Neutrino helicity flips via electroweak interactions

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Electroweak mechanisms via which neutrinos may flip helicity are examined in detail. Exact and approximate expressions for a variety of flip processes relevant in astrophysics and cosmology, mediated by W, Z, and γ exchange, including their interference, are derived for both Dirac and Majorana neutrinos (with emphasis on the former). It is shown that in general flip and nonflip cross sections differ by more than just a multiplicative factor of $m_\nu^2/4E_\nu^2$ contrary to what might be expected and that this additional dependence on helicities can be significant. It is also shown that within the context of the standard model with massive neutrinos, for $\nu e \to \nu e$ scattering, $\sigma_Z^{\text{flip}}/\sigma_\gamma^{\text{flip}} \approx 10^4$, independent of particle masses and energies to a good approximation. As an application, using some general considerations and the fact that the observed $\bar{\nu}_e$ burst from SN 1987A lasted several seconds, these weak-interaction flip cross sections are used to rule out μ and τ neutrino masses above 30 keV. Finally, some other consequences for astrophysics in general and supernovae in particular are briefly discussed.

I. INTRODUCTION

The masslessness of neutrinos in the standard electroweak model[1-3] results from the assumed absence of the $SU(2)_L$ singlet ν_R^i ($i=e,\mu,\tau$), which prohibits the usual Dirac mass term, and the fact that the simple Higgs structure of the theory leads to a lepton-number-conserving global symmetry which disallows a Majorana mass term of the form $\bar{\nu}_L^c \nu_L$. Given current experimental bounds from particle physics [$m_{\nu_e} \le 18$ eV, $m_{\nu_\mu} \le 0.25$ MeV, $m_{\nu_\tau} \le 35$ MeV (Ref. 4)], the possibility of nonzero neutrino masses is clearly an open one. More stringent bounds on the masses of stable neutrinos have been derived from big-bang cosmology.[5-7] The neutrinos we consider here would reasonably be expected to decay via mixing and other modes. Bounds from cosmology and astrophysics on the masses and lifetimes of unstable neutrinos have been derived in Refs. 8–10. In this paper we study electroweak mechanisms arising as direct consequences of these masses which result in the neutrino flipping its helicity. Within the context of the standard theory, a neutrino chirality eigenstate is a superposition of helicity (λ) eigenstates $\nu_\pm$, where $\lambda = \boldsymbol{\sigma}\cdot\mathbf{p} = \pm 1$. For a relativistic particle this translates into the statement that a ν_L (the left-handed neutrino) is predominantly in the $\lambda=-1$ state and a ν_R (the corresponding right-handed partner) is predominantly in the $\lambda=+1$ state, with small admixtures (of order m/E_ν) of the opposite helicity. A spin flip via any mechanism will consequently reduce the effective weak-interaction cross sections drastically, rendering a relativistic (Dirac) neutrino sterile or noninteracting in its passage through matter. Such behavior may have significant consequences in astrophysical and cosmological settings.

The discussion of such consequences in the literature has focused on helicity flips arising due to the anomalous magnetic moment of the neutrino and its consequent coupling to the photon. Here we note that within the context of the standard model, the neutrino is expected to have a tiny magnetic moment[11] directly proportional to its mass:[2]

$$\mu_\nu = \frac{3eG_F m_\nu}{8\pi^2\sqrt{2}} \approx 10^{-19}\mu_B \left[\frac{m_\nu}{1\ \text{eV}}\right], \tag{1}$$

