Stellar energy loss and anomalous contributions to $\gamma\gamma \to \nu\bar{\nu}$

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We compute the stellar energy-loss rate for several anomalous contributions to $\gamma\gamma \to \nu\bar{\nu}$ and for massive neutrinos. Except for the cases where the mass scale of these new contributions is small, or the neutrino masses are large, these reactions are not important as a mechanism for stellar energy loss.

Many years ago Pontecorvo and Chiu and Morrison[1] suggested that the process $\gamma\gamma \to \nu\bar{\nu}$ might play an important role as a mechanism for stellar cooling. Gell-Mann[2] subsequently showed that this process is forbidden in a local $(V-A)$ theory. However, it can occur at the one-loop level which has been computed by Levine[3] for the intermediate-boson $(V-A)$ theory, and the stellar energy-loss rate through $\gamma\gamma \to \nu\bar{\nu}$ was found to be smaller than the rates for competing process (pair annihilation $e^+e^- \to \nu\bar{\nu}$ and photoneutrinos $\gamma e \to e\nu\bar{\nu}$). This result is not modified when the cross section of the above process is computed in the Weinberg-Salam theory as was shown by Dicus.[4]

The conditions for which the $\gamma\gamma \to \nu\bar{\nu}$ reaction is suppressed[2] are weakened if we have more exotic couplings between photons and neutrinos, one of the photons is virtual, or the neutrinos are massive. Many extensions of the standard model lead naturally to new interactions between photons and neutrinos, and there is not any fundamental principle implying that neutrinos should be massless. The existence of massive neutrinos have such far-reaching consequences that the determination and understanding of neutrino masses are major issues for particle physics. Under this point of view it is interesting to investigate whether new contributions to $\gamma\gamma \to \nu\bar{\nu}$ are astrophysically or cosmologically relevant.

New interactions between photons and neutrinos may originate in different contexts (for instance, composite models, models with right-handed neutrinos, etc.). Here, without committing to a specific model we will compute the cross sections and energy-loss rates of new contributions to $\gamma\gamma \to \nu\bar{\nu}$. We compare them to the standard-model contribution, verifying that for some specific conditions they may be larger than this one. The effective Lagrangians that we will consider are

$$\mathcal{L}_{\rm I} = \frac{1}{2\Lambda}\bar{\Psi}\sigma_{\mu\nu}\Psi F^{\mu\nu} , \tag{1a}$$

$$\mathcal{L}_{\rm II} = \frac{i}{2\Lambda^2}\bar{\Psi}\gamma_\mu \overleftrightarrow{\partial}_\nu \Psi F^{\mu\nu} , \tag{1b}$$

$$\mathcal{L}_{\rm III} = \frac{1}{2\Lambda^2} Z^\mu[(\partial_\nu \tilde{F}_{\mu\lambda})F^{\nu\lambda} + \tilde{F}^{\nu\lambda}(\partial_\nu F_{\mu\lambda})] , \tag{1c}$$

$$\mathcal{L}_{\rm IV} = \frac{1}{4\Lambda^3}\bar{\Psi}\Psi F_{\mu\nu}F^{\mu\nu} , \tag{1d}$$

$$\mathcal{L}_{\rm V} = \frac{1}{\Lambda^4}(C_L\bar{\Psi}_L\gamma^\mu\Psi_L + C_R\bar{\Psi}_R\gamma^\mu\Psi_R)\epsilon_{\mu\nu\alpha\beta}(\partial^\nu F^{\alpha\lambda})F^\beta_{\ \lambda} , \tag{1e}$$

where Ψ are neutrino fields, $F^{\mu\nu}$ the electromagnetic field strength tensor, Z^μ the weak neutral field, and Λ is the mass scale of the new interactions. Most of these have been studied in the case of charged-fermion fields.[5] $\mathcal{L}_{\rm I}$ has already been considered as a source of stellar energy loss through plasmon decay[6] $(\Gamma \to \nu\bar{\nu})$, and its coupling is constrained experimentally[6,7] as well as astrophysically.[6,8]

