Fresh inflation: a warm inflationary model from a zero temperature initial state

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Abstract

A two-component mixture fluid which complies with the gamma law is considered in the framework of inflation with finite temperature. The model is developed for a quartic scalar potential without symmetry breaking. The radiation energy density is assumed to be zero when inflation starts and remains below the GUT temperature during the inflationary stage. Furthermore, provides the necessary number of e-folds and sufficient radiation energy density to GUT baryogenesis can take place near the minimum energetic configuration.

Pacs: 98.80.Cq

I. INTRODUCTION AND MOTIVATION

Since its creation, the inflationary Universe scenario has become an integral part of the standard model of cosmology. An ideal inflationary scenario should arise naturally from quantum cosmology without fine tuning. Inflation is needed because it solves the horizon, flatness, and monopole problems of the very early universe and also provides a mechanism for the creation of primordial density fluctuations. For inflation, naturality has played an important role. This is understandable since for phenomena that cannot be directly observed, one attempts a description starting with the most natural expectations. The importance of naturality principles is to provide guidance from more familiar analogies with the hope of gaining predictability. For inflation we can understand naturality as both macroscopic and microscopic. Macroscopically, we would like a description that rests with common-day experience. Microscopically, it should be consistent with the standard model of particle physics.

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Quantum fluctuations of matter field \(\xi\) and thermal fluctuations \(\xi\) play a prominent role in inflationary cosmology. They lead to density perturbations that would be responsible for the origin of structures in the present-day universe \(\xi\). Structure formation scenarios can receive important restrictions based on the measured background anisotropy temperature \(\Delta \theta/\theta = 1.1 \times 10^{-5}\).

A successful model of inflation must satisfy the following requirements.

a) An unstable primordial matter-field perturbation with an energy density nearly \(M^4_p\) must lead to inflation.  
b) The growing of the scale factor during inflation must grow at least by 60 e-foldings, to solve the horizon problem to give a sufficiently globally flat universe.  
c) Inflation must generate density perturbations of temperature with amplitude constrained by the Cosmic Microwave Background (CMB) density fluctuations.  
d) During all the inflationary era, the temperature must remain below the GUT critical temperature \(\theta_{\text{GUT}} \approx 10^{15}\) GeV, to avoid monopole and domain walls proliferation.  
e) At the end of inflation the universe must be with sufficient radiation energy density to make possible the particles creation in the framework of GUT baryogenesis.

The standard slow-roll inflation model separates expansion and reheating into two distinguished time periods. It is first assumed that exponential expansion from inflation places the universe in a supercooled second order phase transition. Subsequently thereafter the universe is reheated. Two outcomes arise from such a scenario. First, the required density perturbations in this cold universe are left to be created by the quantum fluctuations of the inflaton. Second, the scalar field oscillates near the minimum of its effective potential and produces elementary particles. These particles interact with each other and eventually they come to a state of thermal equilibrium at some temperature \(\theta\). This process completes when all the energy of the classical scalar field transfers to the thermal energy of elementary particles. The temperature of the universe at this stage is called the reheating temperature \(\theta\).

The issue of reheating after inflation in these theories, where the universe is revived from the frozen vacuum state, has remained relatively unexplored. This is precisely one area where the full symmetry and particle content of the underlying theory is likely to be crucial \(\xi\). For GUT baryogenesis the large quantum fluctuations may cause the monopole and domain wall problems due to nonthermal symmetry restoration \(\xi\). Hence, if all the particles are created by means of this mechanism at the end of inflation, the universe could come very inhomogeneous. Furthermore, reheating requires globally coherent radiation waves on the scale of the inflated universe. A globally coherent heating process requires a large-scale radiator, which in standard inflation scenario is the inflaton. The question here is how the random inflaton field configuration before and during inflation attains quantum coherence at the end of inflation.

