17-keV neutrino in a left-right model

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We examine the possibility of embedding a 17-keV neutrino in a left-right model with three generations of left- and right-handed neutrinos. In particular, we study scenarios where the 17-keV neutrino is a pseudo Dirac neutrino and its mass is generated radiatively. Two heavy right-handed neutrinos have their mass scale set by the $SU(2)_R$-breaking scale, and the two light-neutrino masses are then determined by a seesaw mechanism. This scenario requires new Higgs bosons beyond those necessary in a minimal left-right model. The left-right symmetry leads to relations between the charged- and neutral-lepton mass matrices. We examine the constraints on the model from charged-lepton masses, leptonic flavor-changing neutral currents, neutrino data, cosmology, and astrophysics.

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I. INTRODUCTION

Evidence in recent years suggests the possible existence of a neutrino in β-decay spectra which is part electron neutrino (at the 1% level) and has a mass of 17 keV [1,2]. Although there have also been null experiments [3], this possibility has generated much theoretical interest as to how such a neutrino might occur [4]. In one scheme, proposed by Glashow [5], there are three left-handed and three right-handed neutrinos; exactly two right-handed neutrinos achieve large Majorana masses and through a seesaw mechanism drive two left-handed neutrino masses to be quite small. The remaining two neutrino states have masses close to the geometric mean of the small and large neutrino masses. They are identified as the two components of the 17-keV neutrino, with a left-handed active component and a right-handed sterile component. Limits from neutrinoless double-β decay require that this neutrino be a pseudo Dirac neutrino; i.e., the two components must be nearly degenerate in mass [6]. Several recent theoretical investigations adopt the basic ideas of this scenario as a starting point [7,8]. In one such model, the 17-keV mass is radiatively generated, which explains why it is much smaller than the charged-lepton masses [7].

In this paper we examine the possibility of using this modified seesaw mechanism with a radiatively generated neutrino mass scale in the context of a left-right model. We will attempt to construct such a model naturally, i.e., without assuming fine-tuning of parameters. Our main guide will be to maintain the "naturalness" of the neutrino masses, which we take to mean that no couplings or vacuum expectation values (VEV's) shall be much smaller than the others unless there is a symmetry which causes it to vanish altogether, or an approximate symmetry which suppresses it. Naturalness also assumes that there shall be no fortuitous cancellation of parameters. We do not attempt to explain the generational hierarchy of charged-fermion masses, but rather take their observed masses as input for constructing a model for neutrino masses. The 17-keV scale arises at the one-loop level and must also be protected from the seesaw mechanism. These requirements then imply that additional symmetries and Higgs fields must be included in the model.

In determining the basic structure of the neutrino masses we must also take into account the many phenomenological constraints on the model. These include limits on neutrino oscillations and flavor-changing neutral currents (FCNC's), and cosmological constraints on the decay of the 17-keV neutrino. We also would like the model to include the possibility of explaining solar-neutrino data via neutrino oscillations. Because of the left-right symmetry, the charged-lepton and Dirac neutrino masses are closely related, and some of the constraints can be quite severe. We find that a judicious choice of additional symmetries and Higgs structure gives neutrino mass matrices that satisfy most of the phenomenological constraints, but that ultimately some fine-tuning is necessary.

To begin, we present an outline of the model, and introduce the set of additional symmetries and Higgs multiplets that are required. We assume the standard set of left- and right-handed fermions

$$q_{Li}=\begin{pmatrix} u_{Li} \\ d_{Li}\end{pmatrix}, \quad q_{Ri}=\begin{pmatrix} u_{Ri} \\ d_{Ri}\end{pmatrix}, \qquad \psi_{Li}=\begin{pmatrix} \nu_{eLi} \\ e_{Li}\end{pmatrix}, \quad \psi_{Ri}=\begin{pmatrix} \nu_{eRi} \\ e_{Ri}\end{pmatrix}, \tag{1}$$

where the latin index stands for generation number. The transformation properties of the fermions under the gauge symmetry $SU(2)_L\times SU(2)_R\times U(1)_{B-L}$ are given in Table I.

We include in the Higgs sector the usual Higgs multiplets that exist in a standard left-right model with a seesaw mechanism:

$$\Phi_s=\begin{pmatrix} \phi_{s1}^0 & \phi_{s1}^- \\ \phi_{s2}^- & \phi_{s2}^0\end{pmatrix}, \quad \Delta_L=\begin{pmatrix} \Delta_L^+/\sqrt{2} & \Delta_L^{++} \\ \Delta_L^0 & -\Delta_L^+/\sqrt{2}\end{pmatrix}, \qquad \Delta_R=\begin{pmatrix} \Delta_R^+/\sqrt{2} & \Delta_R^{++} \\ \Delta_R^0 & -\Delta_R^+/\sqrt{2}\end{pmatrix}. \tag{2}$$

TABLE I. Transformation properties of the fermions and Higgs bosons under the local symmetry $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ and the global symmetry $U(1)' \times U(1)''$.

