Probing the reheating temperature of the universe with a gravitational wave background

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Abstract. The thermal history of the universe after big bang nucleosynthesis (BBN) is well understood both theoretically and observationally, and recent cosmological observations also begin to reveal the inflationary dynamics. However, the epoch between inflation and BBN is scarcely known. In this paper we show that the detection of the stochastic gravitational wave background around 1 Hz provides useful information about thermal history well before BBN. In particular, the reheating temperature of the universe may be determined by future space-based laser interferometer experiments such as DECIGO and/or BBO if it is around $10^6$–$10^9$ GeV, depending on the tensor-to-scalar ratio $r$ and dilution factor $F$.

Keywords: gravity waves/theory, inflation, physics of the early universe

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1. Introduction

Recent cosmological observations have determined cosmological parameters with unprecedented accuracy [1]. From measurements of the cosmic microwave background (CMB) anisotropy, the baryon content of the universe $\Omega_b$ is determined and it agrees with the prediction of big bang nucleosynthesis (BBN). This confirms the standard thermal history of the universe at least for $T \lesssim O(1) \text{ MeV}$, where BBN begins. On the other hand, the primordial power spectrum of the density perturbation has been revealed to be nearly scale invariant. This has not only provided further evidence of inflation in the early universe [2] but also given observational clues for studying its details such as the shape of the scalar potential driving inflation [3].

The thermal history of the universe between the inflationary era and BBN epoch, however, is the least known observationally. In general, the universe is dominated by the coherent oscillation of the inflaton soon after the inflation ends, and finally the inflaton decays into standard model particles, which realizes the hot radiation dominated universe. Thus understanding this ‘reheating’ stage of the universe is very important. In a simple standard inflation scenario, the reheating effect is characterized by one parameter, the reheating temperature $T_R$ [4], which is defined as

$$T_R = \left( \frac{10}{\pi^2 g_*(T_R)} \right)^{1/4} \sqrt{\Gamma_\phi M_G}, \quad (1)$$

where $\Gamma_\phi$ is the decay rate of the inflaton, $M_G = (8\pi G)^{-1/2} = 2.4 \times 10^{18} \text{ GeV}$ is the reduced Planck scale and $g_*$ denotes the relativistic effective degrees of freedom. However, $T_R$ is hardly constrained from cosmological observations now. The only requirement is that
$T_R$ must be larger than a few MeV [5] in order that light elements must be created in a usual way.

However, the reheating temperature $T_R$ contains rich information from the viewpoint of particle physics. First, as is obvious from equation (1), $T_R$ is determined by the inflaton decay rate, and it depends on the inflaton properties, such as its mass, potential and interaction strength with standard model particles. Thus determining $T_R$ may have impacts on choosing realistic inflation models and inflaton candidates. Moreover, in supersymmetry (SUSY) [6], which is well-motivated physics beyond the standard model, theoretical upper bounds on the reheating temperature are imposed from the so-called gravitino problem [7]–[10]. The abundance of the gravitino, which is the superpartner of the graviton, is bounded from above in order not to destroy light elements through its decay processes for an unstable gravitino, or not to overclose the universe for a stable gravitino. This gives upper bounds on $T_R$, since gravitinos are produced efficiently in the reheating era and its abundance is proportional to $T_R$.

Then, how can we probe such an early stage of the universe? As is well known, observed CMB photons come from the last scattering surface located at redshift $z \sim 1100$. The universe is opaque to photons beyond the last scattering surface, and hence we cannot directly look over the universe of $z \gtrsim 1100$ via observations of photons. Even the cosmic background neutrinos cannot be used as a probe of the early universe with $T \gtrsim 1$ MeV, since neutrinos couple with nucleons strongly in such a high temperature environment.

The universe, however, is transparent to gravitational waves up to the Planck epoch in principle. This opens up the window for probing the very early universe, in particular the reheating epoch, with observations of gravitational waves. Inflation dilutes, along with many species of unwanted relics, pre-existing gravitational waves at the Planckian epoch. Instead, quantum mechanical gravitational waves are produced in the inflationary era with an almost scale-invariant power spectrum [11]. Although the amplitude of gravitational waves is constant in the super-horizon regime, once a mode enters the horizon, it is reduced as the universe expands. Since the expansion rate depends on the equation of state of the universe, the corresponding thermal history of the universe is imprinted in the gravitational wave spectrum at present. Although much work has been done which treats the spectrum and detection possibility of the stochastic gravitational wave background of inflationary origin in the literature [12]–[30], the previous work did not take into account the modification of the gravitational wave spectrum around the reheating epoch and assumed a nearly flat spectrum above the observationally interesting frequency $\sim1$ Hz, except for in a few works where it was claimed that the equation of state of the early universe can be probed by looking at the spectrum of the gravitational wave background [21] (this kind of study was also done and extended in [24]). Recently, we have pointed out that taking into account the reheating effect on the gravitational wave spectrum is essential both for the detection itself and for the purpose of probing the thermal history of the universe between BBN and inflation, with an emphasis on its impacts on particle physics [31].

In this paper, we give more detailed analyses of reheating effects on the gravitational wave background and show that future mission concepts based on a space laser interferometer, like Japan’s DECIGO [32] and NASA’s Big Bang Observer (BBO), may have a chance to detect a primordial gravitational wave background produced in the inflationary era and determine the thermal history of the universe before the BBN epoch.
Probing the reheating temperature of the universe with a gravitational wave background

in particular the reheating temperature $T_R$. We also discuss that some ‘non-standard’ cosmological scenarios, including late-time entropy production, can be probed.

This paper is organized as follows. In section 2 the spectrum of primordial gravitational wave backgrounds, taking into account the modification around the reheating epoch, is derived. Using this result, we discuss future prospects for the determination of the reheating temperature with space-based laser interferometers in section 3. We discuss the case of late-time entropy production in section 4. In section 5, the impacts of determining the reheating temperature on particle physics models are summarized. Section 6 is devoted to our conclusions.

2. Primordial gravitational wave spectrum

Primordial gravitational waves produced in the inflationary era have a nearly scale-invariant spectrum and have effects on both large- and small-scale cosmological observations [12]. On large scales at observable frequencies $\sim 10^{-16}$ Hz, the tensor metric perturbation generates $B$-mode polarization anisotropy of the cosmic microwave background (CMB) for small multipole $l$, and such a signal may be probed by next-generation ground-based or satellite observations of the CMB (polarization) anisotropy [33]–[35]. On small scales, gravitational waves may be detected by future space-based interferometer experiments with frequency around $\sim 1$ Hz. Thus primordial gravitational waves give much information on a wide range of cosmological scales. In particular, direct detection of gravitational waves may reveal the state of the universe well before BBN, as we will see.

