Natural double inflation in supergravity

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We propose a natural double inflation model in supergravity. In this model, chaotic inflation first takes place by virtue of the Nambu-Goldstone-like shift symmetry, which guarantees the absence of the exponential factor in the potential for the inflaton field. During chaotic inflation, an initial value of the second inflation (new inflation) is set. In this model, the initial value of new inflation can be adequately far from the local maximum of the potential for new inflation due to the small linear term of the inflaton in the Kähler potential. Therefore, the primordial fluctuations within the present horizon scale may be attributed to both inflations; that is, the first chaotic inflation produces the primordial fluctuations on the large cosmological scales while the second new inflation on the smaller scales. The successive decay of the inflaton for new inflation leads to a reheating temperature low enough to avoid the overproduction of gravitinos in a wide range of the gravitino mass.

I. INTRODUCTION

The fascinating feature of inflation is a generation of primordial density fluctuations in addition to solving the flatness and the horizon problems of the standard big bang cosmology. A single inflation model generally predicts adiabatic fluctuations with a nearly scale-invariant spectrum. It is, however, pointed out that a standard cold dark matter model with a nearly scale-invariant spectrum cannot account for several observational results. The predicted primordial fluctuations have too much power on small scales after they are normalized by the cosmic background explorer (COBE) data. The power spectrum derived from the APM redshift survey has a break around the scale $k \sim 0.05h$ Mpc^{-1} though there is large uncertainty [1]. Furthermore, while the recent observations of anisotropies of the cosmic microwave background (CMB) by the BOOMERANG experiment [2] and the MAXIMA experiment [3] confirmed a flat universe with an inflationary scenario, they found the peculiar feature, that is, a relatively low second acoustic peak. Thus, the demand for the primordial fluctuations with a non-trivial shape is increasing.

Such primordial fluctuations are realized in a double inflation model. In fact, double inflation models have been considered to reconcile the predicted spectra with the observations [4–6]. Many of them were discussed in a simple toy model with two massive scalar fields. Though the model predicts an interesting nontrivial spectrum, a more realistic double inflation model should be constructed on the basis of the particle physics theory. Supersymmetry (SUSY) is one of the most powerful extensions of the standard model of particle physics, which is also indispensable to constructing inflation models. Because SUSY guarantees the flatness of the inflaton and gives a natural solution to the hierarchy problem between the inflationary scale and the electroweak scale. Therefore, it is very important to construct double inflation models in the context of SUSY, especially, its local version, supergravity (SUGRA) [7,8]. Until now, a few double inflation models in SUSY or SUGRA have been proposed [5,6,9]. First of all, double hybrid inflation models in SUSY [10] or SUGRA [11] have been proposed in order to solve the initial

condition problem of hybrid inflation [12]. Later, it is shown that they can produce the primordial fluctuations with the non-trivial feature as observed [5]. But, generally speaking, the reheating temperature after hybrid inflation is so high that it causes the overproduction of gravitinos. On the other hand, Izawa *et al.* [13] proposed preinflation as a solution to the initial condition problem of new inflation, which straightforwardly predicts the low reheating temperature [14]. If hybrid inflation is taken as preinflation, double inflation is realized through the cross term of the superpotential. Later, it has been discussed in the context of large scale structure, CMB [6], and primordial black holes (PBHs) formation [9].

In all of the earlier double inflation models in SUSY or SUGRA, hybrid inflation is adopted as the first inflation responsible for the COBE scale. However, such hybrid inflation is known to suffer from the severe initial condition problem [12]. Hence, another hybrid inflation is required, which takes place near the Planck scale and the produced fluctuations cannot be observed [10,11].[1] After all, triple inflation is required. On the other hand, chaotic inflation [15] is free from any initial condition problems including the flatness problem because it starts around the Planck scale. Nonetheless, one of the reasons why hybrid inflation has often been considered as an inflation model in SUGRA lies in difficulty realizing chaotic inflation in SUGRA. The minimal SUGRA potential for scalar fields has an exponential factor of the form $\exp(|\phi|^2/M_G^2)$, which prevents any scalar field ϕ from taking a value much larger than the reduced Planck scale $M_G \simeq 2.4 \times 10^{18}$ GeV. Several chaotic inflation models have been proposed so far by use of the functional degrees of freedom in SUGRA [16,17]. Rather specific Kähler potentials were adopted without symmetry reasons so that we need the fine tuning. Recently, Kawasaki *et al.* [18] proposed a natural model of chaotic inflation in SUGRA, where the

