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Gravitational-wave mediated preheating

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ABSTRACT

We propose a new preheating mechanism through the coupling of the gravitational field to both the inflaton and matter fields, without direct inflaton–matter couplings. The inflaton transfers power to the matter fields through interactions with gravitational waves, which are exponentially enhanced due to an inflation–graviton coupling. One such coupling is the product of the inflaton to the Pontryagin density, as in dynamical Chern–Simons gravity. The energy scales involved are constrained by requiring that preheating happens fast during matter domination.

1. Introduction

Inflation is the paradigm wherein the universe undergoes exponential expansion, resolving the horizon, entropy and structure formation problems that plague the standard big bang scenario. It is usually believed that inflation ends once the inflaton reaches the bottom of its potential, at which point a new mechanism must act to transfer the inflaton's kinetic energy into a process that leads to particle creation. One such mechanism is *preheating* [1–4]: the inflaton enters a phase of parametric resonance, as it oscillates around the minimum of its potential, and through a direct matter–inflaton coupling, it leads to particle creation. There is a large number of possible direct couplings between the inflaton and the standard model, and one must usually pick one somewhat arbitrarily.

But what if the coupling between the inflaton and the matter fields were *indirect*? Because of the equivalence principle, the graviton will interact with all matter fields and its coupling will be non-arbitrary. Let us then consider the inflaton coupling to matter fields through a graviton intermediary. That is, consider the inflaton at the end of inflation depositing its kinetic energy in the graviton, which due to a direct graviton–inflaton coupling becomes parametrically excited, and then deposits its energy in the matter fields. This can happen if there are new couplings between the inflation field and the graviton so that when the inflaton oscillates it causes resonances in the graviton's modulation. We will show that such couplings are possible via Chern–Simons modified general relativity. As we shall see, no direct inflaton–matter coupling will be necessary to obtain preheating in this gravitational-wave mediated scenario.

2. Action and evolution equations

Many inflationary paradigms exist, but for concreteness consider the following Chern–Simon extension to Chaotic Inflation [5,6]

$$S = \int d^4x\sqrt{-g}\left[\mathcal{L}_{\rm EH} + \mathcal{L}_{\rm int} + \mathcal{L}_\phi + \mathcal{L}_\chi\right], \tag{1}$$

$$\mathcal{L}_{\rm EH} = \frac{R}{16\pi G}, \qquad \mathcal{L}_{\rm int} = \frac{\alpha}{4}\phi R\widetilde{R}, \tag{2}$$

$$\mathcal{L}_\phi = -\frac{1}{2}\left(g^{\mu\nu}\partial_\mu\phi\partial_\nu\phi + m_\phi^2\phi^2\right), \tag{3}$$

where ϕ is the inflaton, χ is the matter field, R is the Ricci scalar associated with the metric $g_{\mu\nu}$, $R\widetilde{R}$ is the Pontryagin density, i.e. the contraction of the Riemann tensor with its dual, and α is a coupling constant with dimensions of inverse mass (we work here in natural units $c = 1 = h$). Except for the interaction term $\mathcal{L}_{\rm int}$, Eq. (1) is just a simple model for inflation with a quadratic inflaton potential, arising from a Taylor expansion about its minimum.

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Many possible graviton–inflaton couplings could be considered, but the one presented above, $\mathcal{L}_{\rm int}$, is well-motivated. Such a coupling arises naturally in a variety of frameworks: (i) in heterotic string theories upon 4-dimensional compactification and a low-energy expansion [7,8]; (ii) in loop quantum gravity when the Barbero Immirzi parameter is promoted to a field and coupled to fermions [9,10]; (iii) in effective field theories of inflation [11]; (iv) in dynamical Chern–Simons gravity [12]. Let us emphasize, however, that the gravitational wave-mediated preheating mechanism proposed here does not depend on this particular coupling.

