Generating the curvature perturbation with instant preheating

Tomohiro Matsuda
Laboratory of Physics, Saitama Institute of Technology, Fusaiji, Okabe-machi, Saitama 369-0293, Japan
E-mail: matsuda@sit.ac.jp

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Abstract. A new mechanism for generating the curvature perturbation at the end of inflation has been investigated. The dominant contribution to the primordial curvature perturbation may be generated during the period of instant preheating. The mechanism converts isocurvature perturbation related to a light field into curvature perturbation, where the ‘light field’ is not the inflaton field. This mechanism is important in inflationary models where kinetic energy is significant at the end of inflation. We show how one can apply this mechanism to various brane inflationary models.

Keywords: string theory and cosmology, inflation, cosmological applications of theories with extra dimensions

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1. Introduction

In the standard scenario of the inflationary Universe, the observed density perturbations are assumed to be produced by a light inflaton that rolls down its potential. This ‘standard scenario’ has been investigated by many authors [1]. Despite the extensive work, we know that in supersymmetric and superstring theories there is a serious problem called the ‘η-problem’. This problem is famous in supersymmetric theories, but it can appear whenever higher corrections lift the inflaton potential [2]. Although it may be possible to construct some inflationary scenarios where the flatness of the inflaton potential is protected by symmetry, it is not always the norm to find a situation where the symmetry appears naturally and all the required conditions for inflation are satisfied without any fine-tuning1. Recently, a new inflationary paradigm has been developed where the conventional slow-roll picture does not play an essential role in generating the curvature perturbation. Along the lines of this ‘new inflationary paradigm’, we consider in this paper a scenario where isocurvature perturbation of a light field is converted into the curvature perturbation at the time of instant preheating when the inflaton kinetic energy is significant at the end of inflation. It is significant to note that we will consider a scenario where a ‘light field’ is not identified with the inflaton. The ‘light field’ is decoupled from the inflationary dynamics, but it plays a significant role at the end of inflation. The idea of a ‘new light field’ has already been investigated by many authors. The most famous examples of this would be the curvaton models [4,5]. In the curvaton models, the origin of the large-scale curvature perturbation in the Universe is the late decay of a massive scalar field that is called the ‘curvaton’. The curvaton paradigm has attracted quite a bit of attention because it is thought to have obvious advantages. For example, since the curvaton is independent of the inflaton field, there was a hope [6] that the curvaton scenario, especially in models with a low inflationary scale, could cure serious fine-tunings associated with the inflation models.

1 It would be very fascinating if slow-roll inflation could be embedded in MSSM [3].
Many attempts have been made to construct realistic models with low inflationary scale [7]. Cosmological defects may play an essential role in low-scale inflationary scenarios [8], but the ultimate solution has not been found yet. Fast-roll inflation in the standard-model throat could be one of the successful models following this approach, as we will discuss later in this paper. However, Lyth has suggested [2] that there is a strong bound for the Hubble parameter during inflation, even when the curvatons are introduced to a model. The bound obtained by Lyth [2] was a critical parameter in the inflationary model with a low inflation scale, but it was later suggested by Matsuda [9] that the difficulty could be avoided if an additional inflationary expansion or a phase transition was present [10]. The idea of a ‘light field’ was investigated most recently by Lyth [11], where an argument was presented that the density perturbations can be generated ‘at the end of inflation’ by the fluctuations of the number of e-foldings induced by the fluctuations related to a light scalar field other than the inflaton\(^2\). The generation of the curvature perturbation is now due to the fluctuations related to a light field, which is independent of the inflaton dynamic details, and thus there could be a similarity between the curvatons and Lyth’s new mechanism in the usage of a light field. However, there is an important advantage in the new mechanism in that, unlike the curvatons, one does not have to worry about the stringent conditions that come from the requirement of the successful late-time dominance and reheating. Using this new idea, we studied the generation of the curvature perturbation without using the slow-roll approximations [13,14]. The situation we considered is very common in the brane inflationary models\(^3\). We considered a brane inflationary model where a light field appears ‘at a distance’ from the moving brane. This idea fits in well with the generic requirements of the models with ‘throat’ inflaton. To be more precise, we demonstrated that the correction lifting the moduli space in the bulk is not a serious problem if there is a symmetry enhancement ‘at the tip’ of the KS throat [16]. However, it is still not clear if the mechanism can work in a model where the inflationary brane has (ultra-) relativistic velocity, such as in the cases of DBI inflation or trapped inflaton. Naively using Lyth’s mechanism in these cases is inappropriate, since the hybrid-type potential does not play a significant role. Moreover, it is unclear what happens if the expectation value of the light field (i.e. the impact parameter) is so large that the inflaton cannot directly hit the waterfall region of the hybrid-type potential. If the brane motion is ultra-relativistic, or the kinetic energy of the inflaton is significant, the so-called ‘instant preheating’ should be important in the analysis. Our new mechanism for generating the curvature perturbation is very useful because this mechanism can be used to discuss the generation of the curvature perturbation in inflationary models ‘without’ branes. Thus, what we will consider in this paper is a generic mechanism for generating the curvature perturbation at the end of inflation that works with instant preheating and a light field. Construction of a successful brane inflationary model is one of our motivations, but there is no reason we have to stick to the brane Universe, as our mechanism is applicable to generic inflationary models.

