Journal of Cosmology and
Astroparticle Physics

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Journal of Cosmology and Astroparticle Physics
An IOP and SISSA journal

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Received July 7, 2016
Revised October 27, 2016
Accepted October 29, 2016
Published November 9, 2016

Abstract. We consider Parker particle production in the Ekpyrotic scenario (in particular in the New Ekpyrotic model) and show that the density of particles produced by the end of the phase of Ekpyrotic contraction can be sufficient to lead to a hot state of matter after the bounce. Hence, no separate reheating mechanism is necessary.

Keywords: alternatives to inflation, physics of the early universe

ArXiv ePrint: 1605.04670

doi:10.1088/1475-7516/2016/11/024

JCAP11(2016)024

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1 Introduction

There is overwhelming observational evidence that there was a early phase in the evolution of the universe when Standard Model matter was in thermal equilibrium (see e.g. [1–3] for modern textbooks on cosmology). The best quantitative evidence for such an early phase comes from the black body nature of the cosmic microwave background [4, 5] and from the abundances of light nuclei (see e.g. [6]).

In Standard Big Bang cosmology it is assumed that the Universe begins in a hot and dense thermal state. However, this scenario cannot explain the isotropy of the cosmic microwave background, and it cannot explain the origin of observed inhomogeneities on length scales which were larger than the Hubble radius at the time t_{eq} of equal matter and radiation.

The inflationary universe scenario [7] (see also [8–13]) provides a solution to these problems of Standard Big Bang cosmology, and at the same time yields a causal theory for the origin of the inhomogeneities which are now explored through cosmological observations [14] (see also [15]). The inflationary scenario posits a phase of almost exponential expansion of the early universe. This phase leaves behind a vacuum state of the Standard Model matter fields. Hence, in order to make contact with the late time cosmology, a new phase must be posited during which Standard Model matter fields approach a hot thermal state. This phase is called *reheating phase*, and is an essential part of the inflationary scenario (see e.g. [16, 17] for initial studies of this phase). Reheating in the inflationary scenario is a semiclassical effect. It relies on the squeezing of quantum vacuum perturbations left behind at the end of the inflationary phase. It is a consequence of coupling terms between the scalar field ϕ which generated inflationary expansion and the Standard Model fields. As first studied in [18, 19] and worked out in more detail in [20, 21] (see [22, 23] for recent reviews on reheating), the energy transfer from ϕ to Standard Model fields proceeds via a parametric instability in the equation of motion for the Standard Model fields induced by the time dependence of ϕ. As a consequence, it is expected that most of the energy density ρ_R at the final time t_R of the phase of exponential expansion ends up in the hot plasma of the post-inflationary universe.[1]

Inflation is not the only early universe scenario which is compatible with cosmological observations (see e.g. [24–26] for recent reviews of alternatives). Bouncing cosmologies with an initial matter-dominated phase of contration (see [27] for a recent review of this *matter bounce* scenario) also produce a spectrum of fluctuations compatible with observations, as does *string gas cosmology* [28, 29] (see also [30–33] for reviews) and the *Galileon genesis scenario* [34], two early universe scenarios of emergent type. However, as emphasized in [35],

[1]Note that the state of Standard Matter fields after preheating does not have a thermal distribution of momenta.

bouncing cosmologies suffer from an instability of the contracting phase to the growth of anisotropies. The Ekpyrotic scenario [36] is an alternative to cosmological inflation which is free of the anisotropy problem [37].

It is of interest to study how the late time hot thermal state emerges in models alternative to inflation, and whether it is necessary to introduce a new mechanism analogous to inflationary reheating. In *string gas cosmology* the hot thermal state of matter emerges directly after the initial stringy Hagedorn phase [28, 29], and no additional physics must be added to the system. Defrosting at the end of the Galileon phase was studied in [38], and was shown to rely on analogous coupling terms as are used in inflationary cosmology. On the other hand, it was shown in [39] that Parker particle production [40, 41] is sufficient to produce a hot early universe in the *matter bounce* scenario. Here, we study matter particle production in the *New Ekpyrotic* scenario [42] (see also [44–47]), a version of the Ekpyrotic scenario which generates an approximately scale-invariant spectrum of cosmological perturbations. We show that the analysis of [39] carries over and that Parker particle production is sufficient to lead to a hot thermal state of matter.[2]