Field	$SU(2)_L$	$SU(2)_R$	$U(1)_{B-L}$	$U(1)'$	$U(1)''$
$\Psi_{L1},\Psi_{L2},\Psi_{L3}$	$\frac{1}{2}$	0	-1	0	-1
Ψ_{R1},Ψ_{R2}	0	$\frac{1}{2}$	-1	0	1
Ψ_{R3}	0	$\frac{1}{2}$	-1	1	1
q_{Lj}	$\frac{1}{2}$	0	$\frac{1}{3}$	0	0
q_{Rj}	0	$\frac{1}{2}$	$\frac{1}{3}$	0	0
Φ	$\frac{1}{2}$	$\frac{1}{2}$	0	0	2
Φ'	$\frac{1}{2}$	$\frac{1}{2}$	0	1	2
Δ_L	1	0	2	0	2
Δ_R	0	1	2	0	-2
Φ_s	$\frac{1}{2}$	$\frac{1}{2}$	0	0	0
η^-	0	0	-2	1	6

The Φ_s multiplet gives quark masses and Δ_L and Δ_R give Majorana neutrino masses. The Φ_s multiplet does not couple to leptons. The additional Higgs fields needed to realize the desired neutrino mass scheme are

$$\Phi = \begin{bmatrix} \phi_1^0 & \phi_1^+ \\ \phi_2^- & \phi_2^0 \end{bmatrix}, \quad \Phi' = \begin{bmatrix} \phi_1'^0 & \phi_1'^+ \\ \phi_2'^- & \phi_2'^0 \end{bmatrix}, \quad \eta^- . \tag{3}$$

The Φ and Φ' provide charged-lepton masses and the η is needed to generate radiative Dirac-type masses for the neutrinos. To ensure that the proper neutrino mass scale arises at one-loop level and that the τ neutrino is protected from the seesaw mechanism, we impose two global symmetries $U(1)'$ and $U(1)''$. The $U(1)''$ symmetry also implies that Φ_s does not contribute to lepton masses. The transformation properties of the fermions and Higgs bosons under the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ and the global symmetries $U(1)'$ and $U(1)''$ are given in Table I.

The Yukawa couplings may be broken into two groups: those that give Dirac-type masses,

$$L = \bar{\psi}_{Li}(h_1^{ij}\Phi + h_2^{ij}\tilde{\Phi} + h_1'^{ij}\Phi' + h_2'^{ij}\tilde{\Phi}')\psi_{Rj} + \bar{q}_{Li}(h_{q1}^{ij}\Phi_s + h_{q2}^{ij}\tilde{\Phi}_s)q_{Rj} + \text{H.c.}, \tag{4}$$

and those that give Majorana masses for the neutrinos,

$$L = h_M^{ij}(\psi_{Li}^T C \Delta_L \psi_{Lj} + \psi_{Ri}^T C \Delta_R \psi_{Rj}) + \text{H.c.}, \tag{5}$$

where C is the charge-conjugation matrix. The summation over generation indices is assumed and $\tilde{\Phi} = \tau_2 \Phi^* \tau_2$. The neutrino masses can be described by a 6×6 mass matrix

$$M_\nu = \begin{bmatrix} 0 & m \\ m^T & M \end{bmatrix}, \tag{6}$$

where m is the 3×3 Dirac mass matrix which comes from Eq. (4) and M is the 3×3 Majorana mass matrix for the right-handed neutrinos which comes from Eq. (5). In our model the left-handed Majorana mass matrix vanishes, and M has rank 2 as required by the Glashow scenario.

The global symmetries cause the h_1 and h_1' couplings in Eq. (4) to vanish, so m is proportional to the VEV's $\langle \phi_2^0 \rangle$ and $\langle \phi_2'^0 \rangle$, while the charged-lepton masses are proportional to $\langle \phi_1^0 \rangle$ and $\langle \phi_1'^0 \rangle$. The VEV's $\langle \phi_2^0 \rangle$ and $\langle \phi_2'^0 \rangle$ are zero at the tree level (how this occurs will be discussed in Sec. II), so m must be generated radiatively; this may be done by one-loop graphs involving the new Higgs field η^-. This general technique was used in Ref. [7] to generate small neutrino masses.

That the 17-keV neutrino mass is not substantially altered by the seesaw mechanism is a natural consequence of the global symmetries. The right-handed τ neutrino transforms nontrivially under the $U(1)'$ symmetry, while the right-handed electron and muon neutrinos do not, which implies that the matrix M has rank 2. The right-handed electron and muon neutrinos acquire a large Majorana mass, which then forces the left-handed electron and muon neutrino masses to be very small via the seesaw mechanism. The τ neutrino may then be identified as the 17-keV neutrino. Another consequence of the $U(1)'$ symmetry is that the mixing between the electron neutrino and the 17-keV neutrino cannot come from the Yukawa couplings of Φ, but instead must come from Φ'. This implies that the charged-lepton masses also come from two distinct VEV's, which leads to FCNC's in the lepton sector. While most of these FCNC's are eliminated by a partial Cabibbo mechanism, the mass matrix requires some fine-tuning to avoid these limits.

There are further limits from cosmological and astrophysical constraints [6]. In this model the lifetime of the 17-keV neutrino is approximately 10^{11} sec, and the dominant decay mode is light neutrino plus massless Majoron. This evades most of the cosmological bounds. However, it is in conflict with the more model-dependent limit obtained from structure formation in the Universe. There also exists a very severe bound from an analysis of primordial nucleosynthesis on the mass splitting of a pseudo Dirac 17-keV neutrino with a sterile component. While this is a problem for other models of this type, in our model the generational hierarchy of the lepton masses naturally leads to a very small splitting.

