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The Salam–Weinberg model of weak interactions and stellar energy loss rates by neutrino processes

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Abstract. The energy loss mechanism from stellar bodies through neutrino emission processes has been discussed in the framework of the Salam–Weinberg unified symmetric model of weak and electromagnetic interactions. As a test of the model we have considered the effect of neutral currents in $\nu\bar{\nu}$ pair production for the process $\gamma+\mathrm{p}\to\mathrm{p}+\nu+\bar{\nu}$; it is shown that although the cross section for this process at high energies (total centre of mass energy > 2 BeV) is dependent on the structure of the vertices $\langle p'|\mathscr{J}_\mu^{\mathrm{Z}}|p\rangle$ and $\langle p'|J_\mu^{\mathrm{A}}|p\rangle$, where $\mathscr{J}_\mu^{\mathrm{Z}}$ and J_μ^{A} are the weak neutral and electromagnetic currents respectively, the energy loss rate is not governed by them. The present calculation for the rate of energy loss by the γp reaction is in agreement with the rough estimate given by Dicus. However, it is an order of magnitude smaller than that calculated by Desrosiers and O'Donnell using local interactions of hadrons with leptons through weak neutral currents.

1. Introduction

The rate of energy loss from stellar bodies has been studied by many authors using various forms of weak interaction theories. The fact that the interactions of neutrinos with matter are insignificant at all energies has been made as an interesting starting point of all these investigations. The neutrino cross section is of the order of $10^{-44}x^2$ cm^2, where x is the neutrino momentum in MeV/c, which becomes of the order of 10^{-38} cm^2 for neutrinos of momentum 1 BeV/c (Chiu 1966). This implies that the neutrino has a mean free path of the order of 10^{22} cm in matter; thus, once a neutrino is formed in the interior of a star which has a radius of about 10^{13} cm, it escapes from the star carrying the energy away.

There are various neutrino production processes which lead to loss of energy from stellar bodies. For example, when the temperature of the stellar interior becomes of the order of 6×10^9 K $= m_\mathrm{e}c^2/k$ (m_e = mass of the electron, k = Boltzmann's constant) electron pairs are produced in equilibrium with the radiation. These pairs annihilate into photons in most cases; only in one out of 10^{22} annihilations is a neutrino pair $\nu_\mathrm{e}\bar{\nu}_\mathrm{e}$ created. Notwithstanding this fact the pair annihilation process $\mathrm{e}^++\mathrm{e}^-\to\nu_\mathrm{e}+\bar{\nu}_\mathrm{e}$ plays a major role in the dissipation of stellar energy, when the temperature and the density of the plasma are of the order of 10^{10} K and 10^7 g cm^{-3} respectively. This has been shown in detail by Chiu and Morrison (1960), Chiu and Stabler (1961), Chiu (1961) and Beaudet *et al* (1967a) in the point-interaction theory of weak hamiltonians (Sudershan and Marshak 1958, Feynman and Gell-Mann 1958). Other important energy loss mechanisms through neutrino production processes, eg, $\gamma+\mathrm{e}\to\mathrm{e}+\nu+\bar{\nu}$ (photoneutrino of electron) (Pontecorvo 1959, Chiu and Stabler 1961, Ritus 1962, Beaudet *et al* 1967a, 1967b), and plasmon $\rightarrow \nu+\bar{\nu}$, have also been discussed by a number of authors. The plasma–neutrino reaction in the ordinary point-interaction theory involves a virtual (e^-e^+) pair, which in turn produces a $\nu\bar{\nu}$ pair. That this is the case can be seen by considering the propagation of a photon through a degenerate electron gas. In such a gas the states up to the Fermi energy are all filled. Thus the propagation of the photon through this Fermi vacuum will modify its mode of propagation. This effect is taken into account by considering the propagation of the photon through a dielectric medium (Adams *et al* 1963, Zaidi 1965, Tsytovich 1961, 1964) and its subsequent formation of a $\nu\bar{\nu}$ pair.

