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Supernova 1987A and the secret interactions of neutrinos

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By using SN1987A as a "source" of neutrinos with energy ~ 10 MeV we place limits on the couplings of neutrinos with cosmic background particles. Specifically, we find that the Majoron–electron-neutrino coupling must be less than about 10^{-3}; if neutrinos couple to a massless vector particle, its dimensionless coupling must be less than about 10^{-3}; and if neutrinos couple with strength g to a massive boson of mass M, then g/M must be less than 12 MeV^{-1}.

I. INTRODUCTION

Supernova 1987A in the Large Magellanic Cloud[1] produced a pulse of neutrinos which was detected by underground neutrino detectors.[2,3] The great distance to the supernova, $D=55\pm15$ kpc, and the concommitant long travel time, affords a unique opportunity to place limits on the properties of neutrinos, limits that in some instances cannot be matched by terrestrial experiments. Limits on neutrino mass,[4] lifetime,[5] and mixing angles[5] have been set using the information obtained from SN1987A. In this paper we consider the limits which can be placed on "secret" interactions of neutrinos with cosmic background particles (CBP's). By secret interactions, we mean interactions not shared by charged particles, i.e., interactions beyond those in the $SU(3)\times SU(2)\times U(1)$ model.

Although the interactions of neutrinos with "matter" (electrons, protons, neutrons, nuclei, etc.) are weak, it is possible that neutrinos have "stronger than weak" interactions with other unknown particles (e.g., Majorons[6,7]), or with themselves.[6,7] If a particle is stable and weakly interacting, it should be present today as a CBP. The detection of neutrinos from SN1987A requires that the mean free path of neutrinos through the CBP's is comparable to or greater than the distance to the supernova. This results in limits to the cross sections of neutrinos with themselves and with other particles ($\sigma \lesssim 10^{-25}$ cm^2).

The neutrino events detected in the Kamiokande II[2] and Irvine-Michigan-Brookhaven[3] (IMB) underground detectors are in qualitative agreement with the predicted neutrino flux from a type-II supernova. The data can best be fit if the bulk of the events are $\bar{\nu}_e$ captures, $\bar{\nu}_e p \rightarrow e^+ n$, with an incident $\bar{\nu}_e$ flux of the order of 10^{10} cm^{-2} (Ref. 8). Since this is about what is predicted, any substantial decrease in the $\bar{\nu}_e$ flux either because of the decay of neutrinos in flight or because of the scattering of neutrinos in flight can be ruled out. Although the data also strongly suggests the existence of some $\nu_i e^- \rightarrow \nu_i e^-$ ($\nu_i = \nu_e, \bar{\nu}_e, \nu_\mu, \bar{\nu}_\mu, \nu_\tau, \bar{\nu}_\tau$) scatterings in addition to $\bar{\nu}_e$ capture, the identity of the incident neutrinos in such processes cannot be ascertained, so we must focus on limits to the interactions of $\bar{\nu}_e$'s.

II. NEUTRINO MEAN FREE PATH

To start, let us assume that $\bar{\nu}_e$'s with energy $E=(|\mathbf{p}|^2+m^2)^{1/2} \gg m_{\bar{\nu}_e}$ are emitted from the supernova and scatter off a background of particles (denoted as X) whose phase-space density is $f_X(\mathbf{p})$. The Boltzmann equation for the evolution of the neutrino phase-space density $f(\mathbf{p})$ including the effects of $\nu(p)X(p_x) \rightarrow \nu(p')X(p'_X)$ scattering is

$$2E\frac{df(\mathbf{p})}{dt}=-f(\mathbf{p})\int\frac{d^3\mathbf{p}_X}{2E_X}f_X(\mathbf{p}_X)\int\frac{d^3\mathbf{p}'_X}{2E'_X}[1\pm f_X(\mathbf{p}'_X)]\int\frac{d^3\mathbf{p}'}{2E'}[1-f_\nu(\mathbf{p}')]\frac{(2\pi)^4}{(2\pi)^9}\delta^4(p+p_X-p'-p'_X)$$
$$\times|M(\nu X\to\nu' X')|^2 . \tag{1}$$

Note that X could also be a neutrino. In Eq. (1), the $+$ ($-$) sign applies if X is a boson (fermion).

