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Abstract

In the light of renewed evidence for the existence of a 17 keV neutrino, we study the possible mass patterns for the charged and the neutral leptons, in the context of a generalized "seesaw"-type of model, which implements a horizontal $U(1)_A$ Peccei-Quinn symmetry. Under some general assumptions concerning the structure of the mass matrix we find that the mass hierarchy between the first two generations of charged leptons and the third one is explained in terms of the natural scales of the model. At the same time, with the additional assumption of the proportionality of Majorana- and Dirac-type couplings, the spectrum of the neutral leptons contains two very light Majorana neutrinos, such as required by the MSW interpretation of the solar neutrino deficit, and the 17 keV "Simpson" neutrino. A cosmologically consistent decay mode of this neutrino is into a ν_e and the axion.

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1 Introduction

The recent β-decay experiments of Simpson and Hime [1] have revived the possibility of a 17 keV neutrino state with a 10% admixture in the amplitude of ν_e:

$$\nu_e = \cos\theta\,\nu_1 + \sin\theta\,\nu_S\,, \tag{1}$$

where ν_S is the 17 keV component with sin$\theta \sim 0.1$ and ν_1 is one or more light components with a mass less than ten eV. The particularity of the so-called Simpson neutrino has stimulated a large activity in the construction of models for neutrino masses [2-11]. As has been pointed out in the literature, ν_S must be either a Dirac or a pseudo-Dirac state, otherwise its contribution to the "effective" mass of the electron neutrino in neutrinoless double beta decays would be $\sim m_S \sin^2\theta \sim 170$ eV, in contradiction with the experimental upper bound of a few eV [12]. From neutrino oscillation experiments [13] one also knows that ν_S cannot be the dominant component of ν_μ. Moreover, choosing the resonant oscillation mechanism of Mikheyev-Smirnov-Wolfenstein (MSW) [14] as a good solution for the solar neutrino problem [15], one should expect that, in general, at least two neutrinos, say ν_1, the dominant component of ν_e, and ν_2, are lighter than $\sim 10^{-2}$ eV. Therefore, ν_2 cannot be the Simpson (tau) neutrino and may be identified, say, with the dominant component of ν_μ. Hence, the most natural scenario for implementing both, the Simpson neutrino and the MSW mechanism, is one with very light electron and muon neutrinos and a Dirac tau neutrino of 17 keV mass.

For Majorana neutrinos, the seesaw mechanism [16] provides an explanation for the smallness of their mass compared to the masses of the charged fermions by introducing into the Standard Model extra heavy iso-singlets, *e.g.*, the right-handed neutrinos. The suppression of the neutrino mass can be easily understood from the perturbative diagonalization of the neutrino mass matrix and works obviously only if the mass of the heavy state is non-zero. Correspondingly, in the case of more than one generation, it was noticed that part of the spectrum can be protected from getting seesaw suppressed masses by assuming a singular mass matrix for the heavy singlet states [17]. This idea has been recently used by some authors for implementing the rather heavy Simpson neutrino in their models [2,3,9].

Similarly, in a two-generation toy version of the Universal Seesaw Model (USM) [18-22], which is essentially the generalization of the seesaw idea to all fermions, the singularity of the mass matrix for the extra heavy singlets (for details, see Section 2) was used to protect the second-generation fermion from having a seesaw-suppressed mass [18]. This provided a natural explanation of the mass hierarchy between the two generations of charged fermions. However, it was noticed that this protection mechanism did not work for the two neutrinos, for which the mass scale turned out to be the same. This different behaviour between the two sectors is due to the presence of Majorana masses for the neutral leptons.

In this paper we study the lepton sector of the USM for the realistic case of three generations assuming the singularity of the mass matrix for the heavy singlets. In section 2 we review the main features of the USM. In sections 3 and 4 we discuss the mass spectrum and the hierarchy pattern of the charged and the neutral leptons in the most general scenario. Then, in section 5, assuming that Majorana- and Dirac-type couplings scale proportionally, we obtain a quite different neutrino spectrum, which contains the Simpson neutrino and the two light states required by the MSW solution of the solar neutrino problem. Finally, in section 6 we derive the lifetime of the 17 keV neutrino and check the consistency of this spectrum with cosmology.

