New model for the neutrino mass matrix

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I suggest a model based on a softly broken symmetry $L_e - L_\mu - L_\tau$ and on Babu's mechanism for the two-loop radiative generation of neutrino masses. The model predicts that one of the physical neutrinos (ν_3) is massless and that its component along the ν_e direction (U_{e3}) is zero. Moreover, if the soft-breaking term is assumed to be very small, then the vacuum oscillations of ν_e have almost maximal amplitude and solve the solar-neutrino problem. New scalars are predicted in the 10 TeV energy range, and a breakdown of e-μ-τ universality should not be far from existing experimental bounds.

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In a model without right-handed (singlet) neutrinos, the three weak-interaction-eigenstate neutrinos ν_e, ν_μ, and ν_τ may acquire $|\Delta I|=1$ Majorana masses given by the following term in the Lagrangian:

$$\mathcal{L}_{\rm mass}^{(\nu)} = \frac{1}{2}(\nu_e^T \;\; \nu_\mu^T \;\; \nu_\tau^T) C^{-1} \mathcal{M} \begin{pmatrix} \nu_e \\ \nu_\mu \\ \nu_\tau \end{pmatrix} - \frac{1}{2}(\overline{\nu_e} \;\; \overline{\nu_\mu} \;\; \overline{\nu_\tau}) C \mathcal{M}^* \begin{pmatrix} \overline{\nu_e}^T \\ \overline{\nu_\mu}^T \\ \overline{\nu_\tau}^T \end{pmatrix}. \tag{1}$$

Here C is the Dirac-Pauli charge-conjugation matrix and $\mathcal{M}$ is a 3×3 symmetric mass matrix. One may diagonalize $\mathcal{M}$ with the help of a unitary matrix U in the following way:

$$U^T \mathcal{M} U = \mathrm{diag}(m_1, m_2, m_3), \tag{2}$$

where m_1, m_2, and m_3 are real and non-negative. The physical neutrinos ν_1, ν_2, and ν_3 are given by

$$\begin{pmatrix} \nu_e \\ \nu_\mu \\ \nu_\tau \end{pmatrix} = U \begin{pmatrix} \nu_1 \\ \nu_2 \\ \nu_3 \end{pmatrix}. \tag{3}$$

Then,

$$\mathcal{L}_{\rm mass}^{(\nu)} = \frac{1}{2}\sum_{i=1}^{3} m_i (\nu_i^T C^{-1} \nu_i - \overline{\nu_i} C \overline{\nu_i}^T). \tag{4}$$

Experiment indicates that two linearly independent squared-mass differences among the three physical neutrinos differ by a few orders of magnitude. Indeed, $\Delta m^2_{\rm atm}$ is of order 10^{-3} eV^2, while $\Delta m^2_\odot$ may be either of order 10^{-5} eV^2 in the case of the Mikheev-Smirnov-Wolfenstein (MSW) solution for the solar-neutrino puzzle, 10^{-7} eV^2 in the case of the LOW solution, or 10^{-10} eV^2 in the case of the vacuum-oscillation (''just so'') solution. It is customary to identify ν_3 as the neutrino which has a mass much different from the masses of the other two, viz.,

$$|m_2^2 - m_1^2| = \Delta m^2_\odot \ll |m_3^2 - m_1^2| \approx |m_3^2 - m_2^2| \approx \Delta m^2_{\rm atm}. \tag{5}$$

Then the negative result of CHOOZ's search for ν_e oscillations [1] is interpreted as $|U_{e3}| \leqslant 0.217$, which is valid for $\Delta m^2_{\rm atm} \geqslant 2\times 10^{-3}$ eV^2.

