Restoring the sting to metric preheating

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The relative growth of field and metric perturbations during preheating is sensitive to initial conditions set in the preceding inflationary phase. Recent work suggests this may protect super-Hubble metric perturbations from resonant amplification during preheating. We show that this possibility is fragile and extremely sensitive to the specific form of the interactions between the inflaton and other fields. The suppression is naturally absent in two classes of preheating in which either (1) the critical points (hence the vacua) of the effective potential during inflation are deformed away from the origin, or (2) the effective masses of fields during inflation are small but during preheating are large. Unlike the simple toy model of a $g^2\phi^2\chi^2$ coupling, most realistic particle physics models contain these other features. Moreover, they generically lead to both adiabatic and isocurvature modes and non-Gaussian scars on super-Hubble scales. Large-scale coherent magnetic fields may also appear naturally.

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I. INTRODUCTION

Standard, isentropic*, inflationary models must end with a phase of reheating during which the inflaton, $\phi$, transfers its energy to other fields. Reheating itself may begin with a violently nonequilibrium “preheating” era, when coherent inflaton oscillations lead to resonant particle production (see [2] and refs. therein). Until recently, preheating studies implicitly assumed that preheating proceeds without affecting the spacetime metric. In particular, causality was thought to be a “silver bullet,” ensuring that on cosmologically relevant scales, the non-adiabatic effects of reheating could be ignored.

In fact, exciting, super-Hubble effects are possible during preheating, and metric perturbations may be resonantly amplified on all length scales [3–5]. Causality is not violated [3] precisely because of the huge coherence scale of the inflaton immediately after inflation (see also [6]). Strong preheating (with resonance parameter $q \gg 1$; see [2,3] for overviews and notation) typically leads to resonant amplification of scalar metric perturbation modes $\Phi_k$, including those on super-Hubble scales (i.e., $k/aH \ll 1$, where $a$ is the scale factor and $H$ the Hubble rate). One of the aims of this Letter is to answer the question “how typical is typical?”

The answer is crucial since preheating can lead to distortions in the anisotropies in the cosmic microwave background (CMB). Observational limits rule out those models that produce unbridled nonlinear growth, but models which pass the metric preheating test on COBE scales may nevertheless leave a non-adiabatic signature of preheating in the CMB. Hence one can no longer universally avoid consideration of reheating when analyzing inflationary predictions for cosmology, even if the final effect of reheating in some particular models is small.

In this vein, it has been argued recently [7,8] that metric perturbations on super-Hubble scales are in fact immune to metric preheating in the archetypal 2-field potential typically used in earlier studies [2,3]. The claim arises because the initial value of the fluctuations in the created bosonic field $\chi$ at the start of preheating is much smaller than that used in [3]. The basic argument is as follows. For the coupling $\frac{1}{2}g^2\phi^2\chi^2$, strong preheating typically† requires $q \equiv g^2\phi^2/m^2 \gg 1$, which increases the effective $\chi$ mass relative to the Hubble rate during inflation, $m_{\chi,\text{eff}} \sim g\phi \sim H \sim m$, where $m$ is the inflaton mass. This leads to an exponential suppression $\propto a^{-3/2}$ of both $\chi$ and $\delta\chi_k$ during inflation; hence these fields would have values at the start of preheating around $\sim 10^{-36}$ smaller than those used in all previous simulations.

Since $\Phi_k$ depends on $\dot{\chi}\delta\chi_k$ [see Eq. (3) below], this would stifle any growth in the small-$k$ modes of $\Phi$ until late times. Initial conditions for large-$k$ modes, in contrast, are claimed to be unaffected, so that they would grow nonlinear first. Their resulting backreaction would then end the resonance before any interesting effects occur on cosmologically significant scales [7,8]‡. Note that

*Note that alternative scenarios with continuous entropy production, such as warm inflation, do exist [1].

†Exceptions exist in which $q$ is small but metric preheating is strong [5].