2 The Universal Seesaw Model

As indicated in the Introduction, one of the attempts to understand the mass-hierarchy between the three generations of quarks and leptons is the Universal Seesaw Model (USM) [18-22], first proposed by Davidson and Wali [19], which is essentially an extension of the well known neutrino-seesaw mechanism [16] to all fermions. It is naturally implemented in the context of left-right symmetric theories based on the $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ gauge group, by introducing, for each left-and right-handed fermion $f_{L(R)}$, a new heavy $SU(2)_L \otimes SU(2)_R$ singlet $F_{L(R)}$, which mimics the role of the right-handed neutrino of the standard (neutrino) seesaw scenario. In other words, in the USM each ordinary fermion (f) has a singlet heavy partner (F), with the same electric and color charges.

The Higgs sector is kept minimal; in fact, at the one-generation level there are only one left- and one right-handed Higgs doublets [1], $\phi_L(2,1)_1$ and $\phi_R(1,2)_1$ with vacuum expectation values v_L and v_R, which couple the ordinary fermions to the heavy singlets. The heavy singlets themselves get their masses from an extra singlet Higgs field, $\sigma(1,1)_0$, whose vacuum expectation value, χ, can also induce the spontaneous breaking of the left-right symmetry à la Chang-Mohapatra-Parida [23]. In the multi-generation case, χ can be identified with the Peccei-Quinn scale [24]. Moreover, as was shown in ref.[20], the Higgs potential allows, within a certain range of parameters, an absolute minimum which corresponds to a hierarchy of the VEV's: $v_L \ll v_R \ll \chi$. The diagrams which generate fermion masses at tree level

[1]We adopt the standard notation, where $(i,j)_k$ label respectively $SU(2)_L$, $SU(2)_R$ representations, and k is the quantum number of $U(1)_{B-L}$.

in the context of the USM are shown in Fig.1 of ref.[21]. Notice that the standard source for the generation of Dirac masses in usual left-right symmetric theories [25], namely the bi-doublet Higgs field, $\phi(2,2)_0$, as well as the standard sources of Majorana masses for the neutrinos, the left- and the right-handed triplet Higgs fields $\Delta_L(3,1)_{-2}$ and $\Delta_R(1,3)_{-2}$, are absent at the lowest order. All these fields, behaving as bilinears of the Higgs doublet fields present in the model, are however recovered as effective fields at low energy.

As a consequence of our choice for the Higgs sector, the tree-level mass matrix is of the seesaw-type, with Dirac-type mass terms in all sectors plus Majorana-type mass terms in the neutral lepton sector. It is then easy to check, that for one generation, while the masses of the ordinary up- and down-quark and the charged lepton are $m_f \sim v_L v_R/\chi$ and therefore are suppressed with respect to the weak scale ($v_L \simeq M_{W_L}$) through the smallness of the v_R/χ ratio, the standard neutrino mass $m_{\nu L} \sim v_L^2/\chi$ is even more suppressed. In particular one recovers the usual relation of the ordinary seesaw model $m_{\nu L}\, m_{\nu R} \sim m_e^2$, which explains the suppression of the neutrino mass with respect to the electron mass, through the largeness of the mass of the right-handed neutrino, $m_{\nu R} \sim v_R^2/\chi$.

When generalizing this model to more than one generation, care has to be taken of those fermions, like the top quark, whose mass is of the order of the electroweak scale. In order to protect them from getting a seesaw-suppressed mass, in ref.[18] an axial (global) $U(1)_A$ horizontal symmetry - broken at a very high scale - was implemented, having the nice feature that it could be identified with the Peccei-Quinn symmetry, which gives a solution to the strong CP problem [24]. The role played by this new symmetry is to distinguish the different families, by assigning to each of them a different quantum number: x_i for the ordinary fermions and y_i for the singlets of the *i-th* generation. Furthermore, since now a Higgs field and its charge-conjugated field, having opposite $U(1)_A$ charges, are distinguished, the Higgs sector is doubled, and one has two left- and two right-handed doublet Higgs fields, whose vacuum expectation values are assumed to be equal in the left and the right sectors [26] respectively [2], and their $U(1)_A$ quantum numbers can be fixed, without any loss of generality, to be ± 1 (see Tab.1 of ref.[20]). Then, the $U(1)_A$ charge of the Higgs singlet field σ is, for a suitable choice of the $U(1)_A$-invariant Higgs potential [21], equal to ± 2. In order to determine the non-zero entries of the mass matrix it is useful to introduce the matrix Q of the $U(1)_A$ quantum numbers,

