Tachyon warm inflationary universe models

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Abstract. Warm inflationary universe models in a tachyon field theory are studied. General conditions required for these models to be realizable are derived and discussed. We describe scalar perturbations (in the longitudinal gauge) and tensor perturbations for these scenarios. We develop our models for a constant dissipation parameter $\Gamma$ in one case and one dependent on $\phi$ in the other case. We have been successful in describing such inflationary universe models. We use recent astronomical observations for constraining the parameters appearing in our model. Also, our results are compared with their analogues found in the cool inflationary case.

Keywords: cosmological perturbation theory, inflation, cosmology of theories beyond the SM, physics of the early universe
1. Introduction

It is well known that many long-standing problems of the big bang model (horizon, flatness, monopoles, etc) may find a natural solution in the framework of the inflationary universe model [1,2]. One of the successes of the inflationary universe model is that it provides 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 [3,4]. In standard inflationary universe models, the acceleration of the universe is driven by a scalar field (inflaton) with a specific scalar potential, and the quantum fluctuations associated with this field generate the density perturbations seeding the structure formation at late time in the evolution of the universe. The standard inflationary model is divided into two regimes: the slow roll and reheating epochs. In the slow roll period the universe inflates and all interactions between the inflation scalar field and any other field are typically neglected. Subsequently, a reheating period is invoked to end the period of inflation. After reheating, the universe is filled with radiation [5] and thus the universe is connected with the radiation big bang phase.

Warm inflation is an alternative mechanism for having successful inflation and avoiding the reheating period [6]. In warm inflation, dissipative effects are important during inflation, so that radiation production occurs concurrently with the inflationary expansion. The dissipating effect arises from a friction term which describes the processes of the scalar field dissipating into a thermal bath via its interaction with other fields. Also, warm inflation shows how thermal fluctuations during inflation may play a dominant role in producing the initial perturbations. In such models, the density fluctuations arise from thermal rather than quantum fluctuations [7]. These fluctuations have their origin...
in the hot radiation and influence the inflaton through a friction term in the equation of motion of the inflaton scalar field [8]. Among the most attractive features of these models, warm inflation ends when the universe heats up to become radiation dominated; at this epoch the universe stops inflating and ‘smoothly’ enters a radiation dominated big bang phase [6]. The matter components of the universe are created by the decay of either the remaining inflationary field or the dominant radiation field [9].

On the other hand, implications of string/M-theory for Friedmann–Robertson–Walker cosmological models have recently attracted great attention, in particular, those related to brane–antibrane configurations such as space-like branes [10]. The tachyon field associated with unstable D-branes might be responsible for cosmological inflation in the early evolution of the universe, due to tachyon condensation near the top of the effective scalar potential [11] which could also add some new form of cosmological dark matter at late times [12]. In fact, historically, as was emphasized by Gibbons [13], if the tachyon condensate starts to roll down the potential with small initial velocity, 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.

Usually, in any model of warm inflation, the scalar field which drives inflation is the standard inflaton field. As far as we know, a model in which warm inflation is driven by a tachyonic scalar field has not yet been studied. The main goal of the present work is to investigate the possible realization of a tachyonic warm inflationary universe model. In this way, we study cosmological perturbations, which are expressed in terms of different parameters appearing in our model. These parameters are constrained from the WMAP three-year data [4].

In section 2, the dynamics of the tachyon warm inflationary models is obtained. In section 3, the cosmological perturbations are investigated. In section 4, we use an exponential scalar potential in the high dissipation regime. There, we distinguish two cases: in the first scenario, we consider a constant dissipation rate \( \Gamma \) and in the second one, the dissipation rate is considered to be a function of the tachyonic field, \( \phi \). Here, we study possible solutions of the approximated field equations. In section 5 we develop, in the same approximation, i.e. in the high dissipation regime, models which are characterized by a power law scale factor. Finally, in section 6, we give some conclusions.

## 2. Warm inflationary model

As was noted by Gibbons [13], a rolling tachyon condensate in a spatially flat Friedmann–Robertson–Walker (FRW) cosmological model is described in terms of an effective fluid with energy–momentum tensor \( T_{\mu}^{\nu} = \text{diag}(-\rho_\phi, p_\phi, p_\phi, p_\phi) \), where the energy density, \( \rho_\phi \), and pressure, \( p_\phi \), associated with the tachyon field are defined by

\[
\rho_\phi = \frac{V(\phi)}{\sqrt{1 - \dot{\phi}^2}} \tag{1}
\]

and

\[
p_\phi = -V(\phi) \sqrt{1 - \dot{\phi}^2}, \tag{2}
\]
respectively. Here, $\phi$ denotes the tachyon field (with unit $1/m_p$, where $m_p$ represents the Planck mass) and $V(\phi) = V$ is the effective potential associated with this tachyon field.

The dynamics of the FRW cosmological model in the tachyonic case, in the warm inflationary scenario, is described by the equations

$$H^2 = \kappa \left[ \frac{V}{\sqrt{1 - \dot{\phi}^2}} + \rho_\gamma \right],$$

(3)

$$\dot{\rho}_\phi + 3H(\rho_\phi + p_\phi) = -\Gamma\dot{\phi}^2 \implies \frac{\dot{\phi}}{(1 - \dot{\phi}^2)} + 3H\dot{\phi} + \frac{V_{,\phi}}{V} = -\frac{\Gamma}{V\sqrt{1 - \dot{\phi}^2}}$$

(4)

and

$$\dot{\rho}_\gamma + 4H\rho_\gamma = \Gamma\dot{\phi}^2,$$

(5)

where $H = \dot{a}/a$ is the Hubble factor, $a$ is a scale factor, $\rho_\gamma$ is the energy density of the radiation field and $\Gamma$ is the dissipation coefficient, with unit $m_p^5$. The dissipative coefficient is responsible for the decay of the tachyon scalar field into radiation during the inflationary regimen. In general, $\Gamma$ can be assumed as a function of $\phi$, and thus $\Gamma = f(\phi) > 0$ by the second law of thermodynamics. Dots mean derivatives with respect to time, $V_{,\phi} = \partial V(\phi)/\partial \phi$, $\kappa = 8\pi/(3m_p^2)$ and we use units in which $c = \hbar = 1$.

