We study models in which the universe exits reheating at temperatures in the MeV regime. By combining light element abundance measurements with cosmic microwave background and large scale structure data we find a fairly robust lower limit on the reheating temperature of $T_{RH} \gtrsim 4$ MeV at 95% C.L. However, if the heavy particle whose decay reheats the universe has a direct decay mode to neutrinos, there are some small islands left in parameter space where a reheating temperature as low as 1 MeV is allowed. The derived lower bound on the reheating temperature also leads to very stringent bounds on models with $n$ large extra dimensions. For $n = 2$ the bound on the compactification scale is $M \gtrsim 2000$ TeV, and for $n = 3$ it is 100 TeV. These are currently the strongest available bounds on such models.

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I. INTRODUCTION

The standard big bang model has been tested thoroughly up to temperatures around 1 MeV where big bang nucleosynthesis occurred. At much higher temperatures the universe is assumed to have undergone inflation, during which the primordial density perturbations are produced.

Towards the end of inflation the inflaton potential steepens so that slow roll is violated, and the universe enters the reheating phase. During this phase all particles which are kinematically allowed are produced, either by direct decay or from the thermal bath produced by the inflaton decay.

Finally the universe enters the radiation dominated phase at a temperature $T_{RH}$, which is a function of the inflaton decay rate. The only certain bound on this reheating temperature comes from big bang nucleosynthesis, and has in several previous studies been found to be around 1 MeV [1, 2].

It should be noted that even if the reheating temperature after inflation is much higher there can still be subsequent "reheating" phases, in the sense that reheating is defined to be a period where the energy density is dominated by an unstable non-relativistic particle species. In standard reheating this is the inflaton, but in supersymmetric models it could for instance be the gravitino.

In the present paper we update previous calculations of this reheating phenomenon, using data from cosmic microwave background and large scale structure observations. Furthermore we extend the analysis to include the possibility of having a direct decay mode of the heavy particle into light neutrinos. If the heavy particle is a scalar this decay is normally suppressed by a factor $(m_\nu/m_\phi)^2$ because of the necessary helicity flip. However, the heavy particle could either be a non-scalar particle, or it could be a pseudo-scalar like the majoron which couples only to neutrinos. Even though such models are slightly contrived it is of interest to study whether the temperature bound on reheating is significantly affected by the possibility of direct decay into neutrinos.

In section II we discuss the set of Boltzmann equations necessary to follow the evolution of all particle species. In section III present results of the numerical solution of these equation, and in section IV we compare model predictions with observational data. Finally, section V is a review of other astrophysical constraints on heavy, decaying particles, and section VI contains a discussion.

II. BOLTZMANN EQUATIONS

We follow the evolution of all particles by solving the Boltzmann equation for each species

$$\frac{\partial f}{\partial t} - H_p \frac{\partial f}{\partial p} = C_{coll},$$

where $C_{coll}$ is the collision operator describing elastic and inelastic collisions.

A. Neutrinos

Neutrinos interact with the electromagnetic plasma via weak interactions. A comprehensive treatment of this can for instance be found in Ref. [3]. The collision integrals can be written as [3]

$$C_{coll,i}(f_1) = \frac{1}{2E_1} \int \frac{d^3p_2}{2E_2(2\pi)^3} \frac{d^3p_3}{2E_3(2\pi)^3} \frac{d^3p_4}{2E_4(2\pi)^3}$$
where \( S[M]_{12 \rightarrow 34,i} \) is the spin-summed and averaged matrix element including the symmetry factor \( S = 1/2 \) if there are identical particles in initial or final states. The phase-space factor is \( \Lambda(f_1, f_2, f_3, f_4) = f_3 f_4 (1 - f_1)(1 - f_2) - f_1 f_2 (1 - f_3)(1 - f_4) \).

This collision integral can be reduced to 2 dimensions using the method developed in Ref. [3]. However, if Pauli blocking and interactions involving only neutrinos are neglected the integrals can in fact be reduced to 1 dimension, as described in Ref. [3]. In the following we use this method. The quantitative error resulting from this is quite small.