[2] This assumption can be justified by the study of minima of a general (left-right symmetric) Higgs potential involving four Higgs doublets (as in our model), with additional discrete symmetries. The minima occur only when $v_{L1} = v_{L2}$, $v_{R1} = v_{R2}$ [26]. In this paper we adopt the following definitions: $v_{L1} = v_{L2} = v_L/\sqrt{2} = (4\sqrt{2}G_F)^{-1/2} \simeq 123$ GeV, and $v_{R1} = v_{R2} = v_R/\sqrt{2}$.

which in the basis $(f_i,\ F_j)$, has the general form:

$$Q = \begin{pmatrix} X & L \\ R & Y \end{pmatrix}, \tag{2}$$

where $X_{ij} = x_i + x_j$, $L_{ij} = x_i + y_j$, $R_{ij} = y_i + x_j$, $Y_{ij} = y_i + y_j$. Then, the non-vanishing mass terms in the mixed $f - F$ sectors, M_L and M_R, will correspond to the entries $L_{ij} = \pm 1$ ($R = L^T$). This is the only case where we have an allowed, *i.e.*, U(1)-invariant, fermion-Higgs coupling. As already stated above, the absence of a Higgs field transforming as a bi-doublet under the left-right symmetric group, tells us that there cannot be Dirac masses involving the ordinary fermions only. On the other hand, non-vanishing entries of the mass matrix for the singlet fermion fields will correspond to $Y_{ij} = \pm 2$, since the $U(1)_A$ charge of σ is ± 2.

From this discussion it follows that the mass matrix for the charged fermions is of the seesaw-type:

$$\mathcal{M} = \begin{pmatrix} 0 & M_L \\ M_R & M_\sigma \end{pmatrix}, \tag{3}$$

where M_σ is the mass matrix for the heavy singlet fermions and $M_{L(R)}$ are the corresponding matrices coupling the ordinary charged fermions with the heavy singlets. With the conditions specified so far, there are still several possible forms for the sub-matrices M_L and M_R. In what follows, we restrict ourselves to the case where both M_L and M_R are of the Fritzsch-type [27], since this *Ansatz* has proved successful in the study of the Kobayashi-Maskawa mixing matrix in the quark sector [20,21]. Then, for three generations:

$$M_L = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & Y_{12}v_L & 0 \\ Y_{21}v_L & 0 & Y_{23}v_L \\ 0 & Y_{32}v_L & Y_{33}v_L \end{pmatrix}, \tag{4}$$

and $M_R = M_L(v_L \leftrightarrow v_R)$. In equ.(4) we have taken into account the fact that in two-doublet models with both Higgs fields developing the same VEV, one has $v_1 = v_2 = v/\sqrt{2}$, v being the "effective" vacuum expectation value responsible for the mass of the gauge bosons, $M_W = gv/\sqrt{2}$, where $v = ({v_1}^2 + {v_2}^2)^{1/2}$. The Yukawa couplings, Y_{ij}, are assumed to be real, left-right symmetric, and symmetric in the generation space, *i.e.*, $Y_{ij} = Y_{ji}$. The matrix of the $U(1)_A$ quantum numbers, Q, is then uniquely determined in terms of a single integer parameter x, which can be chosen such that the determinant of M_σ is either zero (singular case) or different from zero (non-singular case).

The non-singular three-generation model for quarks was studied in refs.[20,21], where the mixing angles were determined in terms of the quark masses. In particular, it was shown [21] that all flavour changing neutral current processes (FCNC) mediated by the two neutral gauge vector bosons, ${Z_1}^0$ and ${Z_2}^0$, being proportional to $v_L^4/(v_R\chi)^2$, are highly suppressed. However, as was first suggested in ref.[18], the so-called singular models ($det\ M_\sigma = 0$) are more suitable for understanding the origin of the generation mass hierarchy. We shall therefore restrict our discussion of the three-generation lepton sector to this case only.

3 The charged lepton sector

As we have shown in the previous section, in the USM model the charged lepton mass matrix is of the seesaw-type, equ.(3). For three generations the submatrices M_L and M_R, which are of Fritzsch-type, equ.(4), can be written in the form:

$$M_L = \begin{pmatrix} 0 & L & 0 \\ L & 0 & aL \\ 0 & aL & bL \end{pmatrix}, \tag{5}$$

and

$$M_R = M_L(L \leftrightarrow R),$$

respectively. The mass matrix for the heavy singlets, M_σ, which must be singular in order to give rise to a mass hierarchy between the generations, depends on a single integer parameter x, which characterizes the Peccei-Quinn fermion charges [20,21]. In view of the interesting results that we obtain in the neutral lepton sector, section 5, we choose the model corresponding to $x = 5$, which gives the following form for M_σ:

$$M_\sigma = \begin{pmatrix} 0 & 0 & K \\ 0 & 0 & 0 \\ K & 0 & 0 \end{pmatrix}. \tag{6}$$

