but is not a realistic inflation model in that it sill suffers from Polonyi problem which is as serious a problem as gravitino problem. We give the first realistic supergravity inflation model in this sense where nonthermal production of gravitinos as well as moduli problem were explicitly analyzed.

The paper is structured as follows. In section II, we explain our choice of superpotential and discuss how inflation develops in our model. We then discuss how supersymmetry breaking field evolves and calculate its radiative correction and its modification to minimal Kähler potential arising from this loop correction. In section III, we see if our model leads to any cosmological crisis, namely, moduli and gravitino problems. We give the conclusion and discussion at the end.

II. SETUP

We consider the superpotential [1] consisting of the inflaton sector and hidden sector of O'Raifeartaigh type [13] which are gravitationally coupled to each other.

$$W = \Delta^2 \, \frac{(\Sigma - M)^2}{M_p} + \Phi_1 \left(\kappa \Phi_2^2 - \mu^2\right) + \lambda \Phi_2 \Phi_3 + C, \tag{2.1}$$

where superfield Σ includes inflaton scalar component σ and Φ_1 includes O'Raifeartaigh scalar field ϕ_1. $\Delta \sim 10^{-4} M_p$ from COBE normalization and M is set to $M_p (= 2.436 \cdot 10^{18}\mathrm{GeV})$ so that the inflaton potential keeps the flatness around the origin (i.e. for $\frac{\partial}{\partial\sigma} V(0) \simeq 0, \frac{\partial^2}{\partial\sigma^2} V(0) \simeq 0$). The dimensionless parameter κ is of order unity while the other mass parameters μ and λ are of intermediate scale ($\sim 10^{-8} M_p$). C is the constant term to cancel the cosmological constant at the vacuum. We start with the discussion for the evolution of scalar fields in inflaton sector and hidden sector which are only gravitationally coupled to each other and, for the sake of clarity, we first treat each sector separately followed by the discussion including the coupling with non-minimal Kähler potential arising from the radiative correction. We assume gauge singlets in the potential for simplicity in the following.

A. Inflaton Sector

The superpotential of inflaton sector is given as

$$W_{inflaton} = \Delta^2 \, \frac{(\Sigma - M_p)^2}{M_p}. \tag{2.2}$$

The general expression for supergravity potential becomes, for Kähler potential, K, and superpotential, W,

$$V = m_i \; {K^{-1}}_j{}^i \; m^j - 3 \; M_p^{-2} |m|^2 \tag{2.3}$$

with

$$K^i{}_j \equiv \frac{\partial^2 K}{\partial \phi_i \partial \phi^j}, \quad m \equiv \mathrm{e}^{\frac{\mathrm{K}}{2\mathrm{M_P^2}}} \mathrm{W} \tag{2.4}$$

$$m^i \equiv D^i m \equiv \partial^i m + \frac{1}{2M_p^2} \left(\partial^i K\right) m \; . \tag{2.5}$$

If we consider, for the moment, the minimal form of Kähler potential, $K = \Sigma^\dagger \Sigma$, the effective supergravity potential from $W_{inflaton}$ for the real part of scalar component, σ, becomes (in natural units)

$$V_{inflaton} = \mathrm{e}^{\mathrm{K}} \left(\left| \frac{\partial \mathrm{W}}{\partial \Sigma} + \Sigma^\dagger \mathrm{W} \right|^2 - 3 |\mathrm{W}|^2 \right) \tag{2.6}$$

$$= \Delta^4 \; \mathrm{e}^{\sigma^2/2} \left(1 - \frac{\sigma^2}{2} - \sqrt{2}\sigma^3 + \frac{7}{4}\sigma^4 - \frac{1}{\sqrt{2}}\sigma^5 + \frac{\sigma^6}{8} \right). \tag{2.7}$$

We here point out the absence of linear and quadratic terms which enables the potential to keep the flatness around the origin. We can obtain the value of $\Delta \sim 10^{-4} M_p$ from COBE normalization condition [15] *,

$$\left(\frac{V}{\epsilon} \right)^{\frac{1}{4}} \simeq 0.027 M_p (1 - 3.2\epsilon + 0.5\eta), \tag{2.8}$$

which should be evaluated at the horizon exit. Scale of inflaton field when the inflation ends and the cosmological scales leave the horizon are obtained from slow-roll conditions, $\epsilon \lesssim 1$, $\eta \lesssim 1$, and 60 e-folding condition,

$$N(\sigma_{exit}) \simeq \int_{\sigma_{end}}^{\sigma_{exit}} \frac{V}{V'} d\sigma \simeq 60. \tag{2.9}$$

We also note the scale of inflation is of the order Δ^4 and the mass of the inflaton is of the order Δ^2/M_p with its decay width $\Gamma_\sigma \simeq m_\sigma^3/M_p^2 = \Delta^6/M_p^5$ assuming gravitational strength coupling to ordinary fields.

B. Hidden Sector

The supersymmetry breaking sector is that of O'Raifeartaigh model,

$$W_{hidden} = \Phi_1 \left(\kappa \Phi_2^2 - \mu^2\right) + \lambda \Phi_2 \Phi_3 + C \quad . \tag{2.10}$$

This is a familiar example of supersymmetry breaking due to non-vanishing F-term from Φ_1, $|F| = \mu^2$, in the vacuum. Therefore we add $C = \frac{\mu^2}{\sqrt{3}} M_p$ (compare with $\frac{-3|W|^2}{M_P}$ term in eqn(2.7)) for the vanishing cosmological constant at the vacuum.† We should, however, expect

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* [1] gives an order of estimates $10^{-4} M_p \leq \Delta \leq 10^{-3.5} M_p$ from the constraints on gravitino abundance and proton decay.