Some other less important processes are perhaps the $\nu\bar{\nu}$ pair production through the reactions: (*a*) $\gamma+p \rightarrow p+\nu+\bar{\nu}$; and (*b*) $\gamma+\gamma \rightarrow \nu+\bar{\nu}$. Desrosiers and O'Donnell (1970) have calculated the energy loss rate for process (*a*) in the ordinary weak interaction theory by introducing weak neutral currents. In principle, process (*b*) should be very effective in transferring the interior energy of stellar bodies away; for, the initial radiation energy confined in a star, because of the small mean free path of the photon, is passed on to the carriers for which the star may be almost transparent. This process is, however, prohibited in the lowest order of the ordinary weak interaction theory in the current–current picture (Gell-Mann 1961). In the intermediate vector boson theory (Lee and Yang 1960) the process, on the other hand, can proceed in the lowest order (Levine 1967, Campbell 1968, Abak 1971). It should be mentioned here that one of the serious defects of the conventional weak interaction theory is that all higher order diagrams are divergent (Ioffe 1967, Ioffe and Shabalin 1967, 1968) and these divergences cannot be removed systematically due to the nonrenormalizability of the theory. To avoid this difficulty various renormalizable weak interaction theories have also been proposed (Kummer and Segrè 1965, Christ 1968). In these models it is assumed that weak interactions are mediated by a number of scalar bosons. The process $\gamma+\gamma \rightarrow \nu+\bar{\nu}$ has been calculated (Biswas *et al* 1973) in the scalar boson exchange model of weak interactions and it has been pointed out that the rate of energy loss due to this process shows no appreciable difference in temperature dependence from that due to the vector boson theory. Although the intermediate scalar boson theory is renormalizable, nevertheless, most of the calculations are beset with many masses of the new types of particles needed in the renormalizable model.

Recently, it has been shown by Salam (1968) and Weinberg (1967, 1971) that a unified renormalizable theory of weak and electromagnetic interactions may be constructed from a theory of Yang and Mills (1954) with spontaneous breaking of gauge invariance. Among the many consequences of the theory the one which is of direct relevance to us is the existence of a neutral intermediate vector boson (Z boson), coupled weakly to a neutral current. The effects of this neutral vector boson in various semi-leptonic processes have been calculated by Weinberg (1972). He has also succeeded in obtaining the weak neutral current form factors in terms of the weak charged current form factors and electromagnetic form factors. Unfortunately, the present data neither confirm nor refute the results of his calculations. In view of this, in the present paper we consider the effect of the neutral weak currents in the $\nu\bar{\nu}$ pair production processes in astrophysics. It is to be noted that the reaction $\gamma+p \rightarrow p+\nu+\bar{\nu}$ receives contributions only from the neutral currents in the lowest order, and therefore, one may hope that the study of this process may supply some useful information regarding the existence of weak neutral currents. The important energy loss processes like $e^++e^- \rightarrow \nu+\bar{\nu}$, and $\gamma+e \rightarrow e+\nu+\bar{\nu}$, however, receive contributions from both neutral and charged currents in this model (Dicus 1972). For the production process $\gamma+\gamma \rightarrow \nu+\bar{\nu}$ we only note that in addition to charged W boson contributions we have also the contributions from the neutral Z boson. It is interesting to point out that due to gauge invariance requirements the Z boson contributions vanish identically and we are left only with the W boson contributions. As shown by Levine (1967) this calculation leads to a convergent result in the lowest order since the divergences due to charged W boson contributions cancel among themselves. Regarding the convergence criterion (Weinberg 1971, t'Hooft 1971, Lee B W 1972) of the various results in this renormalizable model we may mention here that a number of authors have calculated higher order diagrams for different processes and shown that the divergences cancel out exactly if the renormalization procedure is followed systematically (Lee S Y 1972).

This symmetric model of weak and electromagnetic interactions has been applied in the present paper to calculate the energy loss rate for the neutrino emission process $\gamma + \mathrm{p} \rightarrow \mathrm{p} + \nu + \bar{\nu}$. Dicus (1972) has also made a rough estimate of this process in the Salam–Weinberg model. He has assumed that the proton couples to the Z boson and the photon with the same coupling constants as the electron, and has thus neglected the various other form factors occurring in the γpp and Zpp vertices. In particular, he has neglected the contribution from the magnetic moment of the proton to its interactions with the electromagnetic current. Further, the momentum dependent structures of the form factors have not been taken into account in his calculations. On the other hand, we take into account the various form factors associated with the matrix elements of $\mathscr{J}_\mu^Z$ and J_μ^A between the two proton states (Weinberg 1972). For the momentum dependence of the form factors we assume Sachs type structures which decrease rapidly with the square of the momentum transfer of the protons. These structures are important in determining the values of the cross section when the total centre of mass (CM) energy is greater than 2 BeV. At low energies, however, the cross section varies as the fourth power of the CM energy of the photon and is independent of the momentum dependent structures of the form factors. Next we calculate the rate of loss of energy from the square of the matrix element multiplied by the energy taken away by the neutrino pair, by using the number densities of the particles given by Bose and Fermi statistics. In doing so we note that the contributions from the high energies are extremely small at relevant temperatures and mass densities of the stellar bodies, and so we are allowed to calculate it in the nonrelativistic limit. In this limit the momentum dependent structures of the form factors are not very relevant. We find that the energy loss rate is quite small compared to other dominant energy emission processes, and therefore, will not play an important role in the evolution of stars. This model predicts a value lower than the previously estimated value (Desrosiers and O'Donnell 1970) in the ordinary point-interaction theory with a neutral current by a factor of five. However, our results are in agreement with the estimates given by Dicus (1972). Furthermore, the process $\gamma + \mathrm{p} \rightarrow \mathrm{p} + \nu + \bar{\nu}$ dominates over the other known minor processes (Levine 1967, Hieu and Shabalin 1963) and so greater refinement of stellar evolution models may find the difference between this theory and the point-interaction theory with neutral currents.

