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Curvaton reheating in tachyonic inflationary models

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Abstract

The curvaton reheating in a tachyonic inflationary universe model with an exponential potential is studied. We have found that the energy density in the kinetic epoch, has a complicated dependencies of the scale factor. For different scenarios, the temperature of reheating is computed. These temperature result to be analogous to those obtained in the standard case of the curvaton scenario.

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Inflationary universe models [1] have solved many problems of the Standard Hot Big-Bang scenario, for example, the flatness, the horizon, and the monopole problems, among others. In addition, its has provided a causal interpretation of the origin of the observed anisotropy of the cosmic microwave background (CMB) radiation, and also the distribution of large scale structures. In the standard inflationary universe models, the acceleration of the universe is driven by a scalar field ϕ (inflaton) with an specific scalar potential, and the quantum fluctuations associated to this field generate the density perturbations seeding the structure formations at late time in the evolution of the universe. To date, the accumulating observational data, especially those coming from the CMB observations of WMAP satellite [2] indicate the power spectrum of the primordial density perturbations is nearly scale-invariant, just as predicted by the single-field inflation in the context of "slow-roll" over.

At the end of inflation the energy density of the universe is locked up in a combination of kinetic and potential energies of the scalar field, which drives inflation [3]. One path to defrost the universe after inflation is known as reheating [4]. Elementary theory of reheating was developed in [5] for the new inflationary scenario. During reheating, most of the matter and radiation of the universe are created usually via the decay of the scalar field that drives inflation, while the temperature grows in many orders of magnitude. It is at this point where the universe coincides with the Big-Bang model.

Of particular interest is a quantity known as the reheating temperature. The reheating temperature is associated to the temperature of the universe when the Big-Bang scenario begins, that is when the radiation epoch begins. In general, this epoch is generated by the decay of the inflaton field, which leads to a creation of particles of different kinds.

The stage of oscillations of the scalar field is a essential part for the standard mechanism of reheating. However, there are some models where the inflaton potential does not have a minimum and the scalar field does not oscillate. Here, the standard mechanism of reheating does not work [6]. These models are known in the literature like non-oscillating models, or simply NO models [7,8]. The NO models correspond to runaway fields such as module fields in string theory which are potentially useful for inflation model-building because they presents flat directions which survive the famous η-problem of inflation [9]. This problem is related to the fact that between the inflationary plateau and the quintessential tail there is a difference of over a hundred orders of magnitude. On the other hand, an important

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use of NO models is quintessential inflation, in which the tail of the potential can be responsible for the accelerated expansion of the present universe [10].

The first mechanism of reheating in this kind of model was the gravitational particle production [11], but this mechanism is quite inefficient, since it may lead to certain cosmological problems [12,13]. An alternative mechanism of reheating in NO models is the instant preheating, which introduce an interaction between the scalar field responsible for inflation an another scalar field [7].

An alternative mechanism of reheating in NO models is the introduction of the curvaton field [14]. The decay of the curvaton field into conventional matter offers an efficient mechanism of reheating, and does not necessarily introduce an interaction between the scalar field responsible of inflation and another scalar field [8]. The curvaton field has the property whose energy density is not diluted during inflation, so that the curvaton may be responsible for some or all the matter content of the universe at present.