In our scenario, as we have discussed previously [14], it is better to lower the scale of inflation so that one can obtain a large number of e-foldings without slow-roll. On the other hand, the requirements from the generation of the curvature perturbation place a

\(^2\) See also [12], where different models generating a contribution to the curvature perturbation were proposed.

\(^3\) Riotto and Lyth gave another useful discussion of this point [15].
lower bound on the inflationary scale. This problem can be solved if there is a preceding short period of inflationary expansion, the solution of which is similar to the mechanism discussed by us for the curvatons [9]. Assuming that the fluctuation of a light field is generated during the first inflationary expansion with the potential $V_1 \sim M_4^4$, and also that the instant preheating occurs at the end of secondary inflation with the potential $V_2 \sim M_4^2$, the curvature perturbation is enhanced by a factor of $O(M_4^2/M_4^1)$, which makes it possible to have low inflationary scales in both the first and the second inflationary epochs.

Aside from the generation of the curvature perturbation, one might think that the dangerous overproduction of the unwanted cosmological relics might be critical. For example, since the enhanced isometries are known to make the Kaluza–Klein modes completely stable, these ‘stable’ relics may put an unavoidable lower bound on the inflationary scale in the brane Universe [17]. Although this problem was already discussed in a similar context [14], it should be better to revisit this problem again in this paper to complete our analyses.

2. Generating the curvature perturbation with instant preheating

We consider a multi-field inflationary model where the background inflaton fields $\phi_i(t)$ evolve according to the system of coupled differential equations

$$\ddot{\phi}_i + 3H\dot{\phi}_i + \frac{\partial V}{\partial \phi_i} = 0, \quad i = 1, \ldots, n$$

and

$$H^2 = \frac{8\pi}{3M_p^2} \left[ \sum_i \frac{\dot{\phi}_i^2}{2} + V \right].$$

Without losing general applicability, we can discuss our mechanism with two fields, $\phi_1$ and $\phi_2$, where $\phi_1$ is the conventional inflaton field and $\phi_2$ is the ‘light field’. The potential $V(\phi_1, \phi_2)$ is characterized by a hierarchy between the masses of the fields, and can be modelled by

$$V(\phi_1, \phi_2) = \frac{m_1^2}{2} \phi_1^2 + \frac{m_2^2}{2} \phi_2^2,$$

where $m_1 \sim O(H)$ and $m_2 \ll m_1$.\(^4\)

We consider the instant preheating model [18] as the process through which the inflaton decays into lighter particles. The typical coupling to the preheat field $\chi$ is written as

$$\mathcal{L} = \frac{g^2}{2} (\phi_1^2 + \phi_2^2) \chi^2,$$

\(^4\) A different approach has been given by Kolb et al [22], who discussed a new general mechanism to generate curvature perturbation after the end of the slow-roll phase of inflation. This model is based on the simple assumption that the potential driving inflation is characterized by an underlying global symmetry which is slightly broken. See also table 1, where the mechanisms for generating the curvature perturbation at the end of inflation are categorized in four groups. We are mainly concerned about the $\eta$-problem and did not use the slow-roll approximations.
The generated energy density is proportional to the comoving number density produced during the decay process of the preheat field \( \chi \).

\[
\delta \phi \approx \delta \phi_2(t_*) \sim \frac{\phi_2(t_*)}{\phi_1(t_*)} \[
\]

where \( t_* \) is the time when the inflaton \( \phi_1 \) reaches its minimum potential at \( \phi_1 = 0 \) and where the light field \( \phi_2 \) may still have an expectation value \( \phi_2(t_*) \neq 0 \). We used \( \phi_2 = 0 \), \( \delta \phi_2 = 0 \) and \( \phi_1(t_*) = 0 \) to derive equation (2.5). To obtain an estimate of the curvature perturbation produced during the decay process of the preheat field \( \chi \), it is sufficient to note that the generated energy density is proportional to the comoving number density \( n_\chi \). Assuming a smooth decay process of the preheat field, the curvature perturbation \( \zeta \) generated during the instant preheating is

\[
\zeta \approx \frac{\alpha}{n_\chi} \frac{\delta n_\chi}{n_\chi},
\]

where \( \alpha \) is a constant whose numerical value depends on the redshift of the particle produced. Since the field \( \phi_2 \) is approximately massless during inflation, the value of the fluctuation is given by \( \delta \phi_2 \sim H_1 \). In the simplest case of a single-stage inflationary model, the curvature perturbation \( \zeta \) is approximately

\[
\zeta \approx \frac{\alpha g |\phi_2(t_*)|^2}{m |\phi_1(t_*)| |\phi_2(t_*)|} \sim \frac{H_1}{m} \left( \frac{\alpha g}{m} \right)
\]

where \( t_i \) is the time when the inflaton \( \phi_1 \) starts fast-rolling, and \( m_2 \) is used to derive the equation. Since the field \( \phi_2 \) is very light, it is possible to have \( \phi_2(t_i) \approx \phi_2(t_*) \). If there is the ‘eta-problem’, the natural scale of the mass \( m \) would be \( m^2 \geq O(H_1^2) \).

