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Topological curvatons

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Topological defects may play the role of the curvatons. I propose a new mechanism of generating density perturbations from cosmological defects in inflationary models. I show several examples in which defects play crucial role in generating density perturbations.

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I. INTRODUCTION

In the standard scenario of the inflationary Universe, the observed density perturbations are produced by a light inflaton field that rolls down its potential. At the end of inflation, the inflaton oscillates about the minimum of its potential and decays to reheat the Universe. Adiabatic density perturbation is generated because the scale-invariant fluctuations of the light inflaton field are different in different patches. On the other hand, one may construct a model in which the ‘‘light field’’ is not identified with the inflaton. For example, the origin of the large-scale curvature perturbation in our Universe might be induced by the late decay of a massive scalar field, the curvaton [1,2]. The curvaton is assumed to be light during a period of cosmological inflation so that it acquires scale-invariant fluctuations with the required spectrum. After inflation, the curvaton starts to oscillate in a radiation background. During this period, the energy density of the curvaton grows so that it accounts for the cosmological curvature perturbation when it decays. The curvaton can generate the curvature perturbation after inflation if its density becomes a significant fraction of the total energy density of the Universe. The curvaton paradigm has attracted a lot of attention because it was thought to have an obvious advantage. To be more precise, since the curvaton is independent of the inflaton field, there was a hope [3] that the serious fine-tunings of the inflation models could be cured by the curvaton scenario, especially in models of low-scale inflation [4]. More recently, however, it has been suggested [5] that there is a strong bound for the Hubble parameter during inflation. The bound obtained in Ref. [5] seemed to be crucial for inflationary models of a low inflation scale. Although the difficulty can be evaded [6,7] if there was an additional inflationary expansion (or a phase transition), it is still important to find another framework in which the curvaton may accommodate low inflation scale [6,7]. It should be helpful to note again that the original idea of the curvaton is;

(1) The field (curvaton) other than inflaton acquires a perturbation with an almost scale-invariant spectrum and their density becomes a significant fraction of the total.
(2) The curvaton decays into thermalized radiation so that the initial isocurvature density disappears.

In the conventional curvaton scenario, generation of the curvature perturbation is due to the oscillation of the curvaton that persists for many Hubble times, which makes it possible for the curvaton to dominate the energy density before it decays.[1] We agree with the original idea of the curvaton scenario, however in any case one cannot simply ignore the possibility of generating cosmological defects that might dominate the energy density of the Universe. Therefore, it is quite natural and important to investigate the possible scenario of ‘‘topological curvatons’’, in which the domination by the ‘‘curvaton’’ is not due to the simple oscillation of the curvaton field, but due to the evolution of the cosmological defects that decay after they have become a significant fraction. Along the lines of the above arguments, we propose a new mechanism for generating density perturbations. In the case that the cosmological defects dominate the energy density and then decay to reheat (again) the Universe, the spatial fluctuations of their ‘‘nucleation rate’’, ‘‘nucleation time’’ and ‘‘decay rate’’ may lead to the fluctuations of the reheat temperature. The important point is that the scale-invariant perturbation is due to a ‘‘light field’’ that obtains scale-invariant fluctuations during inflation. Therefore, the ‘‘nucleation rate’’, ‘‘nucleation time’’ or ‘‘decay rate’’ of the topological curvatons must depend on the ‘‘light field’’. Although our

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[1] The idea that a field other than inflaton acquires a perturbation with an almost scale-invariant spectrum can be used in different situations. For example, one may construct a model [8] in which the spatial fluctuations of the decay rate of the inflaton lead to the fluctuations of the reheat temperature. This scenario is different from the curvaton scenario in a sense that the reheating is still induced by the inflaton decay. In this case, unlike the curvaton scenario, the ‘‘light field’’ that obtains scale-invariant fluctuations during inflation never dominates the energy density of the Universe. The scale-invariant fluctuations of the ‘‘light field’’ $\delta\phi$ is transmitted to the curvature perturbation by the ϕ-dependent coupling constant and $\Gamma(\phi)$. The second paper in Ref. [8] deals with superhorizon fluctuations of the mass of a heavy particle. These ideas may be similar to our scenarios of the monopoles, however the decay mechanism is qualitatively different.

idea is mainly based on the original idea of the curvaton scenario, there exists a crucial difference from the original. In the original scenario of the curvaton, the "light field" is identified with the curvaton that dominates the energy density of the Universe. In our model, however, the "light field" itself does not play the role of the curvaton, which may lead to some relaxation of the conditions in the original model. In this case, the mechanism that the scale-invariant fluctuations of the "light field" are transmitted to the topological curvatons is important.

