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Gravitino Overproduction in Inflaton Decay

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Abstract

Most of the inflation models end up with non-vanishing vacuum expectation values of the inflaton fields ϕ in the true vacuum, which induce, in general, nonvanishing auxiliary field G_ϕ for the inflaton potential in supergravity. We show that the presence of nonzero G_ϕ gives rise to inflaton decay into a pair of the gravitinos and are thereby severely constrained by cosmology especially if the gravitino is unstable and its mass is in a range of $O(100)$ GeV $\sim O(10)$ TeV. For several inflation models, we explicitly calculate the values of G_ϕ and find that most of them are excluded or on the verge of being excluded for the gravitino mass in that range. We conclude that an inflation model with vanishing G_ϕ , typically realized in a chaotic inflation, is favored in a sense that it naturally avoids the potential gravitino overproduction problem.



1 Introduction

The recent Wilkinson Microwave Anisotropy Probe (WMAP) three year data [1] is consistent with generic predictions in inflation theories and hence strongly supports the basic idea of inflationary universe. Now we are reaching the stage to select the specific inflation model that is favored by the observation.

In most of inflation models in supergravity (SUGRA) an inflaton field ϕ has a non-vanishing expectation value at the potential minimum [2]. This implies that the Kähler potential for the inflaton field ϕ contains linear terms of the inflaton field in the true vacuum, even if the minimal Kähler potential is assumed at the beginning. That is, the inflaton Kähler potential is written at the potential minimum as $K = c\phi + c^*\phi^\dagger + \phi\phi^\dagger + \dots$. With the linear terms the inflaton field ϕ generically has a nonvanishing auxiliary field G_ϕ , where $G = K + \ln|W|^2$ (here and in what follows a subscript i denotes a derivative with respect to the field i). Here, K and W are the Kähler potential and superpotential, respectively. Recently, Ref. [3] has pointed out that, in the context of the moduli problem [4], the nonvanishing auxiliary field enables the decay into a pair of the gravitinos to proceed with a rate much higher than previously thought [5]. We find this decay process is (more) important in the reheating process of the inflaton, and that a stringent constraint on the G_ϕ of the inflaton potential must be satisfied to avoid an overproduction of the gravitino keeping the success of the standard cosmology, especially when the gravitino is unstable and its mass is in the range of $O(100)$ GeV – $O(10)$ TeV. For several inflation models, we explicitly calculate the values of G_ϕ to exemplify how severe the bound is. Among the known inflation models, a class of chaotic inflation models naturally avoids the potential gravitino overproduction problem, since G_ϕ vanishes in the vacuum.

The new constraint on the inflaton potential G_ϕ depends on the gravitino mass. In this letter, we restrict our discussion to the case of the gravitino mass $O(1)$ TeV to make our point clear. We see how severe the new constraint is compared with the cosmological constraint on thermally produced gravitinos [6]. Discussion on more general cases of the gravitino mass is straightforward, which will be given in [7].

2 Inflaton decay into a pair of gravitinos

We estimate the decay rate of an inflaton field ϕ into a pair of gravitinos. To be concrete, we adopt the gravity-mediated SUSY breaking scenarios, in which the gravitino mass is almost constant during and after inflation. The decay process we consider is identical to that recently calculated for the modulus decay [3]. The relevant interactions are [8]

$$e^{-1}\mathcal{L} = -\frac{1}{8}\epsilon^{\mu\nu\rho\sigma}\left(G_\phi\partial_\rho\phi - G_{\bar{\phi}}\partial_\rho\phi^\dagger\right)\bar{\psi}_\mu\gamma_\nu\psi_\sigma - \frac{1}{8}e^{G/2}\left(G_\phi\phi + G_{\bar{\phi}}\phi^\dagger\right)\bar{\psi}_\mu[\gamma^\mu, \gamma^\nu]\psi_\nu, \quad (1)$$

where ψ_μ is the gravitino field, and we have chosen the unitary gauge in the Einstein frame with the Planck units, $M_P = 1$. The real and imaginary components of the inflaton field have the same decay rate at the leading order [3]:

$$\Gamma_{3/2} \equiv \Gamma(\phi \rightarrow 2\psi_{3/2}) \simeq \frac{1}{288\pi} \frac{|G_\phi|^2}{g_{\phi\bar{\phi}}} \frac{m_\phi^5}{m_{3/2}^2 M_P^2}, \quad (2)$$

where we have assumed that the inflaton mass is much larger than the gravitino mass: $m_\phi \gg m_{3/2}$, and $g_{ij^*} = K_{ij^*}$ is the Kähler metric. Note that $g_{\phi\bar{\phi}} = 1$ for a canonically normalized inflaton field. Thus the decay rate is enhanced by the gravitino mass in the denominator, which comes from the longitudinal component of ψ as emphasized in Ref. [3].

It should be noted that the above expression for the decay rate cannot be applicable for $H > m_{3/2}$ [9]. The decay proceeds only if the Hubble parameter H is smaller than the gravitino mass, since the chirality flip of the gravitino forbids the decay to proceed otherwise. Intuitively, the gravitino is effectively massless as long as $H > m_{3/2}$.

We should clarify another important issue: what is the longitudinal component of the gravitino (i.e. goldstino) made of? Similar issue was discussed in the context of the non-thermal ‘gravitino’ production during preheating [10], and it was concluded that the inflatino, instead of the gravitino in the low energy, was actually created [11] ^{#1}. The reason is that the ‘gravitino’ production occurs in a rather early stage of the reheating just after the inflation ends, during which the energy stored in the inflationary sector

^{#1}It should be noted, however, that the inflatinos produced during preheating may be partially converted to the gravitinos in the low energy, since G_ϕ is generically nonzero in the true minimum as we will show later. This effect can further constrain the inflation models, but the detailed discussion is beyond the scope of this letter.

significantly contributes to the total SUSY breaking. In our case, however, the situation is completely different; the decay into the gravitinos becomes effective, since we consider a cosmological epoch, $H < m_{3/2}$, when the SUSY breaking contribution of the inflaton is subdominant. Thus the gravitinos produced by the inflaton through the above decay process should coincide with those in the low energy. The gravitinos produced by inflaton decay are genuine, and thereby the gravitino overproduction problem is present.

3 Cosmological constraint on G_ϕ

We restrict ourselves to the case of unstable gravitino with mass of $O(1)$ TeV. This mass region of the gravitino is an interesting region for gravity-mediated SUSY breaking scenarios. In this case we have already a cosmological constraint on the reheating temperature T_R to avoid the overproduction of gravitinos through thermal scattering [6], that is #2,

$$T_R < O(10^6 \sim 10^8) \text{ GeV}, \quad (3)$$

depending on the gravitino mass and the hadronic branching ratio B_h . This may be easily achieved with sufficiently small couplings relevant for the inflaton decay. We assume that this bound on T_R is satisfied, which means the inflaton decay rate into the standard-model particles must satisfy #3

$$\Gamma_{SM} \simeq \left(\frac{\pi^2 g_*}{90} \right)^{\frac{1}{2}} \frac{T_R^2}{M_P} < O(10^{-6} \sim 10^{-2}) \text{ GeV}, \quad (4)$$

where $g_* \sim 200$ counts the relativistic degrees of freedom. We point out that too much gravitinos may be produced by the inflaton decay via the interaction (1), even if the above inequality is satisfied.

The gravitino-to-entropy ratio is given by

$$Y_{3/2} \simeq 2 \frac{\Gamma_{3/2}}{\Gamma_{SM}} \frac{3 T_R}{4 m_\phi}, \quad (5)$$

where we have neglected the gravitino production from thermal scattering. To keep the success of big-bang nucleosynthesis (BBN), the gravitino abundance must satisfy [6]

$$m_{3/2} Y_{3/2} < O(10^{-14} \sim 10^{-11}) \text{ GeV}. \quad (6)$$

#2 The expansion rate at the decay time is much smaller than the gravitino mass $O(1)$ TeV and hence the condition $H < m_{3/2}$ is satisfied.

