Supersymmetric topological inflation model

M. Kawasaki and Masahide Yamaguchi
Research Center for the Early Universe, University of Tokyo, Tokyo 113-0033, Japan

(Received 6 December 2001; published 13 May 2002)

We propose a topological inflation model in supergravity. In this model, the vacuum expectation value (VEV) of a scalar field takes a value much larger than the gravitational scale $M_G \simeq 2.4 \times 10^{18}$ GeV, which is large enough to cause topological inflation. On the other hand, expansions of the Kähler potential and the superpotential beyond the gravitational scale are validated by the introduction of a Nambu-Goldstone-like shift symmetry. Thus, topological inflation inevitably takes place in our model.

DOI: 10.1103/PhysRevD.65.103518 PACS number(s): 98.80.Cq, 04.65.+e, 11.27.+d

I. INTRODUCTION

Inflation is the most powerful extension to the standard big bang theory because it gives solutions to the flatness problem, the horizon problem, the origin of density fluctuations, and so on [1]. Recent observation of anisotropies of the cosmic microwave background radiation (CMB) by the Boomerang [2], the MAXIMA [3], and the DASI [4] experiments found the first acoustic peak with a spherical harmonic multipole $l \sim 200$, which implies the standard inflationary scenario. Up to now, many types of inflation models have been proposed. Among them, chaotic inflation [5] is very attractive because it does not suffer from any initial condition problem, especially, the flatness (longevity) problem. All other models which occur at low energy scales sustain this problem. That is, why does the universe live so long up to the low energy scale from the gravitational scale? However, this problem is evaded if the universe is open at the beginning. While new inflation and hybrid inflation have another severe problem, namely, the initial value problem [1,6], a fine-tuning of the initial value of the inflaton is not needed for topological inflation [7,8]. Thus, topological inflation is still attractive if the universe is open at the beginning. In fact, a possibility is pointed out that the quantum creation of an open universe can take place with appropriate continuation from the Euclidean instanton [9]. Furthermore, the spectrum of density fluctuations predicted by topological inflation becomes a tilted one, which may be testable in galaxy surveys and CMB observations.

Supersymmetry (SUSY) is one of the most powerful extensions to the standard model of particle physics because it stabilizes the electroweak scale against radiative corrections and realizes the unification of the standard gauge couplings. Therefore, it is important to consider inflation models in the context of SUSY and its local version, i.e., supergravity (SUGRA). In the context of SUGRA, topological inflation is again favorable because it straightforwardly predicts the reheating temperature low enough to avoid the overproduction of gravitinos for a wide range of the gravitino mass. This is mainly because topological inflation occurs at a low energy scale and the inflaton has only gravitationally suppressed interactions with standard particles to keep the flatness of the potential. Furthermore, it is attractive in the scheme of superstring theories[1] because superstring theories compactified on $(3+1)$-dimensional space-time have many discrete symmetries in the low-energy effective Lagrangian [12], which is very useful to cause topological inflation.

