New D-term chaotic inflation in supergravity and leptogenesis

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(Dated: May 23, 2008)

Abstract

We present a new model of D-term dominated chaotic inflation in supergravity. The F-flat direction present in this model is lifted by the dominant D-term, which leads to chaotic inflation and subsequent reheating. No cosmic string is formed after inflation because the U(1) gauge symmetry is broken during inflation. The leptogenesis scenario via the inflaton decay in our D-term chaotic inflation scenario is also discussed.

PACS numbers: 98.80.Cq
I. INTRODUCTION

Among many types of inflation models proposed so far, chaotic inflation is special in that it can take place around the Planck time and make the universe large enough to avoid the recollapse [1], unless the universe is open at the beginning. On the other hand, other inflation models occur typically at much later times so that they suffer from the flatness (longevity) problem [2], namely, why the universe lives so long starting around the Planck scale till the low energy scale without the collapse. Furthermore, other types of inflation except chaotic and topological inflation also suffer from the initial condition problem [2, 3], that is, why the inflaton field is homogeneous over the horizon scale and takes a value which leads to a successful inflation.

For the analysis of a chaotic inflation model, one would need to consider the supergravity which would govern the dynamics of the early universe [4]. It, however, has been long considered challenging to realize chaotic inflation in supergravity simply because the F-term potential of a scalar field in supergravity has an exponential growth which prevents an inflaton from having an initial value much larger than the reduced Planck scale $M_p \simeq 2.4 \times 10^{18}$ GeV. Chaotic inflation is still possible in supergravity, however, by allowing a non-minimal Kähler potential [5, 6], even though it would be hard in general to justify a specific form of Kähler potential unless some symmetry such as the Nambu-Goldstone-like shift symmetry is introduced as done by Kawasaki, Yanagida and one of the present authors (M.Y.) for the natural chaotic inflation in supergravity [7].

Another possibility to circumvent the above difficulty stemming from the F-term is to use the D-term dominated inflaton potential because a D-term does not have an aforementioned dangerous exponential factor. In the conventional models of D-term inflation [12, 13], the energy density is sourced by the constant Fayet-Iliopoulos (FI) term $\xi$ in D-term while the inflaton trajectory follows the F-flat direction whose slope is induced by the one-loop corrections. Since the one-loop corrections cannot exceed the tree level potential energy density, the whole potential energy density is at most $\xi^2 \ll M_p^4$ so that the inflation cannot start from the Planckian energy scale in such conventional D-term inflation models dominated by the FI term. They thus necessarily give rise to hybrid inflation. On the other hand, in the model of D-term chaotic inflation proposed by two of the present authors (K.K. & Y.M) [14], the gauge non-singlet inflaton field identified with the (almost) F-flat direction is automatically lifted by the dominant D-term. Such a direction can thus naturally lead to D-term dominated chaotic inflation. It however cannot lead to the successful reheating in the original model because the gauge symmetry under which the inflaton is charged is unbroken and the charge conservation prohibits the coupling between the inflaton and the standard model fields. In this paper, we present a simpler model of the D-term chaotic inflation in supergravity which can induce the successful reheating where the gauge symmetry is spontaneously broken by the non-vanishing vacuum expectation value (VEV) of the inflaton. As a concrete illustration of a new D-term chaotic inflation scenario, we discuss the leptogenesis [15] by the inflaton decay into the right-handed Majorana neutrinos [8, 16] accompanied by the sufficient reheating.

In the next section, we present a new model of chaotic inflation in supergravity and discuss its dynamics and primordial fluctuations. In Sec. III we discuss the leptogenesis via

\[ \text{\footnotesize 1}\] The application of such a shift symmetry enables us to realize not only chaotic inflation [7, 8] but also its variants in supergravity [4, 11] and superstring [11].
the inflaton decay. The final section is devoted to the summary and discussion.

II. NEW D-TERM CHAOTIC INFLATION MODEL IN SUPERGRAVITY

We introduce three superfields $S, X, \overline{X}$ charged under $U(1)$ gauge symmetry and (global) $U(1)_R$ symmetry. The charges of the superfields are listed in Table I, which ensure our model is anomaly free [17, 18].

The general renormalizable superpotential for these fields is then given by

$$W = \lambda S(X\overline{X} - \mu^2),$$

where we can set the constants $\lambda$ and $\mu$ to be real and positive for simplicity. Note that non-renormalizable terms $S(X\overline{X})^n$ can appear in the superpotential. However, as shown later, $|X\overline{X}| \sim \mu^2$ during inflation so that such higher terms are negligible as long as $\mu \ll M_p$.

