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# RECONCILING COLD DARK MATTER WITH COBE/IRAS PLUS SOLAR AND ATMOSPHERIC NEUTRINO DATA

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## Abstract

We present a model where an unstable MeV Majorana tau neutrino can naturally reconcile the cold dark matter model (CDM) with cosmological observations of large and small scale density fluctuations and, simultaneously, with data on solar and atmospheric neutrinos. The solar neutrino deficit is explained through long wavelength, so-called *just-so* oscillations involving conversions of  $\nu_e$  into both  $\nu_\mu$  and a sterile species  $\nu_S$ , while atmospheric neutrino data are explained through  $\nu_\mu$  to  $\nu_e$  conversions. Future long baseline neutrino oscillation experiments, as well as some reactor experiments will test this hypothesis. The model is based on the spontaneous violation of a global lepton number symmetry at the weak scale. This symmetry plays a key role in generating the cosmologically required decay of the  $\nu_\tau$  with lifetime  $\tau_{\nu_\tau} \sim 10^2 - 10^4$  seconds, as well as the masses and oscillations of the three light neutrinos  $\nu_e$ ,  $\nu_\mu$  and  $\nu_S$  required in order to account for solar and atmospheric neutrino data. It also leads to the invisibly decaying higgs signature that can be searched at LEP and future particle colliders.

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# 1 Introduction

A tau neutrino with a mass in the MeV range is an interesting possibility to consider for two different reasons. On experimental side such a neutrino is within the range of the detectability, for example at a tau-charm factory [1, 2]. On the other hand, if such neutrino decays before the matter dominance epoch, its decay products could then add energy to the radiation thereby delaying the time at which the matter and radiation contributions to the energy density of the universe become equal. Such delay would allow one to reduce the density fluctuations at the smaller scales [3] purely within the standard cold dark matter scenario [4], and could reconcile the large scale fluctuations observed by COBE [5] with the earlier observations such as those of IRAS [6] on the fluctuations at smaller scales. An MeV  $\nu_\tau$  may, however, conflict with the big-bang nucleosynthesis picture [7]. This conflict can be avoided in two ways.

If the tau neutrino has a strong coupling to a light spin zero Goldstone boson - called majoron and denoted  $J$  - with a typical strength  $g_J \gtrsim 10^{-4}$ , then the annihilation of  $\nu_\tau$  pairs to majorons could reduce their number density sufficiently so as to be consistent with the nucleosynthesis bound [8]. In spite of this reduction, subsequent  $\nu_\tau$  decays could increase the energy density of the radiation enough to reconcile COBE with IRAS data. This possibility has been recently investigated [8] in the context of a specific doublet majoron model [9], where an upper bound on the required  $\nu_\tau$  life of  $10^6$  seconds was set.

Alternatively, even if the majoron coupling to the tau neutrino is not so strong it may be possible to reconcile the nucleosynthesis constraints with the MeV  $\nu_\tau$  hypothesis if its decay involves the electron neutrino. The decay  $\nu_e$  could be captured by neutrons so as to reduce the resulting yield of primordial helium. In this case, a  $\nu_\tau$  with mass of a few MeV and lifetime in the range 10-50 seconds has been advocated in ref. [10].

An MeV tau neutrino, though cosmologically interesting, does not obviously fit with the data on solar and atmospheric neutrinos [11], for which neutrino oscillations with quite small mass differences provide the most plausible solutions [12, 13]. However, so far all attempts [8, 14] to obtain an MeV  $\nu_\tau$  with the lifetime required to revive the cold dark matter picture have ignored solar as well as atmospheric neutrino data. It is desirable to develop a coherent model which not only fits COBE and IRAS observations, but also provides solutions to the solar and atmospheric neutrino problems.

In this note we realize the MeV tau neutrino hypothesis in a model that can naturally reconcile the cosmological data on primordial density fluctuations with an explanation of the solar and atmospheric neutrino deficits through neutrino oscillations. The simplest way to

do this is to assume the few MeV  $\nu_\tau$  to be a majorana neutrino and not a Dirac neutrino as assumed in [10]. Indeed, an MeV  $\nu_\tau$  can not pair up with  $\nu_\mu$  or with  $\nu_e$  in order to form a Dirac state, because of the laboratory bounds on the masses of  $\nu_\mu$  or  $\nu_e$ . As a result, an MeV Dirac state would be obtained only by pairing off the two-component  $\nu_\tau$  with a sterile neutrino state of the same mass. In this case at least three other light neutrino species, one of which should also be sterile, would be required in order to fit together the data on solar and atmospheric neutrino oscillations. This follows from the fact that in this case the oscillations of  $\nu_\mu$  or  $\nu_e$  into the MeV  $\nu_\tau$  cannot solve the atmospheric or the solar neutrino problems and, on the other hand, the  $\nu_e - \nu_\mu$  oscillations could solve either but not both, since they require quite different values for the corresponding (mass)<sup>2</sup> differences [12].