†This additional constant term C, strictly speaking, should be modified if we include the radiative correction and the coupling between inflaton and hidden sectors. We, however, stick to this value of C for simplicity because this modification essentially does not change our discussion.

that, when the fields are far away from the vacuum during the inflation, there are additional F-terms from other fields in the effective scalar potential and these F-terms lead to additional 'cosmological constant' Λ^4 [3]. Adding $e^K\Lambda^4$ in the potential indicates us that this cosmological constant term during the inflation gives an additional effective mass of order $\frac{\Lambda^4}{M_p^2}$ to each field in the model. Fields ϕ_2 and ϕ_3 do not posses the linear terms and these two fields quickly roll down to the origin (i.e. to their minimum) during the inflation. The scalar field ϕ_1, however, has a liner term and its minimum shifts according to the evolution of inflaton field as we shall see in the next section. Because we are interested in the particle production after the inflation, we focus on the evolution of O'Raifeartaigh field ϕ_1 among the fields in this supersymmetry breaking sector. Moreover, the F-term at the vacuum has the contribution only from the scalar field ϕ_1 which turns out to have flat direction at the tree level. Because of this flat potential with respect to ϕ_1 at the tree level, the global supersymmetry radiative corrections have a significant effect on the effective potential and consequently give non-negligible modification to the minimal Kähler potential. The radiative corrections of local supersymmetry are always Planck mass suppressed and we do not consider them here. We calculated this non-minimal Kähler potential from the calculation of loop correction [16,17]. Setting the parameter range to be $2\kappa\mu^2 < \lambda^2$ to make ϕ_2 and ϕ_3 stay at the origin in the vacuum leads to the following one-loop correction due to ϕ_1,

$$V_{one\,loop} = \frac{1}{64\pi^2}\left(\sum_{i=1}^{4}(M_i^2)^2\left[\log\left(\frac{M_i^2}{\lambda^2}\right)-\frac{3}{2}\right]\right. \quad (2.11)$$

$$\left. -2\sum_{i=1}^{2}(N_i^2)^2\left[\log\left(\frac{N_i^2}{\lambda^2}\right)-\frac{3}{2}\right]\right), \quad (2.12)$$

where we have defined

$$M_1^2 = \frac{1}{2}(A_1 - A_2), M_2^2 = \frac{1}{2}(A_1 + A_2),$$
$$M_3^2 = \frac{1}{2}(A_3 - A_4), M_4^2 = \frac{1}{2}(A_3 + A_4),$$
$$N_1^2 = \frac{1}{2}(B_1 - B_2), N_2^2 = \frac{1}{2}(B_1 + B_2),$$
$$A_1 = 2\lambda^2 - 2\kappa\mu^2 + 4\kappa^2|\phi_1|^2, \quad A_3 = 2\lambda^2 + 2\kappa\mu^2 + 4\kappa^2|\phi_1|^2,$$
$$A_2 = \sqrt{4\kappa^2\mu^4 + 16\lambda^2\kappa^2|\phi_1|^2 - 16\kappa^3\mu^2|\phi_1|^2 + 16\kappa^4|\phi_1|^4},$$
$$A_4 = \sqrt{4\kappa^2\mu^4 + 16\lambda^2\kappa^2|\phi_1|^2 + 16\kappa^3\mu^2|\phi_1|^2 + 16\kappa^4|\phi_1|^4},$$
$$B_1 = 2\lambda^2 + 4\kappa^2|\phi_1|^2, \quad B_2 = \sqrt{16\lambda^2\kappa^2|\phi_1|^2 + 16\kappa^4|\phi_1|^4}. \quad (2.13)$$

We have used the $\overline{MS}$ scheme and taken the renormalization scale to be λ. We are concerned with the regime $|\phi_1| \lesssim \lambda/\kappa$, as we shall show in the next section. In this case, and for $2\kappa\mu^2 \ll \lambda^2$, we can approximate the one loop potential as

$$V_{one\,loop} = C_1 + \frac{\kappa^2\mu^4}{8\pi^2}\left(\frac{\kappa^2|\phi_1|^2}{\lambda^2} + \ldots\right), \quad (2.14)$$

where ... are terms of order $\kappa^4|\phi_1|^4/\lambda^4$ and higher and C_1 is a small constant term ($\ll \mu^4$) which can be absorbed into the constant part of the superpotential. We can find the following one loop correction to the Kähler potential ($K = K_{minimal} + K_{correction}$), by comparing eqn(2.14) with eqn(2.3),

$$K_{correction} = -\frac{\kappa^2}{32\pi^2}\left(\frac{\kappa^2\Phi_1^2\Phi_1^{\dagger 2}}{\lambda^2}\right). \quad (2.15)$$

We shall use this non-minimal Kähler potential in our analysis for non-thermal production of gravitinos. We note that this radiative correction enhances the coupling to the longitudinal component of gravitino and raises the mass of ϕ_1 which was massless at the tree level to the intermediate scale

$$m_{\phi_1}^2 = \frac{\kappa^4\mu^4}{4\pi^2\lambda^2} = \frac{\alpha_\kappa^2\mu^4}{\lambda^2} \text{ with } \alpha_\kappa \equiv \frac{\kappa^2}{4\pi}, \quad (2.16)$$

which turns out to be crucial to evade the moduli problem.

III. HARMFUL RELICS

We are now in a position to discuss the fate of inflaton and O'Raifeartaigh fields in the coupled effective potential with non-minimal Kähler potential to see if our model leads to any cosmological crisis. We first briefly review the resolutions of so-called Polonyi or moduli problem and we further discuss the non-thermal production of gravitinos.