First, we should answer the question asking what the origin of a scale-invariant perturbation is, in the case that cosmological defects play the role of the curvaton. In the scenario of the topological curvatons, we have two options. The simplest idea is that a light field (or moduli) that obtains the scale-invariant fluctuations during inflation leads to the scale-invariant fluctuations of the "decay rate" of the defects [8]. In this case, however, the fluctuations must remain until the defects dominate and then decay to reheat the Universe. This possibility is useful in the scenarios of the string overproduction and the unstable domain walls, about which we will discuss later. The important point is that the fluctuations of the decay rate cannot survive until late epoch because it must be damped when the "light field" starts oscillation. In the scenarios of monopoles, we show that the scale-invariant fluctuations in the annihilation rate lead to the fluctuations in the initial (number) density when monopoles are frozen.

In this paper, we consider the case in which the scale-invariant fluctuations of the vacuum expectation value of a light field $\sim\delta\phi/\phi$ leads to the similar fluctuations of a coupling constant $\sim\delta g(\phi)/g(\phi)$. This idea is quite similar to the one discussed in Refs. [8,9]. Therefore, we will assume that the form of the field-dependent coupling is similar to the one that has been used in the references. Please be sure that we are not considering any peculiar situation on this point. Then, this may result in the fluctuations of "nucleation rate", "nucleation time", "decay rate" or "annihilation rate". Obviously, the nucleation always occurs much earlier than the decay. Therefore, the mechanisms that lead to fluctuations in the initial density are useful if the required energy domination occurs much later than the light-field oscillation. Based on the above arguments, we consider a scenario in which the scale-invariant fluctuations of the coupling constants lead to the similar fluctuations of the initial density of the topological curvatons. This scenario is discussed for the annihilating monopoles and the string overproduction. It is important to note that the initial density fluctuations cannot survive during the scaling epoch of the conventional strings and domain walls. However, we will show that in some cases the evolution before the scaling epoch may depend on the initial conditions. Therefore, even in the cases of the strings or the domain walls, the scale-invariant fluctuations of a "light field" may lead to the required density perturbation, if the defects decay safe before it starts scaling. We will discuss this issue in more detail for the string overproduction. On the other hand, if the topological defects starts to dominate and decay safe to "reheat again" the Universe soon after inflation, the fluctuations in their decay rate $\delta\Gamma/\Gamma$ should play significant role [8]. It should be noted that the latter idea might be conceptually different from the original idea of the curvatons.

The advantage of our scenario is obvious. In the conventional curvaton scenario, the scale of inflation is bounded below by the requirements for the fluctuations and the thermalization [5]. To avoid this difficulty, one should extend the original scenario [6] so that it includes additional phase transition that makes curvaton heavy at later epoch. To be precise, the curvaton must be light during inflation so that it acquires the required scale-invariant fluctuations, while it must become heavy in order to decay safe. In our scenario, however, the "light field" that obtains scale-invariant fluctuations during inflation does not decay to "reheat again" the Universe. The secondary reheating that is required for the curvaton scenario is induced by the topological curvatons. The important ingredient of our scenario is that **the decay of the topological curvaton is due to the characteristic feature of the symmetry breaking or the characteristic evolution of the resultant topological defects, which are independent of the light field that obtains scale-invariant fluctuations during inflation.**

II. TOPOLOGICAL CURVATONS

Although our scenario is rather different from the conventional curvatons, we still think it is helpful to start with the analogy with the original scenario. The key ingredient of the original curvaton scenario is that the scale-invariant fluctuations of the number density[2] of the curvaton field, which have been generated at the time when curvaton field ϕ starts oscillation at the Hubble parameter $H = H_{\rm osc} \simeq m_\phi$, remains until it dominates the Universe. Although the thermalization is induced by the late decay of the curvaton, the initial fluctuations still remain and lead to the fluctuation of the reheat temperature.

On the other hand, in the case that topological defects play the role of the curvatons, it is not clear how the scale-invariant perturbations are generated by a "light field" that is not identified with the curvatons. It is therefore quite useful to start with referring to the alternative idea that has been advocated in Ref. [8]. The idea is quite simple and useful. Since light fields (ϕ_i) can acquire scale-invariant fluctuations during inflation, ϕ_i-dependent coupling constants $g_j(\phi_i)$ may have the similar scale-invariant fluctuations. These massless fields start oscillation when the

[2]Of course this expression is not precise. However, for later discussions we note here that the evolution of the oscillating curvaton looks quite similar to nonrelativistic matter.

Hubble constant becomes smaller than their mass ($H_{\rm osc} \simeq m_{\phi_i}$), damping to their true vacuum expectation values. Therefore, if the decay of the inflaton starts **before** the damping, the fluctuations of the coupling constants lead to the fluctuations of the reheat temperature. In this case, $\delta T/T \sim \delta\Gamma/\Gamma \sim \delta\phi_i/\phi_i$ is expected (Numerical constants are neglected.) Obviously, this scenario is useful for our scenario of the topological curvatons, as we have commented in the previous section.