Hence, the most economical way to reconcile COBE and IRAS observations with the cold dark matter picture and with solar and atmospheric neutrino data should involve the presence of just one very light sterile neutrino and just one two-component MeV state: the majorana tau neutrino.

To construct a consistent and appealing theoretical model is a non trivial task. Apart from naturally relating the required smallness of the sterile neutrino mass to a suitable symmetry, one has to obey a number of constraints:

(i) The mass and mixing of the very light sterile neutrino  $\nu_S$  should not conflict with the constraints coming from the nucleosynthesis [15] and supernovae [16].

(ii) The model should contain three different mass scales namely  $m_{\nu_\tau} \sim MeV$  to account for the tau neutrino mass,  $\Delta_S \sim 10^{-6} \text{ eV}^2$  or  $10^{-10} \text{ eV}^2$  in order to account for the solar neutrino deficit and  $\Delta_A \sim 10^{-2} \text{ eV}^2$  to solve the atmospheric neutrino problem.

(iii) The  $\nu_\tau$  should decay with lifetime in the range  $10^3 - 10^8$  seconds [8] or  $10 - 50$  seconds [10]. In the former case, it should couple strongly to majoron, while in the latter case its decay products should produce  $\nu_e$ .

(iv) The couplings of the majoron should be strong enough to satisfy (iii), but this coupling should not result in excessive energy loss through majoron emission inside a supernova [17].

We present below a model which successfully meets all these requirements and discuss its most salient features.

## 2 The model

Our model is based on the triplet plus singlet majoron scheme [18] and contains three right handed neutrinos, one of which is kept light by the imposed global symmetry. This way of keeping the sterile neutrino light has been already used in variety of models which tried to accommodate the possible existence of a 17 keV neutrino state [19] or which try to solve the solar, atmospheric and the dark matter problems simultaneously [12, 20] ‡.

We replace the lepton number symmetry of the original singlet majoron model by a generation dependent global symmetry  $U(1)_G$  under which the various fields transform in the manner shown in table 1. The generation dependent symmetry serves two purposes. It keeps the sterile  $\nu_e^c$  light and it leads to the decay of the tau neutrino into lighter neutrinos plus a majoron [22]. The  $SU(2) \times U(1) \times U(1)_G$  invariant couplings of the neutrinos are given by:

$$\begin{aligned} \mathcal{L}_Y = & \frac{1}{2} \frac{m}{\langle T^0 \rangle} \nu_{\tau L}^T C \nu_{\tau L} T^0 - \frac{\phi^0}{\langle \phi^0 \rangle} [m_1 \bar{\nu}_{eL} \nu_{\mu R} + m_2 \bar{\nu}_{\mu L} \nu_{\mu R} + m_3 \bar{\nu}_{\tau L} \nu_{\tau R}] \\ & + \frac{\mu}{\langle \sigma \rangle} \nu_{eR}^T C \nu_{\tau R} \sigma + \frac{1}{2} M [\nu_{\mu R}^T C \nu_{\tau R} + \nu_{\tau R}^T C \nu_{\mu R}] + H.c. \end{aligned} \quad (1)$$

The above Yukawa interactions lead to the following neutrino masses:

$$\mathcal{L}_m = \frac{1}{2} \nu_L^T C M_\nu \nu_L + H.C. \quad (2)$$

where  $M_\nu$  is a  $6 \times 6$  matrix having the following form in the left-handed basis  $\nu \equiv (\nu_e^c, \nu_e, \nu_\mu, \nu_\tau, \nu_\mu^c, \nu_\tau^c)^T$ :

$$M_\nu = \begin{pmatrix} 0 & 0 & 0 & 0 & 0 & \mu \\ 0 & 0 & 0 & 0 & m_1 & 0 \\ 0 & 0 & 0 & 0 & m_2 & 0 \\ 0 & 0 & 0 & m_{\nu_\tau} & 0 & m_3 \\ 0 & m_1 & m_2 & 0 & 0 & M \\ \mu & 0 & 0 & m_3 & M & 0 \end{pmatrix} \quad (3)$$