On the other hand the warm inflation scenario takes into account separately, the matter and radiation energy fluctuations. In this scenario the fluctuations of the matter field lead to perturbations of matter and radiation energy densities \(\xi\), which are responsible for the fluctuations of temperature. In this scenario the matter field interacts with particles which are in a thermal bath with a mean temperature smaller than the GUT critical temperature. This scenario was introduced by Berera \(\xi\). Other approaches for this scenario also were developed \(\xi\). In principle, a permanent or temporary coupling of the scalar field \(\phi\) with other fields might also lead to dissipative processes producing entropy at different eras of
the inflationary stage. The warm inflation scenario appears to be very promising, but the problem with it is that, initially, the thermal bath is unjustified introduced in the framework of chaotic initial conditions needed to give a natural beining to the universe. In this sense, chaotic inflation \[15\] provides a more natural scenario to describe the initial conditions in the universe. A successful model of inflation for an inflaton that interacts with particles in a thermal bath initially must be with \(\rho_r(t = t_o) = 0\) because the universe is, in principle, created from a quantum fluctuation of the vacuum. In this framework, the temperature of the universe must be increasing during inflation to baryogenesis can take place.

The aim of this work is to propose a model of inflation that incorporates some characteristics of both, standard (chaotic) and warm inflationary scenarios to give a natural inflationary scenario in the framework of chaotic initial conditions.

II. GENERAL CONSIDERATIONS

I consider an homogeneous and isotropic universe, described by a flat Friedmann - Robertson - Walker (FRW) metric

\[ ds^2 = -dt^2 + a^2(t)dr^2, \]

where \(a(t)\) is the scale factor of the universe. The early universe can be represented by the Lagrangian density

\[ \mathcal{L} = \frac{1}{2} \partial^\mu \phi \partial_\mu \phi - V(\phi) + \mathcal{L}_{\text{int}}, \]

where \(\mathcal{L}_{\text{int}} \sim -g^2 \phi^2 \Psi^2\) describes the interaction between the scalar field \(\phi\) and the other \(\Psi\)-scalar fields of a thermal bath. The minimally coupled scalar field and the fields of a thermal bath interchange energy during the rapid expansion of the universe. The various outcomes are a result of specially chosen Lagrangians. In the most of cases the Lagrangian is unmotived from particle phenomenology. Clear exceptions are the Coleman - Weinberg potential with coupling constant which arise from GUT theories and SUSY potentials \[16\].

The Einstein’s equations for this system are

\[ 3H^2 = 8\pi G \left[ \frac{\dot{\phi}^2}{2} + V(\phi) + \rho_r \right], \]

\[ 3H^2 + 2\dot{H} = -8\pi G \left[ \frac{\dot{\phi}^2}{2} - V(\phi) + \frac{\rho_r}{3} \right], \]

where \(H = \frac{\dot{a}}{a}\) is the Hubble parameter, \(G = M_p^{-2}\) is the gravitational constant, \(M_p = 1.2 \times 10^{19}\) GeV is the Planckian mass and \(\rho_r\) is the radiation energy density. The equation of motion for \(\phi\) in an interacting system is

\[ \ddot{\phi} + 3H \dot{\phi} + V'(\phi) = -\delta, \]

where \(\delta = \dot{\rho}_r + 3\gamma H \rho_r\) describes the interaction between both subsystems; the inflaton field and the bath. Furthermore, \(\rho_r\) is the radiation energy density. The eq. \([5]\) can be written as two equations.
\[ \ddot{\phi} + 3H \dot{\phi} + V' (\phi) + \frac{\delta}{\dot{\phi}} = 0, \]  
(6)

\[ \dot{\rho}_r + 3\gamma H \rho_r - \delta = 0, \]  
(7)

where \( \gamma \) takes the value \( 4/3 \) in warm inflationary models. Some phenomenological interaction terms as \( \delta \propto \dot{\phi}^2 \) \(^{13} \), \( \delta \propto \dot{\phi}^2 \phi^2 \) \(^{14} \) or \( \delta \propto \dot{\phi}^d \phi^{5-2d} \) \(^{17} \), has been proposed in the literature. Very much work has been made in this framework for a de Sitter expansion of the universe \(^{18} \). However, more interesting and complicated is to work in a quasi - de Sitter expanding universe \(^{12} \), where the fluctuations of the matter field are considered. One of the more interesting consequences that arise from these fluctuations are the metric fluctuations \(^{19} \) on a FRW background metric.

