

# The KeV Majoron as a Dark Matter Particle

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## ABSTRACT

We consider a very weakly interacting KeV majoron as a dark matter particle (DMP), which provides both the critical density  $\rho_{cr} = 1.88 \times 10^{-29} h^2 g/cm^3$  and the galactic scale  $M_{gal} \sim m_{Pl}^3/m_J^2 \sim 10^{12} M_\odot (m_J/1KeV)^{-2}$  for galaxy formation. The majoron couples to the leptons only through some new "directly interacting particles", called DIPS, and this provides the required smallness of the coupling constants. If the masses of these DIPS are greater than the scale  $V_s$  characterizing the spontaneous violation of the global lepton symmetry they are absent at the corresponding phase transition ( $T \sim V_s$ ) and the majorons are produced during the phase transition, never being in thermal equilibrium during the history of the universe. In the alternative case  $m_{DIP} < V_s$  the majorons can be for a short period in thermal equilibrium. This scenario is not forbidden by nucleosynthesis and gives a reasonable growth factor for the density fluctuations compatible with the recent restrictions from the COBE experiment. It provides as a possible signature the existence of an X-ray line at  $E_\gamma = \frac{m_J}{2}$ , produced by the decay  $J \rightarrow \gamma + \gamma$ . A particle physics model which provides the required smallness of the majoron couplings is described. It realizes the possibility of the KeV majoron as a DMP in a consistent way and may also lead to observable rates for flavour violating decays such as  $\mu \rightarrow e\gamma$  and  $\mu \rightarrow 3e$ , testable in the laboratory.

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# 1 Introduction

The basic scale among the Universe structures, the galactic mass,  $M_{gal} \sim 10^{12} M_{\odot}$ , needs a weakly interacting particle with mass  $m_X \sim (m_{Pl}^3/M_{gal})^{1/2} \sim 1$  KeV, where  $m_{Pl} = 1.2 \times 10^{19}$  GeV is the Planck mass. This particle must be electrically neutral and long-lived,  $\tau_x > t_0$  where  $t_0 = 2.06 \times 10^{17} h$  is the age of Universe. The majoron can be such a particle.

The KeV majoron has been recently discussed in the paper [1]. The novel element of this work is that the B-L symmetry is broken explicitly by the gravitational interaction considered as dimension 5 operators inversely proportional to the Planck mass  $m_{Pl}$ . This makes the majoron massive. The coupling constants in this model are such that the majoron is an unstable particle with the lifetime  $\tau_J \ll t_0$ . The basic parameter of the model, the vacuum expectation value  $V_{BL}$  of the complex scalar  $\sigma$  which splits into a heavy scalar field  $\rho$  and the pseudoscalar majoron  $J$ . Limits on  $V_{BL}$  were derived from the decay of the majoron, nucleosynthesis and galaxy formation (the COBE observations).

The cosmological aspects of the problem were studied in a more detailed way in ref. [2] and especially in ref. [3]. In particular the novel feature, the majoron strings, is discussed in the last work.

Some attention has also been given in these papers to the phenomenology of the long-lived majoron, though no elementary particle physics model for this case was specified.

In this letter following [1, 2, 3] we consider the singlet complex scalar  $\sigma$  with  $L \neq 0$  or  $B - L \neq 0$  as the "parent" field for the majoron. The L or B-L global symmetry is broken by  $\langle \sigma \rangle = V_s/\sqrt{2}$ . Further on we shall discuss the case  $V_s < V_{EW}$  (where  $V_{EW} \approx 250$  GeV) to avoid the problem of washing away the baryon asymmetry [4], see also the detailed discussion in [2]. As a result  $\sigma$  splits into the real scalar field  $\rho$  and the majoron according to  $\sigma = \frac{1}{\sqrt{2}}(V_s + \rho) \exp(iJ/V_s)$ .

This phase transition in the Universe occurs at temperature  $T_c \sim V_s$ , while at  $T < T_c$  the majorons did not exist (In the papers [1] and [2] there is an inaccurate statement that the majoron acquires mass at  $T \gg V_s$ ). We assume that the majorons interact directly with some heavy particles, X, with a "usual" coupling constant  $g_{JXX} \sim 10^{-2} - 10^{-3}$ , while the interaction with the ordinary particles (left-handed neutrinos,  $\nu$ , and charged leptons,  $l$ ) can only be mediated by X-particles and has effective coupling constants  $h_{\nu}$  and  $h_l$ , which are extremely small. They may arise either at the tree level or via radiative corrections as considered in the present paper. Further on we shall refer to the X-particles

as Directly Interacting Particles (DIPS).

