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COSMOLOGICAL AND EXPERIMENTAL CONSTRAINTS ON THE TAU NEUTRINO

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A B S T R A C T

Recent experimental results, together with a cosmological argument concerning the survival of primordially synthesized deuterium, exclude an unstable tau neutrino heavier than ~ 1 MeV. The cosmological bound on stable neutrinos then restricts it further to be lighter than ~ 100 eV, in the absence of any available fast decay mode. These constraints are also relevant to other heavy neutrinos and unstable supersymmetric particles.

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1. INTRODUCTION

No stable neutrino can have a mass between ~ 100 eV and ~ 2 GeV without contributing to the energy density of the universe in excess of observational limits [1,2]. Thus the tau neutrino which is experimentally known to be lighter than 143 MeV [3], must be unstable if it is heavier than ~ 100 eV.

The standard $SU(2)_L \times U(1)$ model (as well as $SU(5)$) has only massless, left-handed neutrinos, but its successful phenomenology would be unaffected by the addition of right-handed neutrinos as isosinglets and/or additional Higgs bosons to violate lepton number conservation, thus allowing a Dirac or Majorana neutrino mass. Massive Majorana neutrinos arise more naturally in theories with intrinsic left-right symmetry and lepton-hadron symmetry, such as $SO(10)$ and its subgroup $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, while extended $SO(10)$ models may admit massive Dirac neutrinos [4]. However, if neutrinos do have masses, the flavour eigenstates $\nu_e, \nu_\mu, \nu_\tau \dots$ need not coincide with the mass eigenstates $\nu_1, \nu_2, \nu_3 \dots$, but would, in general, be related through a unitary transformation. Such a generalized Cabbibo-like mixing,

$$\nu_\ell = \sum_i U_{\ell i} \nu_i, \quad \ell \equiv e, \mu, \tau \dots, \quad i \equiv 1, 2, 3 \dots, \quad (1)$$

could result in neutrino oscillations when mass differences are small, and in decays [5,6] of heavy neutrinos into lighter ones.

The consequences of such decays in astrophysical sites of neutrino production and in the early universe have been widely investigated [7], particularly with regard to the radiative decay mode $\nu_i \rightarrow \nu_j, \gamma$. Cowsik [8] has argued that astrophysical and cosmological constraints combined with limits inferred indirectly from laboratory experiments require " ν_τ " (i.e. ν_3) to be stable towards radiative decays. In fact, this decay mode is severely suppressed by a leptonic Glashow-Iliopoulos-Maiani (GIM) mechanism and should, in any case, have a very long lifetime [5]

$$\tau(\nu_3 \rightarrow \nu_1 \gamma) > 2.4 \times 10^{14} \text{ sec} [m_{\nu_3}/\text{MeV}]^{-5}. \quad (2)$$

Nieves [5] has shown that the higher-order process $\nu_i \rightarrow \nu_j, \gamma\gamma$ is not GIM-suppressed and may therefore dominate over $\nu_i \rightarrow \nu_j, \gamma$. For $m_{\nu_3} \lesssim m_e$,

$$\begin{aligned}\tau(\nu_3 \rightarrow \nu_1 \gamma \gamma) &\approx (552960 \pi^5 / \alpha^2 G_F^2) (m_e^4 / m_{\nu_3}^9 |U_{e1}^* U_{e3}|^2) \\ &\gtrsim 1.1 \times 10^{12} \text{ sec} [m_{\nu_3} / \text{MeV}]^{-9} |U_{e3}|^{-2}.\end{aligned}\quad (3)$$

However, we find that when the neutrino mass significantly exceeds that of the charged lepton mediating the decay, this process too becomes rather slow, viz. for $m_{\nu_3} \gg m_e$,

$$\begin{aligned}\tau(\nu_3 \rightarrow \nu_1 \gamma \gamma) &\approx 384 \pi^5 / \alpha^2 G_F^2 m_{\nu_3}^5 |U_{e1}^* U_{e3}|^2 \\ &\gtrsim 1.1 \times 10^{10} \text{ sec} [m_{\nu_3} / \text{MeV}]^{-5} |U_{e3}|^{-2}.\end{aligned}\quad (4)$$