A gravitational wave is described by the tensor perturbation on the metric, $h_{ij}$, which is defined as

$$ds^2 = a^2(t)(-dt^2 + (\delta_{ij} + 2h_{ij})dx^i dx^j),$$

where $h_{ij}$ is symmetric under the exchange of $i$ and $j$, and satisfies the traceless and transverse condition, $h_{ii} = 0$, $h_{ij,j} = 0$. Thus it has two physical degrees of freedom, which we denote as $+$ and $\times$. Conformal time $\tau$ is defined as $d\tau = dt/a(t)$. Since tensor perturbation is gauge invariant as it is, we do not need to concern ourselves with gauge ambiguity as long as only the tensor mode is considered. The tensor perturbation $h_{ij}$ is expanded in the Fourier space as

$$h_{ij} = \sqrt{8\pi G} \sum_{\lambda=+,\times} \int \frac{d^3k}{(2\pi)^3/2} h_{ij}^{\lambda} e_{ij}^{(\lambda)} e^{i\mathbf{kx}},$$

where the superscript $\lambda$ denotes each polarization degree, $\lambda = +/\times$, and $e_{ij}^{(\lambda)}$ is the polarization tensor, which satisfies $e_{ij}^{(\lambda)} e^{ij(\lambda')} = \delta^{\lambda\lambda'}$.

Primordial gravitational waves are produced in the inflationary epoch with an almost scale-invariant spectrum. The amplitude of the gravitational wave produced is proportional to the Hubble scale during inflation, $H_{\text{inf}}$. In terms of the dimensionless power spectrum, it is given by [36]

$$\Delta_h^{(p)}(k)^2 = 64\pi G \left(\frac{H_{\text{inf}}}{2\pi}\right)^2 \left[ 1 - 2\epsilon \ln \frac{k}{k_*} + 2\epsilon(\eta - \epsilon) \left(\ln \frac{k}{k_*}\right)^2 \right],$$

Journal of Cosmology and Astroparticle Physics 06 (2008) 020 (stacks.iop.org/JCAP/2008/i=06/a=020) 4
counting two polarization states of the gravitational wave, where \( \epsilon \) and \( \eta \) denote the slow-roll parameters during inflation defined as \( \epsilon = M_{\text{Pl}}^2(V'/V)^2/2 \) and \( \eta = M_{\text{Pl}}^2(V''/V) \) with an inflaton potential \( V[\phi] \), and its first and second derivative with respective to the inflaton field \( \phi \), \( V' \) and \( V'' \). We take the pivot scale as \( k_\ast = 0.002 \text{ Mpc}^{-1} \). Here slow-roll parameters should be evaluated at the epoch when the CMB scale leaves the horizon.

On the other hand, the spectrum of the curvature perturbation is given by

\[
\Delta^2_R(k) = \frac{4\pi G}{\epsilon} \left( \frac{H_{\text{inf}}}{2\pi} \right)^2 \left[ 1 + (-6\epsilon + 2\eta) \ln \frac{k}{k_\ast} + (6\epsilon^2 - 4\epsilon\eta + 2\eta^2 - \xi) \left( \ln \frac{k}{k_\ast} \right)^2 \right],
\]

where \( \xi \equiv M_{\text{Pl}}^4V''V''/V^2 \). Its normalization is already well determined observationally, as \( \Delta^2_R \sim 2.0 \times 10^{-9} \) on the CMB scale [1]. Thus the slow-roll parameter \( \epsilon \) is related to the tensor-to-scalar ratio \( r \equiv \Delta^2_h/\Delta^2_R(k_\ast) \) through the relation

\[
r = 16\epsilon.
\]

This shows that \( r \) is in general a small parameter and the tensor contribution is also expected to be small compared to the density perturbation. Notice that measuring \( r \) fixes the inflation scale as

\[
V_{\text{inf}} = 3M_{\text{Pl}}^2H_{\text{inf}}^2 \simeq (3.2 \times 10^{16} \text{ GeV})^4 r.
\]

Then let us derive the observable power spectrum of the inflationary gravitational wave background. The evolution of the gravitational wave is described by the following equation:

\[
\ddot{h}_k^\lambda + 3H \dot{h}_k^\lambda + \frac{k^2}{a^2} h_k^\lambda = 0.
\]

The amplitude of the gravitational wave with comoving wavenumber \( k \) remains constant when the mode lies outside the horizon. However, once it enters the horizon, its amplitude begins to damp. For the mode which enters the horizon in the matter dominated regime, the solution is written as

\[
h_k^\lambda(\tau) = h_k^{\lambda(p)} \left( \frac{3j_1(k\tau)}{k\tau} \right),
\]

where \( j_k \) denotes the \( \ell \)th spherical Bessel function. In general, the solution in a power law background \( a(t) \propto t^p \) can be written in the form

\[
h_k(\tau) \propto a(t)^{(1-3p)/2p} J_{(3p-1)/(2(1-p))}(k\tau),
\]

with Bessel function \( J_n(x) \). Another damping factor comes from the fact that the relativistic degrees of freedom \( g_* \) do not remain constant and the expansion rate is modified from a simple power law \( a(t) \propto T^{-1} \) in the early universe. This effect gives the damping factor [28, 29]

\[
\left( \frac{g_s(T_{\text{in}})}{g_{*0}} \right)^{4/3} \left( \frac{g_{*s0}}{g_{*s}(T_{\text{in}})} \right)^{4/3}
\]

\(^5\) Here we neglect the anisotropic stress. Including an anisotropic stress modifies the gravitational wave spectrum for the present frequency \( \lesssim 10^{-9} \text{ Hz} \) due to the neutrino free streaming effect [37, 29], but our concern is the frequency around \( \sim 1 \text{ Hz} \). Hence the following arguments are not modified.
on the power spectrum of the gravitational wave, where $T_{\text{in}}$ denotes the temperature at which the corresponding mode enters the horizon, given by

$$T_{\text{in}} \simeq 5.8 \times 10^6 \text{ GeV} \left( \frac{g_{*s}(T_{\text{in}})}{106.75} \right)^{-1/6} \left( \frac{k}{10^{14} \text{ Mpc}^{-1}} \right)^3.$$  

