NO curvatons or hybrid quintessential inflation

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Abstract. We consider a curvaton scenario in which the late-time domination and the generation of the curvature perturbation is achieved by a non-oscillatory (NO) curvaton potential. Instead of considering the conventional curvaton oscillation, we consider ‘weak trapping’ after preheating, which modifies the evolution of the curvaton density after preheating. The primordial isocurvature perturbation related to the curvaton is first converted into the fluctuation of the number density of the preheat field through inhomogeneous preheating. Then the evolution of the curvatons and the preheat field is controlled by the preheat-field number density. The density of these fields decreases slightly slower than the standard matter density, which suggests that these fields will grow with time. Finally, the preheat field decays to reheat the Universe leaving behind the curvature perturbation. In our scenario the task of the standard curvaton is not executed solely by the curvaton itself but is partially shared with the preheat field. NO curvatons can be considered as the hybrid version of the quintessential inflationary model.

Keywords: cosmological perturbation theory, cosmological phase transitions, inflation, physics of the early universe
Recent observation of the anisotropy of the cosmic microwave background radiation (CMBR) confirms that the present structure of the Universe should have originated from the primordial density perturbation. The traditional explanation for this primordial density perturbation is that during inflationary expansion the perturbation related to the light field (inflaton) generated the curvature perturbation during the inflation. The traditional inflationary scenario, in which the primordial curvature perturbation is generated solely by the inflaton field, is very simple but inevitably puts extra conditions on the inflationary mechanism and the form of the inflaton potential.

Recently, alternatives to the traditional inflationary scenario have been suggested by many authors [1]–[8]. The common thread of all these alternatives is that the isocurvature perturbation related to a light field, other than the inflaton, is converted into the curvature perturbation after (or at the end of) inflation. Since the seed perturbation is generated during primordial inflation, the spatial length scale of the resultant curvature perturbation is very large compared with the correlation length at the time when the curvature perturbation is generated. There are many scenarios related to such conversion mechanisms that characterize the concepts behind these alternative models.

Among these alternatives, we will consider a mixed scenario of curvatons [1,2] and inhomogeneous preheating [7,8]. In the standard curvaton scenario the curvaton energy density is supposed to be negligible just after inflation and the spacetime is unperturbed at that time. The curvaton energy density becomes significant during some era of radiation domination or the era when the kinetic energy of the inflaton dominates the Universe [12]. In the case where radiation dominates the Universe, the radiation energy density will evolve as $\rho_{\text{tot}} \propto a^{-4}$ while the curvaton energy density is (1) almost constant before its oscillation ($\rho_{\text{curv}} \propto a^0$) and (2) evolves as $\rho_{\text{curv}} \propto a^{-3}$ after it starts to oscillate. Therefore, the ratio of the curvaton energy density to the background radiation will grow with time, and finally the curvaton will become the dominant component of the Universe. The curvaton could generate the curvature perturbation if the curvaton oscillation lasts for a long time before the decay. The same thing will happen if the inflaton does not decay but continues to oscillate with a quartic potential. During this oscillatory period the

\[ \text{There is a possible scenario ('heavy curvatons') in which curvatons might become heavy and start to dominate the Universe just after inflation [2].} \]
inflaton energy density will evolve as $\rho_{\text{osc}} \propto a^{-4}$ [10]. Moreover, if the inflaton potential is non-oscillatory, the kinetic energy of the inflaton will evolve as $a^{-6}$ [11]. In this case the domination by the curvaton will occur much earlier and more easily than in the standard curvaton scenario [12].