A. Moduli Problem

We here start with the discussion on well-known potentially dangerous problems, Polonyi problem or, in general, moduli problem. There are two aspects which we should worry about before our discussion on decay of moduli into gravitinos. One is the case when the moduli decay very late (i.e. during or after the nucleosynthesis) which can jeopardize nucleosynthesis predictions because of ultra-relativistic decay products directly from moduli fields destroying the light elements, in particular, ^{4}He and D nuclei. The other is when the entropy release due to its decay is so big that it can over-dilute the baryon asymmetry well below its acceptable amounts (so-called 'entropy crisis'). Our model does not have either of these problems because the radiative correction raises its mass to as much as intermediate scale. Its decay width is indeed enhanced up to $\Gamma_{\phi_1} \simeq m_{\phi_1}^5/|F|^2 \simeq \frac{(\mu^2\alpha_\kappa/\lambda)^5}{\mu^4} \simeq \alpha_\kappa^5\mu$ with $\alpha_\kappa \equiv \kappa^2/4\pi \simeq O(10^{-1})$ and this is of order$10^{-13}M_p$. So O'Raifeartaigh field decays around $10^{13}M_p^{-1}$ in our model which is much before the nucleosynthesis starts around $\sim 10^{40}M_p^{-1}$ and even well before the reheating starts due to the inflaton decay around $1/\Gamma_\sigma \sim M_p^5/\Delta^6 \sim 10^{25}M_p^{-1}$. We however need a great care about the possibility of decay products with long life-time, especially gravitinos. Gravitino decay rate is of order $\Gamma_{m_{3/2}} \sim m_{3/2}^3/M_p^2 \sim \mu^6/M_p^5 \sim 10^{-48}M_p$ and its relativistic decay products, especially ultra-relativistic photon/photino, can destroy the light elements in nucleosynthesis (photo-dissociation process) as we just mentioned. The possible abundant gravitino production from O'Raifeartaigh fields can be caused by the energy release stored during the inflation by the shift of the minimum of SUSY breaking field potential as inflaton evolves. If this energy release is too big, it could lead to large amount of gravitinos and upset the nucleosynthesis predictions. We can see this is not the case for our model as follows [1,4,20]. During the inflation, due to the coupling to the inflaton ($\sigma \sim M_p$), O'Raifeartaigh field amplitude is around the intermediate scale of order $\phi_1 \simeq \frac{\mu^2}{\Delta^2}M_p$ at the minimum of its potential. Therefore we can estimate the energy stored in this O'Raifeartaigh field when it starts oscillation (i.e. $t_{\phi_1} \simeq m_{\phi_1}^{-1}$) to be at most of order

$$\rho(t_{\phi_1}) \simeq m_{\phi_1}^2 \frac{\mu^4}{\Delta^4} M_p^2 \tag{3.1}$$

and its number density n_{ϕ_1} in this coherently oscillating O'Raifeartaigh field is at most

$$n_{\phi_1}(t_{\phi_1}) \simeq m_{\phi_1} \frac{\mu^4}{\Delta^4} M_p^2. \tag{3.2}$$

We can now estimate its number density to entropy ratio at the time of reheating for the gravitinos through the decay of ϕ_1 (at $t = t_r$, say) in an adiabatically expanding universe. Assuming, for the upper bound, ϕ_1 solely decays into gravitinos with 100% branching ratio and using $s \sim \frac{2\pi^2}{45} g_* T^3$ (with g_* effective degree of freedom) and $a^3 \sim t^2$ for matter domination [21] due to the coherently oscillating inflaton field which dominates the energy in the universe,

$$\frac{n_{3/2}(t_r)}{s(t_r)} \simeq \frac{n_{\phi_1}(t_{\phi_1})\left(\frac{a(t_{\phi_1})}{a(t_r)}\right)^3}{0.44 g_* T_{RH}^3} \simeq \frac{n_{\phi_1}(t_{\phi_1})\left(\frac{t_{\phi_1}}{t_r}\right)^2}{0.44 g_* T_{RH}^3}. \tag{3.3}$$

This can lead to the estimate of $n_{3/2}/s$ after reheating by substituting $1.66 \cdot T_{RH}^2\sqrt{g_*}/M_p \sim H \sim t_r^{-1}$ for t_r,

$$\frac{n_{3/2}}{s} \simeq \frac{n_{\phi_1}(t_{\phi_1}) t_{\phi_1}^2 T_{RH}(1.66)^2}{0.44 M_p^2} \simeq \frac{(1.66)^2 \mu^4 T_{RH}}{0.44 m_{\phi_1}\Delta^4}. \tag{3.4}$$

We can now compare this value with one corresponding to the gravitinos produced by the scattering in the thermal bath in reheating era obtained in MSSM [18,19],

$$n/n_{rad}(T \ll 1\mathrm{MeV}) \simeq 1.1 \cdot 10^{-11}\left(\frac{T_{RH}}{10^{10}\mathrm{GeV}}\right). \tag{3.5}$$

Using $s = 1.8 \cdot g_* n_{rad}$ and $g_*(\ll MeV) \simeq 3.36$, we obtain

$$n/s \simeq 1.8 \cdot 10^{-12}\left(\frac{T_{RH}}{10^{10}\mathrm{GeV}}\right). \tag{3.6}$$

flow of the gauge coupling. We can now transform eqn(3.4) by substituting (2.16) to the following form,

$$n/s \simeq 3.5 \cdot 10^{-14}\left(\frac{T_{RH}}{10^{10}\mathrm{GeV}}\right). \tag{3.7}$$

This is smaller than the thermal production of gravitino (3.6) by two orders of magnitude. The radiative correction therefore induces intermediate mass scale for O'Raifeartaigh field and it consequently makes the O'Raifeartaigh field energy released through the decay into gravitinos after the inflation small enough to evade the abundant gravitinos. Hence our model does not suffer from moduli problem as far as the constraint from thermal gravitino production is satisfied.

