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Remarks on the process  $\gamma\gamma \rightarrow \nu\bar{\nu}$  in astrophysics

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We discuss the astrophysical consequences of the process  $\gamma\gamma \rightarrow \nu\bar{\nu}$ , when this reaction is mediated by either Majorons or composite neutral leptons, and present the constraints on their coupling coming from stellar energy loss arguments. We also discuss the effect of nuclear absorption when this reaction is mediated by pions, and show that no significant output of energy is provided in this case. Finally, we comment on the importance of these processes in cosmology.

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## I. INTRODUCTION

The process  $\gamma\gamma \rightarrow \nu\bar{\nu}$  has been thought of for a long time as an important reaction in astrophysics and cosmology since it is able to produce energy loss, modifying the stellar evolution. Several aspects of this reaction were studied since it was first discussed by Chiu and Morrison [1]. In particular, Gell-Mann showed that it is forbidden in a local  $V-A$  theory [2]. This scattering can occur, at the one loop level, in the framework of the Weinberg-Salam theory as shown by Dicus *et al.* [3] although, in this case, the energy loss is smaller than the one due to competing processes (e.g.,  $e^+e^- \rightarrow \nu\bar{\nu}$  and  $\gamma e \rightarrow e\nu\bar{\nu}$ ). Nevertheless, the reaction  $\gamma\gamma \rightarrow \nu\bar{\nu}$  can be important if there exist exotic scalar or pseudoscalar weak interactions [4].

In this work, we study the contribution to the reaction  $\gamma\gamma \rightarrow \nu\bar{\nu}$  of a 1 keV Majoron ( $J$ ), which has been proposed as a dark matter candidate [5]. The reaction  $\gamma\gamma \rightarrow J \rightarrow \nu\bar{\nu}$  is particularly important when the new pseudoscalar particle is light enough to be on resonance at the temperatures of known stars. An interesting feature here is that stellar energy loss constrains simultaneously the Majoron coupling to photons and to neutrinos, which, in principle, can be extended to any other particle with similar properties.

Many composite models [6] predict the existence of excited states of the usual leptons, and also new interactions between neutrinos and photons. We also compute in this paper the effect of new neutral fermions of spin  $\frac{1}{2}$  and  $\frac{3}{2}$ , predicted by these models, on the reaction  $\gamma\gamma \rightarrow \nu\bar{\nu}$ . We establish constraints on the compositeness scale and on the mass of these new states based on stellar energy loss arguments.

In our conclusions, we evaluate the effect of pion ab-

sorption in the nuclear media, which leads to a strong suppression of the process  $\gamma\gamma \rightarrow \pi^0 \rightarrow \nu\bar{\nu}$  [7,8], showing that it destroys the possibility of obtaining limits on  $\pi^0$ - $\nu$  couplings. Finally, we trace some comments on the importance of these processes in cosmology analyzing the production of neutrinos by the photons from the cosmic thermal background.

II. MAJORON CONTRIBUTION TO  $\gamma\gamma \rightarrow \nu\bar{\nu}$ 

The existence of pseudoscalar particles that couple to photons and to neutrinos and that are light enough to be produced in the known stars is predicted by many Majoron models. In this class of models, the effective coupling of the Majoron to photons,  $h_{J\gamma\gamma} J F_{\mu\nu} \tilde{F}^{\mu\nu}$ , is usually generated at the one-loop level, and its effective coupling to neutrinos is  $-iJ\bar{\nu}h_{J\nu\nu}\gamma_5\nu$ . The neutrino mass eigenstates are generated according to the "seesaw" mechanism, and we assume that at least one of these states is light ( $m_\nu \leq m_J/2$ ).

The model proposed by Berezhinsky and Valle [5] predicts the existence of an  $\sim 1$  keV Majoron which is a dark matter candidate. In this model, the coupling of the neutrinos to the Majoron appears naturally at a higher-loop level, and the Majoron couples to charged leptons only through some new directly interacting particles.