## 2 Cosmological Scenario

There are several possible mechanisms for the production of the majorons in the early Universe.

The most natural mechanism is the thermal one [1]. For the case of a long-lived majoron,  $\tau_J > t_0$  where  $t_0 = 2.06 \times 10^{17} h^{-1}$  sec is the age of the Universe, the majoron must have extremely small effective coupling constants with the neutrinos and leptons. Therefore the majoron can be in thermal equilibrium only due to the interactions with DIP, such as right neutrinos [1] or charged particles as in the model considered below. This is possible for DIP masses smaller than the scale of the phase transition  $m_{DIP} < V_s$ . Then, at the time just after the phase transition,  $t \sim t_J$ , the space density of the majorons is

$$n_J(t_J)/n_\gamma(t_J) = 1/2 \quad (1)$$

The second plausible mechanism is connected with majoron strings, unstable due to evaporation of the majorons [2].

We would like to add to this list a third mechanism which operates effectively when  $m_{DIP} > V_s$ . In this case at the moment of the phase transition  $T_J \sim V_s$ , the DIPS are not present, while the interactions with other particles are too weak to keep the majoron in thermal equilibrium. The majorons are born in a phase transition at  $t = t_J$ , each majoron being accompanied by one  $\rho$ -boson.

Before the phase transition the electrically neutral part of the scalar potential connected with  $\sigma$ - field was

$$V(\sigma) = \lambda_s(\sigma^* \sigma - V_s^2/2)^2 \quad (2)$$

with  $\lambda_s \sim 1$ .

After phase transition,

$$\sigma = \frac{1}{\sqrt{2}}(\langle \sigma \rangle + \rho + iJ) \quad (3)$$

The potential energy of the  $\sigma$  field  $V(0)$  is converted mostly into the energy of the  $\rho$  bosons  $\rho_\rho(t_J) \sim \lambda V_s^4/4$  and therefore the space density of both  $\rho$  bosons and majorons becomes  $n_J(t_J) \sim n_\rho(t_J) \sim \lambda V_s^4/4m_\rho$ . If we denote  $m_\rho = V_s/\xi$ , where  $\xi \gtrsim 1$ , then using

$n_\gamma \sim T_s^3 \sim V_s^3$  one has for  $\lambda \sim 1$  that, at the moment of the phase transition,

$$\frac{n_J(t_J)}{n_\gamma(t_J)} \sim \xi \quad (4)$$

The density of majorons at  $t = t_0$  is given by

$$\frac{n_J(t_0)}{n_\gamma(t_0)} = \frac{n_J(t_J)}{n_\gamma(t_J)} \frac{43/11}{N_i} \quad (5)$$

where  $N_i$  is the number of degrees of freedom at  $t = t_J$  ( $N_i = 427/4$  for the particle content of the standard model and  $43/11$  is the ratio of degrees of freedom  $N(\gamma + e + \nu)N(\gamma)/N(\gamma + e)$ ).

Once produced at the phase transition the  $\rho$ -bosons will immediately decay to majorons. For example, the the width for the  $\rho$  decay to two majorons can be parametrized by

$$\Gamma(\rho \rightarrow JJ) = \frac{\sqrt{2}G_F}{32\pi} M_\rho^3 g_{\rho JJ}^2 \quad (6)$$

where the corresponding coupling  $g_{\rho JJ}$  is given in terms of the relevant electroweak and global symmetry breaking vacuum expectation values.

Therefore the present energy density of these majorons is given by

$$\frac{\rho_J(t_0)}{\rho_{cr}(t_0)} = 1.5\xi h^{-2} m_J \quad (7)$$

where  $\xi = 1/2$  for the thermal mechanism of production and  $\xi > 1$  for the production at the phase transition. Here and everywhere else we  $m_J$  is the mass of the majoron in KeV.

The majorons are also coupled (through DIPS) both to neutrinos as well as to the charged leptons with effective coupling constants which we denote by  $h_\nu$  and  $h_l$ , respectively. The coupling  $h_\nu$  determines the majoron lifetime as

$$\tau_{J \rightarrow \nu\nu} = \frac{16\pi}{h_\nu^2 m_J} \quad (8)$$

In order to be DMP the majoron lifetime must be longer than the age of the universe.