In addition to standard weak interactions we allow for a direct decay of \( \phi \) to neutrinos, \( \phi \rightarrow \nu \bar{\nu} \). If \( \phi \) is non-relativistic then each neutrino is born with momentum \( m_\phi/2 \) and in this case the collision integral is

\[
C_{\phi \rightarrow \nu_i \bar{\nu}_i} = b_{\nu_i} \frac{2\pi^2}{(m_\phi/2)^2} \Gamma_\phi n_\phi \delta(p_\nu - m_\phi/2),
\]

where \( b_{\nu_i} \) is the branching ratio into neutrino species \( i \), and \( \Gamma_\phi \) is the decay rate of the heavy particle. For simplicity we assume equal branching ratios into all neutrino species. Even is this is not the case the neutrino distribution functions will be almost equilibrated by oscillations. This means that \( b_{\nu e} \simeq b_{\nu \mu} \simeq b_{\nu \tau} \simeq b_{\nu}/3 \).

Note that if one assumes that neutrinos are in kinetic equilibrium so that they can be described by a single temperature \( T_\nu \) it is in fact possible to solve the Boltzmann equation semi-analytically. However, this is a very poor approximation for the case when there is a direct decay mode \( \phi \rightarrow \nu \bar{\nu} \).

\section*{B. \( \phi \)}

We assume the heavy particle to be completely non-relativistic. If that is the case then the Boltzmann equation can be integrated to give the following equation for the evolution of the energy density

\[
\dot{\rho}_\phi = -\Gamma_\phi \rho_\phi - 3H \rho_\phi,
\]

i.e. there are no inverse decays. This is a good approximation for all the cases covered in the present work.

We only work with masses which are low enough that there are no hadronic decay channels open. This of course severely restricts the possible models. However, if there is a hadronic branching ratio then the minimum allowed reheating temperature increases dramatically, and we are investigating what the lowest possible reheating temperature is.

\section*{C. Electromagnetic plasma}

The evolution of the photon temperature can then be found from the equation of energy conservation

\[
\frac{d\rho_T}{dt} = -3H (\rho_T + P_T),
\]

where \( \rho_T \) and \( P_T \) are the total energy density and the total pressure respectively. This equation can be rewritten as an evolution equation for \( T_\gamma \)

\[
\frac{dT_\gamma}{dt} = - \frac{1}{3H} \left[ \frac{4H \rho_\gamma + 3H (\rho_e + P_e) + 4H \rho_\nu + d \rho_\nu}{\partial \rho_T/\partial T_\gamma + \partial \rho_e/\partial T_\gamma} \right]
\]

\section*{D. Scale factor}

Finally we solve the Friedmann equation to find the scale factor as a function of time

\[
H = \frac{\dot{a}}{a} = \sqrt{\frac{8\pi G \rho_T}{3}}
\]

Altogether we solve Eq. (4) together with Eq. (1) for each neutrino species, Eq. (2) for \( \phi \), and Eq. (6) for the photon temperature, to obtain \( a(t), T_\gamma(t), \rho_\phi(t) \), and \( f_{\nu_i}(t) \).

\section*{E. Initial conditions}

Following convention we define the reheating temperature of the universe to be when

\[
\Gamma_\phi = 3H(T_{RH})
\]

To a reasonable approximation the universe is radiation dominated at this point so that

\[
H = \left( \frac{g_* \pi^2}{90} \right)^{1/2} \frac{T_{RH}^2}{M_{Pl}},
\]

where \( M_{Pl} = 2.4 \times 10^{18} \) GeV is the reduced Planck mass and \( g_* \) is the number of degrees of freedom.

This means that there is a one to one correspondence
between $\Gamma_\phi$ and $T_{RH}$,

$$T_{RH,\text{MeV}} \simeq 0.7 \Gamma_\phi^{1/2},$$

(10)

where $g_* = 10.75$ has been used. Note that the constant of proportionality is somewhat arbitrary (although it should always be of order 1), and just gives a rough idea about the thermal temperature when the universe enters the standard radiation dominated phase.

As long as the initial time is set so that $t_i \ll t(T_{RH})$ and $T_{max} \gtrsim T_{D,\nu}$, where $T_{max}$ is the maximum temperature reached by the plasma after time $t_i$ and $T_{D,\nu}$ is the neutrino decoupling temperature then the final outcome is independent of initial conditions. The universe starts out being strongly matter dominated and the final neutrino energy density, as well as the light element abundances depend only on $\Gamma_\phi$, $m_\phi$, and $b_\nu$. The initial time is found from the Friedmann equation by assuming complete domination of $\phi$ so that $t_i = \frac{2}{3} \pi G \rho_\phi, i / [8]^{-1/2}$.