In order to be a dark matter candidate, the Majoron lifetime into photons ( $\tau_{J\gamma\gamma}$ ) should be large enough to guarantee that the energy density of the produced photons does not exceed the observed energy density. The coupling of the Majoron to photons is given by the electromagnetic anomaly

$$h_{J\gamma\gamma} = \frac{\alpha h_l}{2\pi m_l}, \quad (1)$$

where  $h_l$  is the Majoron coupling to the lepton with mass  $m_l$ , and the Majoron lifetime into two photons is

$$\tau_{J\gamma\gamma} = \frac{64\pi^3}{\alpha^2 h_l^2} \left( \frac{m_l}{m_J} \right)^2 \frac{1}{m_J}. \quad (2)$$

Therefore the coupling to the heaviest lepton ( $\tau$ ) must be constrained to

$$h_\tau \leq 6 \times 10^{-13} H^{-1} \left( \frac{m_J}{1 \text{ keV}} \right)^{-3/2}, \quad (3)$$

where  $H$  is the Hubble constant.

On the other hand, the Majoron lifetime into neutrinos must be larger than the age of the Universe, which implies that [5]

$$h_{J\nu\nu} \leq 1.3 \times 10^{-17} \left( \frac{m_J}{1 \text{ keV}} \right)^{-1/2} H. \quad (4)$$

We can compute the energy output due to  $\gamma\gamma \rightarrow J \rightarrow \nu\bar{\nu}$  from a red giant in order to establish bounds on the

Majoron coupling to photons and neutrinos. The cross section for the process  $\gamma\gamma \rightarrow J \rightarrow \nu\bar{\nu}$ , in the limit  $m_\nu \ll m_J$ , is given by

$$\sigma_{\gamma\gamma \rightarrow \nu\bar{\nu}} = \frac{1}{4\pi} \frac{h_{J\gamma\gamma}^2 h_{J\nu\nu}^2 s^2}{(s - m_J^2)^2 + m_J^2 \Gamma_J^2}, \quad (5)$$

where the effect of the neutrino, as long as it is lighter than the Majoron, is negligible. In Eq. (5) we assumed that the Majoron decay preferentially into two photons, i.e.,  $\Gamma_J \simeq \Gamma_{J\gamma\gamma} \simeq 3.1 \times 10^{-42}$  eV, which is the maximum Majoron width consistent with Eqs. (2) and (3). In this way, we will be underestimating the bound in the Majoron couplings  $h_{J\gamma\gamma}$ , and  $h_{J\nu\nu}$ .

Following the results of Ref. [8], the output of energy in the form of neutrinos is

$$Q_J = \frac{h_{J\gamma\gamma}^2 h_{J\nu\nu}^2}{8\pi^5} (kT)^3 m_J^4 I(\tau, \gamma), \quad (6)$$

where  $I(\tau, \gamma)$  is the integral

$$I(\tau, \gamma) = \int \int \frac{dx_1}{e^{x_1} - 1} \frac{dx_2}{e^{x_2} - 1} (x_1 + x_2) \left\{ 8x_1 x_2 \tau^2 (x_1 x_2 \tau^2 + 1) + \frac{(3 - \gamma^2)}{2} \ln \left[ \frac{(1 - 4x_1 x_2 \tau^2)^2 + \gamma^2}{1 + \gamma^2} \right] \right. \\ \left. + \frac{(1 - 3\gamma^2)}{\gamma} \left[ \arctan(1/\gamma) - \arctan\left( \frac{1 - 4x_1 x_2 \tau^2}{\gamma} \right) \right] \right\}, \quad (7)$$

with  $\tau = kT/m_J$ ,  $\gamma = \Gamma_J/m_J$ , and  $\Gamma_J$  being the total decay width.

A red giant with an average temperature  $T = 10^8$  K and density  $\rho = 10^4$  g/cm<sup>3</sup> has a thermal energy  $kT \sim 10$  keV, and, consequently, the process  $\gamma\gamma \rightarrow \nu\bar{\nu}$  can be on resonance for a Majoron with a mass up to a few keV. The red giant emissivity is limited by the nuclear energy generation rate [9]

$$Q_{\text{RG}} \leq 10^6 \text{ erg cm}^{-3} \text{ s}^{-1}, \quad (8)$$

which imposes a bound on Eq. (6). In the case of a red giant, for the 1 keV Majoron, the comparison of the numerical evaluation of Eq. (6) with Eq. (8) leads to

$$h_{J\gamma\gamma} h_{J\nu\nu} \leq 6.1 \times 10^{-28} \text{ erg}^{-1}. \quad (9)$$

If we consider  $h_{J\gamma\gamma}$  (1), and take into account the limit (3) on  $h_\tau$ , we have  $h_{J\gamma\gamma} \leq 2.5 \times 10^{-13} \text{ erg}^{-1}$ . Therefore, we can impose a bound on the Majoron-neutrino coupling of  $h_{J\nu\nu} \leq 2.5 \times 10^{-15}$ .