‡Note that discussion of the unrenormalized and unregularized variances $\langle \chi^2 \rangle$ should be treated with care, as they are
irrespective of super-Hubble behavior, non-perturbative effects are vital on smaller scales [7,11], and this in itself is a major departure from the old theory of perturbation evolution through reheating. It leads to interesting possibilities, such as significant primordial black hole formation [3] (see also [11,7]).

Returning to super-Hubble scales, \( k/aH \ll 1 \), we will show that the above suppression mechanism is highly sensitive to the particular form of interaction Lagrangian, while metric preheating is not. Indeed, the suppression of \( \chi \) and \( \delta \chi_k \) at the start of reheating argued for in [7,8] is absent for models in either of the following two classes:

Class I - Models in which the vacuum expectation value (vev) of \( \chi \) is nonzero during inflation, i.e., \( \langle \chi \rangle \neq 0 \) (see Sec. II).

Class II - Models in which the \( \chi \) effective mass is small during inflation but undergoes a transition and becomes large during preheating (see Sec. III).

Since these possibilities arise naturally in a variety of realistic particle physics models, we conclude that the suppression mechanism proposed recently [7,8] is fragile, i.e., unstable to small changes in the potential. On the other hand, resonant growth of super-Hubble metric perturbations in preheating is robust, since it persists under small changes of the potential.

Sec. II presents a general potential for Class I models. Sec. IIA gives analytical results showing how initial conditions are unsuppressed, while Sec. IIB presents full numerical simulations of one model in Class I. Sec. III describes realizations of Class II models while Sec. IV discusses new cosmological implications of metric preheating in general.

Background and perturbations

We envisage a model consisting of \( N \) minimally-coupled scalar fields, \( \phi_I \), schematically representing the particle content of the inflationary and preheating eras. More realistic models should consider the gauge group [e.g. \( SU(5) \), \( SO(10) \)], non-minimal coupling and fermionic effects, and of course an accurate phenomenology of metric preheating must begin to study these issues [12,9,10]. However, since we are interested only in essential conceptual points, this simple picture will suffice for now.

The inflaton, \( \phi \equiv \phi_1 \), is treated as a classical field and the other fields corresponding to \( l = 2, \ldots, N \) are assumed to be in their vacuum states near the end of inflation. Each of these fields is split into its homogeneous part and its gauge-invariant fluctuation (in the longitudinal gauge) as \( \phi_I (t, \mathbf{x}) = \phi_I (t) + \delta \phi_I (t, \mathbf{x}) \). The background equations are

\[
H^2 = \frac{1}{8} \kappa^2 \left[ V + \frac{1}{2} \sum k^2 \phi_I^2 \right] , \quad \dot{\phi}_I + 3H \dot{\phi}_I + V_I = 0 ,
\]

where \( \kappa^2 \equiv 8\pi/M_P^2 \) and \( V_I \equiv \partial V/\partial \phi_I \). The linearized equations of motion for the Fourier modes of field \( \langle \delta \phi_I k \rangle \) and scalar metric fluctuations \( \Phi_k \) are

\[
(\delta \phi_I k)^{\prime \prime} + 3H (\delta \dot{\phi}_I k) + (k^2/\ell^2) \delta \phi_I k = - \sum V_{IJ} \delta \phi_J k + 4 \dot{\phi}_I k \Phi_k - 2V_I \Phi_k ,
\]

\[
\Phi_k + H \Phi_k = \frac{1}{2} \kappa^2 \sum \dot{\phi}_I \delta \phi_I ,
\]

This system is subject to the constraint

\[
\left[ \frac{k^2}{a^2} - \frac{1}{2} \kappa^2 \sum \phi_I^2 \right] \Phi_k = - \frac{1}{2} \kappa^2 \sum \phi_I^2 (\delta \phi_I k/\dot{\phi}_I)^\prime ,
\]

which we use to check the accuracy of our numerical integrations of Eqs. (2) and (3) and to set \( \Phi_k \) initial conditions.