The elastic reaction $\nu X\to\nu X$ does not lead to a decrease in the neutrino flux; however, if the background X is light ($M \ll E$) the reaction can lead to substantial energy loss of the neutrino. The relevant detectors[2,3] have threshold energies ~ 10–20 MeV, and a low-energy final-state neutrino will have energy much less than the detection threshold, and is effectively removed from the "detectable" flux. Since the CBP's are expected to have temperatures $\sim 10^{-4}$ eV, scattered neutrinos will often lose significant energy and be removed from the detectable flux. Also, because the initial X has such low energy, the production of a neutrino of momentum $|\mathbf{p}|$ from the collision of an incident neutrino of momentum $|\mathbf{p}'|$ has been in Eq. (1).

If we assume that $f_X(\mathbf{p}_X)$ and $f(\mathbf{p}')$ are much less than one (i.e., no Bose condensation or Fermi degeneracy), the occupancy factors $(1\pm f)$ can be neglected. With the usual definition of the cross section

$$\sigma(s)=\frac{1}{|\mathbf{v}_X-\mathbf{v}_\nu|}\frac{1}{2E}\frac{1}{2E_X}\int\frac{d^3\mathbf{p}'}{2E'_X}\int\frac{d^3\mathbf{p}'}{2E'}\frac{(2\pi)^4}{(2\pi)^6}\delta^4(p+p_x-p'-p'_x)|M(\nu X\to\nu' X')|^2 , \tag{2}$$

the Boltzmann equation becomes

$$-\frac{1}{f}\frac{df(\mathbf{p})}{dt}=\int\frac{d^3\mathbf{p}_X}{(2\pi)^3}f(\mathbf{p}_X)|\mathbf{v}_X-\mathbf{v}_\nu|\sigma(s) , \tag{3}$$

where $s=(p+p_X)^2$. If we consider the evolution of $f(\mathbf{p})$ for $\mathbf{p}$ in the direction from the source to the detector

$$-\frac{1}{f}\frac{df}{dy}\equiv\lambda^{-1}=\int\frac{d^3\mathbf{p}_X}{(2\pi)^3}f_X(\mathbf{p}_X)\frac{|\mathbf{v}_X-\mathbf{v}_\nu|}{|\mathbf{v}_\nu|}\sigma(s) , \tag{4}$$

where $y(\equiv t\,|\mathbf{v}_\nu|)$ is the distance from the source in the direction of the detector. The right-hand side (RHS) of Eq. (4) is simply the inverse of the mean free path λ.

With $|\mathbf{v}_X|=0$ and $\sigma(s)=\text{const}\equiv\sigma_0$, the usual result $\lambda^{-1}=n_X\sigma_0$ is obtained where the number density of X particles is

$$n_X=\int\frac{d^3\mathbf{p}_X}{(2\pi)^3}f_X(\mathbf{p}_X) . \tag{5}$$

However if $|\mathbf{v}_X|\neq 0$ or if $\sigma(s)$ depends upon s, Eq. (4) must be integrated to find the mean free path.

The mass of the electron neutrino is known to be less than about 20 eV (Ref. 9), and the energies of the detected neutrinos are 10 MeV or greater, so the mass of the neutrino can be ignored. In this limit

$$s\to M^2+2EE_X\left[1-\frac{|\mathbf{p}_X|}{E_X}z\right],\quad |\mathbf{v}_X-\mathbf{v}_\nu|\to\left[1+\frac{|\mathbf{p}_X|^2}{E_X{}^2}-2\frac{|\mathbf{p}_X|}{E_X}z\right]^{1/2} , \tag{6}$$

where $z\equiv\cos\theta$ and θ is the incident $\nu-X$ angle. The mean free path is then given by

$$\lambda^{-1}=\frac{1}{4\pi^2}\int_0^\infty d\,|\mathbf{p}_X||\mathbf{p}_X|^2f(\mathbf{p}_X)\int_{-1}^{1}dz\left[1+\frac{|\mathbf{p}_X|^2}{E_X{}^2}-2\frac{|\mathbf{p}_X|}{E_X}z\right]^{1/2}\sigma(s) . \tag{7}$$