#3 Throughout this letter we require $\Gamma_{3/2} \ll \Gamma_{SM}$, otherwise the standard cosmology would be upset.

From (4), (5) and (6), we obtain

$$\Gamma_{3/2} < O(10^{-23} \sim 10^{-17}) \left(\frac{m_{3/2}}{1 \text{ TeV}} \right)^{-1} \left(\frac{m_\phi}{10^6 \text{ GeV}} \right) \text{ GeV}, \quad (7)$$

or equivalently,

$$|G_\phi| < O(10^{-4} \sim 10^{-1}) \left(\frac{m_{3/2}}{1 \text{ TeV}} \right)^{\frac{1}{2}} \left(\frac{m_\phi}{10^6 \text{ GeV}} \right)^{-2}. \quad (8)$$

We show this constraint in Fig. 1 together with predictions of new and hybrid inflation models to be derived in the following sections. We can see that the hybrid inflation model is excluded, while the new inflation model is on the verge. It should be noted that, in deriving the constraint (8), we have substituted the upper bound on the reheating temperature (3). Therefore, the constraint on G_ϕ becomes severer for lower T_R , proportional to $T_R^{1/2}$.

4 Single field inflation model

In this section we estimate G_ϕ for an inflation model in which the inflaton sector consists of a single chiral superfield ϕ with nonzero vacuum expectation value (VEV) in the potential minimum. In the following we assume that the inflaton mass m_ϕ is much larger than the gravitino mass. Since G_ϕ corresponds to the fractional contribution to the SUSY breaking, it is estimated by minimizing the scalar potential with the hidden sector responsible for the SUSY breaking. We assume that the hidden sector contains a chiral superfield z and the superpotential is written as

$$W = W(z) + W(\phi), \quad (9)$$

where $W(z)$ and $W(\phi)$ are superpotentials for z and ϕ , respectively. For simplicity we take the minimal Kähler potential,

$$K = |z|^2 + |\phi|^2. \quad (10)$$

Then the scalar potential is given by

$$V = e^G (|G_z|^2 + |G_\phi|^2 - 3). \quad (11)$$

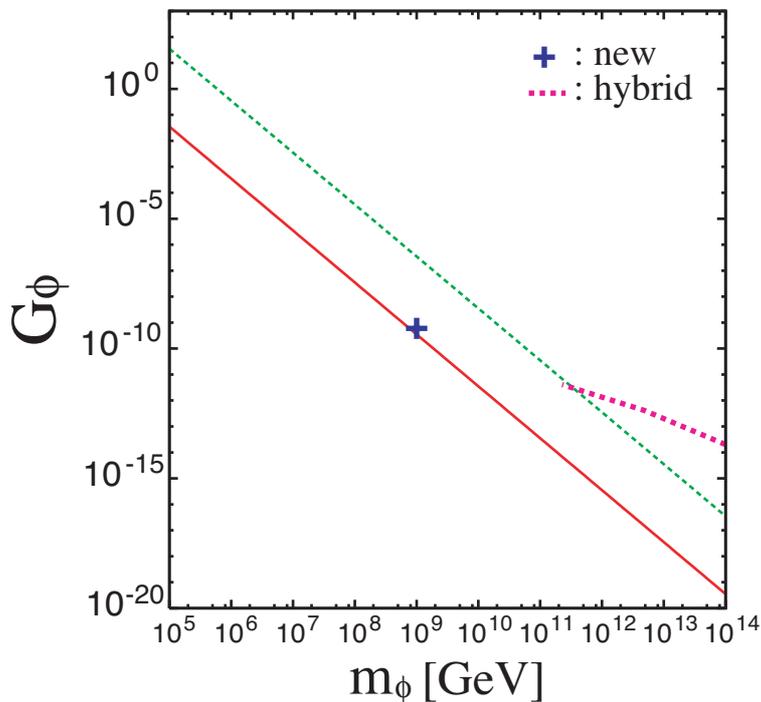


Figure 1: Upper bound on the auxiliary field of the inflaton G_ϕ as a function of the inflaton mass m_ϕ . The solid and dotted lines are for the hadronic branching ratio $B_h = 1$ and 10^{-3} , respectively. We set $m_{3/2} = 1$ TeV. The typical values of G_ϕ and m_ϕ for the new and hybrid inflation models discussed in the text are also shown.