II. GENERAL POTENTIALS IN CLASS I

We now consider the inflaton \( \phi(t, \mathbf{x}) = \phi(t) + \delta \phi(t, \mathbf{x}) \), coupled to a massless scalar field \( \chi(t, \mathbf{x}) = X(t) + \delta \chi(t, \mathbf{x}) \). The often-used interaction term \( \frac{1}{2} g^2 \phi^2 \chi^2 \) is not the only coupling appropriate to preheating, but rather one is used for which resonance occurs. As we show below, additional couplings linear in \( \chi \), as well as quadratic couplings in which \( g^2 < 0 \), provide a mechanism for escaping the inflationary suppression claimed in [7,8]. Essentially, these alternatives produce a nonzero attractor \( X \neq 0 \), to which inflaton drives the \( \chi \) field, so that the initial values of \( X \) and \( \delta \chi, k-0 \) at preheating are not suppressed. These possibilities are incorporated in the effective potential

\[
V = \frac{1}{2} m^2 \phi^2 + \frac{1}{4} \lambda \phi^4 + \frac{1}{4} \lambda \chi \left( \chi^2 - s^2 \right)^2 + \frac{1}{4} \epsilon g^2 \phi^2 \chi^2 + g^2 \kappa^{-3} \phi^n \chi,
\]

for the constants \( \epsilon = \pm 1 \) and \( n = 2, 3 \). The \( \lambda, \lambda \chi \) terms ensure that \( V \) is bounded from below.

The various terms in this potential are phenomenologically well-motivated:

- In theories where supersymmetry (SUSY) is softly broken, the potential will only acquire logarithmic radiative corrections and the suppression may apply. However, in
realistic models with gravity, susy is replaced by supergravity (SUGRA), and SUGRA models often contain couplings of the form $\phi^n X$, $n = 2, 3$ [14].

- Even if $g = 0$ initially, if the $\chi$ field exhibits symmetry-breaking ($\sigma \neq 0$), shifting the field $\chi \to \chi - \sigma$ generates the linear term $g^2 \phi^2 \chi$ via the quadratic coupling. The possible importance of symmetry breaking of this sort has long been noted [15] for its role in generating single-body inflaton decays and hence complete inflaton decay. If we choose $\sigma$ to correspond to the grand unified theory (GUT) scale, then $\sigma/M_{pl} \sim g^2/g^2 \sim 10^{-3}$.

- Negative coupling instability (NCI) models ($\epsilon = -1$) are dominated by the coupling $-g^2 \phi^3 \chi^2$ and the $\chi$ field is driven to a nonzero VEV during inflation [16].

- A fermionic coupling $\bar{\psi} \chi \psi$, would lead to a driving term $h(\bar{\psi} \phi \psi)$ in the $\chi$ equation of motion. This would have a similar effect of giving a nonzero VEV for $\chi$. However, since we do not have the full perturbed Einstein equations for this case, we will not consider it further here.

Note that the potential (5) exhibits a $Z_2 \times Z_2$ symmetry when $g = 0$. If this symmetry is completely broken (and it is partially broken when $g \neq 0$) domain walls produced during inflation or through non-thermal symmetry restoration at preheating are unstable [19]. This offers the attractive possibility of solving the domain wall and monopole problems via defect-defect interactions [20].

### A. Unsuppressed initial conditions

We now present analytical arguments (assuming for simplicity that $\sigma = 0$) to show that the new couplings avoid the claimed suppression of super-Hubble $\chi$ fluctuations. By Eq. (1), the background $X$ field obeys

$$\ddot{X} + \frac{3}{2} H \dot{X} + \epsilon g^2 \varphi^2 X + \lambda X^3 = -\frac{g^2}{\kappa} \kappa^{-3} \varphi^n, \quad (6)$$

and by Eq. (2), its super-Hubble fluctuations satisfy

$$(\delta \chi_k)^{'} + 3H(\delta \chi_k) + [\epsilon g^2 \varphi^2 + 3 \lambda X^2] \delta \chi_k = 4X \Phi_k - [2 \epsilon g^2 \varphi X + n g^2 \kappa^{-3} \varphi^{-1}] \delta \phi_k \approx 2 \epsilon g^2 \varphi X + g^2 \kappa^{-3} \varphi^n + \lambda X^3 \Phi_k. \quad (9)$$