Regardless of the motivation, the theory described above should be considered *effective*, a truncated low-energy expansion of a more fundamental theory that is thus valid only up to some energy cut-off Λ. The effective theory ceases to be a valid description when the interaction term $\mathcal{L}_{\rm int}$ becomes comparable to the Einstein–Hilbert term $\mathcal{L}_{\rm EH}$. The former can be written as a total derivative if ϕ is constant. Therefore, to estimate its size we should first integrate by parts, moving a derivative from $R\widetilde{R}$ to ϕ. The interaction term then becomes comparable to $\mathcal{L}_{\rm EH}$ when $\alpha\dot{\phi} \sim M_p^2 (h_0 f)^{-1}$, where $M_p = G^{-1/2}$ is the Planck mass, f is the gravitational wave frequency and h_0 is the gravitational wave amplitude. Saturating α at Λ^{-1}, $\dot{\phi}$ at HM_p, h_0 at unity and f at H, where H is the Hubble parameter, one finds $\Lambda = (H/M_p)^2 M_p$, which of course satisfies $\Lambda \ll M_p$. Another consequence of the truncation of the effective theory at this order is that the terms neglected in the expansion, such as $(\partial\phi)^4/\Lambda^4$, are indeed small and ignorable. A consequence of all of this is that the interaction term $\mathcal{L}_{\rm int}$ acts as a small perturbation to whichever inflationary mechanism one wishes to consider, and thus, it does not spoil (or really affect) inflation, until inflation ends and the inflaton reaches the bottom of its potential.

Variation of the action in Eq. (1) with respect to all degrees of freedom leads to the field equations [12]

$$G_{\mu\nu} + 16\pi G\, \alpha\, C_{\mu\nu} = 8\pi G\, T_{\mu\nu}, \tag{4}$$

$$\Box\phi - m_\phi^2 \phi + \frac{\alpha}{4} R\widetilde{R} = 0, \tag{5}$$

$$\Box\chi - m_\chi^2 \chi = 0, \tag{6}$$

where $\Box$ is the curved wave-operator, $G_{\mu\nu}$ is the Einstein tensor, $T_{\mu\nu}$ is the sum of the stress-energy tensors of the ϕ and χ fields, and

$$C^{\mu\nu} \equiv \nabla_\alpha \phi\, \epsilon^{\alpha\beta\gamma(\mu} \nabla_\gamma R^{\nu)}_\beta + \nabla_{(\alpha}\nabla_{\beta)}\phi \widetilde{R}^{\beta(\mu\nu)\alpha}, \tag{7}$$

with $\widetilde{R}_{\beta(\mu\nu)\alpha}$ the dual Riemann tensor with indices symmetrized.

3. Order reduction and perturbation theory

Let us expand the equation for the metric tensor and the inflaton about a fixed background: $g_{\mu\nu} = \bar{g}_{\mu\nu} + \lambda h_{\mu\nu}$ and $\phi(t,\mathbf{x}) = \bar{\phi}(t) + \lambda\delta\phi(t,\mathbf{x})$, where λ is an order-counting parameter. The background $\bar{g}_{\mu\nu}$ and $\bar{\phi}(t)$ will be taken to be the flat Friedmann–Lemaître–Robertson–Walker (FLRW) metric and a homogeneous and isotropic background field respectively, while $h_{\mu\nu}$ and $\delta\phi(t,\mathbf{x})$ are first-order perturbations.

The FLRW metric satisfies the background field equations exactly for any homogeneous and isotropic background inflaton field. The Hubble parameter is sourced by the energy density and pressure of this background field and the matter fields (for reasons that will become clear later, we do not decompose χ). The background inflaton field satisfies the homogeneous and isotropic wave equation on an FLRW background with a mass potential. The Pontryagin density does not contribute to the background evolution of the inflaton, because this quantity vanishes exactly when evaluated for any spherically symmetric metric.

