CERN-TH.5632/90

LIMITS ON THE MUON AND TAU NEUTRINO MASSES FROM SN1987A

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ABSTRACT

When neutrinos have a Dirac mass, the right-handed partner can be emitted from the supernova. We calculate the emission rate due to all the important production mechanisms, and use SN1987A data to constrain the neutrino mass. We obtain that $m(\nu_\mu) \leq 14 \text{ KeV}$ and that either $m(\nu_\tau) \geq 34 \text{ MeV}$ or $m(\nu_\tau) \leq 14 \text{ KeV}$. These limits are reliable within a factor of 3.

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January 1990

The supernova SN1987A event has been extensively used to constrain some of the interactions and properties of models of particle physics^[1]. The basic idea is very simple. On the one hand, the observation^[2] of the neutrino pulse from the stellar collapse has allowed one to infer that at least an energy of 2×10^{53} erg was depleted from the star in the form of neutrinos. On the other hand, models of stellar collapse give a maximum possible gravitational energy of 4×10^{53} erg. Taking into account that the diffusion time extends over 1 - 10 s leads to a limit for any additional energy drain from the SN^[3]:

$$Q(\text{exotic}) \leq 2 \times 10^{53} \text{ erg/s} . \quad (1)$$

This bounds the production strength of light particles.

Exotic energy drain can be carried by any new light particle species produced in the SN core (where $T \sim 60$ MeV). Once produced, these light particles might leave the star without scattering if they are very weakly interacting. This is the so-called free-streaming regime. However, the particles can have stronger interactions with matter, be trapped in the dense stellar interior, and slowly diffuse out from the star core. This diffusion regime is the one that applies in the case of the ordinary neutrinos. The existence of the two (extreme) different regimes leads in general to an excluded range for the effective coupling strength associated with the new particle species.

In Refs. [4-6], the implications of the existence of a Dirac neutrino mass for the SN energetics have been discussed. In the presence of a Dirac mass term for the neutrino in the weak Lagrangian, right-handed (RH) neutrinos can be produced in the hot SN core. Once produced, they scatter off the stellar matter via the very same helicity flip mechanism. The SN constraints give a forbidden band of neutrino masses. In Refs. [4, 5], the considered process is scattering of thermal (ordinary) neutrinos on nucleons where RH states are produced. In Ref. [6], the emission rate of RH neutrinos from electron-positron annihilation has been calculated. The purpose of the present letter is to calculate the emission rate including all relevant processes and, in particular, nucleon-nucleon bremsstrahlung which has not been considered before. Nucleon-nucleon bremsstrahlung is the main source of SN left-handed neutrino (LH) emission and hence is expected to be important here also.

We start with neutron-neutron bremsstrahlung of RH neutrinos in the free-streaming regime. The emission rate (energy per unit volume and unit time) is

given by

$$q_{nn}(\nu_R) = 2 \int \prod_{i=1}^4 \frac{d^3 p_i}{(2\pi)^3 2\epsilon_i} \prod_{j=1}^2 \frac{d^3 q_j}{(2\pi)^3 2E_j} S|M|^2 (2\pi)^4 \delta^4(P_i - P_f) E_1 f_1 f_2 \quad (2)$$

where $|M|^2$ is the $nn \rightarrow nn\nu_R\nu_L$ matrix element squared summed over initial and final spins, S is a symmetry factor, $p_i(\epsilon_i)$ are neutron momenta (energies), $q_1(E_1)$ are the RH neutrino momentum (energy), $q_2(E_2)$ are the LH neutrino momentum (energy), P_i and P_f are the total initial and final four-momenta, and f_i is the Fermi-Dirac statistics factor for the initial neutrons, given by (μ is the chemical potential)

$$f_i = \frac{1}{e^{(\epsilon_i - \mu)/T} + 1} . \quad (3)$$

We work in the free-streaming regime and thus there is no statistics factor for the ν_R species. Also, the statistical factor corresponding to ν_L has a negligible effect. This is due to the fact that the LH neutrinos are in thermal equilibrium, which implies $< E_2 > \sim 3T$, and therefore the Pauli blocking factor is very close to one.

