On adiabatic perturbations in the ekpyrotic scenario

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In a recent paper, Khoury and Steinhardt [1] proposed a way to generate adiabatic cosmological perturbations with a nearly flat spectrum in a contracting Universe. To produce these perturbations they used a regime in which the equation of state exponentially rapidly changed during a short time interval. Leaving aside the singularity problem and the difficult question about the possibility to transmit these perturbations from a contracting Universe to the expanding phase, we will show that the methods used in [1] are inapplicable for the description of the cosmological evolution and of the process of generation of perturbations in this scenario.

I. INTRODUCTION

Inflationary theory provides a simple solution to many cosmological problems [2–5]. It also provides a simple mechanism for the generation of adiabatic perturbations of the metric with flat spectrum, which are responsible for large-scale structure formation and for the observed CMB anisotropy [6, 7]. Many predictions of inflationary cosmology are already confirmed by observations. There were many attempts to propose an alternative, equally compelling cosmological theory, but so far they have all been unsuccessful.

Nothing illustrates this statement better than the history of the development of the ekpyrotic/cyclic scenario [8]. This scenario was supposed to solve all cosmological problems without using a stage of inflation. However, the original version of this scenario did not work. The large mass and entropy of the universe remained unexplained, instead of solving the homogeneity problem this scenario only made it worse, and instead of the big bang expected in [8], there was a big crunch [9, 10].

As a result, this scenario was replaced by the cyclic scenario, which postulated the existence of an infinite number of periods of expansion and contraction of the universe [11]. However, when this scenario was analyzed, taking into account the effect of particle production in the early universe, a very different cosmological regime was found [12, 13]. In the latest version of this scenario, the homogeneity problem is supposed to be solved by the existence of an infinitely many stages of low-scale inflation, each of which should last at least 56 e-folds [14].

The most difficult problem facing this scenario is the problem of the cosmological singularity. Originally there was a hope that the cosmological singularity problem will be solved in the context of string theory, but despite the attempts of the best experts in string theory, this problem remains unsolved. A recent attempt to solve this problem was made in the so-called ‘new ekpyrotic scenario’ [15–18]. This scenario involved many exotic ingredients, such as a violation of the null energy condition in a model which combined the ekpyrotic scenario [8] with the ghost condensate theory [19]. However, this model contained ghosts [20]. The negative sign of the higher derivative term in this scenario makes it hard to construct a ghost-free UV completion of this theory [20]; see also [21] for a discussion of another attempt to improve the new ekpyrotic scenario. Even if one eventually succeeds in inventing a UV complete generalization of this theory, such a generalization is expected to obey the null energy condition, in which case the new ekpyrotic scenario will not work anyway.

Leaving all of these unsolved problems aside, the authors of the ekpyrotic scenario continue investigating generation of cosmological perturbations in various versions of this scenario. In inflationary theory, these perturbations are generated after the singularity, but in the ekpyrotic scenario they are supposed to be generated in a collapsing universe before the singularity and preserved on the way through the singularity. Since the singularity problem is not solved, one cannot really say much about it, unless one makes additional assumptions. After a few years of debate, most of the experts agreed, following [24], that adiabatic perturbations with a flat spectrum are not produced in the simplest one-field versions of this theory, contrary to the original claims by the authors of the ekpyrotic/cyclic scenario [8, 11, 25]. Therefore more complicated versions of the ekpyrotic scenario were proposed [20], in which the adiabatic perturbations are produced by a mechanism similar to the inflationary
curvaton mechanism\textsuperscript{[27]}\textsuperscript{[27]}. The simplest versions of the
ekpyrotic scenario of this type were ruled out in\textsuperscript{[28]}\textsuperscript{[28]}.

That is why it is important to examine the recent pa-
paper by Khoury and Steinhardt\textsuperscript{[1]}\textsuperscript{[1]} who have proposed a
new mechanism of generation of adiabatic cosmological
perturbations with a nearly flat spectrum at the stage
prior to the stage of the ekpyrotic collapse. This is the
main goal of our paper. In section III we will examine
the unperturbed background cosmological solution
describing the model of Ref.\textsuperscript{[1]}\textsuperscript{[1]}. As we will see, this solution
requires investigation of the cosmological evolution
at super-planckian values of curvature, when the meth-
ods used in\textsuperscript{[1]}\textsuperscript{[1]} are not valid. A similar conclusion will be
reached in section III where we will show that the theory
of cosmological perturbations used in\textsuperscript{[1]}\textsuperscript{[1]} is also not valid.

