Revised Version

DESY 06-242
FTPI-MINN-06/41
UMN-TH-2533/06

DESY 06-242
FTPI-MINN-06/41
UMN-TH-2533/06

hep-ph/0612263
December, 2006

The Decay of the Inflaton in No-scale Supergravity

Motoi Endo[1], Kenji Kadota[2], Keith A. Olive[2], Fuminobu Takahashi[1] and T. T. Yanagida[3,4]

[1] *Deutsches Elektronen Synchrotron DESY, Notkestrasse 85, 22607 Hamburg, Germany*

[2] *William I Fine Theoretical Physics Institute, University of Minnesota, Minneapolis, MN 55455, USA*

[3] *Department of Physics, University of Tokyo, Tokyo 113-0033, Japan*

[4] *Research Center for the Early Universe, University of Tokyo, Tokyo 113-0033, Japan*

Abstract

We study the decay of the inflaton in no-scale supergravity and show that decay due to the gravitational interactions through supergravity effects is highly suppressed relative to the case in minimal supergravity or models with a generic Kähler potential. We also show that decay to gravitinos is suppressed. We demonstrate that decay and sufficient reheating are possible with the introduction of a non-trivial gauge kinetic term. This channel may be dominant in no-scale supergravity, yet yields a re-heating temperature which is low enough to avoid the gravitino problem while high enough for Big Bang Nucleosynthesis and baryogenesis.

1 Introduction

Although cosmological inflaton provides an attractive solution to several problems including the horizon/flatness problem [1, 2], the realization of an exponentially expanding phase in the early universe in the framework of realistic particle physics models still remains an open question. As we expect supersymmetry and supergravity to play a fundamental role in these models, it is of interest to study inflation in this context as well [3]. However, in supergravity models with a minimal Kähler potential, all scalar fields generically obtain a supergravity mass correction of order the Hubble scale, and the slow-roll condition for the inflaton is in general violated, though specific models [4] can be realized. The slow-roll problem found in generic models of supergravity can be avoided if one uses a non-minimal Kähler potential of the no-scale form [5]. Several models of slow-roll inflation in no-scale supergravity have been considered [6, 7, 8, 9].

No-scale supergravity models based on a non-compact Kähler manifold, with a maximally symmetric coset space $SU(N,1)/[SU(N)\times U(1)]$, have attracted substantial interest from string theory because such a Kähler potential typically appears in the compactification of higher-dimensional superstring models [10] as well as from particle physics model building [11]. While local supersymmetry is broken by the no-scale structure of the Kähler potential, there is a residual global supersymmetry leading to semi-positive definite scalar potential [5, 11]. Consequently, one finds that the tree-level cosmological constant at the global minimum vanishes, in contrast to more generic supergravity models for which the global minimum possesses negative vacuum energy density [12]. Indeed, the tree-level potential is flat for the supersymmetry breaking, Polonyi-like field in no-scale models, and so, even though local supersymmetry is broken, the gravitino mass scale is undetermined. From the viewpoint of particle phenomenology, this offers an interesting explanation of the hierarchy problem through the radiative determination of weak/supersymmetry-breaking scale. A non-trivial gauge kinetic function can generate a gaugino mass which breaks global supersymmetry. Radiative corrections set the scale of the gaugino mass, and the form of the gauge kinetic function determines the ratio between the gravitino and gaugino masses.

In addition to the slow-role criteria, another essential feature of any successful inflationary model is sufficient reheating without the overproduction of gravitinos. So long as there is no direct superpotential coupling of the inflaton to matter fields, i.e. $W = W(\phi^1) + W(\phi^i)$, where ϕ^1 is the inflaton and the ϕ^i are matter superfields, the minimal decay rate proceeds via a 3-body gravitational decay with rate $\Gamma \propto (\langle\phi^1\rangle /M_P)^2\, m_{\phi^1}^3/M_P^2$ [13, 14], where $M_P/\sqrt{8\pi} \simeq 2.4 \times 10^{18}\,\mathrm{GeV}$ denotes the reduced Planck mass. This decay leads to a lower limit on the reheat temperature in generic supergravity models of order $T_{RH} \sim 10^6$ GeV for an inflaton mass of order 10^{12} GeV and an inflaton vacuum expectation value (VEV) $\langle\phi^1\rangle \simeq M_P$. While the thermal production of gravitinos at this temperature is not problematic, the direct decay of the inflaton imposes strong constraints on inflation models [15, 16]. These results also hold in large field inflation models such as chaotic [17] and hybrid [18] inflationary models.