Since the cosmological constant should vanish in the true vacuum,

$$|G_z|^2 + |G_\phi|^2 = 3. \quad (12)$$

The gravitino mass is given by $m_{3/2} = e^{G/2} \simeq |W|$. The potential minimum (=vacuum) is determined by a condition,

$$V_\phi = e^G (G_{z\bar{\phi}} G_{\bar{z}} + G_{\phi\bar{\phi}} G_{\bar{\phi}} + G_\phi) = 0, \quad (13)$$

where we have used Eq. (12) and $G_{i\bar{j}} = \delta_{ij}$, ($i, j = \phi, z$). Since z is responsible for the SUSY breaking, $|G_z| \simeq \sqrt{3}$ and $|W_z| \sim |W| \simeq m_{3/2}$. On the other hand, $|G_\phi| \ll 1$ for the inflaton field. Assuming $|G_\phi| = |\phi^\dagger + W_\phi/W| \lesssim |\phi|$, we have

$$W_\phi \sim \phi W. \quad (14)$$

The order estimation of Eq. (13) leads to

$$\phi + \left(\frac{m_\phi}{m_{3/2}} - \phi^2 \right) G_\phi + G_\phi \sim 0, \quad (15)$$

where we have omitted the coefficients of order unity, and used $W_{\phi\phi} \simeq m_\phi$. Thus we obtain the formula for G_ϕ in the single field inflation model,

$$\langle G_\phi \rangle \sim \frac{m_{3/2}}{m_\phi} \langle \phi \rangle, \quad (16)$$

for $m_\phi \gg m_{3/2}$. Note that this result satisfies the assumption $|G_\phi| \lesssim |\phi|$, so our analysis is consistent. Thus, the inflaton decay into gravitinos can set a constraint on a single field inflation model when the inflaton ϕ takes a non-vanishing expectation value after inflation.

For a concrete example, here we study a new inflation model [12, 13, 14]. In the new inflation model the superpotential of the inflaton sector is written as

$$W = v^2 \phi - \frac{\phi^5}{5}, \quad (17)$$

where $v = 10^{-7} \sim 10^{-6}$ for producing the observed density fluctuations. After inflation, the inflaton ϕ takes the expectation value $\sim \sqrt{v}$. In this model the gravitino mass is related to v as $m_{3/2} \sim v^{5/2}$, and the inflaton mass is given by $m_\phi \sim v^{3/2}$. Thus, Eq. (16) leads to $G_\phi \sim v^{3/2} \sim 6 \times 10^{-10}$ for $m_{3/2} = 1$ TeV and $m_\phi = 10^9$ GeV, which is close to the constraint shown in Fig. 1 ^{#4}.

5 Hybrid inflation model

In the previous section we estimate G_ϕ for single field inflation models. However, when the inflaton sector contains multiple superfields, Eq. (16) cannot be applied. Here, we consider a hybrid inflation model as a representative example.

The hybrid inflation model contains two kinds of superfields: one is ϕ which plays a role of inflaton and the others are waterfall fields Ψ and $\tilde{\Psi}$ [16, 17, 18]. After inflation

^{#4}The new inflation is also realized for $W = X(v^2 - g\phi^4)$ [15] for which the constraint can be relaxed. This is because the universe after inflation is dominated by ϕ with suppressed G_ϕ , while G_X is not suppressed. This should be contrasted with the case of the hybrid inflation model.

ends, ϕ as well as Ψ oscillates around the potential minimum and dominates the universe until the reheating.