During the inflationary era the energy density associated with the tachyonic field is the order of the potential, i.e. $\rho_\phi \sim V$, and dominates over the energy density associated with the radiation field, i.e. $\rho_\phi > \rho_\gamma$. Assuming the set of slow roll conditions $\dot{\phi}^2 \ll 1$ and $\ddot{\phi} \ll (3H + \Gamma/V)\dot{\phi}$ [6], the Friedmann equation is reduced to

$$H^2 = \kappa V$$

(6)

and equation (4) becomes

$$3H [1 + r] \dot{\phi} = -\frac{V_{,\phi}}{V} = -(\ln(V))_{,\phi},$$

(7)

where $r$ is the rate defined as

$$r = \frac{\Gamma}{3HV}$$

(8)

and parametrizes the dissipation of our model. For the high (or weak) dissipation regimen, $r \gg 1$ (or $r < 1$), i.e. the dissipation coefficient $\Gamma$ is much greater (or less) than the product $3HV$.

If we also consider that during tachyon warm inflation radiation production is quasi-stable, then $\dot{\rho}_\gamma \ll 4H\rho_\gamma$ and $\dot{\rho}_\gamma \ll \Gamma\dot{\phi}^2$. From equation (5) we obtain that the energy density of the radiation field becomes

$$\rho_\gamma = \frac{\Gamma\dot{\phi}^2}{4H} = \sigma T_r^4,$$

(9)

where $\sigma$ is the Stefan–Boltzmann constant and $T_r$ is the temperature of the thermal bath. Using equations (7)–(9) we get

$$\rho_\gamma = \sigma T_r^4 = \frac{r}{12\kappa (1 + r)^2} \left[ \frac{V_{,\phi}}{V} \right]^2.$$
Introducing the dimensionless slow roll parameter
\[ \epsilon \equiv -\frac{\dot{H}}{H^2} = \frac{1}{6\kappa} \left( \frac{1}{1+r} \right) \left( V_{,\phi} \right)^2 \frac{1}{V}, \] (11)

it is possible to find a relation between the energy densities \( \rho_\gamma \) and \( \rho_\phi \) given by
\[ \rho_\gamma = \frac{r}{2(1+r)} \epsilon \rho_\phi. \] (12)

Here, the energy density of the tachyonic field during inflation corresponds to the potential energy, i.e. \( \rho_\phi \sim V \).

The second slow roll parameter \( \eta \) becomes
\[ \eta \equiv -\frac{\ddot{H}}{H \dot{H}} \simeq \frac{1}{3\kappa(1+r)V} \left[ V_{,\phi\phi} - \frac{1}{2} \left( \frac{V_{,\phi}}{V} \right)^2 \right]. \] (13)

The condition under which the tachyon warm inflation epoch could take place can be summarized with the parameter \( \epsilon \) satisfying the relation \( \epsilon < 1 \). This condition is analogous to the requirement in which \( \ddot{a} > 0 \). The condition given above is rewritten in terms of the densities by using equation (12) as
\[ \rho_\phi > 2 \frac{(1+r)}{r} \rho_\gamma, \] (14)

which describes the epoch where tachyon warm inflation occurs.

On the other hand, inflation ends when the universe heats up, at a time when \( \epsilon \simeq 1 \), which implies
\[ \rho_\phi \simeq 2 \frac{(1+r)}{r} \rho_\gamma. \] (15)

The number of e-folds at the end of inflation is given by
\[ N(\phi) = -3 \kappa \int_{\phi_i}^{\phi_f} \frac{V^2}{V_{,\phi}} (1+r) \, d\phi'. \] (16)

In the following, the subscripts \( i \) and \( f \) are used to denote the beginning and the end of inflation.

### 3. The perturbations

In this section we will describe scalar perturbations in the longitudinal gauge, and then we will continue describing tensor perturbations.

#### 3.1. Scalar perturbations

By using the longitudinal gauge in the perturbed FRW metric, we write
\[ ds^2 = (1+2\Phi) \, dt^2 - a^2(t) (1-2\Psi) \delta_{ij} \, dx^i \, dx^j, \] (17)
where $\Phi = \Phi(t, x)$ and $\Psi = \Psi(t, x)$ are gauge-invariant variables introduced by Bardeen [14]. In momentum space, for the Fourier components $e^{ikx}$, with $k$ being the wavenumber, the following set of equations (which are derived from the perturbed Einstein field equations; we omit the subscript $k$ here) arises:

$$\dot{\Phi} + H\Phi = \frac{4\pi}{m_p^2} \left[ \frac{-4\rho_{\gamma} a v}{3k} + \frac{V\dot{\phi}}{\sqrt{(1 - \dot{\phi}^2)}} \right],$$

$$\frac{(\delta\phi^\gamma)}{1 - \dot{\phi}^2} + \left[ 3H + \frac{\Gamma}{V} \right] (\delta\phi) + \left[ \frac{k^2}{a^2} + (\ln(V))_{,\phi} + \dot{\phi} \left( \frac{\Gamma}{V} \right)_{,\phi} \right] \delta\phi$$

$$= \left[ \frac{1}{1 - \dot{\phi}^2} + 3 \right] \dot{\Phi} + \left[ \phi \frac{\Gamma}{V} - 2(\ln(V)), \phi \right] \Phi,$$

$$(\delta\rho_{\gamma}) + 4H\delta\rho_{\gamma} + \frac{4}{3}k a \rho_{\gamma} v - 4\rho_{\gamma} \dot{\Phi} - \dot{\phi}^2 \Gamma_{,\phi} \delta\phi - \Gamma\dot{\phi}[2(\delta\phi) - 3\dot{\phi}\Phi] = 0$$

and

$$\dot{v} + 4Hv + \frac{k}{a} \left[ \Phi + \frac{\delta\rho_{\gamma}}{4\rho_{\gamma}} + \frac{3\Gamma\dot{\phi}}{4\rho_{\gamma}} \delta\phi \right] = 0,$$

where $v$ appears from the decomposition of the velocity field $\delta u_j = -(iak_j/k)v e^{ikx} (j = 1, 2, 3)$ [14].

Note that in the case of the scalar perturbations the tachyon and the radiation fields are interacting. Therefore, isocurvature (or entropy) perturbations are generated, besides the adiabatic ones. This occur because warm inflation can be considered as an inflationary model with two basic fields [15, 16]. In this context, dissipative effects themselves can produce a variety of spectra ranging between red and blue [7, 15], thus producing the running blue to red spectrum suggested by WMAP three-year data [4]. We will return to this point soon.