In this letter, we will study the corresponding questions of reheating and gravitino production in no-scale supergravity models. In fact, we find that the decay of the inflaton is highly suppressed in no-scale supergravity, which is directly related to the special structure of the Kähler manifold in no-scale models. In particular, the tree-level gravitational interactions of the inflaton due to supergravity effects vanish exactly at the global minimum. As a consequence, without the direct coupling of the inflaton to matter (which may be problematic for model building) the inflaton is stable at the tree-level. Therefore, the suppression of the gravitational interactions of the inflaton due to the symmetry of Kähler manifold is of prime importance in inflation model building in supergravity.

We show, however, that the introduction of a non-trivial gauge kinetic function can lead to the decay of the inflaton with successful reheating [19] but without the overproduction of gravitinos. Because the dominant decay channel is specified by the gauge kinetic function in no-scale models, we can constrain the gauge kinetic function by the reheating temperate constraints. Interestingly, since the gauge kinetic function relates the gravitino and gaugino masses, the reheating temperature constraints provide an upper bound on the gravitino mass.

2 Setup: No-scale Supergravity

We consider a no-scale model with a Kähler potential of the form [5, 11][a] :

$$K = -3\ln\left[z + z^\dagger - \frac{1}{3}\left(\sum_{i=1}|\phi^i|^2\right)\right] \tag{1}$$

with a supersymmetry breaking field z, an inflaton ϕ^1, matter fields $\phi^i (i = 2, 3, ...)$ and as noted earlier, the superpotential is assumed to include no direct coupling between an inflaton and the other fields, $W = W(\phi^1)+W(\phi^i)$. We also assume that the superpotential does not contain z, so that the tree-level potential for z remains flat at the minimum (one of the notable features in a no-scale supergravity model). Unless explicitly noted, we will adopt Planck units so that $M_P/\sqrt{8\pi} = 1$. The total Kähler potential G is defined as $G \equiv K + F + F^\dagger$ with $F \equiv \ln W$.

The relevant bosonic kinetic terms are derived from $G_a^b(\partial_\mu\phi^a)(\partial^\mu\phi_b^*) = K_a^b(\partial_\mu\phi^a)(\partial^\mu\phi_b^*)$ for all scalar fields including z, the inflaton, and matter. Indices on the Kähler potential refer to derivatives with respect to the fields, $G_a = \partial G/\partial\phi^a$, $G^b = \partial G/\partial\phi_b^*$, etc. After some algebraic rearrangements, these can be written as

$$\frac{1}{12}(\partial_\mu K)^2 + e^{K/3}|\partial_\mu\phi^i|^2 - \frac{3}{4}e^{2K/3}|\partial_\mu(z - z^*) + \frac{1}{3}(\phi_i^*\partial_\mu\phi^i - \phi^i\partial_\mu\phi_i^*)|^2 \tag{2}$$

The vector field kinetic terms are specified by an additional function $f_{\alpha\beta}$ as

$$-\frac{1}{4}(Ref_{\alpha\beta})(F_\alpha)_{\mu\nu}F_\beta^{\mu\nu} \tag{3}$$

and the scalar potential is given as

$$V = e^G[G_i(G^{-1})^i_j G^j - 3] + \frac{1}{2}Ref^{-1}_{\alpha\beta}D^\alpha D^\beta = e^G e^{-K/3}F_i F^{i\dagger} + \frac{1}{2}Ref^{-1}_{\alpha\beta}D^\alpha D^\beta \tag{4}$$

where the D-term is given by $D^\alpha = gG^iT_i^{\alpha j}\phi_j$ with a gauge coupling constant g and the generator of the gauge group T^α. Derivatives of the superpotential are denoted by $F_i = W_i/W = \partial\ln W/\partial\phi^i$. It should be noted that there are no soft masses for scalar fields in the scalar potential. Further, it is known that even anomaly-mediated SUSY breaking effects vanish in no-scale supergravity models [20]. This feature enables us to have an

[a] We require that the argument of the logarithmic function be positive since otherwise the kinetic terms of the matter fields have wrong signs, or equivalently, unitarity is broken.

arbitrarily large gravitino mass by choosing an appropriate gauge kinetic function [8]. We will come back to this point in Sec. 5.