This leads to the following scalar potential consisting of the F-term $V_F$ and D-term $V_D$, along with the canonical Kähler potential $K(\Phi_i, \Phi_i^*) = \sum_i |\Phi_i|^2$ and the minimal gauge kinetic function $f_{ab}(\Phi_i) = \delta_{ab}$,

$$V = V_F + V_D,$$

$$V_F = \lambda^2 e^K \left[ |X\overline{X} - \mu^2|^2 (1 - |S|^2 + |S|^4) + |S|^2 \left\{ |\overline{X} + X^* (X\overline{X} - \mu^2)|^2 + |X + \overline{X}^* (X\overline{X} - \mu^2)|^2 \right\} \right],$$

$$V_D = \frac{g^2}{2} \left( |X|^2 - |\overline{X}|^2 \right)^2,$$

where $g$ is the coupling constant of the U(1) gauge interaction and we do not introduce the FI term for simplicity. Here and hereafter we set the reduced Planck scale $M_p$ to be unity and use the same symbols for the superfields and corresponding scalar fields unless stated otherwise.

The minima of the F-term (the F-flat condition) are given by

$$X\overline{X} - \mu^2 = 0, \quad S = 0,$$

and the minima of the D-term (the D-flat condition) are given by

$$|X| = |\overline{X}|.$$  

The global minima of the potential hence are given by

$$S = 0, \quad X = \mu e^{i\theta}, \quad \overline{X} = \mu e^{-i\theta},$$

|   | $S$ | $X$ | $X^*$ | $N_i$ | $H_u$ | $L_i$ |
|---|-----|-----|-------|------|------|------|
| $U(1)$ | 0   | +1  | −1    | 0    | 0    | 0    |
| $U(1)_R$ | 2   | 0   | 0     | +1   | +1   | 0    |

TABLE I: The $U(1) \times U(1)_R$ charge assignments for the superfields.
where the phase $\theta$ can be set to zero by the $U(1)$ gauge transformation.

One should notice that this superpotential and the corresponding scalar potential are the same as those of the conventional F-term hybrid inflation \cite{19} in which the gauge singlet field $S$ plays the role of an inflaton while $X$ and $\overline{X}$ remain zero during the inflation and then roll down to the global minima after the inflation. In order for the hybrid inflation to start, the field $S$ has to be relatively large but smaller than the reduced Planck scale $M_p$ due to the exponential factor in the F-term while $X$ and $\overline{X}$ almost vanish \cite{20}.

The notable difference between the new D-term chaotic inflation and the conventional F-term inflation is the initial condition. In the new model, inflation occurs when $|X| \gtrsim 1$ or $|\overline{X}| \gtrsim 1$ with $S \sim 0$ and $X\overline{X} \sim \mu^2$ which almost satisfy the F-flat condition. When the universe starts around the Planck scale, the potential energy as well as the kinetic energy is expected to be of order the Planck energy density. This requires the almost F-flat condition because all the fields quickly roll down to the global minimum because of the exponential factor, as will be seen below. The almost F-flat direction is thus naturally realized around the Planck scale. The potential is consequently dominated by the D-term potential, and thus chaotic inflation takes place.

We now investigate the dynamics of the new D-term chaotic inflation model in details. Despite the $e^K$ factor of F-terms, due to the presence of the relatively small but non-vanishing D-terms, the actual inflaton trajectory is slightly deviated from the exact F-flat direction and given by solving the equations (1) $\partial V / \partial S = \partial V / \partial |X| = 0$ or (2) $\partial V / \partial S^* = \partial V / \partial |X|^* = 0$, depending on the initial conditions ((1) for $|X| \gg 1$ or (2) for $|\overline{X}| \gg 1$). Note that the system is invariant under the interchange of $X$ and $\overline{X}$ so that both solutions (trajectories) lead to the essentially same dynamics. We therefore concentrate on the first trajectory given by $\overline{X} = \overline{X}(X), S = 0$, and we call this trajectory $T$.