Moreover, since we are considering a case where kinetic energy is significant at the time of preheating,

\[
\dot{\phi}_1(t_*)^2 \sim m^2 |\phi_1(t_*)|^2 \sim H_i^2 M_p^2
\]

which gives a mass \( m_\chi = g\sqrt{\phi_1^2 + \phi_2^2} \) to the preheat field. Applying the result obtained in [18], the comoving number density \( n_\chi \) of the preheat field \( \chi \) produced during the first half-oscillation of \( \phi_1 \) becomes

\[
n_\chi = \frac{(g|\phi_1(t_*)|^3/2)}{8\pi^3} \exp \left[ -\frac{\pi g |\phi_2(t_*)|^2}{|\phi_1(t_*)|} \right],
\]

Table 1. Mechanisms for generating the curvature perturbation at the end of inflation.

| \( m_1 \gg m_2 \) | Lyth [11] | This paper |
| \( m_1 \sim m_2 \) | Matsuda [13] | Kolb et al [22] |

See the appendix for more details.

The mass of a light field \( \phi_2 \) is protected by symmetry, as we will discuss later.
Generating the curvature perturbation with instant preheating is a natural consequence. We must also consider the condition for the efficient production of the preheat field $\chi$, which is written as

$$m_\chi^2 \simeq g|\phi_2(t_*)|^2 < |\dot{\phi}_1(t_*)|^2.$$  \hfill (2.10)

As we will discuss later in section 3, the effective mass of the preheat field $\chi$ is suppressed by a Lorentz factor in DBI inflation, where the motion of the inflation becomes (ultra-) relativistic. We postpone the analyses of the ultra-relativistic motion.

Putting these conditions into (2.8), we obtained an order estimation

$$\zeta < \sqrt{H_I/M_p}.$$  

The above condition is not favourable for the brane Universe with a low scale. Is it possible to remove this condition without introducing fine-tunings? As we discussed in section 1, a similar problem has already been discussed in the curvaton models. A lower bound for the inflationary scale appeared in the curvaton model and was discussed by Lyth [2], and a solution was suggested by us [9]. Following the solution [9], we have attempted to introduce another inflationary stage that has the Hubble constant $H'_I > H_I$. Denoting the typical scale of the inflationary potential by $M$ and $M'$ for each inflation, and assuming that the relevant perturbation of the light field $\phi_2$ is generated during the preceding inflation with the Hubble constant $H'_I$, we obtained in a straightforward manner the order estimation of the curvature perturbation

$$\zeta \sim \left| \frac{\phi_2(t_*)}{\phi_1(t_i)} \right| \left( \frac{M'^2}{M^2} \right),$$  \hfill (2.11)

where $m \simeq H_I$ is assumed. The factor $(M'/M)^2$ helps the generation of the curvature perturbation in low-scale inflationary models.

3. Brane universe

The origin of a ‘light field’ could be the isometries in the KS throat, which is a useful idea that we have discussed in section 1. At least in the present model of the KS throat there are isometries at the tip of the throat, where the target brane is waiting for the inflationary (moving) brane.

If brane inflation is induced by a hybrid-type brane–antibrane potential, and also the kinetic energy of the inflationary brane is not significant, inflation ends with the waterfall of the conventional hybrid-type potential. In this case, one can use Lyth’s mechanism for generating the curvature perturbation at the end of inflation. Since the waterfall is essential to this model, one has to introduce mild fine-tunings in the initial conditions. To be more precise, the ‘light field’ should not have its expectation value larger than the string scale at the tip, which corresponds to the radius of the ‘waterfall area’ of the hybrid-type potential. The situation is shown schematically in figure 1. Moreover, the motion of the ‘inflaton field’ must not be ultra-relativistic so that the hybrid-type potential can play an essential role at the end of inflation.

On the other hand, we know it is possible to construct inflationary models in the brane Universe, which do not satisfy the above criteria. For example, as discussed by Alishahiha and Tong [19] the ultra-relativistic motion of the inflationary brane may play a crucial role in some inflationary models in the brane Universe. Another kind of inflaton was discussed by Kofman et al [20], in which trapping of a moving brane may induce weak inflationary expansion. It is important to note that the fast motion of an inflationary brane does play an essential role in these models, but the hybrid-type potential does not. Then, we have
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Figure 1. Brane–antibrane inflation ends with reheating induced by the waterfall field (tachyon), if kinetic energy is not significant. In this case, the mechanism advocated by Lyth [11] may play a significant role in generating the curvature perturbation at the end of inflation, where the fluctuation $\delta \phi_2$ is converted into $\delta N$. On the other hand, brane–brane collision or brane–antibrane collision with significant kinetic energy cannot be explained by conventional waterfall.

a natural question as to whether it is possible to generate the curvature perturbation at the end of inflation in these alternative inflation models. In this section we give an answer to the above question.