Since what we need are the non-decreasing adiabatic and isocurvature modes on the large scale $k \ll aH$ (which turn out to be weak time dependent quantities), without loss of generality we may consistently neglect $\dot{\Phi}$ and those terms containing two-times derivatives and, combining with the slow roll conditions, the equations for $\Phi$, $\delta\phi$, $\delta\rho_{\gamma}$ and $v$ are reduced to

$$\Phi = \frac{4\pi}{m_p^2} \left( \frac{V\dot{\phi}}{H} \right) \left[ 1 + \frac{\Gamma}{4HV} + \frac{\Gamma_{,\phi}}{48H^2V} \right] \delta\phi,$$

$$\frac{3H + \frac{\Gamma}{V}}{} (\delta\phi) + \left[ (\ln(V))_{,\phi} + \dot{\phi} \left( \frac{\Gamma}{V} \right)_{,\phi} \right] \delta\phi = \left[ \phi \frac{\Gamma}{V} - 2(\ln(V)), \phi \right] \Phi,$$

$$\delta\rho_{\gamma} \simeq \frac{\dot{\phi}^2}{4H} \left[ \Gamma_{,\phi} \delta\phi - 3\Gamma\Phi \right] \Rightarrow \frac{\delta\rho_{\gamma}}{\rho_{\gamma}} \simeq \frac{\Gamma_{,\phi}}{\Gamma} \delta\phi - 3\Phi$$

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$$\frac{3H + \frac{\Gamma}{V}}{} (\delta\phi) + \left[ (\ln(V))_{,\phi} + \dot{\phi} \left( \frac{\Gamma}{V} \right)_{,\phi} \right] \delta\phi = \left[ \phi \frac{\Gamma}{V} - 2(\ln(V)), \phi \right] \Phi,$$

$$\delta\rho_{\gamma} \simeq \frac{\dot{\phi}^2}{4H} \left[ \Gamma_{,\phi} \delta\phi - 3\Gamma\Phi \right] \Rightarrow \frac{\delta\rho_{\gamma}}{\rho_{\gamma}} \simeq \frac{\Gamma_{,\phi}}{\Gamma} \delta\phi - 3\Phi$$
and
\[ v \simeq -\frac{k}{4aH} \left[ \Phi + \frac{\delta \rho_\gamma}{4\rho_\gamma} + \frac{3\Gamma}{4\rho_\gamma} \delta \phi \right], \tag{26} \]
respectively.

The above equations can be solved taking \( \phi \) as the independent variable instead of \( t \). With the help of equation (7) we find
\[ \left( 3H + \frac{\Gamma}{V} \right) \frac{d}{dt} = \left( 3H + \frac{\Gamma}{V} \right) \dot{\phi} \frac{d}{d\phi} = -\left( \ln(V),\phi \right) \frac{d}{d\phi} \tag{27} \]
and introducing an auxiliary function \( \varphi \) given by
\[ \varphi = \frac{\delta \phi}{\left( \ln(V),\phi \right)} \exp \left[ \int \frac{1}{3H + \frac{\Gamma}{V}} \left( \frac{\Gamma}{V} \right) d\phi \right], \tag{28} \]
we obtain the following equation for \( \varphi \):
\[ \frac{\varphi_{,\phi}}{\varphi} = -\frac{9}{8} \frac{\left( \Gamma/V + 2H \right)}{\Gamma/V + 3H} \left[ \frac{\Gamma + 4HV - \frac{\Gamma,\phi(\ln(V),\phi)}{12H(3H + \Gamma/V)}}{V} \right] \tag{29} \]
Solving equation (29) and using equation (28) we find that
\[ \delta \phi = C \left( \ln(V),\phi \right) \exp[\Im(\phi)], \tag{30} \]
where \( \Im(\phi) \) is given by
\[ \Im(\phi) = -\int \left[ \frac{1}{3H + \frac{\Gamma}{V}} \left( \frac{\Gamma}{V} \right) + \frac{9}{8} \frac{\left( \Gamma/V + 2H \right)}{\Gamma/V + 3H} \right] \left[ \frac{\Gamma + 4HV - \frac{\Gamma,\phi(\ln(V),\phi)}{12H(3H + \Gamma/V)}}{V} \right] d\phi \tag{31} \]
and \( C \) is a constant of integration.

In this way, the expression for the density perturbations becomes (see [17])
\[ \delta H = \frac{2m_p^4}{5} \frac{\exp[-\Im(\phi)]}{\left( \ln(V)\right)_{,\phi}} \delta \phi. \tag{32} \]
We note here that in the absence of the dissipation coefficient, \( \Gamma = 0 \), we see that equation (32) is reduced to \( \delta H \sim V\delta \phi/(H\phi) \sim H\delta \phi/\phi \), which coincides with the expression obtained for cool inflation.

In the case of high dissipation, the dissipation parameter \( \Gamma \) is much greater than the product between the rate expansion \( H \) and the scalar potential, i.e. \( r = \Gamma/3HV \gg 1 \), and equation (32) becomes
\[ \delta_H^2 = \frac{4}{255} \frac{m_p^4 \exp[-2\Im(\phi)]}{H^2r^2} \delta \phi^2, \tag{33} \]
where now \( \Im(\phi) := \Im(\phi)|_{r \gg 1} \) becomes
\[ \Im(\phi) = -\int \left[ \frac{1}{3Hr} \left( \frac{\Gamma}{V} \right),_{\phi} + \frac{9}{8} \left[ 1 - \frac{\left( \ln(V),\phi(\ln(V)),\phi \right)}{36rH^2} \right] \right] \left( \ln(V)\right)_{,\phi} d\phi. \tag{34} \]
In the first approximation, and during the slow roll phase, the relation between the density matter fluctuation, $\delta \rho$, and the metric perturbation, $\Phi$, is given by

$$
\delta \rho \simeq V,_{\phi} \delta \phi \simeq -2 [1 + r] V \left[ 1 + \frac{\Gamma}{4HV} + \frac{\Gamma,_{\phi}\dot{\Phi}}{48H^2V} \right]^{-1} \Phi,
$$

(35)

where we have used equation (23). Note that in the absence of dissipation, i.e. for $\Gamma = 0$, we recover the usual relation $\delta \rho/\rho \simeq -2\Phi$ valid for cool inflation, in which $\rho \simeq V$.

The fluctuations of the tachyon field are generated by thermal interaction with the radiation field, instead of quantum fluctuations. Following reference [18], we may write, in the case $r \gg 1$,

$$
(\delta \phi)^2 \simeq \frac{k_F T_r}{2 m_\phi^2 \pi^2},
$$

(36)

where the wavenumber $k_F$ is defined by $k_F = \sqrt{H/V} = H \sqrt{3r} \geq H$ and corresponds to the freeze-out scale at which dissipation damps out the thermally excited fluctuations. The freeze-out wavenumber $k_F$ is defined at the point where the inequality $(\ln(V))_{,\phi} < \Gamma H/V$, is satisfied [18].