We here confirm that inflation indeed occurs along this (almost F-flat) field trajectory. For this purpose, we first evaluate the mass terms of the field $S$ along the trajectory $T$, $V_{ij}|_T \phi_i^* \phi_j$ with $V_{ij} \equiv \partial^2 V / (\partial \phi_i^* \partial \phi_j)$. Here $\phi_i$ represents $S$, $X$, or $\overline{X}$. The suffix $T$ represents the evaluation along the trajectory $T$. Then, the mass matrix of the fields $S$, $V_{ij}|_T$, is given by

$$
V_{SS}|_T \simeq \lambda^2 e^K (|X|^2 + |\overline{X}|^2),
V_{SX}|_T = 0, \quad V_{s\overline{X}}|_T = 0.
$$

(6)

The effective squared mass of the field $S$ is much larger than the Hubble parameter squared $H^2 \simeq g^2 |X|^4 / 2$ unless the constant $\lambda$ is exponentially small, which makes $S$ quickly go to the zero. As a result, we can safely set $S$ to be zero and we can discuss the dynamics of the inflaton based on the following potential,

$$
V_{\text{eff}}(X, \overline{X}) \equiv V(X, \overline{X}, S = 0) = \lambda^2 e^K |X\overline{X}| - \mu^2 |X|^2 + \frac{g^2}{2} (|X|^2 - |\overline{X}|^2)^2.
$$

(7)

By use of the $U(1)$ gauge symmetry, we can, for instance, make the field $X$ real without loss of generality, so that the $\text{Im} \overline{X}$ rapidly goes to the zero because the effective mass squared of the imaginary part of $\overline{X}$ is given by $m_{\text{im} \overline{X}}^2 \simeq \lambda^2 e^K X^2$. We therefore consider the following effective potential, by redefining the fields $X \equiv \sqrt{2} \text{Re} X$, $\overline{X} \equiv \sqrt{2} \text{Re} \overline{X}$ (we take both $X$ and $\overline{X}$ to be positive for definiteness) and $\mu' \equiv \sqrt{2} \mu$,

$$
V_{\text{eff}}(X, \overline{X}) = \frac{\lambda^2}{4} e^K (X \overline{X} - \mu'^2)^2 + \frac{g^2}{8} \left( X^2 - \overline{X}^2 \right)^2.
$$

(8)
with $K = (X^2 + \overline{X}^2)/2$ and the canonical kinetic terms.

The dynamics of the inflation can be discussed based on the above potential. As mentioned before, we consider the case that $X \gg 1$ initially, which implies $\overline{X} \ll 1$ for $\mu' \ll 1$. For such a case, the trajectory $(T)$ is characterized by the condition $\partial V / \partial \overline{X} = 0$ which is equivalent to

$$X \overline{X} - \mu'^2 = e^{-K} \frac{g^2 (X^2 - \overline{X}^2)}{\overline{X} X + \frac{\lambda^2}{4} (X \overline{X} - \mu'^2)}$$

$$\simeq e^{-K} \frac{g^2 \overline{X}}{\lambda^2 X} (X^2 - \overline{X}^2).$$

Here we have used the fact that $X \overline{X} - \mu'^2 \sim O(e^{-K})$ and $X \gg \overline{X}$. The F-term contribution to the potential is then estimated as

$$V_F = \frac{X (X \overline{X} - \mu'^2)}{2X + \overline{X} (X \overline{X} - \mu'^2)} \frac{g^2}{2} (X^2 - \overline{X}^2)$$

$$< g^2 (X^2 - \overline{X}^2) \ll V_D = \frac{g^2}{8} (X^2 - \overline{X}^2)^2$$

for $X \gg 1$ and $\overline{X} \ll 1$. The potential thus is dominated by the D-term during inflation.

We also consider the mass matrix of the fields $X$ and $\overline{X}$, $V_{ij} | T(V_{ij} \equiv \partial^2 V_{\text{eff}}/(\partial \phi_i \partial \phi_j))$, whose elements are given by

$$V_{XX} | T \simeq \frac{\lambda^2}{2} e^K X^2 + g^2 (X^2 - \overline{X}^2) \overline{X}^2 + \frac{g^2}{2} \left(3X^2 - \overline{X}^2\right),$$

$$V_{X\overline{X}} | T \simeq \frac{\lambda^2}{2} e^K X \overline{X} + \frac{g^2}{2} (X^2 - \overline{X}^2) \overline{X} \left(1 + X^2 + \overline{X}^2\right) - g^2 X \overline{X},$$

$$V_{\overline{X}\overline{X}} | T \simeq \frac{\lambda^2}{2} e^K \overline{X}^2 + g^2 (X^2 - \overline{X}^2) \overline{X}^2 + \frac{g^2}{2} \left(3\overline{X}^2 - X^2\right),$$

up to the order of $O((e^K)^0)$. The effective squared masses, say $m^2$, of the fields $X$ and $\overline{X}$ are given as the solutions of the following equation

$$m^4 - (V_{XX} + V_{X\overline{X}})m^2 + V_{XX} V_{\overline{X}\overline{X}} - V_{X\overline{X}}^2 = 0,$$