Slow-roll conditions must be imposed to assure nearly de Sitter solutions for an amount of time, which must be long enough to solve the problems of the hot big bang. If \( p_t = \frac{\dot{\phi}^2}{2} + \rho_r - V (\phi) \) is the total pressure and \( \rho_t = \rho_r + \frac{\dot{\phi}^2}{2} + V (\phi) \) is the total energy density, hence the parameter \( F = \frac{\rho_t + p_t}{\rho_r} \) which describes slow roll conditions is \(^{20} \)

\[ F = -2 \frac{H}{3} = \frac{\dot{\phi}^2 + 4 \rho_r}{\rho_r + \frac{\dot{\phi}^2}{2} + V}. \]  
(8)

The requirement to assure slow-roll conditions is that \( F \ll 1 \) and, from eq. (8) one obtains the following equations

\[ \dot{\phi}^2 \left( 1 - \frac{F}{2} \right) + \rho_r \left( \frac{4}{3} - F \right) - F V (\phi) = 0, \]  
(9)

\[ H = \frac{2}{3} \int F dt. \]  
(10)

Furthermore, due to \( \dot{H} = H' \dot{\phi} \), from the first equality in (8) we obtain the time dependence for \( \phi \)

\[ \dot{\phi} = -\frac{3H^2}{2H'} F, \]  
(11)

such that from eqs. (8) and (11) the radiation energy density can be written as a function of \( V (\phi) \) and the Hubble parameter

\[ \rho_r = \left( \frac{3F}{4 - 3F} \right) V - \frac{27}{8} \left( \frac{H^2}{H'} \right)^2 \frac{F^2 (2 - F)}{(4 - 3F)}. \]  
(12)

Furthermore, replacing (11) and (12) in eq. (8), one obtains the scalar field potential as a function of \( H \), for a given parameter \( F \)

\[ V (\phi) = \frac{3}{8\pi G} \left[ \left( \frac{4 - 3F}{4} \right) H^2 + \frac{3\pi G}{2} F^2 \left( \frac{H^2}{H'} \right)^2 \right]. \]  
(13)

If \( F \ll 1 \) is a constant, the evolution for the Hubble parameter \( H = \frac{\dot{a}}{a} \) being given from eq. (10)
\[ H(t) = \frac{2}{3F} t^{-1}, \]  
\[ \text{such that} \]
\[ a(t) \sim t^{\frac{1}{3}}. \]

Finally, the number of e-folds \( N = \int_{t_s}^{t_e} H dt \) (\( t_s \) and \( t_e \) are the time when inflation start and ends), is given by

\[ N = \frac{2}{3F} \ln\left( \frac{t_e}{t_s} \right). \]

With Planckian unities (\( G^{-1/2} \equiv M_p = 1 \)) inflation starts when \( t_s = G^{1/2} = 1 \). Hence, for \( t_e \approx 10^{10} G^{1/2} \), one obtains \( N > 60 \) for \( F < 0.55 \).

### III. THE MODEL

The idea of fresh inflation is that the universe is reheating during inflation and the temperature always remains below \( \theta_{GUT} \). This point is very important to avoid the magnetic monopole restoration. If when inflation starts the inflaton field is not thermalized, the radiation component of energy density must be zero at this moment, so that \( \rho_r(t_s) = 0 \). For a model with finite temperature, the effective potential must be dependent of the temperature \( \theta = K T \) (here \( K \) is the Boltzmann constant).

#### A. The effective potential

I will consider a nongauge theory, invariant under a global group \( O(n) \), involving a single \( n \)-vector multiplet of scalar fields \( \phi_i \). At zero temperature it is

\[ V(\phi_i) = \frac{\mathcal{M}_o^2}{2} \phi_i \phi_i + \frac{\lambda^2}{4} (\phi_i \phi_i)^2, \]  
\[ \text{where} \quad \mathcal{M}_o^2 > 0 \quad \text{and} \quad \lambda^2 > 0. \quad \text{Note that I am considering the model without symmetry breaking. With the notation} \quad (\phi_i \phi_i)^{1/2} \equiv \phi, \quad \text{the effective finite temperature potential} \]

\[ V_{eff}(\phi, \theta) = V(\phi) + \rho_r(\phi, \theta) \]  
\[ \text{can be written as} \]

\[ V_{eff}(\phi, \theta) = \frac{\mathcal{M}^2(\theta)}{2} \phi^2 + \frac{\lambda^2}{4} \phi^4, \]  
\[ \text{where} \]

\[ \mathcal{M}^2(\theta) = \mathcal{M}^2(0) + \frac{(n+2)}{12} \lambda^2 \theta^2. \]  

