Late-time Entropy Production from Scalar Decay and Relic Neutrino Temperature

Paramita Adhya and D. Rai Chaudhuri
Department of Physics, Presidency College,
86/1, College Street,
Calcutta 700073, India

Steen Hannestad
Department of Physics, University of Southern Denmark,
Campusvej 55, 5230 Odense M, Denmark
and NORDITA, Blegdamsvej 17, 2100 Copenhagen, Denmark
(Dated: 13 May 2003)

Abstract

Entropy production from scalar decay in the era of low temperatures after neutrino decoupling will change the ratio of the relic neutrino temperature to the CMB temperature, and, hence, the value of $N_{\text{eff}}$, the effective number of neutrino species. Such scalar decay is relevant to reheating after thermal inflation, proposed to dilute massive particles, like the moduli and the gravitino, featuring in supersymmetric and string theories. The effect of such entropy production on the relic neutrino temperature ratio is calculated in a semi-analytic manner, and a recent lower bound on this ratio, obtained from the WMAP satellite and 2dF galaxy data, is used to set a lower bound of $\sim 1.5 \times 10^{-23}$ Gev on the scalar decay constant, corresponding to a reheating temperature of about 3.3 Mev.
I. INTRODUCTION

The question of neutrino equilibration and subsequent decoupling in the early universe has become increasingly important over the past few years with the advent of precision cosmology [1].

The canonical textbook result for neutrino decoupling is that neutrinos decouple prior to electron-positron annihilation, leading to a final neutrino temperature which is related to the photon temperature by

\[ r = \frac{T_{\nu_0}}{T_0} = \left(\frac{4}{11}\right)^{1/3}, \]

where \( T_{\nu_0}, T_0 \), are, respectively, the relic neutrino and CMBR temperatures.

There are quite a few factors which may cause slight departure from the standard value of \( r = \left(\frac{4}{11}\right)^{1/3} \). First, the neutrinos, too, are slightly heated during the \( e^+e^- \) annihilations [2], leading to an overall increase in neutrino energy density of roughly 1%. Finite temperature QED effects lead to an additional slight heating of the neutrinos [3]. However, all of these results are within the standard model and lead to an overall increase in neutrino energy density of a little more than 1%, an effect which should be included in practical calculations.

At present this small effect is not detectable, but it has been estimated that future high-precision measurements of the CMB anisotropy could reach the required level of sensitivity [1].

In the present paper we discuss an additional factor which has effect on the value of \( r \). Various long-lived, massive fields like the gravitino, the Polonyi, the moduli, and the dilaton, which figure in supersymmetric and string theory models [4], pose cosmological problems, because their decay must affect \( \eta \) and nucleosynthesis [5]. To dilute them away, the proposal of thermal inflation [6] has been mooted. A scalar field, the flaton, is used to generate inflation at late times, typically at temperatures of about \( 10^7 \) Gev. The inflation stops when the temperature falls to the flaton mass \( \sim 10^3 \) Gev. Such a particle will go on decaying into the era of Mev-scale temperatures, and affect nucleosynthesis and the CMBR.

The effect on neutrino decoupling and nucleosynthesis has been studied [7, 8, 9, 10]. The studies do not depend materially on the details of flaton phenomenology, and apply to the Mev-scale decay of any scalar \( \phi \).

In this manuscript, we are trying to see how the outpouring of entropy, as \( \phi \) decays, heats up the \( e^-, e^+, \gamma \) plasma, and affects the ratio \( r \). The assumption of current models is that the \( \phi \) should not decay into neutrinos directly [7]. So, the decay of \( \phi \) operates in a direction...
opposite to neutrino heating, and may offset conclusions drawn on the basis of neutrino heating. The effect of changing the lower bound on the scalar decay constant $\Gamma$ should be an interesting input in current calculations of nuclear abundances and CMBR anisotropies.

