The effect of electromagnetic properties of neutrinos on the photon-neutrino decoupling temperature

S.C. İnan and M. Köksal

Department of Physics, Cumhuriyet University, 58140, Sivas, Turkey

Abstract

We examine the impact of electromagnetic properties of neutrinos on the annihilation of relic neutrinos with ultra high energy cosmic neutrinos for the $\nu\bar{\nu} \rightarrow \gamma\gamma$ process. For this process, photon-neutrino decoupling temperature is calculated via effective lagrangian model beyond the standard model. We find that photon-neutrino decoupling temperature can be importantly reduced below the QCD phase transition with the model independent analysis defining electromagnetic properties of neutrinos.
I. INTRODUCTION

Neutrinos and photons are the most abundant particles in the universe. The universe is filled with a sea of relic neutrinos that decoupled from the rest of the matter within the first few seconds after the Big Bang. Unlike the relic photons, relic neutrinos have not been yet observed because of the interactions of their cross sections with matter are overwhelmingly suppressed. It is very important to detect relic neutrinos which have played a crucial role in Big Bang the nucleosynthesis, structure formation and the evolution of the universe. Nevertheless, some indirect evidences of the relic neutrinos may be observed, such as, the UHE neutrinos may interact with relic neutrinos via the $\nu_{\text{cosmic}} + \bar{\nu}_{\text{relic}} \to Z \to n + \gamma$ reactions occurring on the Z resonance [1]. If relic neutrinos do exist, the existence of their mass spectrum may be reveal with detectors of UHE neutrinos, such as Icecube [2], ANITA [3], Pierre Auger Observatory [4], ANTARES [5].

The $\gamma\nu \to \nu\nu$ process has been extensively studied in literature [6–10]. When the neutrinos are massless, the $\nu\bar{\nu} \to \gamma\gamma$ process implies a vanishing cross section from Yang’s theorem [11, 12] due to the vector-axial vector nature of the weak coupling. The cross section of the $\nu\bar{\nu} \to \gamma\gamma$ process can be given to be of order $G_F^2\alpha^2\omega^2(\omega/m_W)^4$ [7, 13]. This situation continues to until center of mass energies $\sqrt{s} \sim 2m_W$ where $m_W$ is the W boson mass. The dimension-8 effective lagrangian induced from loop contributions of SM particles can be given as follows [14]

$$L_{\text{eff}} = i \frac{g^2\alpha}{32\pi m_W^4} A [\bar{\psi} \gamma_\mu(1 - \gamma_5)(\partial^\mu \psi) - (\partial^\mu \bar{\psi})\gamma_\mu(1 - \gamma_5)\psi] F_{\mu\nu} F^{\nu\lambda}$$

where $F_{\mu\nu}$ is the electromagnetic field strength tensor, $g$ is the electroweak gauge coupling, $\alpha$ is the fine structure constant and $A$ is given by

$$A = \left[ \frac{4}{3} \ln \left( \frac{m_W^2}{m_e^2} \right) + 1 \right].$$

It is shown that the equation (1) can be rewritten in the following form [14],

$$L_{\text{eff}} = \frac{1}{8\pi m_W^4} A T^\nu_{\alpha\beta} T^{\gamma\alpha\beta}$$

2
where $T_{\alpha\beta}^{\nu}$ and $T^{\gamma}_{\alpha\beta}$ are the stress-energy tensor of the neutrinos and photons which are given by,

\[
T_{\alpha\beta}^{\nu} = \frac{i}{8} \bar{\psi} \gamma_\alpha (1 - \gamma_5) (\partial_\beta \psi) + \bar{\psi} \gamma_\beta (1 - \gamma_5) (\partial_\alpha \psi) \\
- (\partial_\beta \bar{\psi}) \gamma_\alpha (1 - \gamma_5) \psi - (\partial_\alpha \bar{\psi}) \gamma_\beta (1 - \gamma_5) \psi, \quad (4)
\]

\[
T^{\gamma}_{\alpha\beta} = F_{\alpha\lambda} F^\lambda_{\beta} - \frac{1}{4} g_{\alpha\beta} F^\lambda_{\lambda}. \quad (5)
\]

The photons and neutrinos decouple for the $\nu \bar{\nu} \rightarrow \gamma \gamma$ process is calculated at a temperature $T_c \sim 1.6$ GeV, or approximate one micro second after the Big Bang [8]. If the photon-neutrino interaction can be increased, then decoupling temperature is lowered to the QCD phase transition ($\Lambda_{QCD} \sim 200$ MeV). Therefore, some remnants of the photons circular polarization can possibly be retained in the cosmic microwave background [14] which can be considered as an evidence for the relic neutrino background. Increasing the cross section of $\nu \bar{\nu} \rightarrow \gamma \gamma$ process can be achieved with using models beyond the SM. In this sense, effect of the large extra dimensions [14], unparticle physics [15] and excited neutrinos [16] have been calculated. They have found that the photon-neutrino decoupling temperature can be significantly brought down.