Combining equations (32), (33) and (36) we find

$$
\delta_H^2 \approx \frac{\sqrt{3}}{75 \pi^2} \exp[-2\tilde{\Theta}(\phi)] T_r \frac{r^{1/2}}{\tilde{\varepsilon}}
$$

(37)

or equivalently by using equations (9) and (10) we obtain

$$
\delta_H^2 \approx \frac{\sqrt{3}}{15 \pi^2} \exp[-2\tilde{\Theta}(\phi)] \left[ \left( \frac{1}{\tilde{\varepsilon}} \right)^3 \frac{1}{2r^2 \sigma k^2 V} \right]^{1/4},
$$

(38)

where the dimensionless slow roll parameter in the high dissipation period becomes

$$
\tilde{\varepsilon} \approx \frac{1}{6\kappa r} \left[ \frac{V,_{\phi}}{V} \right]^2 \frac{1}{V}.
$$

(39)

The scalar spectral index $n_s$ is given by

$$
n_s - 1 = \frac{d \ln \delta_H^2}{d \ln k},
$$

(40)

where the interval in wavenumber is related to the number of e-folds by the relation $d \ln k(\phi) = -dN(\phi)$, and using equations (38) and (40), it becomes

$$
n_s \approx 1 - \left[ \frac{3\tilde{\eta}}{2} + \tilde{\varepsilon} \left( \frac{2V}{V,_{\phi}} \left[ 2\tilde{\Theta},_{\phi} - \frac{r,_{\phi}}{4r} \right] - \frac{5}{2} \right) \right],
$$

(41)

where the second slow roll parameter $\eta$ (for $r \gg 1$) is given by

$$
\tilde{\eta} \approx \frac{1}{3\kappa r V} \left[ \frac{V,_{\phi} V}{V} - \frac{1}{2} \left( \frac{V,_{\phi} V}{V} \right)^2 \right].
$$

(42)

One of feature of the three-year data set from WMAP is that it hints at a significant running in the scalar spectral index $dn_s/d\ln k = \alpha_s$ [4]. Dissipative effects themselves can
produce a rich variety of spectra ranging between red and blue [7, 15]. From equation (41) we obtain that the running of the scalar spectral index for our model becomes

\[ \alpha_s = \frac{\mathrm{d}n_s}{\mathrm{d} \ln k} \approx \frac{2V}{V_\phi} \left[ \frac{3\tilde{\eta}_{,\phi}}{2} + \frac{\tilde{\varepsilon}_{,\phi}}{\tilde{\varepsilon}} \left( n_s - 1 + \frac{3\tilde{\eta}}{2} \right) + 2\tilde{\varepsilon} \left( \left( \frac{V}{V_\phi} \right)_{,\phi} \right) \tilde{\varepsilon}_{,\phi} - \frac{\left( \ln r \right)_{,\phi}}{4} \right] \]

(43)

In models with only scalar fluctuations, the marginalized value for the derivative of the spectral index is approximated as \( \mathrm{d}n_s/\mathrm{d} \ln k = \alpha_s \approx -0.05 \) for WMAP-3 only [4].

In the next section we will study the specific cases in which the dissipation parameter, \( \Gamma \), is either \( \Gamma = \text{const} \) or \( \Gamma = \Gamma(\phi) \).

### 3.2. Tensor perturbations

As it is mentioned in [19], the generation of tensor perturbation during inflation would produce stimulated emission in the thermal background of the gravitational wave. This process changes the power spectrum of the tensor modes through an extra temperature dependence factor \( \coth(k/2T) \). The spectrum in this case is given by

\[ A_{g}^2 = \frac{16\pi}{m_p^2} \left( \frac{H}{2\pi} \right)^2 \coth \left[ \frac{k}{2T} \right] \approx \frac{32V}{3m_p^4} \coth \left[ \frac{k}{2T} \right], \]

(44)

where the spectral index \( n_g \) results as

\[ n_g = \frac{\mathrm{d}}{\mathrm{d} \ln k} \ln \left[ \frac{A_{g}^2}{\coth[k/2T]} \right] = -2\varepsilon. \]

(45)

Here, we have used that \( A_{g}^2 \propto k^{n_g} \coth[k/2T] \) and equation (11) [19].

From expressions (37) and (44), we may write the tensor–scalar ratio as

\[ R(k_0) = \left. \left( \frac{A_{g}^2}{P_R} \right) \right|_{k=k_0} = \frac{240\sqrt{3}}{25m_p^2} \left[ \frac{(1 + r\varepsilon)H^3}{r^{1/2}T_r} \exp[2\tilde{\varepsilon}(\phi)] \coth \left( \frac{k}{2T} \right) \right] \]

where we have used that \( \delta_H \equiv 2P_{R,2}^{1/2}/5. \) For \( r \gg 1 \), this relation reduces to

\[ R(k_0) \approx \frac{240\sqrt{3}}{25m_p^2} \left[ \frac{r^{1/2}\varepsilon H^3}{T_r} \exp[2\tilde{\varepsilon}(\phi)] \coth \left( \frac{k}{2T} \right) \right] \]

(46)

where \( k_0 \) is referred to as the pivot point.

From the combination of WMAP three-year data [4] with the SDSS large scale structure surveys [20], there is found an upper bound \( R(k_0) = 0.002 \text{ Mpc}^{-1} < 0.28 \) (95% CL), where \( k_0 = 0.002 \text{ Mpc}^{-1} \) corresponds to \( l = \tau_0k_0 \approx 30 \), with the distance to the decoupling surface \( \tau_0 = 14400 \text{ Mpc} \). SDSS measures galaxy distributions at red-shifts \( a \approx 0.1 \) and probes \( k \) in the range \( 0.016h \text{ Mpc}^{-1} < k < 0.011h \text{ Mpc}^{-1} \). The recent WMAP three-year results give the values for the scalar curvature spectrum \( P_R(k_0) \approx 25\delta_2^2(k_0)/4 \approx 2.3 \times 10^{-9} \) and the scalar–tensor ratio \( R(k_0) = 0.095 \). These values allow us to find constraints on the parameters of our model.
4. Exponential potential in the high dissipation approach

The tachyonic effective potential \( V(\phi) \) is one that satisfies \( V(\phi) \rightarrow 0 \) as \( \phi \rightarrow \infty \). It has been argued that the qualitative tachyonic potential of string theory can be described via an exponential potential of the form [11]
\[
V(\phi) = V_0 e^{-\alpha \phi}, \tag{47}
\]
where \( \alpha \) and \( V_0 \) are free parameters. In the following we will take \( \alpha > 0 \) (with unit \( m_p \)). The parameter \( \alpha \) is related to the tachyon mass [21]. An estimation of these parameters is given for the cool inflationary case in [12], where \( V_0 \sim 10^{-10} m_p^4 \) and \( \alpha \sim 10^{-6} m_p \).