The limits on the Majoron coupling to photons and neutrinos are akin to the ones obtained imposing that this particle should be a dark matter candidate. However, the result of Eq. (9) is independent of this condition, and it could be even stronger if we had a heavier Majoron. This limit can also be extended to any similar model, becoming the more stringent the larger is the mass of the pseudoscalar.

### III. COMPOSITE MODELS AND $\gamma\gamma \rightarrow \nu\bar{\nu}$

It is argued by many authors that the standard model could be only the low energy limit of a more fundamental interaction, characterized by a large mass scale  $\Lambda$ . An example of such new interaction is provided by composite models [6], which also predicts the existence of excited states of the usual fermions. We will consider here the contribution of an excited neutral lepton of spin  $\frac{1}{2}$  or  $\frac{3}{2}$  to the reaction  $\gamma\gamma \rightarrow \nu\bar{\nu}$ .

For our present purpose, we choose a model [10] where the excited spin  $\frac{1}{2}$  electron and its neutrino are assumed to form a weak doublet ( $L^*$ ) which couples to the usual fermion doublet ( $l$ ) through the effective Lagrangian

$$\mathcal{L}_{1/2} = \frac{gf}{\Lambda} \bar{L}^* \sigma^{\mu\nu} \frac{\vec{\tau}}{2} l_L \partial_\mu \vec{W}_\nu \\ + \frac{g'f'}{\Lambda} \bar{L}^* \sigma^{\mu\nu} Y l_L \partial_\mu B_\nu + \text{H.c.}, \quad (10)$$

where  $f$  [ $f'$ ] is a scale factor of the SU(2) [U(1)] coupling constants. In a similar way, we can construct an interaction Lagrangian for the spin  $\frac{3}{2}$  excited particle [11]:

$$\mathcal{L}_{3/2} = \frac{gf}{\Lambda} \bar{L}_\mu^* \gamma_\nu \frac{\vec{\tau}}{2} l_L \vec{W}^{\mu\nu} + \frac{g'f'}{\Lambda} \bar{L}_\mu^* \gamma_\nu Y l_L B_{\mu\nu} + \text{H.c.}, \quad (11)$$

where  $L_\mu^*$  denotes a doublet of vector-spinor fields.

The total cross section for  $\gamma\gamma \rightarrow \nu\bar{\nu}$  due to a spin  $\frac{1}{2}$  composite lepton exchange is

$$\sigma_{1/2} = \frac{1}{128\pi} \left( \frac{e\Delta f}{\Lambda} \right)^4 (M_{1/2})^2 \times \left[ \frac{x}{6} - 1 - \frac{4}{x} + \frac{1}{x^2}(4+3x)\ln(1+x) \right], \quad (12)$$

while, for exchange of a spin  $\frac{3}{2}$  neutral lepton, we have

$$\sigma_{3/2} = \frac{1}{96^2\pi} \left( \frac{e\Delta f}{\Lambda} \right)^4 (M_{3/2})^2 \left[ \frac{x^3}{10} - 3x^2 - 73x - 156 - \frac{84}{x} + \frac{6}{x^2}(14+5x)(1+x)^2\ln(1+x) \right], \quad (13)$$

where we defined  $x \equiv s/M_{1/2(3/2)}^2$  and  $\Delta f = f - f'$ . For  $s \ll M_{1/2(3/2)}^2$  these expressions become very simple:

$$\sigma_{1/2} = \frac{1}{2560\pi} \left( \frac{e\Delta f}{\Lambda} \right)^4 \frac{s^3}{M_{1/2}^4} \quad \text{and} \quad \sigma_{3/2} = \frac{1}{9}\sigma_{1/2}. \quad (14)$$