We assume the bare mass  $M$  to be much greater than other scales appearing in eq. (3), as in the seesaw model. These two heavy states can be diagonalized out leading to four light

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‡In principle, these problems can be solved without invoking a light sterile state if neutrinos are almost degenerate [20, 21]. However, in our present case one is obliged to introduce a light sterile neutrino if the mass of  $\nu_\tau$  is to lie in the MeV range.

states, one of which is massless, from the form of eq. (3). The effective  $4 \times 4$  neutrino mass matrix obtained in the seesaw approximation takes the following form:

$$m_{eff} = \begin{pmatrix} 0 & \beta \cos \theta & \beta \sin \theta & 0 \\ \beta \cos \theta & 0 & 0 & \alpha \cos \theta \\ \beta \sin \theta & 0 & 0 & \alpha \sin \theta \\ 0 & \alpha \cos \theta & \alpha \sin \theta & m_{\nu_\tau} \end{pmatrix} \quad (4)$$

where,

$$\beta^2 \equiv \left(\frac{\mu}{M}\right)^2(m_1^2 + m_2^2) \quad \alpha^2 \equiv \left(\frac{m_3}{M}\right)^2(m_1^2 + m_2^2); \quad (5)$$

and

$$\tan \theta \equiv \frac{m_2}{m_1} \quad (6)$$

Note that the  $\nu_{eR}$  is not allowed to receive a large mass and is kept light after the seesaw mechanism. As already mentioned, one of the four light states described by  $m_{eff}$  is in fact massless,  $\lambda_1 = 0$ . The other three are massive with eigenvalues  $\lambda_{2,3,4}$  approximately given by:

$$\lambda_2 \approx -\beta + \frac{\alpha^2}{2m_{\nu_\tau}} \quad (7)$$

$$\lambda_3 \approx \beta + \frac{\alpha^2}{2m_{\nu_\tau}}$$

$$\lambda_4 \approx m_{\nu_\tau} - \frac{\alpha^2}{m_{\nu_\tau}} \quad (8)$$

These eigenvalues nicely reproduce the hierarchical scales required for our purposes. Because of the chosen quantum numbers with respect to  $U(1)_G$ , the  $\nu_\tau$  is the only state to receive the mass from the  $SU(2)$  triplet fields and could be in the MeV range. The parameters  $\alpha$  and  $\beta$  could be much smaller than  $m_{\nu_\tau}$  if the seesaw mass scale  $M$  is chosen appropriately large. In this case, the matrix  $m_{eff}$  itself has a seesaw structure. The effective matrix describing the mixing of  $\nu_{e,\mu}$  and the sterile neutrino ( $\nu_s \equiv \nu_{eR}$ ) is easily seen to possess an approximate  $L_s - L_\mu - L_e$  symmetry. This leads to a pair of almost degenerate states  $\lambda_{2,3}$ . Their mass provides the (mass)<sup>2</sup> difference

$$\Delta_A \equiv \lambda_2^2 - \lambda_1^2 \approx \beta^2 \quad (9)$$

The approximate symmetry and hence degeneracy among two of the neutrinos is broken by the corrections  $\mathcal{O}\left(\frac{\alpha^2}{m_{\nu_\tau}}\right)$  arising out of the vacuum expectation value (VEV) of the  $\sigma$  field which breaks the global symmetry and generates the majoron. These corrections are

naturally small because of the second seesaw mechanism and they provide another (mass)<sup>2</sup> difference

$$\Delta_S \equiv \lambda_2^2 - \lambda_3^2 \approx \frac{2\beta\alpha^2}{m_{\nu_\tau}} \quad (10)$$

For the values  $\alpha \sim \beta \ll m_{\nu_\tau}$  one naturally obtains the hierarchy  $\Delta_S \ll \Delta_A$ .