In the following we will restrict ourselves to the high dissipation regimen in which \( r \gg 1 \).

4.1. \( \Gamma = \Gamma_0 = \text{const case} \)

With \( \Gamma = \Gamma_0 = \text{const.} \) and using the exponential potential given by equation (47), we find that the slow roll parameters become
\[
\tilde{\varepsilon} = \tilde{\eta} = \frac{1}{6\kappa r V_0 e^{-\alpha \phi}}. \tag{48}
\]
The Hubble parameter is given by
\[
H(\phi) = \sqrt{\kappa V_0 e^{-\alpha \phi}/2}, \tag{49}
\]
while the rate \( r \) becomes
\[
r = \frac{m_p \Gamma_0}{\sqrt{24\pi}} \frac{1}{V_0^{3/2}} e^{3\alpha \phi/2} \gg 1. \tag{50}
\]
The evolution of the \( \dot{\phi} \) during this scenario is governed by the expression
\[
\dot{\phi} = -\frac{V_{,\phi}}{3\eta H V}, \tag{51}
\]
and using equations (49)–(51), we find that the evolution of the tachyonic field as a function of time is given by
\[
\phi(t) = \frac{1}{\alpha} \ln \left[ e^{\alpha \phi_i} + \frac{\alpha^2 V_0}{\Gamma_0} t \right], \tag{52}
\]
where \( \phi(t = t_i = 0) = \phi_i \). Substituting this last equation into equation (49), we obtained for the scale factor
\[
a(t) = \exp \left[ \frac{2\Gamma_0}{\alpha^2} \sqrt{\kappa / V_0} \left[ e^{\alpha \phi_i} + \frac{\alpha^2 V_0}{\Gamma_0} t - e^{\alpha \phi_i/2} \right] \right]. \tag{53}
\]
It is not hard to see that \( \dot{a}(t = t_c) = 0 \) (or equivalently \( \epsilon_f \approx 1 \)), where we get that
\[
t_f = \frac{\alpha^2}{4\kappa \Gamma_0} - \frac{\Gamma_0 e^{\alpha \phi_i}}{\alpha^2 V_0}. \tag{54}
\]
Thus, \( t_f \) represents the time at which inflation ends. Substituting equation (79) into equation (52), we find that the tachyonic field at the end of inflation becomes
\[
\phi_f = -\frac{1}{\alpha} \ln \left[ \frac{4\kappa \Gamma_0^2}{\alpha^4 V_0} \right]. \tag{55}
\]
The energy density of the radiation field becomes
\[
\rho_{\gamma} = \frac{\sqrt{6\pi}}{6\kappa} \left[ \frac{\alpha^2}{\Gamma_0} \right] V_0^{3/2} e^{-3\phi/2}
\]  
and, in terms of \( \rho_{\phi} \), it is given by
\[
\rho_{\gamma} = \frac{\sqrt{6\pi}}{6\kappa} \left[ \frac{\alpha^2}{\Gamma_0} \right] \rho_{\phi}^{3/2}.
\]  
Using equation (16), the total number of e-folds at the end of warm inflation results as
\[
N_{\text{total}} = -3\kappa \int_{\phi_i}^{\phi_f} V_r \, d\tilde{\phi} = \frac{2\Gamma_0}{\alpha^2} \left[ \frac{\kappa}{V_0} \right]^{1/2} \left[ \exp(\alpha\phi_f/2) - \exp(\alpha\phi_i/2) \right],
\]  
where the initial tachyonic field satisfies \( \phi_i < \phi_f \), since \( V_i > V_f \).

Rewriting the total number of e-folds in terms of \( V_f \) and \( V_i \) and using equation (55) we find
\[
N_{\text{total}} = \left[ 1 - \left( \frac{V_i}{V_f} \right)^{1/2} \right].
\]  
Note that \( V_i > V_f \) and thus we obtain \( N_{\text{total}} \lesssim 1 \) e-fold. This problem it is not surprising since this lies in the slow roll approximation which by virtue of equation (48) is valid when \( \phi \ll \ln r \), i.e., when \( \phi \) is not too large for the case in which \( r \gg 1 \). This implies, from equation (52), that this solution is valid only at early times and therefore it is not asymptotic. Therefore, the problem of small e-folds is due to the slow roll approximation. In section 5.1, we will again consider the case \( \Gamma = \text{const} \), but there a power law scale factor is assumed.

4.2. \( \Gamma \) as a function of \( \phi \)

In this case we consider \( \Gamma \) as a function of \( \phi \). We take this to be of the form,
\[
\Gamma = f(\phi) = c^2 V(\phi) = c^2 V_0 e^{-\alpha \phi},
\]  
where \( c^2 > 0 \) (with unit of \( m_p \)). The slow roll parameters are equal, i.e.
\[
\tilde{\varepsilon} = \tilde{\eta} = \frac{\alpha^2 e^\alpha \phi/2}{2c^2 (\kappa V_0)^{1/2}},
\]  
and the dissipation parameter \( r \) is given by
\[
r = \frac{c^2 e^\alpha \phi/2}{3(\kappa V_0)^{1/2}}.
\]  
In this case, we obtain that the evolution of the tachyon field is given by
\[
\phi(t) = \phi_i + \frac{\alpha}{c^2} t
\]  
and the scale factor becomes
\[
\frac{a(t)}{a_i} = \exp \left[ \frac{2c^2}{\alpha^2} \sqrt{\kappa V_0 e^{-\alpha \phi_i/2}} \left( 1 - e^{-\alpha^2 t/2c^2} \right) \right].
\]
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The value of the potential at the end of inflation is given by
\[ V_f = V_0 e^{-\alpha \phi} = \frac{\alpha^4}{4 \kappa c^4} \]  
and the total number of e-folds is related to \( V_f \) and \( V_i \) through
\[ V_i = (2 N_{\text{total}} - 1)^2 V_f. \]  
Note that, unlike in the previous case, the \( N_{\text{total}} \) parameter could assume appropriate values (60 or so) for solving the standard cosmological puzzles. To do that, we need the following inequality to be satisfied: \( V_i > 10^4 V_f \).