where $\mu_B = e/2m_e$ is the electron Bohr magneton. (In various extensions of the standard model the magnetic moment may not be proportional to the mass. However our discussion throughout is confined to the standard model, with only the assumption of zero neutrino masses relaxed.) Since this results in a small flip cross section, any useful application inevitably requires the assumption of a nonstandard magnetic moment, several orders of magnitude larger than the value in (1). One of the aims of this paper is to show that the assumption of nonzero neutrino mass, without any additional nonstandard input, leads to flip cross sections via Z^0 and W exchange which are 4 orders of magnitude higher than those obtained from (1). (This refers to quantum scattering processes. In the presence of a strong, coherent magnetic field and in the absence of significant quantities of matter, the magnetic moment, although small, can be important, as discussed in Ref. 11. We do not consider this situation here.) This result, to a good approximation, is independent of particle masses and energies. The plan of this paper is as follows. In Sec. II we focus on flips arising through weak processes, and calculate cross sections for the following interactions, which may be of relevance in supernovae, neutron, and helium stars and the early Universe:

$$\nu^i_- e^- \to \nu^i_+ e^- , \quad \bar{\nu}^i_+ e^- \to \bar{\nu}^i_- e^- ,$$
$$\nu^i_- N \to \nu^i_+ N , \quad \bar{\nu}^i_+ N \to \bar{\nu}^i_- N . \tag{2}$$

Here $i=e$, μ, τ; N=nucleon; while the subscripts $\pm$ denote the helicity eigenvalues, as before. Although the focus is on Dirac neutrinos, we briefly discuss the modifications necessary if neutrinos are Majorana particles.

In Sec. III we analyze helicity flips occurring via the magnetic moment, giving both exact and approximate forms for the cross sections. We also consider the interference of this process with the process in Sec. II.

Finally, Sec. IV focuses on comparisons of the various results, conclusions, and a discussion of possible applications. In particular it uses the weak flip processes and the several second duration of the neutrino burst from SN 1987A to rule out μ and τ (Dirac) neutrino masses above 30 keV.

II. HELICITY FLIPS VIA W AND Z EXCHANGE

As a first step in obtaining the cross sections in (2), we focus on a specific process, mediated by Z^0 exchange:

$$\nu_\mu(k_1,\lambda_1)+e^-(p_1)\to \nu_\mu(k_2,\lambda_2)+e^-(p_2) . \tag{3}$$

Here the k_i and p_i are the particle momenta and λ_i is the neutrino helicity. The amplitude for (3) is given by

$$\mathcal{A}=\frac{G_F}{\sqrt{2}}\bar{u}(k_2,\lambda_2)\gamma^\mu(1-\gamma^5)u(k_1,\lambda_1)\cdot\bar{u}(p_2) \times\gamma_\mu(c_V-c_A\gamma^5)u(p_1) . \tag{4}$$

The u are the usual Dirac spinors and use has been made of the fact that when the processes (2) occur in typical astrophysical settings the center-of-mass energies are at most a few GeV; hence the amplitude may be written in its effective four-fermion form. The electron helicity indices have been suppressed since they will be summed over. The differential cross section in the center-of-mass frame is then given by

$$\frac{d\sigma}{d\Omega}=\frac{1}{64\pi^2 s}\frac{1}{2}\sum_{\text{el. spins}}|\mathcal{A}|^2 . \tag{5}$$

Here s is the square of the center-of-mass energy. We next write the matrix element squared in the form

$$\tfrac{1}{2}\sum_{\text{el. spins}}|\mathcal{A}|^2=\frac{G_F^2}{4}N^{\mu\nu}E_{\mu\nu} , \tag{6}$$

where

$$N^{\mu\nu}=\tfrac{1}{4}\mathrm{Tr}[(\not{k}_2+m)(1+\gamma^5\not{s}_2)\gamma^\mu(1-\gamma^5) \times(\not{k}_1+m)(1+\gamma^5\not{s}_1)\gamma^\nu(1-\gamma^5)] \tag{7}$$

and

$$E^{\mu\nu}=\mathrm{Tr}[(\not{p}_2+M)\gamma^\mu(c_V-c_A\gamma^5)(\not{p}_1+M) \times\gamma^\nu(c_V-c_A\gamma^5)] . \tag{8}$$