In this study, we have calculated that effect of the electromagnetic properties of neutrinos on the photon-neutrino decoupling temperature for the $\nu \bar{\nu} \rightarrow \gamma \gamma$ process.

II. $\nu \bar{\nu} \rightarrow \gamma \gamma$ PROCESS INCLUDING ELECTROMAGNETIC PROPERTIES OF NEUTRINOS

In the SM, there is no interaction between neutrinos and photons. Besides, minimal extension of the SM with massive neutrinos yields couplings of $\nu \bar{\nu} \gamma$ and $\nu \bar{\nu} \gamma \gamma$ by means of radiative corrections [17-21]. There are a lot of models beyond the SM estimating large enough $\nu \bar{\nu} \gamma$ and $\nu \bar{\nu} \gamma \gamma$ couplings, although minimal extension of the SM give rise to very small couplings. For this reason, it is important to investigate electromagnetic properties of the neutrinos in effective lagrangian methods. Electromagnetic behavior of the neutrinos have significant effects on astrophysics, cosmology and particle physics. In this motivation, we have examined to effect of the Dimension-6 and Dimension-7 effective lagrangians on photon-neutrino decoupling temperature.
A. Dimension-7 Effective Lagrangian

The dimension-7 effective lagrangian defining $\nu\bar{\nu}\gamma\gamma$ coupling can be given by \[\text{[21–26]}\]

\[L = \frac{1}{4\Lambda^3} \bar{\nu}_i \left( \alpha_{R1}^i P_R + \alpha_{L1}^i P_L \right) \nu_j \tilde{F}_{\mu\nu} F^{\mu\nu} + \frac{1}{4\Lambda^3} \bar{\nu}_i \left( \alpha_{R2}^i P_R + \alpha_{L2}^i P_L \right) \nu_j F_{\mu\nu} F^{\mu\nu} \]  

(6)

where $F_{\mu\nu}$ is the electromagnetic field tensor, $\tilde{F}_{\mu\nu} = \frac{1}{2} \epsilon_{\mu\nu\alpha\beta} F^{\alpha\beta}$, $P_{L(R)} = \frac{1}{2} (1 \mp \gamma_5)$, $\alpha_{Lk}^{ij}$ and $\alpha_{Rk}^{ij}$ are dimensionless coupling constants. Latest experimental bounds on neutrino-two photon coupling are obtained from rare decay $Z \to \nu\bar{\nu}\gamma\gamma$ \[\text{[26]}\] and the analysis of $\nu_\mu N \to \nu_s N$ conversion \[\text{[25]}\]. The experimental model independent upper limit for $Z \to \nu\bar{\nu}\gamma\gamma$ decay has been found from the LEP data as follows \[\text{[26]}\],

\[
\left[ \frac{1 \text{GeV}}{\Lambda} \right]^6 \sum_{i,j,k} \left( |\alpha_{Rk}^{ij}|^2 + |\alpha_{Lk}^{ij}|^2 \right) \leq 2.85 \times 10^{-9}.
\]  

(7)

In the external Coulomb field of the nucleus $N$, the model dependent searches of the Primakoff effect on $\nu_\mu N \to \nu_s N$ conversion founds about two orders of magnitude more restrictive bound than LEP data. The potential of photon induced reactions at the LHC to probe electromagnetic properties of the neutrinos has also been studied in the literature for $\Lambda = 1$ GeV \[\text{[27, 28]}\]. It was shown that future experimental researches at the LHC will place more stringent bounds. We have used the model independent bound which was obtained from the LEP data. The contribution of the SM to the $\nu\bar{\nu} \to \gamma\gamma$ process have been calculated in Refs. \[\text{[7, 8]}\] with using equation (1). The squared amplitude for the SM ($|M_1|^2$) can be found from this effective Lagrangian in terms of Mandelstam invariants $s$ and $t$ as below

\[
|M_1|^2 = -16 \left( \frac{g^2 \alpha A}{32\pi M_W^4} \right)^2 t (s^3 + 2t^3 + 3ts^2 + 4t^2s).
\]  