The theory is completely defined once G and the gauge kinetic function $f_{\alpha\beta}$ are specified. For now, we will take $f_{\alpha\beta} = \delta_{\alpha\beta}$. We will consider a non-trivial form for the gauge kinetic function in Sec. 4 when we discuss the possible inflaton decay channel through the terms involving this function.

As one can see from Eq. (4), the scalar potential takes a form reminiscent of globally supersymmetric models. Indeed, it can be rewritten as

$$V = e^{2K/3} W_i W^{i\dagger} + \frac{1}{2}\delta_{\alpha\beta} D^\alpha D^\beta \tag{5}$$

The above scalar potential is semi-positive definite and, from now on, we assume $< F^i > = < F_i^\dagger > = < D^\alpha > = 0$ at the minimum to ensure the vanishing of the cosmological constant.

3 Inflaton mass eigenstate

From the form of the Kähler potential and the scalar kinetic terms in Eq. (2), it is clear that we have defined the theory in a basis with non-minimal kinetic terms. In discussing the decay of an inflaton field, it will be useful to define the inflaton mass eigenstate in a basis with canonically normalized fields. The canonically normalized scalar fields can be read off from Eq. (2)

$$Z_R = \sqrt{\frac{1}{6}} K \tag{6}$$

$$iZ_I = e^{<K>/3}\sqrt{\frac{3}{2}}(z - z^* + \frac{1}{3}\phi^*_{01}\delta\phi^1 - \frac{1}{3}\phi^1_0 \delta\phi^*_1) \tag{7}$$

$$\Phi^i = e^{<K>/6}\phi^i \tag{8}$$

$$A'^{\alpha}_\mu = < Ref_{\alpha\beta} >^{1/2} A^\beta_\mu \tag{9}$$

where A_μ is a gauge boson and we assumed that the scalar components of z and ϕ^1 have finite vacuum expectation values (VEVs), z_0 and ϕ^1_0 respectively ($\delta\phi^1$ is the fluctuation around ϕ^1_0). We assume that all of the other ϕ^i have vanishing VEVs due to some symmetries (e.g. gauge symmetries) [b].

[b] The general procedure to obtain the mass-eigenstate basis was presented in Ref. [19].

It is straight forward to calculate the scalar mass matrix in terms of these canonically normalized fields. The scalar mass matrix elements involving the canonically normalized inflaton field are

$$
\begin{aligned}
< \frac{\partial^2 V}{\partial \Phi^1 \partial \Phi_j^*} > &= < \frac{\partial \phi^1}{\partial \Phi^1} \frac{\partial \phi_j^*}{\partial \Phi_j^*} \frac{\partial^2 V}{\partial \phi^1 \partial \phi_j^*} > \\
&= e^{<G>} e^{<-2K/3>} < F_{1k} F^{\dagger kj} > \qquad (10)
\end{aligned}
$$

where we have used $\langle F_i \rangle = 0$ for all i, and $\langle K_i \rangle = 0$ for $i \neq 1$ (note that subscripts indicating the derivatives are with respect to the model (un-normalized) fields). This leads to

$$
\langle \frac{\partial^2 V}{\partial \Phi^1 \partial \Phi_j^*} \rangle = e^{<K/3>} < W_{1k} W^{*kj} > \qquad (11)
$$

which vanishes in the absence of direct coupling terms between an inflaton and the other fields in the superpotential, $< W_{1j} >= 0$ for $j \neq 1$. Note that we also have

$$
< \frac{\partial^2 V}{\partial \Phi_1^* \partial \Phi_j^*} >=< \frac{\partial^2 V}{\partial \Phi^1 \partial \Phi^j} >=< \frac{\partial^2 V}{\partial \Phi_1^* \partial Z_{R,I}} >=< \frac{\partial^2 V}{\partial \Phi^1 \partial Z_{R,I}^*} >= 0 \qquad (12)
$$

which are obtained from $< F^i >=< F_i^* >= 0$ at the minimum.

Hence, if there is no direct coupling between the inflaton and the other fields in W, the canonically normalized inflaton field is the inflaton mass eigenstate.