Ref. [7] deals directly with the neutrino distribution function and proceeds by solving the Boltzmann equation numerically. We try, on the other hand, to adhere as far as possible to the macroscopic entities like temperature and entropy so as to proceed analytically and keep the physical processes transparent.

Of course, this means that our calculation does not reach the level of precision of a numerical solution to the Boltzmann equation. However, it does provide a reasonable bound on the involved parameters.

The plan of the paper is as follows. Section 1 is this Introduction. In section 2, the entropy production due to scalar decay, since neutrino decoupling, is estimated, and the basic relation whereby it affects the relic neutrino temperature ratio is set out. In section 3, the values of this ratio for different values of the scalar decay constant are calculated. Section 4 uses a lower bound on the relic neutrino temperature ratio from the WMAP satellite data, the 2dF galaxy survey and other related data to find a lower bound on the scalar decay constant, and discusses the results.

II. CALCULATION OF THE ENTROPY PRODUCTION

Kawasaki et al [7] find that the electron and muon/taon neutrino distribution functions show a deficit from the thermalised F.D. distribution as the reheating temperature for the scalar (and, hence, its decay constant) falls. This leads to a decrease in the $\nu_e$ and $\nu_{\mu,\tau}$ energy densities with concomitant effect on the weak interaction rates and freeze-out times. The threshold reheating temperature $T_R$ is around 7 Mev, below which the authors find the effective number of neutrino types $N_{\text{eff}}$ falls below the value 3.

So, even before decoupling, the neutrinos cannot be assigned the photon temperature. Can they be assigned a temperature at all, in particular a lower temperature which will approximately reproduce the decreased distribution? Strictly speaking, the shape of the distribution does not permit this. Ref. [7] defines $\bar{T}_\nu = [2\pi^2n_\nu/(3\zeta(3))]^{1/3}$ and $\bar{R}_E = (\rho_\nu/n_\nu)/(3.151\bar{T}_\nu)$, and finds that $\bar{R}_E$ takes values 1.00, 1.03, 1.50 for $T_R$ values 10, 3, 1 Mev, respectively, while a F.D. distribution should give $\bar{R}_E = 1$. On the other hand, these
results indicate that unless the decay constant $\Gamma$ is much smaller than that required to save standard BBN, the assumption of a neutrino temperature $T_\nu$ a little smaller than the photon temperature $T$, even before decoupling, is not a bad one, especially if one is interested only in locating the parameter region where $\Gamma$ does not affect the standard picture. After decoupling, the neutrinos are no longer affected by the effects of scalar decay, and the momentum of the neutrinos, assumed here to be massless, redshifts as $1/a$, so that one can consider the usual relic neutrino temperature-like parameter which redshifts as $1/a$ from the neutrino decoupling temperature $T_{\nu i}$.

However, the main difference between $T_\nu$ and $T$ must arise from the $e^-, e^+$ annihilations, after the decoupling era, and we are interested in seeing how this is affected by scalar decay. In what follows, we will neglect the difference between $T_{\nu i}$ and $T_i$, just at the epoch of decoupling $t = t_i$, while recognising that the decoupling temperature remains uncertain to some extent. To take this into account, we will consider a decoupling temperature range of $1-3 \text{ Mev}$.

We have to deal with four epochs. $i$ is an epoch just after electron neutrino, $\nu$, decoupling, when we take $T_{\nu i} = T_i$, and

$$g_{Si}^* \approx g_i^* = 11/2. \quad (1)$$

$f$ is an epoch after $e^-, e^+$ annihilation, which we can take to be the present, with neutrino temperature $T_\nu 0$ and photon temperature $T_0$, and

$$g_{Sf}^* = g_f^* = 2. \quad (2)$$

$a$ is the epoch of $e^-, e^+$ annihilation, when the photon temperature is $T_a$, and, although in this epoch, the effective number of degrees of freedom is actually changing, we will take

$$g_{Sa}^* \approx g_a^* \approx g_{Si}^* \approx g_i^* \approx 11/2. \quad (3)$$

A basic assumption is that at some earlier era, $\rho_\phi$ dominated the energy density. We will need a fiducial era for $\phi$, when this $\phi$-domination of the universe ends and radiation domination begins. This, we assume, happens sometime before decoupling of the three families of neutrinos, at a temperature $T_E$, with $g_{SE}^* \approx g_E^* = 43/4$.