Using the slow-roll relation $\Phi \approx -\delta \phi/\varphi$ (for $\frac{1}{2} m^2 \phi^2$ inflation) to simplify the right-hand side yields

$$\ddot{\delta \chi_k} + 3H(\delta \chi_k) + [\epsilon g^2 \varphi^2 + 3 \lambda X^2] \delta \chi_k = 4X \dot{\Phi} + [(2 - n) g^2 \kappa^{-3} \varphi^{-1} + 2 \lambda X^3/\varphi] \delta \phi_k \quad (7)$$

We now consider the two separate cases with similar results:

**Case 1:** $\epsilon = 1, \bar{g} > 0$

Using the fact that $\varphi, H \approx$ constant during inflation, we see that while the solution of the homogeneous part of Eq. (6) decays rapidly towards zero as $a^{-3/2}$, the particular solution arising from the inhomogeneous term is approximately constant. For small $\lambda$, it follows that $\chi$ emerges at the end of inflation $(t = t_0)$ with homogeneous part

$$X(t_0) \approx - (\bar{g}/g)^2 \frac{\varphi}{\kappa} \kappa^{-3} |\varphi(t_0)|^{n/2}, \quad (8)$$

where $\varphi(t_0) \approx 0.3 M_{pl}$.

Similarly, the fluctuations have a decaying transient solution, but also a nontransient solution arising from the driving term on the right of Eq. (7). For $n = 2$, we need to include the small term $\dot{X} \Phi_k$, which is not straightforward to evaluate, but for $n = 3$ we can neglect this term, and Eq. (7) implies (again neglecting the $\lambda$ term):

$$\delta \chi_k(t_0) \approx - (\bar{g}/g)^2 \delta \phi_k(t_0). \quad (9)$$

Thus the super-Hubble $\chi$ fluctuations emerge from inflation unsuppressed, though smaller than the inflaton fluctuations by a factor $(\bar{g}/g)^2$.

**Case 2:** $\epsilon = -1, \bar{g} = 0$

An NCI coupling gives rise to a non-zero VEV, since Eq. (6) has an attractor solution ($\dot{X} \to 0$)

$$X = g \varphi/\sqrt{\lambda \kappa}. \quad (10)$$

Once again, $X$ does not decay exponentially to zero. The fluctuations governed by Eq. (7) have an approximate attractor solution

$$\delta \chi_k \approx (g/\sqrt{\lambda \kappa}) \delta \phi_k, \quad (11)$$

where we have used Eq. (10).

For consistency, inflation should be dominated by the $\frac{1}{2} m^2 \phi^2$ term in the potential, and the super-Hubble fluctuations should be dominated by adiabatic inflaton fluctuations. The background inflaton is governed by

$$\ddot{\varphi} + 3H \dot{\varphi} + m^2 \varphi + \lambda \varphi^3 = - \epsilon g^2 X^2 \varphi - n g^2 \kappa^{-3} X \varphi^{-1}, \quad \text{and standard} \frac{1}{2} m^2 \phi^2 \text{ slow-roll will be secured if } \lambda \text{ is negligible and } |\epsilon g^2 \bar{X}^2 n g^2 \kappa^{-3} \varphi^{-2} X| \ll m^2. \text{ Super-Hubble inflaton fluctuations obey} \quad (10)$$

$$(\delta \phi_k)^{'} + 3H(\delta \phi_k) + [\epsilon g^2 \varphi^2 + 3 \lambda X^2] \delta \phi_k = - [2 \epsilon g^2 \varphi X + n g^2 \kappa^{-3} \varphi^{-1}] \delta \chi_k \approx [2m^2 + \epsilon g^2 X^2 + n(3 - n) g^2 \kappa^{-3} \varphi^{-2} X - 2 \lambda \varphi^2] \delta \phi_k, \quad (7)$$

on using the slow-roll approximation. The above solutions for $X$ and $\delta \chi_k$ ensure that the inflaton fluctuations are not appreciably affected by the $\chi$ field during inflation.