In the next section we write down the lagrangian of the symmetric unified model of weak and electromagnetic interactions, and apply it to calculate $\nu\bar{\nu}$ pair production from the photoproduction process of the proton. In § 3 we calculate the cross section for this process by assuming that the protons form a nondegenerate Fermi gas. We have shown that this assumption is justifiable for the relevant temperatures and mass densities of the plasma. In § 4 we obtain the expression for the rate of loss of energy in the nonrelativistic limit. Finally, we discuss briefly the energy loss mechanisms due to other processes in the last section.

2. The process $\gamma+\mathrm{p} \rightarrow \mathrm{p}+\nu+\bar{\nu}$ in the neutral current model of Salam and Weinberg

The part of the interaction lagrangian of the renormalizable theory of Salam (1968) and Weinberg (1967, 1971) which is relevant for the description of the photoneutrino process

$$\gamma(k)+\mathrm{p}(p) \rightarrow \mathrm{p}(p')+\nu(q)+\bar{\nu}(q') \tag{1}$$

is given by

$$\mathscr{L}_{\mathrm{I}} = \tfrac{1}{4}\mathrm{i}(g^2+g'^2)^{1/2}\{\bar{\nu}_{\mathrm{e}}\gamma_\lambda(1+\gamma_5)\nu_{\mathrm{e}}+(\mathrm{e}\rightarrow\mu)\}Z_\lambda+\tfrac{1}{2}(g^2+g'^2)^{1/2}Z_\lambda\mathscr{J}^{\mathrm{Z}}_\lambda-eA_\lambda J^{\mathrm{A}}_\lambda \tag{2}$$

where J^{A}_λ and $\mathscr{J}^{\mathrm{Z}}_\lambda$ are the electromagnetic and weak neutral currents, respectively. The pair of independent coupling constants g and g' occurring in (2) are related to the weak coupling constant G_β, the electric charge e, and the mass m_{Z} of the neutral boson by the following expressions:

$$G_\beta = (g^2+g'^2)/8m_{\mathrm{Z}}^2 \qquad \text{and} \qquad e = gg'(g^2+g'^2)^{-1/2}. \tag{3}$$

In order to evaluate the matrix element for the process (1) in this model we note that the weak and electromagnetic form factors of protons are required (figure 1). We

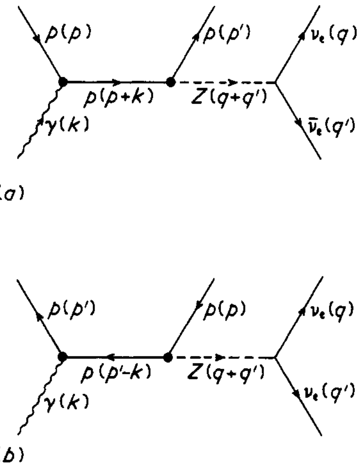

Figure 1. The lowest order Feynman diagrams for the process $\gamma+\mathrm{p}\rightarrow\mathrm{p}+\nu_{\mathrm{e}}+\bar{\nu}_{\mathrm{e}}$. The symbols in parentheses denote the momenta of the particles.

first define the following vertices (Weinberg 1972) which occur in our matrix element:

$$\langle p(p')|\mathscr{J}^{\mathrm{Z}}_\lambda|p(p)\rangle = \bar{u}(p')\{g^0_{\mathrm{V}}\gamma_\lambda+g^0_{\mathrm{A}}\gamma_\lambda\gamma_5+\mathrm{i}f^0_{\mathrm{V}}(p+p')_\lambda+\mathrm{i}h^0_{\mathrm{A}}(p-p')_\lambda\gamma_5\}u(p) \tag{4}$$

and

$$\langle p(p')|J^{\mathrm{A}}_\lambda|p(p)\rangle = \bar{u}(p')\{G^{\mathrm{p}}\gamma_\lambda+\mathrm{i}F^{\mathrm{p}}(p+p')_\lambda\}u(p) \tag{5}$$

where g^0_{V}, g^0_{A}, f^0_{V}, h^0_{A}, G^{p}, and F^{p} are real dimensionless functions of the invariant squared momentum transfer $(p'-p)^2$.