Related to the brane–brane or brane–antibrane collision at the end of inflation, two points should be clarified. One is that on instant preheating the mass of the $\chi$ field might reach a negative value if $\chi$ is the tachyon that appears in a brane–antibrane collision. If so, the collision does not correspond to the instant preheating scenario we referred to in section 2. The tachyon mass becomes negative when the distance between brane and antibrane becomes shorter than the critical length, hence the preheating must be calculated with the negative tachyon mass if instant preheating occurs ‘within’ the region $\phi_2 < M_s$, where $M_s$ denotes the string scale at which the mass of the $\chi$ field becomes negative. Hence, the underlying condition for the standard instant preheating scenario is (1) the collision is brane–brane (not brane–antibrane), or (2) the minimum length between branes at the preheating is longer than the critical length if the collision is brane–antibrane. Once the preheat field is produced at the collision, they will soon obtain huge mass. The other point is that, since we are using light open strings on the branes as the preheat field, the DBI action should be the starting point to see how the $\chi$ field is generated in instant preheating. More specifically, the concern is whether the higher derivative terms in the DBI kinetic term may or may not play distinguishable role in instant preheating. Although this issue seems interesting, we will not deal with this difficult issue in this paper. We will make a simplification to capture the characteristics of the brane collision [20, 21].
3.1. Trapped inflation

Consider the theory of a real scalar field $\phi_1$ with the effective potential $V_1 \simeq m^2 \phi_1^2/2$, where the mass is larger than the Hubble constant, and thus the inflaton $\phi_1$ falls rapidly to its minimum [20]. We will assume that $\phi_1$ falls from its initial value $\phi_{1,0} = \phi_{1*}(1 + \alpha)$ with vanishing initial speed and gives some boson (preheat field) a mass

$$m_\chi^2 \simeq g^2 ((\phi_1 - \phi_{1*})^2 + \phi_2^2),$$

(3.1)

where $\phi_2$ is a light field protected by symmetry. Then, as $\phi_1$ passes $\phi_{1*}$, it creates $\chi$ particles with number density $n_\chi \simeq ((gv)^{3/2}/8\pi^3) \exp(-\pi g \phi_2^2/v)$, where $v$ denotes the velocity of the field $\phi_1$. After a short time these $\chi$ particles become nonrelativistic and induce effective potential

$$V_{\text{eff}} \simeq \frac{1}{2} m^2 \phi_1^2 + g n_\chi \sqrt{(\phi_1 - \phi_{1*})^2 + \phi_2^2},$$

(3.2)

where $\sqrt{(\phi_1 - \phi_{1*})^2 + \phi_2^2} \simeq (\phi_1 - \phi_{1*})$ is a natural assumption. Subsequent expansion of the Universe dilutes the density of $\chi$ particles as $n_\chi \propto a^{-3}$, which eventually reduces the $n_\chi$-dependent correction to the effective potential. Then, the inflaton $\phi_1$ starts moving down again. Let us calculate the number of e-foldings assuming (for simplicity) $g \sim \alpha \sim 1$.

The number of e-foldings occurring during this ‘trapping’ period is [20]

$$N_e \sim \frac{1}{6} \ln \left( \frac{\phi_1}{m} \right) - \frac{1}{3} \left( \frac{\phi_2^2}{v} \right),$$

(3.4)

where the original assumption in [20] was that $\phi_2 \simeq 0$.

Let us consider a fluctuation of a light field $\phi_2$. The fluctuation of the number of e-foldings is induced by the fluctuation related to the light field $\phi_2$. Assuming $|\phi_2| \gg \delta \phi_2$, we obtained

$$\zeta \simeq \delta N \sim \frac{\phi_2}{v} \delta \phi_2,$$

(3.5)

where the curvature perturbation with a given wavenumber $k$ of cosmological interest is $\zeta(k) = \delta N(k)$.

3.2. DBI inflation

The ‘standard’ adiabatic perturbation generated in DBI inflation has been calculated by Alishahiha et al [19]. The result is

$$\zeta_{\text{DBI}} \simeq \frac{\sqrt{\pi} H^2}{\phi_1(t_i)},$$

(3.6)