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[1]Exactly speaking, even if one considers such pre-hybrid inflation, it suffers from the so-called flatness (longevity) problem, that is, why the universe lives beyond the Planck time (though it is milder than the original one) as long as it starts below the Planck scale.

Kähler potential is restricted by the Nambu-Goldstone-like shift symmetry.

In this paper, we propose a natural double inflation model in SUGRA by use of the Nambu-Goldstone-like shift symmetry. In this model, first of all, chaotic inflation takes place. During chaotic inflation, an initial value of the second inflation (new inflation) is set due to the supergravity effects. A similar model has already been proposed by Yokoyama and Yamaguchi [19] (see also Ref. [20]). But, in the model, the initial value of new inflation is so close to the local maximum of the potential for new inflation, which causes the universe to enter a self-regenerating stage [21,22]. Hence, the primordial fluctuations responsible for the observable scale are produced only during the last inflation (new inflation) so that their spectra are only tilted scale-free ones. Also, you should notice that even if the chaotic inflation proposed in Ref. [18] is taken as preinflation in Ref. [13], the same situation occurs, that is, the second inflation becomes eternal inflation because the superpotential in Ref. [18] vanishes during chaotic inflation. On the other hand, in our model, the initial value of new inflation can be adequately far from the local maximum of the potential due to the small linear term of the inflaton in the Kähler potential. Therefore, the primordial fluctuations responsible for the observable universe may be attributed to both inflations, that is, the first chaotic inflation produces the primordial fluctuations on the large cosmological scales while the second new inflation on the smaller scales. The fact that the last inflation is new inflation is favorable because it straightforwardly predicts sufficiently low reheating temperature to avoid the overproduction of gravitinos. Furthermore, our model is simple in that two inflatons belong to the same supermultiplet, namely, one direction of a complex scalar field drives chaotic inflation while another drives new inflation. Also, our model is natural for two reasons. The form of Kähler potential is completely determined by a symmetry, that is, the Nambu-Goldstone-like shift symmetry. Next, though we need the introduction of small breaking parameters, the smallness of parameters is justified also by symmetries. That is, the zero limit of small parameters recovers symmetries, which is natural in the 't Hooft's sense [23].

In the next section, we present our double inflation model in supergravity. In Sec. III, the dynamics is investigated and the primordial density fluctuations are estimated. In the final section, we give discussions and a conclusion.

II. MODEL

We introduce an inflaton chiral superfield $\Phi(x,\theta)$ and a spurion superfield Ξ. We assume that the model is invariant under the following Nambu-Goldstone-like shift symmetry [18]:

$$\Phi \rightarrow \Phi + iCM_G,$$

$$\Xi \rightarrow \left(\frac{\Phi}{\Phi+iCM_G}\right)^2 \Xi, \qquad (1)$$

TABLE I. The charges of the $U(1)_R \times Z_2$ symmetries for the various supermultiplets.

	Φ	X	Ξ	Π
Q_R	0	2	0	0
Z_2	$-$	$-$	$-$	$-$

where C is a dimensionless real constant and M_G is the reduced Planck scale. That is, the combination $\Xi\Phi^2$ is invariant under the shift symmetry. Then, the Kähler potential is a function of $\Phi+\Phi^*$, i.e., $K(\Phi,\Phi^*)=K(\Phi+\Phi^*)$, which allows the imaginary part of the scalar components of Φ to take a value much larger than the gravitational scale. Later we set M_G to be unity.