In eqs.(5) and (6) we have set: $K = Y_{13}^E\chi$, $L = Y_{12}^l v_L/\sqrt{2}$, $R = Y_{12}^l v_R/\sqrt{2}$, $a = Y_{23}^l/Y_{12}^l$ and $b = Y_{33}^l/Y_{12}^l$, the $Y_{ij}^{l,E}$'s being the left-right symmetric Yukawa-type coupling constants in the charged lepton sector.

We now evaluate the mass eigenvalues, to see whether the charged lepton mass hierarchy of the two-generation USM [18], is also present in the realistic three-generation model. For simplicity we assume all mass matrices to be real, disregarding the possibility of spontaneous CP-violation. Using the hierarchy of the VEV's ($\chi \gg v_R \gg v_L$), the secular equation of $\mathcal{M}_l\,\mathcal{M}_l^\dagger$ [3], in leading order, is given by:

$$\lambda^6 \ - \ 2\,\lambda^5 K^2 + \lambda^4 K^4 - \lambda^3 K^4 R^2 (1+a^2) + \lambda^2 K^4 R^2 L^2 (1+a^2)^2$$

[3]Notice that, since $\mathcal{M}_l$ is not a hermitian matrix, one must consider the properties of $\mathcal{M}_l\,\mathcal{M}_l^\dagger$, whose eigenvalues are the fermion masses squared.

$$- \; 2\lambda K^2 L^4 R^4 (1+a^2)(2a^2+b^2+2a^4) \; + \; L^6 R^6 b^4 \; = \; 0 \, . \tag{7}$$

The roots of this equation can be found perturbatively using the fact that the coefficients of the different powers of λ scale differently with the VEV's. For example, as far as the heavy sector is concerned, equ.(7) shows that there are only two such heavy states, $\lambda_5 = \lambda_6 \simeq K^2$, consistent also with the fact that the trace of $\mathcal{M}_l\mathcal{M}_l^\dagger$ is equal to $2K^2$. Analogously, for the light eigenvalues ($\lambda \ll K^2$), equ.(7) reduces, up to terms of $\mathcal{O}(R^2/K^2)$, to a fourth order equation with only two non-vanishing solutions:

$$\lambda_3 \simeq (1+a^2)\, L^2 \ , \qquad\qquad \lambda_4 \simeq (1+a^2)\, R^2 \, . \tag{8}$$

This means that the remaining two solutions are much smaller than L^2. They can be found by substituting the approximate values for $\lambda_3, ..., \lambda_6$ in the expression for the determinant of $\mathcal{M}_l\mathcal{M}_l^\dagger$ and its fifth-order invariant: $\Delta_5 \simeq \lambda_6 \cdot ... \cdot \lambda_3(\lambda_2 + \lambda_1)$. We then obtain the two light eigenvalues:

$$\lambda_{1,2} = \left(\frac{LR}{K}\right)^2 \left(2a^2 + c \mp 2a\sqrt{a^2+c}\right) \, , \tag{9}$$

where $c = b^2/(1+a^2)$.

These results show that the mass of the two lightest states is set by the scale $v_L v_R/\chi$, while there is a mass gap of order v_R/χ to the next mass eigenstate, which is proportional to v_L. There are in addition three heavy mass eigenstates, which correspond to the singlet states, one of order v_R and two of order χ. Hence, the "singular" version of the three-generation USM yields a physical spectrum which is the generalization of the one obtained in the simplified two-generation model [18]. In other words, as a result of the singularity of M_σ, one can naturally (*i.e.*, without the need of any hierarchy among the Yukawa coupling constants) generate a mass hierarchy of order v_R/χ between the first two generations and the third one. On the other hand, the hierarchy between the first and the second generation requires certain assumptions on the ratios of the Yukawa couplings. For example, we can reproduce the mass difference between the electron and the muon by assuming $b/a \ll 1$, in which case, due to the partial cancellation of the two terms in equ.(9), we obtain:

$$m_e \simeq \frac{LR}{K}\left(\frac{c}{2a}\right) \, , \qquad m_\mu \simeq \frac{LR}{K}\,(2a) \, . \tag{10}$$

