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BIG BANG NUCLEOSYNTHESIS CONSTRAINTS ON MASSIVE, UNSTABLE NEUTRINOS*

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Abstract

The tau-neutrino, if sufficiently massive, must be unstable. Big Bang Nucleosynthesis (BBN) can provide constraints on the ν_τ mass and lifetime. The modification to the energy density of the early Universe in the case of a massive τ-neutrino which decays via $\nu_\tau \rightarrow \nu_\mu + \phi$ (where ϕ is a weakly coupled massless scalar) is described and the results of BBN production of the light elements is presented. Consistency with the primordial abundances of D, ^{3}He,^{7}Li and, especially, ^{4}He leads to constraints on the mass (m_{ν_τ}) and lifetime (τ_{ν_τ}) of the tau-neutrino. Very massive ν_τ ($m_{\nu_\tau} \gtrsim$ 5-10 MeV), up to the ARGUS bound of 31 MeV, are only allowed for short lifetimes ($\lesssim$ 40 sec.). Much lighter ($m_{\nu_\tau} \lesssim 0.1$ MeV) ν_τ are permitted for lifetimes longer than ~ 0.01 sec but, $m_{\nu_\tau}(MeV) \lesssim 10\tau_\nu$ (sec) for shorter lifetimes.

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1. Introduction

Big Bang Nucleosynthesis (BBN) has proven to provide a valuable probe of high energy physics (see ref.1 and further references therein). New particles, often suggested by Grand- and Super-Unified extensions of the standard model, may be constrained by BBN even when they are too weakly interacting to have been found at high energy accelerators. As Steigman, Schramm and Gunn[2] showed in particular, the primordial abundance of ^{4}He can provide significant constraints on new, weakly interacting particles. This approach has often been exploited in recent years to constrain, or even exclude, proposed extensions of the standard model. For just one example, the interested reader is referred to the Bertolini & Steigman[3] constraints on models which tried to incorporate a 17 keV neutrino.

There are several reasons for the appeal of this approach to particle physics via BBN. Firstly, in the context of the standard model of cosmology – the hot big bang model – BBN was inevitable. That is, during its early evolution an expanding Universe filled with radiation passed through an epoch where the temperature and density were sufficiently high that complex nuclei could be built from nucleons. Secondly, in the standard model the predicted abundances of the newly synthesized nuclides depend on only one free parameter – the nucleon-to-photon ratio $\eta \equiv N/\gamma$, $\eta_{10} \equiv 10^{10}\eta$. Furthermore, the nuclear reaction network is simple and, the cross sections are measured at lab energies which correspond to the thermal energies at BBN. Finally, for a very narrow range of $\eta(3 \lesssim \eta_{10} \lesssim 4)$ consistency is found[4] between the predictions of BBN and the primordial abundances of D, ^{3}He, ^{4}He and ^{7}Li inferred from astronomical observations. Any "new physics" should not disturb this success of the standard model.

Here, I will present an overview of recent work[5–7] utilizing this approach to constrain the mass and lifetime of the tau-neutrino. It is by now well known (see, e.g., ref 8) that, to avoid a Universe which at present is too young, a stable τ-neutrino must either be very light ($\lesssim 40eV$) or very massive ($\gtrsim 3GeV$). The ARGUS bound[9] ($m_{\nu_\tau} \lesssim 31MeV$) ensures that, if sufficiently massive ($\gtrsim 40eV$), the ν_τ must be unstable. In recent work my colleagues and I have studied the evolution of a massive unstable neutrino and its decay daughters[5–7]. Applying these results to a massive ν_τ which decays into ν_μ and a weakly coupled scalar ϕ (e.g., majoron or familon), we have followed in detail the contributions of ν_τ, ν_μ and ϕ to the total energy density at the epoch of BBN. For a large range of ν_τ masses and lifetimes the universal expansion rate at BBN is modified (compared to the "standard" case of a massless ν_τ) and the predicted abundances of the light elements deviate from the standard model case[7]. In the next section I will describe the evolution of the abundances of ν_τ, ν_μ and ϕ and how they depend on m_{ν_τ} and τ_{ν_τ}. Next, I will show how the total energy density at BBN changes as a function of m_{ν_τ} and τ_{ν_τ} and, these results will be used to understand the modified abundances predicted for BBN. Then, comparing to the observational bounds (primarily on ^{4}He and D + ^{3}He), I derive the allowed/excluded regions in the m_{ν_τ} vs. τ_{ν_τ} plane. I conclude with a summary of the constraints on m_{ν_τ} and τ_{ν_τ} and with some comments regarding alternate decay channels (e.g., $\nu_\tau \to \nu_e + \phi$, $\nu_\tau \to \nu_\mu + \gamma$).

2. Evolution Of An Unstable ν_τ And Its Decay Products

There are three time (or, temperature) scales of relevance to the evolution of the abundances of a massive, unstable τ-neutrino and its decay products (ν_μ, ϕ). For $T \leq T_{NR} \approx m_{\nu_\tau}/3$, the ν_τ becomes non-relativistic and its contribution to the energy density is dominated by its mass rather than its momentum ($\rho_{NR} \sim mn \sim mT^3$, $\rho_{ER} \sim Tn \sim T^4$). For $T \lesssim T_D$, electroweak decoupling occurs in the sense that the rate of such reactions as $\nu_\tau + \bar{\nu}_\tau \leftrightarrow e^+ + e^-$ becomes slow compared to the universal expansion rate. For light ν_τ, $T_D \sim 2$–3 MeV; for heavy ν_τ, $T_D \gtrsim 2$–3 MeV and is mass dependent. When T exceeds T_D, decaying ν_τ are quickly regenerated by two-body reactions among $\nu_e\bar{\nu}_e$, $\nu_\mu\bar{\nu}_\mu$ and $e^\pm$. The third scale of interest is T_{decay}, the temperature at which the bulk of the unstable ν_τ begin decaying. If, for a semi-quantitative estimate, we take the age-temperature relation[8] to be: $t(\text{sec}) \approx T_{MeV}^{-2}$, then for a rest lifetime of $\tau_{\nu_\tau}(\text{sec})$ we have $T_{decay} \approx \tau_{\nu_\tau}^{-1/2}$ for NR decays and $T_{decay} \approx (m_{\nu_\tau}/3\tau_{\nu_\tau})^{1/3}$ for ER decays where, in the Lorentz factor $E_{\nu_\tau}/m_{\nu_\tau}$, we have taken $E_{\nu_\tau} \approx 3T$.