Izawa, Kawasaki, and Yanagida proposed a topological inflation model in supergravity with an *R*-invariant vacuum, which guarantees the vanishing cosmological constant at the end of inflation [13]. However, the model has two weak points, which come from the same source. First of all, the Kähler potential is expanded around the origin with expansion parameter $|\phi|/M_G$. However, the critical value ϕ_c of the vacuum expectation value (VEV) $\langle\phi\rangle$ of the inflaton to cause topological inflation is roughly the gravitational scale M_G. This can be understood from a simple discussion [7,8]. The typical radius $r \sim \langle\phi\rangle/v^2$ of the topological defect is given by equating the gradient energy $(\langle\phi\rangle/r)^2$ and the potential energy v^4. For topological inflation to occur, the typical radius r must be larger than the hubble radius given by $H^{-1} \sim M_G/v^2$, which leads to the rough condition $\langle\phi\rangle \gtrsim M_G$. In fact, Sakai *et al.* found the critical value $\phi_c \simeq 1.7 M_G$ irrespective of the coupling constant for a double well potential [14]. Later, the supergravity model [13] was investigated in detail and it was found that the critical value ϕ_c slightly depends on the slope of the potential and is as small as $0.95 M_G$ at best [15]. Thus, one wonders if the expansion of the Kähler potential is valid. Of course, the *R*-invariant vacuum given by the requirement of $\partial W/\partial\phi_i = W = 0$ for all scalar fields ϕ_i is unchanged irrespective of the form of the Kähler potential. So, the description is still valid that the potential is flat around the origin and the global minima are given by the *R*-invariant vacua. However, as the inflaton approaches the gravitational scale, the expansion of the Kähler potential becomes invalid. Thus, the potential may take a nontrivial shape beyond the gravitational scale and, even worse, a barrier may appear between the flat slope around the origin and the global minima so that inflation becomes the type of old inflation, which results in an inhomogeneous universe and hence does not work as an inflation model.

A similar problem may apply to the superpotential. In Ref. [13], the superpotential is truncated up to the quadratic term ($\lambda' \phi^2$: λ' a real constant) of the inflaton superfield

[1]The superstring inspired models of topological inflation were studied in [10,11].

for simplicity. However, higher order terms $(\mathcal{O}[(\lambda'\phi^2)^n]$ or $\mathcal{O}[(\phi^2)^n]$, $n\geqslant 2)$ may appear, which may drastically change the shape of the potential around the global minima again. Thus, the model proposed in Ref. [13] may not work.

In this paper, we propose a new model of topological inflation in supergravity, where the above problems are evaded by the introduction of a Nambu-Goldstone-like shift symmetry [16]. In the next section, we give our model of topological inflation. For successful topological inflation, we introduce symmetries and, if necessary, spurion fields whose vacuum values softly break the introduced symmetries. In Sec. III, we investigate the dynamics of topological inflation in detail and give a constraint on the parameters associated with the vacuum values of the spurion fields. In the final section, we give the summary of our results.

II. MODEL OF TOPOLOGICAL INFLATION

First of all, we introduce a Z_2 symmetry as an example of a discrete symmetry, which is necessary for producing domain walls. We assume that the inflaton supermultiplet Φ is odd and the other supermultiplets introduced later are even under the above Z_2 symmetry. As the inflaton field φ acquires its vacuum expectation value, the Z_2 symmetry is spontaneously broken.

Next, we introduce a Nambu-Goldstone-like shift symmetry to validate the expansion of the Kähler potential. We assume that the model is invariant under the following Nambu-Goldstone-like shift symmetry [16]: $\Phi\rightarrow\Phi+CM_G$, where C is a dimensionless real constant.[2] Then, the Kähler potential is a function of $\Phi-\Phi^*$, i.e., $K(\Phi,\Phi^*)=K(\Phi-\Phi^*)$, which allows the real part of the scalar components of Φ to take a value larger than the gravitational scale. However, if the shift symmetry is exact, the inflaton cannot have any potential. So, we need to break it softly for successful inflation. For this purpose, we introduce a spurion field Ξ and extend the shift symmetry to include the spurion field Ξ. We assume that the model is invariant under

$$\Phi\rightarrow\Phi+CM_G,$$

$$\Xi\rightarrow\left(\frac{\Phi}{\Phi+CM_G}\right)^2\Xi. \tag{1}$$

That is, the combination $\Xi\Phi^2$ is invariant under the shift symmetry. The vacuum value of Ξ, $\langle\Xi\rangle=u\ll 1$, softly breaks the shift symmetry. Here and hereafter, we set M_G to be unity.