However, for $m_{\nu_3} \gtrsim 2 m_e$, the fast tree-level decay $\nu_3 \rightarrow e^- e^+ \nu_e$ is possible [6], with a lifetime which is scaled from muon decay,

$$\tau(\nu_3 \rightarrow e^- e^+ \nu_e) \approx 2.9 \times 10^4 \text{ sec} [m_{\nu_3} / \text{MeV}]^{-5} |U_{e3}|^{-2} [f(m_e / m_{\nu_3})]^{-1}, \quad (5)$$

where $f(x)$ is a phase-space factor ($\lesssim 1$ for $x \ll 1$). Kolb & Goldman [6] have noted that τ_{ν_3} , and hence the mixing amplitude $|U_{e3}|$, are constrained by two general cosmological arguments: the neutrino decay products should not have (a) excessively distorted the black-body radiation background [9], or (b) generated excessive entropy subsequent to primordial nucleosynthesis [10]. While radiative decays can be restricted further by requiring that primordial helium (^4He) and deuterium (D) should not have been excessively depleted by photofission [8], Kolb & Scherrer [10] have argued that this does not constrain the dominant decay mode $\nu_3 \rightarrow e^- e^+ \nu_e$. They suggest that the rapid thermalization of decay $e^\pm$ by Compton scattering in the early universe, together with the smaller cross section for $e^\pm$ induced fission, result in negligible destruction of primordial D.

However if we demand that ν_3 decays early enough to satisfy constraints (a) and (b) above, then the background photons are themselves sufficiently energetic to be Compton-scattered by the decay $e^\pm$ to energies above the thresholds for ^4He and D photofission. We find that the resultant depletion of primordial D is unacceptably large, thus requiring ν_3 to decay before nucleosynthesis begins. Such short lifetimes are however ruled out by recent laboratory experiments [11-13], forcing the conclusion that $m_{\nu_3} \lesssim 1 \text{ MeV}$. We elaborate on this argument below and then discuss further implications for the mass and lifetime of ν_3 . We also

comment on the relevance of these cosmological constraints to other heavy neutrinos [14] and unstable supersymmetric particles [15].

2. COSMOLOGICAL BOUNDS ON AN UNSTABLE NEUTRINO

The sequence of events involving a massive neutrino in the early universe is shown in fig. 1. Solution of the Boltzmann equation for a neutrino interacting according to the standard electroweak model yields the temperature T_D and number density n_D at "decoupling" [7], as a function of the mass $m_\nu^{(*)}$. Subsequently, the energy density of non-relativistic neutrinos decreases more slowly ($\rho_\nu \propto m_\nu T_\nu^3$) than that of relativistic particles ($\rho_r = N_F a T_Y^4$) and the universe becomes matter-dominated [9] at the temperature^(**),

$$T_m = (8/11) m_\nu n_D / N_F a T_D^3, \quad (6)$$

where N_F is the effective number of relativistic degrees of freedom (≈ 1.45 for $kT_Y < m_e$, including the contribution of ν_e and ν_μ) and photon heating by annihilation of thermal $e^\pm$ at $kT_Y \sim m_e$ makes $T_\nu = (4/11)^{1/3} T_Y$.

The massive neutrino must subsequently decay in order not to contribute excessively to the observed energy density today. Following Dicus et al. [9,10], we assume that all decays occur at time t_d ($\sim \tau_\nu$), making the universe radiation dominated again. A fraction $f = 2/3$ of ρ_ν goes into relativistic electrons and positrons which annihilate and raise the radiation temperature from T_{bd} before decay to T_{ad} after decay, according to $a T_{ad}^4 = a T_{bd}^4 + f \cdot 2 m_\nu n_D [(4/11)^{1/3} T_{bd}/T_D]^3$, so that,

$$T_{ad} = T_{bd}^{3/4} [T_{bd} + f N_f T_m]^{1/4}. \quad (7)$$

(*) We derive n_D from Kolb & Scherrer [10] for $m_\nu \leq 25$ MeV, and from Dicus et al. [10] for $m_\nu > 25$ MeV after correcting for out-of-equilibrium annihilations and additional interaction channels. We thus obtain agreement with Lee & Weinberg [2] for $m_\nu > m_H$. This procedure is necessary since the number of available annihilation channels is a discrete function of m_ν . Note that the widely used [33,37] calculations by Dicus et al. [2,9] erroneously assume fully populated right-handed helicity states, even though a strict V-A interaction is assumed. In fact, at the low energies of interest here, a Dirac neutrino is effectively a 2-component field, since the right-handed states have no gauge couplings.

(**) We use c.g.s. units, except that particle masses are in energy units. (The black-body radiation constant $a \approx 4.722 \times 10^{-9}$ MeV cm⁻³ K⁻⁴).