We mention that non-adiabatic production of moduli fields in pre-heating era could lead to abundant gravitinos [25]. We are, however, not concerned about the parametric resonance effects for moduli fields because scalar coupling terms in Lagrangian in our model are trilinear in hidden sector fields, and those couplings to inflaton field have always Planck mass suppression [5].

The exception where this Planck mass suppression does not occur and preheating effects could be important is the coupling involving longitudinal components of gravitino, which is the subject in the following section.

B. Non-thermal Production of Gravitino

It has been argued recently that parametric resonance mechanism in preheating era for the creation of gravitino can be much more efficient than the thermal one [6–10]. In this nonperturbative mechanism, the gravitinos can be created non-adiabatically through the amplification of vacuum fluctuation via rapid energy transfer from coherently oscillating inflaton field which still dominates the energy density in the universe just after inflation and before the reheating era.

In analyzing the gravitino field equations in the following, we see that the equations for transverse and longitudinal components of gravitino decouple. While transverse component equation has a Planck mass suppressed coupling and thus gravitationally suppressed particle creation, longitudinal component equation is free from Planck mass suppression and it could lead to the abundant gravitino production well above the constraint from thermal production of gravitino. Indeed, gravitino-goldstino equivalence theorem states that the equation for gravitino longitudinal component can be reduced to the equation of goldstino in global supersymmetry in the limit of weak gravitational coupling. This warns us that gravitino longitudinal components could lead to its efficient copious production without Planck mass suppression. Our model however has a Planck scale amplitude for inflaton field after inflation, and it is not obvious if this naive intuitive picture analogous to the goldstinos in global SUSY applies here. Therefore we apply here the formalism developed in [6,7] to calculate the number density of gravitinos created through the nonthermal process.

We first need to describe the evolution of scalar fields and fermion fields and their interactions. It is convenient to work with, among other possible choices, the following rescaled quantities in our numerical analysis,

$$\hat\phi_1 \equiv \frac{\phi_1}{M_p},\quad \hat\phi_2 \equiv \frac{\phi_2}{M_p},\quad \hat\phi_3 \equiv \frac{\phi_3}{\Delta},\quad \hat\sigma \equiv \frac{\sigma}{M_p},\quad \hat\mu \equiv \frac{\mu}{\Delta},$$
$$\hat\lambda \equiv \frac{\lambda}{\Delta},\quad \hat t \equiv t\,\frac{\Delta^2}{M_P},\quad \hat H \equiv H\frac{M_p}{\Delta^2},\quad \hat V \equiv \frac{V}{\Delta^4}, \tag{3.8}$$

where H is Hubble constant, V is a scalar potential from eqn (2.3) with non-minimal Kähler potential obtained in (2.15),

$$K = \Sigma\Sigma^\dagger + \Phi_1\Phi_1^\dagger + \Phi_2\Phi_2^\dagger + \Phi_3\Phi_3^\dagger - \frac{\kappa^2}{32\pi^2}\left(\frac{\kappa^2\Phi_1^2\Phi_1^{\dagger 2}}{\lambda^2}\right). \tag{3.9}$$

In terms of these rescaled quantities, the equations of motion for coherently oscillating scalar fields $\phi(=\sigma,\phi_i)$ read

$$\frac{d^2\hat\phi}{d\hat t^2} + 3\,\hat H\,\frac{d\hat\phi}{d\hat t} + \frac{d\hat V}{d\hat\phi} = 0\ . \tag{3.10}$$

We omit $\hat{}$ in the following discussion as long as it is clear from the contexts. We can concentrate on the field equations for σ and ϕ_1 because the other fields in supersymmetry breaking sector quickly roll down to the origin during the inflaton and stay there[‡]. So we can let the amplitudes of ϕ_2 and ϕ_3 vanish after obtaining the equation of motion for σ and ϕ_1 to see the field evolutions after the inflation.

The Fermion equation follows from the supergravity Lagrangian

$$\begin{aligned} e^{-1}L = &-\frac{1}{2}\,M_p^2\,R - K_i{}^j\,(\partial_\mu\,\phi^i)\,(\partial^\mu\phi_j) - V \\ &-\frac{1}{2}\,M_p^2\,\bar\psi_\mu\,R^\mu + \frac{1}{2}\,m\,\bar\psi_{\mu R}\,\gamma^{\mu\nu}\,\psi_{\nu R} \\ &+\frac{1}{2}\,m^*\,\bar\psi_{\mu L}\,\gamma^{\mu\nu}\,\psi_{\nu L} - K_i{}^j\,[\bar\chi_j\,\not{D}\bar\chi^i + \bar\chi^i\,\not{D}\bar\chi_j] \\ &-m^{ij}\,\bar\chi_i\,\chi_j - m_{ij}\,\bar\chi^i\,\chi^j \\ &+(2\,K_j{}^i\,\bar\psi_{\mu R}\,\gamma^{\nu\mu}\,\chi^j\,\partial_\nu\phi_i + \bar\psi_R\cdot\gamma\upsilon_L + \text{h.c.}) \\ &+(\text{ four fermion and gauge interaction terms }). \end{aligned} \tag{3.11}$$