In summary, our analytical arguments show that by the end of inflation, the $\chi$ field and its super-Hubble fluctuations are not negligibly small; the linear couplings ($\bar{g} > 0, \epsilon = 1$) and the negative quadratic coupling ($\epsilon = -1, \bar{g} = 0$) each provide a mechanism to evade the super-Hubble suppression of $\chi$ fluctuations.

3
B. Numerical simulations

In order to confirm and extend the analytical arguments above, we have performed numerical simulations in one Class I model, with the linear $\delta \phi^2 \chi$ coupling ($\epsilon = 1$). We will assume that $\lambda$ and $\chi$ are extremely small for simplicity. (See [5] for simulations with large $\chi$.)

To avoid subtleties associated with matching inflation to preheating, we numerically integrated Eqs. (1)–(3) starting inside the inflationary era. Our primary interest is in cosmologically relevant scales, so we follow the evolution of a scale that crosses the Hubble radius at $t = t_{in}$, about $55$ e-folds before $t = t_0$, the start of preheating.

The slow-roll approximation gives $N = \kappa^2 (\varphi^2_{in} - \varphi^2_0)/4$ for the number of e-folds before the end of inflation, so we choose $\varphi_{in} = 3M_p$ to get $N \approx 55$. For the background $\chi$ field we use the approximate dominant solution in Eq. (8) and take $X_{in} = -(\ddot{g}/g)^2 \varphi_{in}$. (This initial value is much smaller than that in [7] for the case $\ddot{g} = 0$.) We follow [7] and take the field fluctuations at Hubble-crossing ($k = aH$) as

$$|\delta \phi_{1k}|^2 = H^2/(2\pi^2 \omega_{1k}^2), \quad \langle \delta \phi_{1k} \rangle = \omega_{1k} \delta \phi_{1k},$$

(12)

where $\omega_{1k}^2 = (k/a)^2 + m_\chi^2$, with $m_\chi = g \varphi$. We also take $X_{in} = \omega_{1k} X_{in}$. The initial metric perturbation $(\Phi_0)_{in}$ is then fixed by Eq. (4). The comoving wavenumber is $k \approx m(a(t_0))e^{-N} \kappa \varphi_{in}/\sqrt{6}$. We also take $\ddot{\varphi}/g \leq 10^{-2}$, with $g = \sqrt{4\pi/3} \times 10^{-3}$ and $m = 10^{-6}M_p$. This yields a resonance parameter $q = 3.8 \times 10^5$ which is used for all our simulations here.

As well as tracking a scale that crosses the Hubble radius at $t_{in}$, we consider scales that are well within the Hubble radius at $t_{in}$. Although fluctuations on these small scales are not cosmologically significant, we need to compare their evolution with those on very large scales, since this has a bearing on the question of backreaction. The initial amplitude is given by $a^2_{in}|\delta \phi_{1k}|^2 = 1/(2\omega_{1k})$, and for $k/a_{in} \gg g\varphi_{in} \gg m$ we find that

$$|\delta \phi_{1k}(t_{in})| = 1/(a_{in}\sqrt{2k}).$$

The results are illustrated in Figs. 1 (large scale) and 2 (small scale) as functions of dimensionless time $mt$. The insets show the behaviour of perturbations during inflation. We have also plotted the curvature perturbation

$$\zeta = \Phi - H(\dot{\Phi} + H\Phi)/H = \sum \varphi_i H(\delta \phi_i + \varphi_i\Phi)/\sum \varphi_i^2,$$

(13)

which remains constant in standard reheating on super-Hubble scales in adiabatic models, but which clearly grows exponentially during the non-adiabatic preheating phase in these models.