We consider two limits for M: nonrelativistic (NR) $|\mathbf{p}_X|/E_X\to 0$ and extreme-relativistic (ER) $|\mathbf{p}_X|/E_X\to 1$. In these limits, Eq. (7) gives

$$\lambda^{-1}=\begin{cases}\dfrac{\sqrt{2}}{4\pi^2}\displaystyle\int_0^\infty dE_XE_X{}^2f(E_X)\int_{-1}^{1}dz\,(1-z)^{1/2}\sigma(s=2EE_X(1-z)) & \text{(ER)},\\ n_X\sigma(s=M^2+2EM) & \text{(NR)} .\end{cases} \tag{8}$$

We will also consider two forms for the cross section: $\sigma(s)=a/s$ and $\sigma(s)=as/M^4$, where a is a model-dependent, dimensionless constant. The first form would apply, for instance, if the scattering is mediated by the exchange of a massless particle. The second form describes scattering mediated by the exchange of a boson of mass $M\gg s$.

The mean free path for different choices of $\sigma(s)$ is given in terms of the incident neutrino energy E by

$$\lambda^{-1}=\begin{cases}\dfrac{2an_X}{E}\dfrac{\int_0^\infty dE_X E_X f(E_X)}{\int_0^\infty dE_X E_X{}^2 f(E_X)}\quad [\sigma(s)=a/s]\ ,\\ \dfrac{8an_X}{3M^2}\quad [\sigma(s)=a/M^2]\ ,\\ \dfrac{16}{5}\dfrac{aE}{M^4}n_X\dfrac{\int_0^\infty dE_X E_X{}^3 f(E_X)}{\int_0^\infty dE_X E_X{}^2 f(E_X)}\quad [\sigma(s)=as/M^4]\ .\end{cases}\tag{9}$$

Note that without including the velocity factors, $\lambda\rightarrow 0$ for $\sigma=a/s$.

It is usually the case that the phase-space density of X's can be written in terms of a thermal distribution: $f_X(\mathbf{p}_X)=[\exp(E_X/T)+\epsilon]^{-1}$, where $\epsilon=-1$ for Bose-Einstein statistics, $\epsilon=+1$ for Fermi-Dirac statistics, and $\epsilon=0$ for Maxwell-Boltzmann statistics. It then follows that

$$\lambda^{-1}=\begin{cases}\dfrac{2a}{ET}n_X I_{12}\quad [\sigma(s)=a/s]\ ,\\ \dfrac{8a}{3M^2}n_X\quad [\sigma(s)=a/M^2]\ ,\\ \dfrac{16aET}{5M^4}n_X I_{32}\quad [\sigma(s)=as/M^4]\ ,\end{cases}\tag{10}$$

where

$$I_{mn}=\frac{\int_0^\infty dx\, x^m f(x=E/T)}{\int_0^\infty dx\, x^n f(x=E/T)}\ .\tag{11}$$

The functions I_{12} and I_{32} are given by

$$I_{12}=\begin{cases}\frac{1}{2}\ \ (\text{Maxwell-Boltzmann})\ ,\\ \dfrac{\pi^2}{12\zeta(3)}=0.684\,22\ \ (\text{Bose-Einstein})\ ,\\ \dfrac{\pi^2}{18\zeta(3)}=0.456\,17\ \ (\text{Fermi-Dirac})\end{cases}\tag{12}$$

and

$$I_{32}=\begin{cases}3\ \ (\text{Maxwell-Boltzmann})\ ,\\ \dfrac{\pi^4}{30\zeta(3)}=2.70\ \ (\text{Bose-Einstein})\ ,\\ \dfrac{7\pi^4}{180\zeta(3)}=3.15\ \ (\text{Fermi-Dirac})\ .\end{cases}\tag{13}$$

III. LIMITS ON SECRET INTERACTIONS

It is unavoidable that all limits obtained will be model dependent, and involve assumptions about the number of densities of CBP's. However, they can be obtained in a self-consistent manner. In this section we will first consider two generic cases: neutrino coupling to a massless spin-one boson and neutrino coupling to a massive spin-one boson. We will then consider a specific model: the Majoron model.[6]

We will assume that neutrinos interact with a spin-one boson X through a vector coupling of the form $g_i\bar{\nu}_i\gamma_\mu\nu_i X^\mu$ ($i=e,\mu,\tau,\ldots$). Of course, the interaction must be "secret"; i.e., the X cannot couple (or if it does, only couples *very* weakly) to charged particles. If the boson is massless, then $s\approx 2ET\sim(60\text{ eV})^2$. If the mass of the X exceeds ~ 60 eV, it will be considered "massive;" otherwise it will be considered "massless."