II. BACKGROUND MODEL

The model relies on investigation of a contracting uni-
versal in a theory with the scalar field potential

\[ V = V_0 \left(1 - e^{-c\varphi}\right), \]

(1)

where \(c \gg 1\). To simplify a comparison of the results, we
will use the same notations and units (\(8\pi G = 1\)) as in\textsuperscript{[1]}\textsuperscript{[1]}.
We will start with the homogeneous background solution.
Ignoring the evolution of the universe, the equation of
motion for \(\varphi\) is

\[ \dot{\varphi} + cV_0 e^{-c\varphi} = 0. \]

(2)

It describes evolution of the field \(\varphi\) when \(t\) slowly in-
creases from large negative \(t\). Among other solutions,
this equation has the solution used in\textsuperscript{[1]}\textsuperscript{[1]}:

\[ \varphi = \frac{2}{c} \ln \left(-\sqrt{\frac{V_0}{2c}} \right). \]

(3)

This solution describes a slowly contracting universe with
a nearly unchanging Hubble parameter \(H \approx H_0 < 0\).
One can check that during the pre-ekpyrotic stage, at
\(t_{\text{begin}} < t < t_{\text{end}}\), where

\[ t_{\text{begin}} \approx \frac{1}{H_0}, \quad t_{\text{end}} \approx \frac{1}{H_0 c^2} = \frac{t_{\text{begin}}}{c^2}, \quad H_0 = -\sqrt{\frac{V_0}{3}}, \]

(4)

this solution does not receive significant corrections even
if one takes into account the contraction of the universe and
adds the term \(3H\dot{\varphi}\) to equation (2).

During the time interval from \(t_{\text{begin}}\) to \(t_{\text{end}}\)\textsuperscript{[1]}\textsuperscript{[1]} the scale
factor does not decrease significantly and therefore one

can set it to be \(a \approx 1\). Note that the time variable \(t\)
is negative and its magnitude changes from a large negative
value at the beginning to a small negative value at the end
of this contracting stage. At the same time, the deriva-
tive of the Hubble constant \(\dot{H}\) and, correspondingly, the
equation of state \(w = p/\rho\) change tremendously during
the same time interval. Indeed,

\[ \dot{H} = -\frac{1}{2} \dot{\varphi}^2 \approx -\frac{2}{c^2 t^2}, \]

(5)

hence \(\dot{H}_{\text{begin}} \approx H_0^2 c^2\) and \(\dot{H}_{\text{end}} \approx c^2 H_0^2\). For the equa-
tion of state we have

\[ 1+w = -\frac{2}{3} \frac{\dot{H}}{H^2} \approx \frac{1}{H_0^2 c^2} \approx \frac{1}{c^2} \left(\frac{t_{\text{begin}}}{t}\right)^2 \approx c^2 \left(\frac{t_{\text{end}}}{t}\right)^2. \]

(6)

At the beginning \((1+w)_{\text{begin}} \approx c^{-2} \ll 1\) and at the end
\((1+w)_{\text{end}} \approx c^2 \gg 1\).

According to\textsuperscript{[1]}\textsuperscript{[1]}, the stage described above during
which the spectrum of perturbations is generated is fol-
lowed by a subsequent scaling ekpyrotic stage, when the
Hubble scale evolves as

\[ H \approx H_0 \left(\frac{t_{\text{end}}}{t}\right). \]

(7)

The authors of\textsuperscript{[1]}\textsuperscript{[1]} derive the following necessary con-
tion which must be satisfied if we want the perturbations
generated at the pre-ekpyrotic stage to be on the observ-
ables scales today (see equation (16) in\textsuperscript{[1]}\textsuperscript{[1]}):

\[ \sqrt{V_0} \leq 10^{-30} \sqrt{|H_{\text{end}}|}, \]