On the other hand, however, it is important to note again that the above mechanism cannot work in the scenario of the topological curvatons if the domination starts at late epoch. Therefore, in the case that the thermalization is induced by the "late" decay of the topological curvatons, the scale-invariant fluctuations of the light fields must be converted into another type of fluctuations that will not be damped before "reheating". Obviously, the scenario that uses the fluctuations in the decay rate is not suitable for the late-decay models. In the conventional scenario of the curvatons, one can see that the fluctuation of the light curvaton is converted into the fluctuation of the oscillation, whose evolution is quite similar to the number density of the nonrelativistic matter. Then the ratio of the curvaton density to the background radiation increases with time, while the scale-invariant fluctuations remain until the reheating. From the viewpoints that we have discussed above, it seems interesting to reconsider the curvaton scenario assuming that the curvaton domination is not due to the conventional oscillation of the curvaton field, but due to the evolution of the cosmological defects. Our scenario is also interesting from the viewpoint of baryogenesis, because decaying cosmological defects are sometimes considered as the sources of the baryon-number asymmetry of the Universe.[3] Therefore, the most economical solution of our scenario is that the defects are responsible for both the curvature perturbations and the baryon-number asymmetry of the Universe.

It should be noted that the scenario of the topological curvatons might be inevitable in a wide class of brane inflationary models. To be precise, it has been discussed in Ref. [15] that the naive application of the Kibble mechanism may underestimate the initial density of cosmic superstrings in brane inflationary models. In this case, strings that are formed at the end of brane inflation dominate the Universe soon after inflation, and then they decay to "reheat again" the Universe.[4] In this case, the strings that dominate the Universe after inflation are the natural candidate of the topological curvatons.

A. Fluctuations in initial number density (Monopoles)

First, we will discuss about monopoles. In this scenario, the required fluctuations are generated in the initial number density of the monopoles. To construct the successful scenario of the topological curvatons, we need to discuss about a mechanism that triggers the efficient annihilation of the monopoles [20,21]. This mechanism is often discussed in the context of the monopole problem. Although there had been many possible solutions to the monopole problem [22,23], it is commonly believed that the most attractive solution would be the inflationary Universe scenario, in which the exponential expansion ensures the dilution of the preceding abundance of the unwanted monopoles. However, other solutions are not useless because in some cases monopoles could be produced **after** inflation. For example, apart from the conventional "Grand Unified Theory (GUT) monopoles", monopole is a generic topological defect that appears whenever the required symmetry breaking occurs in the Universe. Solutions of the monopole problem in which the inflationary expansion is not assumed [20,21] are therefore still important in some extended models. In our scenario, since we are considering monopoles as the topological curvatons, the monopoles must become a significant amount of the Universe, and then decay safe by the mechanism of the efficient annihilation [20,21].

If the GUT transition is strongly first order, then the expected relic monopole abundance is

$$\frac{n_M}{s} \simeq \left[\frac{T_c}{M_p}\ln\left(\frac{M_p^4}{T_c^4}\right)\right]^3. \tag{2.1}$$

where we have taken the grand unified transition temperature to be $T_c \sim 10^{14}$ GeV. Here n_M is the monopole number density, and s is the entropy density of the radiation. Since the energy density of the radiation falls as T^4 while that of the monopoles as T^3, the monopole to radiation energy density ratio becomes $\sim 10^2$ by the time of the electroweak transition [21]. However, if there was a confining phase transition [20,21], monopoles will rapidly annihilate to release relativistic particles that will rapidly thermalize to reheat the Universe. Since we are considering a scenario of the topological curvatons, we consider a model in which the transition occurs **after** monopoles have become a significant fraction of the Universe.

1. Efficient production and the origin of the scale-invariant perturbations

Let us consider the "worst" scenario for the standard monopole problem. Here we consider the case in which the nonthermal process leads to the efficient production of the monopoles after inflation [15,24]. In the case that $n_M/s >$

[3] Inflationary models of low fundamental scale are discussed in Refs. [4,10,11]. Scenarios of baryogenesis in such models are discussed in Refs. [12–14], where defects play distinguishable roles.

[4] Brane defects such as monopoles, strings, domain walls and Q-balls has been discussed in Refs. [15–19], where it was concluded that not only strings but also other defects should appear after brane inflation. Production of other topological defects such as monopoles and domain walls is important and cannot be excluded in brane models.

10^{-10}, Preskill [19,25] found that n_M/s is reduced to about 10^{-10} by annihilations. In this case, the fluctuations in the annihilation cross section must play important role in generating number density perturbations. Recently, Bauer, Graesser and Salem(BGS) discussed in Ref. [9] that if the annihilation cross section receives scale-invariant fluctuations, the required entropy perturbation is produced at the freeze-out. In this case, if the particle comes to dominate the energy density of the Universe and subsequently decay, this leads to the required adiabatic density perturbations.[5] In the BGS scenario, the number density of the heavy particle S is given by

$$n_S \simeq \frac{T^3}{m_S M_p \langle \sigma v \rangle}, \tag{2.2}$$

where m_S and $<\sigma v>$ is the mass of the heavy particle and their annihilation rate. In our case, if monopoles are produced by an efficient process and then decoupled from thermal equilibrium after it has become nonrelativistic, the number density of the monopoles at a temperature T after freeze-out is [25]

