

LEFT-RIGHT SYMMETRY AND NEUTRINO STABILITY

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Abstract

We consider a left-right symmetric model in which neutrinos acquire mass due to the spontaneous violation of both the gauged $B - L$ and a global $U(1)$ symmetry broken by the vacuum expectation value (VEV) of a gauge singlet scalar boson $\langle\sigma\rangle$. For suitable choices of $\langle\sigma\rangle$ consistent with all laboratory and astrophysical observations neutrinos will be unstable against majoron emission. All neutrino masses in the keV to MeV range are possible, since the expected neutrino decay lifetimes can be short enough to dilute their relic density below the cosmologically required level. A wide variety of possible new phenomena, associated to the presence of left-right symmetry and/or the global symmetry at the TeV scale, could therefore be observable, without conflict with cosmology. The latter includes the possibility of invisibly decaying higgs bosons, which can be searched at LEP, NLC and LHC.

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1 Introduction

One of the most attractive extensions of the standard electroweak theory is based on the gauge group $G_{LR} \equiv SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ [1, 2]. Apart from offering a possibility of understanding parity violation on the same footing as that of the gauge symmetry, these models incorporate naturally small neutrino masses. The magnitude of these masses is related in these theories to the scale at which the $SU(2)_R$ symmetry gets broken. If this breaking occurs at a low (~ 10 TeV) scale, then the neutrino masses are expected near their present laboratory limits, at least for sizeable values of the Dirac neutrino masses.

Such high values of the neutrino masses may be more than a theoretical curiosity, they may have quite important implications. For example, a tau neutrino with a mass in the MeV range is an interesting possibility first because such a neutrino is within the range of the detectability, for example at a tau-charm factory [3]. On the other hand, if such neutrino decays before the matter dominance epoch, its decay products could then add energy to the radiation thereby delaying the time at which the matter and radiation contributions to the energy density of the universe become equal. This would reduce density fluctuations on smaller scales [4] purely within the framework of the standard cold dark matter model [5], and could reconcile the large scale fluctuations observed by COBE [6] with the earlier observations such as those of IRAS [7] on the fluctuations at smaller scales [8].

It is well known that, if stable, neutrinos would contribute too much to the energy density of the universe if their mass lies in the range [9]

$$60 \text{ eV} \lesssim m_\nu \lesssim \text{few GeV} \quad (1)$$

Thus the interesting possibility of heavy neutrino masses can be consistent with cosmology only if there are new neutrino decay and/or annihilation channels absent in the standard model. Many neutrino decay modes have been suggested but all are quite unlikely to breach the forbidden range given above [10]. For example, neutrino radiative decay modes $\nu' \rightarrow \nu + \gamma$ and $\nu' \rightarrow \nu + \gamma + \gamma$ are disfavored, because they have a very long lifetime [11]. Moreover, such *visible* decays are very constrained by astrophysics [12] as well as laboratory searches [13]. What one needs are *invisible* decays such as $\nu' \rightarrow 3\nu$. It was noted long ago that such decays take place in models where isodoublet and isosinglet mass terms coexist, due to the peculiar structure of the neutral current in these models [14, 15] and this is the case in the left-right models. In contrast to the visible decays, these are almost unconstrained. However, in the simplest models of the seesaw type even if neutrino masses are close to their laboratory limits, the expected lifetimes tend to be too long to allow for sufficient redshift of the heavy neutrino decay products, and thus forbidden by cosmology [15]. Moreover, for $m_{\nu'} \gtrsim 1$ MeV the 3ν decay would also be accompanied by the *visible* channel $\nu' \rightarrow e^+e^-\nu$. This would, in turn suggest a γ -ray burst from a supernova explosion, the photons arising

from subsequent annihilation and/or bremsstrahlung processes. The non-observation of such a burst from SN1987 disfavors this possibility [16].

Although not possible in the simplest models [15], fast invisible neutrino decays can, under certain circumstances, naturally occur in many models where neutrino masses are induced from the spontaneous violation of a global $B - L$ symmetry [17, 18, 10]

$$\nu' \rightarrow \nu + J, \quad (2)$$

where here J denotes the massless Nambu-Goldstone boson, called majoron [19], which follows from the spontaneous nature of lepton number violation. These decays could have important implications in cosmology and astrophysics [10].

Unfortunately this possibility does not arise naturally in the left-right symmetric framework since the global symmetry associated with the conventional majoron is gauged in this case. Thus in order to obtain majoron one needs to impose an additional symmetry which is different from the $B - L$ symmetry, but which nevertheless plays a role in generating the neutrino masses and decays.