We will further assume that the neutrino and the X temperature is that in the standard hot big-bang cosmology,[10] $T=(\frac{4}{11})^{1/3}T_\gamma\simeq 1.9$ K, so that the number density of X's is $n_x=(1.9/2.7)^3 n_\gamma\simeq 139\text{ cm}^{-3}$, where n_γ is the present photon number density. If neutrinos are massive, they may annihilate into X's in the early Universe and not survive as CBP's. This would remove the neutrinos as CBP's and increase the temperature of the X's relative to photons. The factor by which it is increased is model dependent (how many species of massive ν's, etc.), and hence the choice $T=1.9$ K for the X is a *conservative* one. If neutrinos are present we will assume the same temperature, 1.9 K, for them. This results in a number density of 55 cm^{-3} for each type of neutrino ($\nu_e,\bar{\nu}_e,\nu_\mu,\bar{\nu}_\mu,\ldots$).

A. Neutrino coupling to a massless ($M\leq 60$ eV) boson

The scattering of the neutrino with background X's is described by the familiar (ER) Compton form

$$\frac{d\sigma}{dt}=\frac{g^4}{8\pi s^2}\left[\frac{t}{s}+\frac{s}{s+t}+1\right]\ ,\tag{14}$$

where, as usual, $t=(p'-p)^2$, and the final energy of the neutrino is $E'=(1+t/s)E$. The total cross section is found by integrating Eq. (14) over the limits of integration $-s\leq t\leq 0$. Note that there is a logarithmic divergence at $t=-s$. This corresponds to the t-channel exchange of a massless ν. It occurs whenever the final X carries off *all* of the initial neutrino energy. Although the cross section formally diverges, there is a physical cutoff. If the X emerges with $E_X=E$, with high probability it can scatter with another background X ($XX\rightarrow\nu\nu$) to produce another neutrino with energy E. In the opposite limit, $t\rightarrow 0$, the neutrino retains all of its incident energy, and is not removed from the "detectable" flux. The relevant factor is not the total cross section, but the cross section which describes the transport of energy of the incident neutrino by scattering with low-energy particles. This requires a significant energy loss by the initial neutrino, some substantial fraction of s. We can calculate the relevant fraction of the total cross section by taking the limits of integration to be $-s(1-\epsilon)\leq t\leq -\epsilon s$. The relevant part of the cross section then is

$$\bar{\sigma}(\nu X\rightarrow\nu X)=\int_{-s(1-\epsilon)}^{-s\epsilon}dt\,(d\sigma/dt)=\frac{g^4}{16\pi s}\left[(1-2\epsilon)+2\ln\left[\frac{1-\epsilon}{\epsilon}\right]\right]\ .\tag{15}$$

If we choose $\epsilon=\frac{1}{10}$, $\bar{\sigma}(\nu X\rightarrow\nu X)=0.103g^4/s$, or following the notation of the previous section, $a=0.103g^4$.

Assuming that X is the only CBP,

$$\lambda^{-1}=\frac{ag^4}{ET}n_X I_{12}=4.5\times10^{-12}g^4\,\mathrm{cm}^{-1}\ . \tag{16}$$

The requirement $D\lambda^{-1}\lesssim1$ ($D=1.7\times10^{23}$ cm) is satisfied if $g\lesssim1.1\times10^{-3}$.

If there are background neutrinos, the incident $\bar{\nu}_e$ can scatter via $\bar{\nu}_e\bar{\nu}_e\to\bar{\nu}_e\bar{\nu}_e$, $\bar{\nu}_e\nu_e\to XX$, $\bar{\nu}_e\nu_e\to\bar{\nu}_e\nu_e$, and $\bar{\nu}_e\nu_e\to\sum_{i\neq e}\bar{\nu}_i\nu_i$. If other species of neutrinos are also present as CBP's, then $\bar{\nu}_e$ can scatter via $\bar{\nu}_e\nu_i\to\bar{\nu}_e\nu_i$, $\bar{\nu}_e\bar{\nu}_i\to\bar{\nu}_e\bar{\nu}_i$ ($i\neq e$). The differential cross sections $d\sigma/dt$, and effective cross sections, $\bar{\sigma}$, for the various processes are given in Table I. In reactions involving neutrinos other than ν_e, we have assumed that all g_i's are equal: $g_i\equiv g$.