This Lagrangian includes chiral complex multiplets (ϕ_i, χ_i) and the Ricci scalar R. Subscript L and R denote its projection through operators $P_L \equiv (1+\gamma_5)/2$, $P_R \equiv (1-\gamma_5)/2$. Gravitino kinetic term shows up in the form of

$$R^\mu = e^{-1}\,\epsilon^{\mu\nu\rho\sigma}\,\gamma_5\,\gamma_\nu\,D_\rho\,\psi_\sigma\,, \tag{3.12}$$

with covariant derivative

$$D_\mu\psi_\nu = \left(\left(\partial_\mu + \frac{1}{4}\omega_\mu^{mn}\gamma_{mn}\right)\delta_\nu^\lambda - \Gamma^\lambda_{\mu\nu}\right)\psi_\lambda\,. \tag{3.13}$$

The kinetic term for chiral fermion is

$$\begin{aligned} D_\mu\chi_i \equiv &\left(\partial_\mu + \frac{1}{4}\omega_\mu^{mn}\gamma_{mn}\right)\chi_i \\ &+\frac{1}{4M_p^2}\left[\partial_j K\partial_\mu\phi^j - \partial^j K\,\partial_\mu\phi_j\right]\chi_i + \Gamma_i^{j\,k}\chi_j\partial_\mu\phi_k \end{aligned} \tag{3.14}$$

[‡]Once these fields roll down to the origin, they stay at the origin to any higher order because of R-symmetry.

with Kähler connection $\Gamma_i^{j\,k} \equiv K^{-1}{}_i{}^l \partial^j K_l{}^k$ and $\gamma_{m\,n} \equiv [\gamma_m, \gamma_n]/2$. Its mass term reads

$$m^{ij} \equiv D^i D^j m = \left(\partial^i + \frac{1}{2M_p^2}(\partial^i K)\right) m^j - \Gamma_k^{i\,j}\, m^k \,. \tag{3.15}$$

The combination of matter fields gives left-handed component of goldstino

$$\upsilon_L \equiv m^i\, \chi_i + (\not\partial \phi_i)\, \chi^j\, K_j{}^i \,. \tag{3.16}$$

The supersymmetry transformation of goldstino [6]

$$\delta \upsilon = -\frac{3M_p^2}{2}(m_{3/2}^2 + H^2)\, \epsilon,\ m_{3/2} \equiv \frac{|m|}{M_p^2} \tag{3.17}$$

indicates that gravitino mass and Hubble parameter signal supersymmetry breaking. We can obtain the gravitino equation from this Lagrangian,

$$\not D \psi_\mu + m \psi_\mu = \left(D_\mu - \frac{m}{2}\gamma_\mu\right) \gamma^\nu\, \psi_\nu. \tag{3.18}$$

In solving this gravitino equation of motion, we gauge away the goldstino(unitary gauge) and use plane-wave ansatz for the spatial dependence of $\psi_\mu \sim e^{ik\cdot x}$. Moreover it is convenient to decompose the space component of gravitino field into the transverse part ${\psi_i}^T$ and trace parts $\theta \equiv \gamma^i\, \psi_i$ and $k_i \psi_i$ as

$$\psi_i = \psi_i^T + (P_\gamma)_i\; \theta + (P_k)_i\; k_i \psi_i, \tag{3.19}$$

where

$$\begin{aligned}(P_\gamma)_i\; &\equiv \frac{1}{2}\left(\gamma^i - \frac{1}{\vec{k}^2}\, k_i\, (k_j\, \gamma^j)\right),\\ (P_k)_i\; &\equiv \frac{1}{2\, \vec{k}^2}\left(3\, k_i - \gamma_i\, (k_j\, \gamma^j)\right).\end{aligned} \tag{3.20}$$

This leads to the the following succinct form of dynamical field equations which describe the degree of freedom corresponding to transverse and longitudinal components,

$$\left[\gamma^0\, \partial_0 + i\, \gamma^i\, k_i + \frac{\dot a\, \gamma^0}{2} + \frac{a\, \underline{m}}{M_p^2}\right] {\psi_i}^T = 0\,, \tag{3.21}$$

$$\left(\partial_0 + \hat B + i\, \gamma^i\, k_i\, \gamma^0\, \hat A\right) \theta - \frac{4}{\alpha\, a}\, k^2\, \Upsilon = 0\,, \tag{3.22}$$

where

$$\begin{aligned}\Upsilon &= K_j{}^i\, \left(\chi_i\, \partial_0\, \phi^j + \chi^j\, \partial_0\, \phi_i\right)\\ \underline{m} = P_R\, m + P_L\, m^*\,, &\qquad |m|^2 = \underline{m}^\dagger\, \underline{m}\\ \hat A &= \frac{1}{\alpha}\left(\alpha_1 - \gamma^0\, \alpha_2\right)\\ \hat B &= -\frac{3}{2}\, \dot a\, \hat A + \frac{1}{2\, M_p^2}\, a\, \underline{m}\, \gamma^0\, \left(1 + 3\, \hat A\right)\\ \alpha &= 3\, M_p^2\, \left(H^2 + \frac{|m|^2}{M_p^4}\right)\\ \alpha_1 = -M_p^2\left(3\, H^2 + 2\, \dot H\right) - \frac{3}{M_p^2}\, |m|^2,&\ \ \alpha_2 = 2\dot{\underline{m}}^\dagger.\end{aligned} \tag{3.23}$$

We can easily see, reducing the equation into this form, eqn(3.21) describing the transverse component of gravitino ψ_i^T is decoupled from the longitudinal gravitino component, and its coupling to the other fields are Planck mass suppressed. So we hereafter pay our attention to the equation which describes the longitudinal component of gravitino, eqn(3.22). The form of Υ in (3.23) tells us that, in the absence of Kähler terms which mix the various left chiral superfields, we need only worry about the fermionic partners of dynamical scalar fields. Furthermore, for our superpotential, there is no mixing between the fermion associated with ϕ_1 and those of ϕ_2 and ϕ_3, as long as $\phi_2 = \phi_3 = 0$ which is true because they stay at the origin due to R-symmetry once they roll down to the origin during the inflation. Thus, even though the effective masses of the fermions corresponding to ϕ_2 and ϕ_3 are changing, those fermions do not contribute to the goldstino and we can concentrate on the evolution of the other fields for the purpose of our calculation.