Moreover, recent cosmological observations have revealed that the expansion of the present universe is accelerating due to the dominance of an unknown energy density, called dark energy (here we assume the cosmological constant for simplicity). This also gives a suppression factor $\sim (\Omega_{m}/\Omega_{\Lambda})^2$ on the power spectrum [16]. As a result, the present gravitational wave spectrum per log frequency interval is written in the form

$$\Omega_{\text{gw}}(f) = \frac{k^2}{12H_0^2} \Delta_h^2(k),$$  

where

$$\Delta_h^2(k) = \Delta_h^{(p)}(k)^2 \left( \frac{\Omega_m}{\Omega_{\Lambda}} \right)^2 \left( \frac{g_{*s}(T_{\text{in}})}{g_{*s0}} \right) \left( \frac{g_{*s0}}{g_{*s}(T_{\text{in}})} \right)^{4/3} \left( \frac{3j_1(k\tau_0)}{k\tau_0} \right)^2 T_1^2(x_{\text{eq}}) T_2^2(x_R),$$  

with a bar denoting the average over many periods. Here, $T_1(x_{\text{eq}})$ denotes the transfer function, which connects the gravitational wave spectrum of the mode which enters the horizon before and after matter–radiation equality, $t = t_{\text{eq}}$. It is calculated as [16]

$$T_1^2(x_{\text{eq}}) = \left[ 1 + 1.57x_{\text{eq}} + 3.42x_{\text{eq}}^2 \right],$$  

where $x_{\text{eq}} = k/k_{\text{eq}}$ and $k_{\text{eq}} \equiv a(t_{\text{eq}})H(t_{\text{eq}}) = 7.1 \times 10^{-2}\Omega_m h^2 \text{ Mpc}^{-1}.$

On the other hand, $T_2(x_R)$ connects the mode which enters the horizon after and before the reheating ends. Because before the inflaton decays the universe is dominated by the coherent oscillation of the inflaton, the spectrum of the gravitational wave changes for the mode which enters the horizon at the inflaton dominated epoch ($k > k_R$), where

$$k_R \simeq 1.7 \times 10^{13} \text{ Mpc}^{-1} \left( \frac{g_{*s}(T_R)}{106.75} \right)^{1/6} \left( \frac{T_R}{10^6 \text{ GeV}} \right) .$$

In terms of the frequency, this corresponds to

$$f_R \simeq 0.026 \text{ Hz} \left( \frac{g_{*s}(T_R)}{106.75} \right)^{1/6} \left( \frac{T_R}{10^6 \text{ GeV}} \right) ,$$

which is close to the most sensitive frequency band of DECIGO and BBO for $T_R \sim 10^6 \text{ GeV}$. The transfer function $T_2(x_R)$ is obtained by solving simultaneously equation (8) and the following Friedmann equations taking the decay of inflaton into account:

$$\dot{\rho}_\phi + 3H\rho_\phi = -\Gamma_\phi \rho_\phi, \quad \dot{\rho}_r + 4H\rho_r = \Gamma_\phi \rho_\phi, \quad H^2 = \frac{8\pi G}{3} (\rho_\phi + \rho_r),$$

6 Coefficients in the right-hand side of (15) differ from those of [16] reflecting the difference in the definition of $k_{\text{eq}}$.  

Journal of Cosmology and Astroparticle Physics 06 (2008) 020 (stacks.iop.org/JCAP/2008/i=06/a=020)
Probing the reheating temperature of the universe with a gravitational wave background

Figure 1. A schematic picture of the evolution of the Hubble horizon scale $H^{-1}$ and the physical wavelength of some modes. φ.D., R.D., M.D. and Λ.D. denote the inflaton oscillation dominated era, radiation dominated era, matter dominated era, and cosmological constant dominated era, respectively.

where $\rho_\phi$ denotes the energy density of the inflaton coherent oscillation. We find that $T^2(z_R)$ is well approximated by

$$T^2(z_R) = \left[ 1 - 0.32x_R + 0.99x_R^2 \right]^{-1}, \quad (21)$$

where $x_R = k/k_R$. Thus $\Omega_{gw}(f)$ behaves as $\propto f^{-2}(f^0)$ for the mode which enters the horizon in the matter (radiation) dominated regime. In figure 1 we show a schematic picture for the evolution of the Hubble horizon scale $H^{-1}$ by a solid line, and the physical wavelength of various modes, $k > k_R$, $k_{eq} < k < k_R$ and $k < k_{eq}$, by dotted lines. Since the evolution of inflationary gravitational waves is sensitive to the expansion of the universe, $H^{-1}$, figure 1 shows that characteristic feature of the power spectrum can be probed for the mode around $f \sim f_R$.

Before discussing the detection possibility, we briefly mention the effect of the running spectral index. The running spectral index depends on the slow-roll parameters $\epsilon$ and $\eta$. Although $\epsilon$ is related to $r$ through the relation (6), we are free to choose $\eta$ as long as the scalar spectral index $n_s (\simeq 1 - 6\epsilon + 2\eta)$ lies in the favored range from recent WMAP results. In figure 2 the resulting spectrum $\Omega_{gw}(f)$ at $f = 0.1$ Hz is shown for $\eta = 0.01, 0, -0.01$ from upper to lower. Here a sufficiently large reheating temperature $T_R \gtrsim 10^9$ GeV is assumed. As is clearly seen, for $r \gtrsim 0.1$ the gravitational wave amplitude at $f = 0.1$ Hz decreases, since the spectral tilt and its running become large. The difference of $\eta$ yields a slight change in the gravitational wave amplitude for $r \lesssim 0.1$ (cf the current constraint on $r$, $r < 0.20$ (95% C.L.)). Hereafter we set $\eta = 0$ for simplicity.

In figure 3, we show the resulting gravitational wave spectrum for $T_R = 10^9$ and $10^5$ GeV with $r = 0.1$ and 0.001. Together with the predicted gravitational wave spectrum, sensitivities of the future space-based laser interferometer experiments, DECIGO, correlated analysis of DECIGO, ultimate-DECIGO (single) and correlated analysis of ultimate-DECIGO [38] are also presented. It can be seen that theoretical predictions are well above the sensitivity line of DECIGO (correlated) and ultimate-DECIGO for large enough $T_R$ and $r$, while the direct detection is undesirable for $T_R \lesssim 10^4$ GeV.
Figure 2. $\Omega_{gw}(f)$ at $f = 0.1$ Hz for $\eta = 0.01, 0, -0.01$ from upper to lower.

Figure 3. Primordial gravitational wave spectrum for $T_R = 10^9$ GeV and $T_R = 10^5$ GeV, shown by thin and thick lines for $r = 0.1$ and $0.001$. Also shown are the expected sensitivity of DECIGO (green dashed), correlated analysis of DECIGO (blue dot-dashed), ultimate-DECIGO (purple dashed) and correlated analysis of ultimate-DECIGO (red dotted), from upper to lower.