where

$$V_{XX} + V_{X\overline{X}} | T \simeq \frac{1}{2} \lambda^2 e^K \left(X^2 + \overline{X}^2\right) \simeq \frac{1}{2} \lambda^2 e^K X^2,$$

$$V_{XX} V_{X\overline{X}} - V_{X\overline{X}}^2 | T \simeq \frac{3}{4} g^2 \lambda^2 e^K X^4,$$

up to the order of $O(e^K)$. Here, we have used the approximation that $X \overline{X} - \mu'^2 \sim O(e^{-K})$ and $X \gg 1 \gg \overline{X}$. The effective squared masses are then approximately given by

$$m^2 \simeq \frac{1}{2} \lambda^2 e^K X^2, \quad \frac{3}{2} g^2 X^2 \ll H^2 \simeq V_D / 3,$$

$^{2}$ More precisely, we have used the approximation that $|X \overline{X} - \mu'^2| \ll X / \overline{X}, X \overline{X}$. 5
where $H$ is a Hubble parameter. The inflaton field in the new D-term chaotic inflation under discussion corresponds to this effectively massless mode. This light squared mass vanishes for $g = 0$ as expected, reflecting the exact F-flat direction.

Since $V_{,XX} \gg V_{,X}$ and $X \gg \overline{X}$, the inflationary trajectory is given by the minimum of the field $\overline{X}$ represented by $\partial V / \partial \overline{X} = 0$, and we can write the minimum of $\overline{X}$ as a function of $X$, $\overline{X}^m = \overline{X}^m (X)$. The field trajectory governing the inflation dynamics therefore can be parameterized by the field $X$ which we call an inflaton. Then, by inserting the above relation into the effective potential, we define the reduced potential $V_i (X)$ as

$$V_i (X) \equiv V_{\text{eff}} (X, \overline{X}^m (X)) \left( \approx \frac{g^2}{8} X^4 \right).$$

As explicitly shown in Ref. [21], when there is only one massless mode and the other modes are massive, the generation of adiabatic density fluctuations as well as the dynamics of the homogeneous mode is completely determined by the reduced potential $V_i (X)$. Indeed, the equation of motion for the homogeneous mode of the inflaton $X$ along the rolling direction ($T$) is approximated as

$$\ddot{X} + 3H \dot{X} + \frac{\partial V_{\text{eff}}}{\partial X} \bigg|_T = \ddot{X} + 3H \dot{X} + \frac{dV_i}{dX} = 0,$$

where the dot represents time derivative. Thus, the dynamics of the inflation with the inflaton $X$ can be estimated by using the reduced potential $V_i (X)$ as long as the inflation dynamics follows the minimum of $\overline{X}$.

Next, we evaluate the primordial density fluctuations in the longitudinal gauge. The equations of motion for the perturbation $\delta X$ and $\delta \overline{X}$ are given by [22]

$$\ddot{\delta X} + 3H \dot{\delta X} - \frac{\nabla^2}{a^2} \delta X + V_{,XX} \big|_T \delta X + V_{,X \overline{X}} \big|_T \delta \overline{X} = -2 \frac{\partial V_{\text{eff}}}{\partial X} \bigg|_T X + 4 \dot{X} \Phi,$$

$$\ddot{\delta \overline{X}} + 3H \dot{\delta \overline{X}} - \frac{\nabla^2}{a^2} \delta \overline{X} + V_{,XX} \big|_T \delta X + V_{,X \overline{X}} \big|_T \delta \overline{X} = -2 \frac{\partial V_{\text{eff}}}{\partial \overline{X}} \bigg|_T \overline{X} + 4 \dot{\overline{X}} \Phi,$$

where $\Phi$ is the gravitational potential. We hereafter use the same symbols $X$ and $\overline{X}$ for both the homogeneous modes and the full fields for the notational brevity unless stated otherwise.

We are interested only in the adiabatic density fluctuations characterized by the condition

$$\frac{\delta X}{X} = \frac{\delta \overline{X}}{\overline{X}} \quad \iff \quad \delta \overline{X} = \frac{d \overline{X}^m (X)}{dX} \delta X$$

where we have used $\overline{X} / \dot{X} = d \overline{X}^m (X) / dX$. Since the relation $\partial V_{\text{eff}} (X, \overline{X}^m (X)) / \partial \overline{X} = 0$ holds for any $X$ in the regime of our interest, we find

$$\frac{d}{dX} \left[ \frac{\partial V_{\text{eff}}}{\partial \overline{X}} (X, \overline{X}^m (X)) \right] = V_{,XX} \big|_T + V_{,X \overline{X}} \frac{d \overline{X}^m}{dX} \big|_T = 0.$$
Then, using this relation, the equation of motion for the perturbation $\delta X$ can be rewritten as
\[
\ddot{\delta X} + 3H\dot{\delta X} + \frac{\nabla^2}{a^2}\delta X + \left.\frac{V_{,XX}V_{,XX} - V_{,X}^2}{V_{,XX}}\right|_T \delta X = -2\left.\frac{\partial V_{\text{eff}}}{\partial X}\right|_T \dot{\Phi} + 4X\dot{\Phi}.
\] (20)