The six possible orderings of T_{NR}, T_D and T_{decay} determine the dominant processes which influence the evolution of n_{ν_τ}, n_{ν_μ} and n_ϕ (and, their energy densities). In Figure 1 (taken from ref. 7) we show the evolution of number densities with temperature for examples of the six possible orderings. Some brief comments may help clarify the various behaviors in Fig. 1; for further details see refs. 5-7.

1. $T_D > T_{decay} > T_{NR}$: (See Fig.1a) Decays and inverse decays regulate the n_{ν_τ}, n_{ν_μ}, n_ϕ abundances. Since the ν_μ and ν_τ are decoupled, their abundance exceeds the equilibrium value by the same amount that the ν_τ abundance is below the equilibrium value: $n_\phi = \Delta n_{\nu_\mu} = -\Delta n_{\nu_\tau}$. Although decays begin when the ν_τ are ER, the bulk of the ν_τ decays occur when $T \lesssim T_{NR}$ where inverse decays are kinematically suppressed (off scale in Fig. 1a).

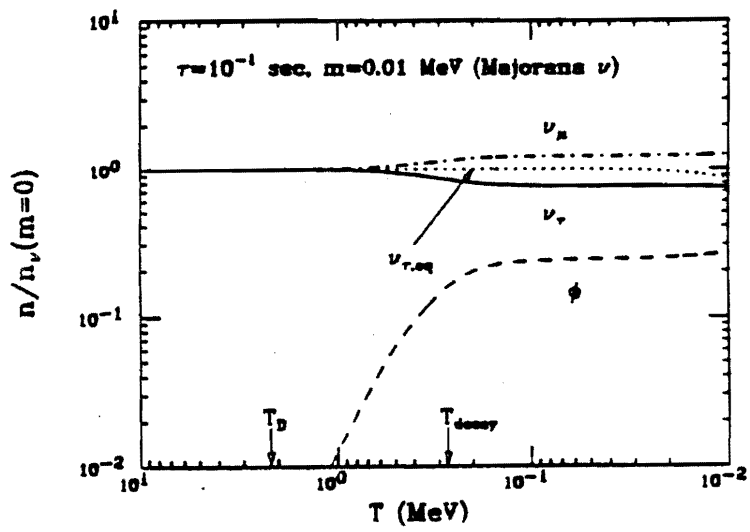

Fig. 1a. The number densities of ν_τ, ν_μ and ϕ, scaled to that of a massless neutrino, as a function of temperature for $T_D > T_{decay} > T_{NR}$.

2. $T_{decay} > T_{NR} > T_D$: (See Fig.1b) The ν_τ begin decaying when they are ER so that inverse decays are important. Since the ν_τ and ν_μ are still coupled to the background plasma ($\nu_e\bar{\nu}_e$, $e^\pm$) the ν_τ and ν_μ abundances are kept close to their equilibrium values while the ϕ abundance increases (decaying ν_τ are regenerated by, e.g., $e^\pm \to \nu_\tau\bar{\nu}_\tau$ and, can continue to decay and produce more ϕ). Since $T_{NR} > T_D$ (albeit not by much in Fig. 1b), annihilations and decays reduce n_{ν_τ} and enhance n_{ν_μ} and n_ϕ for $T \lesssim T_{NR}$.

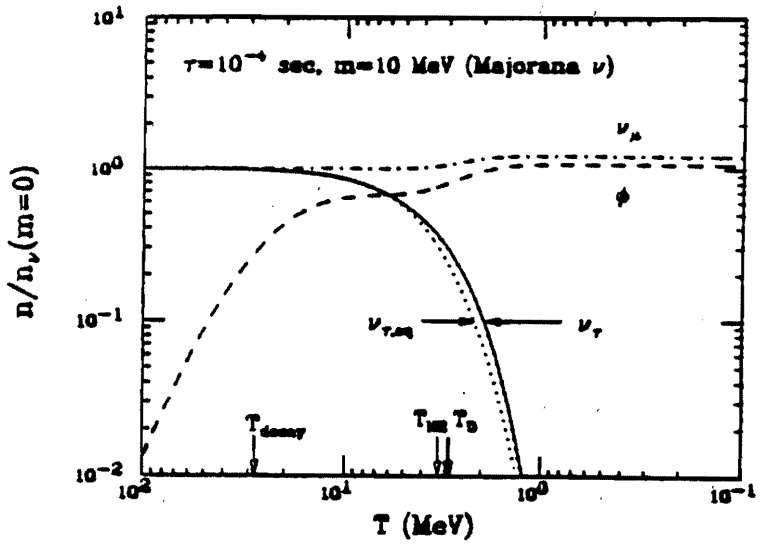

Fig. 1b. Same as 1a. for $T_{decay} > T_{NR} > T_D$.

3. $T_{decay} > T_D > T_{NR}$: (See Fig.1c) This case is very similar to 2. above. Decays begin when ν_τ is ER and, for $T_D > T > T_{NR}$, decays and inverse decays regulate the ν_τ, ν_μ, ϕ abundances. The bulk of the decays occur when $T \lesssim T_{NR}$.