Based on the form of equation of motion involving two chiral superfields, we can infer the terms in the Lagrangian which describe the interactions of the two fields of our interests, namely, θ (longitudinal component of gravitino) and Υ (combination of chiral fermions orthogonal to goldstino υ). Those interaction terms lead to the equation of motion in the following matrix form,

$$\left(\gamma^0\, \partial_0 + i\, \gamma^i\, k_i N + M\right) X = 0\,, \tag{3.24}$$

with vector $X \equiv \binom{\tilde\theta}{\tilde\Upsilon}$ consisting of canonically normalized fields

$$\begin{aligned}\theta &= \frac{2\, i\, \gamma^i\, k_i}{(\alpha\, a^3)^{1/2}}\, \tilde\theta\,,\\ \Upsilon &= \frac{\Delta}{2}\left(\frac{\alpha}{a}\right)^{1/2} \tilde\Upsilon,\end{aligned} \tag{3.25}$$

and diagonal mass matrix M is given by

$$M = \text{diag}\Big(\quad \frac{ma}{2M_p^2} + \frac{3}{2}\left(\frac{ma}{M_p^2}\alpha_1 + \dot{a}\alpha_2\right),$$

$$-\frac{ma}{2M_p^2} + \frac{3}{2}\frac{ma}{M_p^2}\tilde{\alpha_1} + \dot{a}\tilde{\alpha_2} + a(m_{11}+m_{22}) \quad \Big) \qquad (3.26)$$

and matrix N

$$N \equiv \begin{pmatrix} -\tilde{\alpha_1} & 0 \\ 0 & -\tilde{\alpha_1} \end{pmatrix} + \gamma^0 \begin{pmatrix} -\tilde{\alpha_2} & -\Delta \\ -\Delta & \tilde{\alpha_2} \end{pmatrix} \qquad (3.27)$$

for $\tilde{\alpha}_i \equiv \alpha_i/\alpha$ and $\Delta = \sqrt{1-\tilde{\alpha}_1^2-\tilde{\alpha}_2^2}$. The existence of off-diagonal terms warns us the non-trivial mixing of fermion eigenstates.

Once we can reduce the Lagrangian into this form of matrix expression, we can obtain the evolution equations for the mode functions (the function multiplying the creation/annihilation operator) U_r^{ij} and V_r^{ij} (i,j runs over 1 and 2 for two field case and r for helicity $\pm$) and calculate the occupation number of gravitino created from vacuum through these mode functions. In general, one is interested in the physical mass eigenstates (ψ_1, ψ_2) which are non-trivial combinations (with matrix coefficients) of gravitino θ and matter chiral fermions Υ. So, strictly speaking, one would need to diagonalize the Hamiltonian at each moment of field evolution to keep track of the mass eigenstates and their abundance. This diagonalization process is rather involved[§]. Here we use a further simplification for our numerical analysis because we are only interested in the asymptotic value of these abundances. Indeed, since the mixing is small at such late times, we can simply follow the fields of interest (θ, Υ). That is, at late times, these fields are approximate mass eigenstates and, therefore, their occupation numbers correspond to those of (ψ_1, ψ_2). The validity of this approximation will be confirmed if we find that our occupation numbers cease to evolve at the time scales of interest.

With this simplification, we may represent the mode decomposition in the following familiar form,

$$X^i\,(x) = \int \frac{d^3\mathbf{k}}{(2\pi)^{3/2}} e^{i\mathbf{k}\cdot\mathbf{x}} \left[U_r^{ij}\,(k,\,\eta)\, a_j^r\,(k) + V_r^{ij}\,(k,\,\eta)\, b_j^{\dagger r}\,(-k) \right]. \qquad (3.28)$$

We then define the spinor matrix U_- and U_+

$$U_r^{ij} \equiv \left[\frac{U_+^{ij}}{\sqrt{2}}\,\psi_r, \frac{U_-^{ij}}{\sqrt{2}}\,\psi_r \right]^T, \qquad V_r^{ij} \equiv \left[\frac{V_+^{ij}}{\sqrt{2}}\,\psi_r, \frac{V_-^{ij}}{\sqrt{2}}\,\psi_r \right]^T \qquad (3.29)$$

with eigenvectors of the helicity operator $\sigma\cdot\mathbf{v}/|\mathbf{v}|$, $\psi_+ = \binom{1}{0}$ and $\psi_- = \binom{0}{1}$. Using these spinor matrices, the field equation of motion (3.24) has a following simple form in terms of the matrices U_+ and U_-,

$$a(t)\dot{U_\pm} = -i\,k\,U_\mp \mp i\,M\,U_\pm. \qquad (3.30)$$

We can then expand $U_\pm$ in terms of positive and negative frequency solutions,