where $g_s$ is the squared gauge coupling of the corresponding gauge dynamics. The basic idea of the ‘D-cceleration’ is given by the evolution of $\phi_1(t)$ which is fixed by the first-order Friedman equation. The above result is obtained when the dynamics of the probe
D3-brane is captured by the Dirac–Born–Infeld action coupled to gravity [19],
\[
S = \int \frac{1}{2} M_p^2 \sqrt{-g} R + \mathcal{L}_{\text{eff}} + \cdots
\]
where \(\mathcal{L}_{\text{eff}} = -\frac{1}{g_s} \sqrt{-g} \left[ f(\phi)^{-1} \left( 1 + f(\phi) g^{\mu\nu} \partial_\mu \phi \partial_\nu \phi \right)^{1/2} + V(\phi) \right].\) (3.7)
Alishahiha et al suggested in [19] that \(\dot{\phi}_1\) is then given by the formula\(^7\)
\[
\dot{\phi}_1 = -2g_s M_p^2 H'(\phi_1) \frac{\gamma(\phi_1)}{\gamma(\phi_1)} \simeq -f^{-1/2},
\]
where \(\gamma(\phi_1) = \sqrt{1 + 4g_s^2 M_p^4 f(\phi_1) H'(\phi_1)^2}.\) The important point is that for \(\gamma \gg 1\) the ‘speed limit’ dominates the dynamics of the scalar field \(\phi_1\) and leads to an interesting cosmological scenario which differs from the usual slow-roll scenarios. Here \(f(\phi_1)\) is the warp factor of the AdS-like throat which is given by
\[
f(\phi_1) \simeq \frac{\lambda}{\phi_1^2}.
\]
(3.8)

Now let us extend our result to apply to a DBI inflationary model discussed in [19]. Here we assume that inflation ends with instant preheating at \(\phi_1(t_*) > 0\), where a static brane is waiting for an inflationary brane. In backgrounds with \(\phi_1 \neq 0\), there are velocity-dependent corrections to the effective action, which shifts the effective mass
\[
m_{X|\text{DBI}} \propto \frac{m_X}{\gamma},
\]
(3.10)
where \(m_{X|\text{DBI}}\) is the mass of a string stretched from an inflationary brane to another brane waiting at the tip [19]. Hence \(\gamma\) appears in our result as
\[
\zeta_{\text{PR–DBI}} \simeq \frac{\delta n_X}{n_X} \simeq \frac{\alpha |\phi_2||\delta \phi_2(t_*)|}{\gamma^2 \phi_1(t_*)},
\]
(3.11)
where we assumed that \(\gamma\) is a slowly varying function of \(\phi_1\). Here we can express \(\gamma\) as [19]
\[
\gamma^2 \simeq \frac{4g_s^2 M_p^4}{\epsilon_{\text{DBI}}^2 \phi_1^4} \simeq \frac{4g_s \lambda M_p^2 m_i^2}{3\phi_1^2},
\]
(3.12)
where \(\epsilon_{\text{DBI}} \simeq \sqrt{3g_s/\lambda M_p m_i}\) is a ‘slow-roll’ parameter for DBI inflation [19]. Comparing equations (3.11) and (3.6), we find that the condition \(\zeta_{\text{DBI}}/\zeta_{\text{PR–DBI}} \gg 1\) that is needed for the ‘standard’ fluctuation \(\zeta_{\text{DBI}}\) to dominate over \(\zeta_{\text{PR}}\) is written as
\[
\frac{\delta \phi_2}{\phi_2} \gg \frac{\alpha}{\sqrt{g_s} \gamma^2(t_*) \phi_1(t_*)} \phi_1(t_*) \simeq \frac{\alpha}{\sqrt{g_s} \gamma^2(t_*) \phi_1(t_*)^2} \phi_1(t_*)^2 \simeq \frac{\phi_1(t_*)^2}{g_s \gamma^2 \lambda M_p^2 m_i^2}.
\]
(3.13)
\(^7\) See appendix B in [19].
where \( \delta \phi_2 \simeq H/2\pi \) is considered. Following [19] we suppose that \( \phi_1(t_\ast) \simeq m_1 \) and \( \lambda \sim g_s \times 10^{10} \). Then the above condition becomes
\[
\frac{\delta \phi_2}{\phi_2} \gtrsim g_s^{-1/2} \lambda^{-1} \sim g_s^{-3/2} \times 10^{-10}.
\] (3.14)

The crucial factor of \( 10^{-10} \) comes from a normalization of the curvature perturbation obtained in the original DBI inflationary model
\[
\zeta_{\text{DBI}} \simeq g_s^{1/2}/(\epsilon_{\text{DBI}}^{1/2} \lambda^{1/2}) \simeq \lambda^{1/2} m^2/(g_s^{1/2} M_p^2) \simeq 10^{-5}
\] (3.15)
and the assumption \( m \sim 10^{13} \) GeV that was made in [19]. In terms of the gravity dual, \( \lambda \) is given by \( \lambda \sim (R/l_s)^4 \), which is determined by the radius of the AdS space \( R \) and the fundamental string length \( l_s \). The above condition suggests that the ‘standard’ perturbation \( \zeta_{\text{DBI}} \) appears exclusively in the spectrum. Moreover, one can see that the dimensionless parameter \( \lambda \) and the mass scale \( m \) is constrained by the slow-roll condition \( \epsilon_{\text{DBI}} \leq 1 \), which is needed to obtain sufficient number of e-foldings during D-cceleration,
\[
N_c = \epsilon_{\text{DBI}}^{-1} \log \left( \frac{\phi_1(t_\ast)}{\phi_1(t_i)} \right).
\] (3.16)

Actually, the slow-roll condition \( \epsilon_{\text{DBI}} < 1 \) requires \( g_s/\lambda < (m/M_p)^2 \) while the amplitude of tensor modes requires \( m^2/M_p^2 < 10^{-10} \), which leads to the same conclusion without using the normalization of the primordial curvature perturbation. Hence our conclusion is that \( \zeta_{\text{PR-DBI}} \) is generically much smaller than \( \zeta_{\text{DBI}} \) if a sufficient number of e-foldings is generated during DBI inflation.