In the following section we briefly review Parker particle production, and in the third section we summarize the application to the New Ekpyrotic scenario. To set our notation, we use units in which $c = k_b = 1$, and the usual space-time coordinates in which the background metric is

$$ds^2 = dt^2 - a(t)^2 dx^2 \,, \tag{1.1}$$

where t is physical time and x are the comoving spatial coordinates (we take the spatial sections to be flat). Linear metric and matter fluctuations about this background have a complete basis of fundamental solutions which are Fourier modes in comoving spatial coordinates with a time-dependence which depends on the background. The comoving momentum is denoted by k. An important length scale is the Hubble radius

$$l_H(t) \,\equiv\, H(t)^{-1} \,\equiv\, \frac{a(t)}{\dot{a}(t)} \,. \tag{1.2}$$

The Hubble radius separates length scales where the mode evolution is qualitatively different.

2 Review of Parker particle production

Parker particle production [40, 41] (see also [48–50]) is a phenomenon which was first studied in the context of quantum field theory in curved space-time. If we assume that the space-time is Minkowski space-time both at early and at late times, but underwent a period of expansion during an intermediate time interval, then the initial Minkowski vacuum state of a test scalar field χ will evolve non-trivially during the intermediate time interval, and evolve into a final state of χ which is not equal to the final time Minkowski vacuum. From the point of view of the final Minkowski frame, the final state contains χ particles.

The same phenomenon applies to linear cosmological perturbations which evolve in a similar way to test scalar fields on the cosmological background (see e.g. [51, 52] for reviews of

[2] As in the case of preheating after inflation, the state which results from Parker particle production does not have a thermal distribution of excitations, but it is a state which consists of particle excitations as opposed to being a coherent homogeneous condensate, a state which can subsequently thermalize via particle interactions. Whereas preheating after inflation produced mainly infrared excitations, we will see that Parker particle production in the New Ekpyrotic scenario leads to a spectrum of particles peaked in the ultraviolet, and hence closer to a thermal equilibrium state.

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the theory of cosmological perturbations). On length scales smaller than the Hubble radius $H^{-1}(t)$, where $H(t)$ is the expansion rate, fluctuations oscillate as they do in Minkowski space-time, whereas they are squeezed on super-Hubble scales. The squeezing of fluctuations on super-Hubble scales describes the growth of cosmological perturbations, and this corresponds to particle production. If the equation of state of the background cosmology is constant in time, then the squeezing of the cosmological fluctuations occurs at the same rate as the squeezing of test fields (which is the same as the squeezing rate of gravitational waves).

In the case that the background cosmological space-time metric is constant in time both in the far past and in the future, then we know that the wave function of a test scalar field can be written both in the far past and in the future as

$$\chi(x,t) = \frac{1}{\sqrt{2\pi}} V^{-1/2} \int d^3k \left[\alpha_k \psi_k^+ + \beta_k \psi_k^- \right] \tag{2.1}$$

in terms of coefficient functions α_k and β_k. Here, χ_k^+ and χ_k^- are the Minkowski space-time positive and negative frequency wave functions

$$\begin{aligned} \chi_k^+ &= \frac{1}{\sqrt{2k}} e^{ikx} \\ \chi_k^- &= \frac{1}{\sqrt{2k}} e^{-ikx} . \end{aligned} \tag{2.2}$$

If we start the evolution in the vacuum state $|\Psi>_i$ of the system (the harmonic oscillator vacuum of each Fourier mode), then initially $\alpha_k = 1$ and $\beta_k = 0$. This is the Bunch-Davies [53] vacuum. An observer at late times will define a new vacuum state $|\Psi>_f$ in which the coefficient functions have vanishing β_k and $\alpha_k = 1$ at the final time. Because of the time dependence of the background, the state of the system $|\Psi>$ which equals the initial vacuum at early times will evolve non-trivially and turn into a state for which at the final time $\beta_k \neq 0$. The late time coefficients α_k and β_k are called the Bogoliubov coefficients, and they obey the relation (see e.g. [54])

$$\alpha_k^2 - \beta_k^2 = 1 . \tag{2.3}$$

The interpretation of this state for the late time observer is that of a state which contains

$$n_k = |\beta_k|^2 \tag{2.4}$$

particles of comoving wave number k.