In this paper we propose a variant of the left-right symmetric model with an additional spontaneously broken $U(1)$ global symmetry, acting nontrivially on some new isosinglet leptons which mix with the ordinary neutrinos. Thus we extend the fermion sector in order to accommodate the required global symmetry whose spontaneous breaking will yield the majoron. This allows us to incorporate the idea of invisibly decaying neutrinos in the framework of a theory with gauged $B - L$. The additional singlet fermions used in our model may arise in various attempts to unify quarks and leptons in a superstring framework [20].

Majoron decays of neutrinos in the left-right symmetric model has also been considered in ref. [21]. However, our model has a more economic higgs sector and makes use of a different fermion content. Therefore, in a sense, it is complementary to that of [21].

In fact, our model is a left-right embedding of a previously suggested model [18] but has some noticeable differences which we study. We investigate the issue of neutrino stability in this model and demonstrate that, for reasonable choices of the breaking scales, $v_R \gtrsim \omega$ (v_R is the $B - L$ and parity breaking scale while $\omega \equiv \langle \sigma \rangle$ characterizes the breaking of the global $U(1)_G$ symmetry), the neutrino decay amplitude for the majoron decay mode of eq. (2) is of order $(m_D/M)^4$ where $M = g_5 v_R$ with g_5 being the appropriate Yukawa coupling. This is in full agreement with previous studies within the $SU(2) \otimes U(1)$ theory [15, 17]. Nevertheless, for appropriate choices of parameters, this simplest model yields majoron emission ν_τ decay lifetimes which can be fast enough to dilute the relic ν_τ density to acceptable levels for all values of the ν_τ mass. For most typical parameter choices, the ν_μ is light enough as to lie outside the range in eq. (1) and be stable, as required in order to be hot dark matter.

We also propose a very simple variant of this left-right symmetric model where the global $U(1)_G$ symmetry is of the horizontal type, as originally used in ref. [17]. This substantially enhances the neutrino decay amplitude for the majoron decay mode of eq. (2) to order $(m_D/M)^2$. In this case the majoron is a pure gauge singlet, as in the original proposal [19], and therefore both scales $\langle\sigma\rangle$ and v_R may be chosen to be at the TeV scale quite naturally. This opens up a very wide phenomenological potential for left-right extensions of the standard electroweak theory, free of cosmological problems.

2 The simplest model

We consider a model based on the gauge group

$$G_{LR} \equiv SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$$

in which an extra $U(1)_G$ global symmetry is postulated. The matter and higgs boson representation content is specified in table 1. In addition to the conventional quarks and leptons, there is a gauge singlet fermion in each generation ^{||}. These extra leptons might arise in superstring models [20]. They have also been discussed in an early paper of Wyler and Wolfenstein [23]. We will not use the more conventional triplet higgs scalars, which are absent in many of these string models. Instead we will substitute them by the doublets χ_L and χ_R . This could in fact play an important role in unifying this model in $SO(10)$, while keeping left-right symmetry unbroken down to the TeV scale [22].

The Yukawa interactions allowed by the $G_{LR} \otimes U(1)_G$ symmetry are given as

$$\begin{aligned}
-\mathcal{L}_Y &= g_1 \bar{Q}_L \phi Q_R + g_2 \bar{Q}_L \tilde{\phi} Q_R + g_3 \bar{\psi}_L \phi \psi_R + g_4 \bar{\psi}_L \tilde{\phi} \psi_R + \\
&g_5 [\bar{\psi}_L \chi_L S_R^c + \bar{\psi}_R \chi_R S_L] + g_6 \bar{S}_L S_R^c \sigma + h.c.
\end{aligned}
\tag{3}$$

where g_i are matrices in generation space and $\tilde{\phi} = \tau_2 \phi^* \tau_2$ denotes the conjugate of ϕ . This Lagrangean is invariant under parity operation $Q_L \leftrightarrow Q_R$, $\psi_L \leftrightarrow \psi_R$, $S_L \leftrightarrow S_R^c$, $\phi \leftrightarrow \phi^\dagger$ and $\chi_L \leftrightarrow \chi_R$.