$$\begin{aligned} U_+(t) &\equiv \left(1+\frac{M}{\omega}\right)^{1/2} \mathrm{e}^{-\mathrm{i}\int^{\mathrm{t}}\omega\,\mathrm{dt'}}\,\mathrm{A} - \left(1-\frac{M}{\omega}\right)^{1/2} \mathrm{e}^{\mathrm{i}\int^{\mathrm{t}}\omega\,\mathrm{dt'}}\,\mathrm{B} \\ &\equiv \left(1+\frac{M}{\omega}\right)^{1/2}\alpha - \left(1-\frac{M}{\omega}\right)^{1/2}\beta\,, \\ U_-(t) &\equiv \left(1-\frac{M}{\omega}\right)^{1/2} \mathrm{e}^{-\mathrm{i}\int^{\mathrm{t}}\omega\,\mathrm{dt'}}\,\mathrm{A} + \left(1+\frac{M}{\omega}\right)^{1/2} \mathrm{e}^{\mathrm{i}\int^{\mathrm{t}}\omega\,\mathrm{dt'}}\,\mathrm{B} \\ &\equiv \left(1-\frac{M}{\omega}\right)^{1/2}\alpha + \left(1+\frac{M}{\omega}\right)^{1/2}\beta\,, \end{aligned} \qquad (3.31)$$

where diagonal matrix $\omega \equiv \sqrt{k^2+M^2}$. α and β are precisely the generalization of Bogolubov coefficients. Indeed, in the same way as Bogolubov coefficients, we can calculate the occupation number of ith fermion eigenstates from β as

$$N_i\,(t) = \left(\beta^*\beta^T\right)_{ii} \;\; (\text{ no summation for } i\ ). \qquad (3.32)$$

We also keep in our mind that, because of nontrivial mixing of fermion mass eigenstates for the case of coupled field system, we need an extra care about the identification of inflatinos and gravitinos.

We solved the coupled mode equations (3.30) numerically to obtain the occupation numbers for N_θ and N_υ. These are plotted in Figure (1) as a function of comoving momentum at time 1000 in units of inflaton mass $m_\sigma \sim \frac{\Delta^2}{M_p}$ which gives a typical time scale for inflaton oscillations. We have used the typical parameter values $\hat{\mu} = 0.0001, \hat{\lambda} = 0.001$ with an initial inflaton amplitude $0.2M_p$. The O'Raifeartaigh field Φ_1 has an initial amplitude $\hat{\mu}^2 M_p$ and we normalized $a(t)$ to be one at the start of our calculation.

[§] We refer the readers to [7] for the general discussion.

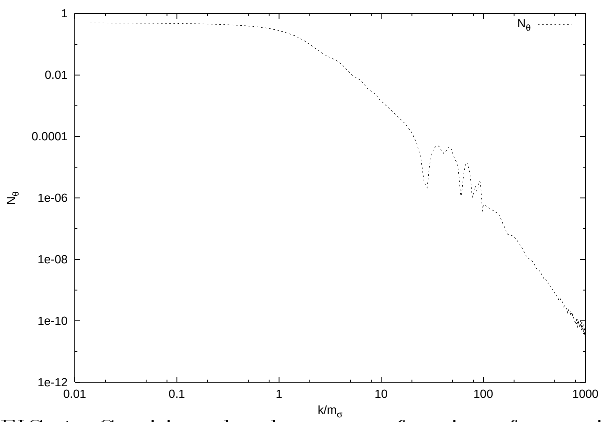

FIG. 1. Gravitino abundance as a function of comoving momentum in units of inflaton mass.

N_θ in our plot corresponds to the field abundance whose mass converges asymptotically to gravitino mass $|m|/M_p^2$ which shows up in the first element of the mass matrix M in eqn(3.26), and it gives an estimate of gravitino abundances produced in preheating era. We point out that, in the presence of time dependent background as usually the case in dealing with cosmological problems, kinetic term of scalar fields can cause the supersymmetry breaking as we can easily see from the supersymmetry transformation of chiral fermion, f_χ(superpartner of scalar χ, say),

$$\delta f_\chi = -\frac{1}{2} P_L \left[m_\chi - \frac{\gamma^0}{\sqrt{2}} \frac{d\chi}{dt} \right] \epsilon \quad , \tag{3.33}$$

where m_χ is defined in eqn(2.5). Therefore we should be aware that the values for N_1 and N_2 at intermediate time (i.e. the time when σ and ϕ_1 are still far from its settlement in the vacuum) does not represent either of the inflatino or gravitino abundance because there exist non-negligible contributions of supersymmetry breaking from both inflaton and hidden sectors.

The time when this SUSY breaking contribution of $\dot{\sigma}$ and m_σ becomes comparable with that of supersymmetry breaking sector is beyond the range of our numerical integration. We checked, however, that in such a small parameter range with so small initial amplitude for ϕ_1 as in our model, N_1 and N_2 converge to asymptotic values and do not change anymore at relatively early stage even when σ still keeps its oscillation. This verifies that Fig.1 should represent a good asymptotic behavior for gravitino abundance. Especially, the cut-off scale ($k \sim m_\sigma$) of the comoving momentum for the number density does not change anymore. This reflects the fact that the oscillation scale of high momentum mode ($k \gtrsim m_\sigma$) is much larger than that of the background field ($m_{3/2}$), so high momentum modes behave adiabatically and do not result in the non-adiabatic amplification anymore in preheating era.

We further comment on the subtle problems in identification of gravitinos and inflatinos in terms of mass eigenstates. Whatever field which dominates the local supersymmetry breaking is considered to be longitudinal components of gravitinos via super Higgs mechanism. The non-thermal production of fermionic fields is efficient just after inflation when the kinetic term of inflaton field governs the supersymmetry breaking in our model. This is nothing but the non-adiabatic field amplifications of longitudinal components of gravitinos which 'eats' the fermionic partner of inflaton field, i.e. *inflatino*. Hence when the fermion preheating is robust just after inflation, it is to-be inflatino, not to-be gravitino of our interests, that is amplified via parametric resonance effects as longitudinal components of gravitinos free from Planck suppression.