Here, \( V_{eff}(\phi, \theta) \) is invariant under \( \phi \to -\phi \) reflexions and \( n \) is the number of created particles due to the interaction of \( \phi \) with the fields of the bath. Furthermore, \( \mathcal{M}^2(0) \) is given by \( \mathcal{M}_o^2 \) plus counterterms [25].
From eq. (13) one obtains the Hubble parameter as a function of $\phi$

$$H(\phi) = 4 \sqrt{\frac{\pi G}{3(4 - 3F)}} M(0) \phi, \quad (20)$$

and the condition of consistence implies that

$$\lambda^2 = \frac{12\pi G}{(4 - 3F)} M^2(0). \quad (21)$$

Since $V_{eff}(\phi, \theta) = V(\phi) + \rho_r$ and $\rho_r = \frac{\pi^2}{30} g_{eff} \theta^4$, from eqs. (18) and (19) one obtains the equality

$$\frac{(n + 2)}{12} \lambda^2 \theta^2 \phi^2 = \frac{\pi^2}{30} g_{eff} \theta^4, \quad (22)$$

such that the number of created particles during inflation will be

$$(n + 2) = \frac{2\pi^2}{5\lambda^2 g_{eff} \theta^2 \phi^2}, \quad (23)$$

where $g_{eff}$ is the effective degrees of freedom for the particles. I consider a Yukawa interaction like [13]

$$\delta(\dot{\phi}, \phi, \theta) = \Gamma(\theta) \dot{\phi}^2, \quad (24)$$

where [23,24]

$$\Gamma(\theta) = \frac{g_{eff}^4}{192\pi} \theta, \quad (25)$$

is the decay width of the particles. Furthermore, the thermal bath is at temperature $\theta \sim \rho_r^{1/4}$ assuming, that $\Psi$ has no self-interaction. The most important difference between standard and warm inflation is that the slope of inflationary does not need to be small, as dissipation will allow slow - roll. From eqs. (14) and (20) we can find the temporal evolution for $\phi$

$$\phi(t) = \lambda^{-1} t^{-1}, \quad (26)$$

such that, if inflation starts at $t_s = G^{1/2}$ and ends at $t_e = 10^{10} G^{1/2}$, hence the respective values for the field will be $\phi_s \simeq 10^7 G^{-1/2}$ and $\phi_e \simeq 10^{-3} G^{-1/2}$, for $\lambda^2 = 10^{-14}$. Replacing (25) in (24), and using the eqs. (26), (14) and (5), we find the temperature as a function of time

$$\theta(t) = \frac{192\pi}{g_{eff}^4 \lambda^2} \left\{ M^2(0) \lambda^2 t + t^{-1} \left[ \lambda^2 \frac{(9F^2 - 18F + 8)}{(4 - 3F)^2} + M^2(0) \pi G \frac{(192F^2 - 72F^3 - 96)}{(4 - 3F)^2} \right] \right\}, \quad (27)$$

which increases linearly with time for $t \gg 1$. If we take $\theta_e = 0.1 \theta_c \simeq 10^{-4} G^{-1/2}$, the number of created particles at the end of inflation is of the order of (I am taking $g_{eff} = 10^2$)

$$n_e \simeq 4 \times 10^8. \quad (28)$$
B. Dynamics for the inflaton field and density fluctuations

In this section I will study the dynamics for the inflaton field to make an estimation for the energy density fluctuations in a globally flat Friedmann - Robertson - Walker (FRW) metric

$$ds^2 = -dt^2 + a^2(t)dx^2.$$  (29)

The dynamics for the spatially homogeneous inflaton field is given by

$$\ddot{\phi} + (3H + \Gamma) \dot{\phi} + V'(\phi) = 0,$$  (30)

where $V'(\phi) \equiv \frac{dV}{d\phi}$. The term $\Gamma \dot{\phi}$ is added in the scalar field equation of motion (30) to describe the continuous energy transferred from $\phi$ to the radiation field. This persistent thermal contact during warm inflation is so finely adjusted that the scalar field evolves always in a damped regime.