$$n_m \simeq \frac{p-1}{A}\frac{m}{CM_p}\left(\frac{T}{m}\right)^{p-1} T^3, \tag{2.3}$$

where we have followed Ref. [25] and used the equation

$$\frac{dn_m}{dt} = -\frac{A}{m^2}\left(\frac{m}{T}\right)^p n_m^2 - 3Hn_m \tag{2.4}$$

and

$$H = \frac{T^2}{CM_p}. \tag{2.5}$$

Here A and p are constants that characterize the annihilation process. We have assumed that inflation itself cannot generate significant density perturbations, which means that there is no significant fluctuation in the temperature at the time of the freeze-out. We have also assumed that there is no fluctuation in the nucleation time t_c, when monopoles are formed. In the original BGS scenario, the decay temperature $T_{\rm dec}$ is determined by the decay rate Γ and $\langle \sigma v \rangle$. On the other hand, in our case $T_{\rm dec}$ is determined by the critical temperature T'_c when the phase transition induces the efficient annihilation of the monopoles. Therefore, in our case $T_{\rm dom}$ and $T_{\rm dec}$ in the BGS scenario becomes [25]

$$T_{\rm dom} \simeq \frac{p-1}{A}\frac{m^2}{CM_p}\left(\frac{T}{m}\right)^{p-1}, \qquad T_{\rm dec} \simeq T'_c, \tag{2.6}$$

where $T_{\rm dom}$ is the temperature when monopoles begin to dominate the energy density. Following the argument in Ref. [9], one can easily find the energy density after monopole annihilation

$$\rho = \left(\frac{\rho(T_{\rm dom})}{\rho(T'_c)}\right)^{1/3} \rho_{\rm rad}, \tag{2.7}$$

where

$$\rho(T_{\rm dom}) \simeq T_{\rm dom}^4 \propto A^{-4} m^{4(3-p)} \qquad \rho(T'_c) \simeq T'^4_c. \tag{2.8}$$

In our case, the coupling to a light field can give rise to the scale-invariant fluctuations in A or m, which result in the scale-invariant density perturbations

$$\frac{\delta\rho}{\rho} = -\frac{4}{3}\frac{\delta A}{A} + \frac{4(3-p)}{3}\frac{\delta m}{m}. \tag{2.9}$$

Once the field-dependent couplings are determined, one can easily estimate the magnitude of the non-Gaussian fluctuations. As is discussed in Ref. [9], it is not difficult to find a model in which the non-Gaussian fluctuations do not violate the observational bound.[6] This highly model-dependent issue does not match the purpose of this paper.

2. Second order phase transition and the origin of the scale-invariant perturbations

If the GUT transition is second order, the expected relic monopole abundance is

$$\frac{n_M}{s} \simeq 10^2 \left(\frac{T_c}{M_p}\right)^3 \ll 10^{-10}. \tag{2.10}$$

Therefore, the fluctuations in the annihilation cross section is not important in this case [19,25], since the monopoles are already "frozen". In this case, the fluctuations in the critical temperature play important role in generating fluctuations in $\delta n_M/n_M$.[7] Here we assume that there is no vacuum-energy domination that induces inflationary expansion or significant entropy production.[8] The critical temperature T_c for the monopole formation is

$$T_c \simeq M_v \simeq ev, \tag{2.11}$$

where M_v is the mass of the vector bosons, and v is the vacuum expectation value of the Higgs boson. Obviously, the critical temperature depends on the coupling constant e and the vacuum expectation value of the Higgs boson v. Assuming that the coefficients in the Higgs potential are homogeneous, one may ignore the fluctuations in v. In this

[5]Our model for monopoles as the topological curvatons is similar to the BGS mechanism, however there is a crucial difference in the decay mechanism. As we have emphasized in the introduction, the crucial ingredient of the topological curvatons is that their decay is induced by the phase transition that is determined by the pattern of the symmetry breaking.

[6]See Refs. [8,9] for more details.

[7]Be sure that the "fluctuations in the critical temperature" does not mean the conventional fluctuations of the background temperature. Here we are considering a peculiar situation where the background temperature distribution is almost homogeneous, but the critical temperature is different in different patches due to the scale-invariant fluctuations of the coupling constant.

[8]The typical example would be thermal inflation [26].

case, $\delta e/e$ dominates the fluctuations. Therefore, one can obtain

$$\delta T_c/T_c \sim \delta e/e. \tag{2.12}$$

In this case, $\delta e/e$ leads to the spatial fluctuations of the initial number density of the monopoles (2.10);

$$\frac{\delta n_M}{n_M} \sim 3\frac{\delta e}{e}. \tag{2.13}$$

The evolution of the nonrelativistic monopoles is quite similar to the evolution of the oscillating curvatons. In this case, the scale-invariant fluctuations that have been imprinted on $\delta n_M/n_M$ are not damped during the evolution. Here we have assumed that the function of the coupling constant $e(\phi)$ has the desired form, where ϕ is the light field that obtains scale-invariant fluctuations during inflation. As we have discussed above, it is not difficult to find a model for $e(\phi)$ in which the non-Gaussian fluctuations do not violate the observational bound. This highly model-dependent issue does not match the purpose of this paper.