It should be mentioned that the process (1) actually involves two reactions, one involving a pair of electron neutrinos and the other, a pair of muon neutrinos. We must consider both to obtain the total rate of loss of energy. Using the lagrangian (2) and the vertex functions (4) and (5) we see that the sum of the matrix elements for the diagrams (1*a*) and (1*b*) is of the following form:

$$T_{fi} = \frac{\mathrm{i}\bar{N}(2\pi)^4\delta^{(4)}(p+k-p'-q-q')}{2(p\,.\,k)(p'\,.\,k)(m_Z^2+2q\,.\,q')}\epsilon_\mu(k)X_{\mu\nu}\{\bar{u}(q)\gamma_\nu(1+\gamma_5)v(q')\} \tag{6}$$

where $\epsilon_\mu(k)$ is the photon polarization,

$$\bar{N} = \frac{e(g^2+g'^2)}{8}\left(\frac{m_p^2 m_\nu^2}{2E_pE_kE_{p'}E_qE_{q'}}\right)^{1/2} \tag{7}$$

and

$$X_{\mu\nu} = \bar{u}(p')[(g_V^0\gamma_\nu+g_A^0\gamma_\nu\gamma_5+2\mathrm{i}f_V^0p'_\nu)\{\gamma\,.\,k(G^p(0)\gamma_\mu+2\mathrm{i}F^p(0)p_\mu)+2p_\mu\}p'\,.\,k$$
$$+\{(G^p(0)\gamma_\mu+2\mathrm{i}F^p(0)p'_\mu)\gamma\,.\,k-2p'_\mu\}(g_V^0\gamma_\nu+g_A^0\gamma_\nu\gamma_5+2\mathrm{i}f_V^0p_\nu)p\,.\,k]u(p). \tag{8}$$

In deriving (6) we have considered only one type of neutrino; the gauge condition $\epsilon\,.\,k = 0$ is used, and plasma effects are neglected.

Since the cross section as well as the energy loss rate involves the square of the matrix element, we square (6), average over the initial spin and polarization states, and sum over the final spin states. The result of our calculations is finally given by

$$\overline{\sum_i}\sum_f\frac{|T_{fi}|^2}{T} = \frac{e^2(g^2+g'^2)^2(2\pi)^4\delta^{(4)}(p+k-p'-q-q')}{2^{12}E_pE_kE_{p'}E_qE_{q'}(p\,.\,k)^2(p'\,.\,k)^2(m_Z^2+2q\,.\,q')^2}\sum_{i=1}^{11}C_iT_i \tag{9}$$

where $C_i \equiv C_i(p, p', k)$ and $T_i \equiv T_i(p, p', k, q, q')$.

The expressions for the scalars T_i and C_i are given in the appendix.

3. The chemical potential of a proton and the cross section for the photoneutrino process of protons

We will assume that the star consists of a completely ionized gas in thermal equilibrium at a temperature T K. The number densities of protons and photons are as usual given by Fermi–Dirac and Bose distributions respectively:

$$n_p = \int \mathrm{d}n_p = \frac{2}{(2\pi)^3}\int\frac{\mathrm{d}^3p}{\exp\{(E_p-\mu)/kT\}+1} \tag{10}$$

$$n_\gamma = \int \mathrm{d}n_\gamma = \frac{2}{(2\pi)^3}\int\frac{\mathrm{d}^3k}{\exp(E_k/kT)-1}, \tag{11}$$

where μ is the chemical potential of a proton including its rest mass. If we neglect the mass of the electrons and positrons then the matter density ρ of the plasma is given by

$$n_p = N\rho/\mu_e \tag{12}$$

where N is Avogadro's number and μ_e is related to the abundance X_j, the nuclear charge Z_j, and the atomic weight A_j of the jth atomic species in the stellar body by

$$\frac{1}{\mu_e} = \sum_j\frac{X_jZ_j}{A_j}. \tag{13}$$

We evaluate the integral (10) for various values of μ and T and show the relationship between ρ/μ_e and μ for different temperatures in figure 2. It should be noted that in the

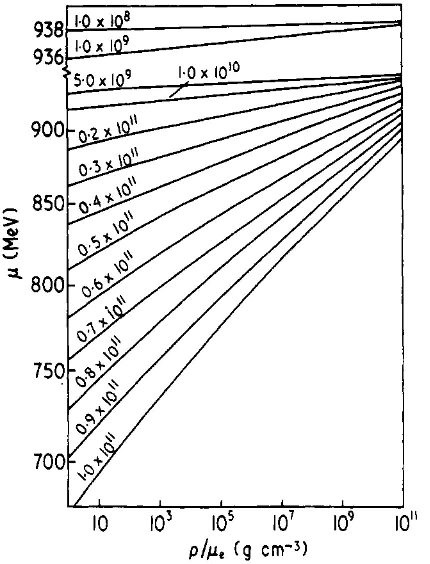

Figure 2. The chemical potential μ (MeV) of a proton as a function of ρ/μ_e (g cm^{-3}) for different temperatures T(K).

usual ranges of temperatures and mass densities of the plasma, the contributions to the integral (10) from the high momentum states are extremely small and therefore one may use the nonrelativistic limit in the calculation of energy loss rate where the distribution functions for the initial and the final particles are occurring.