The effective Lagrangians (1a) and (1b) give rise to the diagrams of Fig. 1(a), Lagrangian (1c) to the diagram of Fig. 1(b), where the Z^0-neutrino coupling is the standard-model one, and (1d) and (1e) have the contribution depicted in Fig. 1(c). To compute the cross sections we will assume that the neutrinos have Dirac masses.[9] We obtain the following results:

$$\sigma_{\rm I} = \frac{\delta f^4}{4}\left[\frac{s}{m_e{}^2}\right]\left[\frac{1\ \text{GeV}}{m_e}\right]^2[3\beta - \tfrac{7}{3}\beta^3 - (1-\beta^2)^2\mathcal{L}] , \tag{2a}$$

$$\sigma_{\rm II} = \frac{\delta}{240}\left[\frac{s}{\Lambda^2}\right]^3\left[\frac{1\ \text{GeV}}{\Lambda}\right]^2(15\beta + 20\beta^3 - 23\beta^5) , \tag{2b}$$

$$\sigma_{\rm III} = 4\pi\alpha\delta\frac{\cos^2\theta_W}{\sin^2\theta_W}\left[\frac{s}{M_W{}^2}\right]^2\left[\frac{m_\nu}{\Lambda}\right]^2\left[\frac{1\ \text{GeV}}{\Lambda}\right]^2\beta , \tag{2c}$$

$$\sigma_{\rm IV} = \delta\left[\frac{s}{\Lambda^2}\right]^2\left[\frac{1\ \text{GeV}}{\Lambda}\right]^2\beta^3 , \tag{2d}$$

$$\sigma_{\rm V} = \delta\left[\frac{s}{\Lambda^2}\right]^2\left[\frac{m_\nu}{\Lambda}\right]^2\left[\frac{1\ \text{GeV}}{\Lambda}\right]^2\beta , \tag{2e}$$

where

$$\delta = \frac{1}{64\pi}\left[\frac{\hbar c}{1\ \text{GeV}}\right]^2 (\text{cm}^2) ,$$

$$\beta^2 = \left[1 - \frac{4m_\nu{}^2}{s}\right], \quad \mathcal{L} = \ln\left[\frac{1+\beta}{1-\beta}\right] ,$$

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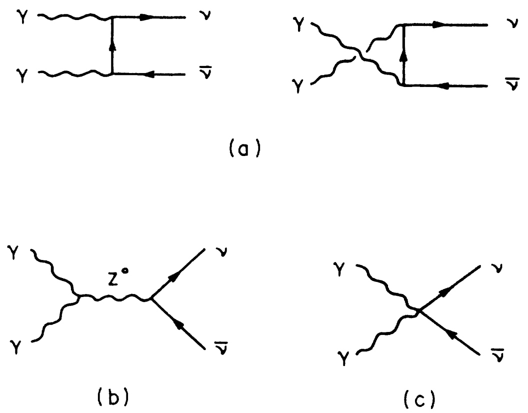

FIG. 1. (a) Diagrams contributing to $\gamma\gamma \to \nu\bar{\nu}$ coming out from the effective interactions (1a) and (1b); (b) the same for the interaction (1c); (c) the same for (1d) and (1e).

and m_ν is the neutrino mass. In (1a) we have replaced (Λ^{-1}) by $(f/2m_e)$, where f is the neutrino magnetic moment in units of electron Bohr magnetons, and in (1e) we assumed, for simplicity, $C_R = 0$ and $C_L = 1$.

The rate at which the energy of the photons (in thermal equilibrium) is converted into neutrino pairs is given by

$$Q = \frac{4c}{(2\pi\hbar c)^6} \int \frac{d^3k_1 d^3k_2 (\omega_1+\omega_2)\,|\mathbf{v}_{\rm rel}|\,\sigma(s)}{\left[\exp\left[\frac{\omega_1}{kT}\right]-1\right]\left[\exp\left[\frac{\omega_2}{kT}\right]-1\right]}, \tag{3}$$

where ω_1 and ω_2 are the photon energies, $\mathbf{k}_i$ are their momenta, $|\mathbf{v}_{\rm rel}|$ is the relative velocity factor, and the cross sections $[\sigma(s)]$ given by (2a)–(2e) can be expressed in terms of ω_1, ω_2, and θ, where θ is the angle between the photons.