In most cosmological models (in particular in inflationary cosmology and in the Ekpyrotic scenario) the metric is neither static initially nor today. However, modes which are relevant to current cosmological observations are inside the Hubble radius both at very early times and today. On scales smaller than the Hubble radius the mode functions of the canonically normalized fields[3] oscillate, and hence the mode wavefunction can be represented in the form (2.1). In the early universe scenarios which have a chance of explaining the origin of the observed structure in the Universe the modes exit the Hubble radius during the initial phase (e.g. the inflationary phase in the case of inflationary cosmology, or the Ekpyrotic phase of contraction in the case of Ekpyrotic cosmology), and re-enter at late times. While the scale is super-Hubble, we can still expand the mode functions at any time t in terms of a local

[3]These are fields with canonical kinetic term. For a test scalar field χ the canonical field is $a(t)\chi$.

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Minkowski frame in the form (2.1) with time-dependent Bogoliubov coefficients. However, the wave function oscillations are frozen out, the state is a squeezed state, adiabaticity is violated and particle number is not a well-defined quantity (see e.g. the discussion in [55]). Hence the Bogoliubov coefficients should not be interpreted as yielding the number of particles. Heuristically, one could say that via (2.4) the Bogoliubov coefficients yield the number of proto-particles, a field state which will admit a particle interpretation once the scale enters the Hubble radius.

3 Parker particle production in the new Ekpyrotic scenario

The Ekpyrotic scenario [36] is an alternative to cosmological inflation which was originally motivated by some ideas in superstring theory, in particular heterotic M-theory [56, 57]. In the Ekpyrotic scenario, our universe emerges from an initial contracting phase which arises when two three-space-dimensional branes (one of which corresponds to our space-time) bounding an extra spatial dimension approach eachother. In the original scenario, the branes collide and this corresponds to a "Big Crunch" singularity after which our universe emerges in a Standard Big Bang phase of expansion.

The effective gravitational theory in our space-time is Einstein gravity coupled to a scalar field ϕ (which is proportional to the logarithm of the brane separation) whose potential was assumed to be a steep negative exponential

$$V(\phi) = -V_0 \exp\left(-\sqrt{\frac{2}{p}}\frac{\phi}{m_{\rm pl}}\right), \tag{3.1}$$

with $0 < p \ll 1$ and $V_0 > 0$, and with $m_{\rm pl}$ denoting the Planck mass. Inserting this into the Friedmann equations we find that the contraction is very slow

$$a(t) \sim (-t)^p . \tag{3.2}$$

This corresponds to an equation of state with

$$w \equiv \frac{p}{\rho} \gg 1 . \tag{3.3}$$

In turn, this equation of state implies that the energy density ρ_ϕ stored in the field ϕ increases as

$$\rho_\phi \sim a^{-2/p} , \tag{3.4}$$

which implies that it increases faster than that in regular cold matter, radiation, curvature and anisotropic stress. In particular, in contrast to contracting phases with usual matter content with $w = 0$ or $w = 1/3$, the contracting phase is safe against the BKL instability [58, 59] of the homogeneous bounce, as shown explictly in [60, 61]. This is a significant advantage of the Ekpyrotic scenario compared to most other bouncing models (see e.g. [35] for a discussion of the instability for regular bouncing models, and the last entry of [24–26] for a recent review of problems of regular bouncing models).

As a consequence of the slow contraction, fixed comoving scales exit the Hubble radius during the period of contraction since the Hubble radius decreases as $(-t)$. Hence, a causal generation mechanism of fluctuations is possible. Although the spectrum of ϕ fluctuations produced during Ekpyrotic contraction is scale-invariant [62], that of curvature perturbations is not [63, 64].