These relations and the one for the tau mass:

$$m_\tau \simeq \sqrt{1+a^2}\, L \, , \tag{11}$$

allow us to fit the charged lepton mass spectrum. If, for simplicity, we choose $a = 1\,$(*i.e.*, $Y^l_{12} = Y^l_{23}$), from equ.(11) we obtain $Y^l_{12} = Y^l_{23} \simeq 10^{-2}$, while from the mass ratio m_μ/m_τ we can determine R/K to be $\simeq 0.04$. Furthermore, from the ratio $b = Y^l_{33}/Y^l_{12} \simeq 2^{3/2}\,(m_e/m_\mu)^{1/2}$ we obtain the value $Y^l_{33} \simeq 2 \cdot 10^{-3}$. Notice that the ratio of the Yukawa couplings scales only as the square root of the lepton masses, so that the Yukawa couplings do not need to share the same level of hierarchy as the masses.

Since, assuming $Y^l_{12} \sim Y^E_{13}$, the mass hierarchy between the second and the third generation, $R/K \simeq 0.04$, reduces essentially to v_R/χ, and because the allowed range of the Peccei-Quinn scale is $10^{10} - 10^{12}$ GeV, (see equ.(36) in the Appendix), also the $SU(2)_R$ symmetry is broken at very high energies. The exact value of χ, and therefore of v_R, will be fixed by studying the neutral lepton sector of the model.

4 The Neutrino Sector (general case)

Because of the presence of Majorana mass terms, the mass matrix for the neutral leptons is a (12×12) symmetric matrix. In the basis:

$$(\nu_{eL},\, \nu_{\mu L},\, \nu_{\tau L},\, \nu^c_{eL},\, \nu^c_{\mu L},\, \nu^c_{\tau L},\, N_{1L},\, N_{2L},\, N_{3L},\, N^c_{1L},\, N^c_{2L},\, N^c_{3L})$$

it takes the form:

$$\mathcal{M}_\nu = \begin{pmatrix} 0 & M_0 \\ M_0^T & M_N \end{pmatrix}, \tag{12}$$

where

$$M_0 = \begin{pmatrix} 0 & dL & 0 & 0 & cL & 0 \\ dL & 0 & L & cL & 0 & bL \\ 0 & L & aL & 0 & bL & eL \\ 0 & cR & 0 & 0 & dR & 0 \\ cR & 0 & bR & dR & 0 & R \\ 0 & bR & eR & 0 & R & aR \end{pmatrix}, \tag{13}$$

and

$$M_N = \begin{pmatrix} 0 & 0 & hK & 0 & 0 & K \\ 0 & 0 & 0 & 0 & 0 & 0 \\ hK & 0 & 0 & 0 & 0 & 0 \\ 0 & 0 & K & 0 & 0 & hK \\ 0 & 0 & 0 & 0 & 0 & 0 \\ K & 0 & 0 & hK & 0 & 0 \end{pmatrix}. \tag{14}$$

M_N is the singular mass matrix for the heavy singlets, corresponding again to the choice $x = 5$. The parameters $a, ..., h$ are ratios of Yukawa-type couplings analogous to the ones defined in the charged lepton sector, but having in general different values.

We proceed now with a semi-analytical evaluation of the neutrino mass eigenvalues exploiting the hierarchical structure of M_ν. We first notice that in the limit $L \to 0$, the rank of the full matrix $\mathcal{M}_\nu$ reduces from twelve to nine, indicating that in this limit three eigenvalues m_1, m_2, $m_3 \to 0$ compared to the scale of R and K. When also $R \to 0$, the rank of the residual matrix is four, leading to four very massive states m_9, m_{10}, m_{11}, $m_{12} \sim \mathcal{O}(K)$. The scale of the various eigenvalues can then be obtained by using some of the invariants of the mass matrix, $\Delta_1 \equiv Tr\, M_\nu$, ..., $\Delta_{12} \equiv Det\, M_\nu$, for which we give only the leading term:

$$\begin{aligned}
\Delta_{12} &\simeq f_{12}\, R^6\, L^6\,, \\
\Delta_{11} &\simeq f_{11}\, K\, R^6\, L^4\,, \\
\Delta_{10} &\simeq f_{10}\, K^2\, R^6\, L^2\,, \\
\Delta_9 &\simeq f_9\, K^3\, R^6\,, \\
\Delta_8 &\simeq f_8\, K^4\, R^4\,,
\end{aligned} \tag{15}$$