The mixing among the four light states is specified by the matrix:

$$U m_{eff} U^T = \text{diag.}(0, m_{\nu_2}, m_{\nu_3}, m_{\nu_4}) \quad (11)$$

where  $m_{\nu_i} \equiv |\lambda_i|$  and  $U$  is defined by

$$U^T \sim \begin{pmatrix} 0 & \frac{1}{\sqrt{2}} & \frac{i}{\sqrt{2}} & \frac{\beta\alpha}{m_{\nu_\tau}^2} \\ -\sin\theta & \frac{1}{\sqrt{2}}\cos\theta & -\frac{i}{\sqrt{2}}\cos\theta & -\frac{\alpha}{m_{\nu_\tau}}\cos\theta \\ \cos\theta & \frac{1}{\sqrt{2}}\sin\theta & -\frac{i}{\sqrt{2}}\sin\theta & -\frac{\alpha}{m_{\nu_\tau}}\sin\theta \\ 0 & \frac{\alpha}{\sqrt{2}m_{\nu_\tau}} & -\frac{i\alpha}{\sqrt{2}m_{\nu_\tau}} & 1 \end{pmatrix} \quad (12)$$

The weak eigenstates are related to the mass eigenstates as follows:

$$\begin{aligned} \nu_s &\approx \frac{1}{\sqrt{2}}(\nu_{2L} + i\nu_{3L}) + \frac{\beta\alpha}{m_{\nu_\tau}^2}\nu_{4L} \\ \nu_{eL} &\approx \cos\theta\frac{1}{\sqrt{2}}(\nu_{2L} - i\nu_{3L}) - \frac{\alpha}{m_{\nu_\tau}}\cos\theta\nu_{4L} - \sin\theta\nu_{1L} \\ \nu_{\mu L} &\approx \sin\theta\frac{1}{\sqrt{2}}(\nu_{2L} - i\nu_{3L}) - \frac{\alpha}{m_{\nu_\tau}}\sin\theta\nu_{4L} + \cos\theta\nu_{1L} \\ \nu_{\tau L} &\approx \nu_{4L} + \frac{\alpha}{\sqrt{2}m_{\nu_\tau}}(\nu_{2L} - i\nu_{3L}) \end{aligned} \quad (14)$$

This gives rise to the following pattern for neutrino oscillations. Consider first the limit in which  $\Delta_S$  is neglected in comparison to  $\Delta_A$ . In this limit, there are no oscillations involving the sterile state  $\nu_s$ . On the other hand the  $\nu_e$  and  $\nu_\mu$  oscillate among themselves with the probability:

$$P_{e\mu} = \sin^2 2\theta \sin^2 \left( \frac{\Delta_A t}{4E} \right) \quad (15)$$

These oscillations could give rise to an explanation of the observed depletion in the atmospheric neutrino flux if the parameters lie in the range [11, 12]

$$\sin^2 2\theta = 0.35 - 0.8; \quad \Delta_A = (0.3 - 2) \times 10^{-2} \text{ eV}^2 \quad (16)$$

When the smaller (mass)<sup>2</sup> difference  $\Delta_S$  is turned on, the  $\nu_e$  and the  $\nu_\mu$  start oscillating into sterile state  $\nu_s$ . The probabilities for the  $\nu_e$  oscillations averaged over the shorter

atmospheric neutrino oscillation scale set by  $\Delta_A$  is given by

$$\begin{aligned}
P_{ee}(t) &= \frac{c^4}{2} \left( 1 + \cos\left(\frac{\Delta st}{2E}\right) \right) + s^4 \\
P_{e\mu}(t) &= \frac{c^2 s^2}{2} \left( 3 + \cos\left(\frac{\Delta st}{2E}\right) \right) \\
P_{es}(t) &= \frac{c^2}{2} \left( 1 - \cos\left(\frac{\Delta st}{2E}\right) \right)
\end{aligned} \tag{17}$$

Since the mixing angle involving  $\nu_2$  and  $\nu_3$  is  $45^\circ$ , the nonadiabatic MSW solution [23] to the solar neutrino problem is not feasible in the present case. However, the relevant mass scale  $\Delta_S$  could be naturally very small so that the solar neutrino flux may get depleted through the long wavelength vacuum oscillations [24]. For example,  $m_{\nu_\tau} = 5$  MeV,  $\beta = 0.1$  eV and  $\alpha = 0.05$  eV would give  $\Delta_S = 10^{-10} eV^2$  and  $\Delta_A = 10^{-2} eV^2$ . One sees that similar values of  $\alpha$  and  $\beta$  naturally result in a very large hierarchy between the solar and atmospheric scales. On the other hand, if one chooses the Dirac masses  $m_{1,2,3}$  in the GeV range, then the required values of  $\alpha, \beta$  can be obtained by choosing the bare mass  $M$  for the right handed neutrino around the intermediate scale  $10^{11}$  GeV. One concludes that the required values of  $\Delta_A$  and  $\Delta_S$  do follow for natural choices of the parameters.