Next, let us discuss the form of the superpotential. We assume that the superpotential is also invariant under the $U(1)_R$ symmetry, which prohibits a constant term in the superpotential. Then, the earlier Kähler potential is invariant only if the R charge of Φ is zero, which compels us to introduce another supermultiplet $X(x,\theta)$ with its R charge equal to two. The general superpotential invariant under the shift and the $U(1)_R$ symmetries is given by

$$W = vX[1 + \alpha_2(\Xi\Phi^2)^2 + \cdots]$$
$$- X[\Xi\Phi^2 + \alpha_3(\Xi\Phi^2)^3 + \cdots], \qquad (2)$$

where v, α_i are complex constants and we have assumed the R charge of Ξ vanishes. Generally speaking, among all complex constants, only one constant can become real by use of the phase rotation of the X field. Later we set v to be real. The shift symmetry is softly broken by inserting the vacuum value into the spurion field, $\langle\Xi\rangle=\lambda$. The parameter λ is fixed at a value whose absolute magnitude is much smaller than unity, representing the magnitude of breaking of the shift symmetry [Eq. (1)]. As long as $|\Phi|\ll|\lambda|^{-1/2}$, higher order terms with α_i of the order of unity become irrelevant for the dynamics of the chaotic inflation. Thus, we can safely neglect them in the following discussion. As shown later, for successful inflation, the constant v must be at most of the order of $|\lambda|$, which is much smaller than unity. Since the constant v is of the order of unity in general, we must invoke the mechanism to suppress the constant v. For the purpose, we introduce the Z_2 symmetry, under which the Φ, X, and Ξ fields are odd.[2,3] (See Table I in which charges for superfields are shown.) Then, the smallness of the constant v is associated with the small breaking of the Z_2 symmetry. That is, we introduce a spurion field Π with the odd Z_2 charge and the

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[2]The fact that the Z_2 charge of Φ is odd is essential for this model because it allows the small linear term of $\Phi+\Phi^*$ in the Kähler potential so that the initial value of new inflation can appropriately deviate from the local maximum of the potential for new inflation.

[3]Note that the spurion field Ξ in fact breaks both the shift symmetry and the Z_2 symmetry at once. So, we expect that the magnitudes of the breaking of both the Z_2 and the shift symmetries are of the same order, that is, $|g|=\mathcal{O}(1)$.

zero R charge, whose vacuum value $\langle\Pi\rangle=v$ softly breaks the Z_2 symmetry. You should notice that though the above superpotential is not invariant under the shift and the Z_2 symmetries, the model is completely natural in 't Hooft's sense [23] because we have enhanced symmetries in the limits $\lambda\rightarrow0$ and $v\rightarrow0$. In the following analysis, we use the superpotential given by

$$W\simeq vX-\lambda X\Phi^2, \qquad (3)$$

$$=vX(1-g\Phi^2), \qquad (4)$$

with $g\equiv\lambda/v$.

The Kähler potential invariant under the shift and the $U(1)_R$ symmetries is given by

$$K=v_2(\Phi+\Phi^*)+\frac{1}{2}(\Phi+\Phi^*)^2+XX^*+\cdots. \qquad (5)$$

Here $v_2\sim v$ is a real constant representing the breaking effect of the Z_2 symmetry. The term $v_3\lambda_3\Phi^2+v_3^*\lambda_3^*\Phi^{*2}$ may appear, where v_3 and λ_3 are complex constants representing the breaking of the Z_2 and the shift symmetries ($|v_3|\sim v$ and $|\lambda_3|\sim|\lambda_4|\sim|\lambda|$). But, these terms are extremely small so that we can safely omit them in the Kähler potential [Eq. (5)]. A constant term is also omitted because it only changes the overall factor of the potential, whose effect can be renormalized into the constants λ and δ_1. Here and hereafter, we denote the scalar components of the supermultiplets by the same symbols as the corresponding supermultiplets.