Furthermore, we introduce the $U(1)_{\mathrm{R}}$ symmetry (R symmetry) because it prohibits a constant term in the superpotential, which ensures vanishing cosmological constant at the end of inflation. Since the Kähler potential is invariant only if the R charge of Φ is zero, another supermultiplet $X(x,\theta)$ with its R charge equal to 2 must be introduced. Then, the superpotential invariant under the Z_2, the shift and the $U(1)_{\mathrm{R}}$ symmetries is given by

$$W=X[\alpha_0+\alpha_1\Xi\Phi^2+\alpha_2(\Xi\Phi^2)^2+\cdots] \tag{2}$$

with $\alpha_i=\mathcal{O}(1)$. As shown later, for successful topological inflation, the coefficient of X, α_0, must be suppressed. For this purpose, we introduce another Z_2 symmetry (named Z_2') and another spurion field Π. Under this Z_2' symmetry, Φ is even and the other superfields are odd. The vacuum value of Π, $\langle\Pi\rangle\equiv v\ll 1$, softly breaks the Z_2' symmetry so that the smallness of $\alpha_0=v$ is associated with the breaking of the Z_2' symmetry. Here you should notice that the Z_2' charge of the spurion field Ξ is also odd. Then, two cases are possible. In the first case, the Z_2 and the shift symmetries are broken at the same time. In this case, we expect $\mathcal{O}(u)=\mathcal{O}(v)$. A similar case is discussed in Ref. [17] in the context of double inflation. In the second case, the Z_2' symmetry is broken first and later the shift symmetry is broken. In this case, we expect $\mathcal{O}(u)\ll\mathcal{O}(v)$. In this paper, we assume the second case. Then, inserting the vacuum values of the spurion fields, the superpotential is written as

$$W=X[v-u\Phi^2+\alpha_2(u\Phi^2)^2+\cdots]. \tag{3}$$

The higher order terms $\alpha_i(u\Phi^2)^i$:$i\geqslant 2$ are negligible because we are interested in only the field values up to the VEV of Φ, that is, $u|\Phi^2|\sim v\ll 1$. After all, we take the following superpotential:

$$W=vX(1-g\Phi^2), \tag{4}$$

with $g\equiv u/v\ll 1$. Here we have assumed that both constants u and v are real and positive for simplicity.

In the same way, the Kähler potential neglecting higher order terms is given by

$$K=-\frac{1}{2}(\Phi-\Phi^*)^2+|X|^2. \tag{5}$$

The higher order terms such as $u'(\Phi^2+\Phi^{*2})$ with $\mathcal{O}(u)=\mathcal{O}(u')$ are negligible because $u\ll 1$ and we are interested in only the field values up to the VEV of Φ, that is, $u|\Phi^2|\sim v\ll 1$. The charges of the supermultiplets are shown in Table I. Here it should be noticed that the model is natural in the sense of 't Hooft [18], that is, the symmetries are recovered if the small parameters u and v are set to be zero.

[2] A shift symmetry introduced by us is possessed by so-called no-scale supergravity theories, which appear in the low energy limit of superstring theories. However, inflation does not take place for a simple no-scale type Kähler potential. So, in order to realize inflation, we will slightly change the form of the Kähler potential keeping the shift symmetry. We hope that our inflaton is one of modulus fields in string theories after we reveal dynamics of the string, particularly, the compactification mechanism, which has not been known yet.

TABLE I. Charges of the $U(1)_{\mathrm{R}}\times Z_2\times Z_2'$ symmetries for the various supermultiplets.

	Φ	X	Ξ	Π	H_uH_d
Q_R	0	2	0	0	0
Z_2	$-$	$+$	$+$	$+$	$+$
Z_2'	$+$	$-$	$-$	$-$	$+$

III. DYNAMICS OF TOPOLOGICAL INFLATION

In this section we investigate the dynamics of topological inflation and give a constraint on the parameters associated with the breaking of the symmetries.

