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Astrophysical and terrestrial constraints on singlet Majoron models

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The general Lagrangian containing the couplings of the Higgs scalars to Majorana neutrinos is presented in the context of singlet Majoron models with intergenerational mixings. The analytical expressions for the coupling of the Majoron field to charged fermions are derived within these models. Astrophysical considerations imply severe restrictions on the parameters of the three-generation Majoron model if the Dirac neutrino mass matrix of the model follows a mass hierarchical pattern dictated by grand unified theories. Bounds that originate from analyzing possible charged lepton-violating decays in terrestrial experiments are also discussed. In particular, we find that experimental searches for muon decays by Majoron emission cannot generally be precluded by astrophysical requirements.

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Astrophysical considerations play an important role in constraining the strength of the coupling of Nambu-Goldstone bosons to matter [1]. Among the various types of such extraordinary light particles (e.g., axions, familons, etc.) [2], which are associated with the spontaneous breakdown of some global symmetry, the Majoron J^0 is a massless pseudoscalar boson arising from the breaking of the baryon-lepton (B-L) symmetry [3–7]. In such scenarios, apart from the standard model (SM) Higgs doublet Φ, an $SU(2)_L \otimes U(1)_Y$ singlet Σ is present which gives rise to $\Delta L = 2$ Majorana mass terms m_M when Σ couples to right-handed neutrinos [3]. Therefore, models with right-handed neutrinos can naturally account for possible lepton-number-violating decays of the Z^0 and Higgs particle [8–11]. Such decays, being forbidden in the minimal SM, can be induced by Majorana neutrinos at the first electroweak loop order. The one-loop vertex function relevant for these lepton-flavor violating decays shows a strong quadratic dependence of the heavy neutrino mass. This nondecoupling physics originating from heavy Majorana neutrinos leads to combined constraints both on neutrino masses and lepton-violating mixings [8,9,11]. Similar nondecoupling effects will, however, occur when one considers the Majoron coupling to charged fermions.

If Majoron models are, in some way, embedded in grand unified theories (GUT's) such as the SO(10) model [12], the Dirac neutrino mass matrix m_D may then be related to the u-quark mass matrix M_U by $m_D = M_U/k$, where $k \sim 1$ represents the running of the Yukawa couplings between the GUT and the low-energy scale [13]. There are also scenarios where m_D can be proportional to the charged lepton mass matrix M_l [12,13]. On the other hand, the B-L scale which is determined by the vacuum expectation value (VEV) of the singlet scalar, i.e., $\langle \Sigma \rangle = w/\sqrt{2}$, depends strongly on the mechanism that the SO(10) gauge group breaks down to $U(1)_{em}$ [14,4] and it has generally the tendency to be much higher than 10 TeV. However, it has been shown in [15] that nonperturbative Planck scale effects on singlet Majoron models can limit w to be less than about 1–10 TeV. This advocates the treatment of regarding w as a free parameter of the theory that should be constrained by our forthcoming considerations. In singlet Majoron models, the masses m_{ν_i} of the ordinary neutrinos can be estimated by the known "seesaw" relation

$$m_{\nu_i} \simeq -m_D \frac{1}{m_M} m_D^T \,, \tag{1}$$

where m_M dictates the mass scale of the heavy Majorana neutrinos N_i.

The mass hierarchical pattern mentioned above should, by analogy, hold for the heaviest family of neutrinos and up-type quarks. This implies that the biggest eigenvalue of m_D will be in the range between 10 and 100 GeV for $100 \le m_t \le 180$ GeV [13]. If astrophysical constraints are now imposed on the J^0-e-e, J^0-u-u, and J^0-d-d couplings, we will then find an upper bound of the parameter,

$$\tan\beta = \frac{\langle \Phi \rangle}{\langle \Sigma \rangle} = \frac{v}{w} \le 10^{-2} \,, \tag{2}$$

for models without interfamily mixings. The situation becomes more involved if mixings between families are considered. This realization will also be illustrated by the present work. Finally, we will discuss the bounds from low-energy experiments [16,17] on the off-diagonal coupling of the Majoron to two different charged leptons.

