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Can oscillating scalar fields decay into particles with a large thermal mass?

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Abstract

We calculate the dissipation rate of a coherently oscillating scalar field in a thermal environment using nonequilibrium quantum field theory and apply it to the reheating stage after cosmic inflation. It is shown that the rate is nonvanishing even when particles coupled to the oscillating inflaton field have a larger thermal mass than it, and therefore the cosmic temperature can be much higher than inflaton's mass even in the absence of preheating. Its cosmological implications are also discussed.

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1. Introduction

In our contemporary understanding, the origin of the primeval fireball whose existence was assumed in the conventional hot big bang cosmology is the reheating processes after inflation—an accelerated cosmic expansion which has made the universe homogeneous and spatially flat with small density fluctuations that eventually grow to the observed large-scale structure [1,2]. The universe is reheated through the dissipation of coherent oscillation of the zero-mode of the *inflaton*, the scalar field whose potential energy drives inflation. While the initial stage of reheating could be rather complicated due to an explosive particle production induced by parametric resonance, which is dubbed as preheating [3], the final stage is dominated by perturbative decay. The latter process determines the reheat temperature, T_R, the temperature at the outset of the radiation domination [4].

Note, however, that in general T_R is much lower than the highest temperature the universe has ever experienced after inflation even in the case only perturbative decay operates to reheat the universe [5]. This means that in the late stage of the reheating processes, the inflaton decays not in a vacuum but in a thermal medium. About this point an interesting claim has been made in [6] that if the would-be decay products of the oscillating inflaton acquire a thermal mass larger than the inflaton mass in the thermal background, it cannot decay into these particles, and that reheating is suspended for some time, based on the observation that the phase space would be closed for the mass of the decay product being larger than half the inflaton mass. The decay width of the inflaton ϕ with mass m_ϕ into two massive particles with mass m reads

$$\Gamma_\phi = \Gamma_{\phi 0}\left(1 - \frac{4m^2}{m_\phi^2}\right)^{1/2}, \tag{1}$$

where $\Gamma_{\phi 0}$ is the decay rate in the case $m = 0$. So if we simply replace m with a thermal mass $m(T) \sim gT$ and if it is larger than $m_\phi/2$, the phase space is closed and inflaton decay is apparently forbidden. Here g is some coupling constant of the would-be decay product. Then thermal history after inflation would be drastically changed. That is, the highest temperature in this era cannot exceed $\sim m_\phi/g$ if preheating is inoperative, and also the reheat temperature is bounded from above by m_ϕ/g and is independent of the decay rate of the inflaton in case the conventional calculation gives a larger value. The former would change the abundance of supermassive particles and the latter affects the gravitino abundance [6], because the gravitino-entropy ratio after inflation is proportional to the temperature at the onset of radiation domination.

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Furthermore, this situation is not specific to the reheating stage after inflation but may apply in any epoch when significant amount of entropy is produced out of the decay of oscillating scalar field with a relatively small mass. Indeed the above possibility was first pointed out by Linde [7] in the context of Affleck–Dine baryogenesis [8] where the Affleck–Dine scalar field oscillates with a mass of order of 10^{2-3} GeV in a medium with a much higher temperature. The final magnitude of baryon asymmetry changes if this suspension of decay is operative [7].

The above naive picture, however, may be too simplistic because a thermal mass is different from the intrinsic mass and because coherent field oscillation is different from a collection of particles. Thus it is very important to analyze this problem from a more fundamental point of view, for it has a profound implication not only to the cosmology of the early Universe but also to particle physics in that it strongly affects various species of the particles produced in the early universe as mentioned above. In this Letter, extending our previous work [9], we analyze this problem in terms of a nonequilibrium quantum field theory at finite temperature [10]. Inclusion of thermal masses of the would-be decay products of the inflaton is achieved by adopting a resummed propagator when we calculate the effective action for ϕ. As a result of resummation the self energy of the decay product acquires not only real part, which appears as a high-temperature correction to the mass, but also an imaginary part. The latter plays a crucial role in determining the dissipation rate of the inflaton. Consequently, we find that the inflaton can dissipate its energy even when its would-be decay products have a larger thermal mass than the inflaton itself.