The superpotential and the Kähler potential given in the previous section lead to the Lagrangian density $\mathcal{L}(\Phi,X)$ for the scalar fields Φ and X,

$$\mathcal{L}(\Phi,X)=\partial_\mu\Phi\partial^\mu\Phi^*+\partial_\mu X\partial^\mu X^*-V(\Phi,X), \tag{6}$$

where the scalar potential V is given by

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

Here and hereafter, we denote the scalar components of the supermultiplets by the same symbols as the corresponding supermultiplets.

We decompose the scalar field Φ into the real component φ and the imaginary component χ,

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

Then the Lagrangian density $\mathcal{L}(\varphi,\chi,X)$ is written as

$$\mathcal{L}(\varphi,\chi,X)=\frac{1}{2}\partial_\mu\varphi\partial^\mu\varphi+\frac{1}{2}\partial_\mu\chi\partial^\mu\chi+\partial_\mu X\partial^\mu X^*-V(\varphi,\chi,X) \tag{9}$$

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

$$\begin{aligned}V(\varphi,\chi,X)=v^2\exp[\chi^2+|X|^2]\Bigg[\Bigg\{\left(1-\frac{g}{2}\varphi^2\right)^2&+\chi^2\left[g+\frac{g^2}{4}(2\varphi^2+\chi^2)\right]\Bigg\}(1-|X|^2+|X|^4)\\&+\left|X\right|^2\Bigg\{2g^2(\varphi^2+\chi^2)+4g\chi^2\left[1+\frac{g}{2}(\varphi^2+\chi^2)\right]\\&+2\chi^2\left[1-g(\varphi^2-\chi^2)+\frac{g^2}{4}(\varphi^2+\chi^2)^2\right]\Bigg\}\Bigg].\end{aligned} \tag{10}$$

Due to the exponential factor $e^{\chi^2+|X|^2}$, χ and $|X|$ are at most of the order of unity. On the other hand, φ can take a value much larger than unity without costing exponentially large potential energy. Then, φ may take a value $\varphi=\mathcal{O}(1/\sqrt{gv})=\mathcal{O}(1/\sqrt{u})$ in some regions of the universe, another value $\varphi=\mathcal{O}(1/\sqrt{g})$ in other regions, and so on. The cosmic history is quite different according to the initial value of φ. For example, in the region with the value $\varphi=\mathcal{O}(1/\sqrt{u})$, the term $(u^2/4)\varphi^4$ dominates the scalar potential, which causes chaotic inflation. Furthermore, if we fine-tune the initial value of φ, it passes through the global minimum $\langle\varphi\rangle=\pm\sqrt{2/g}$ after chaotic inflation and stops near the local maximum $\varphi=0$ so that new inflation may take place. Because of the peculiar nature of new inflation, primordial black holes may be produced. Thus, chaotic-new inflation may take place. However, the case with the similar potential has been already discussed in Ref. [19]. So we concentrate on another interesting region in this paper.

In other regions, there are some places with the values $\varphi\simeq\sqrt{2/g}$ and $\varphi\simeq-\sqrt{2/g}$. Then, topological defects (domain walls) may form if the energy density of the universe dropped enough. In fact, if the VEV of φ is larger than the gravitational scale M_G $(g\lesssim 1)$, the topological defect is unstable and the universe expands exponentially, that is, topological inflation takes place [7,8]. In this paper, we investigate the dynamics of this topological inflation in detail. Here, one should notice that in other regions chaotic inflation takes place just below the Planck scale and the universe expands enough so that our topological inflation model can be free from the flatness problem too. The observations such as spectral index of density fluctuations and gravitational waves decide which region the present universe belongs to.

The effective mass squared of χ, m_χ^2 during topological inflation is given by

$$m_\chi^2\simeq 6H^2, \tag{11}$$

where H is the hubble parameter given by $H^2\simeq v^2/3$. Thus, once topological inflation takes place, χ rapidly oscillates around the origin and the amplitude decays in proportion to $a^{-3/2}$ (a : the scale factor). Therefore, we can safely set χ to be zero at least classically.