Furthermore, the fluctuations of the field $\delta \phi(\vec{x}, t) \equiv \psi(\vec{x}, t)$ are described by the equation of motion

$$\ddot{\psi} - \frac{1}{a^2} \nabla^2 \psi + (3H + \Gamma) \dot{\psi} + V''(\phi)\psi = 0.$$  (31)

Here, the additional second term appears because the fluctuations $\psi$ are spatially inhomogeneous. The eq. (31) can be mapped in a Klein - Gordon one, making $\chi = \psi e^{3/2 \int (H + \Gamma) dt}$

$$\ddot{\chi} - \frac{1}{a^2} \nabla^2 \chi - \frac{k_o^2}{a^2} \chi = 0,$$  (32)

where $k_o$ is the wavenumber that separates the ultraviolet and infrared sectors, such that

$$k_o^2 = a^2 \left[ \frac{9}{4} \left( H + \frac{\Gamma}{3} \right)^2 + 3 \left( \dot{H} + \frac{\dot{\Gamma}}{3} \right) - V''[\phi(t)] \right].$$  (33)

The redefined coarse-grained matter field fluctuations $\chi_{cg} = \psi_{cg} e^{3/2 \int (H + \Gamma/3) dt}$, can be written as a Fourier expansion of the form

$$\chi_{cg} = \frac{1}{(2\pi)^{3/2}} \int d^3k \theta(\epsilon k_o - k) \left[ a_k e^{ik \cdot \vec{x}} \xi_k(t) + a_k^* e^{-ik \cdot \vec{x}} \xi_k^*(t) \right],$$  (34)

where the asterisk denotes the complex conjugate and $\epsilon \ll 1$ is a dimensionless parameter. However, classicality conditions require that the modes for $k < \epsilon k_o$ be real [20]. The requirement to inflation holds is that $k_o^2 > 0$. The infrared sector describes the spectrum with wavelengths much bigger than the size of the horizon. This sector is unstable and it is very important because the spatial inhomogeneities in this sector should be the responsible of the structure formation during the matter dominated era of the universe.

To calculate $k_o^2$ we need to know the temporal evolution of $H(t), V''[\phi(t)]$ and $\Gamma[\theta(t)]$, which are
\[ H[\phi(t)] = \frac{2}{3F} t^{-1}, \]  
\[ V''[\phi(t)] = \mathcal{M}^2(0) + 3t^{-2}, \]  
\[ \Gamma[\theta(t)]|_{t \gg 1} \simeq \mathcal{M}^2(0) t. \]  

Hence, the parameter of mass \( \mu^2(t) = \frac{k^2}{a^2} \) at the end of inflation will be

\[ \mu^2(t)|_{t \gg 1} \simeq \frac{\mathcal{M}^4(0)}{4} t^2 + \left( \frac{1}{F^2} - 3 \right) t^{-2} + \mathcal{M}^2(0) \left( \frac{1}{F} - \frac{1}{2} \right), \]  

which must be positive during inflation. The equation of motion for the temporal zero-mode of \( \chi \) is

\[ \ddot{\xi}_0 - \mu^2(t)|_{t \gg 1} \xi_0(t) \simeq 0, \]  

which has the general solution

\[ \xi_0(t) = C_1 \sqrt{\frac{t_0}{t}} \mathcal{M} \left[ \frac{F - 2}{4F}, \frac{\sqrt{4 - F(4 + 11F)}}{4F}, \frac{\mathcal{M}^2(0)}{2} \right] t^2 \]
\[ + C_2 \sqrt{\frac{t_0}{t}} \text{W} \left[ \frac{F - 2}{4F}, \frac{\sqrt{4 - F(4 + 11F)}}{4F}, \frac{\mathcal{M}^2(0)}{2} \right] t^2, \]  

where \((C_1, C_2)\) are arbitrary constants and \((\mathcal{M}, \text{W})\) are the Whittaker functions. The asymptotic super Hubble zero-modes matter field fluctuations are (taking \( C_1 = 0 \))

\[ \langle \psi_{cg}^2 \rangle \simeq \frac{C_2^2 \epsilon^3}{6\pi^2} \left( \frac{t_0}{t} \right) \mu^3(t) e^{-\frac{\mathcal{M}^2(0)}{2} t^2} \left[ \text{W} \left[ \frac{F - 2}{4F}, \frac{\sqrt{4 - F(4 + 11F)}}{4F}, \frac{\mathcal{M}^2(0)}{2} \right] \right]^2. \]  