It has been pointed out [2] that the assumption of an approximate lepton-number symmetry $\bar{L} \equiv L_e - L_\mu - L_\tau$ (where L_e is the electron number, L_μ is the muon number, and L_τ is the tau number) may constitute a good starting point for a model of the neutrino mass matrix. Indeed, if there are no $|\Delta \bar{L}|=2$ mass terms, then

$$\mathcal{M} = \begin{pmatrix} 0 & rb & b \\ rb & 0 & 0 \\ b & 0 & 0 \end{pmatrix}, \tag{6}$$

where b and r may, without loss of generality, be taken to be real and positive. The mass matrix in Eq. (6) yields $m_3=0$, $m_1=m_2=b\sqrt{1+r^2}$, and

$$U = \begin{pmatrix} \frac{1}{\sqrt{2}} & -\frac{i}{\sqrt{2}} & 0 \\ \frac{r}{\sqrt{2(1+r^2)}} & \frac{ir}{\sqrt{2(1+r^2)}} & \frac{1}{\sqrt{1+r^2}} \\ \frac{1}{\sqrt{2(1+r^2)}} & \frac{i}{\sqrt{2(1+r^2)}} & -\frac{r}{\sqrt{1+r^2}} \end{pmatrix}. \tag{7}$$

This is good for the following reasons: (1) The negative result of CHOOZ's search for ν_e oscillations gets explained through $U_{e3}=0$; (2) since $|U_{e1}|=|U_{e2}|$, vacuum oscillations of ν_e with maximal amplitude would occur were $m_1 \neq m_2$, opening way for the LOW or ''just so'' solutions of the solar-neutrino problem to apply; (3) it is intuitive to expect r to be close to 1. Now, if $r=1$, then ν_μ-ν_τ mixing is maximal, and this explains the atmospheric-neutrino anomaly.

On the other hand, $\bar{L}$ must be broken, because $m_1=m_2$ does not allow for oscillations between ν_1 and ν_2 and a solution of the solar-neutrino puzzle. A good choice, in order to avoid unpleasant majorons, would be to have $\bar{L}$ to be softly broken; this would, moreover, permit a natural explanation for $\Delta m^2_{\odot} \ll \Delta m^2_{\rm atm}$. This option has been suggested in various papers [3]; however, in those models there is no predictive power for the form of the mixing matrix U, a fact which impairs their immediate interest and experimental testability.

In this paper I put forward a simple model with softly broken $\bar{L}$ which maintains some predictive power. The model is based on Babu's mechanism for two-loop radiative generation of the neutrino masses [4]. I recall that, in general, Babu's mechanism leads to one neutrino remaining massless; however, whereas that general mechanism cannot predict the ν_e, ν_μ, and ν_τ components of the massless neutrino, the specific model that I shall put forward retains the exact-$\bar{L}$ prediction $U_{e3}=0$. Moreover, in my model there is a rationale for the ν_e oscillations of maximal amplitude and for the tiny mass difference $\Delta m^2_{\odot}$, which allow a LOW or "just so" explanation of the solar-neutrino deficit; that rationale is provided by the naturalness of the assumption that the term which breaks $\bar{L}$ softly is very small.

In my model I just introduce in the scalar sector, above and beyond the usual standard-model doublet $\phi = (\varphi^+ \varphi^0)^T$, one singly charged singlet f^+ with $\bar{L}=0$, together with two doubly charged singlets g^{2+} and h^{2+} and their Hermitian conjugates. The difference between g^{2+} and h^{2+} lies in that the former field has $\bar{L}=0$ whereas h^{2+} has $\bar{L}=-2$. The Yukawa couplings of the leptons are $\bar{L}$ invariant and are given by

$$\begin{aligned}\mathcal{L}_Y^{(1)} = & -\frac{m_e}{v}(\overline{\nu_{eL}} \ \ \overline{e_L})\begin{pmatrix}\varphi^+\\ \varphi^0\end{pmatrix} e_R - \frac{m_\mu}{v}(\overline{\nu_{\mu L}} \ \ \overline{\mu_L})\begin{pmatrix}\varphi^+\\ \varphi^0\end{pmatrix}\mu_R \\ & -\frac{m_\tau}{v}(\overline{\nu_{\tau L}} \ \ \overline{\tau_L})\begin{pmatrix}\varphi^+\\ \varphi^0\end{pmatrix}\tau_R + f^+[f_\mu(\nu_{eL}^T C^{-1}\mu_L \\ & - e_L^T C^{-1}\nu_{\mu L}) + f_\tau(\nu_{eL}^T C^{-1}\tau_L - e_L^T C^{-1}\nu_{\tau L})] \\ & + e_R^T C^{-1}[g^{2+}(g_\mu \mu_R + g_\tau \tau_R) + h^{2+} h_e e_R] + \text{H.c.},\end{aligned} \tag{8}$$

where f_μ, f_τ, g_μ, g_τ, and h_e are complex coupling constants. Notice that, in Eq. (8), I have already taken, without loss of generality, the Yukawa couplings of ϕ to be flavor diagonal; v denotes the vacuum expectation value of φ^0.