First, let us briefly describe the low-energy structure of the singlet Majoron model. The scalar potential of this model, which should be invariant under the $SU(2)_L \otimes U(1)_Y$ and $U(1)_{B\text{-}L}$ group, is given by [5,6]

$$-\mathcal{L}_V = \mu_\Phi^2(\Phi^\dagger\Phi) + \mu_\Sigma^2(\Sigma^\dagger\Sigma) + \frac{\lambda_1}{2}(\Phi^\dagger\Phi)^2 + \frac{\lambda_2}{2}(\Sigma^\dagger\Sigma)^2 + \delta(\Phi^\dagger\Phi)(\Sigma^\dagger\Sigma) \,. \tag{3}$$

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If all the stability conditions in this model are satisfied (i.e., $\lambda_1, \lambda_2 > 0$, and $\lambda_1\lambda_2 > \delta$) [5], the above potential can always be minimized by the Higgs-field configurations

$$\Phi = \begin{bmatrix} G^+ \\ \frac{v}{\sqrt{2}} + \frac{\phi^0 + iG^0}{\sqrt{2}} \end{bmatrix}$$

and

$$\Sigma = \frac{w}{\sqrt{2}} + \frac{\sigma^0 + iJ^0}{\sqrt{2}} . \tag{4}$$

After the spontaneous breakdown of the $SU(2)_L \otimes U(1)_Y$ gauge symmetry and the anomaly-free global symmetry $U(1)_{B\text{-}L}$, the diagonalization of the resulting Higgs boson mass matrix yields two CP-even Higgs fields (denoted by H^0 and S^0) and one physical massless CP-odd scalar, the Majoron J^0. The weak eigenstates ϕ^0 and σ^0 are then related to the corresponding physical mass eigenstates through the orthogonal transformation

$$\begin{bmatrix} \phi^0 \\ \sigma^0 \end{bmatrix} = \begin{bmatrix} \cos\theta & -\sin\theta \\ \sin\theta & \cos\theta \end{bmatrix} \begin{bmatrix} H^0 \\ S^0 \end{bmatrix} , \tag{5}$$

where

$$\tan 2\theta = -\frac{2\delta \tan\beta}{\lambda_2 - \lambda_1 \tan^2\beta} . \tag{6}$$

The Yukawa sector containing all the relevant Higgs couplings to neutrinos reads

$$\begin{aligned} -\mathcal{L}_Y^{\text{Higgs}} = {} & \bar{\nu}_{L_i}^0 m_{D_{ij}} \nu_{R_j}^0 \frac{\phi^0}{v} + \bar{\nu}_{R_i}^0 m_{D_{ij}}^\dagger \nu_{L_j}^0 \frac{\phi^0}{v} \\ & + \tfrac{1}{2} \bar{\nu}_{R_i}^{0C} m_{M_{ij}} \nu_{R_j}^0 \frac{\sigma^0 + iJ^0}{w} \\ & + \tfrac{1}{2} \bar{\nu}_{R_i}^0 m_{M_{ij}}^\dagger \nu_{R_j}^{0C} \frac{\sigma^0 - iJ^0}{w} . \end{aligned} \tag{7}$$

In Eq. (7) we have assumed the absence of Higgs triplets [18], which seem now to be ruled out by the present from the CERN data e^+e^- collider LEP on the invisible Z^0 width. The interactions of J^0, H^0, and S^0 with the $2n_G$ Majorana neutrinos n_i (n_G denotes the number of generations) are generally described by the Lagrangians

$$\begin{aligned} \mathcal{L}_{\text{int}}^J = -\frac{i g_W t_\beta}{4M_W} J^0 \bar{n}_i [& \gamma_5 (m_{n_i} + m_{n_j}) (\tfrac{1}{2}\delta_{ij} - \text{Re}C_{ij}) \\ & + i(m_{n_i} - m_{n_j}) \text{Im}C_{ij}] n_j , \end{aligned} \tag{8}$$

$$\begin{aligned} \mathcal{L}_{\text{int}}^H = {} & -\frac{g_W}{4M_W} (c_\theta - s_\theta t_\beta) H^0 \bar{n}_i \\ & \times \left[(m_{n_i} + m_{n_j}) \left[\text{Re}C_{ij} + \frac{t_\beta s_\theta \delta_{ij}}{2(c_\theta - s_\theta t_\beta)} \right] \right. \\ & \left. + i\gamma_5 (m_{n_j} - m_{n_i}) \text{Im}C_{ij} \right] n_j , \end{aligned} \tag{9}$$

$$\begin{aligned} \mathcal{L}_{\text{int}}^S = {} & \frac{g_W}{4M_W} (s_\theta + c_\theta t_\beta) S^0 \bar{n}_i \\ & \times \left[(m_{n_i} + m_{n_j}) \left[\text{Re}C_{ij} - \frac{t_\beta \delta_{ij}}{2(t_\theta + t_\beta)} \right] \right. \\ & \left. + i\gamma_5 (m_{n_j} - m_{n_i}) \text{Im}C_{ij} \right] n_j , \end{aligned} \tag{10}$$

where we have used the abbreviations $s_x = \sin x$, $c_x = \cos x$, $t_x = \tan x$, and defined