(8)

where

\[ H_{\text{end}} \approx H_0 \left(\frac{t_{\text{end}}}{t_{\text{ekp}}}\right), \]

(9)

is the Hubble constant not at \(t = t_{\text{end}}\), but at the end
of the ekpyrotic stage, at \(t = t_{\text{ekp}}\). Taking into account
that \(H_0 \approx -\sqrt{V_0}\) we obtain from (5) that this necessary
condition can be rewritten as

\[ |t_{\text{ekp}}| \leq 10^{-60} H_0^{-1} |t_{\text{end}}| \approx 10^{-60} \frac{1}{c^2 H_0^2}. \]

(10)

Here time is expressed in units of Planck time \(t_P \sim M_P^{-1} \sim 10^{-43}\) s. On the other hand, according to\textsuperscript{[1]}\textsuperscript{[1]} (see
also below), the amplitude of the generated perturbations
is \(O(cH_0)\). Since this amplitude should be \(O(10^{-5})\), the
ekpyrotic stage should end at

\[ |t_{\text{ekp}}| < 10^{-50} t_P. \]

(11)

In other words, we are dealing here with frequencies
which are 50 orders of magnitude greater than the Planck
mass! This completely invalidates the use of the classi-
cal equations of motion in the investigation of the back-
ground solution in\textsuperscript{[1]}\textsuperscript{[1]}. Moreover, any mechanism ad-
ressing the singularity problem (if this mechanism exists
at all) should operate on this time scale.

This fact was not noticed in\textsuperscript{[1]}\textsuperscript{[1]} because the authors
concentrated on the energy density at the end of the
ekpyrotic stage,

\[ \rho \approx H^2 \approx \frac{4}{c^2 t_{\text{ekp}}} \approx \frac{10^{100}}{c^4}, \]

(12)
in Planck units. They studied two different regimes, $c = 10^{28}$ and $c = 10^{40}$. In both cases $\rho \ll 1$, so one could think that we are safely in the sub-Planckian regime. However, at the ekpyrotic stage with $c^2 \gg 1$ one has $p \sim c^2 \rho \gg \rho$, so from the Einstein equation $\dot{H} = - (\rho + p)/2$ it follows that

$$p \approx -2\dot{H} \sim c^2 H^2 \approx \frac{4}{c^2 H^2_{\text{ekp}}} \sim \frac{10^{100}}{c^2}, \quad (13)$$

i.e. pressure $p$, as well as $\dot{H}$ and the curvature $R_t$, are more than 20 orders of magnitude greater than the Planck density for both of the values of the parameters discussed in [1].

The situation with higher-order curvature invariants is even worse. For instance,

$$\sqrt{|R_{\alpha\beta}R^{\alpha\beta}|} \sim \dot{H} \sim \frac{1}{c^2 H^2_{\text{ekp}}} \sim \frac{10^{150}}{c^2}.$$  

This means that they exceed the Planckian values by 94 orders of magnitude for $c = 10^{28}$, and by 70 orders of magnitude for $c = 10^{40}$. If one considers even higher invariants, the situation becomes more and more troublesome. This is a direct consequence of the abnormal smallness of the time at the end of the stage of the ekpyrosis [11], which happens in this scenario for all values of its parameters.

One can easily check that the time $t_{\text{end}}$, when the pre-ekpyrotic stage ends and the stage of the ekpyrosis begins in this scenario, is also many orders of magnitude smaller than the Planck time, for the values of the parameters used in [1]. Moreover, one can show that this conclusion is valid for all possible values of parameters compatible with the requirement $cH_0 \sim 10^{-3}$ unless there is a long stage of inflation after the ekpyrotic big crunch.

Note that we are not talking here about hypothetical processes which are supposed to lead to a bounce and prevent the cosmological singularity in this scenario. The pre-ekpyrotic stage, as well as the ekpyrotic stage, are described in [1] by standard methods of classical field theory and general theory of relativity, which are not valid at $|t| < t_p$ and at super-planckian values of curvature. Our results show therefore that the suggested scenario is completely unreliable even as a background model. In the next section we will show that even if one ignores all of the problems discussed above, the calculation of the spectrum of metric perturbations in [1] is also unreliable.