3. Entropy production

If the phase transition is first order or there is an inflationary expansion, one needs to consider an additional entropy production that is induced by the reheating just after the phase transition.[9] The entropy density increases by a factor of $\sim\Delta = (T_c/T_2)^3$, where T_c is the critical temperature and T_2 is the transition temperature. In this case, the fluctuations in the coupling constants may lead to the fluctuations in the entropy increase factor $\delta\Delta/\Delta$ unless the fluctuations in T_c cancels out the fluctuations in T_2. Therefore, the two mechanisms (fluctuations in the entropy increase factor and the topological curvatons) compete in this case.[10]

4. Baryogenesis

As we have mentioned above, the most economical realization of our scenario is to produce the baryon number asymmetry from the decaying topological curvatons. Dixit and Sher argued [21] that as the Universe passes through the phase transition, the abundance of monopoles will be depleted to acceptable levels, and at the same time the process of annihilation can generate the observed baryon asymmetry. The authors have considered an alternative version of the Langacker-Pi scenario and discussed the scenario in which monopoles annihilate slightly after electroweak phase transition. When they annihilate, their energy is released as relativistic particles that rapidly thermalize. At this time, combined with our idea of topological curvatons, the fluctuations (2.13) lead to the fluctuations of the reheat temperature. Then the resulting monopole abundance is reduced to $n_M/s \simeq 10^{-46}$, which is of course negligible today. In this scenario, the monopole annihilation will significantly increase the entropy of the Universe. Therefore, unless there had been efficient mechanism of baryon number production such as Affleck-Dine mechanism [28], the observed baryon asymmetry must be generated after the monopole annihilation. The authors discussed [21] the generation of baryon number asymmetry due to the annihilation of monopoles, which becomes

$$\frac{n_B}{s} \sim \epsilon \frac{T_R}{M_X}, \tag{2.14}$$

where $T_R \sim 10^3$ GeV is the temperature of the radiation at the time of the annihilation, $M_X \sim 10^{16}$ GeV is the mass of a heavy particle, and ϵ is the average net baryon number produce in a single annihilation. Therefore, in the case that the fluctuations of the light fields lead to the fluctuations of the number density of the monopoles, our scenario of topological curvatons works in this model.

Let us consider more generic constraint that bounds the mass scale of the monopoles. If the phase transition is strongly first order, the ratio of the monopole density to the background radiation is given by

$$\frac{\rho_M}{\rho_R} \simeq 10\left(\frac{M_X}{M_p}\right)^3 M_X T^{-1}. \tag{2.15}$$

In our scenario, monopoles must dominate the energy density of the Universe before annihilation. Therefore, the bound $\rho_M/\rho_R \gg 10^{-5}$ must be fulfilled before the time of their annihilation. Assuming that the annihilation occurs before the nucleosynthesis at $T \simeq 10$ MeV, we can obtain a bound

$$M_X \gg 10^{11}\ \text{GeV}, \tag{2.16}$$

which is obtained under stringent conditions. If one considers nonthermal production of monopoles [15,24], the above constraint can be lowered.

Obviously, monopoles can play the required role of the topological curvatons. It is possible to construct models in which monopoles decay safe after they have dominated the

[9] This reheating is not due to the decay of the topological curvatons, but is induced by the vacuum energy of the Higgs field.

[10] The generating mechanism of the scale-invariant perturbation from the entropy increase factor Δ is a novel mechanism, although it may be regarded as a variation of the conventional curvatons. However, it is always difficult to calculate the transition temperature T_2 for the strongly first order phase transitions. On the other hand, we know a practical example of the entropy production. In the scenario of thermal inflation [26], Δ is simply determined by the effective mass of the inflaton field near the origin. In the case that the effective mass obtains the scale-invariant fluctuations, which is of course due to a light field that obtains scale-invariant fluctuations during inflation, the resultant fluctuations of Δ may play significant role. We will discuss this mechanism in a separate publication [27].

energy density of the Universe. The imprint on n_M that has been induced by the scale-invariant fluctuations of the coupling constant can survive during the evolution, and leads to the required density perturbation of the Universe.

B. String overproduction

The scenario of topological curvatons could be important in a wide class of brane inflationary models. It has been discussed in Ref. [15] that the naive application of the Kibble mechanism underestimates the initial density of cosmic superstrings in brane inflationary models. The authors argued that strings that are formed at the end of brane inflation should dominate the Universe soon after inflation. In this case, radiation becomes subdominant soon after inflation, then strings decay to "reheat again" the Universe. The mechanism of string overproduction is nonthermal, however their density just before the reheating may depend on coupling constants. Therefore, the strings produced by overproduction may play the role of the topological curvatons. In this scenario, there could be three chances to transmit the scale-invariant fluctuations of a "light field" to the strings.