3. Prospects for the determination of $T_R$ with future space-based laser interferometer experiments

The spectrum of the primordial gravitational wave background generated during inflation crucially depends on the reheating temperature $T_R$ after inflation, as can be seen from figure 3. Conversely, this opens up the possibility that future experiments devoted to detecting gravitational wave background will probe the reheating stage of the universe.

The important parameters that determine the primordial gravitational wave spectrum are the tensor-to-scalar ratio $r$ and the reheating temperature $T_R$. The tensor-to-scalar ratio determines the overall normalization of the spectrum, and $T_R$ fixes the frequency above which the spectrum is significantly suppressed. The important point is that if the bending point of the gravitational wave spectrum determined by $T_R$ around the frequency given by (16) lies above the sensitivity of detectors, $T_R$ can be determined by observations of gravitational waves.
The non-zero value of the tensor-to-scalar ratio $r$ will be probed with measurements of the $B$-mode CMB polarization [34]. Since the $B$-mode is generated by the tensor perturbation only for large angular scale, $\ell \lesssim 100$, detection of the $B$-mode polarization indirectly confirms the existence of the gravitational wave background. Notice that this effect is seen in very large-scale anisotropy with wavelength comparable to the present horizon scale, so the detection of the $B$-mode polarization is somewhat complementary to the direct detection of the gravitational wave on a small scale with wavelength of the order of the Earth radius. The Planck satellite [39], scheduled to be launched in October 2008, will measure $r$ up to $\sim 0.1$. Ongoing ground-based or balloon experiments devoted to detecting CMB polarization, such as the Q/U Imaging ExperimenT (QUIET) [40] and Clover [41] will detect $r$ up to $\sim 0.01$. Ultimately, a future space mission will be dedicated to detecting the primary $B$-mode signal up to $r \gtrsim 10^{-3}$ [34]. According to recent works [42, 43], the observed value of the scalar spectral index, $n_s \sim 0.961 \pm 0.017$, implies a somewhat large tensor contribution, $r \gtrsim 10^{-3}$ for various inflation models. Thus it seems plausible that non-zero $r$ will be confirmed by CMB polarization experiments.

Once $r$ is measured from those experiments, we may have a chance to detect signals of gravitational wave background of inflationary origin. The planned laser interferometer space antenna (LISA) [44] cannot reach the required sensitivity even if the largest possible value of $r$ allowed by WMAP5 and high enough $T_R$ are assumed. However, future mission concepts like DECIGO and BBO, are likely to detect them or put a meaningful constraint on $T_R$. Therefore our main interest is in demonstrating the observable range of the reheating temperature through such future experiments for various values of $r$.

Before discussing results, however, it should be noted that stochastic noise coming from some astrophysical processes must be taken into account for the purpose of detecting the primordial gravitational wave background. In particular, gravitational waves from white-dwarf binaries are considered to completely hide the primordial ones for the frequency $f \lesssim 0.1$ Hz [45]. But for the frequency range 0.1–10 Hz, where the DECIGO and BBO are most sensitive, foregrounds from astrophysical objects are separable.

In figure 4, we show the parameter region on the $r$–$T_R$ plane which will be probed with the correlated analysis of DECIGO, ultimate-DECIGO (single) and ultimate-DECIGO (correlation). We have cut the sensitivity of these projects off below $f \lesssim 0.1$ Hz, taking into account the stochastic noise from white-dwarf binaries. The gravitational wave background can be detected in the light blue shaded region. Furthermore the dark blue shaded region shows the parameter region where the value of $T_R$ can be determined with signal-to-noise ratio 5. It is seen that for $10^{-3} \lesssim r \lesssim 1$, direct detection of the gravitational wave background can determine the reheating temperature $T_R$, if it lies in the range $T_R \sim 10^6$–$10^8$ GeV.

If higher $T_R$ is realized in Nature, the inflationary gravitational wave background can be detected by DECIGO or BBO, but the value of $T_R$ remains undetermined. In this case, the lower bound on $T_R$ can be read off from these figures. For example if $r = 0.1$, the detection of the gravitational wave background means $T_R \gtrsim 10^8$ GeV, which provides useful constraints on some particle physics models, as we will see in section 5.

7 Gravitational waves from the collapses of Population III stars may be a dominant contribution around the decihertz band [46]. However, it crucially depends on the early star formation rate, which is very uncertain now, and such a signal would be separable if we adopted reliable abundance of Population III stars and took the duty cycle and their angular distribution into account.
Figure 4. In the outer light shaded region the gravitational wave background can be detected, and the inner blue shaded region shows the region where $T_R$ can be determined with signal-to-noise ratio 5 by correlated analysis of DECIGO, ultimate-DECIGO (single) and ultimate-DECIGO (correlation) from upper to lower.

4. Late-time entropy production

So far, we have assumed that there were no late-time entropy production processes after the completion of reheating after inflation. However, this may be too simplified an assumption. Let us consider the case where some scalar field $\chi$ other than the inflaton dominates the universe after the inflaton decays and $\chi$ eventually decays releasing huge entropy. Examples are Polonyi [47] or moduli fields [48, 49], a scalar partner of the axion [50], and others. Such late-decaying particles are interesting since they dilute
Probing the reheating temperature of the universe with a gravitational wave background

Figure 5. A schematic picture of evolution of the Hubble horizon scale $H^{-1}$ and physical wavelength in the presence of late-time entropy production from $\chi$. $\chi$-D. represents the $\chi$-dominated era.

cosmologically harmful gravitinos. We denote the dilution factor as $F$, defined by

$$F = \frac{s(T_{\chi})a^3(T_{\chi})}{s(T_R)a^3(T_R)} = \frac{T_R}{T_{\chi}} \left( \frac{\rho_\chi}{\rho_{\phi}} \right)_{T_R},$$

where $T_{\chi}$ is the decay temperature of $\chi$, which must be larger than a few MeV, and $\rho_\chi$ denotes the energy density of the $\chi$-field coherent oscillation. For a simple model where the $\chi$-field has an initial amplitude $\chi_i$ and begins to oscillate when the Hubble parameter becomes equal to the mass of the $\chi$, it is estimated as

$$F = \frac{T_R}{T_{\chi}} \frac{\chi_i^2}{3M_{\odot}^2}.$$

The abundance of all dangerous cosmological relics produced in the reheating era after inflation, such as the gravitino discussed below, are diluted by this factor.

Importantly, such non-standard cosmological evolution scenarios are imprinted in the present gravitational wave spectrum. In the presence of such a late-decaying particle, the additional $\chi$-matter dominated era suppresses the gravitational wave amplitude for the frequency which re-entered the horizon during or before the time when $\chi$ began to dominate the universe. (Figure 5 shows a schematic picture.) The spectrum now becomes

$$\Omega_{gw}(f, F) = \Omega_{gw}(f) \times T_2^2(x_\chi)T_1^2(x_{\chi R}),$$

where $\Omega_{gw}(f)$ is given by equation (13) with $k_R$ replaced by $k_R(F) = k_R F^{-1/3}$.