### III. PERTURBATIONS

The perturbed homogeneous flat universe in the conformal-Newtonian coordinate system is described by the metric

$$ds^2 = a^2(\eta) \left[ (1 + 2\Phi) d\eta^2 - (1 - 2\Phi) \delta_{ik} dx^i dx^k \right], \quad (14)$$

where $\Phi$ is the gravitational potential induced by the perturbations of the scalar field $\delta \phi$. Linear perturbation theory is reliable only if $\Phi$, as well as the energy density perturbations $\delta T^0_0/T_0^0$, are simultaneously much smaller than unity. Otherwise the higher-order terms in perturbation theory dominate over linear terms, and the perturbative expansion fails.

Before going any further, let us illustrate the statements made above using inflationary perturbations as an example. During inflation, the large-scale energy density perturbations and perturbations of metric are small $\delta T^0_0/T_0^0 \sim -2\Phi < 10^{-5}$.

At the transition to the post-inflationary stage, the value of $\Phi$ grows to about $10^{-5}$ and stabilizes at this level. Meanwhile the energy density perturbations $\delta T^0_0/T_0^0$ continue to grow. Once these perturbations become large, $\delta T^0_0/T_0^0 \sim O(1)$, galaxies begin to form and separate from the cosmological expansion. The subsequent evolution of the large-scale structure of the universe, as well as the related evolution of $\Phi$, cannot be described by perturbation theory. However, this happens long after the end of inflation, which is why the standard theory of generation of inflationary perturbations is reliable.

Ref. [1] did not contain any investigation of the perturbations $\delta T^0_0/T_0^0$, which is necessary to examine validity of their results. As we will see, in their model, the perturbations $\delta T^0_0/T_0^0$ are large already at the pre-ekpyrotic stage, and therefore one cannot apply perturbation theory for the investigation of the generation of perturbations in the metric and in the energy density in this scenario.

To quantize scalar perturbations one introduces the canonical variable [7]

$$v = a \left( \delta \phi + \sqrt{3(1+w)} \Phi \right), \quad (15)$$

which for every Fourier mode $k$ satisfies the equation

$$v''_k + \left( k^2 - \frac{z''}{z} \right) v_k = 0, \quad (16)$$

where prime means the derivative with respect to conformal time $\eta$, and $z = a \sqrt{3(1+w)}$. In the case under consideration

$$\frac{z''}{z} \approx \frac{2}{l^2}, \quad (17)$$

during the time interval $t_{\text{begin}} < t < t_{\text{end}}$. During the pre-ekpyrotic stage $a \approx 1$, but the second derivative of $a$ can be large and hence before identifying conformal and physical time we have to perform all differentiation in our equations, and only after that it is safe to set $t \approx \eta$. For the short-wavelength perturbations with $k |t| \gg 1$

$$v_k \approx \frac{1}{\sqrt{k}} e^{ik\eta} \approx \frac{1}{\sqrt{k}} e^{ikt}, \quad (18)$$
where the amplitude is fixed by the requirement that the perturbations are the minimal vacuum fluctuations. In the long-wave limit for \( k |t| \ll 1 \) the solution of (16) is

\[
v_k \simeq C_1 z + C_2 z \int \frac{dn}{z^2} \simeq \frac{1}{\sqrt{k^2 t}} ,
\]

(19)

where we have fixed the constant of integration requiring the continuity of the solution at the moment of “effective horizon” crossing, \( t_k \simeq 1/k \), and have neglected the decaying mode. It follows from here that for the long-wavelength perturbations the “conserved” \( \zeta \) is equal

\[
\zeta_k \equiv \frac{v_k}{z} \simeq \frac{c}{\sqrt{k} t_b} \simeq \frac{cH_0}{k^{3/2}} .
\]

(20)

Therefore at the end of the pre-ekpyrotic stage the spectrum

\[
\zeta_k k^{3/2} \simeq cH_0
\]

(21)

is scale-invariant in the range of wavelengths \( |t_{\text{end}}| < \lambda < |t_{\text{begin}}| \). For \( cH_0 \simeq 10^{-5} \) the amplitude of perturbations is of the right order of magnitude. This result was the reason for the claim in [1] that the pre-ekpyrotic stage can produce the perturbations similar to those produced by inflation. In this case the flat spectrum appears because of the very fast change of the equation of state in a contracting universe.