In Sec. II we briefly discuss the Higgs sector, with particular emphasis on the two massless Majorons in the model and the radiative generation of the 17-keV mass scale. In Sec. III we examine the constraints on the neutrino masses and mixing, which arise from the known spectrum of charged-lepton masses, FCNC's, neutrino oscillations, and solar-neutrino data. In particular we examine which of these constraints may be met naturally and which, if any, require fine-tuning. We also examine the constraints on neutrino decay properties from cosmology and astrophysics. In Sec. IV we discuss our results.

II. HIGGS SECTOR

We will write all neutral Higgs bosons as the sum of a real and an imaginary part; for example,

$$\phi_1^0 = \frac{\phi_{1r} + i\phi_{1i}}{\sqrt{2}}, \tag{7}$$

and similarly for ϕ_2^0, $\phi_1'^0$, $\phi_2'^0$, ϕ_{s1}^0, ϕ_{s2}^0, Δ_L^0, and ΔR^0. The VEV's are defined as

$$\begin{aligned}
&\langle \phi_{1r} \rangle = \kappa_1, \langle \phi_{2r} \rangle = \kappa_2 \ , \\
&\langle \phi'_{1r} \rangle = \kappa'_1, \langle \phi'_{2r} \rangle = \kappa'_2 \ , \\
&\langle \phi_{s1r} \rangle = \kappa_{s1}, \langle \phi_{s2r} \rangle = \kappa_{s2} \ , \\
&\langle \Delta_{Lr} \rangle = v_L, \langle \Delta_{Rr} \rangle = v_R \ .
\end{aligned} \tag{8}$$

In Ref. [9] it is argued that the VEV's must be relatively real in the standard left-right model if we assume that no parameters in the Higgs potential are unnaturally small. A similar argument can be made in the more complicated model considered here, and we assume that all of the VEV's in Eq. (8) are real.

Given the most general Higgs potential allowed by the symmetries in Table I, it is possible to choose the VEV's κ_2, κ'_2, and v_L to be zero at the tree level. This is the simplest choice which is consistent with the minimization conditions and also allows the Glashow scheme with a radiative 17-keV mass. The nonzero VEV's κ_1, κ'_1, κ_{s1}, and κ_{s2} break $SU(2)_L$ symmetry and therefore must have a scale no larger than the standard model VEV; v_R breaks $SU(2)_R$ and sets the scale for the extra W_R and Z_R bosons. It is convenient to define $\kappa_\chi \equiv (\kappa_1^2 + \kappa_1'^2)^{1/2}$, $\kappa_s \equiv (\kappa_{s1}^2 + \kappa_{s2}^2)^{1/2}$, and $\kappa_L \equiv (\kappa_\chi^2 + \kappa_s^2)^{1/2}$; then the standard W mass is $g_L \kappa_L / 2$. We will assume that the right-handed scale is of order 1 TeV or larger.

The tree-level physical-Higgs-boson mass spectrum is in general quite complicated; here we summarize the results [10]. After symmetry breaking there are 30 degrees of freedom in the physical Higgs sector: two doubly charged Higgs bosons, six singly charged Higgs bosons, eight scalars, and six pseudoscalars. Most of these states acquire a mass proportional to the right-handed breaking scale v_R, which we henceforth will designate as "heavy." There are two singly charged and three neutral scalars that acquire a mass at the ordinary weak scale (i.e., proportional to κ_1, κ_{s1}, κ_{s2}, or κ'_1). Finally, there are two massless Majorons in the model:

$$\begin{aligned}
&\chi_1 = (\kappa'_1 \phi_{1i} - \kappa_1 \phi'_{1i}) / \kappa_\chi \ , \\
&\chi_2 = (\kappa_1 \kappa_s \phi_{1i} + \kappa'_1 \kappa_s \phi'_{1i} - \kappa_\chi^2 \phi_{si-}) / \kappa_\chi \kappa_L \ ,
\end{aligned} \tag{9}$$

where $\phi_{si-} = (\kappa_{s1} \phi_{s1i} - \kappa_{s2} \phi_{s2i}) / \kappa_s$. The interaction of the Majoron χ_1 with leptons has important consequences which are discussed in Sec. III F.

Although they are zero at the tree level, κ_2 and κ'_2 will acquire the nonzero VEV from graphs such as those depicted in Fig. 1. The scale of these contributions are expected to be of order [7]

$$\kappa_2 \simeq \frac{\gamma_i \gamma_j \gamma_k}{16\pi^2} \frac{v_R^2 \kappa_1'^2 \kappa_1}{m_{\phi_2}^2 m_{\delta_R}^2} \ , \quad \kappa'_2 \simeq \frac{\gamma_i \gamma_j \gamma_k}{16\pi^2} \frac{v_R^2 \kappa_1^2 \kappa'_1}{m_{\phi 2}^2 m_{\delta R}^2} \ , \tag{10}$$

where γ_i, γ_j, and γ_k are couplings in the Higgs potential (expected to be of the order of a gauge coupling) and the Higgs-boson masses m_{ϕ_2} and m_{δ_R} in the denominator of

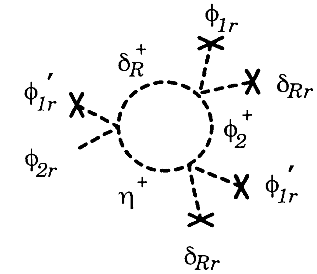

FIG. 1. Typical one-loop diagram which gives a nonzero vacuum expectation value to ϕ_2^0. The similar diagram for $\phi_2'^0$ can be found by interchanging ϕ and ϕ'.