Implications of string/M-theory to Friedman–Robertson–Walker cosmological models have recently attracted great attention, especially those related to brane–antibrane configurations as spacelike branes. The tachyon field associated with unstable D-branes, might be responsible for cosmological inflation at early evolution of the universe, due to tachyon condensation near the top of the effective scalar potential [15,16] which could add some new form of cosmological dark matter at late times [13]. In fact, historically, as was empathized by Gibbons [17], if the tachyon condensate starts to roll down the potential with small initial $\dot{\phi}$, then a universe dominated by this new form of matter will smoothly evolve from a phase of accelerated expansion (inflation) to an era dominated by a non-relativistic fluid, which could contribute to the dark matter specified above. We should note that during tachyonic inflation, the slow-roll over condition becomes $\dot{\phi}^2 < 2/3$ which is very different from the condition for non-tachyonic field $\dot{\phi}^2 < V(\phi)$. In this way, the tachyonic field should start rolling with a small value of $\dot{\phi}$ in order to have a long period of inflation [13]. In this way, the basic field equations for tachyon inflation become $3H\dot{\phi} + (1/V)(dV/d\phi) \approx 0$ and $3H^2 \approx \kappa_0 V$, where ϕ denotes a homogeneous tachyonic field (with unit $1/m_p$–1/energy–m_p is the Planck mass, so that $\dot{\phi}$ becomes dimensionless). $V = V(\phi)$ is the tachyonic potential, H is the Hubble factor and $\kappa_0 = 8\pi m_p^{-2}$. We have used units in which $c = \hbar = 1$. Dots mean derivatives with respect to time. These expressions should be compared with those corresponding to the standard case, where we have $3H\dot{\phi} + dV/d\phi \approx 0$ and $3H^2 \approx \kappa_0 V$. Thus, we observe a clear difference in the scalar field evolution equations, meanwhile the Friedman equation remains practically the same. Certainly, this modifications has important consequences, for instance, the slow-roll over parameters become quite differents [18].

In the following, we explore the curvaton reheating in tachyonic inflationary models with an exponential potential (i.e., a NO model). We follow a similar procedure described in Ref. [12]. As the energy density decreases, the tachyonic field makes a transition into a kinetic energy dominated regime bringing inflation to the end. Following Liddle and Ureña [12], we considered the evolution of the curvaton field through three different stages. Firstly, there is a period in which the tachyonic energy density is the dominant component, i.e., $\rho_\phi \gg \rho_\sigma$, even though the curvaton field survives the rapid expansion of the universe. The following stage, i.e., during the kinetic epoch [19], is that in which the curvaton mass becomes important. In order to prevent a period of curvaton-driven inflation, the universe must remain tachyon-driven until this time. When the effective mass of the curvaton becomes important, the curvaton field starts to oscillate around at the minimum of its potential. The energy density, associated to the curvaton field, starts to evolve as non-relativistic matter. At the final stage, the curvaton field decays into radiation and then the standard Big-Bang cosmology is recovered afterwards. In general, the decay of the curvaton field should occur before nucleosynthesis happens. Other constraints may arise depending on the epoch of the decay, which is governed by the decay parameter, Γ_σ. There are two scenarios to be considered, depending on whether the curvaton field decays before or after it becomes the dominant component of the universe.

In the first stage the dynamics of the tachyon field is described in the slow-roll over approach [13]. Nevertheless, after inflation, the term $V^{-1}\partial V/\partial\phi$ is negligible compared to the friction term. This epoch is called 'kinetic epoch' or 'kination' [19], and we will use the subscript 'k' to label the value of the different quantities at the beginning of this epoch. The kinetic epoch does not occur immediately after inflation, may exist a middle epoch where the tachyonic potential force is negligible respect to the friction term [20].

The dynamics of the Friedman–Robertson–Walker cosmology for the tachyonic field in the kinetic regimen, is described by the equations (see [20])

$$\frac{\ddot{\phi}}{1-\dot{\phi}^2} + 3H\dot{\phi} = 0, \tag{1}$$

and

$$3H^2 = \kappa_0 \rho_\phi. \tag{2}$$

The associated energy density of the tachyonic field, ρ_ϕ, is given by the expression

$$\rho_\phi = \frac{V(\phi)}{\sqrt{1-\dot{\phi}^2}}. \tag{3}$$

The tachyonic potential $V(\phi)$ is such that satisfies $V(\phi) \to 0$ as $\phi \to \infty$. It has been argued [21] that the qualitative tachyon dynamics of string theory can be describe by an exponential potential of the form

$$V(\phi) = V_0 e^{-\alpha\sqrt{\kappa_0}\phi}, \tag{4}$$

where α and V_0 are free parameters. In the following we shall take $\alpha > 0$. We should note that $\alpha\sqrt{\kappa_0}$ represents the tachyon mass [22]. An estimation of these parameters are given in Ref. [13], where $V_0 \sim 10^{-10} m_p^4$ and $\alpha \sim 10^{-5} m_p^2$.