On the other hand, there might be a possibility that the required number of e-foldings is generated during precedent inflation in the bulk, while reheating is induced by a relativistic collision at the tip. Of course this scenario is not a DBI inflationary model since in this case inflation is not due to D-cceleration. One may think that this kind of scenario is more generic than the original DBI inflationary model. In this case, to obtain \( \zeta_{\text{PR-DBI}} \sim 10^{-5} \) one must choose \( \zeta_{\text{PR-DBI}} \simeq \phi_2 m_1/(\lambda^{1/2} M_p^2) \sim 10^{-5} \) for \( \dot{\phi}_1(t_\ast) \simeq f^{-1} \sim m_1^2/\lambda^{1/2} \) and \( \delta \phi_2 \simeq m_1 \) [19], which leads to a condition \( \phi_2 > \lambda^{1/2} M_p \). Hence the condition \( \zeta_{\text{PR-DBI}} \sim 10^{-5} \) requires large \( \phi_2 \) which is larger than the largest allowed value of \( \phi_2 \) \( (\sim g_s^{1/2} M_p) \) in [19]. The value \( \phi_2 > \lambda^{1/2} M_p \) is also larger than the largest impact parameter allowed by instant preheating.

Hence, our conclusion in this section is that the ‘standard’ calculation in DBI inflation cannot be accommodated by the curvature perturbations generated at the end of inflation. This result is true once the D-cceleration becomes crucial during or even after inflation. On the other hand, the situation becomes marginal if brane collision is not ultra-relativistic. In this sense, the previous attempts [11, 13, 22] and our present model may complement each other to give the paradigm of ‘generating the curvature perturbation at the end of inflation’, as is shown in figure 1.

4. Cosmological relics

Because of the \( \eta \)-problem, our new scenario favours low inflationary scale [14]. The question as to whether there are further constraints coming from other cosmological considerations with low inflationary scale requires an answer. Although there is general
knowledge of the various constraints in this direction, the constraints are sometimes based
on some specific conditions that are highly model-dependent. A generic condition that
places a significant bound on the scale of the usual hybrid-type inflationary model is
discussed by Lyth [23], who suggested that the inflaton-waterfall-field coupling induces
a one-loop contribution to the inflaton potential and ruins the slow-roll approximation
as well as the generation of the structure of the Universe. Therefore, one cannot have a
successful slow-roll inflationary model of a hybrid-type potential with a scale lower than
10^9 GeV. Our model does not suffer from this condition. Besides the above condition,
other conditions coming from the fact that the typical mass scale of the brane binding
the maximum value of a scalar field on the brane must be considered. If the fundamental
scale is as low as \( M \sim O(\text{TeV}) \), the inflation field must be a bulk field, and thus one
cannot exclude the serious constraint coming from the decaying Kaluza–Klein mode [24],
since such bulk fields always couple to the Kaluza–Klein states. The crucial point is that
the constraint appears even if one discards the slow-roll approximations. In relation to a
brane inflationary model, it has been suggested [17] that the constraint coming from the
Kaluza–Klein mode becomes much more serious in brane inflationary models with some
isometries in the compactified space, since the isometries stabilizes the unwanted KK
relics. This condition is quite significant in our scenario and excludes low-scale inflation
with \( M < 10^{12} \) GeV. More recently, however, there has been discussion [25] that in
studying a more detailed thermodynamic evolution of the heating process, especially that
of the KK particles, many qualitatively different results compared to the original result
can be found. Following the new result [25], we may consider throat inflation in the
standard-model throat with the inflationary scale \( \sim O(\text{TeV}) \), even if an isometry at the
tip of the throat plays the crucial role in our scenario. One may also worry about dangerous
cosmological ‘defects’ that may be produced during brane annihilation. It is known that
cosmic strings can put an upper bound on the inflationary scale of the last inflation [26],
but the bound is not significant in our scenario. Besides the cosmic strings, monopoles
and walls may be produced as a consequence of brane creation or brane deformation
that may occur during or after the reheating epoch [27]. It is known that these defects
might put a strict upper bound on the inflationary scale, if there is a non-trivial structure
in the compactified space. A natural solution to the domain wall problem in a typical
supergravity model is discussed by us [29], where the required magnitude of the gap in
the quasi-degenerated vacua is induced by \( W_0 \) in the superpotential. The mechanism
discussed [29] is quite natural since the constant term \( W_0 \) is the one that appears to
cancel the cosmological constant. The bound obtained in [28] is severe, but it is also
known that the structure that is required to produce the dangerous defect configuration
of the brane does not appear in the known example of the KS throat. Of course, it is not
clearly understood whether it is possible to obtain the complete standard model (SM) in
the known example of the KS throat. Thus, it is fair to conclude here that this problem is
unsolved and requires further investigation together with the construction of the complete
set of the SM model in the brane Universe.