Note that, although it has been shown that  $\nu_e$  to  $\nu_S$  oscillations where  $\nu_S$  is a sterile state are disfavored by the combined Homestake and Kamiokande data [25], in the present case the situation is more complex, since not only the sterile state  $\nu_S$  but also  $\nu_\mu$  are involved. Since both relevant mixing angles are large <sup>§</sup>, a phenomenologically consistent solution should exist. In order to determine its parameters more sharply a more detailed analysis of the existing solar neutrino data for the present case where both  $\nu_\mu$  and  $\nu_S$  take part in the solar neutrino oscillations would certainly be desirable.

### 3 Cosmology

Let us now turn to the cosmological aspects of the model. A stringent constraint on this scenario comes from primordial big bang nucleosynthesis, which requires that the effective number of degrees of freedom  $g_{eff}(T)$  contributing to the energy density of the universe at the nucleosynthesis time ( $\sim 1$  second) be less than 11.3 [7]. Here  $g_{eff}(T)$  has been defined in terms of the total energy density as  $\rho \equiv \frac{\pi^2}{30} g_{eff}(T) T^4$  which includes the contribution

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<sup>§</sup>A possible restriction on large angle neutrino mixing has been argued to follow from the observed energy spectra of  $\bar{\nu}_e$  from supernova SN1987A [26]. However, ref. [26] considered only the simplest case of two flavour  $\nu_e$  to  $\nu_\mu$  mixing, whereas in the present case one has three neutrino species one of which is sterile, so that those arguments do not directly apply.

of the relativistic species as well as that of the nonrelativistic  $\nu_\tau$ . The latter could violate this bound if it had only the conventional weak interactions. However, the presence of the majoron in our model alters the situation. The majoron in the model is easily seen to be

$$J \approx \sigma_I - \frac{2\omega}{u}T_I + \frac{4\omega^2}{uv}\phi_I \quad (18)$$

where  $u, v$ , and  $\omega$  denote ( $\sqrt{2}$  times) the vacuum expectation values of the  $\sigma, \phi^0$  and  $T^0$  scalar multiplets, respectively, while the suffix  $I$  denotes the corresponding imaginary parts. Note that because of the hierarchy  $\omega \sim MeV \ll u \sim 100GeV$ , the invisible decay of  $Z$  to the majoron is enormously suppressed, in accordance with LEP data [27], unlike in the purely triplet majoron scheme. The  $\nu_\tau$  couples to the majoron dominantly through its triplet admixture. Using eq. (1) and eq. (18) this coupling may be given as follows

$$\mathcal{L}_J \approx -i J \left\{ \frac{m_{\nu_\tau}}{u} \nu_{4L}^T C \nu_{4L} + \frac{\alpha}{\sqrt{2}u} \left[ \nu_{4L}^T C (\nu_{2L} - i\nu_{3L}) + (\nu_{2L} - i\nu_{3L})^T C \nu_{4L} \right] \right\} + H.c. \quad (19)$$

The  $\nu_\tau \sim \nu_4$  has a diagonal coupling to the majoron given by  $g_J \equiv \frac{m_{\nu_\tau}}{u}$ . For sufficiently large  $g_J$  the contribution of  $\nu_\tau$  at the time of nucleosynthesis can be suppressed through  $\nu_\tau$  annihilation. In order for this to happen, one requires [8]

$$T_{EQ1} \equiv \frac{4}{3} m_{\nu_\tau} Y < 0.13 \text{ MeV} \quad (20)$$

which corresponds to  $g_{eff}(T \sim 1MeV) < 11.3$ . Here  $Y$  determines the abundance of the nonrelativistic  $\nu_\tau$ . Using the standard expression [28] for  $Y$  we obtain

$$T_{EQ1} \approx (2.66 \times 10^{-2} MeV) \left( \frac{10 \text{ MeV}}{m_\nu} \right)^2 \left( \frac{u}{100 \text{ GeV}} \right)^4 \left( \frac{x_f}{4} \right)^2 \quad (21)$$

where

$$x_f \equiv \ln z - (n + 1/2) \ln \ln z \quad (22)$$

and

$$z = 0.038(n + 1)g/g_*^{1/2}M_{Pl}m_{\nu_\tau}\sigma_0 \quad (23)$$

where  $n = 1$ ,  $M_{Pl} = 1.2 \times 10^{19} \text{ GeV}$  is the Planck mass,  $g = 2$  and, in our case [8],  $g_* = 10$ . Finally,  $\sigma_0 = \frac{m_{\nu_\tau}^2}{2\pi u^4}$  denotes the  $\nu_\tau$  annihilation cross section.