Before starting the discussion for the suppression of the decay of Φ^1, we make a brief comment on an approximate symmetry of $\delta\Phi^1$ (the fluctuation around its VEV $< \Phi^1 >$). In terms of the canonically normalized fields, the dependence of G on the inflaton Φ^1 appears only in the superpotential. Since the linear term of $\delta\Phi^1$ in the superpotential should vanish from $< \partial W/\partial\Phi^1 >=< \partial W/\partial\delta\Phi^1 >= 0$, the lowest order term in W is quadratic in $\delta\Phi^1$. Noting that the higher order terms in $\delta\Phi^1$ do not affect the decay of the inflaton, G has an approximate Z_2 symmetry of $\delta\Phi^1$. Taking this into account, one can guess that spontaneous inflaton decay should be suppressed (apart from possible Z_2 symmetry breaking via terms including the gauge kinetic function $f_{\alpha\beta}$). In the following, we show explicitly that inflaton decay is indeed suppressed for the terms which are solely determined from G in Sec. 3.1-3.3, followed by the discussion of the possible inflaton decay via terms involving $f_{\alpha\beta}$ in Sec. 4.

3.1 Inflaton coupling terms from the mass matrix expansion

We can consider the expansion of the mass matrix to study the possible decay of Φ^1. For example, the scalar mass matrix $\Phi_i^*(\mathcal{M}_0^2)^i_j\Phi^j(i,j\neq 1)$ can in principle give Φ^1 coupling terms such as $<\partial(\mathcal{M}_0^2)^i_j/\partial\Phi^1>\delta\Phi^1\Phi_i^*\Phi^j$. Such terms however vanish due to the special form of Kähler potential in a no-scale model because

$$\begin{aligned}
\left\langle\frac{\partial}{\partial\Phi^1}(\mathcal{M}_0^2)^i_j\right\rangle\delta\Phi^1\Phi_i^*\Phi^j &= \left\langle\frac{\partial}{\partial\Phi^1}e^Ge^{-2K/3}F_{jk}F^{\dagger ki}\right\rangle\delta\Phi^1\Phi_i^*\Phi^j \\
&= \left\langle\frac{\partial}{\partial\Phi^1}e^{\sqrt{2/3}Z_R}W_{jk}W^{*ki}\right\rangle\delta\Phi^1\Phi_i^*\Phi^j \\
&= \left\langle e^{\sqrt{2/3}Z_R}\frac{\partial\phi^1}{\partial\Phi^1}\frac{\partial}{\partial\phi^1}W_{jk}W^{*ki}\right\rangle\delta\Phi^1\Phi_i^*\Phi^j \\
&= 0
\end{aligned} \tag{13}$$

where the last equality is due to the assumption that there are no terms which directly couple an inflaton to the other fields in W. In the above, we also used the fact that $<\partial D^{\alpha i}/\partial\Phi^1>=0$, so that the D-term does not contribute to inflaton decay either. One also finds that inflaton decay coming from the expansion of the other scalar mass terms vanish by performing the analogous calculations as above, such as $<\partial(\mathcal{M}_0^2)_{Z_R j}/\partial\Phi^1>=0$. Note that the absence of a decay term for the inflaton to scalars is a direct result of the no-scale form of the potential in Eqs. (4) and (5).

Similarly, the expansion of matter Fermion mass matrix terms $\bar{\chi^i}\mathcal{M}_{1/2_{ij}}\chi^j$ as well as the matter Fermion-gaugino mass matrix $\bar{\lambda}_\alpha\mathcal{M}_{1/2}{}^\alpha_i\chi^i$ give a vanishing contribution for inflaton decay. In general, we can write the chiral fermion mass matrix as

$$\bar{\chi^i}\mathcal{M}_{1/2_{ij}}\chi^j = e^{G/2}\bar{\chi^i}\left(G_{ij}+G_iG_j-G^m_{ij}(G^{-1})^n_mG_n\right)\chi^j \tag{14}$$

For the specific case of no-scale supergravity, this can be rewritten as [c]

$$-e^{G/2}\bar{\chi^i}\left(\frac{2}{3}G_iG_j+F_{ij}+\frac{1}{3}F_iF_j\right)\chi^j \tag{15}$$

Subtracting the Goldstino component $\eta=(G_i)\bar{\chi^i}$ in the unitary gauge, the Fermion mass

[c] The following arguments are not affected even if one uses the canonically normalized Fermion fields $\chi'^i=e^{<K>/6}\chi^i, \lambda'^\alpha=<Ref_{\alpha\beta}>^{1/2}\lambda^\beta$.

matrix term becomes

$$- e^{K/2}\bar{\chi}^i \left(W_{ij} - \frac{2}{3}\frac{W_i W_j}{W} \right) \chi^j \tag{16}$$