where f_i are complicated functions of the Yukawa ratios, having small effects compared with the scales of the symmetry-breaking. Since, in the present type of models, the fermion mass hierarchy is not supposed to emerge from large differences in the Yukawa couplings, we shall ignore them. Using the fact that $m_{1,2,3} \ll m_{4,5,6,7} \ll m_{9,10,11,12}$, we can write the following simplified expressions:

$$\begin{aligned}
\frac{\Delta_{12}}{\Delta_{11}} &\simeq \frac{m_1 m_2 m_3}{m_1 m_2 + m_2 m_3 + m_1 m_3} \sim \frac{L^2}{K}\,, \\
\frac{\Delta_{11}}{\Delta_{10}} &\simeq \frac{m_1 m_2 + m_2 m_3 + m_1 m_3}{m_1 + m_2 + m_3} \sim \frac{L^2}{K}\,, \\
\frac{\Delta_{10}}{\Delta_9} &\simeq m_1 + m_2 + m_3 \sim \frac{L^2}{K}\,, \\
\frac{\Delta_9}{\Delta_8} &\simeq \left(m_4^{-1} + m_5^{-1} + m_6^{-1} + m_7^{-1} + m_8^{-1}\right)^{-1} \sim \frac{R^2}{K}\,.
\end{aligned} \tag{16}$$

From the first three relations one finds that the three lightest neutrino states have the same mass scale:

$$m_1 \sim m_2 \sim m_3 \sim \mathcal{O}(\frac{L^2}{K})\,, \tag{17}$$

while from the ratio Δ_9/Δ_8 and the information contained in the determinant (Δ_{12})we find that one eigenvalue, m_4, is of order R^2/K, four eigenvalues, m_5, m_6, m_7, m_8 are of order R, and the remaining four are of order K.

Compared with the mass hierarchy in the charged lepton sector, the absence of a hierarchical structure for the masses of the three standard neutrinos is an interesting and surprising result of the singular USM scenario, entirely due to the presence of Majorana masses for the neutral leptons. We can fix our scales by considering some of the favoured theoretical scenarios for neutrino masses. One possibility would be to account for the MSW solution of the solar neutrino problem, by choosing the mass scale χ in such a way that [14]:

$$\Delta m_{12}^2 = m_2^2 - m_1^2 \sim \left(\frac{L^2}{K}\right)^2 \sim (10^{-3}\,\mathrm{eV})^2 . \tag{18}$$

This would of course imply that all three neutrinos have a mass of order 10^{-3} eV, so that they cannot be a good candidate for the dark matter of the Universe [30]. The other possibility would be to ignore the MSW requirements, assuming a different solution for the solar neutrino problem [15], and choose the mass scale χ such that L^2/K is of the order of $\sim$ 50 eV, corresponding to a cosmologically interesting scenario. In such a case one could then use the freedom in the choice of the Yukawa couplings to satisfy the bounds on the "effective" electron neutrino mass derived from the data on neutrinoless double beta decay [12].

Since in the neutrino sector we have no guidance of how to choose the Yukawa couplings, a natural assumption could be to set them of the same order as the ones we have obtained in the charged lepton sector (*i.e.*, $Y_i^\nu \sim 5 \cdot 10^{-3}$), and then evaluate the Peccei-Quinn scale χ. For the model that naturally satisfies the MSW constraints, χ turns out to be of the order of $\sim 10^{14}$ GeV, which is in contradiction with the present bounds, equ.(36); on the other hand, the cosmologically attractive model gives a compatible value of $\chi \sim 10^{10}$ GeV. One should however stress the fact that this particular choice of the Yukawa couplings is by no means necessary, in the sense that, in general, the Yukawa couplings characterizing the two lepton sectors can be completely independent.

5 A 17 keV Dirac neutrino in the USM

In this section we discuss the neutrino sector in the case where the ratio of the Majorana- to the Dirac-type Yukawa couplings is not arbitrary (as in the previous section), but is constant in generation space. This assumption leads to a more interesting neutrino spectrum consisting of the Simpson neutrino and two very light neutrinos, with a mass consistent with the MSW interpretation of the solar neutrino deficit.