In lowering the inflationary scale it might be thought that it would be difficult to
obtain enough baryon number asymmetry of the Universe (BAU), since the requirement
of the proton stability puts a strict bound on the baryon-number-violating interactions.
This speculation is indeed true. The old scenario of the baryogenesis with a decaying
heavy particle cannot work if the fundamental scale goes as low as the \( O(\text{TeV}) \) scale [30].
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One can solve this problem by introducing cosmological defects that enhance the breaking of the baryon number conservation in the core [31]. There is also a problem in the scenario of Affleck–Dine baryogenesis [32], since the expectation value of a field on a brane cannot become much larger than the typical mass scale of the brane [33]. One can solve this problem by introducing non-trivial defect configuration structures [34]. Moreover, there are more arguments about the mechanism of baryogenesis with low-scale inflationary scale [35]. Although there is no ultimate solution to this problem, it is not incorrect to expect that baryogenesis could be successful even if the inflationary scale is as low as the TeV scale.

5. Conclusions and discussions

A new mechanism for generating the curvature perturbation at the end of inflation is discussed in this paper. The dominant contribution to the primordial curvature perturbation may be generated by this new mechanism, which converts isocurvature perturbation of a light field into curvature perturbation during the period of instant preheating. The light field is ‘not’ the inflaton field. Our new mechanism is important in inflationary models where the kinetic energy is significant at the time of reheating. What we have presented in this paper is a new mechanism for generating the curvature perturbation at the end of inflation, which works even if (1) the kinetic energy is significant, or (2) the inflation potential is not a hybrid type. Thus, our present model complements the original scenarios [11,13,14,22] to give the paradigm of generating the curvature perturbation at the end of inflation.

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Appendix. Instant preheating in two-field models

The key equation in this paper is equation (2.5), which describes the number density of particles produced by instant preheating. However, the original result obtained in [18] was

\[ n_\chi = \frac{(g|\dot{\phi}_1(t_\star)|)^{3/2}}{8\pi^3} \exp \left( -\frac{\pi m_\chi^2}{g|\phi_1(t_\star)|} \right), \]

which applies only to a single-field model. Moreover, in a two-field model there must be some restriction on the assumed value of \( \phi_2 \) and its potential, which was not explicit in the above arguments. Therefore, in this appendix we will present in more detail the derivation of equation (2.5) and also some restriction on \( \phi_2 \) and its potential, which applies to a two-field model. In some cases one may find a sensible amount of non-Gaussianity [36]. A standard adiabatic curvature perturbation generated by the fluctuation in the inflaton

\footnote{A conserved angular momentum may be significant in some other cases [37,38].}
Hence, the comoving number density of the background field dynamics is then given by 9 curvature perturbations after the end of the slow-roll phase of inflation. Let us first consider what happens if both \( \phi_1 \) and \( \phi_2 \) are dynamical at the same time. Here we will assume that \( \phi_1 \) is heavier than \( \phi_2 \), and the \( \phi_1 \)-oscillation starts when \( H \sim m_1 \). If the condition \( \dot{\phi}_1 > \pi g |\phi_2|^2 \) is satisfied at the bottom of the \( \phi_1 \)-potential, the instant preheating terminates the \( \phi_1 \)-oscillation before \( \phi_2 \) starts to oscillate. In this case one may assume that the value of \( \phi_2 \) is a constant during the instant preheating. We have considered this situation in this paper. On the other hand, if the above condition is not satisfied for some time during the \( \phi_1 \)-oscillation, the \( \phi_1 \)-oscillation continues until it loses its energy. The Hubble constant decreases with time as the amplitude of the \( \phi_1 \)-oscillation decreases, and finally \( \phi_2 \) starts to oscillate when \( H \sim m_2 \). Here we have assumed that the decay constant \( \Gamma_{\phi_1} \) is much smaller than \( m_2 \), otherwise the instant preheating cannot play a significant role in reheating. In this peculiar case, both \( \phi_1 \) and \( \phi_2 \) are equally important. The most obvious example in this direction is discussed by Kolb et al. \cite{22} who considered a model with a weakly broken \( U(1) \) symmetry and presented a mechanism to generate curvature perturbations after the end of the slow-roll phase of inflation. Let us make a short review of this scenario to complement our discussions in this paper. Here we assume that the potential takes the simple form

\[
V(\phi_1, \phi_2) = \frac{m^2}{2} \left[ \phi_1^2 + \frac{\phi_2^2}{(1 + x)} \right],
\]