Using $\chi\ll 1$, the scalar potential is approximated as

$$\begin{aligned}V&\simeq v^2\left[\left(1-\frac{g}{2}\varphi^2\right)^2+2g^2\varphi^2|X|^2\right]\\&\simeq v^2(1-g\varphi^2+2g^2\varphi^2|X|^2)\quad\text{for}\quad\varphi\ll 1.\end{aligned} \tag{12}$$

The effective mass squared of X, m_X^2, is given by

$$m_X^2\simeq 2g^2v^2\varphi^2\simeq 6g^2\varphi^2H^2\ll H^2. \tag{13}$$

Thus, X does not oscillate around the origin and instead slow-rolls down along the potential. In fact, during topological inflation,

$$X\sim X_i\exp\left(-\frac{g}{2}\varphi^2\right), \tag{14}$$

where $X_i<1$ is the value of X at the beginning of topological inflation. Thus, $|X|<1$ throughout topological inflation and

$|X|\ll 1$ near the end of topological inflation.[3] The last term in the potential (12) is irrelevant for the dynamics of φ because $g\ll 1$ and $|X|<1$.

Using the slow-roll approximation, the e-fold number acquired for $\varphi>\varphi_N$ is given by

$$N\simeq\int_{\varphi_f}^{\varphi_N}\frac{V}{V'}\simeq\frac{1}{2g}\ln\left(\frac{\varphi_f}{\varphi_N}\right), \tag{15}$$

where the prime represents the derivative with respect to φ and $\varphi_f\sim\sqrt{2/g}$ is the value of φ at the end of topological inflation. φ_N is represented by the e-fold N as

$$\varphi_N\sim\varphi_f e^{-2gN}\sim\sqrt{\frac{2}{g}}e^{-2gN}. \tag{16}$$

Next we evaluate the density fluctuations produced during topological inflation. In this model, there are two effectively massless fields φ and X during topological inflation. However, we can easily show that the metric perturbation in the longitudinal gauge Φ_A can be estimated as [20]

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

$$C_1=H\frac{\delta\varphi}{\dot{\varphi}}, \tag{17}$$

$$C_3=H\left(\frac{\delta\varphi}{\dot{\varphi}}-\frac{\delta X}{\dot{X}}\right)g\varphi^2,$$

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. For simplicity, we deal with X as if it is a real scalar field. You should notice that only φ contributes to the growing adiabatic fluctuations. Then, the amplitude of the curvature perturbation Φ_A on the comoving horizon scale at $\varphi=\varphi_N$ is estimated by the standard one-field formula as

$$\Phi_A\simeq\frac{f}{2\sqrt{3}\pi}\frac{V^{3/2}}{V'}\simeq\frac{f}{2\sqrt{3}\pi}\frac{v}{2g\varphi_N} \tag{18}$$

with $f=3/5$ (2/3) in the matter (radiation) domination. From the Cosmic Background Explorer (COBE) normalization $\Phi_A\simeq 3\times 10^{-5}$ at $N\simeq 60$ [21], the vacuum value v is constrained as

$$v\simeq 1.1\times 10^{-4}\sqrt{g}e^{-gN/2}|_{N=60}\simeq 3.3\times 10^{-5}-6.1\times 10^{-8} \tag{19}$$

for $0.01\leqslant g\leqslant 0.05$. The spectral index n_s is given by

$$n_s\simeq 1-4g. \tag{20}$$

Since the COBE data also show $n_s=1.0\pm 0.2$ [21], the parameter $g=u/v$ is constrained as $g\leqslant 0.05$.[4]