The scalar potential V has a trivial part $V_{\rm trivial}$ which is a quadratic polynomial in $\phi^\dagger\phi$, $f^- f^+$, $g^{2-} g^{2+}$, and $h^{2-} h^{2+}$. Besides, V includes two other terms, with complex coefficients λ and ϵ:

$$V = V_{\rm trivial} + (\lambda f^- f^- g^{2+} + \epsilon g^{2-} h^{2+} + \text{H.c.}). \tag{9}$$

The term with coefficient ϵ breaks $\bar{L}$ softly. I make the following assumptions: this is the only $\bar{L}$-breaking term in the theory, and ϵ is small. These assumptions are technically natural in the sense of 't Hooft [5].[1]

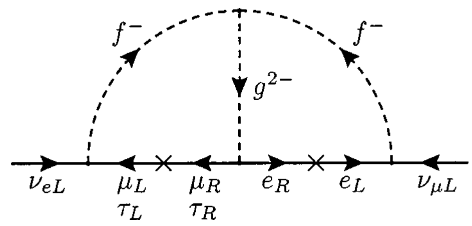

FIG. 1. Two-loop Feynman diagram which generates $\mathcal{M}_{e\mu}$.

From now on I shall assume, without loss of generality, f_μ, f_τ, g_τ, h_e, λ, and ϵ to be real and positive. Only g_μ remains, in general, complex.

The neutrino mass term $\mathcal{M}_{e\mu}$ does not break $\bar{L}$ and is generated at the two-loop level by the Feynman diagram in Fig. 1. A similar diagram generates $\mathcal{M}_{e\tau}$. In both cases, there is in the diagram an inner charged lepton which may be either μ or τ. It is clear that the mass terms thus generated obey the relation

$$r \equiv \frac{\mathcal{M}_{e\mu}}{\mathcal{M}_{e\tau}} = \frac{f_\mu}{f_\tau}. \tag{10}$$

Contrary to what happens in Zee's model [6], this ratio of mass terms is not proportional to a ratio of squared charged-lepton masses [7]. As seen before, in order to obtain maximal ν_μ-ν_τ mixing, one would like to have $r \approx 1$. In the present model, this means that the coupling constants f_μ and f_τ should be approximately equal. In Zee's model, on the other hand, one winds up with the rather unrealistic constraint $f_\mu / f_\tau \approx (m_\tau / m_\mu)^2$.

Let us check whether the diagram in Fig. 1 is able to yield neutrino masses of the right order of magnitude. As we shall see later, we would like to obtain $|\mathcal{M}_{e\mu}| \approx |\mathcal{M}_{e\tau}| \approx \sqrt{\Delta m^2_{\rm atm}} \sim 10^{-2}$–$10^{-1}$ eV. Now, from the diagram in Fig. 1 with an inner τ one obtains

$$\mathcal{M}_{e\mu} = -2\lambda f_\mu f_\tau g_\tau m_e m_\tau \frac{I}{(16\pi^2)^2}, \tag{11}$$

where

$$\begin{aligned} I = & \frac{1}{\pi^4}\int d^4k \, \frac{1}{k^2 - m_f^2}\,\frac{1}{k^2 - m_e^2} \\ & \times \int d^4 q\, \frac{1}{q^2 - m_f^2}\,\frac{1}{q^2 - m_\tau^2}\,\frac{1}{(k-q)^2 - m_g^2} \end{aligned} \tag{12}$$

[1] Notice that the possible $\bar{L}$-breaking term $f^- f^- h^{2+}$ has dimension higher than the one of $g^{2-} h^{2+}$, and therefore the assumption of its absence is natural.