This translates into a constraint on its effective coupling  $h_\nu$

$$h_\nu \leq 1.3 \times 10^{-17} \left( \frac{m_J}{1\text{KeV}} \right)^{-1/2} h \quad (9)$$

On the other hand the couplings to charged leptons,  $h_l$ , are constrained by  $J \rightarrow \gamma\gamma$  decays. Indeed, in the most general case the current connected with the global symmetry whose spontaneous violation gives rise to the majoron is anomalous, and therefore the majoron is coupled to photons through the electromagnetic anomaly of this symmetry. The effective Lagrangean has a form

$$\mathcal{L}_{int}^l = g_{J\gamma\gamma}^l J F_{\mu\nu} \widetilde{F}_{\mu\nu}, \quad (10)$$

where  $g_{J\gamma\gamma}^l$  arises from the diagrams shown in Fig. 1. The effective triangle loop gives

$$g_{J\gamma\gamma}^l = \frac{\alpha_{em}}{2\pi} \frac{h_l}{m_l} \quad (11)$$

and the majoron lifetime

$$\tau_{J\gamma\gamma} = \frac{64\pi^3}{\alpha_{em}^2 h_l^2} \left(\frac{m_l}{m_J}\right)^2 \frac{1}{m_J} \quad (12)$$

The contributions from two-loop diagrams mediated by electrically charged scalar bosons, as in Fig. 1(b), may be of similar magnitude.

The value of  $h_l$  can be roughly constrained by the condition that the energy density of the produced photons must not exceed the observed energy density ( $\sim 10^{-5} \text{eV}/\text{cm}^3$ ) for  $E \lesssim 1 \text{KeV}$ . Using  $\omega_X \sim \frac{\rho_{cr} t_0}{\tau_{J\gamma\gamma}}$  one obtains

$$h_\tau \lesssim 6 \times 10^{-13} h^{-1} \left(\frac{m_J}{1 \text{KeV}}\right)^{-3/2} \quad (13)$$

Notice that for the case  $h_l \sim m_l$  (as in the model discussed below) the coupling to electrons

$$h_e = 1.7 \times 10^{-16} \quad (14)$$

is much smaller than the observational limit from stellar cooling  $h_e \lesssim 3 \times 10^{-13}$  [5].

We now turn to other cosmological requirements for the majoron to be DMP, beyond those given in eq. (9) and eq. (13).

The restriction imposed by big-bang nucleosynthesis has already been discussed in refs [2, 3]. The number of degrees of freedom associated to the majorons at the time of the nucleosynthesis is smaller than the observational limit

$$N = \frac{43/11}{N_i} = 3.7 \times 10^{-2} \quad (15)$$

Let us now discuss the problem of galaxy formation in the light of the COBE results. The basic scale for the majoron corresponds to the galactic Jeans mass

$$M_{gal} \approx m_{Pl}^3/m_J^2 \sim 1.6 \times 10^{12} M_\odot (m_J/1 \text{KeV})^{-2} \quad (16)$$

The corresponding linear scale  $\lambda_{gal} = 2.9 \times 10^{24} h^{-2/3} \text{cm}$ . It enters the horizon at  $t = t_H = 2.8 \times 10^8 h^{-4/3} \text{sec}$ <sup>3</sup>. At this moment the amplitude of fluctuation is equal to

$$\left(\frac{\delta\rho}{\rho}\right)_{\lambda_{gal}} = \delta_H \quad (17)$$

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<sup>3</sup>We use for the equilibrium time  $t_{eq} = 4.49 \times 10^{10} (\Omega_0 h^2)^{-2} T_{2.73}^6 \text{sec}$ . and  $1+z_{eq} = 2.77 \times 10^4 (\Omega_0 h^2) T_{2.73}^{-4}$ , calculated for two light and one heavy neutrinos.

which is the same for all scales entering the horizon, for the Harrison-Zeldovich spectrum. From the moment  $t_H$  up to  $t = t_{eq}$  the amplitude grows as [6]

$$\delta_{\lambda_{gal}}(t) \equiv \left(\frac{\delta\rho}{\rho}\right)_{\lambda_{gal}}(t) = \delta_H(1 + A \ln t/t_H), \quad (18)$$

where  $A$  can be found from the value of  $\delta_{\lambda_{gal}}(t)$  at  $t \leq t_H$ . Since for this interval  $\delta_{\lambda_{gal}}(t) = \delta_H(t/t_H)$ , equating the derivatives at  $t = t_H$  one finds  $A = 1$ . From  $\delta(t) \sim t^{2/3}$  for  $t \geq t_H$  it follows that

$$\left(\frac{\delta\rho}{\rho}\right)_k(t_0) = \delta_H(1 + \ln t_{eq}/t_H)(t/t_{eq})^{2/3} \quad (19)$$

where  $k = 2\pi/\lambda_{gal}$  is the comoving wave number.