At the beginning of inflation the dissipation parameter becomes
\[ r(\phi) = r_1 = \frac{2}{3} \frac{c^4}{(2 N_{\text{total}} - 1) \alpha^2} \gg 1, \]  
resulting in the requirement that \( c^2 \gg \alpha (N_{\text{total}} - 1)^{1/2}. \)

On the other hand, from equation (30) we obtain that
\[ \delta \phi \simeq C \frac{\dot{\phi}}{H} \left[ 1 + \frac{c^2}{3H} \left( \frac{5}{4} + (21 \alpha^2 / 16 c^4) \right) \right] \exp \left[ \frac{c^2 (6 H^2 (4 c^4 - 21 \alpha^2) + H c^2 (8 c^4 - 63 \alpha^2) - 6 c^4 \alpha^2)}{32 H (c^2 + 3H)^2} \right]. \]  
From equation (23), the metric perturbation \( \Phi \) becomes given by
\[ \Phi \simeq -C_1 \frac{\dot{H}}{H^2} \left[ 1 + \frac{c^2}{3H} \left( \frac{5}{4} + (21 \alpha^2 / 16 c^4) \right) \right] \left[ 1 + \frac{c^2}{4H} - \frac{\alpha^2 c^2}{4 H^2 (c^2 + 3H)} \right] \exp \left[ \frac{c^2 (6 H^2 (4 c^4 - 21 \alpha^2) + H c^2 (8 c^4 - 63 \alpha^2) - 6 c^4 \alpha^2)}{32 H (c^2 + 3H)^2} \right], \]  
where the zeroth-order term is related the adiabatic mode and the remaining terms describe the inclusion of dissipation in the fluctuation of the scalar field (the entropy mode) [16].

From equation (40) we obtained that the scalar spectral index of the adiabatic perturbation becomes
\[ n_s \approx 1 - \left[ \frac{3 \tilde{\eta}}{2} - \frac{3}{4} \tilde{\xi} (1 + \tilde{\xi}) \right], \]  
which becomes \( r \) dependent for the case of high dissipation.

From equation (37), when \( r \gg 1 \) or equivalently \( c^2 \gg 3H \), we obtain that the scalar power becomes
\[ P_R(k_0) \approx \frac{k^{10/4}}{2 \pi^2} \left[ \frac{c}{\alpha^2} V(\phi_0)^{5/2} T_r \right] \exp \left[ \frac{3 \alpha^2}{8 c^2 \sqrt{\kappa V(\phi_0)}} \right], \]  
and from equation (46) the tensor–scalar ratio is given by
\[ R(k_0) \approx 3 \pi^2 \left( \frac{\alpha^2}{c \kappa^{11/4}} \right) \left[ \frac{1}{V(\phi_0)^{3/2} T_r} \right] \exp \left[ -\frac{3 \alpha^2}{8 c^2 \sqrt{\kappa V(\phi_0)}} \right] \coth \left[ \frac{k_0}{2 T_r} \right]. \]
where \( V(\phi_0) = V_0 \exp[-\alpha \phi_0] \) and \( \phi_0 \) represents the value of the tachyon field when the scale \( k_0 = 0.002 \, \text{Mpc}^{-1} \) was leaving the horizon.

Using the WMAP three-year data where \( P_R(k_0) \approx 2.3 \times 10^{-9} \), \( R(k_0) = 0.095 \) and choosing the parameters \( c^2 = 10^{10} m_p \), \( T \approx T_\text{r} \approx 0.24 \times 10^{16} \, \text{GeV} \) and \( k_0 = 0.002 \, \text{Mpc}^{-1} \), we obtained from equations (70) and (71) that \( V(\phi_0) \approx 2.43 \times 10^{-15} m_p^4 \) and \( \alpha \approx 1.76 \times 10^{-12} m_p \). We noted that the dissipation coefficient when the scale \( k_0 \) was leaving the horizon is the order of \( \Gamma(k_0) \sim 10^{-5} m_p \). From equation (43) we found that it is necessary to increase the value of \( \alpha \) by several orders of magnitude in order to have a running spectral index \( \alpha_s \) close to a value given by WMAP observations.

5. Power law in the high dissipation approach

An interesting approach is to choose a scale factor of power law type, i.e. \( a \sim t^p \), in which \( p \geq 2 \). This allows us to get the form of the scalar potential \( V(\phi) \) for different values of \( \Gamma \) in the high dissipation approach, i.e. \( \tilde{r} = \Gamma/3H \rho_\phi \gg 1 \). Note that if \( \dot{\phi}^2 \ll 1 \), \( \rho_\phi \sim V \) we get \( \tilde{r} \rightarrow r \).

In the following we will consider some special cases for the parameter \( \Gamma \).

5.1. \( \Gamma = \Gamma_0 = \text{const} \) case

By using equations (4) and (5), we obtained that

\[
\dot{\rho}_\phi + \dot{\rho}_\gamma + 4H \rho_\gamma = 0,
\]

where we have used for the energy density and pressure associated with the tachyon field the relation \( p_\phi = -\rho_\phi (1 - \dot{\phi}^2) \). From equations (3) and (72) and choosing a scale factor of power law type, we find that

\[
\rho_\gamma = \frac{p}{2\kappa t^2}, \quad \rho_\phi = \frac{p(2p-1)}{2\kappa t^2} = \rho_\gamma (2p-1).
\]

From equation (4) (or equivalently from equation (5)) the temporal evolution of the tachyonic field becomes given by

\[
(\phi - \phi_1)^2 = \frac{4(2p-1)p}{\kappa \Gamma_0} t,
\]

where \( \phi_1 \) is a constant of integration. Substituting equation (73) into equation (1) and using equation (74), we obtained that the potential as a function of the tachyonic field \( \phi \)

becomes

\[
V(\phi) = \beta_1 (\phi - \phi_1)^4 \sqrt{1 - \beta_2 (\phi - \phi_1)^6},
\]

where \( \beta_1 \) and \( \beta_2 \) are given by

\[
\beta_1 = \frac{\kappa \Gamma_0^2}{32p(2p-1)} \quad \text{and} \quad \beta_2 = \frac{\kappa^2 \Gamma_0^2}{64p^2(2p-1)^2},
\]

respectively.