The symmetry breaking pattern is specified by the following scalar boson VEVs (as-

^{||}Although the number of such singlets is arbitrary, since they do not carry any anomaly, we add just one such lepton in each generation, while keeping the quark sector as the standard one. Further extensions can be made, as recently discussed in ref. [22].

sumed real):

$$\left\{ \begin{array}{l} \langle \phi \rangle = \begin{pmatrix} k & 0 \\ 0 & k' \end{pmatrix}; \quad \langle \chi_L^0 \rangle = v_L; \quad \langle \chi_R^0 \rangle = v_R; \\ \langle \tilde{\phi} \rangle = \begin{pmatrix} k' & 0 \\ 0 & k \end{pmatrix}; \quad \langle \sigma \rangle = \omega \end{array} \right. \quad (4)$$

The spontaneous violation of the global $U(1)_G$ symmetry generates a physical majoron whose profile in the limit $V^2 \ll v_R^2$ is specified as

$$J = (\omega^2 + \frac{v_L^2 v^2}{V^2})^{-1/2} \{ \omega \sigma_I + \frac{v_L v}{V^2} [v \chi_L^I - \frac{v_L}{v} (k \phi_2 - k' \phi_4)] \} \quad (5)$$

where σ_I , χ_L^I , ϕ_2 and ϕ_4 denote the imaginary parts of the neutral fields in σ , χ_L and the bidoublet ϕ . Here we have also defined the VEVs as $v^2 \equiv k^2 + k'^2$ and $V^2 \equiv v^2 + v_L^2$.

Note that the majoron has no component along the imaginary part of χ_R despite the fact that χ_R is nontrivial under the global symmetry. Clearly, as it must, the majoron is orthogonal to the Goldstone bosons eaten-up by the Z and the new heavier neutral gauge boson present in the model. The latter acquires mass at the larger scale v_R .

The various scales appearing in eq. (5) are not arbitrary. First of all, note that the minimization of the scalar potential dictates the consistency relation [2]

$$\frac{v_R}{v} \sim \lambda \frac{\omega}{v_L} \quad (6)$$

For $\lambda \sim 1$ the singlet VEV is necessarily larger than v_L i.e. $v_L \ll \omega$ and, as a result, the majoron is mostly singlet and the invisible decay of the Z to the majoron is enormously suppressed, unlike in the purely doublet or triplet majoron schemes.

On the other hand, in order that majoron emission does not overcontribute to stellar energy loss one needs to require [24]

$$\frac{v_L^2}{\sqrt{2} \omega V^2} \lesssim 10^{-9} \text{GeV}^{-1}. \quad (7)$$

One sees that eq. (6) and eq. (7) allow for the existence of right-handed weak interactions at accessible levels, provided v_L is sufficiently small, i.e. $\mathcal{O}(\lesssim 100 \text{ keV})$.

Note, however that the astrophysical bound in eq. (7) is hard to reconcile with the low-scale right-handed weak interactions in the case where $v_L \sim v$. For example, $v_L \sim 1 \text{ GeV}$ would require $v_R \gtrsim 10^7 \text{ GeV}$, $\omega \gtrsim 10^5 \text{ GeV}$ with $v^2 \equiv k^2 + k'^2$ fixed by the masses of W and Z bosons. As we will show later, it is possible to avoid this bound altogether in a simple variant of this model (see below).

3 Neutrino Masses and Majoron Couplings

Once all gauge and global symmetries get broken a mass term is generated for the electrically neutral leptons, of the form

$$\frac{1}{2} \left(\bar{\Psi}_L M_\nu \Psi_R + h.c. \right),$$

where $\Psi_L = (\nu_L, \nu_L^c, S_L)$ and $\Psi_R = (\nu_R^c, \nu_R, S_R^c)$. It may be written in block form as

$$M_\nu = \begin{pmatrix} 0 & m_D & \beta \\ m_D^T & 0 & M^* \\ \beta^T & M^\dagger & \mu \end{pmatrix}, \quad (8)$$

where the various entries are specified as

$$\beta = g_5 v_L, \quad M = g_5 v_R \quad m_D = g_3 k + g_4 k', \quad \mu = 2g_6 \omega. \quad (9)$$

Here the matrix m_D is the Dirac mass term determined by the standard higgs bi-doublet VEV $\langle \phi \rangle$ responsible for quark and charged lepton masses, β and M are G and $B - L$ violating mass terms determined by v_L and v_R , while μ is a gauge singlet G -violating mass, proportional to the VEV of the gauge singlet higgs scalar σ carrying 2 units of G charge. Note the zeroes in the first two diagonal entries. They arise because there are no higgs fields to provide the usual Majorana mass terms [20] which would be required in the seesaw mechanism [25, 2, 15].