F. Nucleosynthesis

One of the main observables from the epoch around neutrino decoupling is the abundance of light elements, mainly helium and deuterium. In order to calculate these abundances we have modified the Kawano nucleosynthesis code [8]. First is has been modified to incorporate the modified temperature evolution, and second the subroutines used to calculate weak interaction rates for $n \rightarrow p$ have been modified to incorporate the full numerical electron neutrino distribution coming from the solution of the coupled Boltzmann equations.

This allows us to calculate the abundance of $^4$He and D for the various models.

III. NUMERICAL RESULTS

We have solved the set of coupled Boltzmann equations for all species for the free parameters, $m_\phi$, $\Gamma_\phi$, and $b_\nu$.

The main output from this is the relativistic energy density in neutrinos, parameterized in units of the energy density of a standard model neutrino, $\rho_{\nu_0}$,

$$N_\nu = \frac{\rho_{\nu_e} + \rho_{\nu_\mu} + \rho_{\nu_\tau}}{\rho_{\nu_0}}$$

(11)

A. $b_\nu = 0$

If $b_\nu = 0$ then the equations become independent of $m_\phi$ and this case has already been covered in Ref. [8]. We present this as our first case in order to compare results with those of [8]. Fig. 3 shows the effective number of neutrino species, $N_\nu$, after complete decay of $\phi$. This figure is identical to Fig. 4 in Ref. [8].

We also test whether our results are independent of initial conditions. In Fig. 4 we show $T_\gamma(t)$ and $\rho_\phi(t)$ for $\Gamma_\phi = 6.4$ s$^{-1}$ for two different initial times, $t_i = 1.8 \times 10^{-3}$ s and $t_i = 8.8 \times 10^{-3}$ s. In both cases we assume an initial photon temperature of 2.3 MeV (we could equally well have chosen an initial temperature of 0). While the maximum temperature reached is clearly dependent on $t_i$, $T_\gamma$ and $\rho_\phi$ quickly become indistinguishable, and as long as the temperature where this happens is greater than the neutrino decoupling temperature all final results are independent of $t_i$. Furthermore, as expected [8], the photon temperature scales as $T_\gamma \propto t^{-1/4}$ during the matter dominated period and shifts to the usual $T_\gamma \propto t^{-1/2}$ once the universe becomes radiation dominated (except for a small deviation due to heating by $e^+e^-$ annihilation).

B. $b_\nu \neq 0$

Next we cover the case when $b_\nu \neq 0$. This is much more complicated to solve numerically because of the presence of the delta function $\delta(p_\nu - m_\phi/2)$ and the fact that the solution now depends on both $b_\nu$ and $m_\phi$. In Fig. 5 we show $N_\nu$ for different values of $\Gamma_\phi$ and $b_\nu$.

From this figure is clear that when $b_\nu$ is small the effective number of neutrino species becomes independent of $m_\phi$ and increasing with $\Gamma_\phi$, with $N_\nu \rightarrow 3$ for $\Gamma_\phi \rightarrow \infty$.

For the opposite case when $b_\nu = 1$ (only decay to neutrinos) the situation is the opposite. When $\Gamma \rightarrow \infty$ the limiting value is again $N_\nu = 3$. This corresponds to the case when $\phi$ decays into neutrinos, but the effective neutrino temperature after complete $\phi$ decay is higher than $T_D$.

When $\Gamma \rightarrow 0$ the effective number of neutrino species goes to infinity. This corresponds to the case when $\phi$ decays so slowly that the produced neutrinos never equili-
cause the cross section goes as $E$ very efficient conversion of neutrinos into $e^+e^-$ possible and the result is that $N$ into a sea of low energy produced electrons and positrons are immediately converted also that this effect becomes less pronounced when this rate of neutrinos is much lower. In the case where the magnetic interactions. This means that the production have a very high annihilation cross section to $t$ redistribute with the electromagnetic plasma, leaving only neutrinos.