We also should mention that, because of the couplings, hidden sector may not be the sole cause for supersymmetry breaking even in the vacuum, and in fact, there could be still supersymmetry breaking from inflaton sector in the vacuum at a later time. It, however, can be shown that in the vacuum the supersymmetry contribution from inflaton sector is at most of order $|F|^4$ compared with $|F|^2$ ($\sim \mu^4$) due to hidden sector [23,24], and we still observe the dominant contribution of local supersymmetry breaking from the O'Raifeartaigh field at a later time. We thus expect our plot of N_θ still represents a fairly good overall behavior of gravitino abundance in asymptotic regime.

For the comparison with the gravitino number density constraints from photo-dissociation process in nucleosynthesis for the case of thermal production of gravitinos in thermal bath (3.6), we need to integrate $N_\theta(k)$ over the comoving momentum space. As usually the case with the preheating of fermions, our plot also indicates that occupation number as a function of comoving momentum k can be as large as of order unity at most up to the order of inflaton mass scale, $k_{cutoff} \simeq m_\sigma$ and decreases exponentially for bigger k. So the number density for longitudinal components,

$$n_{3/2} = \frac{1}{\pi^2} \frac{1}{a^3} \int_0^{k_{max}} |\beta_k|^2 k^2 dk \tag{3.34}$$

during the preheating is at most

$$n_{3/2}(t_{pre}) \lesssim k_{cutoff}^3 \simeq m_\sigma^3 \simeq \frac{\Delta^6}{M_p^3} \simeq 10^{-25} {M_p}^3. \tag{3.35}$$

We can now estimate the upper bound of the ratio of gravitino number density $n_{3/2}$ to entropy density in analogy with (3.4).

$$\frac{n_{3/2}(t_r)}{s(t_r)} \simeq \frac{n_{3/2}(t_{pre})\left(\frac{a(t_{pre})}{a(t_r)}\right)^3}{0.44 g_* T_{RH}^3} \simeq \frac{n_{3/2}(t_{pre})\left(\frac{t_{pre}}{t_r}\right)^2}{0.44 g_* T_{RH}^3}, \tag{3.36}$$

and substitution of $1.66 \cdot T_{RH}^2 \sqrt{g_*}/M_p \sim H \sim t_r^{-1}$ for t_r gives us the estimate of $n_{3/2}/s$ after reheating,

$$\frac{n_{3/2}}{s} \simeq \frac{n_{3/2}(t_{pre}) t_{pre}^2 T_{RH} (1.66)^2}{0.44 M_p^2}. \tag{3.37}$$

We expect this efficient gravitino production occurs well within the time range of typical oscillation of supersymmetry breaking field and we can substitute $t_{pre} \sim 1/m_{\phi_1} \sim (\alpha_\kappa \mu)^{-1} \sim 10^9 M_p^{-1}$ and eqn (3.35) in above equation to obtain the upper bound,

$$\frac{n_{3/2}}{s} \lesssim 6.3 \cdot 10^{-15} \left(\frac{T_{RH}}{10^{10}\mathrm{GeV}}\right). \tag{3.38}$$

This upper bound of n/s for the gravitinos from nonthermal production in our model is thus smaller than eqn(3.6) of thermal scattering by at least two orders of magnitude.

We therefore find that our model does not lead to the overproduction of gravitinos due to nonthermal process in preheating period, and reheating temperature constraint due to this effect is less severe than that of gravitinos produced by the scattering in thermal bath during the reheating period. We also point out that the expression given by eqn(3.37) was derived in a general setting and it can be used to obtain, in combination with eqn(3.6), the estimate for the relative significance of the gravitino production in thermal and non-thermal processes once the model of supergravity inflation is given.

IV. CONCLUSION AND DISCUSSION

We showed in this letter a realistic supergravity inflation model which breaks local supersymmetry in the vacuum dominantly via F-term coming from O'Raifeartaigh field in the hidden sector. We emphasized the significance of radiative correction in supersymmetry breaking sector to evade the moduli problem, and subsequently obtained the non-minimal Kähler potential arising from this loop correction. Using this non-minimal form of Kähler potential, we analyzed the possible non-thermal production of gravitinos in preheating era. The emphasis in our analysis was on the longitudinal components of gravitinos which do not suffer from Planck suppression and hence potentially could lead to robust amplification via parametric resonance effects. We showed that the comoving number density of fermionic mass eigenstate converges to its asymptotic value at an early stage of preheating era which is well before the time when the supersymmetry contribution comes from hidden sector fields. Physically, this indicates that the large comoving modes ($k \gtrsim m_\sigma$) is much larger than the typical coherent oscillation of background fields at later times ($k \sim m_{3/2}$). Hence comoving number density for big modes behave adiabatically and we do not expect the parametric amplifications at later times. We also discussed the subtle problems in identification of gravitinos and inflatinos due to the non-trivial mixing of fermionic mass eigenstates. To-be *inflatinos* are longitudinal components of *gravitinos* just after inflation, and its role is replaced by fermionic partner of O'Raifeartaigh fields at later times.

We estimated the upper bound of number density of non-thermally produced gravitinos by integrating out its comoving occupation number over momentum space. Because of the small mass scale of inflaton field, the typical momentum scale of produced gravitinos and consequently the momentum space over which occupation number is integrated out turns out to be small as well. This leads to the relatively small number density of gravitinos and we showed that it gives less significant constraint than that of gravitinos which are produced by thermal scattering.

We point out that the cases involving three and more superchiral fields are rather involved. We can basically follow the formalism discussed in section III B, but we need additional care in interpreting the fermion fields as a superposition of mass eigenstates because we cannot completely gauge away one of fermion fields via unitary gauge as we can do in the two field case [6,7]. The case including the gauge interaction terms and the model of other supersymmetry breaking mechanism besides hidden sector supersymmetry breaking are also to be examined.

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