$$= \frac{1}{2(m_f^2 - m_\tau^2)} \int_0^\infty \frac{dy}{(y+1)(y+x_e)} \times \left[p \ln \frac{y+x_g+1+p}{y+x_g+1-p} - p' \ln \frac{y+x_g+x_\tau+p'}{y+x_g+x_\tau-p'} + (1-x_\tau)\ln x_g + (x_\tau - x_g - y)\ln x_\tau \right]. \qquad (13)$$

Here $x_e = m_e^2/m_f^2$, $x_\tau = m_\tau^2/m_f^2$, $x_g = m_g^2/m_f^2$, and

$$p = \sqrt{(y+x_g-1)^2 + 4y}, \qquad (14)$$

$$p' = \sqrt{(y+x_g-x_\tau)^2 + 4yx_\tau}. \qquad (15)$$

The integral in Eq. (13) is convergent and may be computed numerically.[2] For $m_e, m_\tau \ll m_f$ and $m_g \approx m_f$, one finds I to be of order m_f^{-2}.

In my estimate of $\mathcal{M}_{e\mu}$, I shall therefore set $I \approx m_f^{-2}$. The bounds from e-μ-τ universality in μ decay and in τ decay are $f_\mu/m_f \lesssim 10^{-4}\,\text{GeV}^{-1}$ and $f_\tau/m_f \lesssim 10^{-4}\,\text{GeV}^{-1}$ [4]; if one allows $f_\mu f_\tau/m_f^2$ to be as high as $10^{-8}\,\text{GeV}^{-2}$, then one obtains

$$|\mathcal{M}_{e\mu}| \approx 10^{-15} \lambda g_\tau. \qquad (16)$$

It is reasonable to assume that the Yukawa coupling g_τ is of the same order of magnitude as the Yukawa couplings f_μ and f_τ, and that the dimensionful scalar-potential coupling constant λ is of the same order of magnitude as both m_f and m_g. This leads to $g_\tau/\lambda \sim f_\mu/m_f \sim 10^{-4}\,\text{GeV}^{-1}$. Fortunately, the product λg_τ stays free. In order to obtain $|\mathcal{M}_{e\mu}| \sim 10^{-2}$ eV, it is then sufficient to assume

$$\lambda \approx m_g \approx m_f \sim 10^4 \text{ GeV}, \qquad (17)$$

$$f_\mu \approx f_\tau \approx g_\tau \sim 1. \qquad (18)$$

Extra factors of order 1 may easily enhance $|\mathcal{M}_{e\mu}|$ and bring it up to the desired value 0.06 eV.

The assumption, made in Eq. (18), that the Yukawa couplings are of order 1 may seem unrealistic.[3] However, there are no experimental indications against this possibility when the masses of f^+ and of g^{2+} are assumed to be as high as 10 TeV.[4] For instance, g^{2+} mediates the unobserved decay $\tau^- \to \mu^- e^+ e^-$; however, by comparing that decay with the standard $\tau^- \to \mu^- \overline{\nu}_\mu \nu_\tau$, one easily reaches the conclusion that BR $(\tau^- \to \mu^- e^+ e^-)$ should be at least one order of magnitude below the present experimental bound, when $m_g \approx 10$ TeV and $|g_\mu g_\tau| \approx 1$. A more complicated process is $e^+e^- \to \tau^+\tau^-$, which is mediated by g^{2+} exchange in the t channel. The amplitude A for this process is

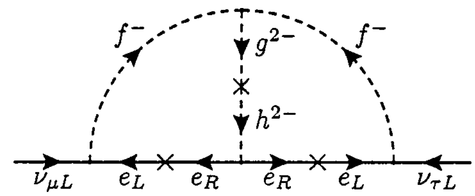

FIG. 2. Two-loop Feynman diagram which generates $\mathcal{M}_{\mu\tau}$.

$$A = \frac{ie^2}{s}[\overline{v}(e)\gamma^\mu u(e)][\overline{u}(\tau)\gamma_\mu v(\tau)] + \frac{ie^2}{3(s-m_z^2)}[\overline{v}(e)\gamma^\mu\gamma_5 u(e)][\overline{u}(\tau)\gamma_\mu\gamma_5 v(\tau)] - \frac{ig_\tau^2}{8(t-m_g^2)}[\overline{v}(e)\gamma^\mu(1+\gamma_5)u(e)] \times [\overline{u}(\tau)\gamma_\mu(1+\gamma_5)v(\tau)]. \qquad (19)$$