Physically,  $T_{EQ1}$  represents the temperature at which the contribution of the  $\nu_\tau$  starts dominating over that of radiation. As would be expected, this should happen after the nucleosynthesis epoch. After  $T_{EQ1}$ , the  $\nu_\tau$  would dominate the energy density of the universe until it decays. This decay would make the universe radiation dominated and this could serve to delay the time at which matter starts dominating again. If one fixes the overall scale of the power spectrum of the density fluctuations in the cold dark matter scenario from the large



scale COBE observations, then one can also fit the small scale density fluctuations provided one chooses  $g_{eff}(T_{EQ2}) \approx \frac{25}{9} \times 3.36$  [8, 10]. Here  $g_{eff}(T_{EQ2})$  determines the total contribution in radiation at the temperature  $T_{EQ2}$  when the radiation and matter densities again become equal. The  $\nu_e, \nu_\mu$ , majoron and photons contribute the amount 3.17 to  $g_{eff}(T_{EQ2})$ . The remaining must come from the decay products of the  $\nu_\tau$ . This can be translated [8] into a constraint on the  $\nu_\tau$  lifetime:

$$\tau_{\nu_\tau} \approx 7.2 \times 10^{19} \left( \frac{T^0}{T_{EQ1}} \right)^2 \text{ sec} \quad (24)$$

where  $T^0 = 2.73$  is the present temperature of the universe. Since  $T_{EQ1} < 0.13$  MeV from nucleosynthesis,  $\tau_{\nu_\tau} > 2 \times 10^2$  sec. On the other hand,  $T_{EQ1}$  should be higher than the temp  $T_{EQ2} \sim \text{few eV}$  of the conventional matter radiation equality, giving an upper bound  $\tau_{\nu_\tau} < 10^{11}$  sec. In fact, requiring that the  $\nu_\tau$  contribution to the present total energy density of the universe is not too large gives a more stringent upper bound on  $\tau_{\nu_\tau}$  as a function of the mass  $m_{\nu_\tau}$  [28, 29]. For example, for  $m_{\nu_\tau} \sim 1$  MeV,  $\tau_{\nu_\tau} \lesssim 10^8$  sec.

The above considerations show that a  $\tau_{\nu_\tau}$  in the range  $\sim 10^2 - 10^8$  seconds would be able to delay matter radiation equality without conflicting with the nucleosynthesis picture.

We now show how the lifetime required by cosmology can naturally occur in the present model. For this note that eq. (19) gives rise to the decay of  $\nu_\tau$  to the lighter states  $\nu_{2,3}$  plus a majoron, with decay width given by

$$\Gamma_{\nu_\tau} = \frac{m_{\nu_\tau}}{4\pi} \left( \frac{\alpha}{u} \right)^2 \quad (25)$$

Using the expressions for  $\Delta_A$  and  $\Delta_S$  this can be written as

$$\tau_{\nu_\tau} = (1.6 \times 10^3 \text{ sec.}) \left( \frac{m_{\nu_\tau}}{10 \text{ MeV}} \right)^{-2} \left( \frac{u}{100 \text{ GeV}} \right)^2 \left( \frac{\Delta_A}{10^{-2} \text{ eV}^2} \right)^{1/2} \left( \frac{10^{-10} \text{ eV}^2}{\Delta_S} \right) \quad (26)$$

It follows from the above equation that one can simultaneously accommodate the solar and atmospheric neutrino deficits and also have a tau neutrino decay with the required lifetime if the singlet VEV  $u$  is chosen around 50 – 200 GeV. In Fig.1 we display the region of  $\nu_\tau$  mass and the singlet VEV  $u$  allowed by the various constraints. The solid curves a and b correspond to  $g_{eff} = 10.86$  and 11.3, respectively. The regions to the right of these curves would be disallowed by nucleosynthesis. For illustration, we give curves c and d corresponding to  $\nu_\tau$  lifetimes  $10^3$  and  $10^4$  seconds, respectively.