Any possible inflaton decay channels to two chiral fermions can be obtained by the expansion of Eq. (16) around the VEV of $< \Phi^1 >$. As one can see

$$\begin{aligned} & - \left\langle \frac{\partial}{\partial \Phi^1} e^{-\sqrt{3/2} z_R} \left(W_{ij} - \frac{2}{3}\frac{W_i W_j}{W} \right) \right\rangle \delta\Phi^1 \bar{\chi}^i \chi^j \\ = & - \left\langle e^{-\sqrt{3/2} z_R} \frac{\partial \phi^1}{\partial \Phi^1} \frac{\partial}{\partial \phi^1} \left(W_{ij} - \frac{2}{3}\frac{W_i W_j}{W} \right) \right\rangle \delta\Phi^1 \bar{\chi}^i \chi^j \end{aligned} \tag{17}$$

and

$$- \left\langle \frac{\partial}{\partial \Phi^k} \frac{\partial}{\partial \Phi^1} e^{-\sqrt{3/2} z_R} \left(W_{ij} - \frac{2}{3}\frac{W_i W_j}{W} \right) \right\rangle \delta\Phi^1 \Phi^k \bar{\chi}^i \chi^j \tag{18}$$

vanish because $< W_i > = < W_1 > = 0$ and we have assumed no direct coupling between inflaton and the other fields in W. Indeed, we can easily see from the above procedures that all the inflaton decay channels due to the expansion of the Fermion mass terms vanish at the tree level.

For completeness, we also consider the inflaton coupling to gauginos and matter Fermions,

$$2igG^i_j (T^\alpha)_{ik} \phi^k \bar{\lambda}_\alpha \chi^i \tag{19}$$

Inflaton decay from this term also vanishes because $\left\langle \partial G^i_j / \partial \Phi^1 \right\rangle = 0$.

Hence all tree-level inflaton decays to scalar and fermion matter fields including gauginos exactly vanish at the global minimum in no-scale supergravity models.

3.2 Derivative coupling

Next we consider possible inflaton decays via kinetic terms. From Eq. (2), we see that the inflaton, ϕ^1, may couple to scalar fields ϕ^i through,

$$- \left\langle \frac{\partial e^{K/3}}{\partial \phi^1} \right\rangle \delta\phi^1 \partial\phi^i \partial\phi_i^\dagger + \text{h.c.} \tag{20}$$

where again $\delta\phi_1$ is the fluctuation around its VEV. However, we can easily see that, by considering the canonically normalized scalar fields, the total decay of the inflaton mass eigenstate should vanish. All of the derivatives of the coefficients in the kinetic terms in Eq. (2) with respect to Φ^1 vanish because $< \partial Z_R/\partial\Phi^1 >= 0$. Hence the Φ^1 decay channel to canonically normalized scalar fields which can be obtained by the expansion of the kinetic terms in Eq. (2) exactly vanish. Couplings such as those in Eq. (20) are absorbed by the canonically normalized mass-eigenstate field $Z_R \equiv K/\sqrt{6}$. The same is true for the matter Fermion kinetic terms

$$- e^{K/3}\sum_{i=1}^{n-1}\bar{\chi}_i \not{\partial}\chi^i \tag{21}$$

Thus the inflaton mass-eigenstate Φ^1 does not decay through the kinetic terms.

3.3 Gravitinos

Finally, let us now consider the decay of the inflaton into a pair of gravitinos. Gravitino overproduction from inflaton decay is quite dangerous because the decay of gravitinos may affect the light-element abundances or the density of the lightest supersymmetric particles (LSPs) produced by gravitino decay may exceed the allowed dark matter abundance. The relevant interactions for gravitino pair production are given by [21]

$$-\frac{1}{8}\epsilon^{\mu\nu\rho\sigma}\left(\frac{\partial G}{\partial\Phi^1}\partial_\rho\Phi^1 - \frac{\partial G}{\partial\Phi_1^*}\partial_\rho\Phi_1^*\right)\bar{\psi}_\mu\gamma_\nu\psi_\sigma - \frac{1}{8}e^{G/2}\left(\frac{\partial G}{\partial\Phi^1}\delta\Phi^1 + \frac{\partial G}{\partial\Phi_1^*}\delta\Phi_1^*\right)\bar{\psi}_\mu\left[\gamma^\mu,\gamma^\nu\right]\psi_\nu \tag{22}$$

where ψ_μ is the gravitino field, and we have chosen the unitary gauge in the Einstein frame. Thus the inflaton couplings with gravitinos are proportional to $\langle\partial G/\partial\Phi^1\rangle = 0$. Once again, for completeness, we display the scalar-gravitino-gaugino coupling

$$-\frac{i}{2}gG^i(T^\alpha)_{ij}\phi^j\bar{\psi}_\mu\gamma^\mu\lambda_\alpha \tag{23}$$

which also can not contribute to inflaton decay because $\langle\partial G^i/\partial\Phi^1\rangle = 0$. Therefore, gravitino pair production rate from inflaton decay vanishes exactly at the tree-level.