From Eq. (1) we find a first integral for $\dot\phi$ in terms of a given by

$$\dot\phi^2=\frac{1}{1+Ca^6},\quad C=\frac{1-\dot\phi_k^2}{\dot\phi_k^2a_k^6}>0, \tag{5}$$

where C is an integration constant, $\dot\phi_k$ and a_k represent values at the beginning of the kinetic epoch for the time derivative of the tachyonic field and the scale factor, respectively. A universe dominated by tachyonic field would go under accelerate expansion if $\dot\phi^2<\frac{2}{3}$. The end of inflation is characterized by the value $\dot\phi^2_{\rm end}=\frac{2}{3}$. The value of $\dot\phi$ at the beginning of the kinetic epoch lies in the range $1\gtrsim\dot\phi_k^2\gtrsim\frac{2}{3}$.

From Eq. (3), after substituting the scalar potential $V(\phi)$ and $\dot\phi^2$ from Eq. (5) into Eq. (2), and considering that $\dot a=(da/d\phi)\dot\phi=(1+Ca^6)^{-1/2}\,da/d\phi$, we get

$$V(\phi)=V\big(\phi(a)\big)=V:=\left[V_0^{1/2}e^{-\alpha\sqrt{\kappa_0}\phi_k/2}-\sqrt{3}\frac{\alpha}{2}C^{1/4}I\right]^2, \tag{6}$$

where I represents the integral

$$\begin{aligned}I(a):=I:=&\int_{a_k}^{a}\frac{a'^{1/2}}{(1+Ca'^6)^{3/4}}\,da'\\ &=\frac{2}{3}a^{2/3}{}_2F_1\left[\frac{1}{4},\frac{3}{4},\frac{5}{4};-Ca^6\right]\\ &\quad-\frac{2}{3}a_k^{2/3}{}_2F_1\left[\frac{1}{4},\frac{3}{4},\frac{5}{4};-Ca_k^6\right],\end{aligned} \tag{7}$$

and ${}_2F_1$ is the hypergeometric function.

Now from Eq. (3) we get an explicit expression for the tachyonic energy density in terms of the scale factor

$$\begin{aligned}\rho_\phi&=\frac{V}{\sqrt{1-\dot\phi^2}}\\ &=\frac{\sqrt{1+Ca^6}}{\sqrt{C}a^3}\left[V_0^{1/2}e^{-\alpha\sqrt{\kappa_0}\phi_k/2}-\sqrt{3}\frac{\alpha}{2}C^{1/4}I\right]^2.\end{aligned} \tag{8}$$

Finally, in the kinetic epoch the tachyonic energy density and the Hubble factor can be written as follows,

$$\rho_\phi=\rho_\phi^k\frac{V}{V_k}\frac{a_k^3}{\sqrt{Ca_k^6+1}}\frac{\sqrt{Ca^6+1}}{a^3}, \tag{9}$$

$$H=H_k\left(\frac{V}{V_k}\right)^{1/2}\frac{a_k^{3/2}}{(Ca_k^6+1)^{1/4}}\frac{(Ca^6+1)^{1/4}}{a^{3/2}}, \tag{10}$$

respectively. Here, $H_k^2=\frac{\kappa_0}{3}\rho_\phi^k$. Note the difference for the standard case, where the energy density is definite by $\rho_\phi^{(std)}=\dot\phi^2/2+V(\phi)$, and has a behavior during the kinetic epoch like stiff matter, i.e., $\rho_\phi^{(std)}=\rho_\phi^k(a_k/a)^6$ [12]. In this way, the Hubble factor in the standard theory follows the law $H^{(std)}=H_k(a_k/a)^3$.

We now study the dynamic of the curvaton field, σ, through different stages. This permits us to find some constraints of the parameters, and thus, to have a viable curvaton scenario. We considered that the curvaton field obeys the Klein–Gordon equation, and for simplicity, we assume that its scalar potential is given by

$$U(\sigma)=\frac{m^2\sigma^2}{2}, \tag{11}$$

where m is the curvaton mass.