We calculate that time evolves with temperature as,

$$t = -[3/8\pi G_N(\rho_\nu + \rho_r)]^{1/2} \int dT_Y/T_Y$$

$$\approx 3.07 \times 10^{20} (\text{sec}) N_F^{-1/2} T_m^{-2} [(T_m/T - 2)(T_m/T + 1)^{1/2} + 2] . \quad (8)$$

The following arguments establish lower limits on T_{bd} and hence upper limits on t_d or, equivalently, the massive neutrino lifetime τ_ν .

2.1 Primordial Nucleosynthesis

The observed ratio of gravitating mass to luminous mass in galaxies [16] and in the universe as a whole [17] requires that baryons contribute a fraction $\Omega_B \equiv \rho_B/\rho_c > 0.01$ to the critical density $\rho_c \equiv 3H_0^2/8\pi G_N \approx 1.05 \times 10^{-2} h_0^2 \text{ MeV cm}^{-3}$. This sets a lower limit to the baryon-to-photon number ratio today, $\eta_0 \approx 2.81 \times 10^{-8} \Omega_B h_0^2 \theta_{2.7}^{-3} > 5.7 \times 10^{-11} (*)$. On the other hand, an upper limit to this ratio during primordial nucleosynthesis, η_{ns} , can be obtained by requiring the synthesized mass fraction of ^4He (which is a slowly increasing function of η_{ns}) to be below the observationally inferred limit [20], $Y_p(^4\text{He}) < 0.25$. Kolb & Scherrer [10] have included the energy density contributed by a massive neutrino in Wagoner's [21] computer code of primordial nucleosynthesis, to obtain such limits to η_{ns} as a function of m_ν . Thus any subsequent generation of entropy, such as from neutrino decay, is constrained by demanding that

$$T_{ad}^3/T_{bd}^3 = \eta_{bd}/\eta_{ad} \leq \eta_{ns}/\eta_0 . \quad (9)$$

Therefore, to achieve concordance between the calculated and observed abundance of ^4He , neutrino decay must occur at a temperature (see eqn. 7),

$$T_{bd} \geq (fN_F) T_m (\eta_0/\eta_{ns})^{4/3} \equiv T_d^{(NS)} , \text{ for } T_{bd} \ll T_m . \quad (10)$$

Kolb & Scherrer [10] have calculated $T_d^{(NS)}$ for the case of radiative neutrino decay ($f = 1/2$) with two assumed values of $\eta_0 = 2 \times 10^{-11}$ and $\eta_0 = 2.4 \times 10^{-10}$. We have scaled their values appropriately for the present case ($f = 2/3$, $\eta_0 = 5.7 \times 10^{-11}$) to obtain $T_d^{(NS)}$ as a function of m_ν (fig. 1). The corresponding lifetime limits the range between $t_d \lesssim 2.5 \times 10^2 - 5 \times 10^5 \text{ sec}$ for $m_\nu \sim 1-25 \text{ MeV}$ (see fig. 3).

(*) The present Hubble expansion parameter H_0 and the black-body background temperature T_0 are parametrized as $h_0 \equiv H_0/100 \text{ km sec}^{-1} \text{ Mpc}^{-1}$ and $\theta_{2.7} \equiv T/2.7^\circ\text{K}$. Conservative observational limits are $0.5 \lesssim h_0 \lesssim 1$ [18] and $2.6 \lesssim T_0 \lesssim 2.9^\circ\text{K}$ [19].

Weinberg [15] has noted that the decay of a massive particle after primordial nucleosynthesis could, in principle, have reheated the universe above $\sim 5 \times 10^9$ °K, thus restoring the weak equilibrium value of the neutron-to-proton ratio and effectively restarting nucleosynthesis. However, massive neutrinos matter-dominate the universe too late for this to happen (see eqn. 7 and fig. 1), so the above bounds cannot be evaded in this manner.

Dicus et al. [10] have obtained similar bounds from consideration of the primordial deuterium (D) abundance, which is a rapidly decreasing function of η_{NS} and thus conceptually a more sensitive probe. Unfortunately, D is easily destroyed in subsequent stellar processing, so the observational lower bound to its primordial mass fraction [20], $X_p(D) \gtrsim 1.5 \times 10^{-5}$, is not restrictive enough to improve the bounds on η_{NS} obtained here. Hence we disagree with the claim by Krauss [22] that the D abundance imposes the constraint $T_{bd} \gtrsim T_m/10$ for all m_ν .