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To solve this problem, the *New Ekpyrotic scenario* was proposed [42] (see also [44–47]). The model involves two scalar fields ϕ and ψ, both with negative exponential potentials

$$V(\phi, \psi) = -V_0 \mathrm{exp}\left(-\sqrt{\frac{2}{p}}\frac{\phi}{m_{\mathrm{pl}}}\right) - U_0 \mathrm{exp}\left(-\sqrt{\frac{2}{q}}\frac{\psi}{m_{\mathrm{pl}}}\right) \tag{3.5}$$

with $p \ll 1$ and $q \ll 1$. Since there are two fields present, it is possible to have entropy fluctuations. In the same way that test scalar fields in an Ekpyrotic background acquire a scale-invariant spectrum of fluctuations, in the two field Ekpyrotic scenario the entropy mode acquires a scale-invariant spectrum, and transmits this spectrum to the curvature fluctuations since any entropy fluctuation induces a growing curvature perturbation.[4]

We are interested in a version of the New Ekpyrotic scenario in which new physics (which involves, from the point of view of Einstein gravity, a violation of the Null Energy Condition) leads to a nonsingular bounce occurring when the background density is

$$\rho_{\mathrm{max}} = M^4 \,, \tag{3.6}$$

where M is the mass scale of the new physics leading to the nonsingular bounce.

In the New Ekpyrotic scenario the background trajectory is given by

$$\begin{aligned} a(t) &\sim (-t)^{p+q} \qquad (3.7) \\ \phi(t) &= \sqrt{2p}m_{\mathrm{pl}}\mathrm{log}\left(-\sqrt{\frac{V_0}{m_{\mathrm{pl}}^2 p(1-3(p+q))}}t\right) \\ \psi(t) &= \sqrt{2q}m_{\mathrm{pl}}\mathrm{log}\left(-\sqrt{\frac{U_0}{m_{\mathrm{pl}}^2 q(1-3(p+q))}}t\right) \,. \end{aligned}$$

The field space is two-dimensional. Flutuations in the field space direction parallel to the background trajectory form the adiabatic mode σ given by [66]

$$\dot{\sigma} = \mathrm{cos}\theta\dot{\phi} + \mathrm{sin}\theta\dot{\psi} \tag{3.8}$$

and has adiabatic fluctuations

$$\delta\sigma = \mathrm{cos}\theta\delta\phi + \mathrm{sin}\theta\delta\psi \tag{3.9}$$

while the orthogonal direction s has perturbations

$$\delta s = -\mathrm{sin}\theta\delta\phi + \mathrm{cos}\theta\delta\psi \,. \tag{3.10}$$

In the following we consider production of χ particles.[5] The canonically normalized entropy field perturbation is

$$v \equiv a\delta s \,, \tag{3.11}$$

and obeys the Fourier space equation

$$v'' + k^2 v - \frac{2}{\eta^2}\left[1 - \frac{3}{2}(p+q)\right] v = 0 \,, \tag{3.12}$$

[4]This is a standard result in the theory of cosmological perturbations. For a recent study see e.g. [65].

[5]Note that it is fluctuations in the same field which yield the curvature perturbations in the new Ekpyrotic scenario. We are interested in the short wavelength modes which could provide "matter" fluctuations at late times after the bounce.

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where η is conformal time related to physical time via $dt = ad\eta$, and where a prime denotes the derivative with respect to η. Note that we are suppressing the index k on the Fourier mode of v. In the following we will study the production of v particles in the contracting phase of the Ekpyrotic scenario.[6]

The initial conditions in Ekpyrotic cosmology are taken to be vacuum, i.e. all fields begin in their Bunch-Davies vacuum state at past infinity. All modes undergo quantum vacuum oscillations until their wavelength crosses the Hubble radius, after which they will be squeezed. The squeezing corresponds to particle production (as discussed in the previous section). The modes on cosmological scales develop into the density perturbations which we observe today, modes on microscopic wavelengths (but larger than the Hubble radius at the end of the contracting phase) become - at Hubble radius re-entry, when the concept of particles becomes well-defined - the particles whose production we are interested in. We are interested in computing the energy density in particles produced by the end of the Ekpyrotic phase of contraction, i.e. when the density reaches $\rho_{\rm max}$.