Taking into account this relation and
\[
d^2V_t/dX^2 = V_{,XX} + 2\frac{dX^m}{dX}V_{,XX} + \left(V_{,XX}ight)^2 \left.\frac{V_{,XX}V_{,XX} - V_{,X}^2}{V_{,XX}}\right|_T \left(\sim \frac{3}{2}g^2X^2\right),
\] (21)
the equation of motion for the perturbation $\delta X$ becomes
\[
\ddot{\delta X} + 3H\dot{\delta X} - \frac{\nabla^2}{a^2}\delta X + \frac{d^2V_t}{dX^2}\delta X = -2\frac{dV_t}{dX}\Phi + 4X\Phi.
\] (22)

The perturbation $\delta X$ is hence completely determined by the reduced potential $V_t(X)$. Note that $d^2V_t/dX^2$ is the effective mass squared along the rolling direction given by $\partial V_{\text{eff}}/\partial X = 0$, and this rolling direction actually coincides with the eigenvector of the effectively massless mode of $\lambda$.

On the other hand, using the adiabatic condition, the gravitational potential is described only by $\delta \phi_1$ as
\[
\left(H - \frac{\nabla^2}{a^2}\right)\Phi = \frac{1}{2M_G^2} \left(\dot{X}\delta X - \dot{\delta X}\right) \left[1 + \left(\frac{dX^m}{dX}\right)^2\right].
\] (23)

Consequently, the relation
\[
\dot{H} = -\frac{\dot{X}^2 + \ddot{X}^2}{2M_G^2} = -\frac{\dot{X}^2}{2M_G^2} \left[1 + \left(\frac{dX^m}{dX}\right)^2\right],
\] (24)
leads to the gravitational potential in the long wave limit
\[
\Phi = \frac{d}{dt}\left(\frac{\delta X}{X}\right).
\] (25)

This expression of the gravitational potential coincides with that of the single field inflation with the reduced potential $V_t(X)$. We can as a result calculate the density fluctuations of our inflation model based on the reduced potential $V_t(X) \sim g^2X^4/8$, which implies that the gauge coupling $g$ should be $g \sim 10^{-6}$ in order to explain the primordial density fluctuations.

After inflation, the inflaton oscillates around the global minimum $X = \overline{X} = \mu'$ and $S = 0$. Note that $S$ remains to vanish. Though the inflaton rolls down almost along the direction of $X$ during inflation, the trajectory after inflation is curved so that both $X$ and $\overline{X}$ oscillate around the global minimum. The mass matrix around the minimum is given by
\[
\begin{align*}
V_{,XX}|_M &= \frac{1}{2}\lambda^2 e^K \mu'^2 + g^2\mu'^2, \\
V_{,X\overline{X}}|_M &= \frac{1}{2}\lambda^2 e^K \mu^2 - g^2\mu^2, \\
V_{,X\overline{X}}|_M &= \frac{1}{2}\lambda^2 e^K \mu'^2 + g^2\mu'^2,
\end{align*}
\] (26)
where the suffix $M$ represents the evaluation at the global minimum $M$. Hence, the effective squared masses $m_\pm^2$ are given by $m_+^2 = \lambda^2 e^K \mu^2$ with the eigendirection $X_+ \equiv (X + \overline{X})/\sqrt{2}$ and $m_-^2 = 2g^2 \mu^2$ with the eigendirection $X_- \equiv (X - \overline{X})/\sqrt{2}$.

### III. REHEATING AND LEPTOGENESIS IN NEW D-TERM CHAOTIC INFLATION

Let us now discuss the issues of the reheating and the baryon asymmetry of the universe in the new D-term chaotic inflation through a concrete example, namely, the non-thermal leptogenesis scenario via the inflaton decay which produces the the heavy Majorana neutrinos $N_i$ non-thermally [16].