Let us find under which conditions the result obtained is valid at the end of the pre-ekpyrotic stage. With this purpose we will calculate the gravitational potential and energy density perturbations. It is convenient to introduce the variable \( u \) related to the gravitational potential as

\[
\Phi = H (1 + w)^{1/2} u ,
\]

(22)

and satisfying the equation

\[
u_k'' + \left( k^2 - \frac{(1/z)''}{1/z} \right) u_k = 0 .
\]

(23)

During the pre-ekpyrotic stage

\[
\frac{(1/z)''}{1/z} \simeq - \frac{H_0}{t} .
\]

(24)

Therefore, for perturbations with \( k^2 |t| \gg |H_0| \), we have

\[
|u_k| \propto e^{ikt} ,
\]

(25)

while for \( k^2 |t| \ll |H_0| \) the solution is

\[
|u_k| \simeq C_1 + C_2 t .
\]

(26)

Using the 0 – i Einstein equation,

\[
\Phi + H \Phi = \frac{1}{2} \dot{\Phi} = \frac{\sqrt{3}}{2} H (1 + w)^{1/2} \delta \Phi ,
\]

(27)

to express \( \delta \phi \) in terms of \( \Phi \) in (15), we obtain [7]

\[
\zeta = \frac{v}{z} = \frac{2 \Phi + H \Phi}{3 H (1 + w)} + \Phi ,
\]

(28)

where the dot denotes the derivative with respect to the physical time \( t = \int \dd{\eta} \). Substituting (22) and using the background equations of motion, according to which

\[
\frac{\dot{H}}{H^2} = - \frac{3}{2} (1 + w) ,
\]

(29)

we find

\[
\zeta = \frac{2}{3} \left( \frac{1}{2} \frac{\dot{\Phi}}{H (1 + w)^2} + \frac{\dot{u}}{H (1 + w) u} + \frac{1}{1 + w} \right) \Phi .
\]

(30)

It follows from here that

\[
\zeta \simeq - \frac{2}{3} \frac{t}{t_{\text{end}}} \Phi ,
\]

(31)

for the perturbations with \( k |t| \ll 1 \) during the pre-ekpyrotic stage. For deriving this result we used (25) and (23). Thus at the end of the pre-ekpyrotic stage \( \zeta \approx \Phi \) and hence the spectrum of the gravitational potential is the same as that for \( \zeta \).

The next thing to do is to determine the amplitude of the energy density perturbations and verify that they also remain small. Using the 0 – 0 Einstein equation, we find

\[
\left( \frac{\delta T_0^0}{T_0^0} \right)_k = - \frac{2k^2}{3H^2 a^2} \Phi_k - \frac{2}{H} \left( \dot{\Phi}_k + H \Phi_k \right) = \left( \frac{2k^2}{3H^2 a^2} + 3 (1 + w) - 2 - \frac{\dot{w}}{H (1 + w)} - \frac{2\dot{u}}{Hu} \right) \Phi_k .
\]

(32)

During the pre-ekpyrotic stage, this relation for perturbations with \( k > H_0 \) reduces to

\[
\left( \frac{\delta T_0^0}{T_0^0} \right)_k \simeq \left( - \frac{2k^2}{3H^2} + \frac{2}{H} \right) \Phi_k .
\]

(33)

This equation implies that at the end of the pre-ekpyrotic stage, at \( t \simeq t_{\text{end}} \)

\[
\frac{\delta T_0^0}{T_0^0} \simeq 2c^2 \Phi
\]

(34)

for \( H_0 < k \ll cH_0 \) and

\[
\left( \frac{\delta T_0^0}{T_0^0} \right)_k \simeq c^2 \Phi_k \left( \frac{k}{cH_0} \right)^2 \gg 2c^2 \Phi_k ,
\]

(35)

for \( k \gg cH_0 \).