Here s_1 and s_2 are the spin four-vectors associated with the incoming and outgoing neutrino, respectively, while m and M are the neutrino and electron masses. If s_1 and s_2 are to represent helicity states in the center-of-mass frame in which the neutrino is relativistic they must satisfy

$$s_i\cdot s_i=-1, \quad s_i\cdot k_i=0, \quad \mathbf{s}_i\|\lambda_i\mathbf{k}_i \quad \text{for } i=1,2 . \tag{9}$$

After some Dirac algebra the neutrino tensor may be rewritten as

$$N^{\mu\nu}=\tfrac{1}{4}\mathrm{Tr}[(\not{k}_2-m\not{s}_2)\gamma^\nu(1-\gamma^5)(\not{k}_1-m\not{s}_1)\gamma^\nu(1-\gamma^5)] . \tag{10}$$

We now introduce two four-vectors associated with the neutrinos:

$$K_i^\mu=k_i^\mu-ms_i^\mu, \quad i=1,2 . \tag{11}$$

Because of the properties in (9) we see that

$$K_i\cdot K_i=0 ; \tag{12}$$

i.e., it is a lightlike vector.

We can now evaluate the traces and the contraction $N^{\mu\nu}E_{\mu\nu}$ in a straightforward way to obtain

$$\begin{aligned} N^{\mu\nu}E_{\mu\nu}=&16(c_V+c_A)^2(p_1\cdot K_1)(p_2\cdot K_2)\\ &+16(c_V-c_A)^2(p_1\cdot K_2)(p_2\cdot K_1)\\ &-16(c_V^2-c_A^2)M^2(K_1\cdot K_2) . \end{aligned} \tag{13}$$

Note that our particular choice of K_i $(i=1,2)$ enables us to write the square of the helicity-flip amplitude in a way completely analogous to the one for the *unpolarized* $\nu e^-\to\nu e^-$ process, with the replacement $k_i\to K_i$.

It is instructive to look at K_i in component form. In the center-of-mass frame, if

$$k^\mu=(E_\nu,|\mathbf{k}|\hat{\mathbf{k}}) , \tag{14}$$

where $\hat{\mathbf{k}}$ is a unit vector along the momentum of the neutrino, we derive from (9) that

$$s^\mu=\lambda m^{-1}(|\mathbf{k}|,E_\nu\hat{\mathbf{k}}) . \tag{15}$$

Here λ is the helicity of the neutrino which can take the values ±1. From this we see that

$$K(\lambda)=(E_\nu-\lambda|\mathbf{k}|)(1,-\lambda\hat{\mathbf{k}})=\eta(\lambda)(1,-\lambda\mathbf{k}) \tag{16}$$

with

$$\eta(\lambda)=E_\nu-\lambda(E_\nu^2-m^2)^{1/2} . \tag{17}$$

Note that for $m\ll E_\nu$ we have

$$\eta(-1)\approx 2E_\nu, \quad \eta(+1)\approx\frac{m^2}{2E_\nu} . \tag{18}$$

We see that for strictly massless neutrinos the "wrong" helicity K is identically equal to zero, as expected.

If we choose a coordinate frame in which the incoming neutrino moves toward the positive z axis and the incoming electron toward the negative z axis with common three-momentum $p=|\mathbf{p}_e|=|\mathbf{p}_\nu|$ and in which the scattered neutrino makes an angle θ with this z axis then we obtain, from Eqs. (13) and (16),

$$N^{\mu\nu}E_{\mu\nu}=16\eta_1(\lambda_1)\eta_2(\lambda_2)[(c_V+c_A)^2(E_e-\lambda_1 p)(E_e-\lambda_2 p)+(c_V-c_A)^2(E_e+\lambda_1 p\cos\theta)(E_e+\lambda_2 p\cos\theta)$$
$$-(c_V^2-c_A^2)M^2(1-\lambda_1\lambda_2\cos\theta)]\ , \tag{19}$$