![FIG. 1. Growth of $|\Phi|$, $|\dot{\varphi}|$, $\sqrt{4\pi/3}|\delta \chi_k|/M_{pl}$ and $\sqrt{4\pi/3}|\delta \phi_k|/M_{pl}$ as functions of dimensionless time $mt$ for the $\delta \phi^2 \chi$ coupling, with $\ddot{\varphi}/g = 10^{-2}$. The mode followed has $k/m_{a_{0}} \approx 10^{-23}$ and crosses the Hubble radius at $t_{in}$, 55 e-folds before the start of preheating; $q = 3.8 \times 10^5$. Inset: The same quantities during inflation.](image1)

![FIG. 2. As in Fig. 1, but for a scale that is well within the Hubble radius at $t_{in}$, with $k/m_{a_{0}} = 10^{-18}$ and $q = 3.8 \times 10^5$. Inset: The same quantities during inflation.](image2)
The results show how the time $t_{\text{nl}}$ strength $\tilde{\varphi}$ in $\Phi$ is affected by changes in the coupling $g_{\chi}$. This change in the $\text{cmb}$ print on the nant growth will typically take place, leading to an imprint on the CMB.

FIG. 3. Comparison of the time to nonlinearity, $t_{\text{nl}}$, of super-Hubble metric and field perturbations as $\tilde{g}/g$ is varied. $q = 3.8 \times 10^5$. Note the synchronisation of $\delta \phi_k, \delta \chi_k$, and $\Phi_k$, which all go nonlinear at the same time. Inset: a zoom of $m_{\text{nl}}$ for $\Phi_k$ with $\tilde{g}/g$ in the range $[10^{-4}, 10^{-3}]$. Note the highly sensitive dependence of $t_{\text{nl}}$ on $\tilde{g}/g$.

An indication of how the strength of the super-Hubble resonance in $\Phi$ is affected by changes in the coupling strength $\tilde{g}/g$ is given in Fig. 3. Here we have plotted the time $t_{\text{nl}}$ when the metric and field fluctuations grow to be nonlinear, e.g. $|\Phi_k(t_{\text{nl}})| = 1$, for $k/m_{\text{nl}} = 10^{-23}$. The results show how $t_{\text{nl}}$ increases in response to the suppression of initial conditions that occurs as $\tilde{g}$ is decreased. Note that synchronization occurs: all fluctuations share roughly the same $t_{\text{nl}}$ values.

### III. CLASS II MODELS

In the models of Class II, the $\chi$ effective mass is simply very small during inflation but then becomes large at preheating. This change in the $\chi$ effective mass occurs naturally in various models.

#### IV. COSMOLOGICAL EFFECTS

Our eventual goal must be to calculate physical quantities such as the power spectrum of $\Phi_k$. Since $P_\Phi = k^3|\Phi_k|^2/2\pi$, one might be concerned that these strong preheating effects at $k \rightarrow 0$ would be made irrelevant by the $k^3$ phase space factor. Perhaps the easiest way to see that this is not so is to look at the evolution of $\zeta_k$. Since $\zeta_k$ is not conserved for small $k$ (see Fig. 1), the standard normalization of the CMB spectrum is drastically increased. This could only take place if the power spectrum of metric fluctuations is strongly affected as $k \rightarrow 0$. This is understandable since preheating acts only as a huge increase in the value of the resonance parameter $q$.

- Consider globally susy hybrid models based on the superpotential $W = aS\tilde{\varphi}^2 - \mu^2 S$ [17]. Here $S$, a singlet, plays the role of the inflaton. The corresponding unbroken potential is

$$V = \alpha^2 |S|^2 (|\varphi|^2 + |\tilde{\varphi}|^2) + |\alpha \varphi \tilde{\varphi} - \mu^2|^2,$$

(together with $D$-terms which vanish along the flat direction $|\varphi| = |\tilde{\varphi}|$). For $S \gg \mu/\sqrt{\alpha}$, inflation occurs with the minimum of the potential at $\langle \varphi \rangle = \langle \tilde{\varphi} \rangle = 0$. However, for $S \leq \mu/\sqrt{\alpha}$, $V$ has a new minimum at $\langle \varphi \rangle = \mu/\sqrt{\alpha}$ and preheating occurs via oscillations around this minimum [18]. Now let us couple $\chi$ not to the inflaton $S$, but to the field $\varphi$ through the term $g^2 \chi^2 |\varphi|^2$. Then the $\chi$ effective mass $g|\varphi|$ vanishes during inflation (up to logarithmic corrections). It only departs strongly from zero once inflation ends and reheating begins leading to a huge increase in the value of the resonance parameter $q$.