$$C_{ij} = \sum_{k=1}^{n_G} U_{ki}^\nu U_{kj}^{\nu *} . \tag{11}$$

The $2n_G \times 2n_G$ unitary matrix U^ν is responsible for the diagonalization of the $2n_G \times 2n_G$ neutrino mass matrix M^ν which is of the "seesaw" form [19]

$$M^\nu = \begin{bmatrix} 0 & m_D \\ m_D^T & m_M \end{bmatrix} . \tag{12}$$

The first n_G eigenvalues of M^ν are identified with the ordinary light neutrinos, ν_e, ν_μ, etc., whereas the remaining n_G Majorana neutrino states N_i are new particles provided by the model and should be heavier than the Z^0 boson in order to escape detection at LEP experiments. The $2n_G$ neutral leptons n_i are related to their weak eigenstates ν_{L,R_i}^0 and ν_{L,R_i}^{0C} through the following unitary transformations (assuming the convention of summation for repeated indices):

$$\begin{bmatrix} \nu_L^0 \\ \nu_R^{0C} \end{bmatrix}_i = U_{ij}^{\nu *} n_{L_j} , \quad \begin{bmatrix} \nu_L^{0C} \\ \nu_R^0 \end{bmatrix}_i = U_{ij}^\nu n_{R_j} . \tag{13}$$

It is now easy to see that in the limit of $\tan\beta \to 0$ [implying also $\theta \to 0$ because of Eq. (6)], the fields S^0 and J^0 decouple fully from matter and only one Higgs field H^0 couples to Majorana neutrinos. This scenario has explicitly been discussed in [20], where for our purposes we will repeat here the interactions of the W^+ and Z^0 boson with the Majorana neutrinos. They are given by the Lagrangians

$$\mathcal{L}_{\text{int}}^{W^\mp} = -\frac{g_W}{2\sqrt{2}} W^{-\mu} \bar{l}_i B_{l_i j} \gamma_\mu (1 - \gamma_5) n_j + \text{H.c.} , \tag{14}$$

$$\mathcal{L}_{\text{int}}^Z = -\frac{g_W}{4\cos\theta_W} Z^{0\mu} \bar{n}_i \gamma_\mu [i\,\text{Im}C_{ij} - \gamma_5 \text{Re}C_{ij}] n_j . \tag{15}$$

Moreover, the interaction Lagrangian of the charged would-be Goldstone bosons $G^\mp$ with the Majorana neutrinos is written down as

$$\begin{aligned} \mathcal{L}_{\text{int}}^{G^\mp} = -\frac{g_W}{2\sqrt{2}M_W} G^- \bar{l}_i [& m_{l_i} B_{l_i j} (1 - \gamma_5) \\ & - B_{l_i j} (1 + \gamma_5) m_{n_j}] n_j + \text{H.c.} , \end{aligned} \tag{16}$$

where

$$B_{l_i j}=\sum_{k=1}^{n_G} V^l_{l_i k} U^{\nu *}_{kj} \ . \tag{17}$$

In Eq. (17) V^l is a unitary matrix relevant for the diagonalization of the charged lepton mass matrix M_l.

At this stage it is important to mention that the presence of Majorana neutrino interactions in the Lagrangians (8)–(10) violates the CP symmetry of the model. In particular, the fact that H^0, S^0, and J^0 couple simultaneously to CP-even $(:\bar{n}_i n_j)$ and CP-odd $(:\bar{n}_i i\gamma_5 n_j)$ operators gives rise to CP-violating transitions between states with different CP quantum numbers [21]. For example, one finds a nonzero contribution when computing self-energy graphs induced by Majorana neutrinos between the CP-even Higgs fields H^0, S^0, and the CP-odd state Z^0. All these transitions turn out to be proportional to the CP-odd combinations $\mathrm{Im}C^2_{ij}$ which change sign when a CP conjugation is applied to the vacuum polarization terms. As an additional byproduct of the violation of the CP symmetry, the general Majoron coupling to two *different* charged leptons can simultaneously possess a scalar and pseudoscalar part.

Armed with the Lagrangians (8)–(10), (14), and (15), it is now straightforward to calculate the coupling $J^0-f_1-f_2$ given by the Feynman graphs shown in Figs. 1(a)–1(c), where $f_{1,2}$ denote charged fermions, i.e., the charged leptons (l_i), e, μ, τ, and the u and d quark. Additional CP-violating diagrams of the Majoron coupling to fermions are depicted in Figs. 2(d) and 2(e). In our analytical calculations we have neglected terms proportional to the small quantities m_f^2/M_W^2. The individual amplitudes contributing to the $J^0-f_1-f_2$ coupling are given by