Here $x_\chi(= k/k_\chi)$ corresponds to the wavenumber which enters the horizon at the decay of $\chi$,

$$k_\chi \simeq 1.7 \times 10^7 \text{ Mpc}^{-1} \left( \frac{g_{*s}(T_{\chi})}{106.75} \right)^{1/6} \left( \frac{T_\chi}{1 \text{ GeV}} \right),$$

Here we assume that $\chi$ begins to oscillate during the inflaton oscillation dominated phase.
and $x_{\chi R} (\equiv k/k_{\chi R})$ corresponds to the epoch when $\chi$-domination begins, given by

$$k_{\chi R}(F) = k_{\chi} F^{2/3}.$$  \hspace{1cm} (27)

We can see that for the mode $k_{\chi R} < k < k_{R}$, which corresponds to the mode which re-enters the horizon in the radiation dominated era before the $\chi$-domination, the energy density of the gravitational waves is suppressed by the factor $\sim (k_{\chi}/k_{\chi R})^2 = F^{-4/3}$ \cite{21}. On the other hand, there are no effects on the large scale with the mode $k < k_{\chi}$.

Thus the gravitational wave spectrum in the presence of late-time entropy production is completely characterized by two additional parameters, the dilution factor $F$ and the decay temperature of $\chi$, $T_{\chi}$. However, in many cases $T_{\chi}$ is expected to be very small, say, $\sim O(1)$ MeV–$O(1)$ GeV, and hence it does not affect the gravitational wave amplitude at the frequency relevant for the direct detection. Thus hereafter we mainly focus on the effect of varying $F$.

Note that non-negligible $F$ affects only the overall amplitude of the gravitational wave background for the mode $k_{\chi R} < k < k_{R}$, and hence there is a degeneracy between $F$ and the tensor-to-scalar ratio $r$ when considering direct detection around 0.1 Hz. However, the tensor-to-scalar ratio should be determined by cosmologically large-scale CMB B-mode anisotropy. Thus if future CMB experiments measure $r$, there does not remain an ambiguity coming from the degeneracy between $r$ and $F$. In other words, if the result of direct detection deviates from the expected signal from the large-scale measurement of $r$, there must be an entropy production process in the early universe.

All of these features are seen in figure 6, where the gravitational wave spectra with $F = 10^2$ and $10^4$ are shown. Here we have fixed $r = 0.1$, $T_{R} = 10^9$ GeV and $T_{\chi} = 1$ GeV. This figure shows how the gravitational wave spectrum is affected by the late-time entropy production.

If 0.1 Hz $\lesssim k_{R}(F) \lesssim 10$ Hz, both $F$ and $T_{R}$ can be determined from the shape of the gravitational wave spectrum. In figure 6 we show future sensitivity for determining the dilution factor $F$ and the reheating temperature $T_{R}$ with fixed tensor-to-scalar ratio $r$, which can be measured by CMB polarization experiments. In figure 7 we show the parameter region in the $F$-$T_{R}$ plane where both parameters can be determined by correlated analysis of DECIGO, ultimate-DECIGO (single), and ultimate-DECIGO.
Figure 7. In the outer light shaded region the gravitational wave background can be detected, and the inner blue shaded region shows the region where $T_R$ can be determined with signal-to-noise ratio 5 by correlated analysis of DECIGO, ultimate-DECIGO (single) and ultimate-DECIGO (correlation) from upper to lower. Here $r = 0.1$ is assumed. (correlation) respectively, by the dark shaded region. The light shaded region shows the parameters where the gravitational wave background is detected, but the value of $T_R$ remains undetermined. Here we have fixed $r = 0.1$ and $T_\chi = 1$ GeV. (The precise value of $T_\chi$ does not matter as long as $T_\chi \lesssim 10^3$ GeV.) On the other hand if $k_R(F) \lesssim 0.1$ Hz, we can measure only the ratio $T_R/F$. However, this ratio contains sufficient information for determining the gravitino abundance including the dilution effect. Finally if $k_R(F) \gtrsim 10$ Hz, only a lower bound on $T_R$ is obtained as $T_R \gtrsim 2 \times 10^9 F^{1/3}$ GeV. Even this case is useful for constraining the gravitino mass, as we will see.
5. Implications for particle physics

We have shown that the reheating temperature of the universe $T_R$ may be determined from future space-based laser interferometer experiments. In this section we discuss some implications of determining $T_R$ for particle physics, in particular SUSY models\(^9\) and baryogenesis mechanisms.

5.1. Supersymmetric models

5.1.1. Thermally produced gravitinos. In SUSY, there exists a superpartner of the graviton, the gravitino. The gravitino mass ranges from $O(1) \text{ eV}$ to $O(100) \text{ TeV}$ depending on SUSY breaking models, while other SUSY particles have the mass of $O(1) \text{ TeV}$. Gravitinos are produced in the early universe through scattering of particles in a thermal bath, and the resultant abundance of the gravitino is proportional to $T_R$ \([8, 51, 52]\),

$$Y_{3/2} \approx 2 \times 10^{-12} \left( 1 + \frac{m_{3/2}^2}{3m_{3/2}} \right) \left( \frac{T_R}{10^{10} \text{ GeV}} \right) \left( \frac{1}{F} \right),$$  \quad (28)

where $Y_{3/2} = n_{3/2}/s$ is the gravitino number-to-entropy ratio, $m_{3/2}$ denotes the gravitino mass and $m_{3/2} (\sim O(1) \text{ TeV})$ is the mass of the gluino, the fermionic superpartner of the gluon. If the gravitino is rather heavy ($m_{3/2} \gtrsim 1 \text{ TeV}$) and unstable, it eventually decays with lifetime typically longer than $1 \text{ s}$, producing high energy photons and hadrons. Those decay processes destroy or overproduce light elements such as $^4\text{He}$, $\text{D}$, $^7\text{Li}$ and $^6\text{Li}$. This constrains the reheating temperature as $T_R \lesssim 10^6 - 9 \text{ GeV}$ depending on the gravitino mass and its hadronic branching ratio \([8]\). On the other hand, if the gravitino is the lightest SUSY particle and stable due to the $R$-parity conservation, it contributes to total matter density of the universe \([10]\). This leads to the constraint

$$T_R \lesssim 7 \times 10^6 \text{ GeV} \left( \frac{m_{3/2}}{1 \text{ GeV}} \right) \left( \frac{m_{\tilde{g}}}{1 \text{ TeV}} \right)^{-2} F,$$  \quad (29)

for $m_{3/2} \sim 10^{-4} - 10 \text{ GeV}$.