where E_e is the center-of-mass energy of the incoming (and outgoing) electron. The subscripts 1 and 2 denote incoming and outgoing helicities for the neutrino. Written in this particular manner the above equation is a general expression from which both flip and nonflip cross sections can be readily obtained. The proportionality of $N^{\mu\nu}E_{\mu\nu}$ to $\eta_1(\lambda_1)\eta_2(\lambda_2)$ in conjunction with (18) indicates that flip and nonflip cross sections will be related by the expected factor of $m^2/4E_\nu^2$, but this is not the only dependence on helicities since the term in square brackets also changes with them. In the case of neutrino electron scattering, for instance, this additional dependence makes a significant difference in the magnitude in the cross section, since the electron three-momentum, which changes sign in (19) depending on the helicity, is comparable in magnitude to the electron energy.

Finally we find, for the differential cross section for flips,

$$\left[\frac{d\sigma}{d\Omega}\right]_{\mathrm{c.m.}}^{\nu_-^\mu e^-\to\nu_+^\mu e^-}=\frac{G_F^2}{16\pi^2 s}m^2[p^2(c_V-c_A)^2(1-\cos^2\theta)-M^2(c_V^2-c_A^2)\cos\theta+M^2(c_V^2+3c_A^2)]\ . \tag{20}$$

Here p^2 (using standard kinematics) is expressible in terms of Lorentz-invariant quantities as

$$p^2=\frac{\beta(s,m^2,M^2)}{4s}\ ,$$
$$\beta(x,y,z)=x^2+y^2+z^2-2xy-2yz-2zx\ . \tag{21}$$

If E_e, $E_\nu \gg m$, M then one may obtain from (20) (on integration and substitution for the vertex factors) the compact form (θ_W = weak mixing angle)

$$\sigma_z^{\mathrm{flip}}\simeq\frac{2G_F^2 m^2\sin^4\theta_W}{3\pi}\ . \tag{22}$$

Although written for the case of νe^- scattering, (20) is also a general one from which all the cross sections in (2) may be deduced by inserting the appropriate vertex factors. The flip cross section for $\nu_-^e e^-\to\nu_+^e e^-$ (which proceeds via both Z and W exchange) is readily obtainable via the standard Fierz transformation, which is equivalent to replacing $c_V\to c_V+1$ and $c_A\to c_A+1$ in (20).

Next, the flip cross section for $\bar\nu_+^\mu e^-\to\bar\nu_-^\mu e^-$ may be obtained by the substitution $c_A\to -c_A$ in (20) which yields

$$\left[\frac{d\sigma}{d\Omega}\right]_{\mathrm{c.m.}}^{\bar\nu_+^\mu e^-\to\bar\nu_-^\mu e^-}=\frac{G_F^2 m^2}{16\pi^2 s}[p^2(c_V+c_A)^2(1-\cos\theta^2)-M^2(c_V^2-c_A^2)\cos\theta+M^2(c_V^2+3c_A^2)]\ . \tag{23}$$

This reduces to the same compact form as in (22) in the event that the particle energies are much greater than their masses.

Finally, all these cross sections have their counterparts where the neutrino (or antineutrino) scatters off a nucleon rather than an electron. Since we did not neglect either of the masses, these cross sections can be directly obtained from (20) and (23) by substituting appropriate couplings c_A and c_V. If the nucleon is a proton, then $c_V=1-4\sin^2\theta_W$ and $c_A=+1$, whereas if the nucleon is a neutron then both c_V and c_A are -1. Of course, M now represents the nucleon mass. Note that all cross sections are quadratic in c_V and c_A hence it is only their relative sign which is important.