Note that \( \langle \psi_{cg}^2 \rangle \) decreases with time for \( t \gg 1 \). The expression for the energy density fluctuations \( \frac{\delta \rho_t}{\rho_t} \simeq \frac{\psi_{cg}^2}{\phi^2 + \frac{1}{4} \rho_t} \) at the end of inflation can be written as a function of \( \phi \) and \( \theta \) because \( \phi(t) = \lambda^{-1} t^{-1} \)

\[ \frac{\delta \rho_t}{\rho_t} \bigg|_{end} \simeq \frac{C_2 \epsilon^{3/2}}{\sqrt{6\pi}} \sqrt{\frac{\phi}{\phi_s}} \mu^3(\phi) e^{-\frac{\mathcal{M}^2(0)}{4\lambda \epsilon} \phi^{2} - \frac{\mathcal{M}^2(0) \phi + \lambda^2 \phi^3}{\lambda^2 \phi^4 + \frac{4\pi}{90} g_{\text{eff}} \phi^4}} \]
\[ \times \text{W} \left[ \frac{F - 2}{4F}, \frac{\sqrt{4 - F(4 + 11F)}}{4F}, \frac{\mathcal{M}^2(0)}{2} \right] \phi^2 \bigg|_{\phi_c, \theta_c}, \]  

where, as was previously calculated, the values at the end of inflation for the matter field and the temperature are \( \phi_c \approx 0.001 \ G^{-1/2} \) and \( \theta_c \approx 10^{-4} \ G^{-1/2} \). Using \( \epsilon = 10^{-3} \) and \( F = 0.3 \), one can calculate the value for the constant \( C_2 \) to obtain the experimental COBE data amplitude for the fluctuations at the end of inflation: \( \frac{\delta \rho_t}{\rho_t} \approx 1 \times 10^{-5} \)

\[ C_2 \approx 2.21 \times 10^{12}, \]  

where the calculations were made using \( \mathcal{M}^2(0) = 10^{-14} \ G^{-1} \).
IV. FINAL COMMENTS

A possible negative aspect of other warm inflationary models is closely related to a possible thermodynamic fine-tuning, because an isothermal evolution of the radiative component is assumed from the very beginning. In the model here developed inflation begins from zero radiation energy density. This is the more significative difference with other warm inflationary models. Note than when inflation starts the model gives $\Gamma \ll H$, but nearly the equilibrium configuration (i.e., for $t > 10^7 G^{1/2}$) the situation is inverse $[H(\phi_e) \ll \Gamma(\phi_e, \theta_e)]$, and the thermal equilibrium is guaranted. Such an effective friction parameter gives a radiation dominated phase at the end of fresh inflation $[\rho_r(\theta_e) \approx (10^{-3} G^{-1/2})^4]$, which is needed for a natural warm inflation - hot big bang transition.

The main difference between standard inflation and the model here studied is that here the universe is heating during the inflationary regime. Furthermore, the dynamics of this heating being perfectly described. Reheating is not a minor phase at the end of standard inflation. Standard inflation may cause the monopole and domain wall nonthermal symmetry restoration [1], which could give a very inhomogeneous universe that disagree with observation. This last possibility is actually rather difficult in simple models of reheating with only two fields [21]. However, this situation can be solved in the case of multiple fields, relevant for GUT models [22]. This is the case here studied, where the number of fields is nearly $n \approx 4 \times 10^8$ for the set of valued parameters of the model. In the model here studied, the particle creation can be justified by means of the interaction between the inflaton with other particles of the thermal bath [13,18,19]. However, the main difference between this model and standard inflation resides in that here there is not oscillation of the inflaton field around the minimum of the potential, due to dissipation is too large at the end of inflation. Here, particles creation becomes during inflation, begining it at zero temperature. Recently, Chung et al. [27] showed that resonant production of particles during inflation from a zero temperature initial state can take place. More recently, A. Berera and R. O. Ramos [28] demonstrated that the zero temperature initial state constitutes a baseline effect that should be revalent in any general statistical state, which gives a strong evidence that dissipation is the norm not exception for an interacting scalar field system. However, must be noted that this formalism is realized in an non-expanding Minkowski spacetime. In this work I attempted to develop a model in which the universe, starting from chaotic initial conditions, expands in an increasing damped regime product of the interaction of the inflaton field with other scalar fields of a zero temperature initial state. The temperature increases with the expansion of the universe because the inflaton transfers radiation energy density to the bath with rate bigger than the expansion of the universe. But the crucial point here is that this model attempt to build a bridge between standard and warm inflationary models, begining from chaotic initial conditions which provides naturality. This is the main difference between the scenario here worked with other. Finally, the model here studied may provides the necessary number of e-folds to explain the flatness/horizon problem for $\frac{\delta \rho_t}{\rho_t} < 0.55$. This inequality assures slow-roll conditions during the inflationary expansion of the universe. The model provides enough postinflationary radiation temperature at the end of inflation to baryogenesis can take place after inflation. Furthermore, the amplitude for energy density fluctuations $[\delta \rho_t/\rho_t]$ decreases with time, being of the order of $|\delta \rho_t/\rho_t| \approx 1 \times 10^{-5}$ when $t$ assumes the value $t_e = 10^{10} G^{1/2}$. 