First of all, it is assumed that the tachyonic energy density, ρ_ϕ, is the dominant component when it is compared with the curvaton energy density, ρ_σ. In the next stage, the curvaton field oscillates around the minimum of the effective potential $U(\sigma)$. Its energy density evolves as a non-relativistic matter, and during the kinetic epoch the universe remains tachyonic-dominated. The last stage corresponds to the decay of the curvaton field into radiation, and then, the standard Big-Bang cosmology is recovered.

In the inflationary regimen is supposed that the curvaton mass satisfied the condition $m\ll H_f$ and its dynamics is described in detail in Refs. [12,23,24]. During inflation, the curvaton would roll down its potential until its kinetic energy is depleted by the exponential expansion and only then, i.e., only after its kinetic energy is almost vanished, it becomes frozen and assumes roughly a constant value i.e., $\sigma_i\approx\sigma_f$. Here the subscripts i and f are used to denote the beginning and the end of inflation, respectively.

The hypothesis is that during the kinetic epoch the Hubble parameter decreases so that its value is comparable with the curvaton mass, i.e., $m\simeq H$. From Eq. (10), we obtain

$$\frac{m}{H_k}=\left(\frac{V_m}{V_k}\right)^{1/2}\frac{a_k^{3/2}}{(Ca_k^6+1)^{1/4}}\frac{(Ca_m^6+1)^{1/4}}{a_m^{3/2}}, \tag{12}$$

where the 'm' label represents the quantities at the time when the curvaton mass is of the order of H during the kinetic epoch.

In order to prevent a period of curvaton-driven inflation, the universe must still be dominated by the tachyonic matter, i.e., $\rho_\phi|_{a_m}=\rho_\phi^{(m)}\gg\rho_\sigma(\sim U(\sigma_f)\simeq U(\sigma_i))$. This inequality allows us to find a constraint on the initial values of the curvaton field in the inflationary epoch. Hence, from Eq. (2), at the moment when $H\simeq m$ we get the restriction

$$\frac{m^2\sigma_i^2}{2\rho_\phi^{(m)}}=\frac{4\pi}{3}\frac{m^2\sigma_i^2}{m_p^2m^2}\ll 1\ \Rightarrow\ \sigma_i^2\ll\frac{3}{4\pi}m_p^2. \tag{13}$$

This value is the same to that found in the standard case [12].

The ratio between the potential energies at the end of inflation is given by

$$\frac{U_f}{V_f}=\frac{4\pi}{3}\frac{m^2\sigma_i^2}{m_p^2H_f^2}\ll\frac{m^2}{H_f^2}. \tag{14}$$

Here, we have used for $V_f=(3/8\pi)H_f^2m_p^2$ and Eq. (13). Thus, the curvaton mass should obey the constraint in the tachyonic model

$$m\ll H_f, \tag{15}$$

which gives from Eq. (14) that $U_f\ll V_f$. We should note that the condition given by Eq. (15) is inherent to the nature of the curvaton field, since the reason $m \ll H_f$ is because only then can be the curvaton superhorizon perturbations of σ be generated during inflation. In this way, the condition $m \ll H_f$ is a fundamental prerequisite for the curvaton mechanism.

After the mass of curvaton field becomes important, i.e., $m \simeq H$, its energy decays like non-relativistic matter in the form

$$\rho_\sigma = \frac{m^2\sigma_i^2}{2}\frac{a_m^3}{a^3}. \tag{16}$$

As we have claimed the curvaton decay could be occur in two different possible scenarios. In the first scenario, when the curvaton comes to dominates the cosmic expansion (i.e., $\rho_\sigma > \rho_\phi$), there must be a moment when the tachyonic and curvaton energy densities becomes equal. From Eqs. (9), (10) and (16) at the time when $\rho_\sigma = \rho_\phi$, which happens when $a = a_{\rm eq}$, we get

$$\begin{aligned} \left.\frac{\rho_\sigma}{\rho_\phi}\right|_{a=a_{\rm eq}} &= \frac{4\pi\sigma_i^2 m^2}{3H_k^2 m_p^2}\frac{V_k a_m^3\sqrt{Ca_k^6+1}}{V_{\rm eq}a_k^3\sqrt{Ca_{\rm eq}^6+1}} \\ &= \frac{4\pi\sigma_i^2}{3m_p^2}\frac{V_m}{V_{\rm eq}}\sqrt{\frac{Ca_m^6+1}{Ca_{\rm eq}^6+1}} = 1. \end{aligned} \tag{17}$$