However, there is a large intermediate region where $N_{\nu} < 3$, even for $b_\nu = 1$. The reason for this unexpected feature can be explained as follows: When high energy neutrinos ($E \gg T$) are produced by direct $\phi$ decay they have a very high annihilation cross section to $e^+e^-$, because the cross section goes as $E^2$. However, the produced electrons and positrons are immediately converted into a sea of low energy $e^+e^-$, and $\gamma$ because of electromagnetic interactions. This means that the production rate of neutrinos is much lower. In the case where the reheating temperature is very high this does not matter because the universe still has time to thermalize completely after $\phi$ decay. However, if $T_{RH} \sim T_D$, this is not possible and the result is that $N_{\nu} < 3$ because of the very efficient conversion of neutrinos into $e^+e^-$. Notice also that this effect becomes less pronounced when $m_\phi$ decreases because neutrinos are born with energies closer to $3T$, and the mismatch between forward and backward rates becomes smaller.

In Figs. 4 and 5 this effect can be seen directly on the distribution functions. In Fig. 4 which shows $b_\nu = 1$, $\Gamma_\phi = 6.4 \, s^{-1}$, and $m_\phi = 120 \, \text{MeV}$, it can be seen that the distribution function is higher than thermal at high energies because of $\phi$ decay. However, there are fewer low energy neutrinos because of the inefficient production via $e^+e^-$ annihilation.

Conversely, in Fig. 5 which shows $b_\nu = 1$, $\Gamma_\phi = 50 \, s^{-1}$, and $m_\phi = 120 \, \text{MeV}$, it can be seen that the decay rate is high enough that neutrinos equilibrate with the electromagnetic plasma, except for a small deviation around $p_\nu = m_\phi/2$. This subsequently leads to $N_{\nu} \simeq 3$ after complete $\phi$ decay.

IV. COMPARISON WITH DATA

In order to constrain the parameters $b_\nu$, $\Gamma_\phi$, and $m_\phi$ we compare the predicted values of $N_{\nu}$, $^4\text{He}$, and $D$ with the observationally determined values. In addition to the parameters directly related to $\phi$ the nucleosynthesis outcome depends crucially on the baryon density, $\eta = n_B/n_\gamma$.

Taken at face value the recent CMB data from the WMAP satellite constrain $\eta$ tightly. However, it has been shown that there is a significant correlation between $\eta$ and $N_{\nu}$ in the CMB data. This means that it is not possible to take CMB constraint on $\eta$ directly and apply it to the nucleosynthesis calculations. Rather a full CMB likelihood analysis for $N_{\nu}$ and $\eta$ must be carried out. This can then be combined with the nucleosynthesis likelihood analysis for $b_\nu$, $\Gamma_\phi$, $m_\phi$, and $\eta$.

First the following subsection covers the current observational status, then the next covers the constraints on decay parameters which can be obtained.

A. Observational data

1. Light element abundances

The primordial helium abundance has been derived by two independent groups. Fields and Olive find the value

$$Y_P = 0.238 \pm 0.002 \pm 0.005,$$  \hspace{1cm} (12)

whereas Izotov and Thuan find

$$Y_P = 0.244 \pm 0.002 \pm 0.005$$  \hspace{1cm} (13)

Because of this inconsistency we blow up the error bars on $Y_P$ and use the value

$$Y_P = 0.238 \pm 0.015,$$  \hspace{1cm} (14)

which encompasses the allowed regions of both observational determinations.

The most recent determination of the primordial deuterium abundance has yielded the value

$$D/H = (2.78 \pm 0.29) \times 10^{-5}$$  \hspace{1cm} (15)
A determination of the primordial lithium abundance has also been performed by several groups. However, this measurement is prone to large systematics and we refrain from using it here.

2. Cosmic microwave background

The CMB temperature fluctuations are conveniently described in terms of the spherical harmonics power spectrum

$$C_l \equiv \langle |a_{lm}|^2 \rangle,$$

where

$$\frac{\Delta T}{T}(\theta, \phi) = \sum_{lm} a_{lm} Y_{lm}(\theta, \phi).$$

Since Thomson scattering polarizes light there are additional power spectra coming from the polarization anisotropies. The polarization can be divided into a curl-free ($E$) and a curl ($B$) component, yielding four independent power spectra: $C_{T,l}, C_{E,l}, C_{B,l}$ and the temperature $E$-polarization cross-correlation $C_{TE,l}$.