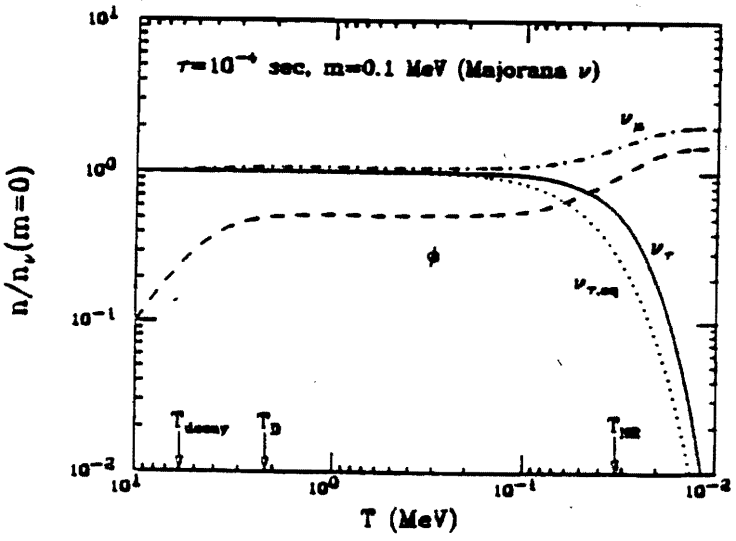

Fig. 1c. Same as 1a. for $T_{decay} > T_D > T_{NR}$.

4. $T_D > T_{NR} > T_{decay}$: (See Fig.1d) Here, the ν_τ decay when NR, after weak decoupling. Since the ν_τ have decoupled before becoming NR, annihilations are inefficient in reducing their abundance (when $T \lesssim T_{NR}$) and, almost an equilibrium (m = 0) abundance of ν_τ is available for decays at $T \lesssim T_{decay}$. Hence the large abundance of ϕ and the large enhancement in ν_μ ($n_\phi \approx n_\nu^{eq}(m=0) \approx \Delta n_{\nu_\mu}$).

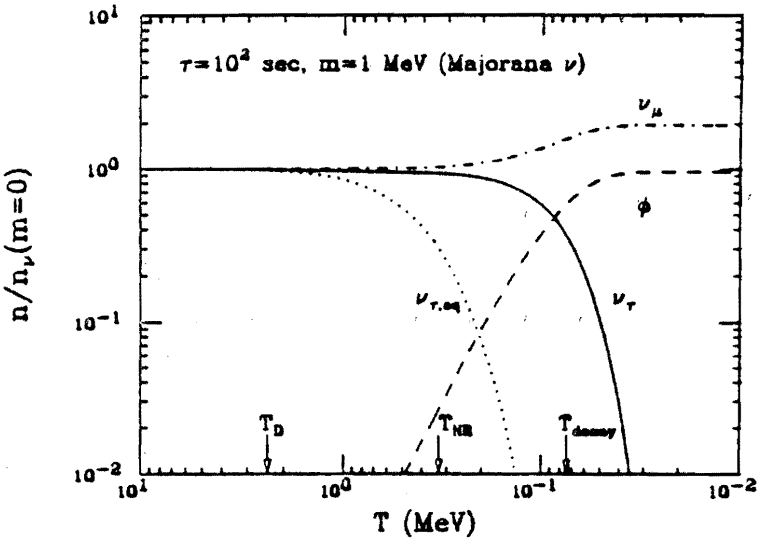

Fig. 1d. Same as 1a. for $T_D > T_{NR} > T_{decay}$.

5. $T_{NR} > T_D > T_{decay}$: (See Fig.1e) First the ν_τ become NR when still coupled. This leads to a reduction in the ν_τ abundance prior to decay with a corresponding suppression of ϕ production and ν_μ enhancement.

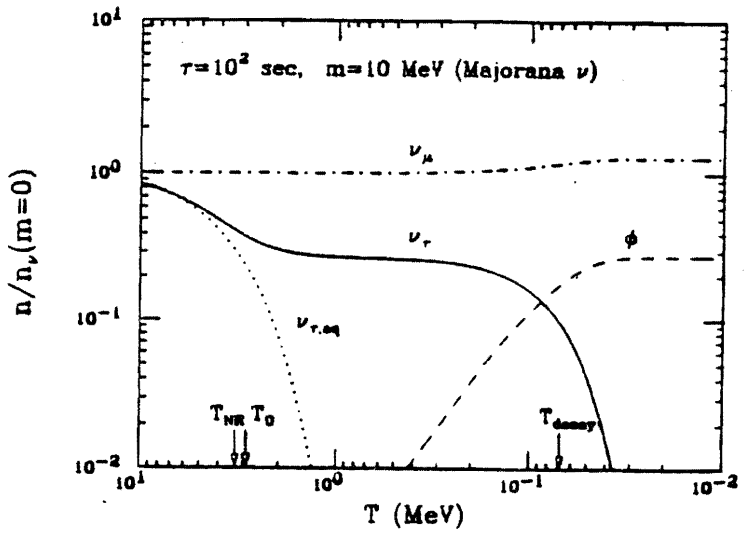

Fig. 1e. Same as 1a. for $T_{NR} > T_D > T_{decay}$.

6. $T_{NR} > T_{decay} > T_D$: (See Fig.1f) Here too, as in 5. above, the ν_τ became NR when still coupled and fewer survive to decay, reducing the ν_μ/ϕ contribution.

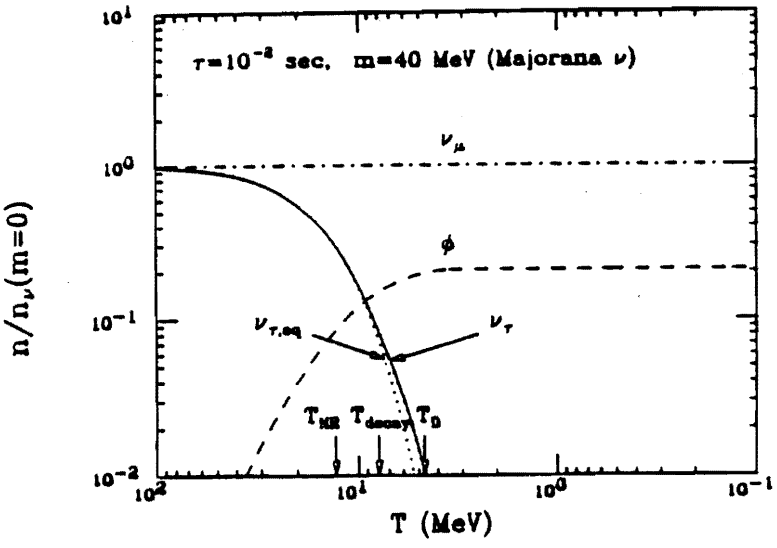

Fig. 1f. Same as 1a. for $T_{NR} > T_{decay} > T_D$.