Perhaps the mixed scenario that we will describe in this paper will be categorized as the curvaton scenario in the sense that the field accompanied by the isocurvature perturbation is (1) initially negligible but after a non-oscillatory period it (and the accompanying preheat field) starts to dominate the energy density of the Universe and (2) the preheat field decays to reheat the Universe generating the curvature perturbation. Note that unlike the standard curvaton the light field (non-oscillatory (NO) curvaton) connected to the primordial isocurvature perturbation neither oscillates nor decays to reheat the Universe. The isocurvature perturbation of the NO curvaton is first converted into the fluctuation of the preheat-field number density. In this sense the NO curvaton seeds the generation of the curvature perturbation that is induced by the accompanying preheat field$^2$. The curvaton has the NO potential but is weakly trapped near the enhanced symmetry point (ESP) because of the effective confining potential that is induced by the preheat field. The trapping mechanism makes the evolution of the curvaton (and the preheat field) energy density completely different from the conventional scenario. Their evolution is determined by the force-balance equation that is written by the NO potential and the number density of the preheat field. We will show that their energy densities evolve as $a^{-3+\epsilon} \ (\epsilon = 3/(n + 1))$ with the constant ratio $\rho_{\text{preheat}}/\rho_{\text{curv}} = n$. The mass of the preheat field grows during the evolution until it finally reaches the critical value where the preheat field decays to reheat the Universe. After the decay of the preheat field, the NO curvaton rapidly roll down the NO potential. In this sense the NO curvaton could be seen as the hybrid version of the quintessential inflation model [11]. However, we do not stick to the idea that the NO curvaton also plays the role of the quintessence, even though the idea seems rather fascinating. The conditions for quintessential inflation are not necessary for the NO curvaton, and vice versa. In this paper we will focus our attention on the NO curvaton.

2. NO curvaton

First we briefly review the basic idea of inhomogeneous preheating. The preheating mechanism that is presented here is basically given by Kofman et al [10]. For simplicity, we consider real scalar fields $\phi_I$, $\sigma$ and $\chi$. Here $\phi_I$ is the inflaton that has a chaotic-type potential. $\sigma$ is the light field accompanied by the primordial perturbation and biases the production of the preheat field $\chi$. $\sigma$ is the ‘curvaton’ in our scenario, although it does not play all the roles played by the standard curvaton. The important roles that characterize the curvaton are (1) isocurvature perturbation, (2) late-time domination and (3) reheating by the decay. In our scenario of the NO curvaton, the first characteristic is related to the ‘curvaton’ $\sigma$, but the second and the third are connected to the preheat field $\chi$. Therefore, although the light field $\sigma$ is called a ‘curvaton’ in this paper, it does not play all the roles that have been played by conventional curvaton as described in

$^2$ A similar scenario has been discussed in [3], where the nucleation rate of a decaying cosmological defect is supposed to be biased by a light field. As in the NO curvaton, the light field that biases the nucleation of the cosmological defects never dominates the energy density of the Universe.
earlier work. We hope this peculiar feature of the NO curvaton does not become a further confusion to our readers. With these scalar fields we will consider a model that contains the inflation sector

\[
\mathcal{L}_I = \frac{1}{2} \partial_\mu \phi_I \partial^\mu \phi_I + \frac{1}{2} \partial_\mu \sigma \partial^\mu \sigma + \frac{1}{2} \partial_\mu \chi \partial^\mu \chi - \frac{g^2}{2} (\phi_I^2 + \sigma^2) \chi^2 - V(\phi_I) - V(\sigma),
\]

(2.1)

where the cosmological constant is disregarded. Since we are considering chaotic-type inflation, the inflaton potential for large $\phi_I$ in this case is given by the generic form

\[
V(\phi_I) = \frac{\lambda_I |\phi_I|^n}{M_I^{n-4}},
\]

(2.2)

while for the curvaton potential we consider the potential

\[
V(\sigma) = \frac{M^{n+4}}{\sigma^n} + \frac{1}{2} m^2 \sigma^2,
\]

(2.3)

which has a run-away behavior if $m = 0$. Here we assume that $m$ is very small (or $m = 0$) so that there is no sensible $\sigma$-oscillation that modifies our scenario. A similar potential could be seen for the repulsive force between branes [4,8].