We will consider a slightly more general setup which can be applied not only to the New Ekpyrotic scenario, but also to others. The equation of motion generalized from (3.12) is

$$v'' + \left(k^2 - \frac{z''}{z}\right) v'' = 0\,, \tag{3.13}$$

where η is the conformal time which in a contracting background goes from $-\infty$ to 0. We posit

$$\frac{z''}{z} = \frac{\nu^2 - 1/4}{\eta^2}\,, \tag{3.14}$$

with

$$\nu \equiv \frac{1+\tilde{p}}{2(1-\tilde{p})}\,, \tag{3.15}$$

where the value of $\tilde{p}$ depends on the specific model. In the case of New Ekpyrotic scenario we have (see (3.12))

$$\nu = \sqrt{\frac{9}{4} - 3(p+q)}\,. \tag{3.16}$$

There is an "effective" Hubble radius which divides modes which oscillate from those which are squeezed. For a mode with comoving wave number k, the conformal time $\eta_H(k)$ of effective Hubble radius crossing is given by

$$k^2\eta_H(k)^2 = \nu^2 - \frac{1}{4}\,. \tag{3.17}$$

At the beginning of the contraction phase all scales we are interested in are inside the effective Hubble radius and hence the k^2 term dominates over the z''/z term, and we start with the Bunch-Davies solution

$$v = v_{\rm BD} \equiv \frac{e^{-ik\eta}}{\sqrt{2k}}\,. \tag{3.18}$$

After effective Hubble radius crossing (which occurs at the time $\eta = \eta_H(k)$), the k^2 is subdominant then we have as solution

$$v = c_1(k)\frac{\eta^{1/2-\nu}}{\eta_H(k)^{1/2-\nu}} + c_2(k)\frac{\eta^{1/2+\nu}}{\eta_H(k)^{1/2+\nu}} \tag{3.19}$$

[6]Here, "particle" has to be interpreted in the sense described at the end of the previous section.

where $c_1(k)$ and $c_2(k)$ are constants which are found by matching v and v' at effective Hubble radius crossing $\eta = \eta_H(k)$. Then

$$
\begin{aligned}
c_1(k) &= \frac{1}{2\nu}\frac{1}{\sqrt{2k}}e^{-ik\eta_H(k)}\left[\nu + \frac{1}{2} + ik\eta_H(k)\right] , \\
c_2(k) &= \frac{1}{2\nu}\frac{1}{\sqrt{2k}}e^{-ik\eta_H(k)}\left[\nu - \frac{1}{2} - ik\eta_H(k)\right] .
\end{aligned} \tag{3.20}
$$

The Bogoliubov coefficient β_k at a time η closer to the bounce can be obtained by expanding the solution (3.19) in terms of the local Bunch-Davies state given by v_{BD} at the time η (where $k|\eta| \ll 1$)

$$
\begin{aligned}
v &= \alpha_k v_{\mathrm{BD}} + \beta_k v^*_{\mathrm{BD}} \\
v' &= \alpha_k v' + \beta_k v'^*_{\mathrm{BD}}
\end{aligned} \tag{3.21}
$$

where the star stands for complex conjugation. Keeping only the growing solution on super-Hubble scales this yields

$$
\beta(\eta) = \frac{c_1(k)\sqrt{2k}}{2}\left(\frac{\eta}{\eta_H(k)}\right)^{1/2-\nu}\left[1 + \frac{1/2-\nu}{ik\eta}\right] . \tag{3.22}
$$

In case of the original Ekpyrotic model we have $\nu \approx 1/2$ and thus

$$
\beta_k = \frac{1}{2}e^{-ik\eta_H(k)}\left[1 + ik\eta_H(k)\right] . \tag{3.23}
$$

Hence

$$
n_k = |\beta_k|^2 = \frac{1}{4}\left[1 + (k\eta_H(k))^2\right] , \tag{3.24}
$$

which for $k\eta_H(k) \ll 1$ is $n_k \approx 1/4$. Hence, there is no significant particle production until the end of the contracting phase in the original Ekpyrotic scenario.