4 Inflaton decay via gauge kinetic function

The exercise of the previous section shows that in the context of no-scale supergravity, not only is there no gravitino problem due to excess reheating, there is virtually no reheating at all. Thus we are presented with a potentially more severe problem for inflation in no-scale supergravity. Finding that the inflaton decay in no-scale supergravity model is indeed highly suppressed, the natural question now would be to find the dominant decay channel of an inflaton to reheat the universe.

The absence of supersymmetry breaking scalar masses in Eq. (4) is one of the features of a no-scale model and one mechanism to mediate supersymmetry breaking to the visible sector can be specified through a z-dependent gauge kinetic function. In this section, we show that terms involving a non-trivial gauge kinetic function can be responsible for the dominant channel of Φ^1 decay [19].

Among the terms involving the gauge kinetic function, the terms of interest here are

$$-\frac{1}{4}(Ref_{\alpha\beta})F^{\alpha}_{\mu\nu}F^{\beta\mu\nu}+\frac{i}{4}(Imf_{\alpha\beta})\epsilon^{\mu\nu\rho\sigma}F^{\alpha}_{\mu\nu}F^{\beta\rho\sigma}+\left(\frac{1}{4}e^{G/2}\frac{\partial f^*_{\alpha\beta}}{\partial\phi^*_j}(G^{-1})^k_jG_k\lambda^\alpha\lambda^\beta+h.c.\right) \quad (24)$$

Other terms involving the gauge kinetic function are the derivative coupling terms, and the inflaton decay rate from those terms are suppressed by the masses of the final state particles. For illustrative purposes, let us take the simplest non-trivial form for $f_{\alpha\beta}$ such that it depends only on the field z

$$f_{\alpha\beta}\equiv\delta_{\alpha\beta}h(z) \quad (25)$$

This simple choice determines the universal Majorana (canonically normalized) gaugino mass at the unification scale for softly broken global supersymmetry

$$m_{1/2}=\left|\frac{1}{2}e^{G/2}\frac{h_z}{Re\ h}(G^{-1})^k_zG_k\right|=\left|\frac{1}{2}e^{(G/2-K/3)}\frac{h_z}{Re\ h}(1-\phi F_{\phi\phi})\right| \quad (26)$$

where $h_z\equiv\partial h/\partial z$ and we used the relation for the projection operator $\alpha_z=(G^{-1})^k_zG_k=-e^{-K/3}(1-\phi F_\phi/3)$[d].

The decay of the inflaton to two gauge bosons can be obtained from the expansion of Eq. (24)

$$\begin{aligned}-\frac{1}{4}\left\langle\frac{\partial}{\partial\Phi^1}h\right\rangle\delta\Phi^1F_{\alpha\mu\nu}F^{\alpha\mu\nu} &= -\frac{1}{4}\left\langle h_z\frac{\partial z}{\partial\Phi^1}\right\rangle\delta\Phi^1F_{\alpha\mu\nu}F^{\alpha\mu\nu}\\ &= -\frac{1}{12}\left\langle h_ze^{-K/3}\Phi^*_1\right\rangle\delta\Phi^1F_{\alpha\mu\nu}F^{\alpha\mu\nu} \quad (27)\end{aligned}$$

[d]Note the gaugino mass term is related to the gravitino mass $m_{3/2}=|e^{G/2}|$, but, in general, the exact ratio can vary depending on the form of $f_{\alpha\beta}$, cf. Eq. (26).

where we have used $\partial z/\partial\Phi^1 = e^{-<K/3>}\Phi_1^*/3$. Thus, the inflaton decay rate to canonically normalized gauge bosons from this coupling becomes of order [e]

$$\Gamma(\Phi_1 \to A_\mu A_\mu) \sim \mathcal{O}(10^{-3}) \times e^{-\frac{<K>}{3}} \left|\left\langle \frac{e^{-K/6}\Phi_1}{M_p} \right\rangle\right|^2 \left|\left\langle \frac{h_z}{Re\ h} \right\rangle\right|^2 \frac{m_{\phi^1}^3}{M_p^2} \tag{28}$$

