Limits on the $\nu_\tau$ mass from nucleosynthesis in the presence of annihilation into majorons

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Abstract

We show that in the presence of sufficiently strong $\nu_\tau$ annihilations into majorons, the primordial nucleosynthesis constraints can not rule out $\nu_\tau$ masses in the MeV range.

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1 Introduction

Despite great experimental efforts the tau-neutrino still remains as the only one that can have a mass in the MeV range. The present experimental limit on its mass is

$$m_{\nu_\tau} < 23 \text{ MeV}$$ \hfill (1)

Further progress will have to wait for the tau-charm or B-factories. On the other hand, many particle physics models of massive neutrinos lead to a tau neutrino with mass in the MeV range. Moreover such a neutrino may have interesting cosmological implications. It is therefore important to examine in more detail the cosmological constraints.

The first comes from the critical density argument. From the contribution of stable $\nu_\tau$ to the present relic density one gets $m_{\nu_\tau} < 92 \Omega h^2 \text{ eV}$, where $h$ is related to the Hubble constant through $h = H_0/(100 \text{ km} s^{-1} \text{ Mpc}^{-1})$. This means that a massive $\nu_\tau$ with mass in the MeV range must be unstable with lifetimes smaller than the age of the Universe. However, it has been shown that in many particle physics models where neutrinos acquire their mass by the spontaneous breaking of a global lepton number symmetry this limit can be avoided due to the existence of fast $\nu_\tau$ decays.

The second constraint comes from primordial nucleosynthesis considerations. In the standard model, these rule out $\nu_\tau$ masses in the range

$$0.5 \text{ MeV} < m_{\nu_\tau} < 35 \text{ MeV}$$ \hfill (2)

It is possible to weaken the constraints on $\nu_\tau$ of Eq. (2), by adding new interactions beyond the standard ones. One may either consider that the $\nu_\tau$'s are unstable during nucleosynthesis or that they possess new channels of annihilation beyond the standard ones. It is this last possibility that we investigate in the case where the $\nu_\tau$'s can annihilate to majorons ($J$).

2 Evolution of $\nu_\tau$ number density

2.1 Before weak decoupling

We will assume that the massive Majorana $\nu_\tau$'s are stable during the nucleosynthesis epoch. They interact with leptons via the standard weak interactions $\nu_\tau \nu_\tau \leftrightarrow \nu_0 \nu_0, e^+e^-$. Moreover the $\nu_\tau$'s annihilate to majorons $\nu_\tau \nu_\tau \rightarrow JJ$ via the diagrams shown in Fig. 1. The t-channel diagram is present in all Majoron models, while the strength of the s-channel scalar exchange diagram is somewhat model-dependent. For the range of $g$ values relevant for our purposes ($g \gtrsim 10^{-5}$) it is quite reasonable to neglect the s-channel process, especially in Majoron models where the breakdown of the global lepton number symmetry happens at the electroweak scale or higher, like the simplest seesaw model. We have shown that for models with low
scale breaking of lepton number, such as given in ref. [8, 9] the bounds obtained by neglecting the s-channel process could be relaxed by at most a factor \( \sim 2 \) or so. Therefore in what follows we will consider that the \( \nu_\tau \)'s annihilate to majorons \( \nu_\tau \nu_\tau \rightarrow JJ \) via the diagonal coupling

\[
\mathcal{L}_{\nu_\tau\nu_\tau J} = i \frac{1}{2} g J \nu_\tau \gamma_5 \nu_\tau
\]

The evolution of the number densities of \( \nu_\tau \) (\( n_\tau \)) and majorons (\( n_J \)) is given by the solution of a set of Boltzmann differential equations

\[
\dot{n}_\tau + 3Hn_\tau = - \sum_{i=e,\nu_0} \langle \sigma_i v \rangle (n_\tau^2 - (n_\tau^{eq})^2) - \langle \sigma_J v \rangle (n_\tau^2 - (n_\tau^{eq})^2) \frac{n_J^2}{(n_J^{eq})^2} \equiv S_{\nu_\tau} \tag{4}
\]

\[
\dot{n}_J + 3Hn_J = \langle \sigma_J v \rangle (n_\tau^2 - (n_\tau^{eq})^2) \frac{n_J^2}{(n_J^{eq})^2} \equiv S_J \tag{5}
\]

where \( \langle \sigma_i v \rangle \) are the thermally averaged cross sections times the \( \nu_\tau \) relative velocity \( v \). The cross sections are for annihilation to majorons

\[
\sigma_J(\eta) = \frac{g^4}{128\pi} \frac{1 - \eta}{m_{\nu_\tau}^2 \eta} \left[ \ln \left( \frac{1 + \sqrt{\eta}}{1 - \sqrt{\eta}} \right) - 2\sqrt{\eta} \right] \tag{6}
\]

and for annihilation into massless fermions

\[
\sigma_i(\eta) = \frac{32G_F^2}{3\pi} \frac{m_{\nu_i}^2 \sqrt{\eta}}{1 - \eta} (b_{Li}^2 + b_{Ri}^2) \tag{7}
\]

where \( b_L^2 + b_R^2 = 1/2 \) for \( i = \nu_0 \) and \( b_L^2 + b_R^2 = 2 \left[ (-1/2 + \sin^2 \theta_W)^2 + (\sin^2 \theta_W)^2 \right] \approx 0.25 \) for \( i = e \). In Eq. (7) we assumed Boltzmann statistics and \( n_i = n_i^{eq} \) for \( i = e, \nu_0 \).