Eq. (10) are heavy. Since $\kappa_1 \simeq \kappa'_1 \simeq 100$ GeV and $v_R \simeq m_{\phi_2} \simeq m_{\delta_2} \simeq 1$ TeV, $\kappa_2 \simeq \kappa'_2 \simeq 10^{-6}\kappa \simeq 100$ keV. These one-loop contributions then set the scale for the 17-keV neutrino mass. A careful check of the Higgs Lagrangian shows that under the symmetries in Table I these are the simplest one-loop diagrams that can give a nonzero value for κ_2 and κ'_2.

III. CONSTRAINTS ON NEUTRINO MASSES AND MIXINGS

A. Diagonalizing the neutrino mass matrix

In order to analyze the constraints on the neutrino masses and mixings, it is useful to diagonalize the neutrino mass matrix M_ν in Eq. (6). If we write M_ν in the $(\nu_{eL}, \nu_{\mu L}, \nu_{\tau L}, \nu_{\tau R}, \nu_{eR}, \nu_{\mu R})$ basis, then

$$m = \frac{1}{\sqrt{2}} \begin{bmatrix} h_2'^{13}\kappa'_2 & h_2^{11}\kappa_2 & h_2^{12}\kappa_2 \\ h_2'^{23}\kappa'_2 & h_2^{21}\kappa_2 & h_2^{22}\kappa_2 \\ h_2'^{33}\kappa'_2 & h_2^{31}\kappa_2 & h_2^{32}\kappa_2 \end{bmatrix} \equiv \begin{bmatrix} m_{13} & m_{11} & m_{12} \\ m_{23} & m_{21} & m_{22} \\ m_{33} & m_{31} & m_{32} \end{bmatrix} \tag{11}$$

and

$$M = \frac{v_R}{\sqrt{2}} \begin{bmatrix} 0 & 0 & 0 \\ 0 & h_M^{11} & h_M^{12} \\ 0 & h_M^{21} & h_M^{22} \end{bmatrix} \equiv \begin{bmatrix} 0 & 0 & 0 \\ 0 & M_{11} & M_{12} \\ 0 & M_{12} & M_{22} \end{bmatrix} . \tag{12}$$

In our notation we have changed the conventional order of the right-handed neutrinos in the matrix M_ν (in order to assist diagonalization), but the numerical subscripts of the m_{ij} and M_{ij} still represent the generation number, i.e., $1 = \nu_e$, $2 = \nu_\mu$, and $3 = \nu_\tau$. Note that since M is a Majorana mass matrix, it is symmetric. We define the charged-lepton mass matrix

$$m_l = \frac{1}{\sqrt{2}} \begin{bmatrix} h_2^{11}\kappa_1 & h_2^{12}\kappa_1 & h_2'^{13}\kappa_1' \\ h_2^{21}\kappa_2 & h_2^{22}\kappa_1 & h_2'^{23}\kappa_1' \\ h_2^{31}\kappa_1 & h_2^{32}\kappa_2 & h_2'^{33}\kappa_1' \end{bmatrix} \equiv \begin{bmatrix} m_{ee} & m_{e\mu} & m_{e\tau} \\ m_{\mu e} & m_{\mu\mu} & m_{\mu\tau} \\ m_{\tau e} & m_{\tau\mu} & m_{\tau\tau} \end{bmatrix} . \tag{13}$$

Since m in Eq. (11) comes from the same Yukawa couplings as m_l, each element of m is related to an element in m_l by a ratio of VEV's, e.g., $m_{22} = m_{\mu\mu}\kappa_2/\kappa_1$ and $m_{33} = m_{\tau\tau}\kappa_2'/\kappa_1'$. In this paper we assume that $\kappa_1 \simeq \kappa_1'$ are at the electroweak scale; then $\kappa_2 \simeq \kappa_2'$ are generated at the one-loop level. Thus we expect $\kappa_2/\kappa_1 \simeq \kappa_2'/\kappa_1'$, which implies that elements of the Dirac part of the neutrino mass matrix m possess the same generational hierarchy as those of m_l. Therefore m_{33} should be the largest element of m, and the eigenvalues of m should exhibit the same hierarchy as the e, μ, and τ masses. This inherited structure for m is another constraint on the model, but it also can account for the narrowness of the mass splitting of the pseudo Dirac neutrino required by primordial nucleosynthesis (see Sec. III F). It is useful to define a matrix ϵ which relates the elements of m to its largest value

$$\epsilon_{ij} = m_{ij}/m_{33} \ ; \tag{14}$$

then the charged-lepton mass matrix is given approximately by $m_l \simeq m_\tau \epsilon$, where m_τ is the τ mass. Henceforth, limits on the charged- and neutral-lepton mass matrices will be expressed in terms of the matrix ϵ.

We may write M_ν in the form

$$M_\nu = \begin{bmatrix} A & B \\ B^T & H \end{bmatrix} , \tag{15}$$

where

$$A = \begin{bmatrix} 0 & 0 & 0 & m_{13} \\ 0 & 0 & 0 & m_{23} \\ 0 & 0 & 0 & m_{33} \\ m_{13} & m_{23} & m_{33} & 0 \end{bmatrix} ,$$

$$B = \begin{bmatrix} m_{11} & m_{12} \\ m_{21} & m_{22} \\ m_{31} & m_{32} \\ 0 & 0 \end{bmatrix} , \quad H = \begin{bmatrix} M_{11} & M_{12} \\ M_{12} & M_{22} \end{bmatrix} . \tag{16}$$