Before ending this section we briefly discuss how the above results would be modified if neutrinos were Majorana particles. The general properties of such neutrinos, which are their own antiparticles, have been discussed by Kayser,[12] where it has been pointed out that CPT invariance forbids such particles from having a nonzero magnetic moment, and that in the limit of zero neutrino masses there appears to be no way to distinguish between Dirac and Majorana neutrinos. Hence although helicity flips via electromagnetic channels are forbidden, they may proceed via the W- and Z^0-exchange mechanisms discussed above.

The analysis proceeds completely analogously to that performed above for Dirac neutrinos, except that the vector part of the neutrino current vanishes, since

$$\bar\nu_M\gamma^\mu\nu_M=(\bar\nu_M)^c\gamma^\mu(\nu_M)^c=-\bar\nu_M\gamma^\mu\nu_M\ , \tag{24}$$

where the subscript M denotes the Majorana nature of the fermion and c indicates charge conjugation. The last part of the equation follows from the required transformation of a fermion current (Dirac or Majorana) under charge conjugation. The neutral-current term thus has the form

$$-\bar u_f\gamma^\mu\gamma^5 u_i+\bar v_i\gamma^\mu\gamma^5 v_f\ , \tag{25}$$

where the second term arises from the fact that the ν_M operator cannot only destroy but also create Majorana neutrinos, unlike the corresponding Dirac operator. Using $v(p,s)=C\bar u^T(p,s)$, one can rewrite the current as

$$j_{\nu_m}^\mu=-2\bar u_f\gamma^\mu\gamma^5 u_i\ . \tag{26}$$

Incorporating this change into the Z^0-exchange amplitude (4) we had for Dirac neutrinos, we obtain

$$\left[\frac{d\sigma^m}{d\Omega}\right]_{\text{c.m.}}^{\nu^\mu_- e_- \to \nu^\mu_+ e_-} = \frac{G_F^2}{64\pi^2 s} \sum_{\text{el. spins}} |\mathcal{A}^m|^2 , \tag{27}$$

where

$$\tfrac{1}{2} \sum_{\text{el. spins}} |\mathcal{A}^m|^2 = 16m^2[p^2(c_V^2+c_A^2)(1-\cos^2\theta) + 2M^2 c_A^2(3+\cos\theta)] . \tag{28}$$

All other cross sections in (2) may be obtained from (28) in exactly the same manner as for Dirac neutrinos.

Finally, we note that the difference between the Dirac and Majorana flip cross sections (20) and (28) is also proportional to the square of the neutrino mass and thus vanishes for massless particles, as expected. Considerations similar to those above have also been discussed in Ref. 13.

III. HELICITY FLIPS VIA THE MAGNETIC MOMENT

As mentioned in the Introduction, a Dirac neutrino may flip its helicity via photon exchange if it has a nonzero magnetic moment. The cross section for this process was derived by Bethe.[14] Here we give a more exact expression and also evaluate the interference between the electromagnetic and weak flip amplitudes, which may be useful if the neutrino magnetic moment is anomalously large compared to the standard-model prediction.

The amplitude for the photon-exchange process is

$$\mathcal{A}^\gamma = \frac{\kappa'}{-q^2}[\bar{u}(k_2,\lambda_2)(i\sigma^{\mu\nu}q_\nu) \times u(k_1,\lambda_1)\cdot\bar{u}(p_2)\gamma_\mu u(p_1)] . \tag{29}$$

Here $(-q^2)$ is the square of the transferred momentum and κ' is the anomalous magnetic moment of the neutrino. On evaluation of the traces and after some standard kinematic manipulation one obtains (in the notation of Sec. II)

$$\tfrac{1}{2} \sum_{\text{el. spins}} |\mathcal{A}_\gamma|^2 = \frac{-2p^2\kappa'^2}{m^2 q^2}\left[(2p^2+m^2+2E_\nu E_e) \times \left[2+\frac{q^2}{2p^2}\right] + 2M^2\right] . \tag{30}$$