The WMAP experiment have reported data on $C_{T,l}$ and $C_{TE,l}$, as described in Ref. [10, 11, 12].

We have performed the likelihood analysis using the prescription given by the WMAP collaboration which includes the correlation between different $C_l$'s [10, 11, 12]. Foreground contamination has already been subtracted from their published data.

3. Large scale structure

The 2dF Galaxy Redshift Survey (2dFGRS) [13] has measured the redshifts of more than 230 000 galaxies with a median redshift of $z_m \approx 0.11$. An initial estimate of the convolved, redshift-space power spectrum of the 2dFGRS has been determined [14] for a sample of 160...
FIG. 4: The distribution function for $\nu_e$ for different values of $T_\gamma$ when $\Gamma = 6.4$ s$^{-1}$, $b_\nu = 1$, and $m_\phi = 120$ MeV. The dotted line is for $T_\gamma = 2.18$ MeV, the dashed for $T_\gamma = 0.42$ MeV, the long-dashed for $T_\gamma = 0.19$ MeV, and the full line for $T_\gamma = 0.01$ MeV. The full grey (red) line is an equilibrium distribution with $T_\nu = T_\gamma$.

FIG. 5: The distribution function for $\nu_e$ for different values of $T_\gamma$ when $\Gamma = 50$ s$^{-1}$, $b_\nu = 1$, and $m_\phi = 120$ MeV. The dotted line is for $T_\gamma = 7.7$ MeV, the dashed for $T_\gamma = 0.93$ MeV, the long-dashed for $T_\gamma = 0.23$ MeV, and the full line for $T_\gamma = 0.01$ MeV. The full grey (red) line is an equilibrium distribution with $T_\nu = T_\gamma$.

Nucleosynthesis is affected both by the expansion rate around $T \sim 0.1-1$ MeV, and by the electron neutrino distribution function. The reason is that electron neutrino enter directly in the weak reactions which interconvert protons and neutrons.

The specific neutrino distributions are therefore found as functions of temperature and used in a modified version of the Kawano BBN code [6]. This is then used to calculate primordial abundances of deuterium and helium.

For calculating the theoretical CMB and matter power spectra we use the publicly available CMBFAST package [27]. As the set of cosmological parameters we choose $\Omega_m$, the matter density, $\Omega_b$, the baryon density, $H_0$, the Hubble parameter, $\tau$, the optical depth to reionization, $Q$, the normalization of the CMB power spectrum, $b$, the bias parameter, and the effective number of neutrino species $N_\nu$, found from the solution of the Boltzmann equations. We assume neutrinos to be almost massless. We restrict the analysis to geometrically flat models $\Omega_m + \Omega_\Lambda = 1$.

For each individual model we calculate $\chi^2$ in the following way: Given a theoretical CMB spectrum the $\chi^2$ of the WMAP data is calculated using the method described in Ref. 12. With regards to the 2dF data we use the data points and window functions from Ref. 28 (http://www.hep.upenn.edu/~max/2df.html). 68% and 95% confidence levels from the data are calculated from $\Delta \chi^2 = 2.31$ and 6.17 respectively.
Freeze-out there are many more protons than neutrons. When $E_\nu \gg m_n - m_p$ the weak absorption cross section is equal on protons and neutrons. This means that additional neutrinos at high energies will have the net effect of converting protons into neutrons, so that in the end more helium is produced. Note that this is the opposite effect of just increasing the weak interaction rates, in which case less helium would be produced. The phenomenon is quite similar to what happens if $\phi$ has a hadronic decay channel. In that case pions and kaons will be produced, which subsequently convert protons to neutrons and lead to overproduction of helium.

In Fig. 6 we show 68%, 95%, and 99.99% confidence exclusion plots for $\Gamma_\phi$ and $m_\phi$, marginalized over $\eta$.

Both when $b_\nu$ is small and when $b_\nu = 1$ the bound on $T_{RH}$ becomes independent of $m_\phi$. In both cases the 95% bound is $T_{RH} \gtrsim 4$ MeV.