The total superpotential W is written as

$$W = W(z) + W(\phi, \Psi, \tilde{\Psi}), \quad (18)$$

where the superpotential $W(\phi, \Psi, \tilde{\Psi})$ for the inflaton sector is

$$W(\phi, \Psi, \tilde{\Psi}) = \phi(\mu^2 - \lambda\tilde{\Psi}\Psi). \quad (19)$$

Here λ is a coupling constant and μ is the inflation energy scale. The potential minimum is located at $\langle\phi\rangle = 0$ and $\langle\Psi\rangle = \langle\tilde{\Psi}\rangle = \mu/\sqrt{\lambda}$ in the SUSY limit. Including the effect of the hidden sector, however, the minimum slightly shifts as shown below. For successful inflation, μ and λ are related as $\mu \simeq 2 \times 10^{-3} \lambda^{1/2}$ for $\lambda \gtrsim 10^{-3}$, and $\mu \simeq 2 \times 10^{-2} \lambda^{5/6}$ for $\lambda \lesssim 10^{-3}$. Moreover, in this type of hybrid inflation there exists a problem of cosmic string formation because Ψ and $\tilde{\Psi}$ generally have $U(1)$ gauge charges. To avoid the problem the coupling should be small as, $\lambda \sim 10^{-4}$ [19].

Now let us estimate G_ϕ and G_Ψ . The conditions for the potential minimum lead to

$$G_{z\phi}G_{\bar{z}} + G_{\phi\phi}G_{\bar{\phi}} + G_\phi + G_{\Psi\phi}G_{\bar{\Psi}} + G_{\tilde{\Psi}\phi}G_{\bar{\tilde{\Psi}}} = 0 \quad (20)$$

$$G_{z\Psi}G_{\bar{z}} + G_{\phi\Psi}G_{\bar{\phi}} + G_\Psi + G_{\Psi\Psi}G_{\bar{\Psi}} + G_{\tilde{\Psi}\Psi}G_{\bar{\tilde{\Psi}}} = 0. \quad (21)$$

Here we do not use the minimization condition for $\tilde{\Psi}$, since it is equivalent to Eq. (21). In the same way as deriving Eq. (14), we assume $|G_\Psi| \simeq |G_{\tilde{\Psi}}| \lesssim |\Psi|$, leading to $W_\Psi/W \sim \Psi$. Together with Eq. (19), we obtain

$$\langle\phi\rangle \sim \frac{m_{3/2}}{\lambda}, \quad (22)$$

where we have used $|W| \simeq m_{3/2}$ and $|\Psi| \simeq |\tilde{\Psi}|$. Then, with use of $G_{\Psi\phi} \sim \lambda\Psi/m_{3/2} - \Psi W_\phi/W$, $G_{z\Psi} \sim \Psi$, $G_{\Psi\Psi} \sim \Psi^2$ and $G_{\tilde{\Psi}\tilde{\Psi}} \sim \lambda\phi/m_{3/2} - \Psi^2$, Eqs. (20) and (21) are written as

$$\frac{W_\phi}{W} + \frac{W_\phi^2}{W^2}G_\phi + G_\phi + \left(\frac{\lambda\Psi}{m_{3/2}} - \Psi\frac{W_\phi}{W}\right)G_\Psi \sim 0 \quad (23)$$

$$\Psi + \left(\frac{\lambda\Psi}{m_{3/2}} - \Psi\frac{W_\phi}{W}\right)G_\phi + G_\Psi + \Psi^2G_\Psi + \left(\frac{\lambda\phi}{m_{3/2}} - \Psi^2\right)G_\Psi \sim 0, \quad (24)$$

where we omitted coefficients of order unity. Assuming $|G_\phi| \lesssim |\phi|$, we obtain

$$\langle G_\phi \rangle \sim \frac{m_{3/2}}{\lambda} \simeq \frac{m_{3/2}}{m_\phi} \langle \Psi \rangle, \quad (25)$$

$$\langle G_\Psi \rangle \simeq \langle G_{\bar{\Psi}} \rangle \sim \frac{m_{3/2}^2}{\lambda^2 \langle \Psi \rangle} \simeq \frac{m_{3/2}^2}{m_\phi^2} \langle \Psi \rangle. \quad (26)$$

Here we have used $W_\phi/W \sim \langle \phi \rangle \sim m_{3/2}/\lambda$ and $m_\phi = \lambda \langle \Psi \rangle$. One can easily check that the above results satisfy the assumptions we made on G_ϕ and G_Ψ . It should be noted that G_ϕ is much larger than G_Ψ . Therefore it is ϕ that produces too much gravitinos.