3. The Expansion Rate At BBN

During the early evolution of the Universe the expansion rate ($H \equiv R^{-1}dR/dt$; R is the scale factor) is determined by the total energy density.

$$H = (8\pi G\rho/3)^{1/2}. \tag{1}$$

Without loss of generality we may write for the total energy density at BBN,

$$\rho = \rho_\gamma + \rho_e + N_\nu \rho_\nu \ (m = 0). \tag{2}$$

In (2) we have counted the contribution from photons and from $e^\pm$ pairs; all other contributions are given by N_ν, in terms of the number of equivalent massless neutrinos. For example, in the standard model with three families of massless, left-handed neutrinos, $N_\nu = 3$ and $\rho/\rho_\nu\ (m = 0) = 43/7$. For a massive, unstable ν_τ, N_ν will deviate from $N_\nu = 3$ and the expansion rate at BBN will differ from its standard value. Indeed, for sufficiently massive, sufficiently long-lived ν_τ, the Universe may be "matter dominated" at BBN (where $\rho \sim mT^3$ rather than $\rho \sim T^4$). This "m/T effect" ensures that, for a wide range of m_{ν_τ} and τ_{ν_τ}, N_ν is temperature dependent at

BBN. In Fig. 2 a-c the N_ν – T relation is shown for three choices of ν_τ lifetime and a range of ν_τ masses. Roughly speaking, those cases where $N_\nu \gtrsim 3.3$ for $T \lesssim 1 MeV$ are "dangerous". However, since N_ν is, in general, temperature dependent, the effect on BBN can only be quantitatively determined by a full BBN calculation.[7]

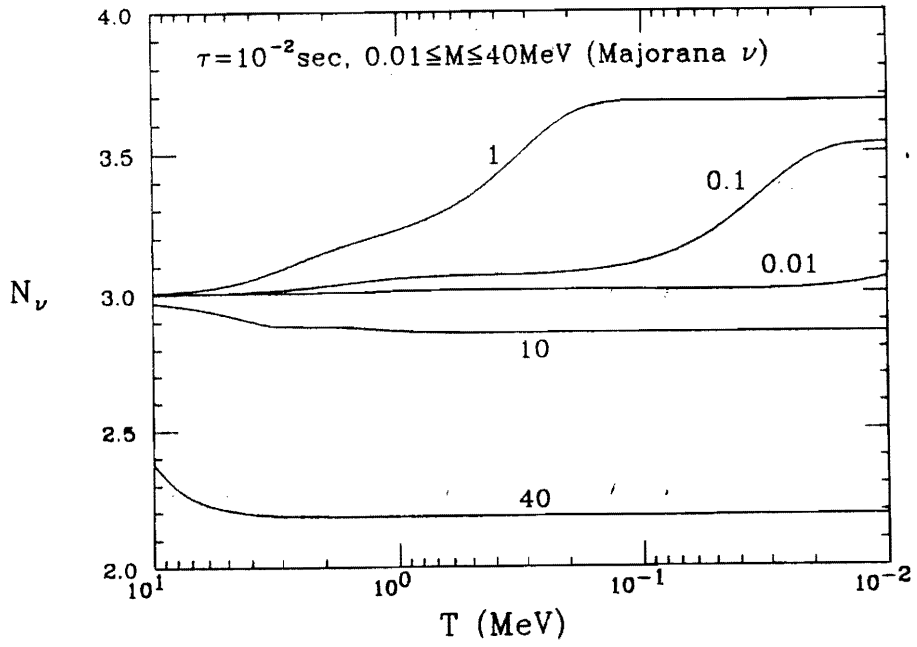

Fig. 2a. The equivalent number of light neutrinos, N_ν as a function of temperature for ν_τ masses in the range 0.01–40 MeV and a lifetime of 0.01 sec.

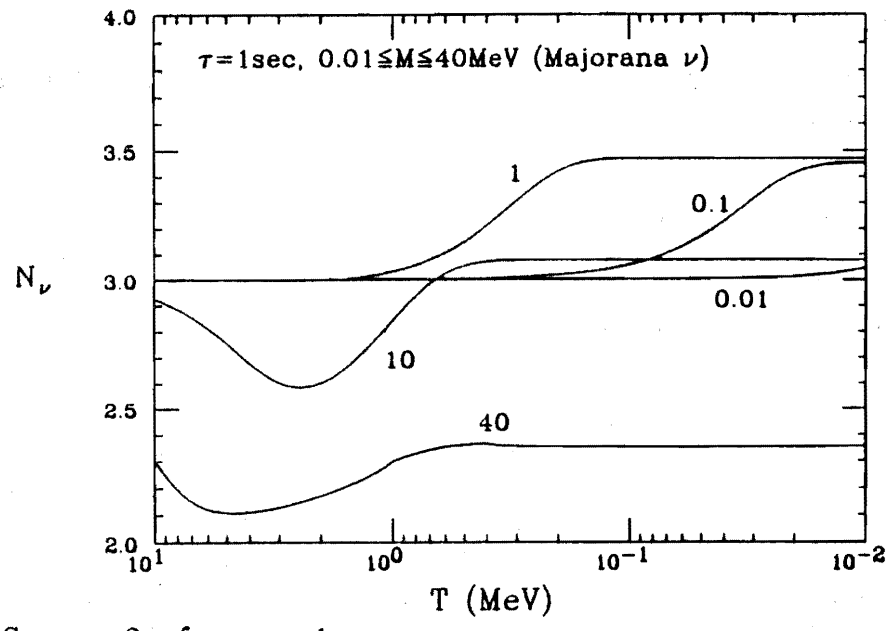

Fig. 2b. Same as 2a. for $\tau_{\nu_\tau} = 1$ sec.