$$\begin{aligned}\mathcal{T}^{l_1 l_2}_a=\Delta^S_{ij}\Big[&-\tfrac{1}{2}(\lambda_i+\lambda_j)(\delta_{ij}-C^*_{ij})I_1(\lambda_i,\lambda_j)+C_{ij}\sqrt{\lambda_i\lambda_j}I_1(\lambda_i,\lambda_j)+\tfrac{1}{2}(\lambda_i-\lambda_j)C^*_{ij}I_3(\lambda_i,\lambda_j)\Big]\\ &+\tfrac{1}{2}\Delta^A_{ij}(\lambda_i-\lambda_j)C^*_{ij}[I_2(\lambda_i,\lambda_j)-I_1(\lambda_i,\lambda_j)] \ ,\end{aligned}\tag{18}$$

$$\begin{aligned}\mathcal{T}^{l_1 l_2}_b=\tfrac{1}{2}\Delta^S_{ij}\{&(C_{UV}-\tfrac{1}{2})[\lambda_i\delta_{ij}-C_{ij}\sqrt{\lambda_i\lambda_j}-\tfrac{1}{2}(\lambda_i+\lambda_j)C^*_{ij}]-(\lambda_i+\lambda_j)(\delta_{ij}-C^*_{ij})L_2(\lambda_i,\lambda_j)\\ &+C_{ij}\sqrt{\lambda_i\lambda_j}[-2L_2(\lambda_i,\lambda_j)+\tfrac{1}{2}(\lambda_j-\lambda_i)I_3(\lambda_i,\lambda_j)]\}\\ &+\tfrac{1}{2}\Delta^A_{ij}(\lambda_i-\lambda_j)[\tfrac{1}{2}(C_{UV}-\tfrac{1}{2})C^*_{ij}-C_{ij}\sqrt{\lambda_i\lambda_j}I_2(\lambda_i,\lambda_j)+C^*_{ij}L_2(\lambda_i,\lambda_j)] \ ,\end{aligned}\tag{19}$$

$$\mathcal{T}^{ff}_c=-2\delta^S(2T^f_z)[\lambda_j C_{ij}(\delta_{ij}-C^*_{ij})-\sqrt{\lambda_i\lambda_j}\mathrm{Re}C^2_{ij}][-\tfrac{1}{2}C_{UV}+L_1(\lambda_i,\lambda_j)] \ , \tag{20}$$

$$\mathcal{T}^{ff}_d=-i\delta^A\mathrm{Im}C^2_{ij}c_\theta(c_\theta-s_\theta t_\beta)(\lambda_i-\lambda_j)\sqrt{\lambda_i\lambda_j}\lambda_H^{-1}[C_{UV}-L_0(\lambda_i,\lambda_j)] \ , \tag{21}$$

$$\mathcal{T}^{ff}_e=-i\delta^A\mathrm{Im}C^2_{ij}s_\theta(s_\theta+c_\theta t_\beta)(\lambda_i-\lambda_j)\sqrt{\lambda_i\lambda_j}\lambda_S^{-1}[C_{UV}-L_0(\lambda_i,\lambda_j)] \ , \tag{22}$$

where

$$\lambda_i=\frac{m^2_{n_i}}{M^2_W} \ , \quad \lambda_H=\frac{M^2_H}{M^2_W} \ , \quad \lambda_S=\frac{M^2_S}{M^2_W} \ , \tag{23}$$

$$C_{UV}=\frac{1}{\varepsilon}-\gamma_E+\ln 4\pi-\ln\frac{M^2_W}{\mu^2} \ , \tag{24}$$

$$\begin{aligned}\Delta^A_{ij}&=\frac{g_W\alpha_W}{16\pi}\tan\beta B^*_{l_1 i}B_{l_2 j}\bar{u}_{l_2}\left[\frac{m_{l_1}}{M_W}(1+\gamma_5)+\frac{m_{l_2}}{M_W}(1-\gamma_5)\right]u_{l_1} \ ,\\ \Delta^S_{ij}&=\frac{g_W\alpha_W}{16\pi}\tan\beta B^*_{l_1 i}B_{l_2 j}\bar{u}_{l_2}\left[\frac{m_{l_1}}{M_W}(1+\gamma_5)-\frac{m_{l_2}}{M_W}(1-\gamma_5)\right]u_{l_1} \ ,\end{aligned}\tag{25}$$

$$\begin{aligned}\delta^A&=\frac{g_W\alpha_W}{16\pi}\frac{m_f}{M_W}\tan\beta\bar{u}_f u_f \ ,\\ \delta^S&=\frac{g_W\alpha_W}{16\pi}\frac{m_f}{M_W}\tan\beta\bar{u}_f\gamma_5 u_f \ .\end{aligned}\tag{26}$$