If we assume that the present number density for all neutrinos is $n_{\nu_i}=n_{\bar{\nu}_i}=55$ cm^{-3}, the inverse mean free path is given by a sum over all the processes from Table I. Assuming two types of neutrinos other than ν_e, λ^{-1} is given by

$$\begin{aligned}\lambda^{-1}&=4.5\times10^{-12}g^4\,\mathrm{cm}^{-1}+\frac{g^4}{ET}n_\nu I_{12}(0.056+0.916+0.685+2\times0.010+4\times0.741)\\&=5.8\times10^{-11}g^4\,\mathrm{cm}^{-1}\ .\end{aligned} \tag{17}$$

The requirement $D\lambda^{-1}\lesssim1$ is satisfied only if $g\lesssim5.6\times10^{-4}$.

If neutrinos couple to a massless spin-one boson and all neutrino species are massless, then the coupling must be less than about 5.6×10^{-4}. If neutrinos are massive they might not be present as CBP's, and in that case the coupling must be less than 1.1×10^{-3}.

We have not been explicit about precisely how much the flux of energy carried by $\bar{\nu}_e$'s can decrease and still be consistent with the experimental results.[2,3] To do so would specify a definite limit to ϵ and (D/λ). However, our results are not very sensitive to the choices made for ϵ or (D/λ), because the limit to g is proportional to $(\lambda/D)^{1/4}$. For instance, if we choose $\epsilon=0.3$, then the limit is $g\lesssim6.6\times10^{-4}$. If we require $\lambda^{-1}D\lesssim3$, then the limit increases by a factor of $3^{1/4}\simeq1.32$.

B. Neutrino coupling to a massive ($M\gtrsim60$ eV) boson

If neutrinos couple to a massive spin-one boson (of mass M) through a vector coupling of the form $g_i\bar{\nu}_i\gamma_\mu\nu_i X^\mu$, the effective low-energy neutrino interactions would be described by a Lagrangian of the form

$$\mathcal{L}_I=\frac{g_i g_j}{M^2}\bar{\nu}_i\gamma_\mu\nu_i\bar{\nu}_j\gamma^\mu\nu_j\ . \tag{18}$$

In the following we will assume $g_i\equiv g$ for all i, and that all neutrino species are present as CBP's. A massive X would decay to $\nu\bar{\nu}$, and would not be present as a CBP.

The various reactions involving an incident $\bar{\nu}_e$ are given in Table II, along with the differential cross sections, and the cross section relevant in the calculation of the mean free path. Again, assuming $n_{\nu_i}=n_{\bar{\nu}_i}\simeq55$ cm^{-3}, Eq. (10) results in

TABLE I. Differential cross sections and effective cross sections for various processes involving a real or intermediate massless vector particle. The quantity $\bar{\sigma}$ is $\int(d\sigma/dt)dt$ evaluated with the limits $-s(1-\epsilon)\le t\le -s\epsilon$. Results for processes with a ν_τ are the same as those for ν_μ.

Process	$(d\sigma/dt)(8\pi s^2/g^4)$	$\bar{\sigma}s/g^4$ $\epsilon=0.1$	$\epsilon=0.3$
$\bar{\nu}_e X\to\bar{\nu}_e X$	$\left[\frac{s+t}{s}+\frac{s}{s+t}\right]$	0.103	0.031
$\bar{\nu}_e\nu_e\to XX$	$-\frac{1}{2}\left[\frac{t}{s+t}+\frac{s+t}{t}\right]$	0.056	0.010
$\bar{\nu}_e\bar{\nu}_e\to\bar{\nu}_e\bar{\nu}_e$	$1+\frac{s^2}{t^2}+\frac{s^2}{(s+t)^2}$	0.916	0.395
$\bar{\nu}_e\nu_e\to\bar{\nu}_e\nu_e$	$2\left[1+\frac{(s+t)^2}{t^2}+\frac{(s+t)^2}{s^2}\right]$	0.685	0.359
$\bar{\nu}_e\nu_e\to\bar{\nu}_\mu\nu_\mu$	$\left[\frac{t^2}{s^2}+\frac{(s+t)^2}{s^2}\right]$	0.010	0.005
$\bar{\nu}_e\nu_\mu\to\bar{\nu}_e\nu_\mu$	$\frac{s^2}{t^2}+\frac{(s+t)^2}{t^2}$	0.741	0.335

TABLE II. Same as Table I, but for a massive vector particle of mass M.