We note that A and B are of the order of the 17-keV scale and H is at the $SU(2)_R$-breaking scale of 1 TeV. To diagonalize M_ν, we follow the procedure of Ref. [7] in which M_ν is block diagonalized into a 2×2 heavy sector and a 4×4 light sector. Then we block diagonalize the 4×4 light sector to get

$$M_\nu^{\text{light}} \simeq \begin{bmatrix} \tilde{B}_{11} & \tilde{B}_{12} \\ \tilde{B}_{21} & \tilde{B}_{22} \end{bmatrix} , \quad M_\nu^{\text{int}} \simeq \begin{bmatrix} \tilde{B}_{33} & m_{33} \\ m_{33} & 0 \end{bmatrix} ,$$
$$M_\nu^{\text{heavy}} \simeq H , \tag{17}$$

where

$$\tilde{B}_{ij} = [m_{i1}m_{j1}M_{22} + m_{i2}m_{j2}M_{11} - (m_{i1}m_{j2} + m_{i2}m_{j1})M_{12}]/(M_{11}M_{22} - M_{12}^2) . \tag{18}$$

The matrix M_ν^{light} describes the two very light neutrinos, M_ν^{int} the two states of the 17-keV neutrino, and M_ν^{heavy} the two heavy neutrinos [10].

B. The 17-keV neutrino mass and mixing

The mass matrix M_ν^{int} describes a pseudo Dirac neutrino with mass given approximately by $(m_{13}^2 + m_{23}^2 + m_{33}^2)^{1/2} \simeq m_{33}$ and mass splitting of order

$$\Delta m_{17} \simeq 2(\tilde{B}_{13}m_{13} + \tilde{B}_{23}m_{23})/m_{33} , \tag{19}$$

where $\tilde{B}_{13}$ and $\tilde{B}_{23}$ are defined in Eq. (18). The mixing coefficient of the 17-keV neutrino with ν_{eL} and $\nu_{\mu L}$ is given approximately by the first and second rows, respectively, of the 2×2 matrix σX^i, where

$$\sigma = \frac{1}{m_{33}^2} \begin{bmatrix} m_{13}m_{33} & m_{33}\tilde{B}_{13} - m_{13}\tilde{B}_{33} \\ m_{23}m_{33} & m_{33}\tilde{B}_{23} - m_{23}\tilde{B}_{33} \end{bmatrix} \tag{20}$$

and X^i is the two-dimensional rotation matrix that diagonalizes M_ν^{int} (with a rotation angle of approximately 45°). The column labels in Eq. (20) correspond to the lower and upper eigenstates of the 17-keV neutrino. Therefore ν_e-ν_{17} mixing is of order ϵ_{13} and ν_μ-ν_{17} mixing is of order ϵ_{23}. Since the $\tilde{B}_{ij}$ are the order of m^2/M, as is M_ν^{light}, the mass splitting of the 17-keV neutrino is then *a priori* of the order of the light-neutrino masses times an e-τ or μ-τ mixing element.

To account for the positive 17-keV neutrino results, we expect

$$\epsilon_{13} \simeq m_{13}/m_{33} \simeq 10^{-1} . \tag{21}$$

Limits on ν_e disappearance from accelerator experiments [11] are barely evaded for this amount of mixing. Similar limits on ν_μ-ν_τ oscillations constrain the ν_μ-ν_τ mixing angle to be approximately 10^{-2} or smaller [11]; this implies

$$\epsilon_{23} \simeq m_{23}/m_{33} < 10^{-2} . \tag{22}$$

There are no limits from ν_e-ν_μ oscillations for the very small mass differences in M_ν^{light}.

C. Constraints due to charged-lepton masses

Because of the relationship that exists between the charged-lepton and neutrino masses, the charged-lepton mass spectrum and mixing place constraints on the possible structure of the matrix ϵ. For example, the determinant of m_l of Eq. (13) must correspond to the product of the eigenvalues

$$\begin{aligned} \text{Det} m_l &\simeq (\epsilon_{11}\epsilon_{22} - \epsilon_{12}\epsilon_{21} + \epsilon_{12}\epsilon_{23}\epsilon_{31} + \epsilon_{21}\epsilon_{13}\epsilon_{32} \\ &\quad - \epsilon_{13}\epsilon_{31}\epsilon_{22} - \epsilon_{23}\epsilon_{32}\epsilon_{11})m_\tau^3 \\ &\simeq m_e m_\mu m_\tau , \end{aligned} \tag{23}$$

where we have used the fact that $\epsilon_{33} \simeq 1$. If we assume that all diagonal terms are the same order of magnitude as one of the eigenvalues, then $\epsilon_{22} \simeq m_\mu/m_\tau, \epsilon_{11} \simeq m_e/m_\tau$, and, barring fortuitous cancellations in Eq. (23), we can infer the approximate limits

$$\epsilon_{12}\epsilon_{21}<2\times10^{-5}, \quad \epsilon_{31}<4\times10^{-3} . \tag{24}$$

There are no significant limits on ϵ_{32} from Eq. (23). The value of the middle eigenvalue, i.e., the muon mass, constrains ϵ_{12} and ϵ_{21} both to be no greater than about m_μ/m_τ.