2.2 The Black-body Radiation Background

The relativistic $e^\pm$ from neutrino decay slow down rapidly by inverse-Compton scattering on thermal photons and then annihilate with each other^(*). This creates, non-thermal gamma rays with an energy density, $\Delta\rho_\gamma$, which may exceed that of the thermal photons,

$$\Delta\rho_\gamma/\rho_\gamma = fN_F (T_m/T_{bd}) . \quad (11)$$

This energy is transferred to the thermal photons via Compton scattering on thermal electrons, thus distorting their spectrum to a Bose-Einstein shape with occupation number $N_{BE} = (e^{x+\mu} - 1)^{-1}$; $x \equiv hv/kT_e$, where μ is a chemical potential and T_e is the (equilibrium) thermal electron temperature. To obtain the Planck form observed today [19] ($N_{pl} = (e^x - 1)^{-1}$), additional photons must be created by radiative processes such as bremsstrahlung and/or double Compton scattering [23-25]. We consider the latter process, which dominates [9] in a low baryon density universe, as demanded by the primordial ${}^4\text{He}$ abundance [10].

(*) The relevant time-scales are (a) Inverse Compton: $t_{e\gamma} \approx 3 m_e^2/4\sigma_T c \rho_\gamma E \approx 6.2 \times 10^{-12}$ sec $(T_\gamma/10^8 \text{ °K})^{-4} (m_\nu/10 \text{ MeV})^{-1}$, (b) non-relativistic $e^\pm$ annihilation $t_{ann} \approx 8/3 \sigma_T c n_{e^\pm} \approx 2.4 \times 10^{-10}$ sec $(T_\gamma/10^8 \text{ °K})^{-3} (m_\nu/10 \text{ MeV})^{2.85}$, where $n_{e^\pm} = 4 n_\nu \approx 1.5 \times 10^9 \text{ cm}^{-3} (m_\nu/\text{MeV})^{-2.85} T_\gamma^3 \gg n_e$ (see eqn. 20). The relativistic $e^\pm$ annihilation time-scale is much longer. Therefore, $t_{e\gamma} \ll t_{ann} \ll t$.

New photons are generated mainly at low frequencies and diffuse upwards in frequency through Compton scattering [23-24], thus gradually smoothing out the spectral distortion according to [25],

$$\frac{d\mu}{dt} = - \frac{8}{3t_{DC}} \frac{f(\mu)}{B(\mu)\phi(\mu)} \frac{I_3}{I_2} \int_{x_{coh}} \frac{dx}{x} \left[1 - \frac{N_{pl}}{N_{BE}} \right]. \quad (12)$$

Here, $x_{coh}(t)$ is the (dimensionless) frequency below which a Planck spectrum has already been attained at time t , and is obtained from [24]

$$x_{coh}^3 / (e^{x_{coh}} - 1) = (2\alpha/3\pi) I_3 (kT_e/m_e) f(\mu). \quad (13)$$

The functions $f(\mu)$, $\phi(\mu)$ and $B(\mu)$ express the deviation from a Planck spectrum [25], and $I_m \equiv \int x^m dx / (e^x - 1)$. By numerical solution of eqns (12-13), we find that the time required for complete thermalization of large distortions, $\Delta\rho_Y/\rho_Y > 1$, compares with the time-scale, t_{DC} , for photon generation by double Compton scattering in an isothermal plasma [24],

$$t_{DC} = (\pi/8\alpha) (m_e/kT_e)^2 t_C. \quad (14)$$

Here, $t_C \equiv (n_e \sigma_T c)^{-1} \approx 9.8 \times 10^{-5} \text{ sec} (T_Y/10^8 \text{ }^\circ\text{K})^{-3} (\Omega_B h_0^2 \theta_{2.7}^{-3})^{-1}$, is the Compton scattering time-scale for a thermal electron density $n_e \approx 10^{-5} \Omega_B h_0^2 (T_Y/T_0)^3 \text{ cm}^{-3}$, and $T_e = T_Y [f(\mu)]^{-1/4} \approx T_Y$ in the present case of a radiation dominated plasma. We note that a similar result is obtained when bremsstrahlung is the dominant photon creation process [23]. Then the requirement that a Planck spectrum be attained by today is,