Another constraint on the model comes from the supernova. The MeV tau neutrino with relatively large couplings to the majoron may cause supernova to rapidly lose energy through majoron emission. The relevant constraints have been looked at in detail in ref. [17]

in the case of the simplest singlet majoron model. In our case the relevant coupling is the first in eq. (19). Since, this coupling is similar (apart from a factor 2) to that of the singlet majoron model we can adopt the analysis of ref. [17]. In our case, the following values of  $u$  and  $m_{\nu_\tau}$  are seen to be disallowed

$$5.7 \times 10^{-3} < \left( \frac{m_{\nu_\tau}}{10 \text{ MeV}} \right) \left( \frac{100 \text{ GeV}}{u} \right)^2 < 0.82 \quad (27)$$

It follows from the above two equations that it is indeed possible to obtain the  $\nu_\tau$  lifetime in the required range and solve the solar and atmospheric neutrino problem for parameter choices lying outside the range forbidden by the supernova.

The time structure of the SN87A antineutrino pulse may also be used to constrain the lifetime of  $\nu_\tau$ . The decay of a massive  $\nu_\tau$  emitted in the supernova on the way could lead to a delayed signal in the detectors. The absence of such a signal has been used in [30] and subsequently in [31] to put an upper bound on the  $\nu_\tau$  lifetime. A rough estimate of the resulting lifetime has been given in ref. [14] to be  $\sim 300$  seconds for an MeV  $\nu_\tau$  decaying by majoron emission. While this is consistent with the bound we have obtained of  $2 \times 10^2$  seconds, there is considerable room to relax this supernova bound. Firstly, only a fraction  $\cos^2 \theta$  of  $\nu_\tau$  decay to  $\nu_e$  in the present case. Moreover, original  $\nu_\tau$  flux may also be suppressed by the Boltzman factor for larger masses  $m_{\nu_\tau} > 10 \text{ MeV}$ . These could result in considerable weakening of the upper bound. Finally, in the derivation of the supernova bound one would have to be careful in differentiating between  $\nu_\tau$  and  $\bar{\nu}_\tau$  decays to  $\bar{\nu}_e$ .

The presence of a light sterile neutrino mixing with ordinary neutrinos is in general constrained by nucleosynthesis [15]. Neutrino oscillations in the early universe would bring the sterile neutrinos into equilibrium at the time of nucleosynthesis with the active ones. The relevant oscillation scale in the present case is  $\Delta_S \sim 10^{-10} \text{ eV}^2$ . The corresponding wavelength is too large to equilibrate the sterile neutrinos in the early universe. In fact the constraints on the relevant mass scale is trivially satisfied in our case.

## 4 Discussion

We have attempted in this paper to provide a coherent explanation of quite distinct phenomena in neutrino physics. The example given here is able to resurrect the cold dark matter picture of structure formation by making it consistent with both COBE and IRAS observations of primordial density fluctuations through a  $\nu_\tau$  of few MeV mass and lifetime in the range  $10^2 - 10^4$  sec. It also explains the solar and atmospheric neutrino problem in terms of mixing among three light neutrinos  $\nu_e$ ,  $\nu_\mu$  and  $\nu_S$ . The scheme presented here

is not unique but is certainly most economical from the point of view of explaining various phenomena mentioned above. More importantly, it contains predictions which can be tested in the near future in long baseline neutrino oscillation experiments, as well as some reactor neutrino experiments. Moreover, due to the presence of the weak scale majoron, the present model allows for a distinctive signature of the invisibly decaying Higgs boson  $h \rightarrow JJ$  [32], which could substantially affect higgs boson search strategies at LEP [33], NLC [34], as well as LHC [35].

Finally, notice that, since  $\nu_\tau$  is much heavier than other neutrinos and has very small mixing with them ( $\mathcal{O}(\alpha/m_{\nu_\tau} \sim 10^{-7})$ ) there are no oscillations involving the  $\nu_\tau$  and experiments such as CHORUS and NOMAD [36] looking for the  $\nu_\mu - \nu_\tau$  oscillations should show negative result. This feature is a generic expectation in any model having  $\nu_\tau$  in the MeV range. It also implies a strong suppression in the neutrinoless double beta decay rate.

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	$SU(2)$	$Y$	$G$
$\ell_{Le,\mu}$	2	-1	-1
$\ell_{L\tau}$	2	-1	1
$e_R, \mu_R$	1	-2	-1
$\tau_R$	1	-2	1
$\nu_{eR}$	1	0	-2
$\nu_{\mu R}$	1	0	-1
$\nu_{\tau R}$	1	0	1
$\phi$	2	-1	0
$T$	3	2	-2
$\sigma$	0	0	1

Table 1:  $SU(2) \otimes U(1)_Y \otimes U(1)_G$  assignments of the leptons and Higgs scalars. Quarks are  $U(1)_G$  singlets.