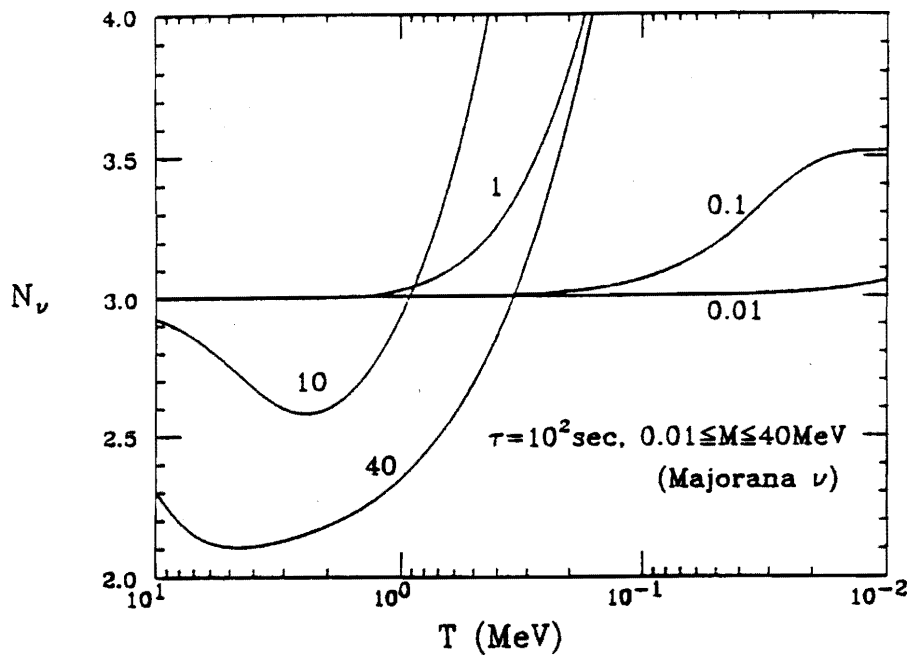

Fig. 2c. Same as 2a. for $\tau_{\nu_\tau} =$ 100 sec.

4. BBN With A Massive, Unstable Tau Neutrino

We have incorporated the modifications of the universal expansion rate due to a massive, unstable ν_τ into the BBN code[7] and have calculated the resulting abundances of the light elements (D, ^{3}He, ^{4}He, ^{7}Li) as a function of m_{ν_τ} and τ_{ν_τ}. A selection of our results is displayed in Fig. 3 a-c where for three choices of τ_{ν_τ} we show the iso-abundance curves for D $+^3$He ($y_{23p} = 10^{-4}$), ^{7}Li ($y_{7p} = 1.4\times10^{-10}$) and ^{4}He($Y_p =$ 0.240) as a function of the ν_τ mass. As expected, in each case the contours for D $+$ ^{3}He and ^{7}Li are relatively insensitive to the deviations in N_ν from 3 (i.e., insensitive to *small* changes in the universal expansion rate at BBN). However, as expected, the ^{4}He yield is very sensitive to N_ν at BBN. In general, as m_{ν_τ} increases, so does N_ν. This would lead to enhanced ^{4}He production.[2] To keep Y_P at 0.240, it is necessary that the nucleon to photon ratio η decrease; hence the shift in the Y_p curve to the left in the $m_{\nu_\tau} - \eta$ plane. Eventually, however, sufficiently massive ν_τ disappear (via annihilation) before BBN and their contribution to N_ν is reduced. Indeed, in the limit of very massive ν_τ, $N_\nu \rightarrow 2$. Thus, as m_{ν_τ} increases the ^{4}He iso-abundance curve swings back towards higher η.

As can be seen in Figures 3, the requirement that the primordial yield of D $+^3$He not be too large[4] ($y_{23p} \leq 10^{-4}$), coupled with an upper bound on ^{4}He ($Y_p \leq 0.240$) defines the allowed regions (the hatched regions in Fig. 3) in the $m_{\nu_\tau} - \eta$ plane for each choice of ν_τ lifetime.

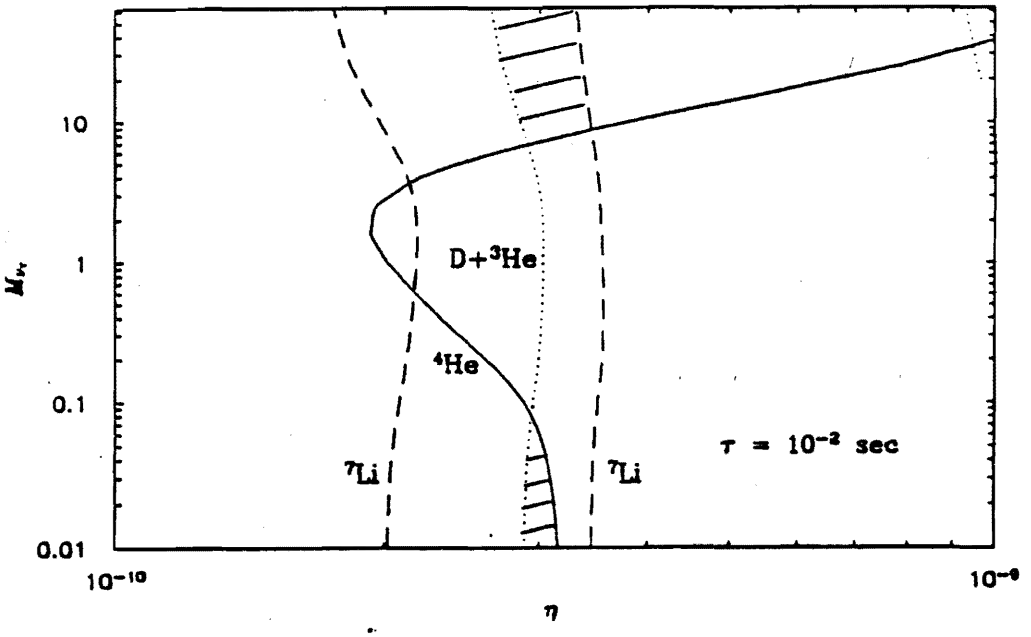

Fig. 3a. Iso-abundance contours for (D +^{3}He)/H = 10^{-4}, $Y = 0.240$ and ^{7}Li/H = 1.4×10^{-10} as a function of the nucleon-to-photon ratio η and tau neutrino mass for ν_τ lifetime 0.01 sec. The hatched regions are allowed by BBN.