Now from Eqs. (10), (12) and (17), we obtain a relation for the Hubble parameter, $H_{\rm eq}$, in terms of curvaton parameters and the ratio of the scale factor at different times, given by

$$\begin{aligned} H_{\rm eq} &= H_k\left(\frac{V_{\rm eq}}{V_k}\right)^{1/2}\frac{a_k^{3/2}}{(Ca_k^6+1)^{1/4}}\frac{(Ca_{\rm eq}^6+1)^{1/4}}{a_{\rm eq}^{3/2}} \\ &= \sqrt{\frac{4\pi\sigma_i^2}{3m_p^2}}\left[\frac{a_m}{a_{\rm eq}}\right]^{3/2} m. \end{aligned} \tag{18}$$

Notice that this result coincides with the one obtained in standard case.

On the one hand, the decay parameter Γ_σ is constrained by nucleosynthesis. For this, it is required that the curvaton field decays before of nucleosynthesis, which means $H_{\rm nucl} \sim 10^{-40} m_p < \Gamma_\sigma$. On the other hand, we also require that the curvaton decay occurs after $\rho_\sigma > \rho_\phi$, and $\Gamma_\sigma < H_{\rm eq}$ so that we get a constraint on the decay parameter,

$$10^{-40} m_p < \Gamma_\sigma < \sqrt{\frac{4\pi\sigma_i^2}{3m_p^2}}\left[\frac{a_m}{a_{\rm eq}}\right]^{3/2} m. \tag{19}$$

Until now, it is interesting to give an estimation of the constraint of the parameters of our model, by using the scalar perturbation related to the curvaton field. During the time the fluctuations are inside the horizon, they obey the same differential equation as the inflaton fluctuations do, from which we conclude that they acquire the amplitude $\delta\sigma_i \simeq H_i/2\pi$. Once the fluctuations are outside the horizon, they obey the same differential equation that the unperturbed curvaton field does and then we expect that they remain constant during inflation. The spectrum of the Bardeen parameter P_ζ, whose observed value is about 2×10^{-9}, allows us to determine the initial value of the curvaton field in terms of the parameter α. At the time when the decay of the curvaton fields occur, the Bardeen parameter becomes [14]

$$P_\zeta \simeq \frac{1}{9\pi^2}\frac{H_i^2}{\sigma_i^2}. \tag{20}$$

The spectrum of fluctuations is automatically Gaussian for $\sigma_i^2 \gg H_i^2/4\pi^2$, and is independent of Γ_σ [14]. This feature will simplify the analysis in the space parameter of our models. Moreover, the spectrum of fluctuations is the same as in the standard scenario.

From Eq. (20) and by using that $H_i^2 = H_f^2(2N+1)$ and $H_f^2 = \alpha^2\kappa_0/6$ [13], we could relate the perturbation with the parameters of the model in such way that we could write

$$\frac{27\pi}{4}\frac{P_\zeta}{(2N+1)}\sigma_i^2 = \frac{\alpha^2}{m_p^2}. \tag{21}$$

This expression allows us to the above equation permit fix the initial value of the curvaton field in terms of the free parameter α. By using Eq. (21), the constraint Eq. (15) becomes

$$\frac{m}{m_p} \ll 3\pi\frac{P_\zeta^{1/2}}{(2N+1)^{1/2}}\frac{\sigma_i}{m_p}. \tag{22}$$

Finally, Eq. (19) restricts the value of the decay parameter Γ_σ, which can be transformed into another constraint upon m and σ_i, so that

$$\frac{m}{m_p}\sqrt{\frac{\sigma_i^2}{m_p^2}} \gg \sqrt{\frac{3}{4\pi}}\times 10^{-40}, \tag{23}$$

where we have used the condition $a_m < a_{\rm eq}$, and Eq. (19).