I have used the convenient approximations $m_e = m_\tau = 0$ and $\sin^2\theta_w = 1/4$ in writing down the standard-model amplitude, and a Fierz transformation in the nonstandard contribution. If one defines $j = 2m_g^2/s$, $z = g_\tau^2/(2e^2)$, and $l = 3(s-m_z^2)/s$, then one finds

$$\frac{d\sigma}{d\cos\theta} \propto \frac{l^2+1}{l^2}(1+\cos^2\theta) + \frac{4}{l}\cos\theta + z\frac{l+1}{l} \times \frac{(1+\cos\theta)^2}{1+j+\cos\theta} + z^2\frac{(1+\cos\theta)^2}{(1+j+\cos\theta)^2}, \qquad (20)$$

where θ is the angle between the momenta of e^- and of τ^- in the center-of-momentum frame. From the differential cross section in Eq. (20), one easily checks that the deviations of both the total cross section and the forward-backward asymmetry from their standard-model predictions are completely negligible when $m_g \sim 10$ TeV, even if g_τ is as large as 1.

Except for $\mathcal{M}_{e\mu}$ and $\mathcal{M}_{e\tau}$, all other matrix elements of $\mathcal{M}$ break $\overline{L}$ and, therefore, they will all be proportional to the $\overline{L}$-breaking parameter ϵ, which is assumed to be small. The matrix elements $\mathcal{M}_{\mu\mu}$, $\mathcal{M}_{\mu\tau}$, and $\mathcal{M}_{\tau\tau}$ arise at two loops from the diagram in Fig. 2. In order to obtain a nonzero $\mathcal{M}_{ee}$ one must go to three loops and use for instance the diagram in Fig. 3. In that diagram there are two inner charged leptons which may be either μ or τ; therefore, there is a contribution to $\mathcal{M}_{ee}$ proportional to m_τ^2, and that matrix element should not be neglected in spite of it only arising at the three-loop level.

The diagram in Fig. 2 clearly leads to the following relation:

[2] It is not possible to use the approximations $m_e = m_\tau = 0$ because they lead to infrared divergences. This is not a problem, since those divergences are logarithmic and $\mathcal{M}_{e\mu}$ in Eq. (11) also includes a factor $m_e m_\tau$.

[3] Notice, however, that, in the standard model, the top-quark Yukawa coupling is also very close to 1.

[4] Concerns about the breakdown of perturbativity are only justified for Yukawa couplings $\gtrsim 4\pi$, i.e., of order 10 or more.

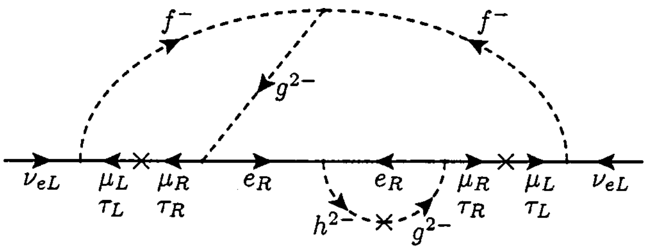

FIG. 3. One of the three-loop Feynman diagrams which generate $\mathcal{M}_{ee}$.

$$\mathcal{M}_{\mu\mu}:\mathcal{M}_{\mu\tau}:\mathcal{M}_{\tau\tau}=f_\mu^2:(f_\mu f_\tau):f_\tau^2=r^2:r:1. \tag{21}$$

One thus obtains that, in the present model,

$$\mathcal{M}=\begin{pmatrix} a & rb & b \\ rb & r^2c & rc \\ b & rc & c \end{pmatrix}, \tag{22}$$

where a, b, and c are complex numbers with mass dimension, while $r=f_\mu/f_\tau$ is a real dimensionless number which should in principle be of order 1. The masses a and c are suppressed relative to b by the soft-breaking parameter ϵ.