After topological inflation ends, the inflaton rapidly oscillates around the global minimum $\langle\varphi\rangle\equiv\pm\sqrt{2/g}$ and decays into standard particles, which reheats the universe. The decays of the inflaton into standard particles can take place if we introduce the following superpotential,[5]

$$W=v'XH_uH_d, \tag{21}$$

where $v'=\alpha\langle\Pi\rangle$ with $\alpha=\mathcal{O}(1)[\mathcal{O}(v')=\mathcal{O}(v)]$ is a constant associated with the breaking of the Z_2' symmetry and H_u,H_d are a pair of the Higgs doublets. If we set the R charge and the Z_2 and the Z_2' charges of H_uH_d to be zero and even, the above superpotential is invariant under all the introduced symmetries before inserting the vacuum value of the spurion field Π. Then the inflaton φ has the interaction with a pair of the Higgs doublets,

$$\mathcal{L}_{int}\simeq gvv'\langle\varphi\rangle\varphi H_uH_d, \tag{22}$$

which gives the decay rate

$$\Gamma\sim g^2v^2v'^2\langle\varphi\rangle^2/m_\varphi \tag{23}$$

with the mass of the φ field, $m_\varphi\simeq 2\sqrt{g}v$. Then the reheating temperature T_R is given by

$$T_R\sim 0.1gvv'\langle\varphi\rangle/\sqrt{m_\varphi}\sim 5\times 10^8\ \text{GeV}-2\times 10^6\ \text{GeV} \tag{24}$$

for $0.01\leqslant g\leqslant 0.05$. The above reheating temperature is low enough to avoid overproduction of gravitinos for a wide range of the gravitino mass [22,23].

IV. SUMMARY

In the present paper, we have proposed a topological inflation model in supergravity. In the model, where a discrete symmetry is spontaneously broken, the vacuum expectation value of the scalar field takes a value much larger than the gravitational scale so that topological inflation can take place. Generally speaking, expansions of the Kähler potential and the superpotential around the origin are invalid beyond

[3] If we take into account a higher order term $-(k_1/4)|X|^2$ ($k_1\gtrsim 1$) in the Kähler potential, the effective mass squared of X is much larger than H^2 so that the amplitude of X rapidly decays and throughout topological inflation we can safely set X to be zero at least classically.

[4] In Ref. [13], the combination of the parameters must be fine-tuned to be nearly equal to unity in order to satisfy the constraint from the spectral index. On the other hand, in our model, the parameter g has only to be small, which originates from the difference of the breaking scales of the shift and the Z_2' symmetries.

[5] The inflaton may also decay into standard particles if we consider higher order terms $u''(\Phi^2+\Phi^{*2})|\psi_i|^2$ in the Kähler potential. Here u'' is a constant associated with the breaking of the shift symmetry with $\mathcal{O}(u)=\mathcal{O}(u'')$ and ψ_i are the standard particles. Then, the reheating temperature T_R becomes 100 GeV–50 MeV for $0.01\leqslant g\leqslant 0.05$.

the gravitational scale. In our topological inflation model, the expansion of the Kähler potential is validated by the introduction of a Nambu-Goldstone-like shift symmetry. Furthermore, one can make the expansion of superpotential valid beyond the gravitational scale by introducing a Z_2' symmetry and combining it with the shift symmetry. Thus, topological inflation inevitably takes place in our model. Our topological inflation model predicts the tilted spectrum of density fluctuations, which may be detectable in future CMB observations and galaxy surveys. Furthermore, the reheating temperature is low enough to avoid overproduction of gravitinos for a wide range of the gravitino mass, especially, that predicted in the gauge mediated SUSY breaking model.

ACKNOWLEDGMENTS

M.K. is supported in part by the Grant-in-Aid, Priority Area ''Supersymmetry and Unified Theory of Elementary Particles''(No. 707). M.Y. was partially supported by the Japanese Grant-in-Aid for Scientific Research from the Ministry of Education, Culture, Sports, Science, and Technology.

[1] See, for example, A. D. Linde, *Particle Physics and Inflationary Cosmology* (Harwood, Chur, Switzerland, 1990).