(1) The fluctuations are produced at the first stage of the overproduction. In this case, there are practical difficulties in estimating the density fluctuations. The mechanism of the overproduction would be highly nonlinear, which means that the estimation of the density fluctuations requires numerical calculation. This possibility is beyond the scope of this paper.

(2) The first stage of the overproduction does not produce density fluctuations, however at the succeeding stage of the efficient loop production the fluctuations in the reconnection probability could lead to the density fluctuations of the small loops. We will investigate this possibility below.

(3) Even if the string density is homogeneous for both the long strings and the string loops, fluctuations can be produced during their decay into radiation. This possibility is similar to the mechanism that has been discussed in Ref. [8], despite the qualitative difference that in our case the reheating is induced by the decay of the topological defects.

In order to examine the above speculation of the topological curvatons in the scenario of the string overproduction, let us assume that the string density that is produced at the first stage of the string overproduction is homogeneous, as we have discussed above. First, we consider the second possibility that we have stated above. The number density of the closed strings that are produced succeedingly from the long strings may depend on the reconnection probability p, which means that the density fluctuations of the closed strings can be produced from $\delta p/p$, even if the string density is homogeneous just after the first stage of the string overproduction. In order to explain our idea, we consider the equations that govern the energy density of the density of long strings (ρ_L) and small loops (ρ_l) [15];

$$\begin{aligned}\dot{\rho}_L &= -2H\rho_L - f\frac{\rho_L}{L} + \mu l\frac{dn_r}{dt}\\ \dot{\rho}_l &= -3H\rho_l + f\frac{\rho_L}{L} - \mu l\frac{dn_r}{dt} - \Gamma_s\rho_l,\end{aligned} \tag{2.17}$$

where dn_r/dt and Γ_s are the rate per unit volume for loops to recombine with long strings and the string decay rate. If the strings are not charged, Γ_s is simply determined by the gravitational radiation. Here $f(p)$ is a function of the reconnection probability of the strings, and the typical length of the strings are denoted by L (long strings) and l(string loops).

In Ref. [15], it has been discussed that starting from high density of the long strings and solving the above equations, the Universe is soon dominated by the string loops but before long the string loops decay into radiation.[11]

Let us first examine the production of the closed string loops from the dense network of the long strings. The above equations suggest that in the limit of the instant production[12] the scale-invariant fluctuations in the reconnection probability p, which is induced by a light field, lead to the scale-invariant density perturbations of the closed strings. To be precise, the long strings decay into small loops due to the second term in r.h.s. ($\sim -f(p)\frac{\rho_L}{L}$), which is proportional to $f(p)$. Here $f(p)$ is the function of the reconnection probability, which will go like $f \sim p^{1/2}$ [15]. Then, string loops dissipate their energy through radiation. Therefore, the string loops that have been dominated the density of the Universe are soon converted into radiation that starts to dominate the Universe after the reheating. Obviously, the reconnection probability of the cosmic strings does depend on the effective coupling constants. For fundamental strings, the reconnection probability is of order g_s^2, then $f(p)$ becomes $f(p) \propto g_s$. In other cases, since light fields or moduli of small compactified dimensions determine the coupling constants in the effective action, the scale-invariant fluctuations in the light fields may lead to the scale-invariant density perturbations of the string loops, $\delta\rho_l/\rho_l \propto \delta f/f$.[13] Therefore, in the case that Γ_s is a homogeneous constant in the Universe, the required density perturbations can be produced from $\delta\rho_l/\rho_l$.

The succeeding process of the reheating from string loops is quite similar to the conventional reheating.

[11]See Fig. 12 in Ref. [15].

[12]In this case, one may assume that $\dot{f} = \dot{L} = 0$ during this period.

[13]Even in the standard inflationary scenario in which inflaton is responsible for the curvature perturbations, the perturbations that have been produced by the inflaton reheating could be modified due to the succeeding stage of the string domination. In both cases the overproduction of cosmological defects could be used to suggest the existence of tachyonic potential and brane inflation, once the precise form of the effective action is determined.

Therefore, the scenario proposed in Ref. [8] is successful in this case. To be precise, the required perturbations can be produced by the fluctuations in the decay rate $\delta\Gamma_s/\Gamma_s$, if Γ_s is a desired function of a light field [8].

The above two mechanisms ($\delta\rho_l/\rho_l$ and $\delta\Gamma/\Gamma$) work independently.

III. CONCLUSIONS AND DISCUSSIONS

In this paper, we have shown that the new mechanism of generating density perturbations can work in inflationary models. We have considered monopoles and strings, and discussed about successful scenarios. In the case of the monopoles, the spatial fluctuations of the initial number density of the monopoles lead to the required fluctuations of the Universe. Strings can play similar role if a dense string network is produced by string overproduction. Other defects may play the role of the topological curvatons, however they are not so simple and require further study. We will comment on the possibility and the difficulty of the domain walls.