$$\int_{T_0}^{T_{ad}} t_{DC}^{-1} (dt/dT_Y) dT_Y \gtrsim 1. \quad (15)$$

Using $dt/dT_Y \approx 6.1 \times 10^{20} f^{1/2} T_Y^{-3}$ for $t > t_d$ (assuming $T_{bd} \ll T_m$), we obtain, $T_{ad} > 10^7 \text{ }^\circ\text{K} (\Omega_B h_0^2 \theta_{2.7}^{-3})^{-1/3} \approx 3.2 \times 10^4 n_0^{-1/3}$, or equivalently^(*)

$$T_{bd} > 3.2 \times 10^4 n_{ns}^{-1/3} \equiv T_d^{(BB)}. \quad (16)$$

(*) This implies $t_d \lesssim 10^4 - 2 \times 10^5 \text{ sec}$ for $m_\nu \approx 1-50 \text{ MeV}$ (fig. 3) and corresponds to a critical red-shift $z_d \equiv T_{ad}/T_0 \approx 3.9 \times 10^6 (\Omega_B h_0^2)^{-1/3}$, below which large distortions cannot be thermalized. Danese and de Zotti [25] have obtained a smaller value, $z_d \approx 1.4 \times 10^6 (\Omega_B h_0^2)^{-1/3}$, which we find is due to their use of an approximation for x_{coh} which is only valid for $x \ll 1$. Thus they underestimate the time necessary for restoration of the complete Planck spectrum, since double Compton scattering restores the Wien region ($x > 1$) much later than the Rayleigh-Jeans region ($x < 1$), as shown by Lightman [24]. Our limits on t_d agree with unpublished calculations by R. Sunyaev [7].

However, for $m_\nu \gtrsim 50$ MeV, ${}^4\text{He}$ may undergo photofission and then its abundance cannot constrain Ω_B and η_{ns} . If $\Omega_B \sim 1$, bremsstrahlung becomes the dominant photon creation process and imposes weaker limits on $T_d^{(\text{BB})}$, which we have derived following Illarianov and Sunayev [23]. Furthermore, the distortions become weak for $m_\nu \gtrsim 100$ MeV (i.e. $T_{\text{bd}} > T_m$) and the limit $T_d^{(\text{BB})}$ follows the decrease of T_m with increasing m_ν , as shown in fig. 1.

2.3 Photofission of Primordial Deuterium

The arguments given above require neutrino decay to occur sufficiently early (see fig. 1) that the thermal photons (average energy $\langle \epsilon \rangle = 2.7 kT_\gamma$) are hot enough to Compton-scatter off the decay $e^\pm$ (average energy $\langle E \rangle \approx m_\nu/3$), in the Klein-Nishina regime, defined by [26] $4\epsilon E/m_e^2 > 1$, i.e.

$$T_{\text{bd}} > \{T_d^{(\text{NS})}, T_d^{(\text{BB})}\} > T_{\text{KN}} \equiv m_e^2 [4(2.7 k)(m_\nu/3)]^{-1} \quad (17)$$

The scattered photon then acquires a large fraction of the $e^\pm$ energy and may be boosted above the photofission thresholds for D ($Q_D = 2.225$ MeV) and ${}^4\text{He}$ ($Q_{\text{He}} \approx 20$ MeV). Indeed, even in the Thomson scattering regime ($T_\gamma \lesssim T_{\text{KN}}$), the average energy of the scattered photon [26], $\langle \epsilon_s \rangle = (4/3)(E/m_e)^2 \langle \epsilon \rangle$, exceeds Q_D above a threshold temperature,

$$T_{\text{th}} \equiv Q_D [(4/3)(m_\nu/3m_e)^2 (2.7 k)]^{-1} \ll \{T_d^{(\text{NS})}, T_d^{(\text{BB})}\}, \quad (18)$$

suggesting that a large fraction of primordial deuterium may be destroyed.

To investigate this, we have computed the total γ -ray generation rate,

$$d^2 n_{\text{tot}} / d\epsilon_s dt = \int dE N(E, t) \int d\epsilon n(\epsilon) (d^2 n_{\epsilon, E} / d\epsilon_s dt), \quad (19)$$

where, $d^2 n_{\epsilon, E} / d\epsilon_s dt$ is the spectrum scattered off monoenergetic photons by relativistic monoenergetic $e^\pm$ (eqn. 44, ref. [27]). The thermal background photons are assumed to have a Planck spectrum, $n(\epsilon) = (8\pi/h^3 c^3) \epsilon^2 / (e^{\epsilon/kT_\gamma} - 1)$. The source spectrum of $e^\pm$ from neutrino decay is,