Since the temperature is much less than the nucleon mass, we can work in the non-relativistic limit. The squared matrix element is found to be, in this limit

$$S|M|^2 = 64 G_F^2 g_A^2 f^4 \left(\frac{m}{m_\pi} \right)^4 \frac{m_\nu^2 E_2}{(E_1 + E_2)^2 E_1} , \quad (4)$$

with $g_A = 1.2$ the axial weak coupling constant, $f = 1$ the strong coupling constant and $m = 0.80$ GeV the effective nucleon mass.

The thermal conditions of the SN core immediately after the collapse are such that a non-degenerate distribution for the nucleons is a very good approximation. We have worked the integral in Eq.(2) in this limit and obtained the emission rate*

$$q_{nn}(\nu_R) = \frac{2^2}{105\pi^{5.5}} m_\nu^2 G_F^2 g_A^2 f^4 (X_n \rho)^2 \left(\frac{m}{m_\pi} \right)^4 \frac{T^{3.5}}{m^{4.5}} \quad (5)$$

Here X_n is the neutron mass fraction. The neutron-proton bremsstrahlung rate is

$$q_{np}(\nu_R) = \frac{2^4}{35\pi^{5.5}} m_\nu^2 G_F^2 g_A^2 f^4 (X_n X_p \rho^2) \left(\frac{m}{m_\pi} \right)^4 \frac{T^{3.5}}{m^{4.5}} , \quad (6)$$

with X_p the proton mass fraction. We neglect p-p bremsstrahlung, since neutrons are more abundant in the SN core.

* The details of very similar calculations can be found in Ref. [7]

Using now Eq. (1), we get

$$\frac{4}{3}\pi R_c^3 [q_{nn}(\nu_R) + q_{np}(\nu_R)] \leq 2 \times 10^{53} \text{ erg/s}, \quad (7)$$

where R_c is the core radius. We take $\rho = 8 \times 10^{14} \text{ g/cm}^3$, $T = 60 \text{ MeV}$, and $R_c = 10 \text{ Km}$, and obtain

$$m_\nu \leq 46 \text{ KeV}, \quad (8)$$

for $X_n = 0.7$.

We now turn our attention to the production of RH neutrinos from the scattering of standard neutrinos off nucleons. In the case where the nucleon is a neutron, $\nu_L(k_1)n(p_1) \rightarrow \nu_R(k_2)n(p_2)$, the emission rate is

$$q_{nn}(\nu_R) = 2 \int \prod_{i=1}^2 \frac{d^3 p_i}{(2\pi)^3 2\epsilon_i} \prod_{j=1}^2 \frac{d^3 k_j}{(2\pi)^3 2E_j} S|M|^2 \delta^4(P_i - P_f) E_2 f_1 g_1 \quad (9)$$

where

$$g_1 = \frac{1}{e^{(E_1 - \mu_\nu)/T} + 1} \quad (10)$$

and μ_ν is the neutrino chemical potential. The matrix element is calculated to be

$$S|M|^2 = 4G_F^2 [(C_V + C_A)^2 K_2 \cdot p_2 K_1 \cdot p_1 + (C_V - C_A)^2 p_2 \cdot K_1 p_1 \cdot K_2 - m^2 (C_V^2 - C_A^2) K_2 \cdot K_1] \quad (11)$$

where $K_{1,2} = k_{1,2} - m_{\nu_{s1,2}}$ and s_1 is the neutrino spin. C_V and C_A are the vector and axial couplings for the neutron.

After integration we get

$$q_{nn}(\nu_R) = \frac{G_F^2}{4\pi^3} m_\nu^2 (C_V^2 + 3C_A^2) \frac{X_n \rho}{m} T^4 F(\mu_\nu/T) \quad (12)$$

with

$$F(y) = \int_0^\infty dx \frac{x^3}{e^{x-y} + 1} \quad (13)$$

The neutrino chemical potential is calculated using the relative density number $Y_\nu \simeq 0.04$ obtained in the SN model of Ref. [8]. We get $\mu_\nu = 95 \text{ MeV}$. Using this, the

values for ρ , T and R_c we mentioned before and combining the νn and νp emission rates, we conclude that the bound in Eq. (1) implies

$$m_\nu \leq 14 \text{ KeV}. \quad (14)$$

The production of RH neutrinos from electron-positron annihilation is subdominant. As has been shown in Ref. [6] the upper limit on the neutrino mass in this case is of the order of 1 MeV . Thus, we do not consider this source.