For each of the cross sections (2a)–(2e), after performing the angular integrations, and defining the variables

$$\omega_1 = m_\nu x, \quad \omega_2 = m_\nu y,$$

$$\gamma = \sqrt{xy}, \quad \eta = \frac{2m_\nu c^2}{kT},$$

and

$$f(x,y,\eta) \equiv \frac{x+y}{\left[\exp\left[\frac{\eta x}{2}\right]-1\right]\left[\exp\left[\frac{\eta y}{2}\right]-1\right]},$$

we obtain

$$Q_{\rm I} = Kf^4 \left[\frac{m_\nu}{m_e}\right]^4 \eta^5 T_9{}^5 J_{\rm I}(\eta), \tag{4a}$$

$$Q_{\rm II} = K \left[\frac{m_\nu}{\Lambda}\right]^8 \eta^5 T_9{}^5 J_{\rm II}(\eta), \tag{4b}$$

$$Q_{\rm III} = 4\pi\alpha K \frac{\cos^2\theta_W}{\sin^2\theta_W} \left[\frac{m_\nu}{\Lambda}\right]^4 \left[\frac{m_\nu}{M_W}\right]^4 \eta^5 T_9{}^5 J_{\rm III}(\eta), \tag{4c}$$

$$Q_{\rm IV} = K \left[\frac{m_\nu}{\Lambda}\right]^6 \eta^5 T_9{}^5 J_{\rm IV}(\eta), \tag{4d}$$

$$Q_{\rm V} = K \left[\frac{m_\nu}{\Lambda}\right]^8 \eta^5 T_9{}^5 J_{\rm V}(\eta), \tag{4e}$$

where

$$J_i(\eta) = \tfrac{1}{16} \int_0^\infty dx \int_{1/x}^\infty dy\, f(x,y,\eta) W_i(\gamma),$$

$$W_{\rm I}(\gamma) = \frac{\gamma\sqrt{\gamma^2-1}}{9}(2\gamma^4+10\gamma^2+3) + \tfrac{1}{3}(1-6\gamma^2)\ln(\gamma+\sqrt{\gamma^2-1}),$$

$$W_{\rm II}(\gamma) = \frac{\gamma\sqrt{\gamma^2-1}}{300}(192\gamma^8+490\gamma^6-728\gamma^4+10\gamma^2+15) + \tfrac{1}{20}\ln(\gamma+\sqrt{\gamma^2-1}),$$

$$W_{\rm III}(\gamma) = \frac{\gamma\sqrt{\gamma^2-1}}{12}(48\gamma^6-8\gamma^4-10\gamma^2-15) - \tfrac{5}{4}\ln(\gamma+\sqrt{\gamma^2-1}),$$

$$W_{\rm IV}(\gamma) = \frac{\gamma\sqrt{\gamma^2-1}}{4}(16\gamma^6-24\gamma^4+2\gamma^2+3) - \tfrac{3}{4}\ln(\gamma+\sqrt{\gamma^2-1}),$$

$$W_{\rm V}(\gamma) = W_{\rm III}(\gamma),$$

$$K = \frac{\delta}{\pi^4}\frac{c}{(\hbar c)^6}(kT_0)^5 \ (\text{erg/cm}^3\,\text{sec}),$$

$$T_9 = \frac{T}{T_0}, \quad T_0 = 10^9\ \text{K}.$$

These equations, (4a)–(4e), can be integrated analytically with the approximation $2m_\nu c^2 < kT$ (or, as we did too, numerically), obtaining

$$Q_{\rm I} = \tfrac{8}{9}Kf^4 \left[\frac{kT_0}{m_e c^2}\right]^4 T_9{}^9 \Gamma(4)\Gamma(5)\zeta(4)\zeta(5), \tag{5a}$$

$$Q_{\rm II} = \tfrac{64}{25}K \left[\frac{kT_0}{\Lambda c^2}\right]^8 T_9{}^{13} \Gamma(6)\Gamma(7)\zeta(6)\zeta(7), \tag{5b}$$

$$Q_{\rm III} = 64\pi\alpha K \frac{\cos^2\theta_W}{\sin^2\theta_W} \left[\frac{kT_0}{\Lambda c^2}\right]^4 \left[\frac{kT_0}{M_W c^2}\right]^2 \left[\frac{m_\nu}{M_W}\right]^2 \times T_9{}^{11} \Gamma(5)\Gamma(6)\zeta(5)\zeta(6), \tag{5c}$$

$$Q_{\rm IV} = 16K \left[\frac{kT_0}{\Lambda c^2}\right]^6 T_9{}^{11} \Gamma(5)\Gamma(6)\zeta(5)\zeta(6), \tag{5d}$$