Let Y_{ij} and $\tilde{Y}_{ij}$ be the Dirac- and the Majorana-type Yukawa couplings, defined by equ.(13) with: $L = \tilde{Y}_{23}v_L/\sqrt{2}$, $R = \tilde{Y}_{23}v_R/\sqrt{2}$, $K = Y_{13}^N\chi$, and $a = \tilde{Y}_{33}/\tilde{Y}_{23}$, $b = Y_{23}/\tilde{Y}_{23}$, $c = Y_{12}/\tilde{Y}_{23}$, $d = \tilde{Y}_{12}/\tilde{Y}_{23}$, $e = Y_{33}/\tilde{Y}_{23}$. The assumption that $Y_{ij}/\tilde{Y}_{ij}$ is constant, so that $e/a = c/d = b$, changes the leading behaviour of the following invariants of the mass matrix $\mathcal{M}_\nu$ with respect to equ.(15):

$$
\begin{aligned}
\Delta_9 &\simeq f_9\, K^3\, L^2\, R^4\,, \\
\Delta_8 &\simeq f_8\, K^2\, R^6\,,
\end{aligned}
\tag{19}
$$

where the coefficients f_i's are generally different from those in equ.(15). Notice that now Δ_9 goes to zero with L, so that the rank of the matrix reduces to eight, instead of nine, implying that four eigenmasses scale with L. As before, when also R goes to zero, the rank further reduces to four, suggesting that four eigenvalues scale with R and four with K. Using these considerations and building the ratios:

$$
\begin{aligned}
\frac{\Delta_{12}}{\Delta_8} &\simeq m_1\, m_2\, m_3\, m_4 \sim \frac{L^6}{K^2}\,, \\
\frac{\Delta_9}{\Delta_8} &\simeq m_1 + m_2 + m_3 + m_4 \sim K\frac{L^2}{R^2}\,, \\
\frac{\Delta_8}{\Delta_4} &\simeq m_5\, m_6\, m_7\, m_8 \sim \frac{R^6}{K^2}\,,
\end{aligned}
\tag{20}
$$

where Δ_4 scales obviously as K^4, it is possible to fully determine the structure of the mass spectrum. From the first equation one sees that two eigenvalues, m_1 and m_2, scale with L^2/K, and two, m_3 and m_4, with L; on the other hand, the second equation suggests that m_3 and m_4 must cancel to order $KL^2/R^2 \ll L$. This implies that the corresponding eigenstates, ν_3 and ν_4, to a good approximation, behave as a single effective Dirac state. Analogously, from the third equation and the vanishing of the trace of $\mathcal{M}_\nu$, one obtains two eigenvalues $\sim R^2/K$, and two $\sim R$, which again combine into an effective Dirac state.

If we now identify the Dirac neutrino in our spectrum, whose mass is proportional to L, with the 17 keV neutrino, we see that the Yukawa couplings, assuming them all to be comparable to each other, are of the order of $\simeq 10^{-7}$. Furthermore, since the two lightest states with mass $\sim L^2/K$ turn out to be essentially ν_e and ν_μ, in order to allow for the MSW mechanism, their mass must be in the range $m_{\nu_e} \sim m_{\nu_\mu} \simeq 10^{-3} - 10^{-4}$ eV, from which we obtain $\chi \simeq 10^9 - 10^{10}$ GeV. In the following we fix $\chi = 10^{10}$ GeV, corresponding to the lower bound of the allowed Peccei-Quinn scale. Using the ratio R/K from the charged lepton sector we then get $v_R = 6 \cdot 10^8$ GeV. Notice that if we would have chosen the electron and the muon neutrino masses such as to account for the dark matter problem , χ would have been too small to be identifiable with the Peccei-Quinn scale.

Our numerical analysis, confirms the previous results on the mass spectrum and suggests that the eigenvectors can be parametrized as shown in Tab.1 [4]. In particular, the two light Majorana states, ν_1 and ν_2, with a mass $\sim 10^{-4}$ eV, lie mainly along the ν_e and the ν_μ direction, respectively [5]. Their right-handed counterparts, ν_5 and ν_6, also Majorana states, have masses in the range 1 - 10 GeV. The remaining part of the spectrum consists of four pairs of mass-degenerate states, which combine into a set of four Dirac neutrinos. As can be seen from Tab.1, the 17 keV Simpson neutrino is the 'plus' linear combination of ν_3 and ν_4, which is essentially ν_{τ_L} with a small admixture of ν_{e_L} of order 10 %, while the 'minus' combination consists of the two singlets N_{2L} and N_{2R}. Correspondingly, $\nu_7 + \nu_8$ gives essentially the right-handed tau neutrino, with a mass $\simeq 60$ GeV. The two heavy Dirac states with masses ~ 1 TeV, are combinations of the remaining four singlet fermions N_{iL} and N_{iR}, $i = 1, 3$. Their mixing with the left- and the right-handed states is of the order $L/K \sim 10^{-8}$ and $R/K \sim 10^{-2}$, respectively, so they essentially decouple from the rest of the spectrum. Due to the smallness of the mixing between the heavy right-handed Majorana states and ν_{eL}, the effective mass in neutrinoless double-beta decay is still consistent with the present experimental limit [12].