We note that equation (72) is valid for

\[
\tilde{r} = \frac{\Gamma}{3H \rho_\phi} \gg 1 \quad \Rightarrow \quad \frac{128p^2(2p-1)^2}{3\kappa^2 \Gamma_0^2} \gg (\phi - \phi_1)^6.
\]
choosing the parameters

\[ k \]

the scalar potential when the scale is \( R \)
and the tensor–scalar ratio is given by

\[
\frac{\rho_\phi + \rho_\gamma}{3(\rho_\phi + \rho_\gamma)} = \exp\left(\frac{2}{p}\right)
\]

and from equations (77) and (79), we have that

\[
\phi_1 = \left[\frac{3p(2p-1)^2}{4\kappa T_0(p-1)}\right]^{1/3}
\]

and from equations (77) and (79), we have that

\[
\phi_0 = \left[\frac{4^p p^2(2p-1)(p-1)}{3\kappa^2 T_0^2}\right]^{1/6} \exp\left(\frac{N}{2p}\right).
\]

On the other hand, the spectrum of scalar perturbations becomes

\[
P_R(k_0) \approx \frac{2}{\pi^2} \left[\frac{V(\phi_0)^{23}}{\beta_1^{23/2}}\right]^{1/4} \left(\frac{T_r}{(\phi_0 - \phi_1)^{13} [4 - 7\beta_2 (\phi_0 - \phi_1)^6]^2}\right)
\]

and the tensor–scalar ratio is given by

\[
R(k_0) \approx \frac{512\pi}{m_p^2} \left[\frac{\beta_1^{23/2}}{V(\phi_0)^{23}}\right]^{1/4} \left(\frac{(\phi_0 - \phi_1)^{17} [4 - 7\beta_2 (\phi_0 - \phi_1)^6]^2}{\Gamma_0^2 T_r}\right) \coth\left(\frac{k_0}{2T}\right).
\]

Using the WMAP three-year data, in which \( P_R(k_0) \approx 2.3 \times 10^{-9} \), \( R(k_0) = 0.095 \) and choosing the parameters \( p = 2 \), \( T \approx T_r \approx 0.24 \times 10^{16} \text{ GeV} \) and \( k_0 = 0.002 \text{ Mpc}^{-1} \), we obtain from equations (81) and (82) that \( (\phi_0 - \phi_1) \approx 2 \times 10^3 m_p \) and \( \Gamma_0 \approx 10^{-11} m_p^5 \). Here the scalar potential when the scale is \( k_0 = 0.002 \text{ Mpc}^{-1} \) is leaving the horizon becomes \( V(\phi_0) \sim 3 \times 10^{-10} m_p^4 \). Note that this value of \( V(\phi_0) \) becomes similar to that found for cool inflation where a chaotic potential is used [19]. Another case of interest is that of \( p = 4 \), where it is found that \( (\phi_0 - \phi_1) \approx 10^7 m_p \), \( \Gamma_0 \approx 10^{-22} m_p^5 \) and \( V(\phi_0) \sim 10^{-17} m_p^4 \). In this case, from equation (43) we found that if we decrease the value of \( (\phi - \phi_1) \) by a few orders of magnitude, then the running spectral index \( \alpha_s \) takes a value which agrees with the value specified by the WMAP observations.

5.2. \( \Gamma \) as a function of \( t \)

In this case, for simplicity we consider that the dissipation coefficient is \( \Gamma = \gamma/t \), where \( \gamma \) is a positive-definite constant (with the unit of \( m_p^4 \)). By using equations (4) and (73), we obtain that the tachyonic field becomes

\[
(\phi - \phi_1) = \beta_3 \ln|t|,
\]
where \( \tilde{\phi}_1 \) is a constant of integration and \( \beta_3 = \sqrt{(2p-1)p/\kappa\gamma} \). We also note that equation (72) is valid for

\[
\frac{\dot{r}}{3H\rho_\phi} \gg 1 \implies \frac{2\kappa\gamma}{3p^2(2p-1)} \gg \exp(-2(\phi - \phi_1)/\beta_3).
\]

(84)

The potential is now given by

\[
V(\phi) = \left( \frac{\gamma\beta_3^2}{2} \right) \exp \left[ \frac{-2(\phi - \tilde{\phi}_1)}{\beta_3} \right] \sqrt{1 - \beta_3^2 \exp \left[ \frac{-2(\phi - \tilde{\phi}_1)}{\beta_3} \right]}.
\]

(85)

From equations (78) and (83), we obtain that the time at the end of inflation is given by

\[
t_f = \frac{(2p-1)}{2} \sqrt{\frac{3p}{\kappa\gamma(p-1)}}.
\]

(86)

The tachyonic field in terms of the \( N \) e-folds parameter becomes

\[
(\phi - \tilde{\phi}_1) = \beta_3 \left[ \ln \left( \frac{(2p-1)}{2} \sqrt{\frac{3p}{\kappa\gamma(p-1)}} \right) - \frac{N}{p} \right].
\]

(87)

Following a scheme similar to those of the previous cases, in order to calculate the spectrum of scalar perturbations and the tensor–scalar ratio, we obtained that \( (\phi_0 - \tilde{\phi}_1) \approx 69.6 m_p \) and \( \gamma \approx 0.02 m_p^4 \), and the potential when the scale \( k_0 = 0.002 \) Mpc\(^{-1} \) is leaving the horizon becomes \( V(\phi_0) \approx 0.2 \times 10^{-14} m_p^4 \). We have chosen the values \( p = 2 \) and \( T \approx T_r \approx 0.24 \times 10^{16} \) GeV. For the values of the parameters given above, and from equation (43), we find that the running spectral index \( \alpha_s \) lies in the range established by WMAP observations.

6. Conclusions

In this paper we have investigated the tachyonic warm inflationary scenario. In the slow roll approximation we have found a general relationship between the radiation and tachyon energy densities. This has led us to a general criterion for tachyon warm inflation to occur; see equation (14).

In relation to the corresponding perturbations, the contributions of the adiabatic and entropy modes were obtained explicitly. Here, it is shown that the dissipation parameter plays a crucial role in producing the entropy mode. Additionally, we have found a general relation for the density perturbation expressed by equation (32). The tensor perturbation are generated via stimulated emission into the existing thermal background (see equation (44)) and the tensor–scalar ratio is modified by a temperature dependent factor.