Process	$(d\sigma/dt)(8\pi M^4/g^4)$	$\bar{\sigma}M^4/g^4s$ $\epsilon=0.1$	$\epsilon=0.3$
$\bar{\nu}_e\bar{\nu}_e\to\bar{\nu}_e\bar{\nu}_e$	$\frac{1}{2}[4s^2+(s+t)^2+t^2]$	5.74×10^{-2}	2.81×10^{-2}
$\bar{\nu}_e\nu_e\to\bar{\nu}_e\nu_e$	$4(s+t)^2+s^2+t^2$	5.68×10^{-1}	3.69×10^{-2}
$\bar{\nu}_e\nu_e\to\bar{\nu}_\mu\nu_\mu$	$(s+t)^2+t^2$	1.93×10^{-2}	8.38×10^{-3}
$\bar{\nu}_e\nu_\mu\to\bar{\nu}_e\nu_\mu$	$(s+t)^2+s^2$	4.15×10^{-2}	1.62×10^{-2}

$$\lambda^{-1}=\frac{16}{5}\frac{g^4}{M^4}ETn_XI_{32}(5.74\times10^{-2}+5.68\times10^{-1}+2\times1.93\times10^{-2}+4\times4.15\times10^{-2})$$
$$=2.92\times10^{-28}\frac{g^4}{(M/\text{MeV})^4}\ \text{cm}^{-1}\ , \qquad (19)$$

using $\epsilon=0.1$. The requirement $D\lambda^{-1}\lesssim1$ implies that $g/(M/\text{MeV})\lesssim12$. If we assume that $g\approx e=0.3$, then $M\gtrsim2.5\times10^4$ eV.

C. The Majoron model

The Majoron model of Gelmini and Roncadelli[6] is a definite and well-studied[7] model in which neutrinos have secret interactions. In the model, Majorana neutrinos ν_e, ν_μ, and ν_τ couple to a massless spin-0 boson (the Majoron) with couplings g_{ii} $(i=e,\mu,\tau)$. The neutrino masses are proportional to the g_{ii}. We will assume all the g_{ii} are equal: $g_{ii}\equiv g$.

Neutrinos with masses in excess of about 10 eV would annihilate into Majorons as the temperature of the Universe drops below the mass of the neutrino.[7,11] This has two effects. It depletes the relevant neutrino species as a CBP, and increases the temperature of the cosmic background Majorons relative to the photons (in the same way that e^+e^- annihilation increases the photon temperature relative to the neutrino temperature). The Majoron temperature will depend upon the number of neutrinos with mass greater than 10 eV. To be conservative we will ignore the possible increase of the Majoron temperature, and we will take the present Majoron temperature to be $T_M=T_\nu=1.9$ K. Since the Majorons are spin 0, $n_M=\frac{1}{2}(1.9/2.7)^3n_\gamma\simeq70\ \text{cm}^{-3}$.

The differential cross section for $\bar{\nu}_eM\to\bar{\nu}_eM$ scattering is

$$\frac{d\sigma}{dt}=\frac{g^4}{32\pi s}\left[\frac{t}{s}+\frac{s}{t+s}+3\right]. \qquad (20)$$

This, of course, has the expected divergence at $t=-s$. Employing the same cutoff as before,

$$\bar{\sigma}(\bar{\nu}_eM\to\bar{\nu}_eM)=\frac{g^4}{64\pi s}\left[\ln\left[\frac{1-\epsilon}{\epsilon}\right]-(1-2\epsilon)\right]$$
$$=4.18\times10^{-2}g^4/s\quad(\epsilon=0.1)\ . \qquad (21)$$

If Majorons are the only CBP's, then

$$\lambda^{-1}=\frac{ag^4}{ET}n_MI_{12}=9.13\times10^{-13}g^4\ \text{cm}^{-1}\ . \qquad (22)$$

The requirement that $D\lambda^{-1}\lesssim1$ leads to the limit $g\lesssim1.6\times10^{-3}$.