The number densities for the outgoing particles are described by

$$n_{p'} = \int \mathrm{d}n_{p'} = \frac{2}{(2\pi)^3}\int \mathrm{d}^3p'\left(1-\frac{1}{\exp\{(E_{p'}-\mu)/kT\}+1}\right) \tag{14}$$

$$n_q = \int \mathrm{d}n_q = \frac{1}{(2\pi)^3}\int \mathrm{d}^3q \tag{15}$$

$$n_{q'} = \int \mathrm{d}n_{q'} = \frac{1}{(2\pi)^3}\int \mathrm{d}^3q'. \tag{16}$$

We note that the factor $1/[\exp\{(E_{p'}-\mu)/kT\}+1]$ in (14) may be neglected in comparison to 1 for densities and temperatures prevailing in stars (figure 2). In other words we may assume that the protons in the stars form a nondegenerate Fermi gas.

Using (9) we have the total cross section of the photoneutrino process (1) for nondegenerate protons:

$$\sigma = \frac{e^2(g^2+g'^2)^2}{(8\pi)^5E_pE_k(p\,.\,k)^2}\int\frac{\mathrm{d}^3p'}{E_{p'}(p'\,.\,k)^2}\frac{1}{\{m_z^2+(p+k-p')^2\}^2}$$
$$\times\iint\frac{\mathrm{d}^3q\,\mathrm{d}^3q'}{4E_qE_{q'}}\delta^{(4)}(p+k-p'-q-q')\sum_{i=1}^{11}C_iT_i. \tag{17}$$

Making use of Lenard's identity (Lenard 1953)

$$\iint\frac{\mathrm{d}^3q\,\mathrm{d}^3q'}{4E_qE_{q'}}\delta^{(4)}(\bar{p}-q-q')q_\mu q'_\nu = \frac{\pi}{24}(\bar{p}^2\delta_{\mu\nu}+2\bar{p}_\mu\bar{p}_\nu) \tag{18}$$

we integrate over the neutrino momentum states and note that the contributions from $C_i T_i$ for $i = 7$ to 11 vanish identically. The p' integration is done in the CM system of the initial particles and the cross section σ_{CM} is evaluated numerically for different values of the CM energy of the photon (figure 3). For the different weak and electromagnetic

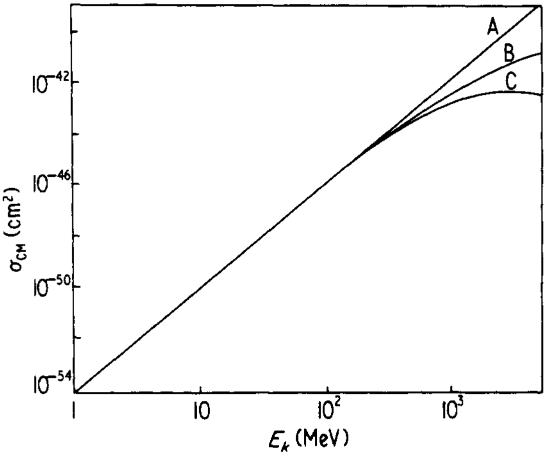

Figure 3. The cross section σ_{CM} (cm^2) for the process $\gamma + p \rightarrow p + \nu_e + \bar{\nu}_e$ as a function of the centre of mass energy E_k (MeV) of the photon for three different values of the form factors $\mathscr{H}_V(q^2) = \mathscr{H}_A(q^2)$; curves A, B and C respectively correspond to $\mathscr{H}_V(q^2) = 1$, $\mathscr{H}_V(q^2) = 1/(1-q^2/m_x^2)$ and $\mathscr{H}_V(q^2) = 1/(1-q^2/m_x^2)^2$ where $m_x = 1$ BeV.

couplings of protons we have used the following values of the form factors (Weinberg 1972):

$$G^p(0) = 2{\cdot}79, \qquad F^p(0) = 1{\cdot}79/2m_N$$
$$g_V^0(q^2) = (2{\cdot}35 - 5{\cdot}58 \sin^2\theta)\mathscr{H}_V(q^2)$$
$$f_V^0(q^2) = (1{\cdot}85 - 3{\cdot}58 \sin^2\theta)\mathscr{H}_V(q^2)/2m_N, \tag{19}$$

and

$$g_A^0(q^2) = 0{\cdot}6\mathscr{H}_A(q^2)$$

where θ is the mixing angle defined by

$$g'/g \equiv \tan\theta \tag{20}$$

and m_N is the nucleon mass.