III. DYNAMICS OF DOUBLE INFLATION AND PRIMORDIAL DENSITY FLUCTUATIONS

Neglecting higher order terms and a constant term in the Kähler potential, the Lagrangian density $L(\Phi,X)$ for the scalar fields Φ and X is given by

$$L(\Phi,X)=\partial_\mu\Phi\partial^\mu\Phi^*+\partial_\mu X\partial^\mu X^*-V(\Phi,X), \qquad (6)$$

where the scalar potential V of the chiral superfields $X(x,\theta)$ and $\Phi(x,\theta)$ is given by

$$V=v^2e^K[|1-g\Phi^2|^2(1-|X|^2+|X|^4)$$
$$+|X|^2]-2g\Phi+(v_2+\Phi+\Phi^*)(1-g\Phi^2)|^2]. \qquad (7)$$

Decomposing the scalar field Φ and the complex constant g into real and imaginary components

$$\Phi=\frac{1}{\sqrt{2}}(\varphi+i\chi), \qquad (8)$$

$$g=g_R+ig_I, \qquad (9)$$

the Lagrangian density $L(\varphi,\chi,X)$ is written as

$$L(\varphi,\chi,X)=\frac{1}{2}\partial_\mu\varphi\partial^\mu\varphi+\frac{1}{2}\partial_\mu\chi\partial^\mu\chi+\partial_\mu X\partial^\mu X^*$$
$$-V(\varphi,\chi,X), \qquad (10)$$

with the potential $V(\varphi,\chi,X)$ given by

$$V(\varphi,\chi,X)=v^2e^{-v_2^2/2}\exp\left[\left(\varphi+\frac{v_2}{\sqrt{2}}\right)^2+|X|^2\right]([1-g_R(\varphi^2-\chi^2)+2g_I\varphi\chi+\tfrac{1}{4}(g_R^2+g_I^2)(\varphi^2+\chi^2)^2](1-|X|^2+|X|^4)$$

$$+|X|^2\{2(g_R^2+g_I^2)(\varphi^2+\chi^2)-(v_2+\sqrt{2}\varphi)\{\sqrt{2}(g_R\varphi-g_I\chi)[2-g_R(\varphi^2-\chi^2)+2g_I\varphi\chi]$$

$$-\sqrt{2}(g_R\chi+g_I\varphi)[g_I(\varphi^2-\chi^2)+2g_R\varphi\chi]\}+(v_2+\sqrt{2}\varphi)^2[1-g_R(\varphi^2-\chi^2)+2g_I\varphi\chi+\tfrac{1}{4}(g_R^2+g_I^2)(\varphi^2+\chi^2)^2]\}). $$

$$(11)$$

A. Chaotic inflation

Though the potential is very complicated, the dynamics is not so. While $\varphi,|X|\lesssim\mathcal{O}(1)$ due to the factor $e^{v_2\varphi+\varphi^2+|X|^2}$ ($v_2\ll1$), χ can take a value much larger than unity without costing exponentially large potential energy. Then, the scalar potential is approximated as

$$V\simeq|\lambda|^2\left(\frac{\chi^4}{4}+2\chi^2|X|^2\right), \qquad (12)$$

with $|\lambda|^2=(g_R^2+g_I^2)v^2$. Thus, the term proportional to χ^4 becomes dominant, which leads to chaotic inflation starting around the Planck epoch. Using the slow-roll approximations, we obtain the e-fold number N_c,

$$N_c\simeq\frac{\chi_{N_c}^2}{8}. \qquad (13)$$

During chaotic inflation, the potential minimum for φ, $\varphi_{\min}$, is given by

$$\varphi_{\min}\simeq-v_2/\sqrt{2}-\frac{g_I}{g_R^2+g_I^2}\frac{4}{\chi^3}, \qquad (14)$$

and the mass squared of φ, m_φ^2, becomes

$$m_\varphi^2\simeq\frac{|\lambda|^2}{2}\chi^4\simeq6H^2\gg\frac{9}{4}H^2, \quad H^2\simeq\frac{|\lambda|^2}{12}\chi^4, \qquad (15)$$

where H is the hubble parameter at that time. Hence, φ oscillates rapidly around the minimum $\varphi_{\min}$ with its amplitude damping in proportion to $a^{-3/2}$ (a: the scale factor of the universe). Thus, at the end of chaotic inflation, φ settles down to the minimum $\varphi_{\min}$.