The basic relation is

$$g_{Si}(2\pi^2/45)a_i^3T_{i}^3 + \Delta S = g_{Sf}(2\pi^2/45)a_f^3T_{f}^3. \quad (4)$$
\( \Delta S \) is the entropy poured into the plasma between \( T_i \) and \( T_f \), in accordance with the assumption that \( \phi \)-decay into neutrinos is not allowed. It is to be calculated from

\[
dS = -d(a^3 \rho_\phi)/T, \tag{5}\]

\( \rho_\phi \) being the scalar energy density at radiation temperature \( T \), radiation including all massless particles in equilibrium with photons. If \( \rho_R \) is the radiation density, temperature is defined from

\[
\rho_R = (\pi^2/30)g^*T^4. \tag{6}\]

Taking the scalar decay rate constant to be \( \Gamma \) \[^{11}\], the equation for the evolution of the scalar energy density is

\[
\frac{\partial}{\partial t} \rho_\phi + 3H \rho_\phi = -\Gamma \rho_\phi. \tag{7}\]

Defining \( \Phi = a^3 \rho_\phi \) and \( R = a^4 \rho_R \), where \( a \) is the scale factor, \((7)\) becomes

\[
\dot{\Phi} = -\Gamma \Phi, \tag{8}\]

with the solution \[^{12}\]

\[
\Phi = \Phi_E e^{-\Gamma(t-t_E)}, \tag{9}\]

where \( \Phi_E \) is the value of \( \Phi \) at \( t = t_E \). So,

\[
\dot{S} = (\Gamma/T)\Phi_E e^{-\Gamma(t-t_E)}. \tag{10}\]

In \[^{10}\], the exponential will dominate at the epochs \( i \) and \( f \), and in between, when \( \rho_R \gg \rho_\phi \). So, in the pre-exponential, we may approximate \( T \) from \((3)\) using the full radiation domination equations \( H = 1/(2t) \) and \( H^2 = 8\pi\rho_R/(3M_{Pl}^2) \). This gives the usual relation

\[
1/T = \alpha \sqrt{t}, \text{with } \alpha^2 = 2.7215 \times 10^{-19} g^{*2} Gev^{-1}. \tag{11}\]

Introducing the variable

\[
y = \Gamma t = \Gamma/(\alpha^2 T^2), \tag{12}\]

we get

\[
\Delta S = \Gamma \Phi_E e^{\Gamma t_E} \int_i^f \alpha \sqrt{t} e^{-\Gamma t} dt
\]

\[
= (1/\sqrt{\Gamma}) \Phi_E e^{\Gamma t_E} \int_i^f \alpha \sqrt{y} e^{-y} dy
\]

\[
= (1/\sqrt{\Gamma}) \Phi_E e^{\Gamma t_E} \alpha [\int_i^a \sqrt{y} e^{-y} dy]
\]

\[
+ (4/11)^{\frac{\gamma}{2}} [\int_a^f \sqrt{y} e^{-y} dy], \tag{13}\]
where $\alpha_i$ is the value of $\alpha$ with $g^* = g_i^*$ and $\alpha_f/\alpha_i = (\frac{4}{11})^{\frac{1}{3}}$. We have assumed $g^*$ to be $g_i^*$ from $t_i$ to $t_a$ and $g_f^*$ from $t_a$ to $t_f$.