### Figure Caption

Fig.1 shows the region of  $\nu_\tau$  mass and lifetimes as a function of the singlet VEV  $u$  which are allowed by the various constraints. The solid curves a and b illustrate the nucleosynthesis constraints corresponding to  $g_{eff} = 10.86$  and  $11.3$ , respectively. The regions to the right of these curves would be forbidden. Curves c and d correspond to  $\nu_\tau$  lifetimes  $10^3$  and  $10^4$  seconds, respectively.

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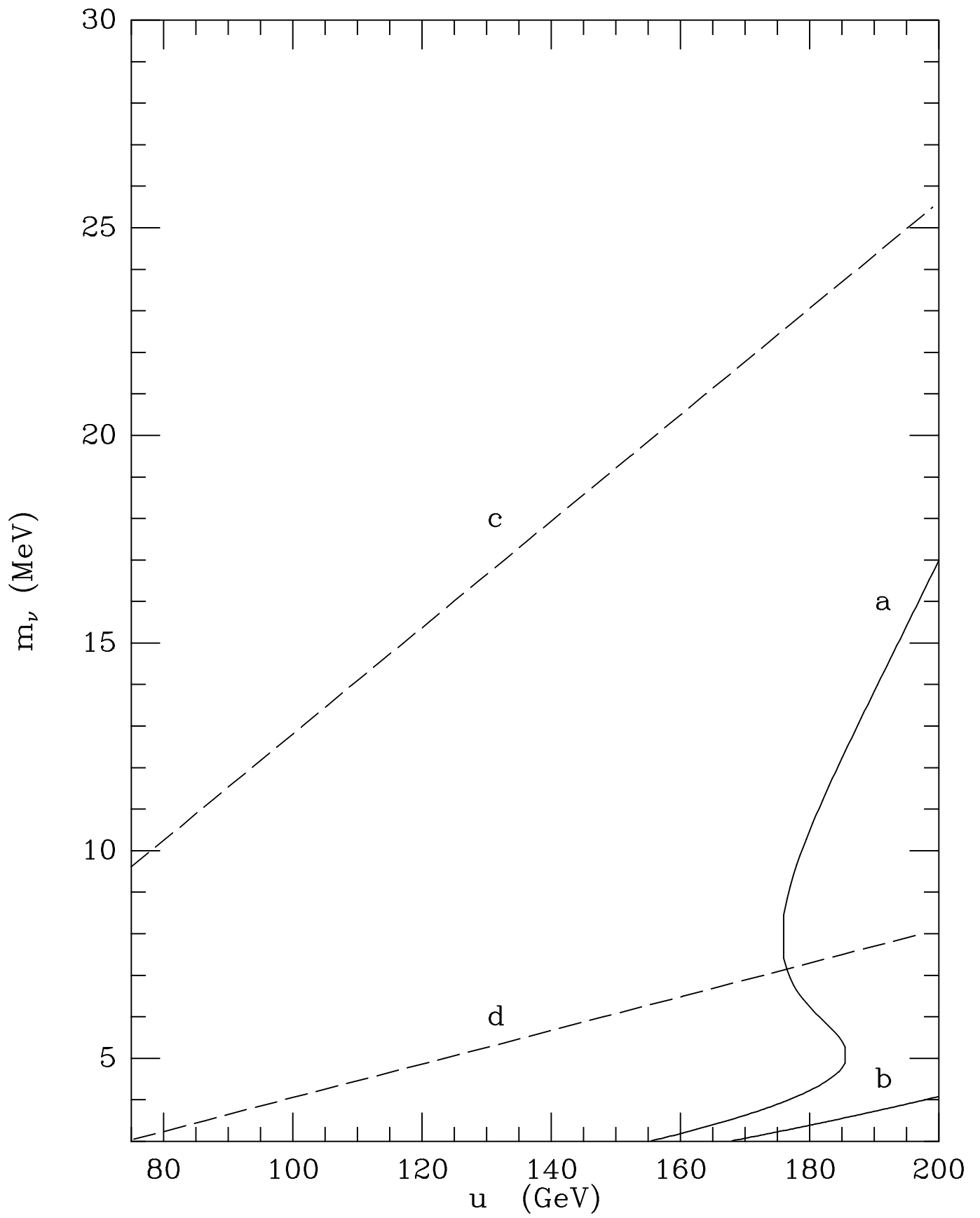


Fig. 1



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