On the other hand, the mass squared of X, m_X^2, is dominated by

$$m_X^2 \simeq 2|\lambda|^2\chi^2 \simeq \frac{24}{\chi^2}H^2, \tag{16}$$

which is much smaller than the hubble parameter squared in the early stage of chaotic inflation so that X also slow-rolls towards the origin. Later we set X to be real and positive making use of the freedom of the phase choice. In this regime, the classical equations of motion for the X and χ fields are given by

$$3H\dot{X} \simeq -m_X^2 X, \tag{17}$$

$$3H\dot{\chi} \simeq -|\lambda|^2\chi^3, \tag{18}$$

which yield

$$X \propto \chi^2. \tag{19}$$

This relation holds actually if and only if quantum fluctuations are unimportant for both χ and X. However, following the same procedure done in Refs. [18,19], we can easily show that for X, quantum fluctuations are smaller than the classical value and X is much smaller than unity throughout chaotic inflation.

We investigate the density fluctuations produced by this chaotic inflation. As shown above, there are the two effectively massless fields, χ and X. Using Eq. (12) and adequate approximations, the metric perturbation in the longitudinal gauge Φ_A can be estimated as [24]

$$\Phi_A = -\frac{\dot{H}}{H^2}C_1 - 16\frac{X^2}{\chi^2}C_3,$$

$$C_1 = H\frac{\delta\chi}{\dot{\chi}},$$

$$C_3 = H\left(\frac{\delta\chi}{\dot{\chi}} - \frac{\delta X}{\dot{X}}\right)\frac{2}{\chi^2}, \tag{20}$$

where the dot represents the time derivative, the term proportional to C_1 corresponds to the growing adiabatic mode, and the term proportional to C_3 the nondecaying isocurvature mode. You should notice that only χ contributes to the growing adiabatic fluctuations. Then, the amplitude of curvature perturbation Φ_A on the comoving horizon scale at $\chi = \chi_{N_c}$ is estimated by the standard one-field formula as

$$\Phi_A(N_c) \simeq \frac{f}{2\sqrt{3}\,\pi}\frac{V^{3/2}}{V'} \simeq \frac{f}{2\sqrt{3}\,\pi}\frac{|\lambda|\chi_{N_c}^3}{8}, \tag{21}$$

where $f = 3/5$ (2/3) in the matter (radiation) domination. Later we consider the case where the comoving scale corresponding to the COBE scale exits the hubble horizon during chaotic inflation. Defining N_{COBE} as the e-fold number corresponding to the COBE scale, the COBE normalization requires $\Phi_A(N_{\mathrm{COBE}}) \simeq 3\times 10^{-5}$ [25], which yields

$$|\lambda| \simeq 4.2\times 10^{-3}\chi_{N_{\mathrm{COBE}}}^{-3}. \tag{22}$$

The spectral index n_s is given by

$$n_s \simeq 1 - \frac{3}{N_{\mathrm{COBE}}}. \tag{23}$$

B. New inflation

As χ becomes of the order of unity, the dynamics becomes a little complicated because the term with $v^2 g_R\chi^2$ or the constant term v^2 may become dominant. Depending on the parameters g_R and g_I, we have a break or no break between chaotic and new inflation.[4]