For $m_{3/2} = 1$ TeV and $\lambda \sim 1 - 10^{-4}$ we obtain $\mu \sim 2 \times 10^{-3} - 10^{-5}$, $G_\phi \sim 4 \times 10^{-16} - 4 \times 10^{-12}$ and $m_\phi \sim 5 \times 10^{15} - 2 \times 10^{11}$ GeV. From Fig. 1, one can see the hybrid inflation model is excluded by the gravitino overproduction. Although the constraint on G_ϕ becomes slightly mild for $\lambda \lesssim 10^{-4}$, it is then disfavored by the WMAP result. This is because the density fluctuation becomes almost scale-invariant for $\lambda \lesssim 10^{-4}$ while the spectral index is $n_s = 0.95 \pm 0.02$ according to the WMAP three year data [1].

So far we have considered the standard hybrid inflation, but the final results Eqs. (25) and (26) also apply to a smooth hybrid inflation model [20], which is favored compared to the hybrid inflation model in a sense that the predicted spectral index is smaller. The constraint on this model, however, is more or less similar to the hybrid inflation, and the smooth hybrid inflation is also excluded.

6 Chaotic inflation model

A chaotic inflation model [21, 22] is based on a Nambu-Goldstone-like shift symmetry of the inflaton chiral multiplet ϕ . Namely, we assume that the Kähler potential $K(\phi, \phi^\dagger)$ is invariant under the shift of ϕ ,

$$\phi \rightarrow \phi + iA, \quad (27)$$

where A is a dimensionless real parameter. Thus, the Kähler potential is a function of $\phi + \phi^\dagger$; $K(\phi, \phi^\dagger) = K(\phi + \phi^\dagger)$. We identify its imaginary part with the inflaton field φ . Moreover, we introduce a small breaking term of the shift symmetry in the superpotential in order for the inflaton φ to have a potential:

$$W = mX\phi, \quad (28)$$

where we introduced a new chiral multiplet X , and $m \simeq 10^{13}\text{GeV}$ determines the inflaton mass. Furthermore, we impose a discrete Z_2 symmetry which forbids the problematic linear term in Kähler potential such as $K = c(\phi + \phi^\dagger) + \dots$ [21, 22]. Then we take the minimal Kähler potential,

$$K = \frac{1}{2}(\phi + \phi^\dagger)^2 + |X|^2. \quad (29)$$

The superpotential (28) and Kähler potential (29) lead to the scalar potential,

$$V(\varphi, X) \simeq \frac{1}{2}m^2\varphi^2 + m^2|X|^2, \quad (30)$$

for $|X| < 1$. For $\varphi \gg 1$ and $|X| < 1$, the φ field dominates the potential and the chaotic inflation takes place [for details see Refs [21, 22]].

In this chaotic inflation model we see $G_\phi = 0$ since $\phi = 0$ and $X = 0$ in the vacuum. Including the hidden sector does not modify the potential minimum due to the Z_2 symmetry. Therefore, the new gravitino problem does not exist.

7 Conclusions

In this letter we have shown that an inflation model with a nonzero VEV $\langle\phi\rangle \neq 0$ generically leads to the gravitino overproduction, which can jeopardize the successful standard cosmology especially when the gravitino is unstable and its mass is in the region of $O(100)\text{GeV}$ and $O(10)\text{TeV}$. We have explicitly calculated G_ϕ , which is an important parameter to determine the gravitino abundance, for several inflation models. To be concrete we have fixed $m_{3/2} = 1\text{TeV}$ and shown that the new inflation is on the verge of being excluded, while the hybrid inflation model is excluded. The most attractive way to get around this new gravitino problem is to have the potential minimum at the origin as a class of chaotic inflation models does. Among the known models, such a chaotic inflation model is favored in a sense that it is free from the potential gravitino overproduction problem.

