Reheating in Chaotic D-Term Inflation

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Abstract
A simple model is discussed to give rise to successful reheating in chaotic D-term inflation
with a quadratic inflaton potential, introducing a trilinear coupling in the Kähler potential.
Leptogenesis through the inflaton decay is also discussed in this model.

1 Introduction

Inflation gives a natural solution to the horizon and the flatness problems as well as the
mechanism to generate primordial density fluctuations [1]. The recent observations like
the cosmic microwave background anisotropy by the Wilkinson Microwave Anisotropy
Probe (WMAP) [2] strongly supports inflationary scenarios. Among many types of
inflation models proposed thus far, chaotic inflation, which was proposed by Linde
[3], is attractive in that it does not suffer from any initial condition problem. Since
supersymmetry is a plausible solution to the hierarchy problem in particle physics, it
is natural to consider inflation models in supergravity. However, it is very difficult to
realize chaotic inflation in supergravity. This is mainly because the F-term potential
in supergravity has the exponential of the Kähler potential as a factor, and in general
prevents scalar fields from having an initial value much larger than the reduced Planck
scale \( M_p \simeq 2.4 \times 10^{18} \) GeV.

One method to circumvent this difficulty is to introduce the shift symmetry [4, 5, 6, 7, 8],
which guarantees the flat directions especially in the Kähler potential and allows
a field along the flat directions to be identified with an inflaton field. Another method
is to use the D-term potential, which does not have the exponential factor as the F-
term potential does [13]. A field along the F-flat direction, which is automatically lifted
by the D-term potential, can play the role of inflaton to give rise to chaotic inflation.

See Refs. [9, 10, 11] for early attempts to realize chaotic inflation in supergravity. See also Ref.
[12], in which the authors attempt to realize chaotic inflation in superstring by use of the monodromy.
The idea was used in the papers [13, 14] to give chaotic inflation models with a quartic potential. Chaotic inflation models with a quartic potential in general seems disfavored by the recent observations [2] and may require additional contributions like curvaton [15, 16, 17] to explain the observations.

Recently, a simple chaotic inflation model with a quadratic potential was proposed in supergravity [18], which is based on $N = 1$ supersymmetric $U(1)$ gauge theory where the Fayet-Iliopoulos parameter replaced by a dynamical field plays the role of inflaton. Since the potential of the model has only the D-term potential and depends on the Fayet-Iliopoulos parameter quadratically, the model gives rise to chaotic inflation with a quadratic potential of the inflaton field. Thus, in the paper [18], the eta-problem was naturally solved. It is well-known that many models in supergravity for inflation suffer from not only the eta-problem but also the overproduction of gravitino. However, to avoid the overproduction of gravitino, one needs to make sure two necessary conditions; the first is to obtain sufficiently low reheating temperature, and the second is to prevent the inflaton field from developing the vacuum expectation value (VEV) after inflation. It is also known that it is extremely difficult to meet these two conditions. Therefore, in this paper, we concentrate on the reheating stage in the model. By introducing interactions of the inflaton with right-handed neutrinos, it will be discussed that one can achieve successful reheating with low reheating temperature and leptogenesis.

In the next section, after briefly reviewing the model of Ref. [18], we introduce an interaction necessary for reheating and investigate its effect on the dynamics of the inflation. In Sec. 3, we evaluate the reheating temperature and show that leptogenesis really works in our scenario. In the final section, we give summary and discussions.

# 2 Model and Chaotic Inflation

The model proposed in the paper [18] is $N = 1$ supersymmetric $U(1)$ gauge theory with a chiral superfield $S$, whose scalar component is denoted by

$$S = \frac{1}{\sqrt{2}} (\rho + i \zeta).$$

The real scalar field $\rho$ plays the role of the Stuckelburg field giving mass to the $U(1)$ gauge multiplet, while the real scalar field $\zeta$ is the inflaton field appearing in the D-term potential as if it was the Fayet-Iliopoulos parameter.