In order to determine the light neutrino masses and majoron couplings we will work in the seesaw approximation, which we define as $M, \mu \gg m_D, \beta$. In this case the mass matrix in eq. (8) can be brought to block diagonal form via a transformation U (with $U^\dagger U = U U^\dagger = 1$),

$$\hat{M}_\nu \equiv U M_\nu U^T = \begin{pmatrix} \hat{M}_1 & 0 \\ 0 & \hat{M}_2 \end{pmatrix}, \quad (10)$$

where

$$\begin{aligned} \hat{M}_1 &= -(m_D, \beta) \begin{pmatrix} 0 & M^* \\ M^\dagger & \mu \end{pmatrix}^{-1} \begin{pmatrix} m_D^T \\ \beta^T \end{pmatrix} = \\ &= \epsilon \mu \epsilon^T - (\beta \epsilon^T + \epsilon \beta^T), \quad \epsilon \equiv m_D M^{\dagger-1} \end{aligned} \quad (11)$$

denotes the effective light neutrino mass matrix, determining the masses of ν_e , ν_μ and ν_τ . Notice that the light neutrino masses are generated due to the interplay of the violation of the global G as well as the gauged $B - L$ symmetry. Due to the relation in eq. (6) the two contributions to the neutrino masses in the last line of eq. (11) will be typically comparable.

The heavy sector is characterized by a 6×6 mass matrix given as

$$\hat{M}_2 \simeq M_2 \equiv \begin{pmatrix} 0 & M^* \\ M^\dagger & \mu \end{pmatrix}.$$

Finally, the matrix \hat{M}_ν is further diagonalized by a block diagonal unitary matrix T

$$T\hat{M}_\nu T^T = M_{diag} = \text{diag}(m_1, \dots, m_9). \quad (12)$$

which can be written as

$$T = \begin{pmatrix} V_1 & 0 \\ 0 & V_2 \end{pmatrix},$$

The total diagonalizing matrix A can then be written as

$$A = TU = \begin{pmatrix} V_1(1 - \frac{1}{2}\rho\rho^\dagger) & -V_1\rho \\ V_2\rho^\dagger & V_2(1 - \frac{1}{2}\rho^\dagger\rho) \end{pmatrix} + \mathcal{O}(\rho^3), \quad (13)$$

where V_1 and V_2 are the matrices that diagonalize the light and heavy neutrino mass matrices respectively, and

$$\rho \equiv (m_D, \beta)M_2^{-1} = (-(\epsilon\mu - \beta)M^{*-1}, \epsilon). \quad (14)$$

Note that $\rho \rightarrow 0$ as $M \rightarrow \infty$. This parameter plays the same role in the present model as the expansion parameter ϵ introduced in the $SU(2) \otimes U(1)$ context in ref. [15]. The relation between weak and mass eigenstates may then be written as

$$(\nu_R^c, \nu_R, S_R^c)^T = A^T(\nu_R^{c'}, \nu_R', S_R^{c'})^T,$$

where the prime refers to the mass eigenstate basis. The majoron-neutrino interaction Lagrangean, obtainable using eqs.(3) and (5), may be written in terms of weak eigenstates as

$$\mathcal{L}_J = \frac{iJ}{\sqrt{2}} \frac{1}{\sqrt{\omega^2 + \left(\frac{v_L v}{V}\right)^2}} \left\{ \frac{1}{2} \bar{S}_L \mu S_R^c + \left(\frac{v}{V}\right)^2 \bar{\nu}_L \beta S_R^c - \left(\frac{v_L}{V}\right)^2 \bar{\nu}_L m_D \nu_R \right\} + h.c. \quad (15)$$

4 Neutrino Decays and Cosmology

Relic neutrinos will overcontribute to the present-day energy density of the universe unless there are decay and/or annihilation channels. The cosmological density constraint on the neutrino decay lifetime for a $m_\nu \lesssim 1$ MeV neutrino is given as [26, 27]

$$m_\nu \left(\frac{\tau_\nu}{t_0}\right)^{1/2} \lesssim 100 h^2 \text{ eV} \quad (16)$$

where t_0 and h are the present age of the universe and the normalized Hubble parameter. The above constraint follows from demanding that an adequate redshift of the heavy neutrino decay products occurs.

In the present model, even though $B-L$ is a gauge symmetry, neutrino masses following from eq. (8) are accompanied by the existence of a massless majoron J given by eq. (5). This will lead to invisible neutrino decays with majoron emission, eq. (2).

To determine the neutrino decay rates we are interested in those majoron couplings to light neutrinos that are nondiagonal in the mass eigenstate basis. These couplings can be determined by rewriting explicitly eq. (15) in terms of mass eigenstates. This procedure is straightforward but subtle. There are, here too, the same tricky cancellations first discovered in the context of the standard $SU(2) \otimes U(1)$ model in ref. [15]. The result is that majoron couplings to light neutrinos are still diagonal to $\mathcal{O}(\epsilon^2)$, and therefore can not induce neutrino decay to this order.

In order to see this more clearly and, at the same time, determine the required nondiagonal couplings we prefer, instead of directly using eq. (15), to use a more general and powerful method based on the use of Noether's theorem for the global G -current. The method was given in the Sec. VI of ref. [15] and subsequently used, *e.g.*, in the first paper of ref. [17]. It has the advantages of being simpler and more systematic.