Here, let us comment on other solutions. Throughout this letter we have assumed that no entropy production occurs after the reheating completes. If the huge entropy is produced at late time [23], the new gravitino problem can be greatly relaxed. Another

even manifest solution is to assume the gravitino mass $m_{3/2} < \text{a few keV}$ ^{#5}. In this case, the produced gravitinos get into thermal equilibrium due to relatively strong interactions with the standard-model particles.

Finally, we shortly note how severe the new gravitino problem is for the gravitino mass other than $O(1)\text{TeV}$. We will see that it becomes milder for a mass region $m_{3/2} < O(1)$ GeV suggested from gauge mediation models [25], since the cosmological constraint is weaker; it comes from the requirement that the gravitino abundance should not exceed the present dark matter abundance. For a mass region $m_{3/2} = O(100)$ TeV suggested from anomaly mediation models [26], we will have a constraint more or less similar to the present result. The reason why we have the stringent constraint on G_ϕ for $m_{3/2} = O(100)$ GeV – $O(10)$ TeV is that the reheating temperature is so severely constrained by the BBN for this mass range. This means that the reheating temperature cannot be arbitrarily small even for the gravitino mass other than $O(1)\text{TeV}$, since the branching ratio of the direct gravitino production increases. The extended analysis including a broad mass region for the gravitino will be given in Ref. [7].

Note added: Very recently, Dine et al [27] have pointed out that the leading term of G_ϕ is cancelled out after taking into account a mass mixing between the inflaton and the hidden-sector field z and the effective G_ϕ^{eff} becomes much smaller if one takes the minimal Kähler potential. We agree with them in this point. However, as pointed out in Ref. [27], if the mass of the hidden sector field z is comparable or larger than the inflaton mass, such a cancellation does not occur. In the case of the new inflation model, the inflaton mass can be comparable to the mass of z , and our constraint in the text is then applicable. In addition, even if the hidden-sector field mass is smaller than the inflaton mass, G_ϕ may not be suppressed if the Kähler potential is non-minimal. In fact, it is quite natural to have a non-minimal Kähler potential in the gravity-mediated SUSY breaking model. If one introduces $\delta K = \kappa/2 |\phi|^2 z^2 + \text{h.c.}$, the effective G_ϕ^{eff} becomes $G_\phi^{\text{eff}} \sim \kappa m_{3/2} \langle \phi \rangle / m_\phi$. Thus, G_ϕ in the text should be replaced by this effective G_ϕ^{eff} . Notice that the hidden-sector field z has no charge in any symmetry and hence the above Kähler term has no

^{#5}Taking account of the constraints from the CMB and the structure formation, the upper bound on the gravitino mass can be reduced to $O(10)\text{eV}$ [24].

reason to be suppressed. We find that there is still a very stringent constraint on the hybrid inflation model unless the κ is extremely small. (For the hybrid inflation model, $G_{\phi}^{\text{eff}} \sim \kappa m_{3/2} \langle \Psi \rangle / m_{\phi}$ results from $\delta K = \kappa/2 |\Psi|^2 z^2 + \text{h.c.}$, where ϕ and Ψ are the inflaton and the water fall field, respectively). The detailed analysis will be given elsewhere.