The energy output from a plasma in thermodynamical equilibrium at temperature  $T$ , due to the  $\gamma\gamma \rightarrow \nu\bar{\nu}$  reaction is given by

$$Q_{1/2(3/2)} = \frac{1}{2\pi^4} \int \frac{\omega_1^2 d\omega_1}{\exp(\omega_1/kT) - 1} \int \frac{\omega_2^2 d\omega_2}{\exp(\omega_2/kT) - 1} \times \int_{-1}^1 (\omega_1 + \omega_2)(1 - \eta) \sigma_{1/2(3/2)}^{\gamma\gamma \rightarrow \nu\bar{\nu}}(s) d\eta, \quad (15)$$

where  $\omega$  is the energy of the initial photons and  $\eta = \cos\theta$ , with  $\theta$  being the angle between the initial photon momenta. For a red giant, we may use the cross section given by Eq. (14), obtaining, for a spin  $\frac{1}{2}$  exchange,

$$Q_{1/2} = \frac{1}{50\pi^5} \left( \frac{e\Delta f}{\Lambda} \right)^4 \frac{(kT)^{13}}{M_{1/2}^4} \Gamma(7)\Gamma(6)\zeta(7)\zeta(6), \quad (16)$$

and  $Q_{3/2} = Q_{1/2}/9$  for a spin  $\frac{3}{2}$  exchange. Therefore, the bound coming from the comparison with Eq. (8) gives

$$\frac{|\Delta f|}{\Lambda_{1/2} M_{1/2}} \leq 2.8 \times 10^2 \text{ GeV}^{-2}$$

and

$$\frac{|\Delta f|}{\Lambda_{3/2} M_{3/2}} \leq 2.6 \times 10^3 \text{ GeV}^{-2}. \quad (17)$$

In the case of a light excited neutral lepton (i.e.,  $s \gg M_{1/2}^2$ ), the cross section for the reaction  $\gamma\gamma \rightarrow \nu\bar{\nu}$  does not depend on  $M_{1/2}$ , and we are able to obtain a bound exclusively on the compositeness scale through the comparison of the red giant energy output of  $|\Delta f|/\Lambda_{1/2} \leq 1.8 \times 10^{-2} \text{ GeV}^{-1}$ . Recent results from the CERN  $e^+e^-$  collider LEP Collaborations [12], assuming that  $B(\nu^* \rightarrow \nu\gamma) = 1$ , impose a bound of  $(f \cot\theta_W + f' \tan\theta_W)/(4\Lambda) \leq 1.5 \times 10^{-4} \text{ GeV}^{-1}$ , almost independently of the excited neutrino mass in the range  $0 < M_{\nu^*} < 90 \text{ GeV}$ . In this way, the LEP results allow

a very light excited fermion as far as it is very weakly coupled to the  $Z$ . For the sake of definiteness, if we take  $f = 1$ ,  $f' = 0$ , we can see that the astrophysics bound for a heavy excited neutral lepton is less stringent than the one that can be obtained from collider experiments. Also for a light excited fermion, the astrophysics limit is about two orders of magnitude weaker than the one coming from LEP measurements.

#### IV. COMMENTS AND CONCLUSIONS

We investigated some consequences of the process  $\gamma\gamma \rightarrow \nu\bar{\nu}$  in astrophysics paying particular attention to two different sources of this reaction, i.e., coming from a Majoron and new composite neutral leptons exchange. Through the study of the Majoron contribution, we were able to obtain strong constraints on its coupling to photons and neutrinos. For a 1 keV Majoron dark matter candidate the limit on its coupling to photons and neutrinos is  $h_{J\gamma\gamma} h_{J\nu\nu} \leq 6.1 \times 10^{-28} \text{ erg}^{-1}$ . In general, this process can constrain the couplings of any light pseudoscalar particle to photons and neutrinos, or any other light and weakly interacting particles. We also obtained limits on the compositeness scale when the reaction  $\gamma\gamma \rightarrow \nu\bar{\nu}$  arises from new composite interactions between photons and neutrinos.