Using it one can easily determine the coupling matrix of the majoron to the light mass eigenstate neutrinos in the present model as

$$g_{ab} = \frac{1}{\langle \sigma \rangle} [m_a R_{ab} + m_b R_{ba}]; \quad a \neq b \quad (17)$$

where m_a denote the light neutrino masses and the matrix R is determined by the three light entries of the 9×9 matrix

$$R \approx A^* Q_1 A^T, \quad (18)$$

where Q_1 is a diagonal matrix related to the G charges of the leptons $\Psi_L = (\nu_L, \nu_L^c, S_L)$. Since only the gauge singlet leptons transform under G , the matrix Q_1 can be written as

$$Q_1 = \text{diag}(0, 0, 1). \quad (19)$$

The matrix A was previously defined as that which diagonalizes the full neutrino mass matrix. Using eq. (13) we can rewrite the part R_L of the matrix R connecting light neutrinos as the 3×3 matrix

$$R_L = V_1^* \rho^* \hat{Q}_1 \rho^T V_1^T = V_1^* \epsilon^* \epsilon^T V_1^T, \quad (20)$$

where the 6×6 diagonal matrix \hat{Q}_1 is defined as $\text{diag}(0,0,0,1,1,1)$. This shows explicitly that the nondiagonal entries of the majoron coupling matrix in eq. (17) arise manifestly at $\mathcal{O}(\epsilon^4)$, in agreement with results found in the $SU(2) \otimes U(1)$ theory [15].

In order to get an idea of the expected neutrino decay rates in this model we first make a crude estimate of the magnitude of the neutrino masses following from eq. (11). Using eq. (6) and eq. (7) one sees that

$$m_{\nu_i} \sim 2 \frac{g_6}{g_5^2} \frac{\omega m_{q_i}^2}{v_R^2} \lesssim 2 \frac{g_6}{g_5^2} \left(\frac{m_{q_i}}{\text{GeV}} \right)^2 \text{ eV} . \quad (21)$$

In estimating this upper limit we have assumed the Dirac neutrino masses to be of the same order as the corresponding up-quark masses. Assuming a reasonable choice of parameters where the ratio $(2g_6/g_5^2)$ lies in the range 1 to 10^3 we get $m_{\nu_\tau} \lesssim 10 \text{ keV}$ to 10 MeV , $m_{\nu_\mu} \lesssim 1 \text{ eV}$ to 1 keV and $m_{\nu_e} \lesssim 10^{-4}$ to 10^{-2} eV . Thus the ν_μ may well be stable, as required in order to be dark matter, while the ν_τ is expected to violate the cosmological limit eq. (1) and has to decay with lifetime obeying eq. (16).

We now make a simple numerical estimate of the ν_τ lifetime. From eq. (20) we can parametrize the nondiagonal coupling responsible for ν_τ decay to the lighter neutrinos plus majoron as

$$g_{3a} = \frac{m_3}{\langle \sigma \rangle} (g_s + g_c), \quad (22)$$

where m_3 denotes m_{ν_τ} and we have set

$$g_s = \frac{1}{2} [(\epsilon\epsilon^T)_{22} - (\epsilon\epsilon^T)_{11}] \sin 2\theta_L , \quad (23)$$

$$g_c = (\epsilon\epsilon^T)_{12} \cos 2\theta_L . \quad (24)$$

In obtaining these formulas we have assumed for simplicity that ν_τ mixes with only one of the two lighter neutrinos, with mixing angle θ_L , and that ϵ is real. For small values of θ_L we may only keep the second term. Assuming a simple scaling ansatz $\epsilon \sim m_D/M$ we get

$$g_c \approx \left(\frac{m_D}{M} \right)^2 \quad (25)$$

One can now easily see that the neutrino decay lifetime becomes

$$\tau(\nu_3 \rightarrow \nu + J) = \frac{16\pi}{g_{3a}^2} \frac{1}{m_{\nu_3}} \approx 3 \times 10^7 (\text{keV}/m_3)^3 \left(\frac{\langle \sigma \rangle}{10^6 \text{ GeV}} \right)^2 \left(\frac{m_D}{M} \right)^{-4} \text{ sec} \quad (26)$$

The lifetime in eq. (26) can be short enough to obey the cosmological constraint in eq. (16). Indeed, from eqs. (21)–(26) one can readily find the following dependence between $\tau(\nu_3)$ and m_{ν_3} in our model:

$$\tau(\nu_3) \approx 1.4 \times 10^8 \left(\sqrt{g_6} \frac{\omega}{\text{GeV}} \right)^4 \left(\frac{\text{keV}}{m_{\nu_3}} \right)^5 \text{ sec} .$$

This dependence is plotted for two illustrative values of $\sqrt{g_6}\omega$ in Fig. 1 alongside with the cosmological bound of eq. (16). It can be seen from this figure that for $\sqrt{g_6}\omega = 35 \text{ GeV}$ cosmology does not constrain the model for $m_{\nu_\tau} > 1 \text{ keV}$. For $\sqrt{g_6}\omega = 10^3 \text{ GeV}$ cosmological

bound excludes values of m_{ν_τ} below 100 keV. Thus in this case cosmological considerations provide a lower limit on the ν_τ mass.