In the case of New Ekpyrotic model $\nu \approx 3/2$. At first sight it looks like the second term on the right hand side of (3.22) will dominate. However, the contribution of this term does not have an interpretation as particles.[7] As mentioned at the end of the previous section, the particle interpretation only applies when scales re-enter the Hubble radius. At that time, the solution for v should be interpreted as a standing wave which decays equally into a right-moving and left-moving wave. Hence, the value of β is given by half of the ratio of the final amplitude of v to the initial amplitude, and we have

$$
\begin{aligned}
\beta_k(\eta) &= \frac{c_1(k)\sqrt{2k}}{2}\left(\frac{\eta}{\eta_H(k)}\right)^{-1} \\
&= \frac{1}{3}e^{-ik\eta_H(k)}\left[1 + i\frac{k\eta_H(k)}{2}\right]\left(\frac{\eta}{\eta_H(k)}\right)^{-1} .
\end{aligned} \tag{3.25}
$$

The corresponding number density of produced particles is

$$
n_k(\eta) = \frac{1}{9}\left[1 + \frac{(k\eta_H(k))^2}{4}\right]\left(\frac{\eta}{\eta_H(k)}\right)^{-2} . \tag{3.26}
$$

[7] Another argument for neglecting this term is the following: once the scales re-enter the Hubble radius and the particle interpretation becomes valid, this term is small compared to the one we are keeping.

We now can compute the energy density $\rho_p(\eta)$ in the produced particles (strictly speaking it is the energy density which the state will have once the scales re-enter the Hubble radius at late times) at the end of the Ekpyrotic phase of contraction, a time we denote by $t_{\rm end}$. We choose to normalize the scale factor such that $a(t_{\rm end}) = 1$, and thus physical and comoving momenta coincide at this time. We have

$$\rho_p(\eta) = \frac{1}{(2\pi)^3} \int_0^{k_H(\eta_{\rm end})} n_k k d^3k \,, \tag{3.27}$$

where the final factor of k represents the energy of the mode once it starts to oscillate. Inserting (3.26) and making use of the relation (3.17) to determine $k_H(\eta)$ in terms of η, and in addition using the fact that for Ekpyrotic contraction $\eta_{\rm end} \simeq t_{\rm end}$ we find

$$\rho_p(t_{\rm end}) \sim t_{\rm end}^{-4} \,, \tag{3.28}$$

which is to be compared with the background density $\rho_{\rm bg}$

$$\rho_{\rm bg}(t_{\rm end}) \sim t_{\rm end}^{-2} m_{\rm pl}^2 \,. \tag{3.29}$$

If the background density at the final time is given by M^4 in terms of a "new physics mass scale" then we find

$$\frac{\rho_p(t_{\rm end})}{\rho_{\rm bg}(t_{\rm end})} \sim \left(\frac{M}{m_{\rm pl}}\right)^4 \,. \tag{3.30}$$

This result implies that if the scale of new physics is high (e.g. between the scale of particle physics "Grand Unification" and the Planck scale), that then a sufficiently high density of particles is produced to lead to post-bounce hot big bang phase beginning at temperatures not much lower than that of Grand Unification. Note that the energy density which at the bounce remains in the Ekpyrotic field rapidly redshifts relative to that in the produced particles after the bounce.

One might wonder if the new Ekyrotic model is compatible with having a higher energy scale than the Grand Unification one. In the literature on the new Ekpyrotic model [42, 43], it is generally assumed that the background density at the end of the Ekpyrotic phase is of order of the grand unification scale in order for the power spectrum of the curvature perturbations to have the correct normalization, namely $\sim 10^{-10}$. However, this is a simplifying assumption for a specific realization of the exit from the Ekpyrotic phase, parametrized by the model dependent quantity $\beta = \arctan(\sqrt{qV_0/pU_0})$. This quantity appears in the power spectrum of curvature perturbations (which is related to the power spectrum of entropy perturbations, P_S):