We now calculate the value of  $\delta_H$  from the quadrupole momentum of microwave radiation found in the COBE experiment [7]  $Q = 17 \times 10^{-6}$  K for the case of the Harrison-Zeldovich spectrum. As shown in ref. [8]  $(\frac{\delta\rho}{\rho})_\lambda$  for the comoving scale  $\lambda$  is related to the dimensionless quadrupole momentum  $a_2$  as

$$\left(\frac{\delta\rho}{\rho}\right)_\lambda = \left(\frac{108}{\pi}\right)^{1/2} a_2 \left(\frac{cH^{-1}}{\lambda}\right)^2 I(\lambda) \quad (20)$$

where  $H$  is the Hubble constant,  $I(\lambda) = 1/6$  for large scales  $\lambda$ , and  $a_2$  is connected to  $Q$  as

$$a_2^2 = \frac{4\pi}{5} \frac{Q^2}{T^2} \quad (21)$$

Using  $(\frac{\delta\rho}{\rho})_\lambda = \delta_H$  for  $\lambda = \frac{2}{3}cH^{-1}$  we obtain

$$\delta_H = 2.15 \times 10^{-5} \quad (22)$$

Finally, for the galactic scale  $\lambda_{gal}$  we obtain from eq. (19)  $(\frac{\delta\rho}{\rho})_{\lambda_{gal}}(t_0) = 3.6$  for  $h=1$  and  $(\frac{\delta\rho}{\rho})_{\lambda_{gal}}(t_0) = 1.2$  for  $h=0.5$ , values which are sufficient for entering the nonlinear stage of the perturbation's growth.

### 3 A Particle Physics Model

From the above discussion we see the need to have very small couplings of the majoron to neutrinos (so as to provide  $\tau_J > t_0$ ) as well as to charged leptons (so as to avoid too many photons from the  $J \rightarrow \gamma\gamma$  decay). These constraints are summarized in eq. (9) and eq. (13). Here we demonstrate that a realistic model can meet such requirements in a natural way. In our example the required smallness of the couplings of the majoron to neutrinos and charged leptons arises from their radiative origin. This is achieved without the need to introduce new mass scales, much larger than that of the standard electroweak model.

The model contains, in addition to the Standard Model particles, one singly charged singlet scalar  $\eta^-$  [9] and one doubly charged singlet scalar  $\chi^{--}$  [10]. It is characterized by the following lepton Yukawa interactions [11]

$$\mathcal{L} = -\frac{\sqrt{2}m_i}{v}\bar{\ell}_i\phi e_{Ri} + f_{ij}\ell_i^T C i\tau_2 \ell_j \eta^+ + h_{ij}e_{Ri}^T C e_{Rj}\chi^{++} + H.c. \quad (23)$$

( $h$  and  $f$  are symmetric and anti-symmetric coupling matrices, respectively). In addition, the scalar boson lagrangean is given by [12, 13]

$$\begin{aligned} V_1 = & -\frac{1}{2}\mu_\phi^2\phi^\dagger\phi + \frac{1}{4}\lambda_\phi(\phi^\dagger\phi)^2 + \frac{1}{2}\mu_\chi^2\chi^\dagger\chi + \frac{1}{4}\lambda_\chi(\chi^\dagger\chi)^2 + \frac{1}{2}\mu_\eta^2\eta^\dagger\eta + \frac{1}{4}\lambda_\eta(\eta^\dagger\eta)^2 \\ & + \frac{1}{2}\epsilon_{\phi\chi}\phi^\dagger\phi\chi^\dagger\chi + \frac{1}{2}\epsilon_{\phi\eta}\phi^\dagger\phi\eta^\dagger\eta + \frac{1}{2}\epsilon_{\eta\chi}\eta^\dagger\eta\chi^\dagger\chi + \epsilon\sigma\chi^{++}\eta^-\eta^- + H.c. \end{aligned} \quad (24)$$

where  $\epsilon$  and  $\lambda$  are dimensionless coupling constants and  $\mu$ 's are mass parameters.