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9
ACKNOWLEDGMENTS

MB acknowledges the support of CONACYT (México).
REFERENCES

[1] A. H. Guth, Phys. Rev. D23, 347 (1981).
[2] A. H. Guth and E. J. Weinberg, Nucl. Phys. B212, 321 (1983).
[3] A. D. Linde, Rep. Prog. Phys. 47, 925 (1984); P. J. Steinhardt, Comments Nucl. Part. Phys. 12, 273 (1984).
[4] A. D. Linde, Phys. Lett. B108, 389 (1982).
[5] B. S. De Witt, Phys. Rev. 160, 1113 (1967); J. A. Wheeler, in: Relativity, Groups, and Topology, edited by C. M. De Witt and J.A. Wheeler, Benjamin, NY (1968).
[6] A. A. Starobinsky, in Fundamental Interactions, ed. V. N. Ponomarev (MGPI Press, Moscow, 1984); M. Bellini, H. Casini, R. Montemayor, P. Sisterna, Phys. Rev. D54, 7172 (1996).
[7] A. Berera, Li-Zhi Fang, Phys. Rev. Lett. 74, 1912 (1995).
[8] A. D. Linde, Particle Physics and Inflationary Cosmology, (Harwood Academic Publishers, NY, 1990); T. Padmanabhan, Structure Formation in the Universe (Cambridge University Press, 1993).
[9] A. D. Dolgov and A. D. Linde, Phys. Lett. 116B, 329 (1982); L. F. Abbott, E. Fahri, and M. Wise, Phys. Lett. 117B, 29 (1982); L. Kofman, A. D. Linde, A. A. Starobinsky, Phys. Rev. D56, 3258 (1997).
[10] J. Baacke, K. Heitmann, and C. Patzold, Phys. Rev. D55, 7815 (1997).
[11] L. Kofman, A. Linde, and A. A. Starobinsky, Phys. Rev. Lett. 76, 1011 (1996).
[12] M. Bellini, Phys. Lett. B232, 31 (1998); M. Bellini, Phys. Rev. D58, 103518 (1998).
[13] A. Berera, Phys. Rev. Lett. 75, 3218 (1995).
[14] H. P. de Oliveira and R. O. Ramos, Phys. Rev. D57, 741 (1998).
[15] A. D. Linde, Phys. Lett. B129, 177 (1983).
[16] H. Georgi and S. L. Glashow, Phys. Rev. Lett. 32, 438 (1974).
[17] A. Albrech, P. J. Steinhardt and M. S. Turner, Phys. Rev. Lett. 48, 1437 (1982).
[18] see for example A. Berera, Phys. Rev. D, 2519 (1996).
[19] M. Bellini, Class. Quantum Grav. 17, 145 (2000); M. Bellini, Nucl. Phys. B563, 245 (1999). E-print: gr-qc/9908063; see also M. Bellini, Gen. Rel. Grav. 33, 127 (2001). E-print: gr-qc/0007083.
[20] J. M. F. Maia, J. A. S. Lima, Phys. Rev. D60, 101301 (1999).
[21] D. Boyanovsky, H. J. de Vega, R. Holman and J. F. J. Salgado, Phys. Rev. D54, 7570 (1996).
[22] B. A. Basset, F. Tamburini, Phys. Rev. Lett. 81, 2630 (1998).
[23] A. Berera, M. Gleiser, R. O. Ramos, Phys. Rev. D58, 123508 (1998).
[24] J. Yokoyama, A. Linde, Phys. Rev. D60, 083509 (1999).
[25] S. Weinberg, Phys. Rev. D9, 3357 (1974).
[26] M. Bellini, Class. Quantum Grav. 16, 2393 (1999).
[27] D. J. H. Chung, E. W. Kolb, A. Riotto, I. I. Tkachev, Phys. Rev. D62, 043508 (2000).
[28] A. Berera, R. O. Ramos, e-print hep-ph/0101043.