$$Q(E, t) = KE(dn_\nu/dt), \quad (20)$$

where $K = 32/m_\nu^2$ is normalized to two helicity states, and $n_\nu \approx (4/11)n_D (T_\gamma^3/T_D^3) e^{-t/\tau_\nu}$ for $kT_\gamma < m_e$. The energy spectrum of $e^\pm$ is governed by the inverse-Compton energy loss rate $\langle \dot{E} \rangle$, according to the equation [26],

$$\partial N(E, t) / \partial t + \partial (N \langle \dot{E} \rangle) / \partial E = Q(E, t), \quad (21)$$

which has the Green's function,

$$G(E, E_0; t, t_0) = |\langle \dot{E} \rangle|^{-1} \delta(t - t_0 - \int_{E_0}^E dE' / \langle \dot{E}' \rangle). \quad (22)$$

By evaluating (using eqn. 46, ref. [27]),

$$\langle \dot{E} \rangle = \int d\epsilon n(\epsilon) \int d\epsilon_s (\epsilon_s - \epsilon) d^2 n_{\epsilon, E} / d\epsilon_s dt \quad (23)$$

we find that, $-\int dE' / \langle \dot{E}' \rangle \ll t$, in the present case, thus greatly simplifying the evaluation of $N(E, t)$ in an expanding universe. Then,

$$\begin{aligned} N(E, t) &= \int dt_0 \int dE_0 G(E, E_0; t, t_0) Q(E_0, t_0) \\ &\approx K(dn_v/dt) E / \langle \dot{E} \rangle, \end{aligned} \quad (24)$$

i.e., at any time t , we need consider only those $e^\pm$ created within the preceding interval $t_{ey} = E / \langle \dot{E} \rangle \ll t$, so that their energy losses due to expansion can be ignored. This result obtains even in the extreme Klein-Nishina limit where $N(E, t)$ is governed by an integro-differential equation [26] rather than eqn. (21).

Following Lindley [8,28], we now calculate the rate of depletion of the deuterium abundance (using eqns. 19, 23, 24),

$$\begin{aligned} [X(D)]^{-1} dX(D)/dt &\approx \int d\epsilon_s (\sigma_D / n_e \sigma_T) d^2 n_{tot} / d\epsilon_s dt \\ &\equiv \eta^{-1} f_D(m_\nu, T_\gamma) e^{-t/\tau_\nu}, \end{aligned} \quad (25)$$

where $\sigma_D(\epsilon_s)$ is the deuterium photofission cross section [29] and $\sigma_T(\epsilon_s)$ is the total cross section [30] for photon energy losses by Compton scattering and pair production^(*) on thermal electrons (of density $n_e \approx 18.1 \eta T_\gamma^3$). We have numerically computed $f_D(m_\nu, T_\gamma)$ for various neutrino masses and find it to be broadly constant over a wide range of temperature (see fig. 2). Thus we can estimate an average $\langle f_D \rangle$ and easily integrate eqn. (25). To allow primordial nucleosynthesis of the light elements (D, ^3He , ^4He , ^7Li) in accordance with observational data, $\eta \equiv n_B/n_\gamma$ must be in the range 10^{-10} – 10^{-9} [10,20]. For such values of η , the depletion rate is

(*) We ignore pair production on thermal photons since this process has a threshold, $\epsilon_s \gtrsim m^2/c$, which, for the temperatures of interest here ($T_\gamma < 10^9$ °K), is much higher than the energies ($\epsilon_s \sim 3$ MeV) at which f_D receives its maximum contribution. However, this process may be important for helium photofission, (considered in ref. [43]), for $T_\gamma > 10^8$ °K. This problem is under investigation.

sufficiently high that an upper limit on τ_ν follows from the very conservative requirement,

$$X(D)_{t=t_{ns}} / X(D)_{t \gg \tau_\nu} \approx \exp[\eta^{-1} \langle f_D \rangle e^{-t_{ns}/\tau_\nu}] < 100 \quad (26)$$

given that the primordial deuterium abundance peaks [21] at $t_{ns} \approx 300$ sec (corresponding to $T_Y \approx 8 \times 10^8$ K). Choosing appropriate values of η for each m_ν [10], we then obtain the stringent constraint, $\tau_\nu < 100$ sec for $5 < m_\nu < 300$ MeV, as shown in fig. 3. Our argument is strengthened by noting that for $T_Y \gg T_{KN}$, the decay $e^\pm$ initiate electromagnetic cascades which would exacerbate the depletion of deuterium [28].