To this we add

$$V_2 = -\frac{1}{2}\mu_\sigma^2\sigma^\dagger\sigma + \frac{1}{4}\lambda_\sigma(\sigma^\dagger\sigma)^2 + \epsilon_{\phi\sigma}(\phi^\dagger\phi)(\sigma^\dagger\sigma) + \epsilon_{\chi\sigma}(\chi^\dagger\chi)(\sigma^\dagger\sigma) + \epsilon_{\eta\sigma}(\eta^\dagger\eta)(\sigma^\dagger\sigma) \quad (25)$$

Minimizing the total potential one finds the Higgs boson vacuum expectation value responsible for electroweak breaking, given by

$$\langle\phi\rangle = \frac{\mu_\phi}{\sqrt{\lambda_\phi}} \quad (26)$$

In addition, the complex singlet  $\sigma$  also acquires a nonzero vacuum expectation value  $\langle\sigma\rangle$  which breaks the global symmetry, while the  $\chi$  and  $\eta$  mass terms are chosen positive in such a way that the electrically charged Higgs bosons do not acquire vacuum expectation values, as required.

The charged lepton masses are generated when the  $SU(2) \otimes U(1)$  symmetry is broken by  $\langle\phi\rangle$ , through the first term. On the other hand, neutrinos acquire masses only radiatively, at the two loop level, by the diagram in Fig. 2 [11]. This follows directly from the spontaneous violation of the lepton number symmetry by the nonzero vacuum expectation value of the  $\sigma$  field.

Notice that in the model summarized above the majoron is introduced at the electroweak scale. Many variant models of this type can be considered [14]. Such models have recently been considered in more detail in connection with the question of electroweak baryogenesis [13] and Higgs boson physics [14]. In all of them, the presence of the neutral singlet  $\sigma$  provides two additional electrically neutral degrees of freedom: an extra CP-even Higgs boson (corresponding to the  $\rho$  field) and a CP-odd scalar boson, i.e. the majoron,  $J$ , which remains massless as a result of the spontaneous nature of the lepton number

breaking. Both CP-even mass eigenstate scalar bosons  $H_1$  and  $H_2$  (corresponding to the standard model Higgs boson and the extra one coming from  $\rho$ ) can decay by majoron emission, either directly, such as  $H_1 \rightarrow JJ$  and  $H_2 \rightarrow JJ$  or indirectly, via  $H_2 \rightarrow H_1 H_1$ . As long as the scale at which the lepton number breaking is sufficiently low, as assumed here, these dark Higgs decay modes will be substantial, suppressing the standard Higgs decays, such as into  $b\bar{b}$ . As discussed above it also implies that the  $\rho$  bosons produced at the phase transition immediately decay to the majorons that form the dark matter.

The implications of this model have recently attracted a lot of attention [14, 15, 16]. In addition to the presence of the invisible Higgs decay modes the production of the two CP-even Higgs bosons through the Bjorken process  $e^+e^- \rightarrow ZH_i$  will be smaller than predicted in the standard model thus weakening the LEP limit on Higgs boson mass [15]. An implication of this fact is that this model bypasses the conflict with electroweak baryogenesis [13] which requires  $m_H \lesssim 40$  GeV [17].

In what follows we show how this model provides a way where all the required conditions eq. (9), eq. (13) can be realized.

First note that the majoron coupling to neutrinos arises from the same two-loop diagram as Fig. 2. An estimate of this graph gives [18]

$$h_\nu^{ij} \approx \frac{\epsilon \sum_{a,b} f_{ai} f_{jb} h_{ab} m_a m_b}{256\pi^4 M_0^2} \quad (27)$$

For natural choices of parameters, consistent with all present observations, e.g.  $f_{e\tau}, f_{\mu\tau}, h_{\tau\tau} \sim 10^{-2}$ ,  $\epsilon^2 \sim 10^{-5}$ , and charged Higgs boson masses of about 200 GeV, these  $h_\nu$  couplings have a typical magnitude  $h_\nu \sim 10^{-17}$  which is quite compatible with what is required by eq. (9).