A. Estimate of $\Phi_E$

Because there are as yet no firm phenomenological values for $\Phi_E$, we have to make an estimate. In the expression for $\dot{S}$ given in (10), the exponential dominates at the epochs we are interested in. So, we estimate $\Phi_E$ in the pre-exponential just to an order of magnitude, defining the epoch $E$ by taking $\rho_R = \rho_\phi$ at $t = t_E$. Just for estimating $\Phi_E$ in the pre-exponential, we use the crude approximation $a \propt 1/T$, although, actually, in the presence of entropy generation there is departure from this type of evolution, and, in fact, in the regime of full $\Phi$ domination, when $\rho_\phi \gg \rho_R$, $T \propto a^{-\frac{4}{5}}$ [12, 13]. In the regime of $\rho_R \gg \rho_\phi$, of course, we expect a closer fit to $a \propt 1/T$.

With this approximation, in the pre-exponential,

$$
\Phi_E = a_E^3 \rho_\phi E \\
= a_E^3 \rho RE \\
= (\pi^2/30) g_E^* a_E^3 T_E^4 \\
= (\pi^2/30) g_E^* a_i^3 T_i^3 T_E.
$$

B. Estimate of $T_E$

$T_E$ is the temperature when $\Phi$-domination passes into $R$-domination, i.e. when $H^2$ passes from $\frac{8\pi}{3M_P^2} \rho_\phi$ to $\frac{8\pi}{3M_P^2} \rho_R$. So, we try to find approximations for $H^2$ closer to the epoch $\rho_\phi = \rho_R$, and on either side of it. We first consider the era before this, i.e. the era of incomplete $\Phi$-domination. From

$$
\frac{\partial}{\partial t}[a^3(\rho_\phi + \rho_R)] + p_R \frac{\partial}{\partial t}a^3 = 0,
$$

and (7),(8), one obtains

$$
\dot{R} = a \Gamma \Phi.
$$

The Friedmann equation can be written as

$$
H^2 = \frac{8\pi \Phi}{3M_P^2 a^3} (1 + \frac{R}{\Phi a}).
$$
In the era of incomplete $\Phi$-domination, the term $\frac{R}{\phi_0}$ on the RHS of (16) is a correction term, and may be evaluated to the approximation

$$R - R_I = \frac{dR}{da} |_{a_I} (a - a_I),$$

(17)

where $t_I$ refers to some initial epoch such that $a \gg a_I$. Also, if it is supposed that the scalar decay produces sufficiently copious radiation, $R_I \ll R$. As a correction term is being dealt with, these approximations should not cause much deviation from the actual evolution. Then, for $t$ sufficiently later than $t_I$, but within the regime under consideration, one may write, in the correction term on the RHS of (16),

$$R \approx \frac{\Gamma}{H} \phi a,$$

(18)

using (15) and (17). As we are hoping to find only an order of magnitude estimate of a lower bound on $\Gamma$, (18) is not a bad approximation. Thus, if we go even to very early times in this era, when $\Phi = \Phi_I \approx \text{constant}$ [13], integration of (15) leads to

$$R \approx \frac{2}{5H} \phi a,$$

using the approximations $a \sim t^{\frac{4}{3}}$, $H = \frac{\dot{a}}{a}$. So, introducing a new evolution variable $x = \frac{\Gamma}{H}$, we write, for use in the correction term on the RHS of (16),

$$x = \frac{\Gamma}{H} \approx \frac{R}{\phi a} = \frac{\rho_R}{\rho_\phi}.$$

(19)

Next, we consider the era of interest to us when $\phi a \ll R$, such that $\phi a / R$ cannot be neglected, but its higher powers can. In this era of incomplete radiation domination, the Friedmann equation is put in the form

$$H^2 = \frac{8\pi R}{3M_{Pl}^2 a^4} (1 + \frac{\phi a}{R}).$$

(20)