D. Flavor-changing neutral currents

FCNC decays of the charged leptons mediated by a Majoron can be a potential source of difficulty for the model since there is no experimental evidence for such decays. The simplest and most dangerous FCNC decays are $\mu\to e\chi$ and $\tau\to e\chi$ or $\mu\chi$. The limits on the branching ratio for these decays are [12]

$$\begin{aligned} &B(\mu\to e\chi)<2.6\times10^{-6} ,\\ &B(\tau\to e\chi)<7.1\times10^{-3} ,\\ &B(\tau\to\mu\chi)<2.2\times10^{-2} , \end{aligned} \tag{25}$$

where it is assumed that the Majoron escapes undetected. The theoretical decay rate for $\mu\to e\chi$ is

$$\frac{\Gamma(\mu\to e\chi)}{\Gamma(\mu\to e\bar{\nu}_e\nu_\mu)}=\frac{12\pi^2h_{e\mu}^2}{G_F^2m_\mu^4} , \tag{26}$$

where $h_{e\mu}$ is the flavor-changing coupling. Similar formulas hold for $\tau\to e\chi$ and $\tau\to\mu\chi$. From the limits in Eq. (25), we find

$$h_{e\mu}<2\times10^{-11}, \quad h_{e\tau}<7\times10^{-7}, \quad h_{\mu\tau}<10^{-6} . \tag{27}$$

Other possible FCNC decays include three-body decays mediated by a virtual Majoron. The rate for $\mu\to e\bar{e}e$, for instance, is proportional to $h_{e\mu}^2h_{ee}^2$. If the diagonal coupling h_{ee} has order of magnitude m_e/M_W (in order to provide the proper size for the electron mass), then the constraint on $h_{e\mu}$ from the experimental limit on $\mu\to e\bar{e}e$ is less severe than the limit in Eq. (27). Similar statements hold for three-body FCNC decays of the τ.

We now see how restrictive the limits of Eq. (27) are on our model. The couplings of the charged leptons to the Majorons in Eq. (4) may be written in the form

$$\begin{aligned} L=&-i\frac{\sqrt{2}}{\kappa_\chi}\overline{e_L^m}m_l^{\text{diag}}e_R^m\left[\frac{\kappa_1'}{\kappa_1}\chi_1+\frac{\kappa_s}{\kappa_L}\chi_2\right]\\ &+i\frac{\kappa_\chi}{\kappa_1}\overline{e_L^m}(V_L^\dagger h_2'V_R)e_R^m\chi_1+\text{H.c.}, \end{aligned} \tag{28}$$

where m_l^{diag} is the diagonalized charged-lepton mass matrix, e_L^m and e_R^m are the mass eigenstates of the charged leptons, and V_L and V_R are the diagonalization matrices for the left- and right-handed charged leptons, respectively.

The first term of Eq. (28) is flavor diagonal, but because the charged-lepton masses do not come exclusively from the h_2' couplings, the second term could have FCNC's. However, because of the form of h_2' (the only nonzero elements are $h_2'^{j3}$), the FCNC's are somewhat suppressed by a partial Cabibbo mechanism. An analytic expansion of the FCNC in inverse powers of m_τ may be found by applying the block diagonalization procedure of Sec. III A. We find that the largest FCNC contributions are [10]

$$\begin{aligned} &(V_L^\dagger h_2'V_R)_{\mu e}\simeq\frac{m_{\tau e}m_{\tau\mu}m_\mu}{m_\tau^2m_W} ,\\ &(V_L^\dagger h_2'V_R)_{\mu e}\simeq\frac{m_{\mu\tau}m_{\tau e}m_\mu^2}{m_\tau^3m_W} ,\\ &(V_L^\dagger h_2'V_R)_{\tau\mu}\simeq-\left[\frac{m_{\tau\mu}}{m_\tau}+10^{-1}\frac{m_{e\mu}}{m_\tau}\right]\frac{m_\tau}{m_W} . \end{aligned} \tag{29}$$

The limits in Eq. (27) can be expressed approximately as

$$\begin{aligned} &(V_L^\dagger h_2'V_R)_{\mu e}\simeq10^{-9}m_\tau/m_W ,\\ &(V_L^\dagger h_2'V_R)_{\tau e}\simeq(3\times10^{-5})m_\tau/m_W ,\\ &(V_L^\dagger h_2'V_R)_{\tau\mu}\simeq(4\times10^{-5})m_\tau/m_W . \end{aligned} \tag{30}$$

Then applying the bounds in Eq. (29) to the rotated Yukawa couplings in Eq. (30) gives the limits

$$\begin{aligned} &\epsilon_{31}<3\times10^{-5}\simeq10^{-1}m_e/m_\tau ,\\ &\epsilon_{32}<4\times10^{-5}\simeq10^{-1}m_e/m_\tau ,\\ &\epsilon_{12}<4\times10^{-4}\simeq m_e/m_\tau . \end{aligned} \tag{31}$$

Although ϵ_{12} is not abnormally small (in the context of the other elements of the matrix), ϵ_{31} and ϵ_{32} require a suppression of approximately one order of magnitude below the next-smallest element (which comes from the electron mass). This indicates some degree of fine-tuning is necessary to completely satisfy the FCNC constraints.

A natural question to ask at this point is whether another choice of global symmetries would eliminate the need for fine-tuning to meet the FCNC constraints. One such example is to assume that the lepton doublet Ψ_{L3} has a U(1)$'$ quantum number of unity (just like the Ψ_{R3} lepton doublet); then $m_{\tau e}$ and $m_{\tau\mu}$ are forbidden by the global symmetries, which would appear to automatically satisfy the first and second limits in Eq. (29). However, the FCNC expression also changes; the limit of Eq. (27) now requires $m_{e\mu}$ to be less than $10^{-2}m_e$. So this alternative scenario requires even more fine-tuning, and it rules out any hope of obtaining appreciable light-neutrino mixing. Other quantum number assignments run into similar difficulties. Therefore it appears that the quantum numbers given in Table I are the most desirable since they can give some light-neutrino mixing with a small amount of fine-tuning.