2. Model and equation of motion

For clarity we adopt a simpler model than [9], that is, we adopt the following Lagrangian.

$$\mathcal{L} = \frac{1}{2}(\partial_\mu \phi)^2 - \frac{1}{2}m_\phi^2\phi^2 + \frac{1}{2}(\partial_\mu \chi)^2 - \frac{1}{2}m_\chi^2\chi^2 - \mathcal{M}\phi\chi^2 - \frac{1}{4}g^2\chi^4. \tag{2}$$

Here ϕ is an oscillating real scalar field and χ is another real scalar field. ϕ can decay into a pair of χ particles through the interaction $\mathcal{M}\phi\chi^2$, if it is energetically allowed. The dimensionful coupling constant $\mathcal{M}$ may be written as $\mathcal{M} = hm_\phi$ in some supersymmetric inflation models where h is a Yukawa coupling [11]. In such models ϕ can also decay into two fermions, which is suppressed by Pauli-blocking at finite temperature. On the contrary, the decay into two bosons is enhanced due to the induced emission. This is the reason we consider the latter decay process.

We analyze the behavior of the above system under the following assumption which mimic the cosmological situations we are interested in, such as the late reheating phase after inflation. First we neglect cosmic expansion since we are interested in the phenomena which occur in a shorter time scale than the expansion time. Second we assume χ is in a thermal state with a specific temperature $T = \beta^{-1}$ and that it acquires a large thermal mass due to the self coupling. Note that χ can easily be thermalized during the reheating stage since its thermalization rate, $\sim g^4 T$, can naturally be much larger than the cosmic expansion rate. Finally the scalar field ϕ is oscillating but we consider the situation the parametric resonance is already terminated with a field amplitude $\mathcal{M}|\phi| < m_\phi^2$.

Due to its coherent nature, scalar field oscillation behaves nearly classically, but its decay is of course a quantum process. So we calculate an effective action for ϕ and derive an equation of motion for its expectation value. For this purpose we should use the in–in or the closed time-path formalisms in which the time contour starting from the infinite past must run to the infinite future without fixing the final condition and come back to the infinite past again in calculating the generating functional [12]. This method has been applied to various cosmological problems by a number of authors [13–15]. The generating functional in the present model is given by

$$Z[J,K] = \mathrm{Tr}\left[T_-\left\{\exp\left[i\int_\infty^{-\infty} dt \int d^3x\,(J_-\phi_- + K_-\chi_-)\right]\right\} \times T_+\left\{\exp\left[i\int_{-\infty}^{\infty} dt \int d^3x\,(J_+\phi_+ + K_+\chi_+)\right]\right\}\rho\right] \equiv e^{iW[J,K]}. \tag{3}$$

where X_+ denotes a field component X on the plus-branch ($-\infty$ to $+\infty$) and X_- that on the minus-branch ($+\infty$ to $-\infty$). The symbol T_+ represents the ordinary time ordering, and T_- the anti-time ordering. $J_\pm$ and $K_\pm$ are the external fields for ϕ and χ, respectively. ρ is the initial density matrix which is assigned according to the assumption above mentioned.

In terms of the components along the plus and the minus branches, the effective action reads

$$\Gamma[\phi_+,\phi_-] = W[J_+, J_-, K_\pm = 0] - \int_{-\infty}^{\infty} dt \int d^3x\left[J_+(x)\phi_+(x) - J_-(x)\phi_-(x)\right], \tag{4}$$

with $\phi_+(x) = \frac{\delta W[J_+,J_-]}{\delta J_+(x)}$ and $\phi_-(x) = \frac{-\delta W[J_+,J_-]}{\delta J_-(x)}$.

Here we consider a one-loop correction depicted in Fig. 1 which includes the essential effect in our analysis for illustration. We also note that instead of ϕ_+ and ϕ_- it is more convenient to use $\phi_c \equiv (\phi_+ + \phi_-)/2$ and $\phi_\Delta \equiv \phi_+ - \phi_-$ and set $\phi_\Delta \to 0$ in the end because ϕ_+ and ϕ_- should be identified with each other eventually. Then the effective action to this order is

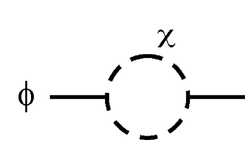

Fig. 1. One-loop Feynman diagram incorporated in the effective action. Solid line denotes ϕ, and broken line χ.