If, well into this epoch, the correction term $\phi a / R$ on the RHS of (20) is neglected, the full radiation domination relations are found:

$$H = \frac{1}{2t}, \text{ and } a = At^{\frac{1}{2}},$$

(21)

$A$ being a constant.
has, as solution, a falling exponential in $t$, viz. $\Phi \sim e^{-\Gamma t}$. Instead of taking the falling exponential in $t$ directly, a suitable approximation to the correction term on the RHS of (20) is first worked out. Let $t_0$ be a sufficiently late epoch, when $\Phi = \Phi_0 \approx 0$. Then, for use only in the correction term on the RHS of (20), one takes

$$\Phi - \Phi_0 = \tilde{\Phi}\left(\frac{1}{t}\right) - \tilde{\Phi}\left(\frac{1}{t_0}\right) = \frac{d\tilde{\Phi}}{dt}\bigg|_{t_0}\left(\frac{1}{t} - \frac{1}{t_0}\right).$$

Neglecting $\Phi_0, 1/t_0$ compared to $\Phi, 1/t$, respectively, an approximation

$$\Phi \approx \frac{B}{t}, \quad (22)$$

will be used only in the correction term on the RHS of (20), i.e. in the correction term, the falling exponential will be approximated by a rectangular hyperbola. $B$ is a constant. A similar approximation is considered for $R$. It ought to be mentioned that $R$ refers to the total radiation present, and not only to that produced by decay. However, the change in $R$ is due to $\phi$ decay and consequent entropy production. In the absence of this decay, $\dot{R} = 0$.

Using (21) and (22) in (15), and, integrating, one obtains, for use only in the correction term on the RHS of (20),

$$R - R_E \approx 2AB\Gamma\left(\frac{t}{2} - t_E^\frac{1}{2}\right),$$

an approximation which corresponds to (17), because of $a \approx At^\frac{1}{2}$. If $t_E$ is sufficiently early compared to $t$, though within the regime under consideration, and there is sufficiently copious radiation production since $t_E$, it is sufficient to take

$$R \approx 2AB\Gamma t^\frac{1}{2}$$

in the correction term on the RHS of (20). This relation, together with (21) and (22), are now used to give, in the correction term on the RHS of (20),

$$x = \frac{\Gamma}{H} = \frac{R}{\Phi a}, \quad (23)$$

once again, as in (19).

Introducing the variable $x$ in (20), we get (20) in the form

$$\frac{\Gamma^2}{x^2} = \frac{8\pi R}{3M_P^2 a^4}(1 + \frac{1}{x}). \quad (24)$$
This is a good equation for incomplete radiation domination when $x \gg 1$. However, we will approximate $T_E$ in (14), i.e. in the pre-exponential of (10), by putting $x = x_E \approx 1$ for $t = t_E$ in (24).

Now, (19) and (23) signify that our approximations are equivalent to taking $\rho_\Phi = \rho_R$ when $\Gamma = H$, to an order of magnitude, in the correction term in $H^2$, and this is our way of bypassing lack of knowledge about the initial value of $\Phi$. This approximation has been explained in the preceding parts of this section. It differs from a common approach to the problem where decay is supposed to occur at $t = 1/\Gamma$ and $\Phi$ is put equal to the scalar mass at this epoch. Here, we allow the scalar to decay over time, but the price of bypassing knowledge about scalar energy density or mass is paid by the approximation inherent in (19) and (23) and our way of estimating $T_E$. The result for $T_E$ is

$$T_E \approx 2^{1/4} \sqrt{\Gamma / \alpha_E}. \quad (25)$$

If we use (19) in the incomplete $\phi$-domination case when $x \ll 1$, equation (16) becomes

$$\frac{\Gamma^2}{x^2} = \frac{8\pi \Phi}{3M_{Pl} a^3}(1 + x). \quad (26)$$

This gives the same value for $T_E$ if this equation is extrapolated to $x = x_E \approx 1$ for $t = t_E$.