These cosmological constraints have been derived assuming Dirac neutrinos. If the neutrino is instead a Majorana particle, the surviving cosmological number density increases by a factor $\sim 3-10$ [22], thus augmenting the cosmological effects and improving our limits on τ_ν .

3. CONSTRAINTS ON THE TAU NEUTRINO

Although we have so far considered only the decay mode $\nu_i \rightarrow e^- e^+ \nu_e$, the cosmological constraints (eqns. 10, 16, 26) apply to any decay mode which creates charged particles and/or photons. We now compare these with recent experimental limits on massive neutrino decays. An experiment at TRIUMF [11] has measured the branching ratio $R_\pi = \pi \rightarrow e\nu/\pi \rightarrow \mu\nu$, which sets upper limits on the mixing $|U_{ei}|^2$ of any ν_i (with $4 \lesssim m_\nu \lesssim 54$ MeV) emitted in π decay. We have extended these limits down to $m_\nu \sim 1$ MeV (using eqn. 3.16 of Shrock [6] and ref. [31]) and then translated them into lower limits on $\tau(\nu_i \rightarrow e^- e^+ \nu_e)$. This is shown in fig. 3, where we have identified ν_i with ν_3 . The CHARM Collaboration [12] have set lower limits on τ_{ν_3} by a direct search for the e^+e^- pairs from decaying tau neutrinos produced by heavy flavour (F) decays in a beam dump experiment. This is also shown in fig. 3 for $25 \lesssim m_{\nu_3} \lesssim 250$ MeV; however the recent measurements of a lower F mass would degrade these limits by a factor ~ 2 and restrict $m_{\nu_3} \lesssim 190$ MeV (C. Santoni, private communication). It is evident that the incompatibility of the experimental and cosmological constraints rules out $\nu_3 \rightarrow e^- e^+ \nu_e$ decays.

However, for $m_{\nu_3} \gtrsim m_\mu + m_e$, we must also consider the possibility of $\nu_3 \rightarrow \mu^- e^+ \nu_e$, with the lifetime [6],

$$\tau(\nu_3 \rightarrow \mu^- e^+ \nu_e) \approx 2.9 \times 10^4 \text{ sec} [m_{\nu_3}/\text{MeV}]^{-5} |U_{\mu 3}|^{-2} [f'(m_\mu/m_{\nu_3})]^{-1} \quad (27)$$

where $f'(x)$ is a phase space factor. A direct search for this decay has been conducted by the BEBC Collaboration [13] in a beam dump experiment and they obtain lower limits on τ_{ν_3} for $120 \lesssim m_{\nu_3} \lesssim 190$ MeV (see fig. 3). In addition, for any ν_i , $\tau(\nu_i \rightarrow \mu^- e^+ \nu_e)$ is restricted by upper limits on $|U_{\mu i}|^2$ obtained at KEK [13] through searches for secondary peaks in the decay spectrum of $K \rightarrow \pi \nu$. These limits are also shown, again identifying ν_i with ν_3 . Again, the conflict of the cosmological bound with experiment, albeit marginal, renders the possibility of $\nu_3 \rightarrow \mu^- e^+ \nu_e$ very unlikely. Furthermore, the MARK II [3] upper limit on m_{ν_3} may soon be lowered to ~ 100 MeV (B. Barish, private communication). (Nieves [5] has suggested the decay $\nu_3 \rightarrow \nu_2 \gamma \gamma$ when $m_{\nu_3} < m_\mu + m_e$, but it is seen from eqn. (3), replacing m_e , U_{e3} , U_{e1} by m_μ , $U_{\mu 3}$, $U_{\mu 2}$, that this cannot ever be fast enough to satisfy the cosmological bounds).