In order to achieve reheating in the model, we also introduce neutral chiral superfields $N_i$ ($i = 1, 2, 3$), which play the role of right-handed (s)neutrinos, and the minimal supersymmetric standard model (MSSM) superfields $\phi^k$. Under the gauge transformation, they transform as

$$V \rightarrow V + i (\Lambda - \bar{\Lambda}), \quad S \rightarrow S + \sqrt{2} M \Lambda, \quad N_i \rightarrow N_i, \quad \phi^k \rightarrow \phi^k,$$

where $V$ is the vector superfield of the gauge field $v_\mu$, and $M$ will turn out to be the
inflaton mass. Since we take the Kähler potential

\[ K = -\frac{1}{2} (S - \bar{S} + \sqrt{2} i M \mathcal{V})^2 - i \frac{g}{M_p} (S - \bar{S} + \sqrt{2} i M \mathcal{V}) \bar{N}_i N_i + \bar{\Phi}_k \Phi^k, \quad (1) \]

with \( g \) a coupling constant, it follows that the Killing potential \( D \) is given by

\[ D = \frac{M}{\sqrt{2}} \left[ -i (S - \bar{S}) + \frac{g}{M_p} \bar{N}_i N_i \right]. \]

Note that such a trilinear coupling between the inflaton and the right-handed neutrinos in the Kähler potential is peculiar to our model because the R-symmetry prohibits such a term for many existing models. Furthermore, we take the superpotential to be

\[ W = \frac{m_i}{2} N_i N_i + \bar{W}(N, \phi^k), \]

where \( \bar{W}(N, \phi^k) \) includes the Yukawa interactions \( y_{ij} N_i L_j H_u \) of the right handed neutrinos \( N_i \) with the left-handed lepton superfields \( L_i \) \((i = 1, 2, 3)\) and the Higgs superfield \( H_u \).

Since it is convenient to describe the model in terms of gauge invariant variables, let us change the variable as

\[ v_{\mu} \to v_{\mu} + \frac{1}{M} \partial_{\mu} \rho \]

to completely eliminate the field \( \rho \) from the action as the gauge degrees of freedom.

One then finds that the potential is given by

\[ V = \exp \left[ \frac{1}{M_p^2} \left( \zeta^2 + |N_i|^2 + \frac{\sqrt{2} g}{M_p} \zeta |N_i|^2 + K(\phi^*_k, \phi^k) \right) \right] \times \]

\[ \left[ \frac{1}{1 + \sqrt{2} g \zeta} |D_k N_i| W|^2 + \frac{1 + \sqrt{2} g \zeta}{1 + \frac{\sqrt{2} g}{M_p} \zeta - \frac{g^2}{M_p^2} |N_i|^2} \left( \frac{\sqrt{2} g}{M_p} W - \frac{g}{M_p} \frac{1}{1 + \frac{\sqrt{2} g}{M_p} \zeta} (m_i N_i^2 + \bar{W}) \right) \right. \]

\[ + g^k N_k \bar{W}(D_k \bar{W})^* - \frac{3}{M_p^2} |W|^2 \] \[ + \frac{M^2}{2} \left( \zeta + \frac{g}{\sqrt{2} M_p} |N_i|^2 \right)^2 + V_{D}^{MSSM}. \quad (2) \]

where \( D_k W \) denotes

\[ D_k W = \frac{\partial}{\partial \phi^k} W + \frac{1}{M^2_p} \left( \frac{\partial}{\partial \phi^k} K \right) W. \quad (3) \]

Note that the first term on the right hand side of Eq. (2) is the F-term potential, and the remaining terms the D-term potential. \( V_{D}^{MSSM} \) represents the contribution coming from the standard model gauge group to the D-term potential.

\(^{2}\)The notation in this paper is slightly different from the one in [18], and in this paper we follow the convention of chapter XXIV and Appendix G in the textbook [19], with trivial exceptions for the Kähler potential and the superpotential.
When the universe starts around the Planck energy density and the inflaton $\zeta$ takes a value much larger than $M_p$, the exponential factor in the F-term potential enforces the conditions $D_N W = D_k W = W = 0$ to set the contribution from the F-term in the potential to zero. Since the directions with $N_i = 0$ satisfy the condition, the system could remain in one of the directions during inflation. In fact, the effective masses squared of the neutrinos at the origin are estimated as

$$m_i^2(\zeta) \simeq \frac{m_i^2}{1 + \frac{\sqrt{2} g}{M_p} \zeta} e^{\frac{\zeta^2}{M_p^2}} + \frac{g \zeta}{\sqrt{2} M_p} M^2.$$  

On the other hand, the Hubble parameter squared in the same period is given by

$$H^2 \simeq \frac{1}{6} \left( \frac{M}{M_p} \right)^2 \zeta^2.$$  

Therefore, for large $\zeta$, the effective masses of the neutrinos are much larger than the Hubble parameter so that $N_i$ remains at the origin during inflation$^3$. Here and hereafter, for simplicity, we will take positive $g$ without loss of generality and assume positive $\zeta$ during inflation. Thus, the potential of the inflaton is dominated by the D-term with the quadratic potential, which leads to chaotic inflation. One thus needs to take the inflaton mass $M$ to be of order $10^{13}$ GeV in order to explain the primordial density fluctuations.