(8)

The new physics contribution with using equation (6) comes from $t$ and $u$ channels diagrams for the $\nu\bar{\nu} \to \gamma\gamma$ process. The polarization summed amplitude dimension-7 effective interaction square ($|M_2|^2$) is given below,

\[
|M_2|^2 = \frac{s^3}{8\Lambda^6} \sum_{i,j,k} \left( |\alpha_{Rk}^{ij}|^2 + |\alpha_{Lk}^{ij}|^2 \right).
\]  

(9)
It has been obtained that there is no contribution from the interference term of the SM and dimension-7 effective interaction to the $\nu \bar{\nu} \rightarrow \gamma \gamma$ scattering. The reason is that the SM interaction contains neutrinos of opposite helicity, dimension-7 effective interaction contain neutrinos of the same helicity. Hence, the total squared amplitude can be found,

$$|M|^2 = |M_1|^2 + |M_2|^2. \quad (10)$$

For $\nu \bar{\nu} \rightarrow \gamma \gamma$ process, the differential cross section can be obtained by using

$$\frac{d\sigma}{dz} = \frac{1}{2!} \frac{1}{32\pi s} |M|^2. \quad (11)$$

Therefore, we get the total cross section ($\sigma_{cm}$) as follows,

$$\sigma_{\nu \bar{\nu} \rightarrow \gamma \gamma} = \frac{s^3}{20\pi} \left( \frac{g^2 \alpha A}{32\pi M_W^4} \right)^2 + \frac{s^2}{256\pi \Lambda^6} \sum_{i,j,k} \left( |\alpha_{Rk}^{ij}|^2 + |\alpha_{Lk}^{ij}|^2 \right). \quad (12)$$

We have showed as a function of the center of mass energy $\sqrt{s}$ for both the SM and total cross sections in fig. (1). During numeric analysis we have assumed to $\Lambda = 1$ GeV to compare our results with current experimental LEP limit. In this figure, $\beta^2 = \sum_{i,j,k} \left( |\alpha_{Rk}^{ij}|^2 + |\alpha_{Lk}^{ij}|^2 \right)$ is taken to be $2.89 \times 10^{-9}$ which is current experimental LEP bound. It has been shown that deviation from the SM increases as the $\sqrt{s}$ decreases. Also, Fig. (2) shows that the SM and total cross sections via the $\beta^2$ for $\sqrt{s} = 5$ GeV. The total cross section is nearly the same as the SM at $\beta^2 \sim 10^{-13}$. This value almost $10^4$ times larger than the current experimental LEP limit. Specific values of the $\beta^2$ and $\sqrt{s}$ total cross section can be easily discerned from the SM cross section. Therefore, dimension-7 effective interaction can effect to the photon-neutrino decoupling temperature.

The temperature at which the $\nu \bar{\nu} \rightarrow \gamma \gamma$ process ceases to take place can be found from the reaction rate per unit volume,

$$\rho = \frac{1}{(2\pi)^6} \int \frac{d^3 p_1}{\exp(E_1/T) + 1} \int \frac{d^3 p_2^*}{\exp(E_2/T) + 1} \sigma |\vec{\nu}|. \quad (13)$$

where $\vec{p}_1$ and $\vec{p}_2^*$ are the momentums of the neutrinos, $E_1$ and $E_2$ are the energies of the neutrinos, $T$ is the temperature, $|\vec{\nu}|$ is the flux. The $\sigma |\vec{\nu}|$ can be obtained in terms of $\sigma_{cm}$ in the center of mass frame by using of invariance of $\sigma |\vec{\nu}| E_1 E_2$.
\[ \sigma|\vec{v}| = \frac{\sigma_{\text{cm}} s}{2E_1E_2} \]  

(14)

\[ \sigma|\vec{v}| = \frac{s^4}{40\pi E_1 E_2} \left( \frac{g^2 \alpha A}{32\pi M_W^4} \right)^2 + \frac{s^3 \beta^2}{512\pi E_1 E_2} \]  

(15)

where \( s = 2E_1E_2 (1 - \cos \theta_{12}) \) and \( \theta_{12} \) is the angle between \( \vec{p}_1 \) and \( \vec{p}_2 \). Then the reaction rate per unit volume can be obtained as follows,

\[ \rho = \frac{g^4 \alpha^2 A^2}{25(2\pi)^7 m_W^8} T^{12} \int_0^\infty \frac{x^5 dx}{e^x + 1} \int_0^\infty \frac{y^5 dy}{e^y + 1} + \frac{\beta^2}{4(2\pi)^5} T^{10} \int_0^\infty \frac{x^4 dx}{e^x + 1} \int_0^\infty \frac{y^4 dy}{e^y + 1} \]  