However, there is an intermediate regime for $b_\nu$ which allows for much lower values of $T_{RH}$. The reason for this can be seen directly from Fig. 8 in KKS, i.e. there is an intermediate range where $N_\nu$ can be kept close to 3, even for low $\Gamma_\phi$. However, for large masses (which is of course by far the most likely) there is no allowed region. The reason is the one given in the previous section: More high energy neutrinos will produce more helium, and this in turn will conflict with observations.

The final outcome is that for almost all values of $m_\phi$ and $b_\nu$ there is a robust lower bound on $T_{RH}$ which is around 4 MeV. However there is a small region where $b_\nu \sim 0.9, m_\phi \lesssim 40$ MeV where a reheating temperature as low as roughly 1 MeV is allowed.

### V. OTHER CONSTRAINTS

If $\phi$ is a scalar then the decay rate to neutrinos is normally suppressed by a factor $m_\phi^2$ because of the necessary helicity flip. Therefore the simplest assumption is that $\phi$ has no branching into neutrinos. If for instance the heavy particle is a pseudo-scalar like the axion, then there is an upper bound on the coupling to photons $g_{\phi\gamma} \lesssim 0.5 \times 10^{-10}$ GeV$^{-1}$ for $m_\phi \lesssim 30$ keV. For higher masses the bound is significantly weaker. However, even if this bound is used together with the decay width $\Gamma_{\phi \to 2\gamma} = g_{\phi\gamma}^2 m_\phi^3/64\pi$ then we find that

$$\Gamma_{\phi \to 2\gamma} \lesssim 50 m_\phi^3,10\text{MeV} \text{ s}^{-1},$$

which is easily satisfied for the parameter space we are considering.

On the other hand, if $\phi$ is a particle like the majoron which couples only to neutrinos then the decay width is

$$\Gamma_{\phi \to \nu\bar{\nu}} = \frac{g_{\phi\nu}^2 m_\phi}{16\pi} \sim 3 \times 10^{19} g_{\phi\nu}^2 m_\phi,\text{MeV} \text{ s}^{-1}$$

The bound on the dimensionless coupling constant comes from BBN as well as supernova considerations and is of
FIG. 7: 68%, 95%, and 99.99% confidence exclusion plot of the parameters $\Gamma_\phi$ and $m_\phi$ using all available data (CMB+LSS+BBN). The top left plot is for $b_\nu = 0.1$, the top right for $b_\nu = 0.5$, the bottom left for $b_\nu = 0.9$, and the bottom right for $b_\nu = 1.0$.

order $10^{-6} - 10^{-5}$ for majorons in the MeV mass range [31, 33]. For more massive majorons the bound weakens. Again it is clear that the decay parameters which we consider here are not excluded by any other astrophysical or experimental data.

The final conclusion is that heavy, decaying particles such as the ones considered here cannot be directly excluded by any current data. Furthermore a branching ratio into neutrinos can be anywhere from 0 to 1.

VI. CONCLUSION

We have carefully calculated constraints on models with extremely low reheating temperature, where a massive particle decays around $T \sim 1$ MeV. By combining constraints on light element abundances with constraints on $\eta$ and $N_\nu$ from CMB and large scale structure we derived a fairly robust limit of

$$T_{RH} \gtrsim 4 \text{ MeV}. \quad (21)$$

This bound is a significant improvement over the previous bound of $T_{RH} \gtrsim 0.7$ MeV, calculated from BBN alone. It is interesting that the lower bound is significantly higher than the $n \leftrightarrow p$ conversion freeze-out temperature, $T \sim 0.8$ MeV, and even higher than the neutrino decoupling temperature $T_D \sim 2$ MeV. This shows that even small residual effects can be measured with present observational data.

Models with reheating temperature in the MeV regime are in general difficult to reconcile with such features as baryogenesis. However, in models with large extra dimensions a low reheating temperature is essential in order to avoid overproduction of massive Kaluza-Klein graviton states. This means that we can use our present bound to derive limits on the compactification scale in such models. For the case of two extra dimensions the bound is $M \gtrsim 2000$ TeV and for $n = 3$ it is $M \gtrsim 100$
TeV. This bound is somewhat stronger than the bound coming from considerations of neutron star cooling and gamma ray emission.

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