A. Domain walls

Domain walls are produced when discrete symmetry is spontaneously broken. In the simplest case of Z_2 walls, it is known that the evolution of the system of the domain walls becomes scale-invariant due to the efficient annihilation and reconnection processes which suggests only one large wall exists per horizon. In this case, the energy density of the walls is given by

$$\rho_w \sim \sigma t^{-1}, \tag{3.1}$$

where σ is the tension of the domain wall. Although the Universe could become wall-dominated at $t_* \sim (G\sigma)^{-1}$, the wall-domination occurs (generally) after scale-invariant evolution starts. This suggests that late-time behavior of the system that could be relevant for our scenario is not sensitive to the choice of the initial state.

In a more complex system of Z_n walls, the domain walls may become frustrated and their density decreases as [29]

$$\rho_w \sim \sigma a^{-1}, \tag{3.2}$$

where a is the scale factor. The density of such domain walls will quickly dominate the Universe before the system starts the scale-invariant evolution. In this case, the initial density fluctuations of the domain walls could remain at the time of wall-domination and also affect the subsequent thermalization. However, unlike the scenario of the string overproduction, it is quite difficult to make a practical calculation during this period.

On the other hand, one may expect that a coupling constant in the potential of the domain walls depends on a light field ϕ that obtains a scale-invariant fluctuation during inflation. Then, in the case that the function $\sigma(\phi)$ has a desired form, $\delta\phi/\phi$ may lead to the density fluctuations of the walls

$$\frac{\delta\rho_w}{\rho_w} \sim \frac{\delta\sigma}{\sigma} \sim \frac{\delta\phi}{\phi}. \tag{3.3}$$

This is a trivial example of the domain-wall curvatons.

Of course, it is possible to produce the required perturbations from the scale-invariant fluctuations of the decay rate of the walls $\delta\Gamma_w/\Gamma_w$, as we have discussed for the strings [8].

The most important ingredient for the domain walls is the mechanism that triggers the decay of the cosmological defects [30]. If the discrete symmetry is explicitly broken to a small extent, the vacua are not degenerated and have slightly different energies. The gap of the energies of the vacua leads to the gap of the pressure, which finally leads to the instant collapse of the network of the domain walls. In general, the magnitude of the explicit breaking term is determined by an intentional higher-dimensional term, except for some peculiar cases [31]. Therefore, to obtain a successful cosmological scenario, a kind of fine-tuning is required for the domain walls.

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[1] S. Mollerach, Phys. Rev. D **42**, 313 (1990); A. D. Linde and V. Mukhanov, Phys. Rev. D **56**, R535 (1997); K. Enqvist and M. S. Sloth, Nucl. Phys. **B626**, 395 (2002).

[2] T. Moroi and T. Takahashi, Phys. Lett. B **522**, 215 (2001); D. H. Lyth and D. Wands, Phys. Lett. B **524**, 5 (2002).

[3] K. Dimopoulos and D. H. Lyth, Phys. Rev. D **69**, 123509 (2004); M. Giovannini, Phys. Rev. D **67**, 123512 (2003).

[4] P. Kanti and K. A. Olive, Phys. Lett. B **464**, 192 (1999); N. Arkani-Hamed, S. Dimopoulos, N. Kaloper, and J. March-Russell, Nucl. Phys. **B567**, 189 (2000); R. N. Mohapatra, A. Perez-Lorenzana, and C. A. de S. Pires, Phys. Rev. D **62**, 105030 (2000); A. Mazumdar, S. Panda, and A. Perez-Lorenzana, Nucl. Phys. **B614**, 101 (2001); A. Mazumdar, Phys. Lett. B **469**, 55 (1999); A. M. Green and A. Mazumdar, Phys. Rev. D **65**, 105022 (2002); D. H. Lyth, Phys. Lett. B **448**, 191 (1999); T. Matsuda, Phys. Rev. D **66**, 107301 (2002); T. Matsuda, Phys. Lett. B **423**, 35 (1998).