We consider the right handed Majorana neutrinos in addition to the Minimal Supersymmetric Standard Model (MSSM) fields in the superpotential

$$W = \lambda S (X \overline{X} - \mu^2) + \sum_i \alpha_i X \overline{X} N_i N_i + \sum_{i, j} h_{ij}^\nu N_i L_j H_u + W^{\text{MSSM}},$$  \hspace{1cm} (27)

where the subscripts $i$ and $j$ represent the generation indices, $h_{ij}^\nu$ is the Yukawa coupling, $L_j$ is the lepton doublet, $H_u$ is the up-type Higgs and $W^{\text{MSSM}}$ is the superpotential of the MSSM. The charges for various supermultiplets are given in Table II. Note that non-renormalizable terms like $N_i N_i (X \overline{X})^n$ and $S(X \overline{X})^n$ can appear in the superpotential but they are negligible as long as $\mu \ll 1$ because $|X \overline{X}| \sim \mu^2$ during inflation.

Taking the canonical Kähler potential, the minimal gauge kinetic function and the vanishing FI term, we can calculate the scalar potential consisting of the F-term $V_F$ and D-term $V_D$

$$V = V_F + V_D,$$

$$V_F = e^K \left[ \lambda (X \overline{X} - \mu^2) + S^* W \right]^2 + \left| \lambda S \overline{X} + \sum_i i \alpha_i \overline{X} N_i N_i + X^* W \right|^2$$

$$+ \left| \lambda S X + \sum_i i \alpha_i X N_i N_i + \overline{X}^* W \right|^2 + \sum_i \left| 2 \alpha_i X \overline{X} N_i + \sum_j h_{ij}^\nu L_j H_u + N_i^* W \right|^2$$

$$+ \sum_k \left| \frac{\partial W}{\partial \xi_k} + \xi_k^* W \right|^2 - 3|W|^2 \right],$$

$$V_D = \frac{g^2}{2} \left( |X|^2 - |\overline{X}|^2 \right)^2 + V_D^{\text{MSSM}}, \hspace{1cm} (28)$$

where $\xi_k$ represents the MSSM field and $V_D^{\text{MSSM}}$ represents the D-term concerning the standard model gauge group. The minimum of the F-term (the F-flat condition) is given by

$$V_F = 0 \iff X \overline{X} - \mu^2 = 0 \quad \& \quad S = 0 \quad \& \quad N_i = 0 \quad \& \quad \text{MSSM F-flat condition.} \hspace{1cm} (29)$$

On the other hand, the minimum of the D-term (the D-flat condition) is given by

$$V_D = 0 \iff |X| = |\overline{X}| \quad \& \quad \text{MSSM D-flat condition.} \hspace{1cm} (30)$$

The global minimum of the potential hence is given by

$$S = 0, \quad X = \mu e^{i\theta}, \quad \overline{X} = \mu e^{-i\theta}, \quad \text{MSSM flat condition} \hspace{1cm} (31)$$
where the phase $\theta$ is arbitrary. As was done in the previous section, one can show that $S$ and $N_i$ go to the zeros dynamically during the inflation so that the effective dynamics is described by the effective potential

$$V_{\text{eff}}(X, \overline{X}) = \frac{\lambda^2}{4} e^K (X \overline{X} - \mu'^2)^2 + \frac{g^2}{8} \left( X^2 - \overline{X}^2 \right)^2,$$

where we have used the redefined fields $X \equiv \sqrt{2} \Re X$ and $\overline{X} \equiv \sqrt{2} \Re \overline{X}$. Thus, the dynamics and primordial fluctuations of the inflation are essentially unchanged even if we include other fields besides $X$ and $\overline{X}$.

Let us now discuss the reheating and leptogenesis in this model. Note that the spontaneous breaking of the gauge symmetry due to the non-vanishing VEV of the inflaton is crucial because, if the inflaton VEV vanishes, the charge conservation would prohibit the inflaton decay in our model. The Majorana masses of right-handed neutrinos $M_i$ are given by

$$M_i = \alpha_i \langle X \overline{X} \rangle / 2 = \alpha_i \mu'^2 / 2.$$  (32)

The produced $N_i$ decay into leptons $L_j$ and Higgs doublets $H_u$ through the Yukawa interactions via

$$W = h_{\nu}^{ij} N_i L_j H_u$$  (35)