(16)

where \( x = E_1/T \) and \( y = E_2/T \). The integration is easily written by

\[ \rho = \frac{g^4 \alpha^2 A^2}{25(2\pi)^7 m_W^8} T^{12} \left[ \frac{31}{32} \Gamma(6) \zeta(6) \right]^2 + \frac{\beta^2}{4(2\pi)^5} T^{10} \left[ \frac{15}{16} \Gamma(5) \zeta(5) \right]^2 \]  

(17)

where \( \zeta(x) \) is the Riemann Zeta function. At temperature \( T \), the interaction rate \( R \) can be found by dividing \( \rho \) by the neutrino density \( n_\nu = 3\zeta(3)T^3/4\pi^2 \),

\[ R = 2.30 \times 10^4 \left( \frac{T}{\text{GeV}} \right)^9 + 2.31 \times 10^{23} \beta^2 \left( \frac{T}{\text{GeV}} \right)^7 \text{sec}^{-1}. \]  

(18)

Multiplying equation (18) by the age of the universe,

\[ t = 1.48 \times 10^{-6} \left( \frac{T}{\text{GeV}} \right)^{-2} \]  

(19)

at least one interaction to occur is \( Rt = 1 \). As a result, the decoupling temperature can be found with solution of the following equation,

\[ 3.40 \times 10^{-2} \left( \frac{T}{\text{GeV}} \right)^7 + 3.42 \times 10^{17} \beta^2 \left( \frac{T}{\text{GeV}} \right)^5 = 1. \]  

(20)

In Fig. (3) we have plotted the solution of the this equation for different values of the \( \beta^2 \). Here, current experimental LEP bound have taken to be maximum value of the \( \beta^2 \).
B. Dimension-6 Effective Lagrangian

The Dimension-6 effective lagrangian for non-standard $\nu\bar{\nu}\gamma$ interaction [26, 29–31] is given by

$$L = \frac{1}{2} \mu_{ij} \bar{\nu}_i \sigma_{\mu\nu} \nu_j F^{\mu\nu}$$

(21)

here $\mu_{ii}$ is the magnetic moment of $\nu_i$ and $\mu_{ij}$ ($i \neq j$) is the transition magnetic moment. In equation (21), new physics energy scale $\Lambda$ is absorbed in the definition of $\mu_{ii}$. We will examine $\nu\bar{\nu}\gamma$ interaction on the $\nu\bar{\nu} \rightarrow \gamma\gamma$ process assuming neutrino magnetic moment matrix is virtually flavor diagonal and only one of the matrix elements is different from zero. Also, the standard relic neutrinos is considered to comprise of the three active neutrinos of the SM. Current experimental bounds on neutrino magnetic moment are stringent. The most sensitive bounds from neutrino-electron scattering experiments with reactor neutrinos are at the order of $10^{-11} \mu_B$ [32–35]. Bounds derived from solar neutrinos are at the same order of magnitude [36]. Bounds on magnetic moment can also be derived from energy loss of astrophysical objects. These give about an order of magnitude more restrictive bounds than reactor and solar neutrino probes [37–43].

The polarization summed amplitude square for the $\nu\bar{\nu} \rightarrow \gamma\gamma$ process is given by the following equation,

$$|M|^2 = -16 \left( \frac{g^2 \alpha A}{32\pi M_W^4} \right)^2 t (s^3 + 2t^3 + 3ts^2 + 4t^2s) + 16\mu^4 tu + 32\mu^2 tus \left( \frac{g^2 \alpha A}{32\pi M_W^4} \right).$$

(22)

Then the total cross section for the $\nu\bar{\nu} \rightarrow \gamma\gamma$ process can be obtained as follows,

$$\sigma_{\nu\bar{\nu} \rightarrow \gamma\gamma} = \int_{-1}^{1} dz \frac{d\sigma}{dz} = \frac{s^3}{20\pi} \left( \frac{g^2 \alpha A}{32\pi M_W^4} \right)^2 + \frac{\mu^2 s}{12\pi} \left( \mu^2 + 2s \left( \frac{g^2 \alpha A}{32\pi M_W^4} \right) \right).$$

(23)