As has been discussed in Ref. [9], a source of uncertainty arises because we do not know which is the equation of state at supernuclear densities. To estimate the uncertainty, we follow that reference and change the central density by a factor of 2. All the other parameters change correspondingly. For example, the core temperature changes by a factor of 1.6, if one assumes an isentropic collapse. Notice that this allows the core temperature to span the conventional range of $T \sim 30 \text{ MeV} - 100 \text{ MeV}$. This implies a final uncertainty in the limits in Eq. (8) and (14) of a factor of 3 in either direction.

We see that the source $\nu_L N \rightarrow \nu_R N$ leads to a slightly better bound than $NN \rightarrow NN \nu_L \nu_R$. Within the reliability of the bounds the two constraints are very similar. We adopt Eq. (14) as our upper limit on Dirac neutrino masses.

The limits in Eq. (8) and (14) assume that we are in the free-streaming regime. To check that this is the case, we calculate the RH-neutrino mean free-path in the dense central core

$$\lambda(\text{core}) = \frac{1}{n_N \sigma} = 8 \times 10^2 \text{ Km} \left(\frac{50 \text{ KeV}}{m_\nu} \right)^2 \quad (15)$$

Here n_N is the nucleon density in the core and σ is the cross-section for the process $\nu_R N \rightarrow \nu_L N$, $N = \text{nucleon}$, which is the main source of opacity.* When $m_\nu < 50 \text{ KeV}$, we have that $\lambda(\text{core}) > R_c$. We are indeed in the free streaming regime when Eqs. (8) and (14) apply. By increasing the neutrino mass, however, we shall reach the region where the RH-neutrino is trapped. Clearly, for sufficiently large masses the diffusion time will be large enough for the emission rate to again be allowed by the generic limit in Eq. (1). In other words, there is a lower bound to the neutrino mass.

* The cross-section is energy-independent in the non-relativistic limit for the nucleon target. This simplifies the thermalization.

To calculate this lower bound, we assume the density profile given in Ref. [9] with the parameters that give the correct standard-neutrino emission rates^[7]. We find the RH-neutrinosphere by calculating the optical depth $\tau(R) = \int_R^\infty dr/\lambda$ and imposing $\tau = 2/3$. The thermal emission can then be easily calculated and the limit in Eq. (1) leads to

$$m_\nu \geq 34 \text{ MeV} . \quad (16)$$

One may wonder if for such mass values RH-neutrinos are produced. The answer is positive, as the following simple argument shows. The mean nucleon energy is $3T/2$. Even for $T = 30 \text{ TeV}$, there is enough energy for the RH-neutrino production, and Eq. (16) makes sense. We estimate this limit to be uncertain by a factor of 3.

Our Eqs. (14) and (16) constitute a forbidden range of Dirac neutrino masses. The consequences of our limits depend on the type of neutrino considered. There are no implications for the electronic type. Since the laboratory upper bound for the muon neutrino is much below the value in Eq. (16) we get the absolute upper limit

$$m(\nu_\mu) \leq 14 \text{ KeV} . \quad (17)$$

In the case of the tau neutrino, the present laboratory constraint is $m(\nu_\tau) \leq 35 \text{ MeV}/^{[10]}$. Given the uncertainty in our SN limits, we believe there is still an allowed window between the laboratory bound and ours:

$$m(\nu_\tau) \leq 14 \text{ KeV} \text{ or } m(\nu_\tau) \geq 34 \text{ MeV} . \quad (18)$$

We stress again that these equations, as well as Eq. (17), are reliable within a factor of 3.

In conclusion, we have calculated the upper bound on Dirac neutrino masses coming from RH-neutrino production by scattering of neutrinos on nucleons. Our result in Eq. (14) confirms the limit of 20 KeV in Ref. [4] and 30 KeV in Ref. [5]. Given the importance of this bound, we feel it is reassuring that we get essentially the same limit using a different model for the supernova. We have shown that nucleon-nucleon bremsstrahlung leads to the similar upper bound in Eq. (8). There is an important discrepancy, however, between Refs. [4] and [5]. Raffelt and Seckel find that there is still an allowed window for $m(\nu_\tau)$, while in Ref. [5] the authors obtain no window. Our results indicate that there is indeed an allowed range between $\sim 10 \text{ MeV}$ and 35 MeV for the tau neutrino mass.

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