One can rewrite the cosmological constraint eq. (16) in terms of the VEVs ω and v_R instead of $\tau(\nu_3)$ and m_{ν_3} : $(\omega/\text{GeV})^{1/2}(v_R/\text{GeV})^3 \lesssim 6 \times 10^{19}A$, where $A \equiv \left(\frac{\sqrt{g_6}}{g_5}h^2\right)$. The astrophysical constraint of eq. (7) can be written using eq. (6) as $(\omega/\text{GeV})(v_R/\text{GeV})^{-2} < 10^{-9}$. These two constraints are plotted in Fig. 2 for $A = 10^5$. The region below both straight lines illustrates what is allowed for this representative choice of parameters. For example, we can see from this figure that ω cannot exceed 6×10^5 GeV, corresponding to a value of $v_R \approx 2.4 \times 10^7$ GeV.

This generalizes to our left-right symmetric model the results obtained in the analysis of the question of neutrino stability in majoron models given in ref. [15]. The tau neutrino is expected to be in the keV to MeV range with a lifetime that can be as short as 1 sec.

5 A Model with Enhanced Neutrino Decays

We now briefly sketch a variant of the previous model with exactly the same particle content, but with the global $U(1)_G$ symmetry G assigned in a nonsequential way. The model may be seen also as a left-right symmetric variant of the original horizontal lepton number models [17].

The G charges of the lepton doublets of the first two generations can be assigned as 1 and -1 respectively. Similarly, under the global $U(1)_G$ symmetry the gauge singlets transform with charges +1 (S_{1L} and S_{2R}^c) and -1 (S_{2L} and S_{1R}^c). Finally the third generation leptons carry no G charge.

Another important difference with respect to the model discussed in the previous sections, insofar as the G assignments of the higgs scalar bosons are concerned, is that now χ_L and χ_R carry no G charge and therefore the resulting majoron will be a pure gauge singlet, as in the original model [19]. This has a very important phenomenological implication, namely that one now avoids the astrophysical constraint of eq. (7), allowing for very low values of the G breaking scale $\langle\sigma\rangle$ which may naturally lie at the electroweak scale.

The quantum numbers are summarized in Table 2. Since the same higgs multiplets are used the mass matrix has the same general structure as in eq. (8). However, as a result of the horizontal assignment of the global charges of the leptons, the entries in eq. (8) now have special textures in generation space.

To find these textures, let us first notice that our present G charge assignment supports the discrete parity symmetry of the model, provided the parity operation for the gauge-singlet

leptons of the first two generations is modified: $S_{Le} \leftrightarrow S_{R\mu}^c$, the rest of the fields transforming as before. The g_5 term in the Yukawa Lagrangean of eq. (3) will now read

$$(g_5)_1[\bar{\psi}_{Le}\chi_L S_{R\mu}^c + \bar{\psi}_{Re}\chi_R S_{Le}] + (g_5)_2[\bar{\psi}_{L\mu}\chi_L S_{Re}^c + \bar{\psi}_{R\mu}\chi_R S_{L\mu}] + (g_5)_3[\bar{\psi}_{L\tau}\chi_L S_{R\tau}^c + \bar{\psi}_{R\tau}\chi_R S_{L\tau}] \quad (27)$$

One can now readily find the entries of the neutrino mass matrix in eq. (8). They are given by diagonal forms for both m_D and M , while the remaining entries β and μ take on the following forms

$$\beta = \begin{pmatrix} 0 & \times & 0 \\ \times & 0 & 0 \\ 0 & 0 & \times \end{pmatrix} \quad (28)$$

and

$$\mu = \begin{pmatrix} \times & \times & 0 \\ \times & \times & 0 \\ 0 & 0 & \times \end{pmatrix}, \quad (29)$$

where in the last equation the 12 and 33 entries are bare masses, allowed by the G symmetry, while the 11 and 22 are proportional to the VEVs of σ^* and σ respectively.