On the other hand, for the second scenario, the decay of the field happens before this it dominates the cosmological expansion, that is, we need that the curvaton field decays before that its energy density becomes greater than the tachyonic one. Additionally, the mass is no-negligible so that we could use Eq. (16). The curvaton decays at a time when $\Gamma_\sigma = H$ and then from Eq. (10) we get

$$\frac{\Gamma_\sigma}{H_k} = \left(\frac{V_d}{V_k}\right)^{1/2}\frac{a_k^{3/2}}{(Ca_k^6+1)^{1/4}}\frac{(Ca_d^6+1)^{1/4}}{a_d^{3/2}}, \tag{24}$$

where 'd' labels the different quantities at the time when the curvaton decays, allowing the curvaton field decays after the mass takes importance, so that $\Gamma_\sigma < m$; and before that the curvaton field dominates the expansion of the universe, i.e., $\Gamma_\sigma > H_{\rm eq}$ (see Eq. (18)). Thus,

$$\sqrt{\frac{4\pi\sigma_i^2}{3m_p^2}}\left[\frac{a_m}{a_{\rm eq}}\right]^{3/2} m < \Gamma_\sigma < m. \tag{25}$$

Notice that the range of Γ_σ is the same that obtained in the standard case.

Now for the second scenario, the curvaton decays at the time when $\rho_\sigma < \rho_\phi$. If we defined the r_d parameter as the ratio between the curvaton and the tachyonic energy densities, evaluated at $a = a_d$ and for $r_d \ll 1$, the Bardeen parameter is given by [14,25]

$$P_\zeta \simeq \frac{r_d^2}{36\pi^2}\frac{H_i^2}{\sigma_i^2}. \tag{26}$$

With the help of Eq. (24) we obtain

$$r_d=\frac{\rho_\sigma}{\rho_\phi}\bigg|_{a=a_d}=\frac{4\pi\sigma_i^2m^2}{3H_k^2m_p^2}\frac{V_ka_m^3\sqrt{Ca_k^6+1}}{V_da_k^3\sqrt{Ca_d^6+1}}=\frac{4\pi\sigma_i^2}{3\Gamma_\sigma^2}\frac{m^2}{m_p^2}\frac{a_m^3}{a_d^3}. \tag{27}$$

From Eq. (25) we obtain that $r_d<(a_{\text{eq}}/a_d)^3$, then from $r_d\ll 1$ we get $(a_{\text{eq}}/a_d)^3\ll 1$. Therefore, the condition $a_{\text{eq}}\ll a_d$ allows us to use expression (26) for the Bardeen parameter.

From expression (26) and (27) we could write

$$\frac{\sigma_i^2}{m_p^2}=\frac{81}{4}\frac{m_p^2}{m^4}\left(\frac{a_d}{a_m}\right)^6\frac{P_\zeta}{(2N+1)}\frac{\Gamma_\sigma^4}{H_f^2}, \tag{28}$$

and thus the expression (25) becomes

$$\sqrt{27\pi}\frac{a_d^3}{a_m^{3/2}a_{\text{eq}}^{3/2}}\frac{P_\zeta^{1/2}}{(2N+1)^{1/2}}\frac{m_p}{m^2H_f}\Gamma_\sigma^2<\frac{\Gamma_\sigma}{m}<1. \tag{29}$$

Even though the study of gravitational waves was developed in Ref. [26] for the tachyonic model, it is interesting to give an estimation of the constraint on the curvaton mass, using this type of tensorial perturbation. Under the approximation give in Ref. [27], the corresponding gravitational wave amplitude in the tachyonic model may be written as

$$h_{\text{GW}}^2\simeq C_1\frac{V_i}{m_p^4},$$

where the constant $C_1\approx 10^{-3}$. Now, using that $V_i=V_f(2N+1)$, we obtain

$$h_{\text{GW}}^2\simeq 3C_1(2N+1)\frac{H_f^2}{8\pi m_p^2}. \tag{30}$$

In this way, from Eqs. (15) and (30) we get that

$$m^2\ll 8\pi m_p^2\frac{h_{\text{GW}}^2}{3C_1(2N+1)}. \tag{31}$$

If we consider that h_{GW} of the order of 10^{-5} and if take the number of e-fold to be $N\simeq 70$ (but in context of the curvaton may be much lower that this value, let say 45 or so, since the inflationary scale can be lower) we find that the above equation gives the following upper limit for the curvaton mass