We have calculated the total cross section with using experimental limits of the neutrino magnetic moments ($\mu_{\nu_i}, i = e, \mu, \tau$) for the $\nu\bar{\nu} \rightarrow \gamma\gamma$ process. These bounds are $\mu_e =$
3.2 \times 10^{-11} \mu_B, \mu_\mu = 6.8 \times 10^{-10} \mu_B \text{ and } \mu_\tau = 3.9 \times 10^{-7} \mu_B \text{[44].} \text{ It has been seen that there are barely contribution from neutrino magnetic moments to the SM cross section of this process and we have not shown results in here. Therefore, this effective interaction must not reduce to photon-neutrino decoupling temperature significantly. This result can be seen with using same procedure as above. Then, the } \rho \text{ and } R \text{ are calculated by,}

\begin{align*}
\rho &= \frac{g^4 \alpha^2 A^2}{25(2\pi)^7 m_W^8} T^{12} \left[ \frac{31}{32} \Gamma(6) \zeta(6) \right]^2 + \\
&\quad \frac{\mu^2}{18\pi^5} \left( 6 \frac{g^2 \alpha A}{32\pi M_W} \right) T^{10} \left[ \frac{15}{16} \Gamma(5) \zeta(5) \right]^2 + \mu^2 T^8 \left[ \frac{7}{8} \Gamma(4) \zeta(4) \right]^2, \quad (24)
\end{align*}

\begin{align*}
R &= 2.30 \times 10^4 \left( \frac{T}{\text{GeV}} \right)^9 + 8.85 \times 10^{13} \mu^2 \left( \frac{T}{\text{GeV}} \right)^7 + 9.71 \times 10^{22} \mu^4 \left( \frac{T}{\text{GeV}} \right)^5 \text{ sec}^{-1}. \quad (25)
\end{align*}

The solution of the following equation gives the decoupling temperature for photon-neutrino coupling,

\begin{align*}
3.40 \times 10^{-2} \left( \frac{T}{\text{GeV}} \right)^7 + 1.31 \times 10^{7} \mu^2 \left( \frac{T}{\text{GeV}} \right)^5 + 1.44 \times 10^{17} \mu^4 \left( \frac{T}{\text{GeV}} \right)^3 = 1. \quad (26)
\end{align*}

From this equation, we have found that the photon-neutrino decoupling temperature almost same the SM \((T_c \sim 1.6 \text{ GeV})\) when we used the experimental bounds on neutrino magnetic moments as we expected.

III. CONCLUSION

If neutrino-photon decoupling temperature can be decreased to below the QCD phase transition \((\Lambda_{QCD} \sim 200 \text{ MeV})\), this could be an evidence for the relic neutrino background. Because some remnant the circular polarization could possibly be sustained in the cosmic microwave background. For reducing decoupling temperature, the total cross section of the photon-neutrino process should be increased. This can be done with contribution of new effective interactions. In this motivation, we have examined the effect of electromagnetic properties of the neutrinos on the photon-neutrino decoupling temperature with interaction of relic neutrinos with UHE cosmic neutrinos via the \(\nu \bar{\nu} \rightarrow \gamma \gamma\) process. First, we have
investigated to dimension-7 effective interaction effect on $\nu\bar{\nu} \rightarrow \gamma\gamma$ process. It is found that this effective interaction contribution to total cross section of the $\nu\bar{\nu} \rightarrow \gamma\gamma$ process is significant depending on the $\beta^2$. Therefore, photon-neutrino decoupling temperature can be reduced below the $\Lambda_{QCD}$ as seen from the Fig. On the other hand, even if $\beta^2$ is eight order of magnitude smaller than current experimental bound, this effective interaction can reduce to $T_c$ below the obtained value of the SM.

Second, we have examined to dimension-6 effective interaction impact on $\nu\bar{\nu} \rightarrow \gamma\gamma$ process. This effective interaction describes neutrino magnetic moment. Current experimental bounds on neutrino magnetic moment are stringent. Therefore, the contribution of the this effective interaction very tiny on the SM cross section $\nu\bar{\nu} \rightarrow \gamma\gamma$. Hence, the photon-neutrino cross section decoupling temperature is not almost changed.

Consequently, we have shown that dimension-7 effective interaction can permit of reduced the decoupling temperature for the $\nu\bar{\nu} \rightarrow \gamma\gamma$ process.