For example, let us consider the situation where $T_R$ is revealed to be larger than $\sim 10^7 \text{ GeV}$ and negligible dilution factor ($F \sim 1$) is confirmed with future space-based laser interferometer experiments. For the unstable gravitino with $m_{3/2} \gtrsim 1 \text{ TeV}$, as is often the case with gravity-mediated SUSY breaking models, its abundance is constrained as $Y_{3/2} \lesssim 10^{-16}$ \([8]\), and hence $T_R$ should be less than around $\sim 10^6 \text{ GeV}$. Thus SUSY breaking models which predict the gravitino mass of $O(1) \text{ TeV}$ will be excluded, even if the gravitino mass might not be determined by accelerator experiments. Also for the light gravitino scenario, the following constraint is obtained:

$$m_{3/2} \gtrsim 1 \text{ GeV} \left( \frac{1}{F} \right) \left( \frac{T_R}{10^7 \text{ GeV}} \right) \left( \frac{m_{\tilde{g}}}{1 \text{ TeV}} \right)^2.$$  \quad (30)

Thus the gravitino mass with $m_{3/2} \lesssim 1 \text{ GeV}$, which can be realized in some classes of gauge-mediated SUSY breaking models \([53]\), is excluded, since otherwise the gravitino abundance exceeds the present dark matter abundance.

\(^9\) SUSY theories contain nearly twice the particle species compared with the standard model and the relativistic effective degrees of freedom at high temperature are doubled, $g_*(T \gtrsim 1 \text{ TeV}) = 228.75$. This leads to slight suppression of the gravitational wave spectrum, but does not much affect the results of section 3.
5.1.2. Non-thermally produced gravitinos. Recently it has been pointed out that a non-thermal production process of gravitinos from the inflaton decay gives a significant amount of gravitino abundance \[ Y_{3/2}^{(NT)} \approx 9 \times 10^{-11} \left( \frac{m_\phi}{10^{13} \text{ GeV}} \right)^2 \left( \frac{\langle \phi \rangle}{10^{15} \text{ GeV}} \right)^2 \left( \frac{10^6 \text{ GeV}}{T_R} \right), \]

where \( m_\phi \) and \( \langle \phi \rangle \) denote the mass and VEV of the inflaton. Since non-thermal contribution is proportional to \( T_R^{-1} \), this gives a lower bound on \( T_R \) for fixed \( m_{3/2} \). Furthermore the inflaton always decays into the MSSM sector appearing in the superpotential in supergravity, which provides a lower limit on the reheating temperature as \[ T_R \gtrsim 10 \text{ TeV} | y_t \left( \frac{228.75}{g_{*}(T_R)} \right)^{1/4} \left( \frac{m_\phi}{10^{13} \text{ GeV}} \right)^{3/2} \left( \frac{\langle \phi \rangle}{10^{15} \text{ GeV}} \right), \]

where \( y_t \) is the top Yukawa coupling.

In figure 8 the allowed parameter region for \( m_{3/2} - T_R \) is shown for the case of a stable gravitino for \( m_\phi = 10^{13} \text{ GeV} \) and \( \langle \phi \rangle = 10^{15} \text{ GeV} \). The upper left and lower right regions are excluded from thermal \((28)\) and non-thermal production \((31)\), respectively.

Gravitinos produced by the decay of the next-to-lightest SUSY particle may also be the dark matter \([54,55]\), but BBN constraints almost exclude such a possibility.

\footnote{Gravitinos produced by the decay of the next-to-lightest SUSY particle may also be the dark matter \([54,55]\), but BBN constraints almost exclude such a possibility.}
Also spontaneous decay of the inflaton gives a lower bound on $T_R$ (32) as denoted by the dotted line. It is seen that $10^5 \text{ GeV} \lesssim T_R \lesssim 10^9 \text{ GeV}$ is favored, and this range is interesting from the viewpoint of direct detection of gravitational waves.

5.2. Baryogenesis mechanism

Another important issue which is deeply related to the reheating temperature is baryogenesis. As is well known, almost all the ordinary matter in the universe consists of baryons, not anti-baryons. It is a long-standing mystery how the sizable amount of baryon asymmetry is generated after inflation. We must rely on some new physics which involves extra baryon number violation, $CP$ violation and non-equilibrium processes in order to create matter–anti-matter asymmetry. One of the most popular mechanisms for creating a correct amount of baryon asymmetry is the thermal leptogenesis scenario using the right-handed neutrino [58]. In order for this mechanism to work well, $T_R \gtrsim 10^9 \text{ GeV}$ is required [59].

Right-handed (s)neutrinos can also be produced non-thermally via right-handed sneutrino condensation [60, 61] or inflaton decay [62, 63], and $T_R \gtrsim 10^6 \text{ GeV}$ is required in order to create the observed amount of baryon asymmetry in these scenarios. Thus non-thermal leptogenesis scenarios can also be favored or disfavored from future space-based gravitational wave detectors.

Besides those leptogenesis scenarios, there are many other baryogenesis mechanisms which we do not list here [64], and many of them predict baryon asymmetry proportional to $T_R$. Thus determining $T_R$ has important implications for baryogenesis mechanisms.

6. Conclusions and discussion

In this paper we have shown that direct detection of a stochastic gravitational wave background of inflationary origin carries rich information on the early universe. In particular, the reheating temperature of the universe after inflation $T_R$ can be determined with future space-based laser interferometer experiments, DECIGO and BBO. If the non-zero tensor-to-scalar ratio is confirmed by CMB polarization observations, DECIGO or BBO will detect signals of the gravitational wave background if $T_R \gtrsim 10^5 \text{ GeV}$. Moreover, they can not only detect gravitational waves, but also determine $T_R$ if $T_R \lesssim 10^9 \text{ GeV}$, through the $T_R$ dependence of the gravitational wave spectrum. Thus combined analysis of direct detection of the gravitational wave background and CMB anisotropy measurements provides a consistency check of inflation models. It also restricts realistic models among the many inflation models, since both the tensor-to-scalar ratio and the reheating temperature crucially depend on the properties of the inflaton—its mass, potential, interaction strength, etc.

We have also discussed implications of determination of the reheating temperature for particle physics. The detection of the gravitational wave background clearly goes beyond just a probe of thermal history before BBN. In the environment of the very early universe with extremely high temperature, currently undiscovered particles predicted by some physics beyond the standard model may be efficiently produced. Thus probing this epoch is directly connected to particle physics. For example, some SUSY breaking models will be excluded or severely constrained from the so-called gravitino problem, if
future observations determine $T_R$. Also some baryogenesis scenarios, including the thermal leptogenesis scenario, may be excluded or disfavored. In SUSY axion models, constraints on $T_R$ may become much more stringent [65,66].