### 3 Reheating and Leptogenesis

So thus, we have seen that the model of the paper $^{18}$ with the neutrino fields $N_i$ still gives rise to successful chaotic inflation. In this section, we would like to discuss reheating and the baryon asymmetry of the universe through the leptogenesis scenario $^{20}$ via the inflaton decay $^{21, 22, 23, 24, 25}$ in our model.

After inflation, the inflaton field starts oscillating about the origin$^4$. Thanks to the interaction term between the inflaton $\zeta$ and the neutrino multiplets $N_i$ in the Kähler potential, the inflaton can decay into the (s)neutrinos through the interactions

$$\frac{g}{\sqrt{2} M_p} \zeta \left[ -M^2 |N_i|^2 + (2m_i^2) |N_i|^2 + \bar{N}_i \partial_\mu \partial^\mu N_i + N_i \partial_\mu \partial^\mu \bar{N}_i - i \bar{\chi}_i \sigma^\mu \partial_\mu \chi_i - i \chi_i \sigma^\mu \partial_\mu \bar{\chi}_i \right],$$  

where the field $\chi_i$ is the spinor component of the chiral superfield $N_i$. Here you should notice that the trilinear coupling of the first term originates from the D-term potential, which is peculiar to our D-term model of chaotic inflation. In the last four terms on

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$^3$ Since the kinetic terms of the neutrinos $N_i$ are not canonical, we must have taken account of it to estimate their effective masses. However, the exponential factor is so large that such an effect is negligible.

$^4$ The kinetic term of the neutrinos become singular at $\zeta = -M_p/(\sqrt{2} g)$. However, since we assume that the inflaton takes a positive value during inflation, it cannot reach this value for $g \ll 1$ during the oscillation phase as well.
the right hand side of Eq. (4), the derivatives can be replaced by the neutrino mass $m_i$ by using the on-shell condition in lowest order perturbation theory.

Assuming that $m_1$ is of order $10^{12}$ GeV, while $m_2$ and $m_3$ are assumed to be larger than the inflaton mass $M$, we only have to consider the decay of the inflaton into $N_1$. Therefore, the first term coming from the D-term potential gives the dominant contribution to the decay rate of the inflaton. Then, one obtains the decay rate of the inflaton $\zeta$

$$\Gamma_\zeta \simeq \frac{1}{16\pi M} \left( \frac{g M^2}{\sqrt{2} M_p} \right)^2.$$ 

Therefore, the reheating temperature is given by

$$T_R \simeq 3.0 \times 10^7 \text{ GeV} \left( \frac{g}{0.1} \right) \left( \frac{M}{10^{13} \text{ GeV}} \right)^{3/2}.$$ 

Inflation models with such a low reheating temperature are viable for the wide range of the gravitino mass [26].

The produced (s)neutrinos $N_i$ decay into (s)leptons $L_j$ and Higgs(ino) doublets $H_u$ through the Yukawa interactions

$$W = h^i_\nu N_i L_j H_u$$

in the superpotential, where we have taken a basis where the mass matrix for $N_i$ is diagonal. We also assume $| (h_\nu)_{i3} | > | (h_\nu)_{i2} | \gg | (h_\nu)_{i1} |$ ($i = 1, 2, 3$). We consider only the decay of $N_1$ with the above assumption that the mass $m_1$ is much smaller than the others. The interference between the tree-level and the one-loop diagrams including vertex and self-energy corrections generates the lepton asymmetry [20, 27, 28, 29],

$$\epsilon_1 \equiv \frac{\Gamma(N_1 \rightarrow H_u + l) - \Gamma(N_1 \rightarrow \overline{H_u} + \overline{l})}{\Gamma(N_1 \rightarrow H_u + l) + \Gamma(N_1 \rightarrow \overline{H_u} + \overline{l})} \approx -\frac{3}{8\pi} \sum_{i=2,3} \text{Im} \left( h_\nu h^1_\nu \right)^2 \frac{m_1}{m_i} \approx \frac{3}{8\pi} \frac{m_1}{\langle H_u \rangle^2} m_{\nu_3} \delta_{\text{eff}} \sim 10^{-4} \left( \frac{m_1}{10^{12} \text{ GeV}} \right) \left( \frac{m_{\nu_3}}{10^{-2} \text{ eV}} \right) \delta_{\text{eff}},$$

with the effective CP violation phase given by

$$\delta_{\text{eff}} \equiv -\frac{\text{Im} \left[ (h_\nu m_\nu^*) h^T_\nu \right]_{11}^2}{m_{\nu_3} \left( h_\nu h^1_\nu \right)_{11}},$$

where $m_{\nu_3}$ is a mass eigenvalue of the left-handed neutrino mass matrix $m_\nu$ and can be estimated by the seesaw mechanism [30, 31] as

$$m_{\nu_3} \simeq \frac{| (h_\nu)_{33} |^2 \langle H_u \rangle^2}{m_3} \sim 10^{-2} \text{ eV} \left( \frac{| (h_\nu)_{33} |}{10^{-1}} \right)^2 \left( \frac{m_3}{10^{13} \text{ GeV}} \right)^{-1}.$$
Here we assumed that the \((h_\nu)_{33}\) is dominant in \(h_\nu\) and \(m_3 \gg m_1\).