The analytical expressions for the one-loop functions I_1, I_2, I_3, L_0, L_1, and L_2 are given in the Appendix A. In Eq. (20), T^f_z stands for the third component of the weak isospin of the charged fermion (f) and takes the values $T^u_z=\frac{1}{2}$, $T^{l,d}_z=-\frac{1}{2}$. The UV divergences in the amplitudes (19)–(22) vanish identically due to the equalities [20,9]

$$\sum_{i=1}^{2n_G} B_{li}C_{ij}=B_{lj} \ , \tag{27}$$

$$\sum_{k=1}^{2n_G} C_{ik}C^*_{jk}=C_{ij} \ , \tag{28}$$

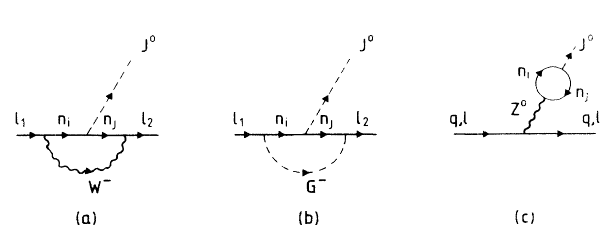

FIG. 1. Feynman graphs responsible for the coupling of the Majoron to charged fermions, $J^0-f_1-f_2$.

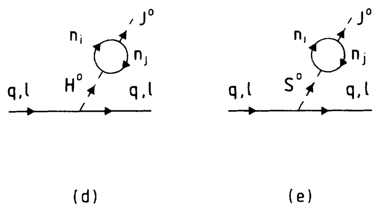

FIG. 2. Additional CP-odd graphs giving a vanishing contribution to the coupling J^0-f-f due to Eq. (30).

$$\sum_{i=1}^{2n_G} m_{n_i} B_{li} C_{ij}^{*} = 0 \ , \tag{29}$$

$$\sum_{k=1}^{2n_G} m_{n_k} C_{ik} C_{jk} = 0 \ . \tag{30}$$

Note also that the scalar J^0-f-f coupling given by the amplitudes $\mathcal{T}_d^{ff}$ and $\mathcal{T}_e^{ff}$ vanishes *completely* after doing some algebra due to the identity (30). We have also checked that similar cancellations of scalar pieces occur in the diagonal coupling J^0-l-l in Eqs. (18) and (19) on account of Eq. (29). This is a nice reciprocity between the Goldstone theorem and the interactions of the Majoron field to fermions which must be of derivative type [2]. As a result, the diagonal coupling of the Majoron to fermions should be of a pure *pseudoscalar* nature.

In order to pin down numerical predictions, we first consider the conservative case of a model with one generation or equivalently a three-generation model without interfamily mixings. Then, the coupling J^0-e-e, g_{Jee}, defined by the relation

$$i\mathcal{T}^{ee} = g_{Jee}\bar{e}i\gamma_5 e \ , \tag{31}$$

takes the simple form

$$g_{Jee} \simeq \frac{g_W\alpha_W}{16\pi} t_\beta \frac{m_e}{M_W}\left[(s_L^{\nu_e})^2 \frac{\lambda_{N_e}^2}{1-\lambda_{N_e}}\left\{1+\frac{\ln\lambda_{N_e}}{1-\lambda_{N_e}}\right\} + \frac{1}{2}\sum_{e,\mu,\tau} (s_L^{\nu_l})^2 \lambda_{N_l}\right] . \tag{32}$$

In Eq. (32) the lepton-violating mixings $(s_L^{\nu_l})^2$ are defined as

$$(s_L^{\nu_l})^2 = \sum_{i=n_G+1}^{2n_G} |B_{li}|^2 \ . \tag{33}$$

In this scenario, the matrix m_D is diagonal and $(s_L^{\nu_l})^2 = m_{D_l}^2/m_{N_l}^2$. Since $m_{D_i} \simeq m_{l_i}$ or m_{u_i} as usually dictated by GUT's [13], the heavy Majorana neutrinos N_l will have large masses between 1 and 10^5 TeV due to Eq. (1). On the other hand, astrophysical constraints arising from helium ignition in red giants or the observational evidence of white dwarf cooling rates lead to an upper bound [1] on

$$g_{Jee} \lesssim (9.-1.4)\times 10^{-13} \ . \tag{34}$$

However, the range $3\times 10^{-13} \lesssim g_{Jee} \lesssim 6\times 10^{-7}$ is excluded from the helium ignition argument mentioned above, if the radius of giant core or dwarf is bigger than the mean free path of the pseudoscalars that these particles require to freely escape from them [1]. It is now obvious that the most stringent constraint on $\tan\beta$ arises from the heaviest family. Thus, for $m_D \simeq m_\tau$ one obtains that

$$g_{Jee} \simeq \frac{g_W\alpha_W}{32}\tan\beta \frac{m_e}{M_W}\frac{m_\tau^2}{M_W^2} \ , \tag{35}$$

yielding, because of Eq. (34),

$$\tan\beta \lesssim 0.4 \ . \tag{36}$$

Of course, if $m_D \simeq m_t/k \simeq 10$ GeV, one finds a much stronger bound, i.e.,

$$\tan\beta \lesssim 10^{-2} \ . \tag{37}$$

Note also that such low-energy realizations make unlikely the invisible decay of massive Higgs bosons into Majoron pairs [6,7].