As a final remark, gravitational waves may also be generated from some other cosmological processes, such as preheating after inflation [67]–[69] and a first-order phase transition followed by subsequent bubble collisions [70, 71]. Although the amplitude and typical frequency from these contributions are highly model dependent, they might give complementary information on inflation models or a cosmological evolution scenario after inflation, if detected independently of primordial gravitational waves analyzed in this paper. Moreover, recently it was pointed out that decay of domain walls associated with gaugino condensation may produce a significant amount of gravitational waves and can be used as a probe of the gravitino mass [72].

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References

[1] Spergel D N et al (WMAP Collaboration), 2007 Astrophys. J. Suppl. 170 377
Komatsu E et al, 2008 Preprint 0803.0547 [astro-ph]
[2] Guth A H, 1981 Phys. Rev. D 23 347 [SPIRES]
Sato K, 1981 Mon. Not. R. Astron. Soc. 195 467
Starobinsky A A, 1980 Phys. Lett. B 91 99 [SPIRES]
For a review, see, for example, Linde A D, 2007 Preprint 0705.0164 [hep-th]
[3] Hawking S W, 1982 Phys. Lett. B 115 295 [SPIRES]
Starobinsky A A, 1982 Phys. Lett. B 117 175 [SPIRES]
Guth A H and Pi S-Y, 1982 Phys. Rev. Lett. 49 1110 [SPIRES]
[4] Kolb E W and Turner M S, 1990 The Early Universe (Reading, MA: Addison-Wesley)
[5] Kawasaki M, Kohri K and Sugiyama N, 1999 Phys. Rev. Lett. 82 4168 [SPIRES]
Kawasaki M, Kohri K and Sugiyama N, 2000 Phys. Rev. D 62 023506 [SPIRES]
Hannestad S, 2004 Phys. Rev. D 70 043506 [SPIRES]
Ichikawa K, Kawasaki M and Takahashi F, 2005 Phys. Rev. D 72 043522 [SPIRES]
[6] For a review, see Martin S P, 1997 Preprint hep-ph/9709356
[7] Khlopov M Y and Linde A D, 1984 Phys. Lett. B 138 265 [SPIRES]
Ellis J R, Kim J E and Nanopoulos D V, 1984 Phys. Lett. B 145 181 [SPIRES]
[8] Kawasaki M, Kohri K and Moroi T, 2005 Phys. Lett. B 625 7 [SPIRES]
Kawasaki M, Kohri K and Moroi T, 2005 Phys. Rev. D 71 083502 [SPIRES]
[9] Jedamzik K, 2006 Phys. Rev. D 74 103509 [SPIRES]
[10] Moroi T, Murayama H and Yamaguchi M, 1993 Phys. Lett. B 303 289 [SPIRES]
[11] Starobinsky A A, 1979 JETP Lett. 30 682
Rubakov V A, Sazin M V and Veryaskin A V, 1982 Phys. Lett. B 115 189 [SPIRES]
Abbott L F and Wise M B, 1984 Nucl. Phys. B 244 541 [SPIRES]
[12] Maggiore M, 2000 Phys. Rep. 331 283 [SPIRES]
[13] Allen B, 1988 Phys. Rev. D 37 2078 [SPIRES]
[14] Sahni V, 1990 Phys. Rev. D 42 453 [SPIRES]
[15] Abbott B and Romano J D, 1999 Phys. Rev. D 59 102001 [SPIRES]
[16] Turner M S, White M J and Lidsey J E, 1993 Phys. Rev. D 48 4613 [SPIRES]
[17] Turner M S, 1993 Phys. Rev. D 48 3502 [SPIRES]
[18] Liddle A R, 1994 Phys. Rev. D 49 3805 [SPIRES]
Liddle A R, 1995 Phys. Rev. D 51 4603 (erratum)
Probing the reheating temperature of the universe with a gravitational wave background