In order to investigate when new inflation starts, we define the following two fields φ' and χ':

$$\begin{pmatrix} \varphi' \\ \chi' \end{pmatrix} = \begin{pmatrix} \cos\theta & -\sin\theta \\ \sin\theta & \cos\theta \end{pmatrix}\begin{pmatrix} \varphi \\ \chi \end{pmatrix}, \tag{24}$$

where θ is a constant characterized by

$$\cos 2\theta = \frac{g_R}{\sqrt{g_R^2+g_I^2}}, \quad \sin 2\theta = \frac{g_I}{\sqrt{g_R^2+g_I^2}}. \tag{25}$$

Then, the potential with $X \simeq 0$ is rewritten as

$$V(\varphi',\chi',X\simeq 0) \simeq v^2 e^{-v_2^2/2}\exp\Bigg[\Bigg(\varphi'\cos\theta + \chi'\sin\theta$$
$$+ \frac{v_2}{\sqrt{2}}\Bigg)^2\Bigg]\Bigg[\Bigg(1-\frac{|g|}{2}\varphi'^2\Bigg)^2 + \chi'^2\Bigg(|g|+\frac{|g|^2}{2}\varphi'^2$$
$$+ \frac{|g|^2}{2}\chi'^2\Bigg)\Bigg]. \tag{26}$$

We find that the global minima are given by $\varphi'^2 = 2/|g|$ and $\chi' = 0$. New inflation can take place if $|g| \gtrsim \cos^2\theta$, that is, $|\lambda| \gtrsim v\cos^2\theta$. Since the mass squared for φ', $m_{\varphi'}^2$, is given by

$$m_{\varphi'}^2 \simeq -|g| + \frac{1}{2}|g|^2\chi'^2 + \cos^2\theta\Bigg(1+|g|\chi'^2+\frac{1}{4}|g|^2\chi'^4\Bigg), \tag{27}$$

new inflation begins when $\chi' \simeq \chi'_{\mathrm{crit}}$, which is defined as

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[4]Exactly speaking, the constant term may become dominant (small hybrid inflation) before new inflation starts.

$$\chi'_{\text{crit}} \equiv \frac{2}{|g|} \sqrt{\frac{|g| - \cos^2\theta}{|g| + 2\cos^2\theta}}. \tag{28}$$

Hereafter, we set $g_I = 0 (\cos\theta = 1, \sin\theta = 0)$ for simplicity. In this case, for $\chi \simeq 0$ and $X \ll 1$, the potential is approximated as

$$V(\varphi, \chi \simeq 0, X \ll 1)$$
$$\sim v^2 [1 - (g_R - 1)\widetilde{\varphi}^2 + 2(g_R - 1)^2 \widetilde{\varphi}^2 |X|^2 + \cdots], \tag{29}$$

with $\widetilde{\varphi} = \varphi - \varphi_{\text{max}}$ and $\varphi_{\text{max}} \equiv v_2 / [\sqrt{2}(g_R - 1)]$. If $g_R \gtrsim 1$,[5] φ slow-rolls down toward the vacuum expectation value $\eta = \sqrt{2/g_R}$ and new inflation takes place.

In our model, φ stays at φ_{min} until new inflation starts. Thus, the initial value of $\widetilde{\varphi}$, $\widetilde{\varphi}_i$, for new inflation is given by

$$\widetilde{\varphi}_i = -\frac{v_2}{\sqrt{2}} \frac{g_R}{g_R - 1}. \tag{30}$$

The e-fold number N_n is estimated as

$$N_n \simeq \frac{1}{2(g_R - 1)} \ln \left| \frac{\sqrt{2}}{v_2} \frac{g_R - 1}{g_R} \right|. \tag{31}$$