The afore-mentioned hierarchical scheme, however, is, in general, not valid if one introduces intergenerational mixings in the singlet Majoron model. This situation seems to be a natural possibility that can be realized by GUT models, since m_D and M_U matrices may get related in such high-energy scenarios [i.e., $m_D(M_X) = M_U(M_X)$ with M_X indicating the grand unification scale] [13]. In addition, it has explicitly been demonstrated in [22,20] that the scale of m_M can be ~ 100 GeV without contradicting experimental bounds on neutrino masses. For instance, family-independent mass matrices for the form of m_D [24] can lead to patterns with such a low scale for m_M [25]. Then, the mixings $(s_L^{\nu_l})^2$ can be treated as purely phenomenological parameters, since Eq. (33) should now read

$$(s_L^{\nu_l})^2 \simeq m_D \frac{1}{m_M^2} m_D^\dagger \tag{38}$$

and cannot therefore be related with the light-neutrino mass matrix of Eq. (1). The mixing angles $(s_L^{\nu_l})^2$ can generally be constrained by a global analysis of a great number of low-energy experiments and LEP data [23]. In this scenario one also makes the remarkable observation that g_{Jee} can severely be suppressed for a certain choice of the mass parameters λ_{N_l} and mixings $(s_L^{\nu_l})^2$. For example, if all heavy neutrino masses m_{N_l} are approximately equal and $\lambda_{N_l} \gg 1$, then the choice

$$(s_L^{\nu_e})^2 \simeq (s_L^{\nu_\tau})^2 \tag{39}$$

leads to $g_{Jee} = 0$. However, even if the Majoron coupling to electrons vanishes, the corresponding coupling to nucleons $\mathcal{N}$, $g_{J\mathcal{N}\mathcal{N}}$, is not zero anymore. The reason is that the destructive first term in the brackets of Eq. (32) does not exist anymore and such a fine-tuning is thus not possible. Since $g_{J\mathcal{N}\mathcal{N}}/g_{Jee} \simeq m_{\mathcal{N}}/m_e \simeq 2\times 10^3$, one may derive useful constraints from the consideration of cooling rates of neutron stars due to the energy-loss mechanism by Majoron emission. In Fig. 3, we present ex-

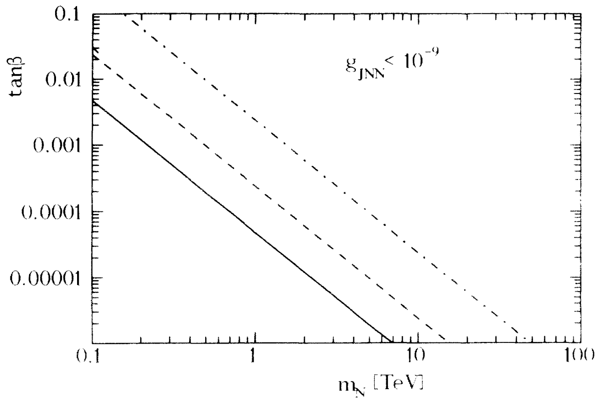

FIG. 3. Exclusion plots from astrophysical requirements. We have considered the values $(s_L^{\nu_l})^2 = 5\times10^{-2}$ (solid line), $(s_L^{\nu_l})^2 = 10^{-2}$ (dashed line), $(s_L^{\nu_l})^2 = 10^{-3}$ (dot-dashed line). The area lying above the curves is excluded by the restriction $g_{J\mathcal{N}\mathcal{N}} < 10^{-9}$. In addition, we assume that all heavy neutrino masses are approximately equal with m_N.