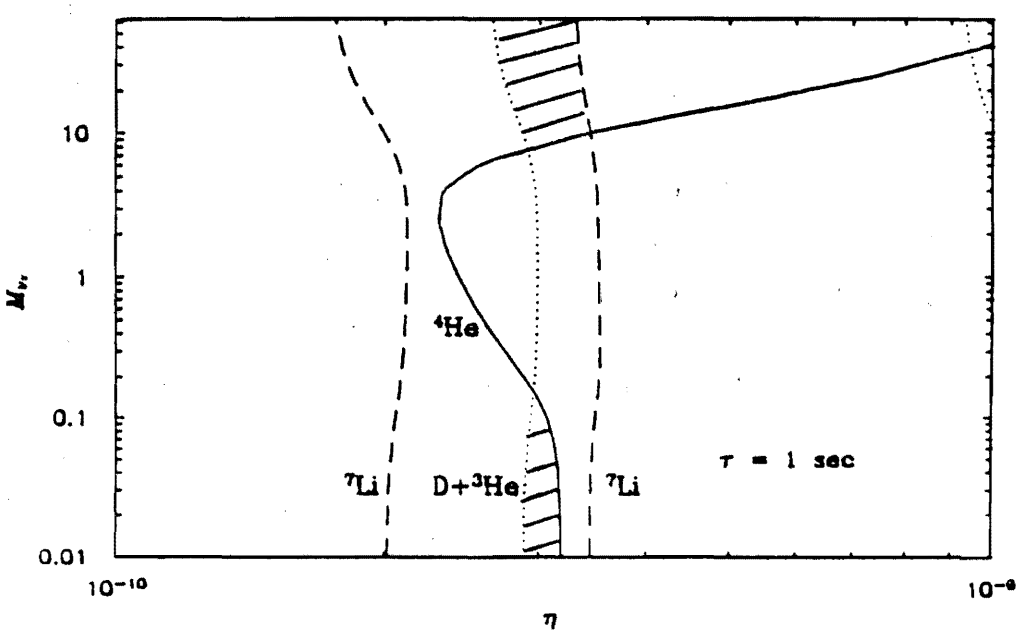

Fig. 3b. Same as 3a. for τ_{ν_τ} = 1 sec.

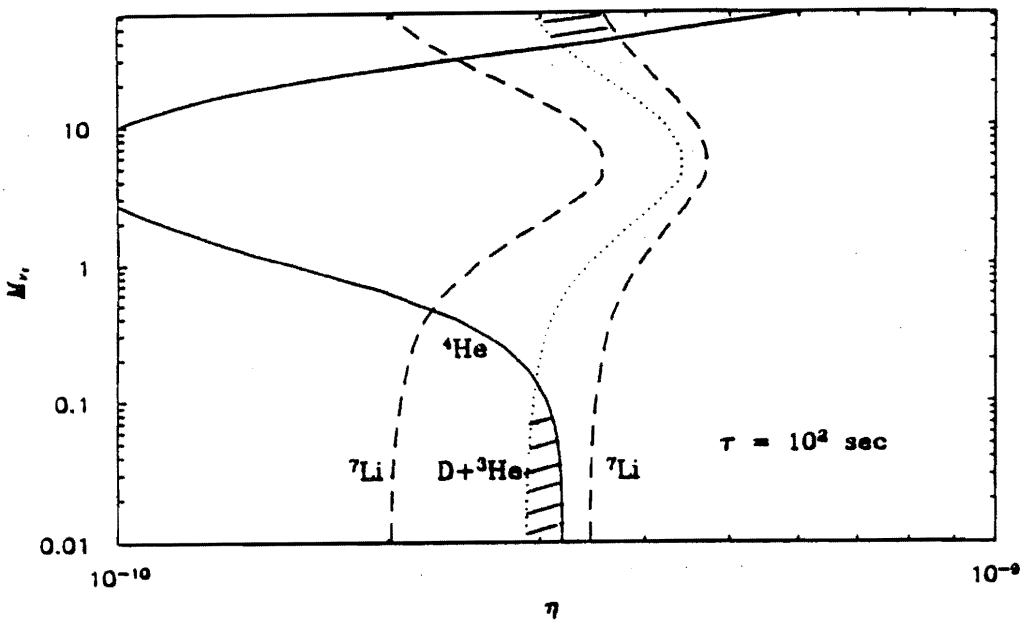

Fig. 3c. Same as 3a. for $\tau_{\nu_\tau} = 100$ sec.

5. Constraints On The ν_τ Mass And Lifetime

My colleagues and I have explored the allowed range of ν_τ masses and lifetimes consistent with the primordial abundances of the light elements[7] and our conclusions are summarized in Figure 4. To explore the sensitivity of our excluded region to the knowledge of the upper bound to the primordial abundance of ^{4}He , we have considered $Y_p \leq 0.235$ and ≤ 0.245 in addition to the "standard" case of $Y_p \leq 0.240$.