The helicity state chosen for the incoming (outgoing) neutrino is negative (positive). In the approximation of low-momentum transfer and if E_e, $E_\nu \gg M$, m one obtains

$$\tfrac{1}{2} \sum_{\text{el. spins}} |\mathcal{A}_\gamma|^2 = \frac{-16p^4\kappa'^2}{m^2 q^2} \tag{31}$$

and

$$\left[\frac{d\sigma}{dq^2}\right]_{\text{c.m.}} = \frac{\pi\alpha^2}{m^2}\kappa'^2\left[\frac{1}{-q^2}\right] . \tag{32}$$

Integration then gives the total cross section

$$\sigma_\gamma^{\text{flip}} \approx \frac{\pi\alpha^2\kappa^2}{M^2}\ln\left[\frac{q^2_{\max}}{q^2_{\min}}\right] , \tag{33}$$

where we have redefined κ by a ratio of the particle masses since this will allow it to be expressed in terms of the more standard electron Bohr magneton, which we shall use in the next section. In this form it is identical to the result in Ref. 14 and is adequate for most applications of astrophysical interest involving the magnetic moment, with the added approximation that the logarithmic term is of order unity. (A more careful evaluation using screening considerations yields a value close to 6, as shown in Ref. 15.)

Next, we address the possibility that the amplitude for Z^0 exchange (4) and magnetic moment (29) might interfere if they happen to be of comparable magnitude: i.e.,

$$\sigma^{\text{total}}_{\nu^\mu_- e \to \nu^\mu_+ e} = \sigma_Z + \sigma_\gamma + \sigma_{\text{int}} . \tag{34}$$

The square of the interference amplitude is given by

$$|\mathcal{A}_{\text{int}}|^2 = \mathcal{A}_\gamma \mathcal{A}_Z^\dagger + \mathcal{A}_\gamma^\dagger \mathcal{A}_Z . \tag{35}$$

Using the transformation properties of bilinear covariants under Hermitian conjugation one obtains from (4) and (29), after the usual trace evaluation and kinematics,

$$\tfrac{1}{2} \sum_{\text{el. spins}} |\mathcal{A}^{\text{int}}|^2 = 4[(E_\nu E_e + p^2)(c_A - c_V) - 2M^2 c_V] . \tag{36}$$

Under the same approximations as those leading to (33) above, one obtains

$$\left[\frac{d\sigma^{\text{int}}}{dq^2}\right]_{\text{c.m.}}^{\nu_- e_- \to \nu_+ e_-} = \frac{-\alpha G_F \kappa}{2\sqrt{2}s}\frac{m}{M} . \tag{37}$$

An order-of-magnitude comparison may now be made with (20) which shows that for neutrinos with a mass in the few keV range, the interference cross section is comparable to the weak flip cross section when $\kappa \approx 10^{-13}$, which is roughly the current upper bound (from the supernova) on this quantity.

IV. DISCUSSION AND DERIVATION OF MASS LIMIT ON μ AND τ NEUTRINOS

In the two preceding sections we have presented a detailed analysis of the various mechanisms via which a neutrino (Dirac or Majorana) may flip its spin. In the case of a Dirac particle this leads to a drastic alteration in its properties in passage through matter. This section focuses on conclusions and some applications that follow from the cross sections derived above.

We first note that as evident in (19), a simple multiplicative factor of $m_\nu^2/4E_\nu^2$ is not the only difference between flip and nonflip amplitudes. In the center-of-mass frame, the additional dependence on helicities (i.e., the presence of λ_i's inside the square brackets) makes a significant difference whenever target (the particle off which the neutrino scatters) energies are high compared to target masses. For the case of electron targets, as mentioned earlier, this changes the cross sections by an order of magnitude. While the flip cross section stays essentially constant if neutrino and electron energies are high compared to their masses [see Eq. (22)], the nonflip cross sections vary, as is well known, with the square of the center-of-mass energy. For these particles it is also interesting to examine the relative magnitudes of the flip cross sections obtained from the two main types of flip mechanisms, i.e., magnetic moment and W, Z exchange. Since the neutrino magnetic moment is not experimentally known, we will consider two very different values, that predicted by theory and the current upper bound.