$$m\ll 10^{-4}m_p\sim 10^{15}\ \text{GeV}, \tag{32}$$

which coincides with the limit reported in Ref. [28]. We note that in this model we have $V''=3\alpha^2H_f^2$, where the prime denotes differentiation with respect to the tachyonic field ϕ, and if $V''>H_f^2m_p^4$ the curvature perturbations can becomes too large compare to the COBE observations [28], so that the inflationary scale cannot be much larger than the scale of grand unification (inflation does not produce perturbations if $\alpha>m_p^2/\sqrt{3}$). This means that $H_f\leqslant 10^{13}$ GeV [6]. Hence, the bounds in Eq. (32) is redundant unless the inflaton does not produce any curvature perturbations.

In order to give an estimation of the gravitational wave, we move to the kinetic epoch in which the energy density of gravitational waves evolves as in Refs. [10,29]

$$\frac{\rho_g}{\rho_\phi}\sim h_{\text{GW}}^2\left(\frac{V_k}{V}\frac{a_k}{a}\right)\sqrt{\frac{Ca_k^6+1}{Ca^6+1}}, \tag{33}$$

where we have used the same approximation than that used in Ref. [27].

On the other hand, when the curvaton field decays, i.e., $(\Gamma_\sigma=m)$ it produces radiation which decays as $1/a^4$. Then we may write

$$\rho_r^{(\sigma)}=\frac{m^2\sigma_i^2}{2}\frac{a_m^3}{a_d^3}\frac{a_d^4}{a^4}. \tag{34}$$

If the radiation produced from the curvaton scalar field is equal to the tachyonic density, i.e., $\rho_\sigma^{(r)}=\rho_\phi$, at the time in which $a=a_{\text{eq}}$, then we could keep the gravitational waves stable, so that

$$\frac{\rho_r^{(\sigma)}}{\rho_\phi}\bigg|_{a=a_{\text{eq}}}=\frac{4\pi m^2}{3m_p^2}\frac{\sigma_i^2}{H_k^2}\frac{V_k}{V_{\text{eq}}}\left(\frac{a_m}{a_k}\right)^3\left(\frac{a_d}{a_{\text{eq}}}\right)\sqrt{\frac{Ca_k^6+1}{Ca_{\text{eq}}^6+1}}=1, \tag{35}$$

and used Eqs. (12), (24) and (33), we obtain a constrains during the kinetic epoch given by

$$\frac{m\sigma_i}{m_p}\gg h_{\text{GW}}H_k\left(\frac{\Gamma_\sigma}{H_k}\right)^{1/3}\left(\frac{V_k}{V_d}\right)^{1/6}\left(\frac{a_k}{a_m}\right)^{3/2}\left(\frac{Ca_k^6+1}{Ca_d^6+1}\right)^{1/12}, \tag{36}$$

where we have used $\rho_g/\rho_r\ll 1$ at the time in which $a=a_{\text{eq}}$.

We note that from Eqs. (13), (32) and (36) we obtain a bound for the m parameter, i.e.,

$$\sqrt{\frac{4\pi}{3}}h_{\text{GW}}H_k\left(\frac{\Gamma_\sigma}{H_k}\right)^{1/3}\left(\frac{V_k}{V_d}\right)^{1/6}\left(\frac{a_k}{a_m}\right)^{3/2}\left(\frac{Ca_k^6+1}{Ca_d^6+1}\right)^{1/12} \ll m\ll 10^{-4}m_p. \tag{37}$$

It is interesting to note that in this case we have obtained a bound from bellow for the m curvaton mass.

In the first scenario, our computes allow to get the reheating temperature as hight as $10^{-9}m_p$, since the decay parameter $\Gamma_\sigma\propto T_{\text{rh}}^2/m_p$, where T_{rh} represents the reheating temperature. Here, we have used Eqs. (19), (21) and (32), $a_m/a_{\text{eq}}\sim 10^{-1}$ and $\alpha\sim 10^{-5}m_p^2$. We should compare this bound with the bound coming from gravitino over-production, which gives $T_{\text{rh}}\leqslant 10^{-10}m_p$ [30].