[5] D. H. Lyth, Phys. Lett. B **579**, 239 (2004).
[6] T. Matsuda, Classical Quantum Gravity **21**, L (2004).
[7] M. Postma, J. Cosmol. Astropart. Phys. 05 (2004) 002; K. Dimopoulos, D. H. Lyth, and Y. Rodriguez, J. High Energy Phys. 02 (2005) 055; Y. Rodriguez, Mod. Phys. Lett. A **20**, 2057 (2005).
[8] G. Dvali, A. Gruzinov, and M. Zaldarriaga, Phys. Rev. D **69**, 023505 (2004); G. Dvali, A. Gruzinov, and M. Zaldarriaga, Phys. Rev. D **69**, 083505 (2004).
[9] C. W. Bauer, M. L. Graesser, and M. P. Salem, Phys. Rev. D **72**, 023512 (2005).
[10] T. Matsuda, Phys. Rev. D **67**, 083519 (2003); T. Matsuda, J. Cosmol. Astropart. Phys. 11 (2003) 003; T. Matsuda, Phys. Rev. D **68**, 047702 (2003).
[11] T. Matsuda, J. Cosmol. Astropart. Phys. 06 (2003) 007; T. Matsuda, Phys. Rev. D **68**, 127302 (2003).
[12] G. R. Dvali and G. Gabadadze, Phys. Lett. B **460**, 47 (1999); A. Masiero, M. Peloso, L. Sorbo, and R. Tabbash, Phys. Rev. D **62**, 063515 (2000); A. Pilaftsis, Phys. Rev. D **60**, 105023 (1999); R. Allahverdi, K. Enqvist, A. Mazumdar, and A. Perez-Lorenzana, Nucl. Phys. **B618**, 277 (2001); S. Davidson, M. Losada, and A. Riotto, Phys. Rev. Lett. **84**, 4284 (2000).
[13] T. Matsuda, Phys. Rev. D **66**, 023508 (2002); T. Matsuda, Phys. Rev. D **66**, 047301 (2002); T. Matsuda, Phys. Rev. D **65**, 107302 (2002); T. Matsuda, J. Phys. G **27**, L103 (2001).
[14] T. Matsuda, Phys. Rev. D **67**, 127302 (2003); T. Matsuda, Phys. Rev. D **65**, 103501 (2002); T. Matsuda, Phys. Rev. D **65**, 103502 (2002); T. Matsuda, Phys. Rev. D **64**, 083512 (2001).
[15] N. Barnaby, A. Berndsen, J. M. Cline, and H. Stoica, J. High Energy Phys. 06 (2005) 075.
[16] T. Matsuda, J. Cosmol. Astropart. Phys. 10 (2004) 014.
[17] T. Matsuda, J. High Energy Phys. 11 (2004) 039; T. Matsuda, J. High Energy Phys. 10 (2004) 042.
[18] T. Matsuda, J. High Energy Phys. 09 (2003) 064.
[19] T. Matsuda, Phys. Rev. D **70**, 023502 (2004); T. Matsuda, J. High Energy Phys. 05 (2005) 015; T. Matsuda, hep-ph/0509061; T. Matsuda, hep-ph/0509062; T. Matsuda, hep-ph/0509064.
[20] P. Langacker and So-Young Pi, Phys. Rev. Lett. **45**, 1 (1980).
[21] V. V. Dixit and M. Sher, Phys. Rev. Lett. **68**, 560 (1992).
[22] E. W. Kolb and M. S. Turner, *The Early Universe* (Addison-Wesley, Reading, MA, 1990).
[23] D. Stojkovic and K. Freese, Phys. Lett. B **606**, 251 (2005).
[24] E. J. Copeland, S. Pascoli, and A. Rajantie, Phys. Rev. D **65**, 103517 (2002); A. D. Dolgov and A. D. Linde, Phys. Lett. **116B**, 329 (1982); L. F. Abbott, E. Farhi, and M. B. Wise, Phys. Lett. **117B**, 29 (1982); J. H. Traschen and R. H. Brandenberger, Phys. Rev. D **42**, 2491 (1990); L. Kofman, A. D. Linde, and A. A. Starobinsky, Phys. Rev. Lett. **73**, 3195 (1994); L. Kofman, A. D. Linde, and A. A. Starobinsky, Phys. Rev. Lett. **76**, 1011 (1996); I. I. Tkachev, Phys. Lett. B **376**, 35 (1996); I. Tkachev, S. Khlebnikov, L. Kofman, and A. D. Linde, Phys. Lett. B **440**, 262 (1998); L. M. Krauss and M. Trodden, Phys. Rev. Lett. **83**, 1502 (1999); J. Garcia-Bellido, D. Y. Grigoriev, A. Kusenko, and M. E. Shaposhnikov, Phys. Rev. D **60**, 123504 (1999); A. Rajantie, P. M. Saffin, and E. J. Copeland, Phys. Rev. D **63**, 123512 (2001); G. N. Felder, L. Kofman, Andrei D. Linde, and I. Tkachev, J. High Energy Phys. 08 (2000) 010.
[25] J. P. Preskill, Phys. Rev. Lett. **43**, 1365 (1979).
[26] D. H. Lyth and E. D. Stewart, Phys. Rev. D **53**, 1784 (1996).
[27] T. Matsuda, to be published.
[28] I. Affleck and M. Dine, Nucl. Phys. **B249**, 361 (1985).
[29] M. Kawasaki and F. Takahashi, Phys. Lett. B **618**, 1 (2005).
[30] A. Vilenkin and E. P. S. Shellard, *Cosmic Strings and other Cosmological Defects* (Cambridge University Press, Cambridge, 2000).
[31] T. Matsuda, Phys. Lett. B **486**, 300 (2000); T. Matsuda, Phys. Lett. B **436**, 264 (1998); T. Matsuda, Phys. Lett. B **423**, 40 (1998).