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 $N_1$ assuming that the mass $M_1$ is much smaller than the others, $M_2$ and $M_3$. The interference between the tree-level and the one-loop diagrams including vertex and self-energy corrections generates the lepton asymmetry $\epsilon_1$,

$$\epsilon_1 \equiv \frac{\Gamma(N_1 \to H_u + l) - \Gamma(N_1 \to \overline{H}_u + \overline{l})}{\Gamma(N_1 \to H_u + l) + \Gamma(N_1 \to \overline{H}_u + \overline{l})} = - \frac{3}{8\pi} \left( h_{\nu} h_{\nu}^{\dagger} \right)_{11} \sum_{i=2,3} \text{Im} \left( h_{\nu} h_{\nu}^{\dagger} \right)_{11} \frac{M_1}{M_i}$$  (33)

The reheating temperature $T_R$ hence becomes

$$T_R \simeq 0.1 \sqrt{\Gamma \langle X_+ \to N_i \rangle} \sim 10^7 \text{ GeV} \left( \frac{\alpha}{0.1} \right) \left( \frac{\lambda}{10^{-4}} \right) \left( \frac{\mu'}{10^{14} \text{ GeV}} \right)^{3/2}$$  (34)

with $\alpha \equiv \sqrt{\sum_i \alpha_i^2}$ where the sum $i$ runs for only the generations of right-handed neutrinos which the inflaton $X_+$ can decay into.
\[
\sum \frac{3}{8\pi} \frac{M_1}{(H_u)^2} m_{\nu_3} \delta_{\text{eff}} \\
\sim 10^{-7} \left( \frac{M_1}{10^9 \text{GeV}} \right) \left( \frac{m_{\nu_3}}{10^{-2} \text{eV}} \right) \delta_{\text{eff}},
\]

where the effective CP violation phase is

\[
\delta_{\text{eff}} \equiv -\text{Im} \left[ h_{\nu} (m_\nu^*) h_T^T \right]_{11}.
\]

\(m_{\nu_3}\) here is a mass eigenvalue of the left-handed neutrino mass matrix \(m_\nu\) estimated by the seesaw mechanism \([26]\) as, based on our assumption that the \((h_\nu)_{33}\) is the dominant entry in \(h_\nu\) and \(M_3 \gg M_1\),

\[
m_{\nu_3} \simeq \frac{|(h_\nu)_{33}|^2 (H_u)^2}{M_3} \\
\sim 10^{-2} \text{eV} \left( \frac{|(h_\nu)_{33}|}{5 \times 10^{-3}} \right)^2 \left( \frac{M_3}{10^{10} \text{GeV}} \right)^{-1},
\]

which is consistent with the mass inferred from the Super-Kamiokande experiments \([27]\) for \(|(h_\nu)_{33}| \sim 10^{-2}\) and \(M_3 \sim 10^{10}\) GeV.

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

\[
\Gamma_{N_1} = \Gamma(N_1 \to H_u + l) + \Gamma(N_1 \to \overline{H}_u + \overline{l}) \\
\simeq \frac{1}{8\pi} \Sigma |(h_\nu)_{11}|^2 M_1 \\
\simeq \frac{1}{8\pi} |(h_\nu)_{13}|^2 M_1 \\
\sim 10 \text{ GeV} \left( \frac{|(h_\nu)_{13}|}{5 \times 10^{-6}} \right)^2 \left( \frac{M_1}{10^9 \text{GeV}} \right).
\]

Thus, for a wide range of parameters, the decay rate \(\Gamma_{N_1}\) is much larger than the decay rate of the inflaton \(\Gamma_{(X_+ - N_1)}\) so that the produced \(N_1\) immediately decays into leptons and Higgs supermultiplets.

Before estimating the lepton asymmetry produced in our model, let us evaluate the lepton asymmetry needed to explain the observed baryon number density. A part of the produced lepton asymmetry (or, exactly speaking, \(B - L\) asymmetry) is converted into the baryon asymmetry through the sphaleron processes, which can be estimated as \([28]\)

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

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

\[
\frac{n_B}{s} \simeq (0.1 - 1) \times 10^{-10},
\]
we need the lepton asymmetry,
\[ \frac{n_L}{s} \simeq -(0.3 - 3) \times 10^{-10}. \] (42)