[2] P. de Bernardis *et al.*, Nature (London) **404**, 955 (2000); A. E. Lange *et al.*, Phys. Rev. D **63**, 042001 (2001).

[3] S. Hanany *et al.*, Astrophys. J. Lett. **545**, L5 (2000); A. Balbi *et al.*, *ibid.* **545**, L1 (2000).

[4] N. W. Halverson *et al.*, astro-ph/0104489; C. Pryke *et al.*, astro-ph/0104490.

[5] A. D. Linde, Phys. Lett. **129B**, 177 (1983).

[6] G. Lazarides, C. Panagiotakopoulos, and N. D. Vlachos, Phys. Rev. D **54**, 1369 (1996); G. Lazarides and N. D. Vlachos, *ibid.* **56**, 4562 (1997); N. Tetradis, *ibid.* **57**, 5997 (1998).

[7] A. Linde, Phys. Lett. B **327**, 208 (1994); A. Linde and D. Linde, Phys. Rev. D **50**, 2456 (1994).

[8] A. Vilenkin, Phys. Rev. Lett. **72**, 3137 (1994).

[9] S. W. Hawking and N. Turok, Phys. Lett. B **425**, 25 (1998); A. Linde, Phys. Rev. D **58**, 083514 (1998).

[10] M. C. Bento and O. Bertolami, Phys. Lett. B **384**, 98 (1996); M. C. Bento, O. Bertolami, and N. J. Nunes, *ibid.* **427**, 261 (1998).

[11] J. Ellis, N. Kaloper, K. A. Olive, and J. Yokoyama, Phys. Rev. D **59**, 103503 (1999).

[12] For example, see J. Polchinski, *String Theory* (Cambridge University Press, Cambridge, England, 1998).

[13] K. Izawa, M. Kawasaki, and T. Yanagida, Prog. Theor. Phys. **101**, 1129 (1999).

[14] N. Sakai, H. Shinkai, T. Tachizawa, and K. Maeda, Phys. Rev. D **53**, 655 (1996); N. Sakai, *ibid.* **54**, 1548 (1996).

[15] M. Kawasaki, N. Sakai, M. Yamaguchi, and T. Yanagida, Phys. Rev. D **62**, 123507 (2000).

[16] M. Kawasaki, M. Yamaguchi, and T. Yanagida, Phys. Rev. Lett. **85**, 3572 (2000); Phys. Rev. D **63**, 103514 (2001).

[17] M. Yamaguchi, Phys. Rev. D **64**, 063502 (2001); **64**, 063503 (2001).

[18] G. 't Hooft, in *Recent Developments in Gauge Theories*, edited by G. 't Hooft *et al.* (Plenum, Cargèse, 1980).

[19] J. Yokoyama, Phys. Rev. D **58**, 083510 (1998); **59**, 107303 (1999); M. Yamaguchi and J. Yokoyama, *ibid.* **63**, 043506 (2001).

[20] D. Polarski and A. A. Starobinsky, Nucl. Phys. **B385**, 623 (1992); Phys. Rev. D **50**, 6123 (1994); A. A. Starobinsky and J. Yokoyama, in *Proceedings of the Fourth Workshop on General Relativity and Gravitation*, edited by K. Nakao *et al.* (Kyoto University Press, Kyoto, 1994), p. 381.

[21] C. L. Bennett *et al.*, Astrophys. J. Lett. **464**, L1 (1996).

[22] Y. Yu. Khlopov and A. D. Linde, Phys. Lett. **138B**, 265 (1984); J. Ellis, G. B. Gelmini, J. L. Lopez, D. V. Nanopoulos, and S. Sarker, Nucl. Phys. **B373**, 92 (1999); M. Kawasaki and T. Moroi, Prog. Theor. Phys. **93**, 879 (1995).

[23] T. Moroi, H. Murayama, and M. Yamaguchi, Phys. Lett. B **303**, 289 (1993).