In the second scenario from Eqs. (25) and (32), we could estimate the reheating temperature to be of the order of $\sim 10^{-3}m_p$ as an upper limit.

As it was reported in Ref. [13] at the end of inflation ρ_ϕ at best could scales as a^{-3}, it is valid irrespectively of the form of the tachyonic potential provides it satisfies $V(\phi)\to 0$ as $\phi\to\infty$. However, this is not in general since in our particular case, we have found that it is possible to get a more complex expression for the dependence of ρ_ϕ in terms of the scale factor. This could be seen from Eq. (8).

On the other hand, the shape of the tachyon condensate effective potential depends on the system under consideration. In bosonic string theory, for instance, this potential has a maximum, $V = V_0$, at $\phi = 0$, where V_0 is the tension of some unstable bosonic D-brane. A local minimum, $V = 0$, generically at $\phi \to \infty$, corresponding to a metastable closed bosonic string vacuum, and a runaway behavior for negative ϕ. An exact classical potential (i.e., exact to all orders in α', but only at tree level in g_s) encompassing these properties has been considered [31],

$$V(\phi) = V_0(1 + \phi/\phi_0)\exp(-\phi/\phi_0), \tag{38}$$

where the parameters V_0 and ϕ_0 in terms of the string length l_s and the open string coupling constant g_s, are given by

$$V_0 = \frac{v_0}{g_s l_s^4 (2\pi)^3}, \qquad \phi_0 = \frac{1}{\alpha\sqrt{\kappa_0}} = \tau_0 l_s, \tag{39}$$

with v_0 and τ_0 dimensionless parameters, such that V_0/v_0 is the tension of a D3-brane and $\tau_0 l_s$ is the inverse tachyon mass [22]. The gravitational coupling in 4 dimensions is given in terms of the stringy parameters by

$$\kappa_0 \equiv 8\pi G_N = \frac{8\pi}{m_p^2} = \pi g_s^2 l_s^2 \left(\frac{l_s}{R}\right)^6 = \frac{g_s^2 l_s^2}{v},$$
$$v = \frac{1}{\pi}\left(\frac{R}{l_s}\right)^6. \tag{40}$$

Here, R is the compactification radius of the compact 6-dimensional manifold, taken to be a 6-torus. For the $D = 4$ effective theory to be applicable, one usually requires that $R \gg l_s$, i.e., $v \gg 1$.

From Eq. (29) we obtain that

$$36\pi \frac{a_d^3}{(a_m a_{\rm eq})^{3/2}} \frac{P_\zeta^{1/2}}{(2N+1)^{1/2}} \frac{\Gamma_\sigma^2}{m^2} < \frac{g_s}{\tau_0 v^{1/2}}. \tag{41}$$

From Eq. (41) and taking $N = 50$, $P_\zeta \sim 10^{-5}$, and $\tau_0 = 1$, we find a constraint (from the reheating scenario) for the parameter g_s coming from string theory, which is given by

$$g_s^2 > 10^{-4} \frac{a_d^6}{(a_m a_{\rm eq})^3} \frac{\Gamma_\sigma^4}{m^4} v, \tag{42}$$

the above expression give to us a lower bound for the string coupling constant. From the amplitude of gravitational waves produced during inflation the upper bound is $g_s^2 \leqslant 10^{-9} v$ [6]. In this way, we have the following constraint for g_s^2/v

$$10^{-4} \frac{a_d^6}{(a_m a_{\rm eq})^3} \frac{\Gamma_\sigma^4}{m^4} < \frac{g_s^2}{v} \leqslant 10^{-9}. \tag{43}$$

Summarizing, we have describe curvaton reheating in tachyonic inflationary model in which we have considered two cases. In the first case the curvaton dominates the universe before it decay. Our results are specified by Eqs. (19) and (22), and we see that they are identical with the standard curvaton scenario [12]. In the second case where the curvaton decaying before domination, we have arrived to Eq. (29), which represents one of the most important constraint by using the curvaton approach.

In conclusion, we have introduced the curvaton mechanism into NO inflationary tachyonic model as another possible solution to the problem of reheating, where there is not need to introduce an interaction between the tachyonic and some auxiliary scalar field.

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