We will consider the generation of the preheat field $\chi$ when the inflaton $\phi_I$ approaches the origin just after chaotic inflation. Since the mass of the preheat field ($m_\chi^2 = g^2 (\phi_I^2 + \sigma^2)$) depends both on $\phi_I$ and $\sigma$, the generation of the preheat field will be biased if there is isocurvature perturbation related to $\sigma$. Let us describe the basic idea of the preheating. Just after the inflation, the mass of the preheat field is very large but will become very small when the inflaton approaches the origin. Therefore, the adiabatic condition is violated if

\[
|\dot{m}_\chi| / m_\chi^2 \sim g |\dot{\phi}_I| / g^2 (\phi_I^2 + \sigma^2) > 1
\]

(2.4)

and nonadiabatic particle production occurs. If the inflaton is the chaotic type, the nonadiabatic region is very narrow compared with the initial amplitude, and thus efficient particle production occurs within a very short period of time when the inflaton passes the ESP. In this case one can neglect the expansion of the Universe during particle production. The number density of the preheat field at the first scattering is calculated in [10] and given by

\[
n_\chi = \frac{(gv)^{3/2}}{(2\pi)^3} \exp \left( - \frac{\pi m_\chi^2}{gv} \right).
\]

(2.5)

Here $v$ is the absolute value of the inflaton velocity at the ESP. We would like to suppose that the coupling constant is large ($g \sim 1$). If the decay of the preheat field is negligible during this period, the preheat field induces the confining effective potential for both $\phi_I$ and $\sigma$, which leads to the trapping of these fields at (or near) the ESP [9]. However, as we are considering the unusual NO potential for the curvaton $\sigma$, the trapping of the NO curvaton cannot be explained by the standard trapping scenario. We will discuss this exceptional trapping mechanism later in this section.

Before discussing the trapping of the NO curvaton, we will describe the origin of the fluctuation that appears for the number density of the preheat field. According to equation (2.5), the production is efficient if $m_\chi^2 < v$. Assuming that the potential energy during inflation is efficiently converted into the kinetic energy of the inflaton, the velocity
of the inflaton becomes \( v \simeq H_1 M_p \). Then the condition for the efficient production puts an upper bound on the initial value of \( \sigma \):

\[
\sigma < \sqrt{H_1 M_p}.
\]  

(2.6)

If the value of \( M \) is very small compared with the inflationary scale, the \( \sigma \)-potential is flat except for the very narrow region near the origin. Therefore, we will assume that the curvaton \( \sigma \) experiences a flat potential during inflation. The flatness condition for the curvaton is \( m_\sigma < H_1 \) during inflation, which becomes

\[
M \left( \frac{M}{H_1} \right)^{2/(n+2)} < \sigma.
\]  

(2.7)

Remember that we are considering the case in which the isocurvature perturbation related to \( \sigma \) plays an important role in biasing the production of the preheat field. Therefore, we will consider the initial value of \( \sigma \)

\[
M \left( \frac{M}{H_1} \right)^{2/(n+2)} < \sigma < \sqrt{H_1 M_p}
\]  

(2.8)

so that \( \sigma \) appears with the primordial isocurvature perturbation and also does not disturb the generation of the preheat field \( \chi \). Considering the biasing effect related to the primordial perturbations \( \delta \sigma \), we obtained the fluctuation of the number density [8]

\[
\frac{\delta n_\chi}{n_\chi} \simeq \frac{2\pi g \sigma \delta \sigma}{v}.
\]  

(2.9)

After the first scattering, the scatterings occur successively where the number densities are given by

\[
n_{j+1}^{k} = b_{j}^{k} n_{k}^{j}.
\]  

(2.10)

Here the coefficient \( b \) is given by

\[
b_{j}^{k} = 1 + 2e^{-\pi \mu^2} - 2\sin \theta^j e^{-\pi \mu^2/2} \sqrt{1 + e^{-\pi \mu^2}},
\]  

(2.11)

where \( \mu^2 = (k^2 + m_\chi^2)/gv \). Here \( \theta^j \) is a relative phase. It is possible to consider \( \delta b_{k}^{j} \), but they are not important in the present scenario because the light field \( \sigma \) is trapped near the ESP by the strong confining potential just after the first scattering. Then the small \( \sigma \) makes the fluctuations appearing in \( \delta b \) completely negligible compared with the one generated at the first scattering. After successive scatterings the kinetic energy of the inflaton is transferred into the excitations of the preheat field, which leads to the number density [9]

\[
n_\chi \sim v^{3/2} g^{-1/2}
\]  