In addition, there is also a non-negligible contribution for the inflaton decay to gauginos via the gaugino mass term in Eq. (24) which from Eq. (26) becomes of order

$$\begin{aligned} &\frac{1}{4}\left\langle e^{G/2-K/3}\frac{\partial}{\partial\Phi^1}\left[\frac{h_z}{Re\ h}(1-\phi^1 F_1/3)\right]\right\rangle \delta\Phi^1\lambda^{\alpha'}\lambda^{\alpha'} \\ =& \frac{1}{12}\left\langle e^{G/2-2K/3}\left[\left(\frac{h_z}{Re\ h}\right)_z \Phi_1^* - \frac{h_z}{Re\ h}\Phi^1 F_{11}\right]\right\rangle \delta\Phi^1\lambda^{\alpha'}\lambda^{\alpha'} \end{aligned} \tag{29}$$

where λ' is the canonically normalized gaugino field $\lambda'^\alpha =< Re\ h >^{1/2} \lambda^\alpha$. If $(h_z/Reh)_z$ is small, the second term in Eq. (29) dominates and we obtain a decay rate to gauginos

$$\Gamma(\Phi_1 \to \lambda'\lambda') \sim \mathcal{O}(10^{-3}) \times e^{-\frac{<K>}{3}} \left|\left\langle \frac{e^{-K/6}\Phi^1}{M_p} \right\rangle\right|^2 \left|\left\langle \frac{h_z}{Re\ h} \right\rangle\right|^2 \frac{m_{\phi^1}^3}{M_p^2} \tag{30}$$

as in Eq. (28), where we have used $m_{\phi^1} = |e^{G/2-K/3}F_{11}|$ (from Eq.(10)).

The reheating temperature in our example is estimated to be of order

$$T_{RH} \sim \mathcal{O}(10^7) \times \left|\left\langle \frac{h_z}{(Re\ h)} \right\rangle\right| \text{GeV} \tag{31}$$

for $m_{\phi^1} \sim 10^{-7}$ (coming from the constraints on the cosmic perturbation amplitude [f]) and $\langle\Phi^1\rangle \simeq M_P$, which is low enough to avoid the gravitino problem [22] unless $< h_z/(Re\ h) >$ is tuned to be much larger than $\mathcal{O}(1)$.

5 Discussion and conclusion

In this letter we have shown that the many inflaton decay channels to the matter fields and gravitinos are highly suppressed when the Kähler potential is of the no-scale form, and that the dominant inflaton decay channels depend on the gauge kinetic function.

[e] For the numerical estimation, we assumed there are a total of 12 gauge bosons as in MSSM.

[f] We are considering here small field inflation models, such as that in [6], with an inflaton VEV $< \phi >\sim \mathcal{O}(1)$. In fact, inflation models with a VEV $< \phi >\ll \mathcal{O}(1)$ would suffer from insufficient reheating in no-scale supergravity.

As a consequence, the decay of the inflaton connects the reheating temperature with the relation between gaugino and gravitino masses. For example, in our example, the reheating temperature is proportional to the ratio of the gaugino mass to the gravitino mass determined by $< h_z/Re\ h >$. This is in contrast to the discussions for the other inflation models in no-scale supergravity previously considered [7, 8] where either explicit inflaton couplings or non-trivial gravitational couplings are assumed for the inflaton decay. The direct coupling of an inflaton and other fields in superpotential, for instance, will spoil the flatness of the inflaton potential [g] even though no-scale Kähler potential can still help suppressing the reheating temperature and the gravitino production.

The low reheating temperature certainly helps resolving the gravitino problem. We note, however, that the gravitino problem is not generic in no-scale supergravity models, as the gravitino mass is a priori undetermined and may be as large as the Planck scale [8]. As we have discussed, in generic inflation models based on no-scale supergravity, a Planck mass gravitino would require a very small value for h_z restoring our initial problem of sufficient reheating. For instance, if we require a reheating temperature T_{RH} larger than 10 MeV [23, 24, 25], $< h_z/Re\ h >$ cannot be smaller than $\mathcal{O}(10^{-9})$ from Eq. (31), which leads to $m_{3/2} \lesssim \mathcal{O}(10^{12})$ GeV for $m_{1/2} = \mathcal{O}(10^3)$ GeV (assuming $e^{-K/3} \sim \mathcal{O}(1)$, see Eq. (26)).