We conclude that if the matrix ϵ has the approximate form

$$\epsilon\simeq\frac{1}{m_\tau}\begin{bmatrix} m_e & m_e & 10^{-1}m_\tau\\ m_e & m_\mu & 10^{-2}m_\tau\\ 10^{-1}m_e & 10^{-1}m_e & m_\tau \end{bmatrix} , \tag{32}$$

then the constraints on the lepton mass matrices and mixings will be met. The off-diagonal entries in Eq. (32) are upper limits, so they could be smaller (except for ϵ_{13}, which is fixed by the amount of ν_e-ν_{17} mixing).

E. Solar-neutrino masses and mixing

Finally, we examine the additional constraint imposed by requiring that the light-neutrino masses and mixing explain the solar-neutrino puzzle [13]. The light-neutrino mass matrix is given approximately by

$$M_\nu^{\text{light}} \simeq C \simeq \frac{m_{17}^2}{\text{Det}H}\left[\begin{bmatrix}\epsilon_{11}^2 & \epsilon_{11}\epsilon_{21}\\ \epsilon_{11}\epsilon_{21} & \epsilon_{21}^2\end{bmatrix}M_{22}+\begin{bmatrix}\epsilon_{12}^2 & \epsilon_{12}\epsilon_{22}\\ \epsilon_{12}\epsilon_{22} & \epsilon_{22}^2\end{bmatrix}M_{11}-\begin{bmatrix}2\epsilon_{11}\epsilon_{12} & \epsilon_{11}\epsilon_{22}+\epsilon_{12}\epsilon_{21}\\ \epsilon_{11}\epsilon_{22}+\epsilon_{12}\epsilon_{21} & 2\epsilon_{21}\epsilon_{22}\end{bmatrix}M_{12}\right]. \quad (33)$$

Given the mixing scheme of Eq. (32), the dominant terms in Eq. (33) are those proportional to M_{11} and lead to solar-neutrino masses of order $(m_{17}m_\mu/m_\tau)^2/M_H$, where M_H is a typical mass for the heavy right-handed neutrinos, and solar-neutrino mixing angle $\arctan(\epsilon_{12}/\epsilon_{22})$. Since the scale M_H derives from the right-handed VEV $\langle\Delta_R\rangle \simeq 1$ TeV, one might expect $M_H \simeq 100$ GeV. If this is so, then the light-neutrino masses are of order 10^{-5} eV, which gives a mass-squared difference of $\delta m^2 \simeq 10^{-10}$ eV^2, appropriate for long-wavelength oscillations [14]. However, $\epsilon_{12}/\epsilon_{22} \lesssim m_e/m_\mu$, so the large mixing required for the long-wavelength scenario is not possible in this model. If, on the other hand, $M_H \simeq 1$ GeV, then the light-neutrino mass is of order 10^{-3} eV and $\delta m^2 \simeq 10^{-6}$ eV^2, appropriate for resonant Mikheyev-Smirnov-Wolfenstein (MSW) oscillations in the Sun. The solar-neutrino mixing parameter $\sin^2 2\theta$ has a maximum value of about 2×10^{-4}, about an order of magnitude too small to provide enough mixing for resonant MSW oscillations in the Sun. Thus the model does not seem to be able to include either of the popular neutrino oscillation explanations to the solar-neutrino deficit.

One might ask if a right-handed neutrino with mass of order 1 GeV should already have been detected. Since it is sterile under $\text{SU}(2)_L\times\text{U}(1)_{B-L}$, it can be produced only via mixing with a light neutrino (which is very small), or by mixing of the standard Z boson with the much heavier Z' boson (which is also small). Furthermore, even if it is produced, the lifetime is long enough that it could very likely leave the detector before decaying, thereby escaping detection. This heavy neutrino is too heavy to be a factor in the mass density of the Universe. Therefore we conclude that such a relatively light right-handed neutrino may exist. The best place to observe such a particle would be in the decays of the Z' boson, where it would contribute a standard neutrino partial width.

F. Other astrophysical constraints

There is a constraint on the mass of a Dirac or pseudo Dirac neutrino from the cooling rate of a supernova. If the right-handed component of the 17-keV neutrino is sterile, it can escape and cool the core too rapidly if the mass is too large. The exact value of this limit is uncertain; values from 1 to 28 keV have been obtained [15]. While the lower limit appears to marginally rule out the pseudo Dirac scenario with a sterile right-handed component, it is close enough to being allowed that we do not believe that models of this type should be abandoned until a more decisive limit is obtained.

There is also a very strong constraint from primordial nucleosynthesis on the mass splitting of a pseudo Dirac ν_{17} which has a sterile component [16]:

$$\Delta m_{17} < 1.8\times 10^{-7}\ \text{eV}\ . \quad (34)$$

This is much more severe than the usual constraint from double-β decay. The mass splitting can be calculated from Eq. (19), using the approximate form for ϵ in Eq. (32). If M is a generic heavy-neutrino mass then

$$\Delta m_{17} \simeq \frac{2m_{12}m_{32}m_{13}}{m_{33}M} \simeq 2\times 10^{-2}\frac{m_e^2}{m_\tau^2}\frac{m_{33}^2}{M} \simeq \frac{0.5\ \text{eV}^2}{M}\ . \quad (35)$$

The bound of Eq. (34) is met without any additional constraint on ϵ if $M > 3$ MeV. Since M is proportional to v_R this is not a problem. Hence, the smallness of the elements m_{12} and m_{32} compared to m_{33} implies that the bound in Eq. (34) does not present any serious difficulties for the model.