Putting the value of $T_E$ from (25) in (14), (13) becomes

$$\Delta S = 2^{1/4} (\pi^2/30) g_E a_3^3 T_i^3 e^{\frac{1}{4}}(\alpha_i/\alpha_E) [\int_a^i \sqrt{y} e^{-y} dy$$

$$+ \left( \frac{4}{11} \right)^{1/4} (\int_a^f \sqrt{y} e^{-y} dy)]. \quad (27)$$

Now, in (1), (27),

$$a_i^3 T_i^3 = a_f^3 T_f^3, \quad (28)$$

because the temperature of the decoupled neutrinos red-shifts as $1/a$. Using (28) and (27), we get

$$r^3 + 2.431 r^3 [\int_i^a \sqrt{y} e^{-y} dy + \left( \frac{4}{11} \right)^{1/4} (\int_a^f \sqrt{y} e^{-y} dy)] = 4/11, \quad (29)$$

where we have put $r = T_o / T_0$, and taken $\alpha_i / \alpha_E = (g_i^* / g_E^*)^{1/4} = (22/43)^{1/4}$.

III. RESULTS

From (12), (11), (1), (2) and (3), we find

$$y_i = \frac{1.56675 \Gamma_0}{(T_i/Mev)^2}, y_f = \frac{2.60 \Gamma_0}{(T_f/Mev)^2}, y_a = \frac{1.56675 \Gamma_0}{(T_a/Mev)^2}. \quad (30)$$
\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|c|c|c|}
\hline
$T_i$ & $T_i = 1\text{Mev}$ & $T_i = 2\text{Mev}$ & $T_i = 3\text{Mev}$ \\
\hline
$\Gamma_0$ & $r$ & $\Gamma_0$ & $r$ & $\Gamma_0$ & $r$ \\
\hline
1.0 & 0.587 & 4 & 0.587 & 10 & 0.599 \\
1.2 & 0.608 & 5 & 0.612 & 12 & 0.620 \\
1.4 & 0.626 & 6 & 0.635 & 13 & 0.630 \\
1.5 & 0.635 & 7 & 0.654 & 14 & 0.639 \\
1.6 & 0.643 & 8 & 0.669 & 15 & 0.648 \\
1.8 & 0.657 & 10 & 0.690 & 16 & 0.656 \\
2.0 & 0.669 & 12 & 0.702 & 20 & 0.680 \\
3.0 & 0.702 & 14 & 0.708 & 24 & 0.695 \\
4.0 & 0.711 & 20 & 0.713 & 28 & 0.703 \\
10 & 0.714 & 25 & 0.714 & 32 & 0.708 \\
25 & 0.714 & 50 & 0.714 & 50 & 0.713 \\
\hline
- & - & - & 75 & 0.714 \\
- & - & - & 100 & 0.714 \\
\hline
\end{tabular}
\caption{Relic Neutrino Temperature Ratio $r$ for different values of the Scalar Decay Parameter $\Gamma_0 = \Gamma/(10^{-24}\text{Gev})$}
\end{table}

$T_i$ is a few Mev. Here, we calculate results for $T_i = 1, 2, 3 \text{Mev}$. $T_f$ being $<< 1\text{eV}$, $y_f$ can be put $\approx \infty$ because of the nature of the incomplete Gamma function $\Gamma(1.5, x)$, where

$$\int_{y_1}^{y_2} \sqrt{ye^{-y}}dy = \Gamma(1.5, y_1) - \Gamma(1.5, y_2).$$

(31)

Also, the properties of the incomplete Gamma function $\Gamma(1.5, x)$ indicate that the results should be insensitive to the precise value of $T_a$ in the range $0.3 < T_a/\text{Mev} < 0.5$. We have checked this numerically for $T_a/\text{Mev} = 0.3, 0.4, 0.5$. The values of the ratio $r$ for different values of $\Gamma_0$ are displayed in Table 1.