[19] Turner M S and White M J, 1996 Phys. Rev. D 53 6822 [SPIRES]
[20] Turner M S, 1997 Phys. Rev. D 55 435 [SPIRES]
[21] Seto N and Yokoyama J, 2003 J. Phys. Soc. Japan 72 3082
[22] Ungarelli C, Corasaniti P, Mercer R A and Vecchio A, 2005 Class. Quantum Grav. 22 S955 [SPIRES]
[23] Smith T L, Kamionkowski M and Cooray A, 2006 Phys. Rev. D 73 023504 [SPIRES]
[24] Boyle L A and Steinhardt P J, 2008 Phys. Rev. D 77 063504 [SPIRES]
[25] Smith T L, Peiris H V and Cooray A, 2006 Phys. Rev. D 73 123503 [SPIRES]
[26] Chongchitnan S and Elstathiou G, 2006 Phys. Rev. D 73 083511 [SPIRES]
[27] Friedman B C, Cooray A and Melchiorri A, 2006 Phys. Rev. D 74 123509 [SPIRES]
[28] Zhao W and Zhang Y, 2006 Phys. Rev. D 77 124001 [SPIRES]
[29] Watanabe Y and Komatsu E, 2006 Phys. Rev. D 73 123515 [SPIRES]
[30] Chiba T, Himemoto Y, Yamaguchi M and Yokoyama J, 2007 Phys. Rev. D 76 043516 [SPIRES]
[31] Nakayama K, Saito S, Suwa Y and Yokoyama J, 2008 Phys. Rev. D 77 124001 [SPIRES] [0802.2452 [hep-ph]]
[32] Seto N, Kawamura S and Nakamura T, 2001 Phys. Rev. Lett. 87 221103 [SPIRES]
[33] Bock J et al, 2006 Preprint astro-ph/0604101
[34] Amari E, Hirata C and Seljak U, 2005 Phys. Rev. D 72 123006 [SPIRES]
[35] Endo M, Hamaguchi K and Takahashi F, 2006 Phys. Rev. D 73 083508 [SPIRES]
[36] Endo M and Takahashi F, 2006 Prog. Theor. Phys. Suppl. 163 204
[37] Friedman B C, Cooray A and Melchiorri A, 2006 Phys. Rev. D 74 123509 [SPIRES]
[38] Zhao W and Zhang Y, 2006 Phys. Rev. D 77 124001 [SPIRES]
[39] Watanabe Y and Komatsu E, 2006 Phys. Rev. D 73 123515 [SPIRES]
[40] Chiba T, Himemoto Y, Yamaguchi M and Yokoyama J, 2007 Phys. Rev. D 76 043516 [SPIRES]
[41] Liddle A R and Lyth D H, 2000 Cosmological Inflation and Large Scale Structure (Cambridge: Cambridge University Press)
[42] Weinberg S, 2004 Phys. Rev. D 69 023503 [SPIRES]
[43] Kudoh H, Haruyama T and Himemoto Y, 2006 Phys. Rev. D 73 064006 [SPIRES]
[44] http://www.rssd.esa.int/index.php?project=Planck
[45] Taylor A C et al, 2004 Preprint astro-ph/0407148
[46] Buonanno A, Sigl G, Raffelt G G, Janka H T and Muller E, 2005 Phys. Rev. D 72 084001 [SPIRES]
[47] Sandick P, Olive K A, Daigne F and Vangioni E, 2006 Phys. Rev. D 73 104024 [SPIRES]
[48] Moroi T, Yamaguchi M and Yanagida T, 1995 Phys. Lett. B 342 105 [SPIRES]
[49] Kawasaki M, Moroi T and Yanagida T, 1997 Phys. Rev. D 56 6704 [SPIRES]
[50] Nakamura S and Yamaguchi M, 2007 Phys. Lett. B 655 167 [SPIRES] [0707.4538 [hep-ph]]
[51] Moroi T and Randall L, 2000 Nucl. Phys. B 570 455 [SPIRES]
[52] Kohri K, Yamaguchi M and Yokoyama J, 2004 Phys. Rev. D 70 043522 [SPIRES]
[53] Kohri K, Yamaguchi M and Yokoyama J, 2005 Phys. Rev. D 72 083510 [SPIRES]
[54] Nagai M and Nakamura K, 2007 Phys. Rev. D 76 123501 [SPIRES]
[55] Endo M, Hamaguchi K and Takahashi F, 2006 Phys. Rev. Lett. 96 211301 [SPIRES]
[56] Nakamura S and Yamaguchi M, 2007 Phys. Rev. D 77 023502 [SPIRES] [0707.4538 [hep-ph]]
[57] Endo M, Hamaguchi K and Takahashi F, 2006 Phys. Rev. D 74 023531 [SPIRES]
[58] Asaka T, Nakamura S and Yamaguchi M, 2007 Phys. Rev. D 74 023520 [SPIRES] [0707.4538 [hep-ph]]
[59] Kim J E, 1991 Phys. Rev. Lett. 67 3465 [SPIRES]
[60] Lyth D H, 1993 Phys. Rev. D 48 4523 [SPIRES]
[61] Hashimoto M, Izawa K I, Yamaguchi M and Yanagida T, 1998 Phys. Lett. B 437 44 [SPIRES]
[62] Endo M and Takahashi F, 2006 Phys. Rev. D 74 063502 [SPIRES]
[63] Kawasaki M and Nakayama K, 2008 Preprint 0802.2457 [hep-ph]
[64] Bolz M, Brandenburg A and Buchmuller W, 2001 Nucl. Phys. B 606 518 [SPIRES]
[65] Bolz M, Brandenburg A and Buchmuller W, 2008 Nucl. Phys. B 790 336 [SPIRES]
[66] Pradler J and Steffen F D, 2007 Phys. Rev. D 75 023509 [SPIRES]
[67] Pradler J and Steffen F D, 2007 Phys. Lett. B 648 224 [SPIRES]
[68] Rychkov V S and Strumia A, 2007 Phys. Rev. D 75 075011 [SPIRES]
[69] Giudice G F and Rattazzi R, 1999 Phys. Rep. 322 419 [SPIRES]
[70] Roszkowski L, Ruiz de Austri R and Choi K Y, 2005 J. High Energy Phys. JHEP08(2005)080 [SPIRES]
Probing the reheating temperature of the universe with a gravitational wave background

Cerdeno D G, Choi K Y, Jedamzik K, Roszkowski L and Ruiz de Austri R, 2006 J. Cosmol. Astropart. Phys. JCAP06(2006)005 [SPIRES]

[55] Steffen F D, 2006 J. Cosmol. Astropart. Phys. JCAP09(2006)001 [SPIRES]

[56] Kawasaki M, Takahashi F and Yanagida T T, 2006 Phys. Lett. B 638 8 [SPIRES]

Endo M, Takahashi F and Yanagida T T, 2008 Phys. Lett. B 658 236 [SPIRES]

[57] Endo M, Kawasaki M, Takahashi F and Yanagida T T, 2007 Phys. Rev. D 76 083509 [SPIRES]

[58] Kawasaki M and Yanagida T, 1986 Phys. Lett. B 174 45 [SPIRES]

[59] Buchmuller W, Peccei R D and Yanagida T, 2005 Ann. Rev. Nucl. Part. Sci. 55 311 [SPIRES]

[60] Murayama H, Suzuki H, Yanagida T and Yokoyama J, 1993 Phys. Rev. Lett. 70 1912 [SPIRES]

[61] Murayama H, Suzuki H, Yanagida T and Yokoyama J, 1994 Phys. Rev. D 50 2356 [SPIRES]

[62] Lazarides G and Shafi Q, 1991 Phys. Lett. B 258 305 [SPIRES]

[63] Asaka T, Hamaguchi K, Kawasaki M and Yanagida T, 1999 Phys. Lett. B 464 12 [SPIRES]

[64] Asaka T, Hamaguchi K, Kawasaki M and Yanagida T, 2000 Phys. Rev. D 61 083512 [SPIRES]

[65] Covi L, Kim H B, Kim J E and Roszkowski L, 2001 J. High Energy Phys. JHEP05(2001)033 [SPIRES]

[66] Kawasaki M, Nakayama K and Senami M, 2008 J. Cosmol. Astropart. Phys. JCAP03(2008)009 [SPIRES]

[67] Khlebnikov S Y and Tkachev I I, 1997 Phys. Rev. D 56 653 [SPIRES]

[68] Garcia-Bellido J and Figueroa D G, 2007 Phys. Rev. Lett. 98 061302 [SPIRES]

[69] Garcia-Bellido J, Figueroa D G and Sastre A, 2007 Preprint 0707.0839 [hep-ph]

[70] Kosowsky A, Turner M S and Watkins R, 1992 Phys. Rev. D 45 4514 [SPIRES]

Kosowsky A and Turner M S, 1993 Phys. Rev. Lett. 69 2026 [SPIRES]

[71] Easther R, Giblin J T, Lim E A, Park W I and Stewart E D, 2008 Preprint 0801.4197 [astro-ph]

[72] Takahashi F, Yanagida T T and Yonekura K, 2008 Preprint 0802.4335 [hep-ph]