The Majoron in the lepton section, χ_1, has many important consequences. In addition to its participation in flavor-changing processes (see Sec. III D), the principal decay of the 17-keV neutrino is to a light neutrino plus χ_1 via a one-loop diagram involving Higgs bosons (see Fig. 2). The approximate lifetime is given by

$$\tau_{17}(\nu_{17}\to\nu\chi_1) \simeq \left[\frac{1}{8\pi}\right]^5\left[\frac{m_{17}\kappa_\chi}{m_\phi^2}\right]^2\left[\frac{\lambda}{2}h_2^{33}h_2'^{13}\right]^2 m_{17} \simeq 10^{11}\ \text{sec}\ , \quad (36)$$

where m_ϕ is the mass scale of $\phi_2^\pm$ and $\phi_2'^\pm$, i.e., of the order of M_W, and λ is a parameter from the Higgs potential which determines the Higgs coupling to χ_1.

The relevant excluded ranges of τ_{17} are [6]

$$\tau_{17} > 8.4\times 10^{11}\ \text{sec} \quad (37\text{a})$$

(energy density) [6],

$$3\times 10^{-4}\ \text{sec} < \tau_{17} < 2\times 10^{8}\ \text{sec} \quad (37\text{b})$$

(supernova pulse) [17],

$$10^{-12}\ \text{sec} < \tau_{17} < 10^{-4}\ \text{sec} \quad (37\text{c})$$

(supernova cooling) [6],

$$\tau_{17} > 7\times 10^{6}\ \text{sec} \quad (37\text{d})$$

(structure formation in Universe) [18]. We see that all of

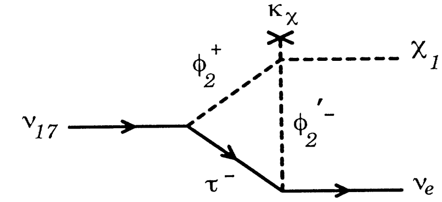

FIG. 2. The dominant diagram for the decay of the 17-keV neutrino.

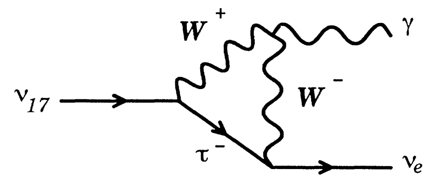

FIG. 3. The dominant diagram for the radiative decay of the 17-keV neutrino.

these limits are evaded except for the last, which is, however, subject to much uncertainty.

There are also constraints on the decay of ν_{17} into a light neutrino plus photon [6]. The measured γ-ray flux from SN 1987A gives the bound

$$\tau(\nu_{17} \to \nu\gamma) > 10^{15} \text{ sec} . \tag{38}$$

In the current model, this decay can occur at the one-loop level; the relevant diagrams, which involve Higgs bosons, are shown in Fig. 3. The radiative lifetime is approximately

$$\tau(\nu_{17} \to \nu\gamma) \simeq \left[\frac{9\alpha G_F^2}{2048} m_{17}^5 \epsilon_{13}^2 \left[\frac{m_\tau}{M_W} \right]^2 \right]^{-1} \simeq 10^{22} \text{ sec} , \tag{39}$$

so that this limit is easily met.

IV. DISCUSSION

In this paper we have attempted to construct a left-right model which includes the 17-keV neutrino in a natural way. The existence of three left-handed and three right-handed neutrinos suggests that one can use a scenario in which the 17-keV neutrino is a pseudo Dirac neutrino, consisting primarily of $\nu_{\tau L}$ and $\nu_{\tau R}$, which avoids the usual seesaw mechanism because the mass matrix of the heavy right-handed Majorana neutrinos has rank 2. The left-handed neutrinos ν_{eL} and $\nu_{\mu L}$ and the right-handed neutrinos ν_{eR} and $\nu_{\mu R}$ undergo the seesaw mechanism and become very light and very heavy, respectively. In order to explain the smallness of the 17-keV scale we must assume that the VEV's which normally set this scale are zero at the tree level, so they must be generated at the one-loop level.

In order to achieve these conditions naturally one must impose two global symmetries on the Lagrangian. These global symmetries in turn require that new Higgs bosons be introduced (in addition to those usually present in a standard left-right model) to give the proper form to the neutrino mass matrix and to allow the appropriate VEV's to achieve a nonzero value at the one-loop level. The existence of more than one set of Higgs bosons implies the existence of flavor-changing neutral currents, and a small amount of fine-tuning is necessary to meet all of the phenomenological bounds. The mass of the light neutrinos is at a scale that in principle could explain the solar-neutrino deficit through neutrino oscillations; however, the FCNC limits appear to make the light-neutrino mixing too small. Therefore the solution to the solar-neutrino problem must come from somewhere else.

Finally, there are many astrophysical and cosmological constraints on the model. Most of these are obeyed by our model; the ones that are not (the limit on τ_{17} from the large-scale structure formation of the Universe and the limit on the pseudo Dirac mass from supernova cooling) are not yet on a firm theoretical footing, but could rule out the model in the future should a more precise calculation confirm the current limits.

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