(2.12)

with the perturbation given by equation (2.9). According to observations, the spectrum of the curvature perturbation on a cosmological scale is \( P_\zeta^{1/2} \simeq 2 \times 10^{-5} \). As in the standard curvaton scenario, the preheat field decays leaving behind the curvature perturbation

\[
\zeta \simeq \alpha \sigma \frac{\delta \rho_\chi}{\rho_\chi}.
\]  

(2.13)
where $r$ is the ratio between $\rho_\chi$ and the total energy density of the Universe. The value of $\alpha$ is $\alpha = 1/3$ in the standard scenario, while in our scenario it is slightly shifted.

Now we will consider $\sigma$-trapping. The idea of ‘trapping’ that is used in this paper is rather different from the original idea that has been used to discuss brane trapping [9] 3. One might think that the NO potential is an effective potential so that it must be regularized near the origin by an explicit cut-off. If so, $\sigma$ will initially be trapped at the ESP and will then start to roll down the NO potential when $n_\chi$ is diluted below some critical value. This speculation is conceivable, but even if there is an explicit cut-off we can use the effective potential as long as we are considering the late-time evolution of the curvaton. Therefore, for the late-time evolution the effective minimum of the $\sigma$-potential is obtained from the force-balance equation

$$gn_\chi(t) - \frac{nM^{n+4}}{\sigma^{n+1}} = 0,$$  

which leads to the expectation value of $\sigma(t)$:

$$\sigma(t) = M \left( \frac{nM^3}{gn_\chi(t)} \right)^{1/(n+1)}.$$  

The above equation suggests that the mass of the preheat field does depend on the cosmological time through the evolution of $n_\chi$. From these equations we can calculate the ratio of $\rho_\chi$ to $V(\sigma)$:

$$\frac{\rho_\chi}{V(\sigma)} = n.$$  

It is now clear that the ratio is given by the constant $n$. Note that the energy density of the preheat field is always larger than the curvaton potential energy. Since the number density $n_\chi$ evolves as $n_\chi \propto a^{-3}$, the energy density of the preheat field $\rho_\chi$ will evolve as

$$\rho_\chi \propto a^{-3(1-1/(n+1))}$$  

keeping the ratio $\rho_\chi/V(\sigma) = n$. Therefore, the evolution of the preheat field density is slightly slower than the standard matter $\rho_m \propto a^{-3}$. Note that the mass of the preheat field $m_\chi$ grows as

$$m_\chi \simeq g\sigma \propto a^{3/(n+1)}.$$  

The decay of the preheat field is determined by its couplings to light fermions $\sim g_\chi_\chi\bar{\psi}\psi$, where the preheat field is supposed to decay at $H \simeq \Gamma_\chi$. If the decay rate is given by

$$\Gamma_\chi \simeq m_\chi^3/M_p^2,$$  

the preheat field will decay when $\sigma$ reaches the value

$$\sigma_d \simeq \left( \frac{M_pM^{(n+4)/2}}{g^3} \right)^{2/(n+6)}.$$  

3 It will be very interesting to compare our scenario with the scenario of ‘trapped quintessential inflation’ [13]. Although the mechanism for generating the curvature perturbation is completely different, our present scenario could be seen as a hybrid version of the model discussed in [13]. Of course we can accommodate the ‘trapped quintessential inflation’ itself by simply introducing an extra light field that induces inhomogeneous preheating at the trapping. This is not the hybrid version of the original scenario. In this case the extra light field may have a conventional (non-NO) flat potential.
Here we have assumed that the Universe is already dominated by the preheat field when it decays (i.e. the Hubble constant is given by $H^2 \approx \rho_\chi/M_p^2 \approx V(\sigma)M_p^2$). Suppose that the inflaton energy density is initially $\rho_1 \approx v^2$ and then evolves as $\rho_1 \propto a^{-4}$, and also that the preheat field is initially $n_\chi^i \approx v^{3/2}$ and then evolves as $n_\chi \propto a^{-3}$. Here the condition for the curvaton domination is $\rho_d \lesssim g\sigma d n_\chi^d$. The preheat-field number density at the decay is determined by the value of $\sigma_d$ and the relation (2.15), and becomes $n_\chi^d \approx M^{n+4}/\sigma_d^{n+1}$. Therefore, the inflaton density at that time is

$$\rho_1^d \approx v^2 \left( \frac{n_\chi^d}{n_\chi^i} \right)^{4/3} \approx M^4 \left( \frac{M^{n+1}}{\sigma_d^{n+1}} \right)^{4/3}.$$  (2.21)