On the other hand the corresponding neutrino masses can be written as

$$m_\nu = h_\nu \langle \sigma \rangle$$

and lie in the  $10^{-6}$  eV range or below, for  $\langle \sigma \rangle \simeq 100$  GeV. Such small neutrino masses might play a role in the understanding of the solar neutrino deficit.

On the other hand, the majoron coupling to charged leptons arises from the two-loop diagrams shown as Fig. 3. Using the estimate from ref. [12]

$$h_l \sim \frac{1}{16\pi^2} \frac{\epsilon^2 \langle \phi \rangle m_l}{M^2} \sum_k (f_{lk}^2 + 2h_{lk}^2) \quad (28)$$

one finds that, for reasonable choices of parameters,  $h_l$  values can lie in a range quite consistent with what is required by eq. (13). For example, for the choice of parameters described above, one gets for  $h_l \sim 2 \times 10^{-13}$  for the case  $l = \tau$ .



Note that, although both  $h_\nu$  and  $h_l$  arise at the same two-loop level, there is enough freedom in the choice of model parameters to obtain the required hierarchy  $h_l \gg h_\nu$ .

As a final comment we note that in this model the decay  $J \rightarrow \gamma\gamma$  also receives contributions from two-loop diagrams mediated by electrically charged scalar bosons (Fig. 1(b)). A rough estimate shows that it may be of similar importance as the effective triangle graph of Fig. 1(a).

## 4 Conclusions

In this paper we outlined the possibility of the KeV majoron playing the role of DMP.

In order to be DMP the majoron lifetime must be longer than the age of the universe and its effective couplings with neutrinos and charged leptons,  $h_\nu$  and  $h_l$ , respectively, must be rather small. This can be naturally realized in models where the majoron interacts with these particles only radiatively, via the exchange of DIPS. We have shown explicitly that in this case these constants can be very small, as in the model discussed in section 3.

We considered two possible scenarios for the cosmological production of the majorons. If the DIPS are lighter than  $V_s$ ,  $m_{DIP} < V_s$ , the majorons can be for a short time period in thermal equilibrium. In the alternative case the DIPS are heavier than  $V_s$  and therefore they are not present at the phase transition ( $T \sim V_s$ ) when the majorons are born. Since both effective couplings,  $h_\nu$  and  $h_l$  are very small, in this case never in the history of the Universe were the majorons in thermal equilibrium.

As we demonstrated, this scenario is not in conflict with nucleosynthesis data and results in a reasonable growth factor for the density perturbations in agreement with the limits of the COBE experiment.

A possible observational signature for our scenario is the existence of a monochromatic line  $E_\gamma = 0.5m_J$  KeV in the extragalactic diffuse flux. Its intensity is determined by the coupling  $h_\tau$  of the majoron to the tau lepton and equals to

$$I_\gamma = \frac{1}{4\pi} \frac{\rho_{cr}}{m_J} \frac{ct_0}{\tau_{J\gamma\gamma}} \simeq 6m_J^2 \left( \frac{h_\tau}{3 \times 10^{-13}} \right)^2 \text{ cm}^{-2} \text{ sec}^{-1} \text{ sr}^{-1} \quad (29)$$

This line may be detected also from the galactic center and nearby cluster of galaxies. We shall discuss this problem somewhere else. As we have already mentioned, the decay  $J \rightarrow \gamma\gamma$  due to electrically charged scalar boson loops (Fig. 1(b)) can compete with the process discussed above.

A particle physics signature for this model (section 3) consists in the possibility of observable rates for flavour violating decays such as  $\mu \rightarrow e\gamma$  and  $\mu \rightarrow 3e$ . The first decay process arises at one-loop, due to singly-charged scalar boson exchange, while the second is present already at the tree level, due to  $\chi^{++}$  exchange. These processes could be tested at laboratory experiments such as MEGA or at PSI.

This work was supported by DGICYT under grant PB92-0084 and by Generalitat Valenciana. We thank A. Dolgov, I. Khriplovich, J. Peltoniemi and A. Smirnov for valuable discussions.

## Figure Captions

**Fig. 1.**

Diagrams leading to the decay  $J \rightarrow \gamma\gamma$ : (a) Effective triangle loop contributions; (b) Scalar boson mediated contributions.

**Fig. 2.**

Diagram generating the small coupling  $h_\nu$  of the majoron to neutrinos. It also generates very small Majorana neutrino masses

**Fig. 3.**

Diagram generating the coupling  $h_l$  of the majoron to charged leptons.

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