\section*{IV. CONCLUSIONS}

A semi-analytical method for calculating the change in the relic neutrino temperature ratio and, hence, in the effective number of neutrino species, resulting from reheating at
very low temperatures, has been presented.

While this calculation obviously does not have the accuracy of a full numerical solution of
the Boltzmann equation, it does have the merit of being transparent and easily reproducible.

In Ref. [7], a lower bound on the reheating temperature was calculated from a considera-
tion of the impact of incomplete neutrino equilibration on big bang nucleosynthesis. While
this bound is powerful, it has the problem that it is flavour dependent in the sense that
incomplete electron neutrino equilibration has a direct impact on the \( n - p \) conversion rate.

Here we use, instead, a bound which relies on energy density only. Recent results [14]
show that an overall best fit for the WMAP \( T \) and \( TE \) data, combined with the Wang et
al compilation of CMB data [15], the 2dFGRS data, the HST key project data on \( H_0 \), and
the SNI-a data on \( \Omega_m \), give bounds on \( N_{\text{eff}} \):

\[
1.9 < N_{\text{eff}} < 7.0 \text{ (95\% confidence) }. 
\] (32)

This corresponds to a lower bound on \( r \):

\[
r > 0.637. \quad (33)
\]

Table I shows that the change in \( r \) due to \( \Gamma \) is sensitive to \( T_i \), i.e. the neutrino decou-
pling temperature. The lower bound on \( \Gamma \) for a lower bound of 0.637 on the relic neutrino
temperature ratio, corresponding to (33), is found to be between 1 and \( 2 \times 10^{-24} \text{ Gev} \), 5 and
\( 7 \times 10^{-24} \text{ Gev} \), 12 and \( 15 \times 10^{-24} \text{ Gev} \), respectively, for \( T_i = 1, 2, 3 \text{ Mev} \). We conclude that to
ensure that the relic neutrino temperature ratio \( r = T_{\nu_0}/T_0 \) does not exceed the lower bound
of 0.637 [14], the scalar decay constant \( \Gamma \) must be greater than about \( 15 \times 10^{-24} \text{ Gev} \). Taking
the definition of the reheating temperature in [7], this corresponds to a reheating tempera-
ture of 3.3 MeV, which is comparable to the result found in Ref. [7]. In their calculation,
\( N_{\text{eff}} = 1.9 \) corresponds roughly to \( T_R \approx 2.2 \text{ MeV} \).

This means that even though the approximations we use are rough, the end result is very
similar to what is found from the full numerical solution. The reason is that we are only
interested in energy density, not in the underlying neutrino distribution function. On the
other hand, for BBN purposes the distribution function of \( \nu_e \) is very important, because
high energy neutrinos have more weight in the \( n - p \) conversion processes. It should also
be noted that the lower bound on \( T_R \) of 0.7 MeV found in Ref. [7] is not based on \( N_{\text{eff}} \)
alone, but rather on a detailed study of primordial abundances. This clearly illustrates
the fact that BBN bounds are highly flavour sensitive, and shows that the much simpler energy density bound from CMB and large scale structure leads to a stronger bound on the reheating temperature.

[1] R.E.Lopez, Scott Dodelson et al, Phys. Rev. Lett. 82, 3952 (1999); R. Bowen, S. H. Hansen, A. Melchiorri, J. Silk and R. Trotta, Mon. Not. Roy. Astron. Soc. 334, 760 (2002) arXiv:astro-ph/0110636.

[2] D.Dicus et al, Phys. Rev D 26, 2694 (1982); M.A.Herrera and S.Hacyan, Astrophys. J. 336, 539 (1989); N.C.Rana and B.Mitra, Phys. Rev. D 44, 393 (1991); S.Dodelson and M.S.Turner, Phys. Rev. D 46, 3372 (1992); A.D.Dolgov and M.Fukugita, Phys. Rev. D 46, 5378 (1992); S.Hannestad and J.Madsen, Phys. Rev. D 52, 1764 (1995); A.D.Dolgov, S.H.Hansen and D.V.Semikoz, Nucl. Phys. B503, 426 (1997); N.Y.Gnedin and O.Y.Gnedin, astro-ph/9712199.