The horizontal nature of the G assignments removes the additional $\mathcal{O}(\epsilon^2)$ suppression in the neutrino decay rate. To see this note that now the matrix R_L of eq. (20) is replaced by

$$R_L = V_1^* \hat{Q}_2 V_1^T \quad (30)$$

where \hat{Q}_2 is a 3×3 matrix given by $\text{diag}(1, -1, 0)$ dictated by the G charge assignments. Note that in the above equation there is no $\rho \sim \epsilon$ suppression.

In summary, the main features of this second model are

1. The majoron is a pure gauge singlet, allowing for the G breaking scale to be as low as the electroweak scale;
2. Since eq. (7) need not hold in this model, the left-right symmetry can be realized at the TeV scale;
3. Majoron emission neutrino decay amplitudes are enhanced to $\mathcal{O}(\epsilon^2)$.

The combined effect of the above features is to provide a tremendous enhancement of the neutrino decay amplitude, leading to a lifetime shorter than eq. (26) by as much as 20 orders of magnitude.

The existence of models such as this opens a very wide phenomenological potential for left-right extensions of the standard model, consistent with all cosmological observations.

6 Discussion

We have examined the issue of neutrino stability in a class of left-right symmetric models where neutrinos may acquire mass from the spontaneous violation of both the gauged $B - L$ symmetry and a global $U(1)$ symmetry broken by the vacuum expectation value of a gauge singlet scalar boson σ . For suitable choices of $\langle\sigma\rangle$ consistent with all laboratory and astrophysical observations neutrinos will be unstable against majoron emission. We have considered two models. In the simplest one the global symmetry is flavour blind, while in the second it distinguishes between leptons of different type. In the first model the tau neutrino may be heavy and unstable against majoron emission, with decay amplitude of order $(m_D/M)^4$. Despite such strong a suppression, neutrino decay rates are consistent with the cosmological requirements in a wide range of the parameters of the model. The parity violation scale can be as low as a few TeV if the left-handed doublet higgs VEV v_L is $\lesssim 100$ keV, but should be $\gtrsim 10^9$ GeV for v_L of the order of the electroweak scale. In the second model, neutrino decay amplitudes are substantially enhanced by a combined effect of low values for both global and left-right symmetry breaking scales $\langle\sigma\rangle$ and v_R . These scales may be naturally chosen to be at the TeV scale. This opens up a very wide phenomenological potential for left-right extensions of the standard electroweak theory, consistent with cosmology. First of all, our models allow neutrino masses in the keV to MeV range, with potential effects related to neutrino masses and mixing, such as enhanced neutrinoless $\beta\beta$ decay rates. Moreover, they allow for the existence of neutral heavy leptons with masses at the weak scale. If lighter than the Z boson, these may give rise to quite striking signatures at LEP [28]. In addition, there are the potential effects due to the presence of right-handed weak currents at the TeV scale, including neutrinoless $\beta\beta$ decays and many other effects. Finally, there may be effects associated to the global symmetry violation at the TeV scale, such as the unusual possibility of an invisibly decaying higgs boson [29] $h \rightarrow JJ$, which can be searched at future colliders such as LEP, NLC and LHC [30].

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	$SU(2)_L \otimes$	$SU(2)_R \otimes$	$U(1)_{B-L} \otimes$	$U(1)_G$
Q_{Li}	2	1	1/3	0
Q_{Ri}	1	2	1/3	0
ψ_{Li}	2	1	-1	0
ψ_{Ri}	1	2	-1	0
S_{Li}	1	1	0	1
ϕ	2	2	0	0
χ_L	2	1	-1	1
χ_R	1	2	-1	-1
σ	1	1	0	2

Table 1: $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L} \otimes U(1)_G$ assignments of the quarks, leptons and higgs scalars.

	$SU(2)_L \otimes$	$SU(2)_R \otimes$	$U(1)_{B-L} \otimes$	$U(1)_G$
Q_{Li}	2	1	1/3	0
Q_{Ri}	1	2	1/3	0
ψ_{Le}	2	1	-1	1
$\psi_{L\mu}$	2	1	-1	-1
$\psi_{L\tau}$	2	1	-1	0
ψ_{Re}	1	2	-1	1
$\psi_{R\mu}$	1	2	-1	-1
$\psi_{R\tau}$	1	2	-1	0
S_{Le}	1	1	0	1
$S_{L\mu}$	1	1	0	-1
$S_{L\tau}$	1	1	0	0
ϕ	2	2	0	0
χ_L	2	1	-1	0
χ_R	1	2	-1	0
σ	1	1	0	2

Table 2: $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L} \otimes U(1)_G$ assignments of the quarks, leptons and higgs scalars in the model of section 5. Notice the nonsequential assignment of the global charge.