Thus we see that for the observed value \( \Phi \sim 3 \times 10^{-5} \) the perturbations of the energy density exceed unity unless \( c \lesssim 10^2 \). Note that the large \( c \) approximation used in [1] requires \( c \gg 1 \). Eq. (35) implies that even in the
marginal case $1 \ll c \lesssim 10^2$ the energy density perturbations exceed unity at all scales smaller than $O(10^{-2})H_0^{-1}$ for $c \lesssim 10^2$. For $c \gtrsim 10^2$ we expect that the linear perturbation theory breaks down at all interesting scales. Let us show that this is what really happens in the case under consideration.

From (27) we find that during the pre-ekpyrotic stage
\[ \delta \phi \simeq c\Phi, \] for perturbations with $k|\phi| \ll 1$. Expansion of the energy momentum tensor around the background contains terms of all orders in $\delta \phi$. In particular, in the expansion of $\delta T_{\theta 0}^0$ we have
\[ \delta T_{\theta 0}^0 = V_{,\phi} \delta \phi + \frac{1}{2} V_{,\phi}^2 \delta \phi^2 + \ldots \] (37)

It is clear that linear perturbation theory is applicable only if the quadratic term here is small compared to the linear term. Let us find when this may happen. Taking into account that in the case under consideration $V_{,\phi} = cV_{,\phi}$ we find that
\[ \frac{V_{,\phi} \delta \phi^2}{V_{,\phi} \delta \phi} = c \delta \phi \simeq c^2 \Phi. \] (38)

This ratio does not exceed unity only if $1 \ll c \lesssim 10^2$, in agreement with the conclusion which we just obtained by a different method. Thus linear perturbation theory is not applicable for $c \gtrsim 10^2$, and it fails completely for $c = 10^{38}$ and for $c = 10^{40}$ studied in [1].

Note that this problem is much more severe than the potentially curable problem of anomalously large non-gaussianity discussed in [1]. The nongaussianity problem appears only for extremely small wavelengths $k^{-1} > c^{-1}H_0^{-1}$, where $c$ can be as large as $10^{28}$ or $10^{40}$. Meanwhile for $c \gtrsim 10^2$ the problems discussed in our paper occur at all wavelengths.

Moreover, as we already mentioned, the spectrum is not satisfactory even in the marginal case $1 \ll c \lesssim 10^2$. Eq. (35) implies, in particular, that in the limiting case of $c = O(1)$, when the large $c$ approximation used in [1] breaks down, the spectrum of perturbations of density is blue with $n_s \sim 5$, and its amplitude becomes greater than unity on scales two orders of magnitude smaller than $H_0^{-1}$. In addition, according to [12], in this case one would be forced to consider the ekpyrotic stage ending at the density exceeding the Planck density by 100 orders of magnitude.

IV. CONCLUSIONS

As we have seen, the new version of the ekpyrotic scenario requires calculations to be performed at curvatures which are exponentially greater than the Planck curvature. Moreover, the theory of generation of metric perturbations used in this scenario is based on calculations which are exponentially far away from the domain of applicability of this theory. These problems appear independently of the major problem of this scenario, which is the problem of the cosmological singularity.

Each time when a new version of the ekpyrotic/cyclic theory is proposed, the authors suggest that its predictions be compared with predictions of inflationary cosmology to check which of the theories is better with respect to observations. In accordance with this tradition, the authors of Ref. [1] say that their new theory "predicts non-gaussianity and a spectrum of gravitational waves that is observationally distinguishable from inflation." The main problem with this suggestion is that one can test any theory only if it is internally consistent and only if it really makes definite predictions.

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It was suggested in [22] that one might avoid the problems of the new ekpyrotic scenario discussed in [20] by using an approach to the effective field theory with higher derivatives proposed in [23]. Ref. [23] suggested to use iterations, initially solving equations of the theory omitting the higher derivative terms, and then substituting the solutions to the higher derivative terms. This proposal works for a large class of theories where omission of the higher-derivative terms does not lead to pathologies. However, it does not help in the case of the new ekpyrotic scenario because this scenario, with the higher derivative being omitted, exhibits a catastrophic gradient instability [24].

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