In principle, the reaction  $\gamma\gamma \rightarrow \pi^0 \rightarrow \nu\bar{\nu}$  [8] could also be an important source of supernova energy loss. However, the effect of nuclear pion absorption [13] alters the pion decay width in a dense medium and spoils the bounds on the exotic  $\pi^0$  decay coming from this reaction. On resonance, the cross section in a dense medium will be proportional to  $\Gamma_\gamma \Gamma_\nu / \Gamma_{\text{abs}}$ , where  $\Gamma_\gamma \equiv \Gamma_{\pi^0 \rightarrow \gamma\gamma}$ ,  $\Gamma_\nu \equiv \Gamma_{\pi^0 \rightarrow \nu\bar{\nu}}$ , and  $\Gamma_{\text{abs}}$  is the absorption rate of pions formed on resonance in the supernovae nuclear medium. This absorption rate is given by  $\Gamma_{\text{abs}} = \sigma_{\text{abs}} n_N v_{\pi N}$ , where  $\sigma_{\text{abs}}$  is the cross section for true pion absorption,  $n_N$  is the nucleon density, and  $v_{\pi N}$  is the relative pion-nucleon velocity. Since the pion formation is resonant, we have  $v_{\pi N} \approx v_N \sim p_F/m_N$ , where  $v_N$  is the nucleon velocity, and  $p_F = (3\pi^2 n_N)^{1/3}$  is the Fermi momentum, with  $n_N = (0.001 - 0.003) \text{ GeV}^3$  [14]. The main source of pion absorption is via  $\pi^0 NN \rightarrow NN$ , and we can estimate the cross section for this reaction assuming that the dense matter is composed by elementary nuclei (e.g., deuteron) [15], so that  $\sigma_{\text{abs}} \approx 1/\lambda_{\text{abs}}(\rho_N/2)$ , where  $\lambda_{\text{abs}} \sim 5 \text{ fm}$  is the pion absorption length for the energies that we are considering [15,16] and  $\rho_N$  is approximately the quasideuteron density. Therefore, we can establish that  $\Gamma_{\text{abs}}$  ranges from 10 to 100 MeV. This leads, however, to a bound on  $\Gamma(\pi^0 \rightarrow \nu\bar{\nu})$  worse than the present experimental result [17], since  $\Gamma_{\text{abs}}/\Gamma_\gamma$  is  $\sim 10^6$ .

The reaction  $\gamma\gamma \rightarrow \nu\bar{\nu}$  can also give rise to some cosmological consequences once photons from the thermal background may produce right-handed neutrinos, influencing the bound on light neutrinos (or equivalent particles), originated from big-bang nucleosynthesis [18]. The reaction  $\gamma\gamma \rightarrow \pi^0 \rightarrow \nu\bar{\nu}$  was studied by Lam and Ng [19], who obtained a strong limit on the decay  $\pi^0 \rightarrow \nu\bar{\nu}$ . It is important to point out that the strong absorption of  $\pi^0$

by a nucleon does not interfere in this process. The nucleon density is given here by  $n_N = n_{N_0}(R_0/R)^3$ , where  $n_{N_0}$  is the present density of the universe,  $R_0$  its radius, and for temperatures of  $O(m_\pi/2)$ , the ratio  $(R_0/R)$  is  $\sim 10^{11}$  [20]. With the baryonic density limited by  $0.01 < \Omega_b H^2 < 0.02$ , we verify that the nucleon density is too small to introduce any sizable effect of nuclear absorption. At the temperature where the pion mechanism is at resonance, the nucleon density is already frozen. However, the strong interaction among pions occurs much faster than the pion decay: the rate of  $\pi - \pi$  scattering,  $\Gamma_{\pi-\pi} \simeq 0.2$  MeV, dominates the pion lifetime in the dense medium, resulting in a suppression of several orders of magnitude in the rate of neutrino production. Therefore, it is very unlikely that the constraints obtained in Refs. [19,21] will survive these effects.

We should point out that, in the case of  $\gamma\gamma \rightarrow J \rightarrow \nu\bar{\nu}$ , no constraints are obtained from cosmology, since for a light Majoron its effect will appear after nucleosynthesis. If we had a heavier Majoron, this process would be

important in a supernovae although the strong absorption problem pointed out above should be investigated for each specific model to verify if it does not spoil the mechanism as in the pion case.

We conclude that only for very particular cases, e.g., a heavy pseudoscalar with mass larger than 1 keV without strong interactions with a dense medium, the reaction  $\gamma\gamma \rightarrow \nu\bar{\nu}$  will play a role in astrophysics.

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