clusion plots of the parameters $\tan\beta$ versus m_N for three different values of $(s_L^{\nu_l})^2$ by considering that [26]

$$g_{J\mathcal{N}\mathcal{N}} \lesssim 10^{-9} . \tag{40}$$

For a discussion of possible uncertainties on the upper bound of the coupling $g_{J\mathcal{N}\mathcal{N}}$ that can arise from various reasons, we refer the reader to [27]. The main uncertainty, however, has been discussed in [28] for the axion coupling to nucleons, $g_{a\mathcal{N}\mathcal{N}}$, and can equally apply to $g_{J\mathcal{N}\mathcal{N}}$. In fact, it has been found in [28] that axion emission rates from SN 1987 A can be consistent with experiments based on neutrino observations if the coupling $g_{a\mathcal{N}\mathcal{N}}$ is sufficiently large for axions so as to be trapped and thermalized in the hot core of supernova. This may relax the upper bound of Eq. (40) by a factor of 10^3 and even larger. Ultimately, we must notice that astrophysical bounds on the Majoron coupling to two photons [1], $C_{J\gamma\gamma}$, are trivially satisfied, since the derivative-type interaction of the Majoron field and the absence of anomaly terms in the theory [29] will imply $C_{J\gamma\gamma}=0$ for on-mass-shell Majorons.

In the following we will focus our attention on bounds resulting solely from terrestrial experiments [16,17] by analyzing lepton-flavor-violating decays, i.e., $l_1 \to J^0 l_2$ with $l_1 \neq l_2$. To the leading order of the heavy neutrino limit one finds from Eqs. (18) and (19) that

$$B(l_1^- \to J^0 l_2^-) \simeq \frac{3\alpha_W}{8\pi} \tan^2\beta |B_{l_1N_i}^* B_{l_2N_i}|^2 \lambda_N^2 \frac{M_W^2}{m_{l_1}^2} . \tag{41}$$

The experimental information we have for the above lepton-violating decays are the upper bounds

$$\begin{aligned} &B(\mu \to J^0 e) \leq 2.6\times10^{-6}\ [16] , \\ &B(\tau \to J^0 e) \leq 7.1\times10^{-3}\ [17] , \\ &B(\tau \to J^0 \mu) \leq 2.3\times10^{-3}\ [17] . \end{aligned} \tag{42}$$

In order to quantitatively estimate the magnitude of the lepton-violating effects that could be constrained by the branching ratios stated in (42), we use the upper bound of the quantity

$$|B_{l_1N_i}^* B_{l_2N_i}| \leq (s_L^{\nu_l})^2 = \max((s_L^{\nu_{l_1}})^2 ,\ (s_L^{\nu_{l_2}})^2) . \tag{43}$$

The exclusion plots implied by these experiments are presented in Fig. 4 for the three different decay channels. For comparison, we have taken the astrophysical bound coming from Eq. (40) into account in Fig. 4, from which one easily concludes that experimental searched for the decay $\mu \to J^0 e$ may not be excluded by astrophysical constraints and can hence be sensitive to new physics beyond the SM.

In conclusion, astrophysical considerations may lead to useful constraints on the parameters of singlet Majoron models with intergenerational mixings. It has been demonstrated that three-generation Majoron models can indeed be constrained if the naturalness condition from GUT's that $m_D \propto M_U$ is taken into account. The latter has allowed our treatment of originally considering m_N and the lepton-violating mixings $(s_L^{\nu_l})^2$ as free parameters of the theory which have been restricted later on by our phenomenological analysis. However, possibilities of how to evade some of the astrophysical constraints have also been discussed. For example, g_{Jee} vanishes for a specific choice of parameters [30]. Furthermore, terrestrial experiments give independently severe restrictions both on lepton-violating mixings and heavy neutrino masses. Aside from rather involved R-parity broken models [31], this minimal extension of the SM, the singlet Majoron model, offers an attractive alternative that can account for possible lepton-flavor-violating signals in precision experiments. We emphasize again the fact that measurements of the TRIUMF Collaboration [16] for exotic decay modes, such as $\mu \to J^0 e$, lie in an area which may not be excluded by astrophysics and have substantial

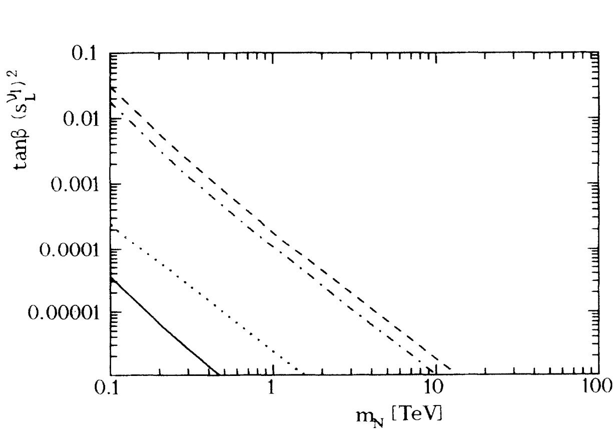

FIG. 4. Exclusion plots originating from the decays: $\mu \to J^0 e$ (solid line), $\tau \to J^0 e$ (dashed line), $\tau \to J^0 \mu$ (dot-dashed line). For comparison, we have considered the astrophysical bound $g_{J\mathcal{N}\mathcal{N}} \leq 10^{-9}$ (dotted line). The areas lying above the curves are excluded by the aforementioned conditions.

chances to establish new physics beyond the SM. Finally, due to the CP-odd interactions that Majorona neutrinos introduce in such models [see, e.g., Eqs. (9), (14), and (15)], one may be motivated to discuss their phenomenological impact of possible CP-violating effects in the decays of the Higgs particle into top-quark, W-, and Z-boson pairs [32].