Consider (33), which, with the approximation mentioned earlier may be expressed as

$$\sigma_\gamma^{\text{flip}} \approx \frac{\pi\alpha^2\kappa^2}{M^2} = \alpha\kappa^2\mu_B^2 = \alpha\mu_\nu^2 \ . \tag{38}$$

The standard model predicts, from (1), a cross section proportional to the square of the neutrino mass. The Z-exchange cross section (22) also exhibits this proportionality. Their ratio, independent of the mass and energy of the particles to a good approximation, is thus

$$\frac{\sigma_Z^{\text{flip}}}{\sigma_\gamma^{\text{flip}}} \approx \frac{64\pi^2}{27\alpha^2}\sin^4\theta_W \approx 10^4 \ . \tag{39}$$

However, experimental upper bounds on the magnetic moment are far above the predicted value. Bounds of $\mu_\nu \le 5\times 10^{-13}\mu_B$ have recently been obtained[16,17] from observations of neutrinos from SN 1987A. Using this value in (38),

$$\sigma_\gamma^{\text{flip}} \approx 6.2\times 10^{-50}\ \text{cm}^2 \ . \tag{40}$$

Since we are now considering a nonstandard moment, the ratio of interest can no longer be assumed to be mass independent. We find that when $m_\nu \simeq 10$ keV is used in (22) one has

$$\frac{\sigma_Z^{\text{flip}}}{\sigma_\gamma^{\text{flip}}} \approx 1 \ . \tag{41}$$

This shows that the neutrino mass required for the Z flip cross section to exceed the maximum allowed magnetic moment flip cross section is relatively low.

The cross sections for nucleon targets, also obtainable from (19) and (20) by the substitutions mentioned earlier, become relevant when one considers, for instance, the core of a collapsed star, as we do below. Again, the flip cross sections are relatively insensitive to neutrino energies as long as these are small compared to the nucleon mass and are typically enhanced over those due to electron targets by a factor of $\approx 20-25$.

Next, we briefly mention possible applications that our results may have in astrophysics and cosmology. Although there has been extensive discussion of these, the focus has been on the effect due to the magnetic moment.[18-20] If the magnetic moment is, however, close to that predicted by the standard theory, then the flip mechanisms discussed in Sec. II dominate for the same neutrino mass and could be important under conditions where the neutrino traverses significant amounts of matter. In supernovae, for instance, the gravitational energy of collapse is equally distributed among all three species of neutrinos. If the μ and τ neutrinos have masses in the keV range, significantly large numbers of these could flip helicity and escape quickly, not undergoing the slow thermal diffusion expected. Using the Z flip cross sections above and typical values for the core density, one finds that a neutrino with mass in the keV range will flip helicity and leave the core in several milliseconds, which is much shorter than the cooling/diffusion time of several seconds. We also note here that direct pair production of flipped neutrinos is possible, which may be a channel for rapid energy loss. The cross sections for such processes may be easily obtained by "crossing" the amplitudes given in Sec. II. The consequent energy drain and cooling could affect the energetics of supernova bursts, currently a topic of active investigation by several groups.[21-25] (These considerations would not be relevant if neutrinos were purely Majorana particles.)

Finally, we can establish a bound on the μ and τ neutrino mass using the SN 1987A neutrino measurement,[26,27] in the same spirit as was done for the e neutrino magnetic moment in Ref. 16. We compare the expected luminosity of flipped neutrinos to the maximum neutrino luminosity from a cooling neutron star in the standard, non-helicity-flip case, about 10^{53} ergs/sec. It is argued that a larger luminosity would result in cooling of the star that is too rapid to explain the signal observed from SN 1987A which, in the standard case, is due to thermal $\bar{\nu}_e$'s emitted from the neutron star's surface and was observed to last 5–10 sec.