Now let us describe our calculations. First we normalized the number densities to the number density of massless neutrinos, \( n_0 \approx 0.181T^3 \), introducing the quantities
\[ r_\alpha \equiv n_\alpha / n_0, \text{ where } \alpha = \nu_\tau, J, \text{ and the corresponding equilibrium functions } r_\alpha^{eq}. \] Then from Eq. (5) we get
\[ \frac{dr_\alpha}{dT} = \left( \frac{S_\alpha}{n_0} - 3Hr_\alpha \right) \frac{1}{T} - \frac{3}{T} r_\alpha \] (8)
On the other hand, the time derivative of the temperature is obtained from Einstein’s equation
\[ \dot{\rho} = -3H(\rho + P) \] (9)
where \( \rho \) is the total energy density and \( P \) is the pressure.

2.2 Past weak decoupling

Once the \( \nu_\tau \)'s decouple from the standard weak interactions, they remain in contact only with the majorons. Then one has two different plasmas, one formed by \( \nu_\tau \)'s and \( J \)'s and the other by the rest of particles, each with its own temperature, denoted by \( T \) and \( T_\gamma \), respectively. We assume that the photon temperature evolves in the usual way, \( \dot{y} = H y \). The evolution equations of the \( \nu_\tau \) and \( J \) number densities are now simplified versions of Eq. (5), because we \( S_{\nu_\tau} = -S_J \),
\[ \dot{n}_\nu + 3Hn_\nu = -S_J \]
\[ \dot{n}_J + 3Hn_J = S_J \] (10)
To solve these equations one must find a relation between \( T \) and \( T_\gamma \). This is obtained using Einstein’s equation for the \( \nu_\tau + J \) plasma.

In order to determine the final \( \nu_\tau \) frozen density which will be relevant during nucleosynthesis we have to solve numerically the appropriate set of coupled differential equations for each pair of values \((m_{\nu_\tau}, g)\). Before weak decoupling these are Eqs. (5) e (9) while after decoupling one should use Eq. (10). The initial conditions are, for sufficiently high temperatures, \( r_{\nu_\tau} = r_{\nu_\tau}^{eq}, r_J = r_J^{eq} \) and \( T_{\nu_\tau} = T_J = T_{\nu_0} \).

3 Nucleosynthesis constraints

The value of \( r_{\nu_\tau}(m_{\nu_\tau}, g) \) is used to estimate the variation of the total energy density \( \rho_{tot} = \rho_R + \rho_{\nu_\tau} \). In \( \rho_R \) all relativistic species are taken in account, including the majorons and two massless neutrinos, whereas \( \rho_{\nu_\tau} \) is the energy density of the massive \( \nu_\tau \)'s.

In order to compare with the standard model situation, we can now express the effect of the \( \nu_\tau \) mass and of the annihilation to majorons in terms of an effective number of massless neutrino species, \( N_{eq} \), which we calculate for each value of \( r_{\nu_\tau}(m_{\nu_\tau}, g) \). To do this, we first numerically calculate the evolution of the neutron fraction, \( r_n \), by varying the value of \( N_{eq} \). Then we incorporate \( \rho_{tot} \) in this numerical code and perform the integration for each pair of \((m_{\nu_\tau}, g)\) values. Comparing
the \( r_n \) obtained in each case at \( T_\gamma \simeq 0.065 \text{ MeV} \) (the moment when practically all neutrons are wound up in \(^4\text{He}\)), we can relate \((m_{\nu_\tau}, g)\) and \(N_{eq}\).

The results are shown in Fig. (2). One can see that for fixed \(N_{eq}\), a wide range of tau neutrino masses is allowed for large enough coupling constants. In fact all masses below 23 MeV are allowed provided the coupling constant \(g\) exceeds a value a few times \(10^{-4}\). For comparison, the dashed line corresponds to the case when \(g = 0\) and no majorons are present. These results can also be expressed in the \(m_{\nu_\tau} - g\) plane, as shown in Fig. (3).

4 Majoron models

There has been a variety of majoron models proposed in the literature[11]. They are attractive extensions of the standard model where neutrinos acquire mass by virtue of the spontaneous violation of a global lepton number symmetry. Apart from a phenomenological interest of their own, majoron models offer the possibility of loosening the cosmological limits on the neutrino masses, either because neutrinos decay or annihilate to majorons.

We have seen above that the restrictions imposed by primordial nucleosynthesis upon a heavy tau neutrino disappear for sufficiently large values of the \(\nu_\tau \nu_\tau\) majoron coupling \(g\). Different models imply different expectations for the coupling \(g\) and for the relation between \(g\) and the \(\nu_\tau\) mass. Just to give a concrete example let us consider the supersymmetric models with spontaneous violation of R parity[9]. In these models we get the relation between \(g\) and \(m_{\nu_\tau}\) depicted in Fig. (4). The
Figure 3: The values of $g(m_{\nu_{\tau}})$ above each line would be allowed by nucleosynthesis if one adopts the $N_{eq}^{max} = 3, 3.4, 3.8, 4.2$ (from top to bottom).

different curves correspond to different values of the parameter $v_R$ of the model. We see that we can obtain for $g$ values of the order of a few times $10^{-4}$, as required in our nucleosynthesis analysis.

5 Conclusions

We have determined the restrictions imposed by primordial nucleosynthesis upon a heavy tau neutrino in the presence of sufficiently strong $\nu_{\tau}$ annihilations into majorons. We have shown that for values of the $\nu_{\tau} \nu_{\tau}$ majoron coupling in excess of $10^{-4}$ all the masses up to the experimental limit of 23 MeV can be allowed. We have shown that in a supersymmetric model with spontaneous violation of R parity such values can be obtained.

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Figure 4: Expected values of $m_{\nu_\tau}$ and $g$ in model of ref [9]

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