where \( x \) is a small constant that measures the symmetry breaking. Given the potential (A.2), the solution for the background dynamics can be computed neglecting the expansion of the Universe. We can neglect the expansion of the Universe since after a first transient the trajectory in field space will be mostly along the radial direction, and also the instant preheating occurs immediately. Introducing a complex field \( \phi \equiv \phi_1 + i\phi_2 = |\phi|e^{i\theta} \), the background field dynamics is then given by

\[
\begin{align*}
\phi_1(t) &= |\phi_0| \cos \theta_0 \cos mt \\
\phi_2(t) &= |\phi_0| \sin \theta_0 \cos \left( \frac{mt}{1 + x} \right),
\end{align*}
\]

where the index 0 is for the initial conditions. Linearizing the trajectory \cite{22}, one can obtain a minimum distance \(|\phi(t_*)|\),

\[
|\phi(t_*)| = \frac{|\phi_0| \pi x}{2\sqrt{2}} \sin 2\theta_0.
\]

After a simple calculation \cite{22}, the velocity of the field is obtained as

\[
\dot{\phi}(t_*) \simeq m|\phi_0| \left( 1 - x \sin^2 \theta_0 \right)^{1/2}.
\]

The preheat particles \( \chi \) generated will be characterized by an effective mass \( m_\chi = g|\phi_*| \). Hence, the comoving number density of \( \chi \) particles produced by the instant preheating is

9 If the \( U(1) \) symmetry is strongly broken, the trajectory becomes a Lissajous curve. The instant preheating occurs when the trajectory satisfies the condition \(|\dot{\phi}(t_*)| < g|\phi(t_*)|^2\) at a minimum distance \( \phi(t_*) \), where \( \phi \) is defined as \( \phi^2 \equiv \phi_1^2 + \phi_2^2 \). Thus, in the case where the \( U(1) \) symmetry is strongly broken, and also if the motion of \( \phi_2 \) is not negligible, the calculation of the curvature perturbation requires numerical methods. This case may be interesting, but it is beyond the scope of this paper.
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given by

\[ n_\chi = \frac{(g|\dot{\phi}(t_*)|)^{3/2}}{8\pi^3} \exp \left[ -\frac{\pi g|\phi(t_*)|^2}{|\phi(t_*)|} \right]. \quad (A.6) \]

If \( \phi_2 \) is so light that it cannot move before the instant preheating, the trajectory of the oscillation becomes linear and one obtains equation (2.5) instead of equation (A.6).

Although in our mechanism slow-roll inflation is not required for generating curvature perturbations, it is important to consider a constraint that the curvature perturbation generated at the end of inflation should dominate over the standard inflaton perturbation at horizon exit. Fast-roll inflation is an obvious example, where the inflaton field has a mass of \( O(H) \). Since the inflaton is not light in fast-roll inflation, one may simply neglect the standard inflaton perturbation. On the other hand, the standard inflaton perturbation may be important if the large number of e-foldings is generated due to a ‘slow-roll’ inflaton. The constraint that the curvature perturbations generated by an instant preheating, which are given by equation (2.11), equation (3.5) or equation (3.11), should dominate over the standard inflaton perturbation \( \zeta_{st} \) is written as \( \zeta/\zeta_{st} \gg 1 \). This constraint depends crucially on the scenario of inflation, and requires further study.

In this paper we have assumed that the impact parameter \( \phi_2 \) is much larger than its fluctuation \( \delta \phi_2 \) that exits the horizon during inflation. This condition is very crucial to obtain the required perturbation spectrum. Once this condition is violated, higher-order terms that have been neglected in the above calculation may become significant, giving rise to a large non-Gaussianity [36]. Let us see more details. The expression for \( \delta n_\chi/n_\chi \), which was given by equation (2.6), has the additional terms proportional to \( (\delta \phi_2)^2 \), which are given by

\[ \frac{\delta n_\chi}{n_\chi} \simeq -\frac{2\pi g|\phi_2(t_*)|^2}{|\phi_1(t_*)|^2} \frac{\delta \phi_2(t_*)}{|\phi_2(t_*)|^2} - \frac{1}{2} \left( \frac{2\pi g|\phi_2(t_*)|^2}{|\phi_1(t_*)|^2} - \frac{4\pi^2 g^2|\phi_2(t_*)|^4}{|\phi_2(t_*)|^2} \right) \frac{\delta \phi_2(t_*)}{|\phi_2(t_*)|^2}. \quad (A.7) \]

The non-Gaussianity parameter, if observable, is

\[ -\frac{3}{5} f_{NL} = \frac{|\dot{\phi}_1(t_*)|^2}{4\pi g \alpha |\phi_2(t_*)|^2} - \frac{1}{2\alpha}. \quad (A.8) \]

Considering a condition \( m_1^2|\phi_1(t_1)|^2 \simeq |\dot{\phi}_1(t_1)|^2 \), non-Gaussianity becomes obviously large when \( \phi_2(t_*) < m_1 < \delta \phi_2(t_*) \).

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