For the q^2 dependence of the vector and axial-vector form factors we have taken the following two forms:

(i) $$\mathscr{H}_V(q^2) = \mathscr{H}_A(q^2) = 1/(1-q^2/m_x^2) \tag{21a}$$

and

(ii) $$\mathscr{H}_V(q^2) = \mathscr{H}_A(q^2) = 1/(1-q^2/m_x^2)^2 \tag{21b}$$

where m_x is a mass parameter. For the usual Sachs' dipole fit (21b) the experiments on electron–proton scattering indicate that $m_x = 0{\cdot}71$ BeV for low values of q^2 (Gasiorowicz 1966). In our calculations, however, we set $m_x = 1$ BeV for both the forms (21a) and (21b)

and find that the cross section is very well described by $\mathscr{H}_V(0)$ and $\mathscr{H}_A(0)$ at low photon energies (figure 3).

For the value of the mixing angle θ we look into the process $\bar{\nu}_e + e^- \rightarrow \bar{\nu}_e + e^-$ and $\nu + p \rightarrow \nu + p$ in Weinberg's model and find that the present experimental evidences indicate $\theta \simeq 30°$ (Weinberg 1972). In the Salam–Weinberg model the lowest value of the mass of the intermediate charged vector boson is $m_W \simeq 37{\cdot}3$ BeV (Weinberg 1971). However, using different models originating from the same gauge symmetry, Schechter and Ueda (1970) and Lee (1971) have shown that m_W comes out to be precisely 37·3 BeV. This value of m_W in turn gives an estimate of m_Z through the relation

$$m_Z^2 = (g^2 + g'^2)m_W^2/g^2.$$

About the cross section in the CM system we may point out that σ_{CM} can be described by the empirical relation

$$\sigma_{CM} \simeq 9{\cdot}32 \times 10^{-55}\, E_k^4 \qquad \text{(in cm}^2\text{)} \tag{22}$$

in the nonrelativistic region. In (22) E_k, the CM energy of the photon, is measured in MeV. We may point out that the value of m_W chosen here is not crucial; we have varied the value of m_W over a large range, including $m_W \simeq 53$ BeV obtained by Schwinger (1973) within gauge models, and found that the result does not change significantly when m_W is sufficiently large, say greater than 10 BeV, and E_k is less than 1 BeV. It may be mentioned that the cross section for the photoneutrino process of electrons has the same energy dependence (E_k^4) in the nonrelativistic region for a nondegenerate electron gas.

4. Energy loss rate

The energy loss rate due to the process $\gamma + p \rightarrow p + \nu_e + \bar{\nu}_e$ is

$$Q = \int dn_p \int dn_k \int (E_p + E_k - E_{p'})\, dn_{p'} \iint dn_q\, dn_{q'} \overline{\sum_i} \sum_f \frac{|T_{fi}|^2}{T} \tag{23}$$

where $\overline{\Sigma}_i \Sigma_f |T_{fi}|^2/T$ is given by equation (9).

As explained earlier we will confine ourselves to the nonrelativistic and nondegenerate case; and hence we may simplify (23) to

$$Q = \int dn_p \int dn_k E_k \sigma. \tag{24}$$

Now using equations (10), (11), (12), and (22) we evaluate the integrals in equation (24) to find

$$\frac{Q}{\rho} \simeq \frac{5{\cdot}51}{\mu_e} T_9^8 \qquad \text{(in erg g}^{-1}\text{ s}^{-1}\text{)}, \qquad (T_9 = T/10^9) \tag{25}$$

which is about one-fifth of the value, $28{\cdot}8\ T_9^8/\mu_e$ erg g^{-1} s^{-1}, given by Desrosiers and O'Donnell (1970) in the point interaction theory with a neutral current. Our result for the energy loss, however, agrees with the value estimated by Dicus (1972) from the expression of the energy loss rate due to the process $\gamma + e \rightarrow e + \nu + \bar{\nu}$ by suitably replacing the mass and the coupling constants.

5. Discussion

Most astrophysically important reactions have a very sharp dependence on temperature above 10^9 K. In a comparison of reactions, therefore, the question of the exact power of T_9 to use is crucial. The present calculation shows that the loss of energy due to the photoneutrino process of protons depends on T_9^8 which is typical of the more important reactions in neutrino astrophysics involving electrons. Although it has almost the same temperature dependence as the other important neutrino emission processes, this process is negligible compared to others because of the heavy mass of the proton compared to that of the electron. This reaction, however, has a dominant energy loss rate over the other minor sources of neutrino emission. As mentioned earlier the photon–photon scattering ($\gamma+\gamma \rightarrow \nu+\bar{\nu}$) rate calculated in this model is exactly the same as in the charged intermediate vector boson theory and in the latter model Levine (1967) has shown that the energy loss rate is unimportant for all temperatures below 10^{13} K. Hieu and Shabalin (1963) have found that the energy loss rate in other similar processes like $\gamma+\gamma \rightarrow \gamma+\nu+\bar{\nu}$ is also unimportant.