Our first specific models are described using an exponential scalar potential and for two different dependences of the dissipation coefficient, \( \Gamma \). In the first case, we took \( \Gamma = \Gamma_0 = \text{Cte} \) and we have found that the number of e-folds in this case is insufficient. For the case in which the dissipation coefficient \( \Gamma \) is taken to be a function of the tachyon field, i.e. \( \Gamma = f(\phi) = c^2 V(\phi) \), it has been possible to describe appropriately tachyonic warm inflationary universe models. Also, in place of choosing a specific form of the scalar
tachyon potential \( a \text{ priori} \), we have assumed a power law for the scale factor. Here, we have described inflationary models for \( \Gamma = \Gamma_0 = \text{const} \) and for \( \Gamma \) variable. In both cases, we have obtained explicit expressions for the corresponding scalar potential. By using the WMAP three-year data, we have found some constraints for the parameters appearing in our model. From the normalization of the WMAP data, the potential becomes of the order of \( V(\phi_0) \sim 10^{-15} m_p^4 \) when it leaves the horizon, at the scale of \( k_0 = 0.002 \text{ Mpc}^{-1} \).

Dissipative effects play a crucial role in producing the entropy mode; they can themselves produce a rich variety of spectra ranging between red and blue. The possibility of a spectrum which runs from blue to red is particularly interesting because it is not commonly seen in inflationary models, which typically predict red spectra. Models of inflation with dissipative effects and models with interacting fields have much more freedom to yield spectra. In particular, we anticipate that the Planck mission will significantly enhance our understanding of \( \alpha_s \) by providing high quality measurements of the fundamental power spectrum over a larger wavelength range than WMAP. Summarizing, we have been successful in describing tachyon warm inflationary models for describing the early epoch of the universe in the high dissipation regime.

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References

[1] Guth A, The inflationary universe: A possible solution to the horizon and flatness problems, 1981 Phys. Rev. D 23 347 [SPIRES]
[2] Albrecht A and Steinhardt P J, Cosmology for grand unified theories with radiatively induced symmetry breaking, 1982 Phys. Rev. Lett. 48 1220 [SPIRES]

A complete description of inflationary scenarios can be found in the book by Linde A, 1990 Particle Physics and Inflationary Cosmology (New York: Gordon and Breach)
[3] Spergel D N, First year Wilkinson Microwave Anisotropy Probe (WMAP) observations: Determination of cosmological parameters, 2003 Astrophys. J. Suppl. 148 175

Peiris H V, First year Wilkinson Microwave Anisotropy Probe (WMAP) observations: Implications for inflation, 2003 Astrophys. J. Suppl. 148 213
[4] Page L et al, Wilkinson Microwave Anisotropy Probe (WMAP) three year results: implications for cosmology, 2006 Preprint astro-ph/0603449
[5] Linde A, A new inflationary universe scenario: a possible solution of the horizon, flatness, homogeneity, isotropy and primordial monopole problems, 1982 Phys. Lett. B 108 389 [SPIRES]

Albrecht A and Steinhardt P J, Cosmology for grand unified theories with radiatively induced symmetry breaking, 1982 Phys. Rev. Lett. 48 1220 [SPIRES]

Hawking S W and Moss I G, Supercooled phase transitions in the very early universe, 1982 Phys. Lett. B 110 35 [SPIRES]
[6] Berera A, Warm inflation, 1995 Phys. Rev. Lett. 75 3218 [SPIRES]

Berera A, Interpolating the stage of exponential expansion in the early universe: A possible alternative with no reheating, 1997 Phys. Rev. D 55 3346 [SPIRES]
[7] Hall L M H, Moss I G and Berera A, Scalar perturbation spectra from warm inflation, 2004 Phys. Rev. D 69 083525 [SPIRES]

Moss I G, Primordial inflation with spontaneous symmetry breaking, 1985 Phys. Lett. B 154 120 [SPIRES]
Tachyon warm inflationary universe models

Berera A and Fang L Z, Thermally induced density perturbations in the inflationary era, 1995 Phys. Rev. Lett. 74 1912 [SPIRES]

Berera A, Warm inflation at arbitrary adiabaticity: A model, an existence proof for inflationary dynamics in quantum field theory, 2000 Nucl. Phys. B 585 666 [SPIRES]

Berera A, Thermal properties of an inflationary universe, 1996 Phys. Rev. D 54 2519 [SPIRES]

Taylor A N and Berera A, Perturbation spectra in the warm inflationary scenario, 1997 Phys. Rev. D 55 3346 [SPIRES]

Sen A, Rolling tachyon, 2002 J. High Energy Phys. JHEP04(2002)048 [SPIRES]

Sen A, Field theory of tachyon matter, 2002 Mod. Phys. Lett. A 17 1797 [SPIRES]

Sami M, Chingangbam P and Qureshi T, Aspects of tachyonic inflation with exponential potential, 2002 Phys. Rev. D 66 043530 [SPIRES]

Gibbons G W, Cosmological evolution of the rolling tachyon, 2002 Phys. Lett. B 537 1 [SPIRES]

Bardeen J, Gauge invariant cosmological perturbations, 1980 Phys. Rev. D 22 1882 [SPIRES]

Starobinsky A and Yokoyama J, Density fluctuations in Brans–Dicke inflation, 1995 Proc. 4th Workshop on General Relativity and Gravitation ed K Nakao et al (Kyoto: Kyoto University) p 381 [gr-qc/9502002]

Starobinsky A and Tsujikawa S, Cosmological perturbations from multifield inflation in generalized Einstein theories, 2001 Nucl. Phys. B 610 383 [SPIRES]

Oliveira H and Joras S, On perturbations in warm inflation, 2001 Phys. Rev. D 64 063513 [SPIRES]

Oliveira H, Density perturbations in warm inflation and COBE normalization, 2002 Phys. Lett. B 526 1 [SPIRES]

Liddle A and Lyth D, 2000 Cosmological Inflation and Large-Scale Structure (Cambridge: Cambridge University Press)

Linsey J, Liddle A, Kolb E and Copeland E, 1997 Rev. Mod. Phys. 69 373 [SPIRES]

Bassett B, Tsujikawa S and Wands D, Inflation dynamics and reheating, 2005 Rev. Mod. Phys. 78 537 [SPIRES]

Taylor A and Berera A, Perturbations spectral in the warm inflationary scenario, 2000 Phys. Rev. D 62 083517 [SPIRES]

Bhattacharya K, Mohanty S and Nautiyal A, Enhanced polarization of CMB from thermal gravitational waves, 2006 Preprint astro-ph/0607049

Tegmark M et al, Cosmological parameters from SDSS and WMAP, 2004 Phys. Rev. D 69 103501 [SPIRES]

Fairbairn M and Tytgat M, Inflation from a tachyon fluid?, 2002 Phys. Lett. B 546 1 [SPIRES]