To first-order in λ, the equations for the metric tensor perturbation become [13]

$$\begin{aligned}\bar{\Box} h'_{ij} = {}& \frac{16\pi G}{a} \epsilon^{lm}_{(i}[(\ddot{\bar{\phi}} - H\dot{\bar{\phi}})\dot{h}'_{j)l} + \dot{\bar{\phi}}\,\bar{\Box} h'_{j)l}]_{,m} \\ & + 16\pi G a^2 p_\chi h'_{ij},\end{aligned} \tag{8}$$

while, neglecting scalar metric perturbations (we are looking at modes shorter than the Hubble scale), the equation for the inflaton perturbation is

$$\delta\ddot{\phi} + 3H\delta\dot{\phi} - \frac{1}{a^2}\bar{\nabla}^2\delta\phi = -m_\phi^2\delta\phi, \tag{9}$$

where $\bar{\nabla}^2$ and $\bar{\Box}$ are the Laplacian and wave operators in a homogeneous and isotropic FLRW background, p_χ is the pressure of the χ field and $h'_{ij} \equiv a^{-2} h_{ij}$. Notice again that the Pontryagin density does not enter the evolution equation of the inflaton perturbation, since it also vanishes identically to linear order in λ.

We can simplify the evolution equation for the metric perturbation through order reduction. As discussed in [14–16], we decompose the metric perturbation into a general relativistic piece $h^{\rm GR}_{\mu\nu}$ and a deformation $\delta h_{\mu\nu}$, namely $h_{\mu\nu} = h^{\rm GR}_{\mu\nu} + \alpha^2 \delta h_{\mu\nu}$. Note that the deformation is proportional to α^2 because Eq. (4) is proportional to $\alpha C_{\mu\nu}$, $C_{\mu\nu}$ is proportional to ϕ, and ϕ is proportional to α due to Eq. (5). Using this decomposition, we can order reduce Eq. (8): the left-hand side is proportional to $\bar{\Box}\delta h_{\mu\nu}$, while the right-hand side is proportional to a differential operator acting on $h^{\rm GR}_{\mu\nu}$. This differential operator will contain one term of the form $\bar{\Box}$, which automatically vanishes when acting on $h^{\rm GR}_{\mu\nu}$ because $R_{\mu\nu}[h^{\rm GR}_{\alpha\beta}] = 0$. Using this and going to the transverse-traceless (TT) gauge [17] and in the left/right-circular polarization basis for the gravitational wave perturbation, Eq. (8) becomes

$$\bar{\Box} h_R = i\frac{16\pi G}{a}\alpha(\ddot{\bar{\phi}} - H\dot{\bar{\phi}})\frac{\partial}{\partial z}\dot{h}_R + 16\pi G a^2 p_\chi h_R. \tag{10}$$

The equation for h_L can be obtained by taking $i \to -i$ and $h_{L/R} \to h_{R/L}$. The amplitudes h_L and h_R are defined by

$$h'_{ij} = \frac{1}{\sqrt{2}}\begin{pmatrix} -(h_L + h_R) & i(h_L - h_R) & 0 \\ i(h_L - h_R) & (h_L + h_R) & 0 \\ 0 & 0 & 0 \end{pmatrix}. \tag{11}$$

Notice that Eqs. (8) and (9) are not coupled and can thus be solved independently, once the evolution of the background fields is obtained.

Let us now discuss the evolution of the matter fields. We anticipate that the matter occupation number will be generated through parametric resonance, so even a small perturbation of size $\mathcal{O}(|h_{\mu\nu}|)$, may have a large effect. We therefore treat the matter field, χ, exactly, without a perturbative decomposition, while the gravitational field is treated to first order in its perturbations, $h_{\mu\nu}$. We obtain the equation to first order in $h_{\mu\nu}$,

$$\ddot{\chi} + 3H\dot{\chi} - \frac{1}{a^2}\bar{\nabla}^2\chi + h^{ij}\partial_i\partial_j\chi = -m_\chi^2\chi. \tag{12}$$

In the TT gauge and in a circular GW polarization basis, this equation becomes

$$\ddot{\chi} + 3H\dot{\chi} - \frac{1}{a^2}\bar{\nabla}^2\chi - \frac{1}{\sqrt{2}a^2}{\rm Re}[h_L\partial_L^2 + h_R\partial_R^2]\chi = -m_\chi^2\chi, \tag{13}$$

where ${\rm Re}[x]$ is the real part of x and we have defined $\partial_{L,R} \equiv (\partial/\partial x \mp i\partial/\partial y)/\sqrt{2}$.