There is another point of view from which the reaction $\gamma+\mathrm{p} \rightarrow \mathrm{p}+\nu+\bar{\nu}$ should be studied. At present there are various models of weak interactions. However, contrary to the normal theories of weak interactions Weinberg's model predicts the existence of neutral currents which contain the neutrino term $\bar{\nu}\gamma_\lambda(1+\gamma_5)\nu$. Certainly, astrophysical neutrino processes can serve as an interesting probe to determine the existence of these neutral currents.

Acknowledgments

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Appendix

The expressions for the T_i and C_i are given by

$$T_1 = -2q\,.\,q'$$

$$T_2 = 2p\,.\,q\;p\,.\,q' - p^2\;q\,.\,q'$$

$$T_3 = 2(p\,.\,q\;p'\,.\,q' + p\,.\,q'\;p'\,.\,q - p\,.\,p'\;q\,.\,q')$$

$$T_4 = 2p'\,.\,q\;p'\,.\,q' - p^2\;q\,.\,q'$$

$$T_5 = 2(k\,.\,q\;p\,.\,q' + k\,.\,q'\;p\,.\,q - p\,.\,k\;q\,.\,q')$$

$$T_6 = 2(k\,.\,q\;p'\,.\,q' + k\,.\,q'\;p'\,.\,q - p'\,.\,k\;q\,.\,q')$$

$$T_7 = 2(p\,.\,q'\;k\,.\,q - p\,.\,q\;k\,.\,q')$$

$$T_8 = 2(k\,.\,q'\;p'\,.\,q - k\,.\,q\;p'\,.\,q')$$

$$T_9 = 2(p\,.\,q'\;p'\,.\,q - p\,.\,q\;p'\,.\,q')$$

$$T_{10} = 2\{p\,.\,p'(k\,.\,q'\;p'\,.\,q - k\,.\,q\;p'\,.\,q') + p'\,.\,k(p'\,.\,q'\;p\,.\,q - p'\,.\,q\;p\,.\,q') + p^2(p\,.\,q'\;k\,.\,q - p\,.\,q\;k\,.\,q')\}$$

$$T_{11} = 2\{p^2(k\,.\,q'\,p'\,.\,q - k\,.\,q\,p'\,.\,q') + p\,.\,k(p'\,.\,q'\,p\,.\,q - p'\,.\,q\,p\,.\,q') + p\,.\,p'(p\,.\,q'\,k\,.\,q - p\,.\,q\,k\,.\,q')\}$$

$$C_1 = 8PQ(h+i) - 16m_p^2 mU + 16V\{PQ(B^2+L^2) - (p\,.\,p' + B - L)m\} + 16m_p\{wQ + P(B-L)(PW - QX) - P(lW - jX)\}$$

$$C_2 = 16PQG(D+G) + 16dEQ - 4M\{j(H+G) + 2aQ\} + 8QGMR - 4v + 4m_p^2 RM^2 + 32LV(LQ - Q^2 - m_p^2 P - p'\,.\,pL) + 32m_p^2 LU(Q - 2L) + 16g_V^0 I \times\{-m_p^2 L(P+Q) - jQ\} - 8m_p^2 PQI^2 + m_p[-16EQg_A^0(B-L) + 4u + 8g_V^0 \times\{-QM(B-L) - 2QG(P-Q) + 2jM + 4LG(2P-Q) - 2LMR\} + 16PQIG + 8m_p^2 IM(P+Q) + 16LQ(X-Y)]$$

$$C_3 = -8PQ(GC + DH + 2GH) - 8EPd + 4aPM + 4rQ + 4jHM - 4QHMR + 4v + 4OMR - 8dEQ - 4G(AR - s) + 16V\{LP(B-L) - PQ(B-L) + m + m_p^2 L(P+Q) + 2p\,.\,p'L^2\} - 16m_p^2 LU(P + Q - 4L) + 8g_V^0\{2PQB(I+J) - PQJ(B-L) - PI(j - m_p^2 L) - QJ(m_p^2 L + l)\} + 8PQIJp\,.\,p' + m_p[4g_V^0\{2BN(P+Q) - Q(B-L)(N - 2C) + 2QH(P-Q) - 2s - 4L(H+G)(2P-Q) + 2L(t + MR) + 2PG(P-Q) - 2BM(P+Q) + PM(B-L) + 2jM\} + 8g_A^0\{EQ(B-L) - 2BE(P-Q) + PE(B-L)\} - 8PQHI + 4OI(P+Q) + 8P^2GJ - 4p\,.\,p'MJ(P+Q) + 8L(P+Q)(Y-X) - 4u]$$