The mass matrix in Eq. (22) immediately leads to two predictions of this model: there is one massless neutrino (ν_3), and its component along the ν_e direction, i.e., U_{e3}, vanishes. Indeed, the diagonalizing matrix U reads

$$U=\begin{pmatrix} \cos\psi & -i\sin\psi & 0 \\ e^{i\alpha}\dfrac{r\sin\psi}{\sqrt{1+r^2}} & e^{i\alpha}\dfrac{ir\cos\psi}{\sqrt{1+r^2}} & \dfrac{1}{\sqrt{1+r^2}} \\ e^{i\alpha}\dfrac{\sin\psi}{\sqrt{1+r^2}} & e^{i\alpha}\dfrac{i\cos\psi}{\sqrt{1+r^2}} & -\dfrac{r}{\sqrt{1+r^2}} \end{pmatrix} \times \mathrm{diag}(e^{i\theta_1},e^{i\theta_2},1); \tag{23}$$

cf. Eq. (7). In the matrix of Eq. (23), $\alpha\equiv\arg[ab^*+bc^*(1+r^2)]$ is a physically meaningless phase. The Majorana phases θ_1 and θ_2 are necessary in order to obtain real and positive m_1 and m_2. The sole physically observable phase is $2(\theta_1-\theta_2)$ [8]. The mixing angle ψ is given by

$$\tan 2\psi=-1/\varepsilon, \tag{24}$$

where

$$\varepsilon=\frac{|c|^2(1+r^2)-|a|^2}{2\sqrt{1+r^2}|ab^*+bc^*(1+r^2)|} \tag{25}$$

is a quantity of order ϵ, just as a/b and c/b, and may therefore be assumed to be very small. Thus ψ is close to 45°. The amplitude of the oscillations of ν_e relevant for the solution of the solar-neutrino problem is $4|U_{e1}U_{e2}|^2=(1+\varepsilon^2)^{-1}$, i.e., almost maximal (it deviates from 1 only by a term of order ε^2). Thus the present model favors the ‘‘just so’’ and the LOW solutions of the solar-neutrino puzzle.

The soft-breaking parameter ϵ should be tiny. Indeed, one finds

$$\frac{\Delta m_\odot^2}{\Delta m_{\mathrm{atm}}^2}\approx 2\,\frac{|ab^*+bc^*(1+r^2)|}{|b|^2\sqrt{1+r^2}}\sim\epsilon; \tag{26}$$

if we want the ‘‘just so’’ solution for the solar-neutrino puzzle to apply, then we must accept ϵ to be of order 10^{-7}; if we need the LOW solution, then $\epsilon\sim 10^{-4}$. Such a tiny soft breaking of $\bar{L}$ may eventually be explained by some new physics at a very high energy scale.

From the nonobservation of neutrinoless double beta decay, one derives the bound $|\mathcal{M}_{ee}|\leqslant 0.2$ eV [9]. This is not a problem to the present model. Indeed, as m_3 is predicted to vanish, m_1 and m_2 should both be very close to $\sqrt{\Delta m_{\mathrm{atm}}^2}\approx 0.06$ eV. Thus, in the approximation $\cos^2\psi=\sin^2\psi=1/2$, one has

$$|\mathcal{M}_{ee}|\approx(0.03\ \ \mathrm{eV})|e^{2i(\theta_1-\theta_2)}-1|<0.2\ \ \mathrm{eV}. \tag{27}$$

Moreover, the phase $2(\theta_1-\theta_2)$ is very close to zero—indeed, it vanishes in the limit of $\bar{L}$ conservation.

In conclusion, the model that I have presented in this paper makes the exact predictions $m_3=0$ and $U_{e3}=0$, while it naturally accomodates maximal amplitude ν_e oscillations and a tiny $\Delta m_\odot^2$. Maximal ν_μ-ν_τ mixing follows from the reasonable assumption that two Yukawa couplings are almost equal. Neutrino masses are small because they are radiatively generated at the two-loop level. Indeed, the fact that two neutrino masses are as *large* as 0.06 eV practically forces the new mass scale, at which the extra scalars lie, to be in the 10 TeV range, while deviations from e-μ-τ universality in μ decay and in τ decay should be close at hand. The model requires some physical mechanism for generating a tiny soft breaking of $\bar{L}$.

Note added. A paper by Kitabayashi and Yasuè [10], suggesting the same model that I have presented here, appeared after I submitted this for publication.

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