[3] A.F.Heckler, Phys. Rev.D 49, 611 (1994); R.Lopez and M.S.Turner, Phys. Rev. D 59, 103502 (1999).

[4] G.F.Giudice and R.Rattazzi, Phys. Rep. 322, 419 (1999).

[5] G.D.Coughlan, W.Fischler, E.W.Kolb, S.Raby, G.G.Ross, Phys. Lett. 131B, 59 (1983); T.Banks, D.B.Kaplan, and A.E.Nelson, Phys. Rev. D 49, 779 (1994); J.Ellis, J.E.Kim, and D.V.Nanopoulous, Phys. Lett. 145B, 181 (1984); J.Ellis, D.V.Nanopoulous, and M.Quiros, Phys. Lett B 174, 176 (1986); G.German and G.G.Ross, Phys. Lett. B 172, 305 (1986); K.Enqvist, D.V.Nanopoulous, and M.Quiros, Phys. Lett. 169B, 343 (1986); O. Bertolami, Phys. Lett. B 209, 277 (1988); B. de Carlos, J.A.Casa, F.Quevedo, and E.Roulet, ibid 318, 447 (1993); T.Banks, M.Berkooz, and P.J.Steinhardt, Phys. Rev. D 52, 705 (1995); M.Dine, L.Randall, and S.Thomas, Phys. Rev. Lett. 75, 398 (1995); M.Kawasaki and T.Yanagida, Phys. Lett. B 399, 45 (1997); J.Hashiba, M.Kawasaki, and T.Yanagida, Phys. Rev. Lett. 79, 4525 (1997); S.Kasuya, M.Kawasaki and F.Takahashi, Phys. Rev. D 65, 063509 (2002); R.Allahverdi, K.Enqvist and A. Mazumdar, Phys.Rev. D 65, 103519 (2002).

[6] D.H.Lyth and E.D.Stewart, Phys. Rev. Lett. 75, 201 (1995); L.Randall and S.Thomas, Nucl. Phys. B 449, 229 (1995); T.Asaka, J. Hashiba, M.Kawasaki and T.Yanagida, Phys. Rev. D 58, 083509 (1998); T.Asaka and M.Kawasaki, Phys. Rev. D 60, 123509 (1999).

[7] M.Kawasaki, K.Kohri, and N.Sugiyama, Phys. Rev. Lett. 82, 4168 (1999); M.Kawasaki,
K. Kohri, and N. Sugiyama, Phys. Rev. D 62, 023506 (2000).

[8] M. Kawasaki, K. Kohri, and T. Moroi, Phys. Rev. D 63, 103502 (2001).

[9] K. Kohri, ibid 64, 043515 (2001).

[10] Paramita Adhya and D. Rai Chaudhuri, hep-ph/0203142 (2002).

[11] A. D. Dolgov and A. D. Linde, Phys. Lett. 116B, 329 (1982); A. Albrecht, P. J. Steinhardt, M. S. Turner, and F. Wilczek, Phys. Rev. Lett 48, 1437 (1982); L. Abbott, E. Farhi, and M. Wise, Phys. Lett. 117B, 29 (1982).

[12] R. J. Scherrer and M. S. Turner, Phys. Rev. D 31, 681 (1985).

[13] D. J. H. Chung, E. W. Kolb, and A. Riotto, Phys. Rev. D 60, 063504 (1999).

[14] Steen Hannestad, arXiv:astro-ph/0303076 (to appear in JCAP).

[15] X. Wang, M. Tegmark, B. Jain and M. Zaldarriaga, arXiv:astro-ph/0212417