For example, considering $n = 4$ and $M = 10^{-16} M_p$ we obtained $\sigma_d \approx 10^3 M$, which suggests that $\rho_1^d / \rho_\chi^d \approx 10^{-20}$. Obviously, it is very easy to satisfy the required condition for curvaton (preheat-field) domination. If there is a cut-off for the effective potential $V(\sigma)$, the curvaton $\sigma$ will be trapped at the origin with (perhaps) a constant value of the potential until the preheat field is diluted below the critical value. Then the evolution is determined by the number density of the preheat field, which does not induce any difference from the above calculation.

After the decay of the preheat field, the curvaton $\sigma$ does not experience the confining potential that had been induced by the preheat field. The effective mass of the curvaton at that time is $m_\sigma \approx M^{n+4}/\sigma_d^{n+2}$. Considering again the case with $n = 4$ and $M = 10^{-16} M_p$, we obtain the ratio between the curvaton mass and the Hubble constant just after the decay; $m_\sigma^2 / H^2 \approx 10^{20}$. This result suggests that the curvaton starts to roll down the potential and rapidly approaches the tracker solution. It is always possible to introduce fine-tuning of the cosmological constant so that our model fits the cosmological observations. One may also consider various kinds of quintessential potential for the NO curvatons, which may or may not explain the dark energy of the Universe depending on the later quintessential scenario. In any case, the fine-tuning of the cosmological constant will remain a serious problem until we can discover the secret related to the radiative corrections, the symmetry breaking and the origin of the cosmological constant.

3. Conclusions and discussions

We considered the curvaton scenario in which the curvaton has a non-oscillatory potential. The primordial isocurvature perturbation of the NO curvaton is converted into the number density of the preheat field when the preheat field is generated through preheating. The evolution of the curvaton and the preheat field is controlled by the trapping mechanism until the preheat field decays to reheat the Universe. Since the densities of the NO curvatons (and the preheat field) evolve slightly slower than the standard matter, the ratio between these fields and the total energy density grows with time. At the same time, the mass of the preheat field grows as the NO curvaton rolls down the potential. A peculiar feature of the scenario is that the late-time domination and the reheating is induced by the preheat field, not by the NO curvaton itself. Our basic idea is (1) the nucleation rate of the cosmological relic density is fluctuated by the primordial isocurvature perturbations,
(2) the rate of the massive relic density to the total energy density of the Universe grows during the cosmological evolution and (3) eventually the massive relic decays to reheat the Universe leaving behind the curvature perturbation. In our present model the light field (‘NO curvatons’) accompanied by the primordial isocurvature perturbation is responsible for neither the domination nor the reheating. This peculiar feature of the NO curvaton is of course very different from the conventional curvaton scenario. In our scenario the isocurvature perturbation of the curvaton is used solely to bias the nucleation rate of the preheat field. The preheat field (not the curvaton itself) is responsible for the late-time domination and the reheating. The situation is very similar to the ‘alternatives’ to the traditional inflationary scenario, in which the role of the traditional inflaton is shared with the additional light field, making the constraints on the inflaton potential looser than the traditional ones. The NO curvaton shares the role of the traditional curvatons with the preheat field, making the constraints on the curvaton potential very different from the traditional curvaton scenario.

Moreover, NO curvatons could be considered as the hybrid version of the quintessential inflation model. In this hybrid model one does not have to look for a potential that has the property of chaotic-type inflaton in the left-hand side while having run-away behavior in the right-hand side. Finding such a peculiar potential has been a serious obstacle for the single-field scenario. In the hybrid model the conventional inflaton potential and the quintessential potential can be separated and are connected at the ESP.

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