The region between the solid curves is excluded if $Y_p \leq 0.240$. As expected, sufficiently light $\nu_\tau(\lesssim 120keV)$ are permitted. However, very light, very short-lived ν_τ are excluded since, for $m_{\nu_\tau}(MeV) \lesssim 10\tau_{\nu_\tau}(\text{sec})$, $N_\nu \gtrsim 3.3$ due to the contributions from the ϕ. Stable ν_τ with $m_{\nu_\tau} \lesssim 50MeV$ contribute too much to N_ν at BBN and are excluded for $Y_p \leq 0.240$. As the ν_τ lifetime decreases, the bound on the ν_τ mass decreases as well. The ARGUS[9] bound of $m_{\nu_\tau} \leq 31$ MeV is reached at $\tau_{\nu_\tau} \approx 40$ sec. Thus, there is a "window" for $\tau_{\nu_\tau} \lesssim 40$ sec and $5 - 10 \lesssim m_{\nu_\tau}$ (MeV) ≤ 31. For a ν_τ in this window the BBN yield of ^{4}He is *less* than 0.240; we return to the consequences of this possibility shortly.

For the more restrictive bound $Y_p \leq 0.235$, the *entire* region below the dash-dot curve in Fig. 4 is excluded - including a *massless* ν_τ! The reason for this is that in the standard BBN (with $N_\nu = 3$) $Y_p > 0.235$ for $y_{23p} \leq 10^{-4}$. A ^{4}He mass fraction as small as 0.235 is inconsistent with standard BBN. Here, too, there is a high mass window since for such massive ν_τ, $N_\nu < 3$.

If, however, the bound on primordial ^{4}He is less restrictive ($Y_p \leq 0.245$), a larger region in the $m_{\nu_\tau} - \tau_{\nu_\tau}$ plane is allowed as shown by the dashed curves in Fig. 4. Nonetheless, if very massive, ν_τ must decay quickly.

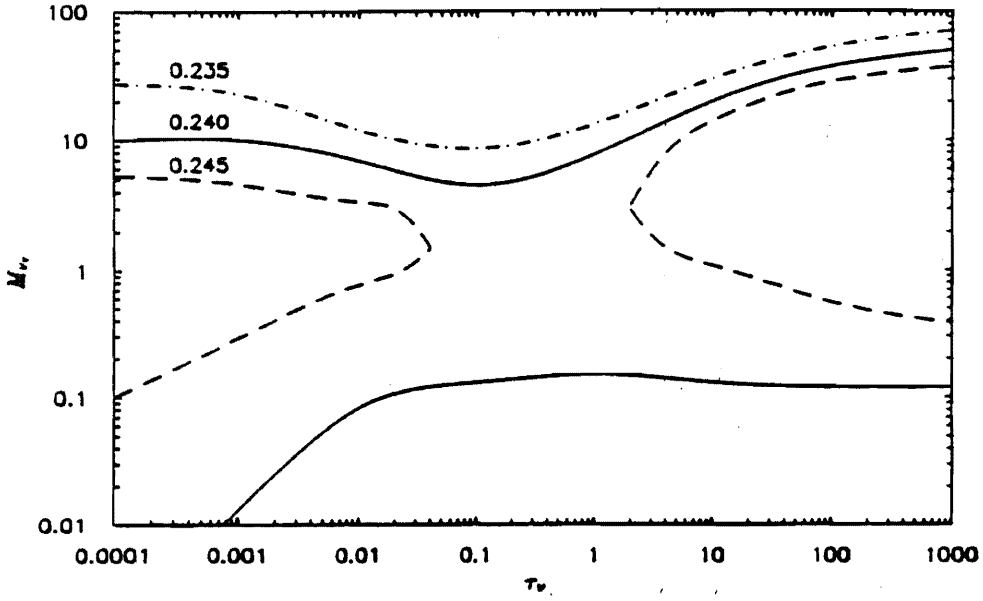

Fig. 4. The allowed/excluded regions in the $m_{\nu_\tau} - \tau_{\nu_\tau}$ plane. For $Y \leq 0.240$ the region between the solid lines is excluded; for $Y \leq 0.245$ the regions between the dashed lines is excluded; for $Y \leq 0.235$ the entire region below the dot-dashed curve is excluded.

Let us return to the $Y_p \leq 0.240$ bound and consider the high mass, short lifetime window. For ν_τ masses/lifetimes in this window $Y_p < 0.240$ is produced. Therefore, the BBN constraint on other, light weakly interacting particles ($\Delta N_\nu \lesssim 0.3$) is relaxed. To see this effect consider a ν_τ which saturates the ARGUS[9] bound ($m_{\nu_\tau} = 31 MeV$). Provided that the lifetime is less than 40 sec, in the absence of any new particles, $Y_p < 0.240$. If the bound $Y_p = 0.240$ is saturated, this would permit "new" particles to contribute to the total energy density at BBN. In Figure 5 we show this contribution - expressed as ΔN_ν - as a function of the ν_τ lifetime. For appropriately short lifetimes a new light scalar ($\Delta N_\nu = 4/7$) or, even a new light neutrino $\Delta N_\nu = 1$) is permitted. Alternately, if the observational data should lead to the firm conclusion that $Y_p < 0.240$, the high mass/short lifetime ν_τ window could rescue a seriously challenged hot big bang model. It is, therefore, crucial that this region of masses – accessible in accelerator experiments – be explored as carefully as possible. Either the window will be closed or, a massive ν_τ discovered!