[1] T. J. Weiler, Phys. Rev. Lett. 49, 234 (1982).
[2] R. Abbasi et al., Phys. Rev. D 83, 092003 (2011).
[3] S. W. Barwick et al., Phys. Rev. Lett. 96, 171101 (2006).
[4] P. Abreu et al., Phys. Rev. D 84, 122005 (2011).
[5] J. A. Aguilar et al., Nuclear Inst. and Methods in Physics Research A 656, 11 (2011).
[6] A. Abbasabadi et al., Phys. Rev. D 59, 013012 (1998).
[7] D. A. Dicus and W. W. Repko, Phys. Rev. D 48, 5106 (1993).
[8] A. Abbasabadi et al., Phys. Rev. D 59, 013012 (1999).
[9] A. Abbasabadi, A. Devoto and W. W. Repko, Phys. Rev. D 63, 093001 (2001).
[10] W. P. Lam and K. W. Ng, Phys. Rev. D 44, 3345 (1991).
[11] C. N. Yang, Phys. Rev. 77, 242 (1950).
[12] M. Gell-Mann, Phys. Rev. Lett 6, 70 (1961).
[13] M. J. Levine, Nuovo Cimento A XLVIII, 67 (1967).
[14] D. A. Dicus, K Kovner and W W Repko, Phys. Rev. D 62, 053013 (2000).
[15] S. Dutta and A. Goyal, Phys. Lett. B 664, 25 (2008).
[16] S. C. İnan and M. Köksal, arXiv: 1203.5881.
[17] B. W. Lee and R. E. Schrock, Phys. Rev. D 16, 1444 (1977).
[18] W. Marciano and A. I. Sanda, Phys. Lett. B 67, 303 (1977).
[19] B. W. Lynn, Phys. Rev. D 23, 2151 (1981).
[20] R. J. Crewther, J. Finjord and P. Minkowski, Nucl. Phys. B 207, 269 (1982).
[21] S. Dodelson and G. Feinberg, Phys. Rev. D 43, 913 (1991).
[22] J. F. Nieves, Phys. Rev. D 28, 1664 (1983).
[23] R. K. Ghosh, Phys. Rev. D 29, 493 (1984).
[24] J. Liu, Phys. Rev. D 44, 2879 (1991).
[25] S. N. Gninenko and N. V. Krasnikov, Phys. Lett. B 450, 165 (1999).
[26] F. Larios, M. A. Perez and G. Tavares-Velasco, Phys. Lett. B 531, 231 (2002).
[27] İ. Şahin and M. Köksal, JHEP 03, 100 (2011).
[28] İ. Şahin, Phys. Rev. D 85, 033002 (2012).
[29] F. Larios, R. Martinez and M. A. Perez, Phys. Lett. B 345, 259 (1995).
[30] M. Maya et al., Phys. Lett. B 434, 354 (1998).
[31] N. F. Bell et al., Phys. Rev. Lett. 95, 151802 (2005).
[32] H. B. Li et al., Phys. Rev. Lett. 90, 131802 (2003).
[33] Z. Daraktchieva et al., Phys. Lett. B 615, 153 (2005).
[34] H. T. Wong et al., Phys. Rev. D 75, 012001 (2007).
[35] H. T. Wong et al., Phys. Rev. Lett. 105, 061801 (2010).
[36] C. Arpesella et al., Phys. Rev. Lett. 101, 091302 (2008).
[37] G. G. Raffelt, Phys. Rep. 320, 319 (1999).
[38] V. Castellani and S. Degl’Innocenti, Astrophys. J. 402, 574 (1993).
[39] M. Catelan, J. A. d. Pacheco and J. E. Horvath, Astrophys. J. 461, 231 (1996).
[40] A. Ayala, J. C. D’Olivo and M. Torres, Phys. Rev. D 59, 111901 (1999).
[41] R. Barbieri and R. N. Mohapatra, Phys. Rev. Lett. 61, 27 (1988).
[42] J. M. Lattimer and J. Cooperstein, Phys. Rev. Lett. 61, 23 (1988).
[43] A. Heger, A. Friedland, M. Giannotti and V. Cirigliano, Astrophys. J. 696, 608 (2009).
[44] K. Nakamura et al., (Particle Data Group) J. Phys. G 37, 075021 (2010).
FIG. 1: The cross sections of $\nu\bar{\nu} \rightarrow \gamma\gamma$ process as a function center of mass energy $s^{1/2}$ when $\beta^2$ parameter is taken to be $2.89 \times 10^{-9}$.

FIG. 2: The SM and total cross sections of $\nu\bar{\nu} \rightarrow \gamma\gamma$ process as a function $\beta^2$ for $s^{1/2} = 5$ GeV.
FIG. 3: The decoupling temperature $T_c$ as a function of $\beta^2$. 