Thus ν_3 must be lighter than ~ 1 MeV. It would then have been relativistic at "decoupling" and should have decayed with a lifetime [9],

$$\tau_\nu \lesssim 2 \times 10^8 \text{ sec} [m_\nu / \text{MeV}]^{-2} \quad (28)$$

in order that the red-shifted energy density in the decay products be below present observational limits (see Pal [35] for a lucid exposition). The radiative decay modes do not satisfy this and are, in any case, observationally constrained to have $\tau_\nu > 10^{24}$ sec [32]. The channel $\nu_i \rightarrow 3\nu_i$, is also GIM suppressed, but may be mediated sufficiently rapidly by Higgs scalars in $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ [33], or through flavour changing neutral currents in extensions of $SU(2)_L \times U(1)$ [34]. However, both these proposals lead to conflict with experiment [35] and also imply unacceptable radiative decay lifetimes [36]. The "invisible" decay $\nu_i \rightarrow \nu_i J$ is possible in models in which neutrinos acquire masses through spontaneous violation of global lepton number, with an associated Goldstone boson, the majoron (J) [37,38]. The lifetime is generally unacceptably long [39], except in a recent model which also postulates simultaneous violation of sequential lepton number [40]. However this model is contrived in as much as it cannot be experimentally tested. It is then reasonable to conclude that ν_3 is stable ($\tau_{\nu_3} > \tau_{\text{universe}} \sim 3 \times 10^{18}$ sec) and therefore subject to the cosmological bound [1], $m_{\nu_3} < 100 \text{ eV}^{(*)}$.

(*) Except in the triplet Majoron model [38] in which cosmological neutrinos do not survive. In this case, astrophysical considerations impose the constraint $m_{\nu_3} < 40 \text{ keV}$ [41].

4. CODA

The above constraints on m_{ν_3} are in conflict with the expectation in $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ that $m_{\nu_3} \sim 10-100$ MeV [4]. We also exclude any massive neutrino in the range $1 < m_\nu < 50$ MeV and require $\tau_\nu < 100$ sec for heavier neutrinos up to $m_\nu \sim 300$ MeV. This is relevant to models with a rather low right-handed mass scale [14].

Our cosmological limits are applicable to any unstable particle which has a cosmological energy density comparable to that of neutrinos. The appropriate limit is that given here for a neutrino with the same value of fT_m (eqn. 6) as that of the hypothetical new particle. For example, the unstable massive photino considered by Kim et al. [15] is so constrained as to imply a restrictive limit on the photino mass, $m_{\tilde{\gamma}} > 150$ MeV, for a Peccei-Quinn symmetry breaking scale, $V_{PQ} > 10^9$ GeV, in their model. Conversely, any unstable particle expected in the lifetime range excluded by our limits must have an energy density sufficiently smaller than that of neutrinos. Such a constraint on massive (~ 100 GeV), unstable ($\tau \sim 10^8$ sec) gravitinos implies a severe restriction on the reheating temperature ($T_{\text{reheat}} \lesssim 10^9-10^{10}$ GeV) in supersymmetric inflationary models (see Ellis et al. [15]). Finally, a light tau neutrino implies that at most one new weakly interacting particle can be relativistic during primordial nucleosynthesis, and thus constrains the properties of light supersymmetric particles [42].

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FIGURE CAPTIONS

- Fig. 1 A neutrino of mass m_ν "decouples" at (radiation) temperature $T_\gamma = T_D$ and (matter) dominates the universe at $T_\gamma = T_m$. It must decay at $T_\gamma > \{T_d^{(NS)}, T_d^{(BB)}\}$ in order not to affect primordial nucleosynthesis or the black-body background. The decay $e^\pm$ undergo Klein-Nishina scattering off background photons for $T_\gamma \gg T_{KN}$. The average scattered photon energy exceeds the deuterium photofission threshold for $T_\gamma > T_{th}$. To prevent excessive destruction of D, neutrino decay must occur before nucleosynthesis begins, at $T_\gamma > 10^9$ °K. The time t_r applies to the radiation dominated era $T_\gamma \gg T_m$.
- Fig. 2 The normalized rate of depletion of the deuterium abundance, $f_D \equiv (n e^{t/\tau_\nu}) \cdot d \ln X(D)/dt$, due to photofission initiated by neutrinos of various masses.
- Fig. 3 Experimental lower limits on the lifetime $\tau(\nu_\mu \rightarrow e^- e^+ \nu_e)$ from: (a) TRIUMF [11], (b) CHARM [12] and on $\tau(\nu_\mu \rightarrow \mu^- e^+ \nu_e)$ from (c) BEBC [13] and (c') KEK [13]. Cosmological upper limits from: (d) deuterium photofission, (e) primordial nucleosynthesis, (f) black-body background. There is no allowed region below (g), the MARK II [3] upper bound on m_{ν_μ} . The limits (a,d-f) apply to any heavy neutrino and rule out the mass range $1 \lesssim m_\nu \lesssim 50$ MeV.

Fig. 1

Fig. 2

Fig. 3