The total decay rate \(\Gamma_{N_1}\) of \(N_1\) is given by

\[
\Gamma_{N_1} = \Gamma(N_1 \to H_u + l) + \Gamma(N_1 \to \bar{H}_u + \bar{l}) \\
\approx \frac{1}{8\pi} |(h_\nu)_{13}|^2 m_1 \\
\approx 10^6 \text{ GeV} \left( \frac{|(h_\nu)_{13}|}{10^{-2}} \right)^2 \left( \frac{m_1}{10^{12} \text{ GeV}} \right).
\]

Thus, the decay rate \(\Gamma_{N_1}\) is much larger than the decay rate of the inflaton \(\Gamma_\zeta\) so that the produced \(N_1\) immediately decays into lepton and Higgs supermultiplets once they are produced from the inflaton.

A part of the produced lepton asymmetry is converted into the baryon asymmetry through the sphaleron processes, which can be estimated \[32, 33\] as

\[
\frac{n_B}{s} \approx -\frac{8}{23} \frac{n_L}{s},
\]

where we have assumed the standard model with two Higgs doublets and three generations. In order to explain the observed baryon number density

\[
\frac{n_B}{s} \approx (0.4 - 1) \times 10^{-10},
\]

it is necessary to achieve the lepton asymmetry

\[
\frac{n_L}{s} \approx -(1 - 3) \times 10^{-10}.
\]

Now we estimate the lepton asymmetry produced through the inflaton decay. For \(m_1 \gtrsim 10^{12} \text{ GeV}\), \(m_1\) is much larger than the reheating temperature \(T_R\). In this case, the produced \(N_1\) is out of equilibrium and the ratio of the lepton number to entropy density can be estimated as

\[
\frac{n_L}{s} \approx \frac{3}{2} \epsilon_1 B_r \frac{T_R}{M} \\
\approx -5 \times 10^{-10} \delta_{\text{eff}} B_r \left( \frac{g}{10^{-1}} \right) \left( \frac{M}{10^{13} \text{ GeV}} \right)^{1/2} \left( \frac{m_1}{10^{12} \text{ GeV}} \right),
\]

where \(B_r\) is the branching ratio of the inflaton \(\zeta\) into \(N_1\). As stated before, we typically take \(m_1\) to be of order \(10^{12}\) GeV, and \(m_2\) and \(m_3\) to be at least larger than \(10^{13}\) GeV, which implies \(B_r = \mathcal{O}(1)\). Then, the typical values of the parameters can explain the baryon number density in the present universe. Thus, sufficient baryon asymmetry can be generated for low reheating temperature without fine-tuning of the parameters in our model.

4 Summary and Discussions

In this paper, we have discussed how to reheat the universe in the model of the paper \[18\]. For this purpose, we introduced the interactions of the inflaton with right-handed
neutrinos in the Kähler potential, which yields the effective trilinear couplings between them originating from the D-term potential. This property is peculiar to the model of D-term chaotic inflation. We have seen that this model leads to successful chaotic inflation even after introducing such an interaction. The inflaton can decay into right-handed (s)neutrinos via the effective trilinear interactions. The produced right-handed (s)neutrinos quickly decay into the MSSM particles through the Yukawa interactions so that the universe is completely reheated. We have found the typical reheating temperature to be $10^7$ GeV for $g \simeq 0.1$, which is low enough to avoid the overproduction of gravitino for a wide range of the gravitino mass. The decay of right-handed (s)neutrinos produces the lepton asymmetry if the CP violation exists. We have shown that the amount of the produced lepton asymmetry is sufficient to explain the present baryon asymmetry for typical parameter values.

The inflaton could also couple to other particles if similar trilinear couplings were introduced in the Kähler potential. However, the branching ratio to the right-handed neutrinos would be dominant in the inflaton decay, as long as the coupling constant to them is not significantly suppressed in comparison to those to the other particles. Therefore, they would hardly affect our results on the reheating temperature.
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