$$k^3 P_\zeta(k) = \frac{4\epsilon\beta^2}{M_{\rm pl}^2} \, k^3 P_S(k) = \beta^2 \frac{H_{\rm end}^2}{2\epsilon M_{\rm pl}^2} \,, \tag{3.31}$$

where, in the new Ekpyrotic model literature, it is assumed that $H_{\rm end}^2/M_{\rm pl}^2 = \rho_{\rm bg}/M_{\rm pl}^4$, where $\rho_{\rm bg} \sim M_{\rm GUT}^4$, with ϵ, the analogous to the slow roll parameter from inflation, of order 10^{-2}. Thus, one obtains the restriction that $\beta \sim \mathcal{O}(1)$, which represents model with a sharp turn in the field trajectory at the end of the Ekpyrotic phase. However, this parameter can be smaller than one still maintaining rapid change in the Ekpyrotic potential at the end of the Ekpyrotic phase necessary in order to have a subsequent bouncing phase. We assume here then that $\beta < 1$, allowing for the energy scale of our model to be higher than the grand unification, and maintaining the correct normalization of the curvature power spectrum.

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4 Conclusions and discussion

We have studied Parker particle production in the contracting phase of the New Ekpyrotic scenario and have found that the process can be sufficiently efficient to lead to a hot thermal expanding universe beginning at temperatures not much lower than that of Grand Unification provided that the new physics scale which yields the bounce is higher than that of Grand Unification. Hence, the scenario does not require a separate physics sector to generate reheating.

Note that the distribution of particles produced by the Parker process does not have a thermal spectrum, as is the case after preheating in inflationary cosmology. What the Parker process does (and similarly preheating) is to convert energy from the homogeneous condensate to particle quanta. These quanta can then interact and thermalize by particle scattering.

We have only considered Parker particle production in the contracting phase of the Ekpyrotic scenario. This leads to a lower bound on the total number of particles produced during the entire cosmological evolution. The reason is twofold. First, particles will also be produced during the bounce phase and in the post-bounce expanding phase, and this will add to the total number of particles produced. However, the bounce phase is expected to be short compared to a Hubble expansion time at the end of the Ekpyrotic phase, and hence we do not expect Parker particle production to be important during the bounce. During the post-bounce phase of radiation-dominated expansion there is no squeezing of the fluctuations, and hence we do not expect Parker particle production during most of the post-bounce period. The second point which supports our statement that we have computed a lower bound on the effectiveness of Parker particle production is that the dominant mode will continue to be squeezed on super-Hubble scales after the Ekpyrotic phase of contration ends, and hence the particles produced during the contracting phase will remain.

Note that the method we used to analyze Parker particle production is applicable to a wide range of models. The application to the *matter bounce* scenario was already considered in [39]. For all contracting models with $\nu \neq 1/2$ Parker particle production may be important.

Note also that Parker particle production is operative in inflationary cosmology, and produces an energy density of the order of H_I^4, where H_I is the Hubble constant during inflation. In the case of inflation, however, we do not have the freedom to raise the value of H_I to increase the energy density of produced particles to that of grand unification, the reason being that the spectrum of gravitational waves is scale-invariant, given by the value of H_I, and bounded from above from observations. In the case of the new Ekpyrotic scenario the spectrum of gravitational waves is a blue vacuum spectrum, and hence gravitational waves do not lead to a constraint on the scale of the energy density of produced particles.

Acknowledgments

One of us (RB) wishes to thank the Institute for Theoretical Studies of the ETH Zürich for kind hospitality. RB acknowledges financial support from the "Dr. Max Rössler and Walter Haefner Foundation" and the ETH Zurich Foundation, and from a Simons Foundation fellowship. The research of RB is also supported in part by funds from NSERC and the Canada Research Chair program. WSHR is grateful for the hospitality of the Physics Department of McGill University. WSHR was supported by Brazilian agencies CAPES (process No 99999.007393/2014-08) at the beginning of this work, and FAPES at the end (BPC No 476/2013). EF and LG acknowledge financial support from CNPq (Science Without Borders). We also wish to thank J. Quintin for useful discussions.

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