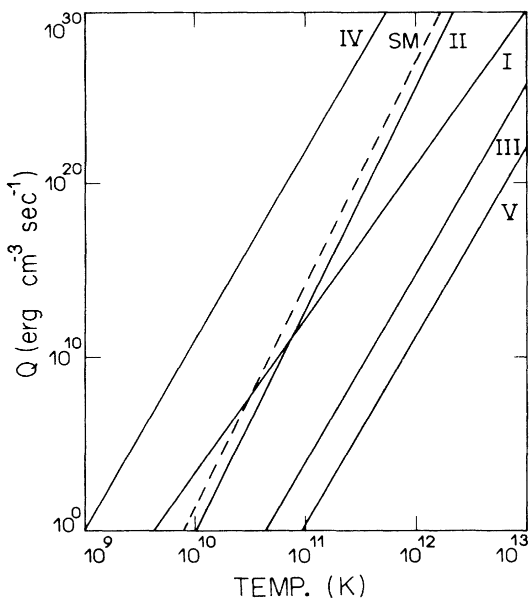

FIG. 2. Results of the numerical integration of Eqs. (4a)–(4e) for $m_\nu = 10$ eV.

$$Q_V = 16K \left[\frac{kT_0}{\Lambda c^2}\right]^6 \left[\frac{m_\nu}{\Lambda}\right]^2 T_9{}^{11}\Gamma(5)\Gamma(6)\zeta(5)\zeta(6) . \tag{5e}$$

Our results are shown in Fig. 2, where we used $m_\nu = 10$ eV, $f = 2.0\times10^{-11}$ which comes out from a superstring-inspired model,[10] $\Lambda = 1$ TeV, $\sin^2\theta_W = 0.226$, and $M_W = 83$ GeV. We have drawn in Fig. 2 the energy-loss rate for $\gamma\gamma\to\nu\bar{\nu}$ of the standard model, Q_{SM} (Ref. 11):

$$Q_{SM} \approx 2\times10^{-12} T_9{}^{13}\ (\text{erg/cm}^3\,\text{sec}) . \tag{6}$$

As noticed before,[3] the reaction $\gamma\gamma\to\nu\bar{\nu}$ can be important only at cosmological temperatures. The interaction (1d) gives an energy output some orders of magnitude larger than the standard-model one. All of these contributions are smaller than the one of $e^+e^-\to\nu\bar{\nu}$ at the same temperature.

In some of the cases we have studied there are no reasons to have a small value of Q, other than the small cross section due to the suppression of many powers in $(1/\Lambda)$, and we have been conservative when we adopted $\Lambda = 1$ TeV. A smaller value of Λ, as is possible in some composite models, would cause a substantial modification of our results. We shall also have another considerable enhancement in the emitted energy if we have larger neutrino masses[12] [see Eqs. (5c) and (5e)].

Comparing the energy output of $\gamma\gamma\to\nu\bar{\nu}$ with the plasmon decay one $(\Gamma\to\nu\bar{\nu})$ (Ref. 13), we see that the first one wins from the second in the temperature range of 10^8–10^{13} K and densities (ρ/μ_e) up to 10^6 g/cm^3, for larger densities (10^6–10^9 g/cm^3) the plasmon decay takes over up to temperatures $\sim10^{11}$ K. The photons are quite sensible to the plasma effect at large densities; introducing these effects into the computation of Q for $\gamma\gamma\to\nu\bar{\nu}$, the temperature dependence will be modified and we expect that it may lead to larger neutrino luminosities. We are currently studying this possibility.

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