$$C_4 = 16PQH(C+H) + 16dEP - 8rP - 8sH + 8HAR - 4v - 4nR + 32LV(PQ - 2BP + PL - Qm_p^2 - Lp\,.\,p') + 32m_p^2 LU(P - 2L) + 16g_V^0\{P^2 J(B-L) - lPJ + m_p^2 LQ(I+J)\} - 8m_p^2 PQJ^2 + m_p[16g_A^0 E\{2B(P-Q) - P(B-L)\} + 8g_V^0\{4BC(P+Q) - 2C(B-L)(P+Q) + (B-L)(NP + 2CQ) - 2PH(P-Q) - 2s + 4LH(2P-Q) - 2tL\} - 16P^2 JH + 16PL(X-Y) + 4u + 8(P+Q) \times(-Nm_p^2 + 2Cp\,.\,p')J]$$

$$C_5 = -8aPG + 4aMR - 8dES + 16p\,.\,p'LV(Q-L) + 16m_p^2 LU(B + Q - 2L) + 8g_V^0 I \times\{p\,.\,p'j + m_p^2 P(B-L) + m_p^2 j\} + m_p\{8yG + 8z + 4g_V^0 MR(B-L) + 16PLW + 8aPI + 8x\}$$

$$C_6 = -8P(i+h) + 8aPH - 4rR + 8dES + 16V\{-P(B-L)^2 + lB + m_p^2 LQ + p\,.\,p'L^2\} - 16m_p^2 LU(B + P - 2L) + 8g_V^0 J\{p\,.\,p'l - p\,.\,p'P(B-L) + m_p^2 l\} + m_p\{-8yH - 8z - 4g_V^0 t(B-L) - 16PLX + 8bPJ - 16w - 8x\}$$

$$C_7 = -32g_V^0 g_A^0 j(B-L) - 32m_p g_V^0 jE$$

$$C_8 = -16P(aE + bF) - 32g_V^0 g_A^0 P(B^2 - L^2) - 32m_p g_V^0\{PE(B-L) + lF\}$$

$$C_9 = 32g_V^0 g_A^0\{m - PQ(B-L)\}$$

$$C_{10} = 8PFH+8Ea-4(Nf+2Cd)-32g_V^0g_A^0(L^2+BP)-8g_A^0J\{P(B-L)-l\}+4m_p \times\{2g_V^0E(B-L)-4g_V^0L(E+F)+g_A^0(B-L)(2C-N)+2g_A^0H(P+Q)\}$$

$$C_{11} = -8PFG-8aE+4dM+32g_V^0g_A^0(L^2+BQ)+8g_A^0jl-4m_p\{2Eg_V^0(B-L)-4g_V^0L \times(E+F)+g_A^0M(B-L)+2g_A^0G(P+Q)\}$$

where

$P = p\,.\,k$, $Q = p'\,.\,k$, $B = G^p(0)Q$, $L = G^p(0)P$, $C = 2F^p(0)g_V^0P$, $D = 2F^p(0)g_V^0Q$,
$E = 2F^p(0)g_A^0P$, $F = 2F^p(0)g_A^0Q$, $G = 2Lf_V^0$, $H = 2Bf_V^0$, $I = 4F^p(0)f_V^0P$,
$J = 4F^p(0)f_V^0Q$, $M = 4Pf_V^0+2C$, $N = 4Qf_V^0$, $A = NP+2CQ$, $O = Np\,.\,p'-2Cm_p^2$
$R = -m_p^2+p\,.\,p'$, $S = m_p^2+p\,.\,p'$, $U = g_V^{02}-g_A^{02}$, $V = g_V^{02}+g_A^{02}$,
$W = Dg_V^0-Fg_A^0$, $X = Cg_V^0+Eg_A^0$, $Y = Dg_V^0+Fg_A^0$, $Z = Cg_V^0-Eg_A^0$,
$a = Dp\,.\,p'+m_p^2C$, $b = -m_p^2D-Cp\,.\,p'$, $d = Fp\,.\,p'-m_p^2E$, $f = -Fm_\rho^2+Ep\,.\,p'$,
$h = -aC+bD$, $i = \mathrm{d}E+fF$, $j = Qp\,.\,p'+m_p^2P$, $l = -Qm_p^2-Pp\,.\,p'$,
$m = lQ-jP$, $n = -m_p^2(N^2+4C^2)+4p\,.\,p'CN$, $r = 2aC+bN$, $s = lN+2jC$,
$t = R(N+2C)$, $u = 8g_A^0E(j-LS)$, $v = -4m_p^2E^2S$, $w = lY-jZ$,
$x = -m_p^2L(W-Z)-p\,.\,p'L(X-Y)$, $y = 2g_V^0P(2L-B)-g_V^0(j-l)$, $z = g_A^0ES(B-L)$.

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