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References

- [1] D. N. Spergel *et al.*, arXiv:astro-ph/0603449.
- [2] D. H. Lyth and A. Riotto, Phys. Rept. **314**, 1 (1999).
- [3] M. Endo, K. Hamaguchi and F. Takahashi, arXiv:hep-ph/0602061;
S. Nakamura and M. Yamaguchi, arXiv:hep-ph/0602081.
- [4] G. D. Coughlan, W. Fischler, E. W. Kolb, S. Raby and G. G. Ross, Phys. Lett. B **131** (1983) 59;
T. Banks, D. B. Kaplan and A. E. Nelson, Phys. Rev. D **49** (1994) 779;
B. de Carlos, J. A. Casas, F. Quevedo and E. Roulet, Phys. Lett. B **318** (1993) 447.
- [5] M. Hashimoto, K. I. Izawa, M. Yamaguchi and T. Yanagida, Prog. Theor. Phys. **100**, 395 (1998).
- [6] M. Kawasaki, K. Kohri and T. Moroi, Phys. Lett. B **625**, 7 (2005); Phys. Rev. D **71**, 083502 (2005).
- [7] M. Kawasaki, F. Takahashi and T.T. Yanagida, in preparation.
- [8] J. Wess and J. Bagger, Supersymmetry and Supergravity, (Princeton University Press, 1992).
- [9] M. Endo, in private communication.

- [10] R. Kallosh, L. Kofman, A. D. Linde and A. Van Proeyen, *Phys. Rev. D* **61**, 103503 (2000);
G. F. Giudice, I. Tkachev and A. Riotto, *JHEP* **9908**, 009 (1999);
G. F. Giudice, A. Riotto and I. Tkachev, *JHEP* **9911**, 036 (1999);
R. Kallosh, L. Kofman, A. D. Linde and A. Van Proeyen, *Class. Quant. Grav.* **17**, 4269 (2000) [Erratum-ibid. **21**, 5017 (2004)].
- [11] H. P. Nilles, M. Peloso and L. Sorbo, *Phys. Rev. Lett.* **87**, 051302 (2001); *JHEP* **0104**, 004 (2001).
- [12] K. Kumekawa, T. Moroi and T. Yanagida, *Prog. Theor. Phys.* **92**, 437 (1994).
- [13] K. I. Izawa and T. Yanagida, *Phys. Lett. B* **393**, 331 (1997).
- [14] M. Ibe, K. I. Izawa, Y. Shinbara and T. T. Yanagida, arXiv:hep-ph/0602192.
- [15] T. Asaka, K. Hamaguchi, M. Kawasaki and T. Yanagida, *Phys. Rev. D* **61**, 083512 (2000);
V. N. Senoguz and Q. Shafi, *Phys. Lett. B* **596**, 8 (2004).
- [16] E. J. Copeland, A. R. Liddle, D. H. Lyth, E. D. Stewart and D. Wands, *Phys. Rev. D* **49**, 6410 (1994).
- [17] G. R. Dvali, Q. Shafi and R. K. Schaefer, *Phys. Rev. Lett.* **73**, 1886 (1994).
- [18] A. D. Linde and A. Riotto, *Phys. Rev. D* **56**, 1841 (1997).
- [19] M. Endo, M. Kawasaki and T. Moroi, *Phys. Lett. B* **569**, 73 (2003).
- [20] G. Lazarides and C. Panagiotakopoulos, *Phys. Rev. D* **52**, 559 (1995).
- [21] M. Kawasaki, M. Yamaguchi and T. Yanagida, *Phys. Rev. Lett.* **85**, 3572 (2000).
- [22] M. Kawasaki, M. Yamaguchi and T. Yanagida, *Phys. Rev. D* **63**, 103514 (2001).
- [23] D. H. Lyth and E. D. Stewart, *Phys. Rev. D* **53**, 1784 (1996);
M. Kawasaki and F. Takahashi, *Phys. Lett. B* **618**, 1 (2005).
- [24] M. Viel, J. Lesgourgues, M. G. Haehnelt, S. Matarrese and A. Riotto, *Phys. Rev. D* **71**, 063534 (2005).
- [25] G. F. Giudice and R. Rattazzi, *Phys. Rept.* **322**, 419 (1999).

- [26] L. Randall and R. Sundrum, Nucl. Phys. B **557**, 79 (1999);
G. F. Giudice, M. A. Luty, H. Murayama and R. Rattazzi, JHEP **9812**, 027 (1998);
J. A. Bagger, T. Moroi and E. Poppitz, JHEP **0004**, 009 (2000).
- [27] M. Dine, R. Kitano, A. Morisse and Y. Shirman, arXiv:hep-ph/0604140.