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APPENDIX: THE LOOP INTEGRALS

We first define the useful functions $B_1(\lambda_i,\lambda_j)$ and $B_2(\lambda_i,\lambda_j)$ as

$$B_1(\lambda_i,\lambda_j)=\lambda_i(1-x)+\lambda_j x\ , \tag{A1}$$

$$B_2(\lambda_i,\lambda_j)=1-y+y[\lambda_i(1-x)+\lambda_j x]\ , \tag{A2}$$

where x and y are Feynman parameters. The loop integrals L_0, L_1, L_2, I_1, I_2, and I_3 are then given by

$$L_0(\lambda_i,\lambda_j)=\int dx\ \ln B_1(\lambda_i,\lambda_j)$$
$$=-1+\tfrac{1}{2}\ln\lambda_i\lambda_j-\frac{\lambda_i+\lambda_j}{2(\lambda_i-\lambda_j)}\ln\frac{\lambda_j}{\lambda_i}\ , \tag{A3}$$

$$L_1(\lambda_i,\lambda_j)=\int dx\ x\ \ln B_1(\lambda_i,\lambda_j)$$
$$=-\frac{1}{4}+\frac{\lambda_i^2}{2(\lambda_i-\lambda_j)^2}\ln\frac{\lambda_i}{\lambda_j}+\tfrac{1}{2}\ln\lambda_j-\frac{\lambda_i}{2(\lambda_i-\lambda_j)}\ , \tag{A4}$$

$$L_2(\lambda_i,\lambda_j)=\int dx\ dy\ y\ \ln B_2(\lambda_i,\lambda_j)$$
$$=\frac{3}{4(\lambda_i-\lambda_j)}\left[\frac{\lambda_j}{1-\lambda_i}-\frac{\lambda_i}{1-\lambda_j}\right]-\frac{3\lambda_i\lambda_j}{4(1-\lambda_i)(1-\lambda_j)}+\frac{1}{2(\lambda_i-\lambda_j)}\left[\frac{\lambda_j^2\ln\lambda_j}{1-\lambda_j}-\frac{\lambda_i^2\ln\lambda_i}{1-\lambda_i}\right]\ , \tag{A5}$$

$$I_1(\lambda_i,\lambda_j)=\int\frac{dx\ dy\ y}{B_1(\lambda_i,\lambda_j)}$$
$$=\frac{\lambda_i\lambda_j\ln(\lambda_i/\lambda_j)+\lambda_j\ln\lambda_j-\lambda_i\ln\lambda_i}{(1-\lambda_i)(1-\lambda_j)(\lambda_i-\lambda_j)}\ , \tag{A6}$$

$$I_2(\lambda_i,\lambda_j)=\int\frac{dx\ dy\ y^2}{B_2(\lambda_i,\lambda_j)}$$
$$=-\frac{1}{2(1-\lambda_i)(1-\lambda_j)}+\frac{1}{2(\lambda_i-\lambda_j)}\left[\ln\frac{\lambda_i}{\lambda_j}-\frac{\ln\lambda_i}{(1-\lambda_i)^2}+\frac{\ln\lambda_j}{(1-\lambda_j)^2}\right]\ , \tag{A7}$$

$$I_3(\lambda_i,\lambda_j)=\int\frac{dx\ dy\ y^2(1-2x)}{B_2(\lambda_i,\lambda_j)}$$
$$=-\frac{1}{2(1-\lambda_i)(1-\lambda_j)}\left[\frac{\ln\lambda_i}{1-\lambda_i}-\frac{\ln\lambda_j}{1-\lambda_j}\right]-\frac{1}{2(\lambda_i-\lambda_j)}\left[\frac{\lambda_j}{1-\lambda_j}+\frac{\lambda_i}{1-\lambda_i}\right]-\frac{\ln(\lambda_j/\lambda_i)}{2(\lambda_i-\lambda_j)^2}\left[\frac{\lambda_j^2}{1-\lambda_j}+\frac{\lambda_i^2}{1-\lambda_i}\right]\ . \tag{A8}$$

The integration interval of the variables x and y is [0,1].

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