The expected luminosity of flipped neutrinos can be expressed as

$$L \simeq 2N_\nu \int \langle \epsilon_\nu \rangle (n\sigma_Z^{\text{flip}} c) n_\nu dV \ , \tag{42}$$

where N_ν is the number of neutrino flavors that flip (the factor of 2 includes the antineutrinos). n_ν and $\langle \epsilon_\nu \rangle$ are the neutrino number density and the average energy, and are about $1.8(2\pi)^{-1}(kT/\hbar c)^3$ and $3.15kT$, respectively, for μ and τ neutrinos which have a thermal distribution. The quantity in parentheses is the rate of neutrino flipping per neutrino. For the present situation we find that the flip rate from nucleons gives an effective cross section approximately 20 times larger than that from electrons, so here n refers to the nucleon density and $\sigma_Z^{\text{flip}} \simeq 3\times 10^{-47}(m/50\ \text{keV})^2\text{cm}^2$. The integral in (45) is proportional therefore to the mass average T^4 in the neutron star:

$$L \simeq 1.4\times 10^{53} N_\nu (m/50\ \text{keV})^2 \times \langle T/30\ \text{MeV} \rangle^4\ \text{ergs/sec} \ , \tag{43}$$

where a $1.4M_\odot$ neutron star is assumed. Setting $L < 10^{53}$ ergs/sec, we find

$$m \le 42 N_\nu^{-1/2} \langle T/30\ \text{MeV} \rangle^{-2}\ \text{keV} \ . \tag{44}$$

For the reference value of $T = 30$ MeV this translates to a limit of ~ 30 keV, which is nearly the same as that mentioned in Ref. 28.

However, if the optical depth, i.e., $\int n\sigma_{\text{flip}} dr$ is unity or greater, neutrinos will still be trapped. The rate of reflipping being the same as that of flipping, this occurs for an average density of 10^{15} g/cm^3 and a neutron star radius of 10 km when $m_\nu = 0.5$ MeV. The relevant quantity, however, is the diffusion time scale, $\approx 3R^2 n\sigma_{\text{flip}}/c$, where R is the star radius. If this is greater than 1–2 sec then flipped neutrinos are effectively trapped on the overall cooling time scales. For the previous numbers this occurs if $m_\nu > 100$ MeV. This is far above the experimental limit for μ and τ neutrinos. Hence we see that reflipping considerations do not provide a window on our limit.

This limit should be viewed as somewhat tentative. A firmer limit will depend on including the cooling due to helicity flips in a consistent way. Note however that it holds for mu, tau, and possible fourth-generation neutrinos, stable or unstable with a lifetime greater than a couple of seconds. (This limit, and the implicit assumption that the neutrino does not decay in the few seconds it takes to diffuse through the supernova core are consistent with the bounds in Refs. 8–10.) Two important differences from the magnetic moment limit calculations should be mentioned. The first is that (44) has a large temperature sensitivity. Second, the helicity flips in this case do *not* lead to neutronization. Leptons must continue to diffuse to the surface, even if the star is cold. It is conceivable that otherwise rapid cooling due to helicity flips is self-regulating and that thermal pairs created by the *e*-neutrino diffusion remains adequate to explain the observed signal from SN 1987A. Nevertheless, the limit in (44) should be of the correct magnitude.

Helicity flips via W,Z exchange could also be significant under the high density and temperature conditions prevailing in the early Universe. The significance of flips has been discussed in Ref. 29 in connection with the magnetic moment mechanism, where their effect on helium synthesis and the expansion rate of the Universe has been discussed. Using arguments similar to those used in that paper to derive upper bounds on μ_ν, one can in an analogous manner obtain a bound on the neutrino mass using the weak flip cross sections. However the mass bound so obtained is roughly 2 orders of magnitude higher than that derived above from supernova cooling constraints.

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