- Consider models which have strong running of coupling constants with energy [21]. If the corresponding beta function of the theory is negative, such as occurs in QCD, the theory is asymptotically free and the coupling increases at lower energies. Perhaps the strongest examples of this are based on S-type dualities, where the coupling $g^2$ is very small during inflation but is very large during reheating which occurs in the strongly coupled phase with dual coupling $\propto 1/g^2 \gg 1$. An example is provided by ‘dual inflation’ [22], based on the susy Seiberg-Witten model. Here $m_{\chi,\text{eff}} \sim \phi\tilde{\phi} < H$, and $\chi$ fluctuations are similar to those in the inflaton, and not strongly suppressed. In fact, it is arguable that models of this sort are needed if preheating is to be viable in non-susy theories, since large $g$ leads to radiative corrections to the potential which may violate the slow-roll conditions for inflation.
scales. Further, because the metric perturbations go non-linear, whether on sub- or super-Hubble scales, the corresponding density perturbations $\delta$ must typically have non-Gaussian statistics. This is simply a reflection of the fact that $\delta \in [-1, \infty)$, so that the distribution of necessity becomes skewed and non-Gaussian. Further, in Class II models, where $\langle \chi \rangle = 0$ during inflation, $\chi$ perturbations in the energy density will necessarily be non-Gaussian ($\chi^2$ distributed), even if $\delta \chi_k$ is Gaussian distributed since stress-energy components are quadratic in the fluctuations (see e.g. [23]). Non-Gaussian effects are therefore an intrinsic part of many metric preheating models (particularly those in Class II), and open up a potential signal for detection in future experiments.

The second new feature we would like to identify is the breaking of conformal invariance. Once metric perturbations become large on some scale, the metric on that scale cannot be thought of as taking the simple Friedmann-Lemaître-Robertson-Walker (FLRW) form, and conformal invariance will be lost. This is particularly important for the production of primordial magnetic fields, which are usually strongly suppressed due to the conformal invariance of the Maxwell equations in a FLRW background. We naively expect the production of photons to be dominant around the scale at which the breaking of conformal invariance is most manifest, namely the scale $k$ at which $\Phi_k$ is largest. This too may yield, after detailed analysis, a specific signature of preheating.

The coherent oscillations of the inflaton during preheating further provide a natural cradle for producing a primordial seed for the observed large-scale magnetic fields [26]. A charged inflaton field, with kinetic term $D_\mu \phi (D^\mu \phi)^*$, will couple to electromagnetism through the gauge covariant derivative $D_\mu = \nabla_\mu - ieA_\mu$. This will naturally lead to parametric resonant amplification of the existing magnetic field, which will produce large-scale coherent seed fields on the required super-Hubble scales without fine-tuning.

V. CONCLUSIONS

This concludes our survey of some simple models which avoid the suppression of metric preheating effects. The suppression discussed in [7,8] is highly sensitive to the form of the particle interactions considered; when couplings are considered which are found in most realistic particle physics models, the effects of [7,8] receive. Instead, in models from either of the two general classes highlighted here, preheating will generically produce a strong amplification of metric perturbations on cosmologically-significant scales. The robustness of this amplification further demonstrates the need to move towards more realistic models of preheating in order to understand the predictions of inflation for observational cosmology.

Preheating yields the possibility of inducing a post-inflationary universe with both isocurvature and adiabatic modes on large scales. If these are uncorrelated and of roughly equal strength, the corresponding Doppler peaks will tend to cancel [24,25]. This mechanism is independent of the one discussed in [3], which requires nonlinearity to persist until decoupling. However, one might suspect that the adiabatic and isocurvature modes would be strongly correlated, leaving the possibility of a “smoking gun” fingerprint of preheating. The challenge remains to distinguish such correlations from those induced in hybrid inflation [25]. These exciting possibilities are under consideration.

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5A tiny seed field must exist during inflation if for no other reason than the existence of the conformal trace anomaly [27] or one-loop corrections to the QED action in a curved spacetime [26]. We thank Giuseppe Pollicino for discussion on this point.