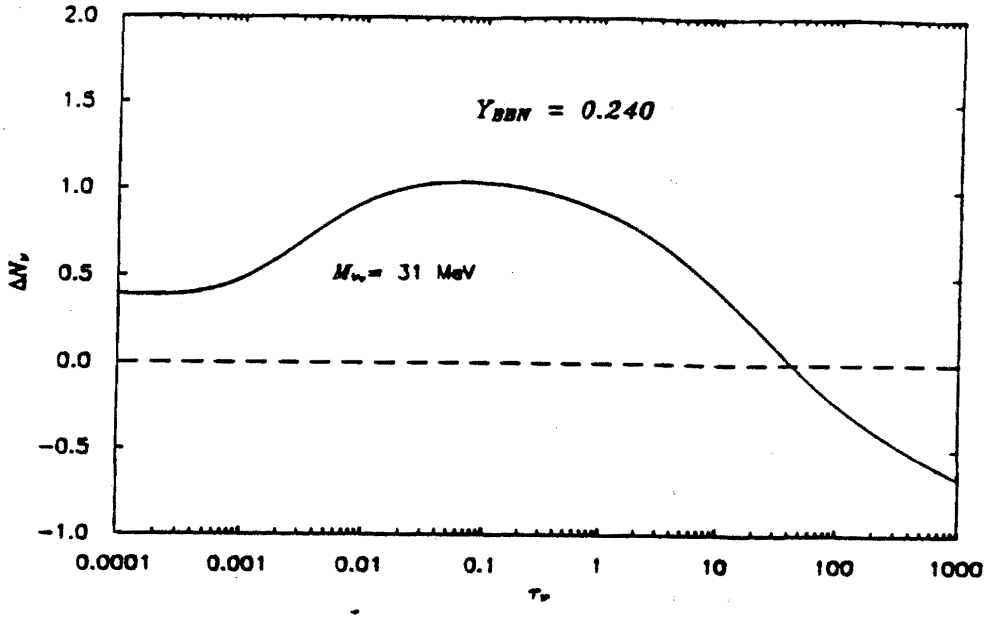

Fig. 5. The "extra" number of equivalent neutrinos, ΔN_ν, allowed for $m_{\nu_\tau} = 31$ MeV and $Y \leq 0.240$, as a function of τ_{ν_τ}.

To put our BBN constraints on the ν_τ mass and lifetime in a broader context, an expanded mass–lifetime plane is shown in Figure 6 (from ref. 7). Our[7] excluded region (for $Y_p \leq 0.240$) lies between the curves labeled (1) and the ARGUS[9] bound

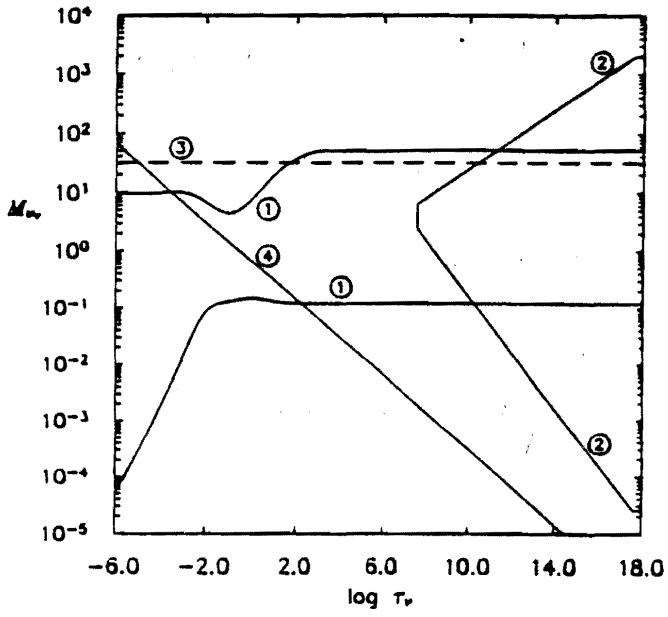

Fig. 6. Constraints on the ν_τ mass and lifetime.

excludes the region above (3). The region to the right of (2) is excluded by the constraint on the age/density of the present Universe ($t_0 \geq 10$Gyr, $\Omega_0 \leq 1$). Curve (4) is a guess at a possible mass–lifetime relation scaled from the bound to $\tau \to \mu + \phi$.

Finally, I end with a comment on alternate decay modes for a massive ν_τ. Generally, radiative decays ($\nu_H \to \nu_L e^+ e^-$, $\nu_H \to \nu_L \gamma$) are more strongly constrained than those into neutral scalars. For example SN1987A has constrained rapid radiative decay[10,11]. If B_γ is the branching ratio for radiative decay, the absence of gamma rays from SN1987A requires that τ_{ν_τ} (sec) $\gtrsim 8 \times 10^{11} B_\gamma m_{\nu_\tau}^{-1}$ (MeV). And, since the ν_τ carry away ~6 orders of magnitude more energy than is visible in the supernova, only a small fraction of ν_τ could decay into γ or $e^\pm$ pairs within the $\sim 3 \times 10^{12}$ cm radius of the supernova[11]. Thus, for $m_{\nu_\tau} \lesssim 10$ MeV, τ_{ν_τ} (sec) $\gtrsim 10^7 B_\gamma m_{\nu_\tau}$ (MeV). For a 10 - 31 MeV ν_τ, this bound is somewhat weakened by the suppression of thermal production of massive ν_τ; however, this suppression factor is not likely more than a factor of $\sim 10^2$ ensuring that radiative decays could contribute only little to BBN[7].

Since $\nu_\mu - \nu_\tau$ mixing is expected to be larger than $\nu_e - \nu_\tau$ mixing, it is expected that $\nu_\tau \to \nu_e \phi$ is less likely than our $\nu_\tau \to \nu_\mu \phi$ decay mode. However, decay produced ν_e could modify the evolving n/p ratio and effect the ^{4}He abundance. For low mass ($\lesssim$ 1 MeV) ν_τ we expect this effect to be small[12] and so the lower boundaries in the $m_{\nu_\tau} - \tau_{\nu_\tau}$ phase should be unmodified. For more massive ν_τ the situation is different – and more complex. Preliminary work[12] suggests that relatively long-lived, massive ν_τ may increase the ^{4}He abundance while shorter-lived, massive ν_τ decrease it. If verified by detailed numerical work, this would "tilt" the upper envelope in the $m_{\nu_\tau} - \tau_{\nu_\tau}$ plane to higher masses at longer lifetimes and slightly smaller masses at shorter lifetimes. This latter effect may widen the window at high masses and short lifetimes making even more urgent accelerator studies of the mass range above a few MeV.

6. Acknowledgements

The results described here have been obtained in close collaboration with my Ohio State colleagues, students and postdoc. I am pleased to acknowledge the valuable contributions of H.S. Kang, M. Kawasaki, P. Kernan, R.J. Scherrer and T. P. Walker and I thank them for permitting me to review our common work here. I also thank K. Enqvist, I. Rothstein, M. Turner and J. Valle for helpful discussions. This work is supported at Ohio State by the DOE (DE-A02-76ER01545).

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