In the new inflation regime, both φ and X acquire large quantum fluctuations because they are effectively massless fields. However, following the same procedure as done in the chaotic inflation regime, it is shown that only φ contributes to the adiabatic fluctuations. The amplitude of curvature perturbation Φ_A on the comoving horizon scale at $\widetilde{\varphi} = \widetilde{\varphi}_{N_n}$ is given by

$$\Phi_A(N_n) \simeq \frac{f}{2\sqrt{3}\pi} \frac{v}{2(g_R - 1)\widetilde{\varphi}_{N_n}}, \tag{32}$$

where $f = 3/5$ (2/3) in the matter (radiation) domination. The spectral index n_s of the density fluctuations is given by

$$n_s \simeq 1 - 4(g_R - 1). \tag{33}$$

Now, let us comment on the amplitude of the density fluctuations in the case with a break. Density fluctuations with comoving wave number corresponding to the horizon scale around the beginning of new inflation (if any, hybrid inflation) are induced during both chaotic and new inflation. However, following the procedure as done in Refs. [6,7], we can show that the density fluctuations produced during chaotic inflation are a little less than newly induced fluctuations

at the beginning of new inflation $[\simeq v/(2\pi\sqrt{3})]$. Furthermore, the fluctuations produced during chaotic inflation are more suppressed for smaller wavelength. Thus, we assume that the fluctuations of φ induced during chaotic inflation can be neglected when we estimate the amplitude of the density fluctuations during new inflation.

After new inflation, φ oscillates around the minimum $\varphi = \eta$ so that the universe is dominated by a coherent oscillation of the scalar field $\sigma \equiv \varphi - \eta$. Expanding the exponential factor $e^{v_2\varphi + \varphi^2}$ in e^K

$$e^{v_2\varphi + \varphi^2} = e^{\eta^2}(1 + 2\eta\sigma + \cdots), \tag{34}$$

we find that σ has gravitationally suppressed linear interactions with all scalar and spinor fields including minimal supersymmetric standard model (MSSM) particles. For example, let us consider the Yukawa superpotential $W = y_i D_i H S_i$ in MSSM, where D_i, S_i are doublet (singlet) superfields, H represents Higgs superfields, and y_i are Yukawa coupling constants. Then, the interaction Lagrangian is given by

$$\mathcal{L}_{\text{int}} \sim y_i^2 \eta\sigma D_i^2 S_i^2 + \ldots, \tag{35}$$

which leads to the decay width Γ given by

$$\Gamma \sim \Sigma_i y_i^4 \eta^2 m_\sigma^3. \tag{36}$$

Here $m_\sigma \simeq 2\sqrt{g_R} e^{\sqrt{2/g_R}} v$ is the mass of φ at the vacuum expectation value $\eta = \sqrt{2/g_R}$. Then, the reheating temperature T_R is given by

$$T_R \sim 0.1 \bar{y} \eta m_\varphi^{3/2}, \tag{37}$$

with $\bar{y} = \sqrt{\Sigma_i y_i^4}$. Taking $\bar{y} \sim 1$, the reheating temperature T_R is given by

$$T_R \sim v^{3/2} \lesssim |\lambda|^{3/2} \sim 10^{-4} \chi_{N_{\text{COBE}}}^{-9/2}. \tag{38}$$

For $N_{\text{COBE}} = \mathcal{O}(10)$, T_R is low enough to avoid the overproduction of gravitinos in a wide range of the gravitino mass [26].

IV. DISCUSSION AND CONCLUSIONS

In this paper, we propose a natural double inflation model in SUGRA. By use of the shift symmetry, first of all, chaotic inflation can take place, which has no initial condition problem. Later, new inflation takes place, which straightforwardly leads to the low reheating temperature enough to avoid the overproduction of gravitinos. In our model, the initial value of new inflation is set during chaotic inflation and is appropriately far from the local maximum of the potential for new inflation so that both inflations are responsible for the primordial fluctuations on the observable scale. In the forthcoming paper, we will investigate the density fluctuations produced during both inflations in detail and compare them with recent observations. Moreover, we will discuss the

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[5]In fact, for $g_R \gtrsim 1$, we can show that there is no local minimum between the local maximum φ_{max} and the global minima.

PBHs formation. Since the second inflation is new inflation in our model, the produced density fluctuations become of the order of unity due to the peculiar property of new inflation. PBHs may be identified with the massive compact halo objects (MACHOs) or be responsible for antiproton fluxes observed by the BESS experiments.

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