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FIGURE CAPTIONS

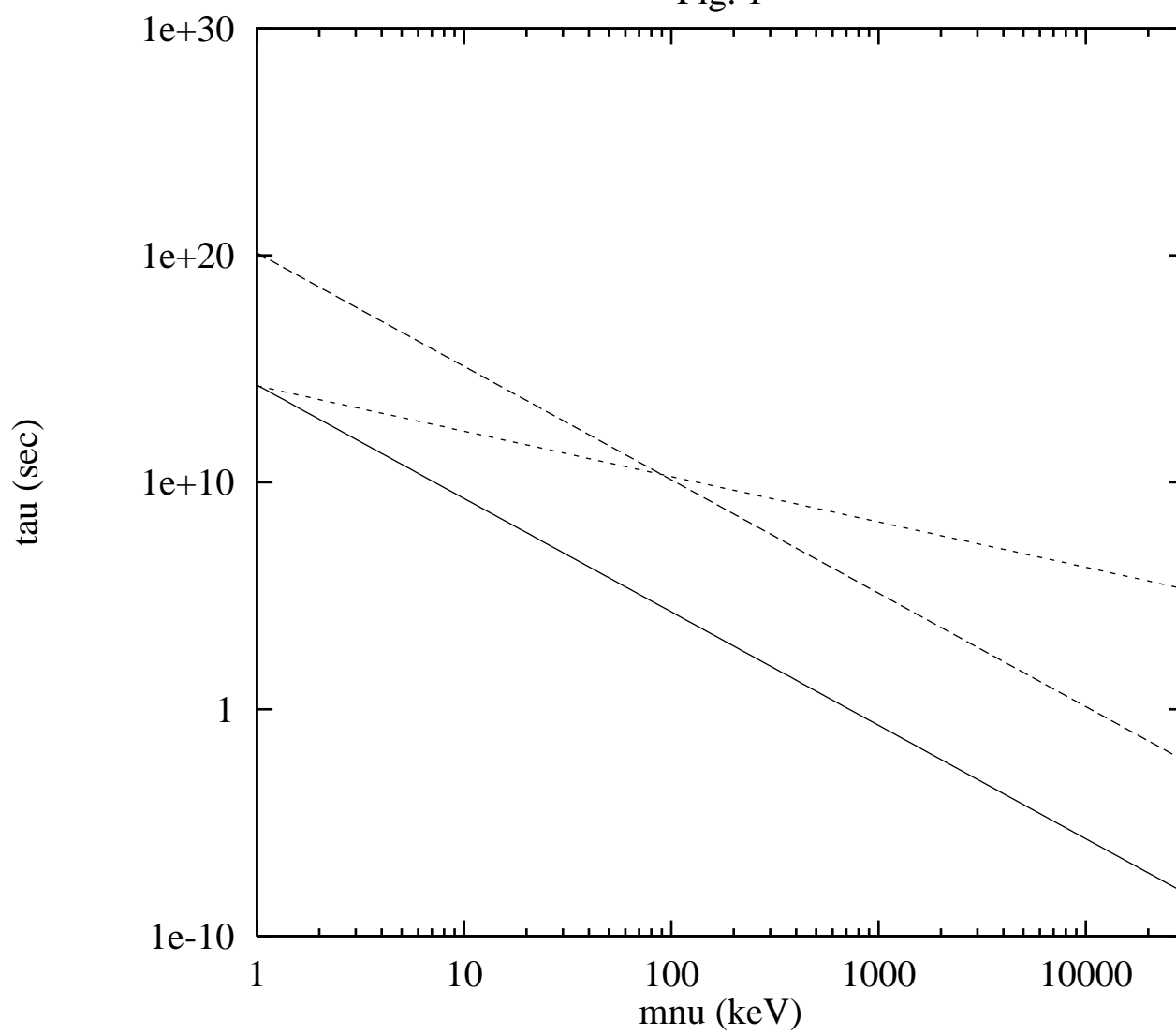
Fig. 1: Typical expectations for the tau neutrino lifetime as a function of the neutrino mass in the model of Sec. 2–4. The dotted line corresponds to the cosmological limit of eq. (16) (the region below the line is allowed). Solid (dashed) line is the relation between the ν_τ mass and lifetime for two typical choices of the parameter $\sqrt{g_6}\omega=35$ (10^3) GeV.

Fig. 2: Constraints on the singlet (ω) and right-handed doublet (v_R) VEVs following from cosmological limit eq. (16) (solid line) and from the red giant constraint eq. (7) (dashed line) for an illustrative choice of the parameter $h^2 \frac{\sqrt{g_6}}{g_5^3} = 10^5$. The region below both lines is allowed.

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Fig. 1



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Fig. 2