If $m_{\bar{\nu}_e}\lesssim10$ eV, $\bar{\nu}_e$'s would survive primordial annihilation and be present as CBP's (Ref. 11). In that case, possible scattering channels for the incident $\bar{\nu}_e$ are $\bar{\nu}_e\bar{\nu}_e\to\nu_e\nu_e,\nu_\mu\nu_\mu,\nu_\tau\nu_\tau$, $\bar{\nu}_e\bar{\nu}_e$, $\bar{\nu}_\mu\bar{\nu}_\mu$, $\bar{\nu}_\tau\bar{\nu}_\tau$, and $\bar{\nu}_e\nu_e\to\nu_e\bar{\nu}_e$, MM. Using $n_{\nu_i}=n_{\bar{\nu}_i}=55\ \text{cm}^{-3}$, we find, from Table III,

$$\lambda^{-1}=9.13\times10^{-13}g^4\ \text{cm}^{-1}+6.71\times10^{-13}g^4\ \text{cm}^{-1}$$
$$=1.58\times10^{-12}g^4\ \text{cm}^{-1} \qquad (23)$$

and $D\lambda^{-1}\lesssim1$ gives $g\lesssim1.4\times10^{-3}$.

Finally, if m_{ν_μ} and m_{ν_τ} have masses less than 10 eV,

TABLE III. Same as Table I, but for the Majoron model.

Process	$(d\sigma/dt)(32\pi s^2/g^4)$	$\bar{\sigma}s/g^4$ $\epsilon=0.1$	$\epsilon=0.3$
$\bar{\nu}_eM\to\bar{\nu}_eM$	$\frac{t}{s}+\frac{s}{t+s}+3$	4.18×10^{-2}	1.82×10^{-2}
$\bar{\nu}_e\nu_e\to MM$	$-\frac{1}{2}\left[\frac{s}{t}+\frac{t}{s+t}+3\right]$	2.97×10^{-3}	2.35×10^{-4}
$\bar{\nu}_e\nu_i\to\nu_e\bar{\nu}_i$	1	7.96×10^{-3}	3.98×10^{-3}
$\bar{\nu}_e\bar{\nu}_i\to\nu_e\nu_i$	1	7.96×10^{-3}	3.98×10^{-3}
$\bar{\nu}_e\bar{\nu}_e\to\bar{\nu}_i\bar{\nu}_i$	1	7.96×10^{-3}	3.98×10^{-3}
$\bar{\nu}_e\bar{\nu}_e\to\nu_i\nu_i$	1	7.96×10^{-3}	3.98×10^{-3}

they also would be present as CBP's and the channels $\bar{\nu}_e\bar{\nu}_\mu \to \nu_e\nu_\mu$, $\bar{\nu}_e\nu_\mu \to \nu_e\bar{\nu}_\mu$, $\bar{\nu}_e\bar{\nu}_\tau \to \nu_e\nu_\tau$, and $\bar{\nu}_e\nu_\tau \to \nu_e\bar{\nu}_\tau$ are open and we find

$$\lambda^{-1} = 1.95\times 10^{-12} g^4 \text{ cm}^{-1} \tag{24}$$

and the limit $g \lesssim 1.3\times 10^{-3}$.

IV. CONCLUSIONS

Interactions of neutrinos with cosmic background particles provide a unique opportunity to constrain the interactions of neutrinos with themselves and/or other particles—limits that cannot be matched in the laboratory setting. We have considered three models. The first model was a neutrino interacting with a massless spin-one boson. In that case the "charge" of the neutrino must be less than about 10^{-3} (cf. the charge of the electron $e = 0.3$). The second model was a neutrino interacting with a massive ($M \gtrsim 60$ eV) spin-one boson. In that case $g/M \lesssim 12$ MeV^{-1}. The final model was the Majoron model.[6] In that model the coupling of the Majoron to ν_e must be less than about 10^{-3}.

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