Now we estimate the lepton asymmetry produced through the inflaton decay. For \( M_1 \gtrsim 10^9 \text{ GeV} \), \( M_1 \) is one hundred times 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
\[
\frac{n_L}{s} \simeq \frac{3}{2} \epsilon_1 B_r \frac{T_R}{m_+} \left( \frac{M_1}{10^9 \text{ GeV}} \right) \left( \frac{T_R}{10^7 \text{ GeV}} \right) \left( \frac{m_+}{10^{10} \text{ GeV}} \right)^{-1} \sim -10^{-10} \delta_{\text{eff}} B_r \left( \frac{\lambda}{10^{-4}} \right)^{-\frac{3}{2}} \left( \frac{\alpha}{0.1} \right) \left( \frac{\alpha_1}{0.1} \right) \left( \frac{\mu'}{10^{14} \text{ GeV}} \right)^{\frac{3}{2}},
\] (43)

where \( B_r \) is the branching ratio of the inflaton \( X_+ \) into \( N_1 \). For \( M_3 \sim M_2 \sim m_+ \sim 10^{10} \text{ GeV} \) with \( \alpha_2 \sim \alpha_3 = \mathcal{O}(1) \), \( m_+ \) is comparable to \( M_2 \) and \( M_3 \) so that the decay into \( N_2 \) and \( N_3 \) are prohibited kinematically or suppressed by the phase space and hence \( B_r \approx \mathcal{O}(1) \). Note also that since \( m_- \sim 10^8 \text{ GeV} \) for \( g \sim 10^{-6} \) and \( \mu' \sim 10^{14} \text{ GeV} \), \( m_- \ll M_1, M_2, M_3 \) so that the decays of \( X_- \) into all the right-handed neutrinos are prohibited kinematically, as assumed previously. In this case, \( \alpha = \alpha_1 \sim 0.1 \) and \( T_R \sim 10^7 \text{ GeV} \) (low enough reheating temperature to avoid the gravitino problem [29]) resulting in \( n_L/s \sim 10^{-10} \delta_{\text{eff}} \) which is consistent with the baryon number density in the present universe.

IV. SUMMARY AND DISCUSSION

In this paper we have proposed a new model of D-term dominated chaotic inflation in supergravity. The F-flat direction present in this model is automatically lifted by the D-term, which leads to chaotic inflation. The superpotential of our model was originally proposed as the F-term hybrid inflation where the gauge singlet field \( S \) plays the role of an inflaton with the vanishing \( X \) and \( \overline{X} \) during inflation. On the other hand, we showed that another initial condition such that \( S \sim 0, |X| \gg 1 \) or \( |X| \gg 1 \) with \( XX \sim \mu^2 \) can naturally occur around the Planck scale for a successful D-term chaotic inflation. In contrast to the previously proposed D-term chaotic inflation model [14], the inflaton can acquire the non-vanishing VEV after inflation which breaks the gauge symmetry spontaneously so that it can decay into the visible sector for the sufficient reheating\(^4\). No cosmic string is formed after inflation because the U(1) gauge symmetry is broken during inflation, while such a cosmic string formation in the conventional D-term inflation is often problematic [31].

Our model leads to the quartic potential chaotic inflation which has the tight constraints from the recent observations [32]. The relaxation of such constraints is possible by, for instance, an appropriate choice of the non-minimal gauge kinetic function such as a form

\(^4\) It was recently pointed out that the supergravity effects induce the inflaton decay into the visible sector when the inflaton acquires the non-vanishing VEV after inflation [30], and an analogous mechanisms could help reheat the universe for a D-term chaotic inflation model as well.
\[ f = 1 + d_X |X|^2 + d_2 |X|^2 \quad (d_X, d_2: \text{constants}) \] which gives a quadratic inflaton potential. One of the present authors (T.K.) also proposed the quadratic potential chaotic inflation by use of the FI field \[ \text{[33]} \] even though a successful reheating needs more care \[ \text{[34]} \]. The consideration of the primordial fluctuations from a MSSM flat direction acting as a curvaton in our model could be another possibility for a viable quartic potential chaotic inflation model \[ \text{[35, 36]} \]. We also mention that, for a toy model using a minimal gauge kinetic function illustrated in our discussion, the gauge coupling \( g \) should be \( g \sim 10^{-6} \) in order to explain the primordial density fluctuations. This value of the gauge coupling is much smaller than the standard gauge couplings. However, this may not be a problem because the gauge symmetry may be a hidden gauge symmetry, or the gauge coupling could be suppressed, for instance, by considering the extra dimensions. This topic would be worth further investigation.

We have also discussed the leptogenesis scenario via the inflaton decay in this chaotic inflation model, which can explain the observed baryon asymmetry for a reasonable parameter set consistent with the data from the neutrino experiments.

**Acknowledgments**

We thank A. D. Linde, R. Kallosh, Y. Shinbara, F. Takahashi, and J. Yokoyama for useful discussions. K.K. is supported by DOE grant DE-FG02-94ER-40823. T.K. was supported in part by a Grant-in-Aid (#19540268) from the MEXT of Japan. M.Y. is supported in part by JSPS Grant-in-Aid for Scientific Research No. 18740157 and No. 19340054.

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