Although the non-minimal couplings between the inflaton ϕ^1 and the supersymmetry breaking field z induce gravitino overproduction in a generic supergravity model, as we saw in Sec. 3.3, inflaton decay into a pair of gravitinos is suppressed in no-scale supergravity . This suppression is quite important. If such suppression did not occur, it would be very difficult to satisfy the BBN constraints [22] for an unstable gravitino of a mass $m_{3/2} = 100\,\mathrm{GeV} - 10\,\mathrm{TeV}$. Further, even for gravitinos heavier than 10 TeV, the abundance of the LSPs produced by gravitino decay might exceed the dark matter abundance [h]. In

[g] For example, it may be hard to prevent the couplings via non-renormalizable operators. The inclusion of a non-renormalizable coupling such as $W \ni \lambda_n \phi^1 {\phi^i}^n$ $(n > 2)$ in addition to the superpotential $W \ni \mu^2(\phi^1 - (\phi^1)^4/4)$ which can lead to inflation in no-scale supergravity [6] would require the fine-tuning of the coupling constant $\lambda_n \ll 10^{-7}$ to insure the flatness of the inflaton potential if the ϕ^i obtain a large (say Planck scale) vev due to the quantum fluctuations during inflation. Even if the ϕ^i happen to have the vanishing vev during inflation so that the inflaton potential around the origin is not affected by such a coupling, we may still derive a limit from the reheating temperature due to the additional decay channels of the inflaton which is of order $\lambda_n \lesssim 10^{-21+7n}$ (assuming the conservative limit of $T_{RH} < 10^8 GeV$).

[h] For gravitinos heavier than $\mathcal{O}(10^{3-4})$ TeV, the resulting temperature after gravitino decay can be high

particular, as long as the decay via the gauge kinetic function is the dominant source for reheating, the reheating temperature is inversely proportional to the gravitino mass. So, heavier gravitinos and a lower reheating temperature as a result of a smaller decay rate of the inflaton via the gauge kinetic function, would lead to a larger branching ratio for gravitino production [15], making the problem more severe. Fortunately, gravitino production is suppressed in no-scale supergravity, and the problems mentioned above are avoided.

The production of a baryon asymmetry is also somewhat constrained. For example, baryo/leptogenesis through out-of-equilibrium decay normally requires the inflaton to decay directly into fields generating the asymmetry. In the no-scale models discussed above, the inflaton decays directly only into gauge bosons and gauginos, thus severely limiting possible mechanisms. Models such as the Affleck-Dine mechanism at low temperature would remain plausible possibilities [26].

As mentioned before, since the modulus z has a flat potential at the tree level, there is a moduli problem associated with the z field. The moduli problem [27] is a prevailing problem in many inflation models, and is quite often a more serious problem than the gravitino problem. Firstly, one needs to stabilize the moduli. For instance, to avoid the run-away minimum during inflation, one may need to modify the Kähler potential [8] or add additional D-term or non-perturbative effects [9]. Secondly, even if we can stabilize the moduli, one still needs to worry about their late-time decay which can jeopardize Big Bang Nucleosynthesis [22]. Furthermore, it should be noted that, once z is stabilized by e.g. introducing non-trivial z-dependence of the superpotential or modifying the structure of the Kähler potential, the anomaly-mediation effects are generically non-negligible. Then, it would be difficult to have the gravitino mass larger than $\mathcal{O}(100)$ TeV on the basis of naturalness.

Finally, we note that, analogous to the discussion of the decay of Φ^1 to gravitinos in Sec. 3.3, the modulus Z_R (which is also a mass eigenstate [i]) couples to gravitinos with a coupling of order $e^{G/2}G_{Z_R}/8$. Recalling that $< G_{Z_R} >= \sqrt{6}$, the modulus decay to

enough for LSPs to annihilate efficiently (or reach thermal equilibrium). Then the overclosure problem can be resolved.

[i]This may not be the case if z is stabilized by modifying its potential, which, however, does not essentially affect the following discussion.

the gravitinos may not be negligible (the possible significance of the modulus decay to the gravitinos was also pointed out in [28]), if the decay is kinematically allowed. One possibility to avoid this problem of late-time moduli decay is the enhancement of the moduli decay at the minimum [29]. We leave the study of the moduli problem in no-scale supergravity for the future work.

Acknowledgments

K.K. and K.A.O would like to thank Tony Gherghetta and John Ellis for useful discussions. The work of K.K. and K.A.O. was partially supported by DOE grant DE–FG02–94ER–40823.

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