By expressing the interaction eigenstates in terms of the mass eigenstates, one can write the three standard neutrinos in the following way:

$$\begin{aligned} \nu_{e_L} &\simeq \cos\theta\nu_1 + \sin\theta\nu_S \,, \\ \nu_{\mu_L} &\simeq \nu_2 \,, \\ \nu_{\tau_L} &\simeq -\sin\theta\nu_1 + \cos\theta\nu_S \,, \end{aligned} \tag{21}$$

where

$$\nu_S = \frac{\nu_3 + \nu_4}{\sqrt{2}} = \sin\theta\nu_{e_L} + \cos\theta\nu_{\tau_L} \tag{22}$$

denotes the "Simpson" neutrino. For our next discussion of the neutrino decays it is also convenient to define:

$$\tilde{N}_2 = \frac{N_{2L} - N_{2R}}{\sqrt{2}} = \frac{\nu_3 - \nu_4}{\sqrt{2}} \,, \tag{23}$$

and

$$N_2 = \frac{N_{2L} + N_{2R}}{\sqrt{2}} = \frac{\nu_7 - \nu_8}{\sqrt{2}} \,. \tag{24}$$

[4]We recall that in our model the fermion mass hierarchy does not arise from a hierarchy in the Yukawa couplings. We have therefore chosen the Yukawa coupling ratios close to one.

[5]The mixing angle between ν_{μ_L} and ν_{e_L} (or ν_{τ_L}) is always smaller than 0.01, while the mixing with the corresponding right-handed states is less than $L/R \sim 10^{-6}$.

Notice that $\tilde{N}_2$ is mass degenerate with the Simpson neutrino, with which it forms an effective Dirac state. The same applies for N_2 and ν_{τ_R}.

We would like now to comment on two points, which have been crucial for obtaining the type of spectrum discussed in this section. The first one, is the requirement that Majorana- and Dirac-type couplings, in general independent, should be proportional. This assumption, which does not find a justification in our model, might have a deeper meaning in the context of some more fundamental theory. The second comment concerns our choice of a particular type of singular mass matrix for the heavy extra singlets, namely, the one which corresponds to x=5, fixing therefore all the Peccei-Quinn charges. Interestingly, we have noticed that the type of spectrum discussed above is a characteristic feature of the x=5 case only, and does not arise in the other (singular) models with different values of x. On the other hand, we have found that the general behaviour shown in section 4 is essentially independent on the parameter x, within, of course, the singular type of models. These observations seem to suggest that, if the Simpson neutrino will turn out to be real, the Peccei-Quinn charges of our model could be fixed.

6 Decay modes and neutrino lifetimes

According to the standard model of cosmology [31], any (standard) neutrino with a mass in the range[6]: $97 \cdot (\frac{2}{g_\nu})\,\mathrm{eV} \leq m_\nu \leq 2(5)\,\mathrm{GeV}$, must decay fast enough such that its decay products do not produce a too large energy density for the universe. In particular, for a weak interacting 17 keV neutrino, the cosmological density bound results in the following constraint on its lifetime:

$$\tau_{\nu_S} \leq \left(\frac{2}{g_\nu}\right)^2 1.5 \cdot 10^{12}\,\mathrm{sec}\,. \tag{25}$$

The possibility that, in some model, the 17 keV neutrino could decay into three light neutrinos at a rate consistent with equ.(25) has been considered in refs.[4,5]. As it is known, the radiative decay into a photon and a lighter neutrino has already been excluded as the dominant decay channel from the SN87A constraints and the lack of distortion of the cosmic microwave background radiation [4]. Another decay mode, which has been proved successful in most models, is into a Majoron and a light neutrino [2,3,5,9]. Since in the present model, $U(1)_{B-L}$ is a local symmetry, its breaking cannot be associated with the Majoron [32]. However, our model

[6]Here g_ν is the number of helicity states, and the five in the parenthesis refers to a Majorana neutrino.

has a global (horizontal) $U(1)_A$ symmetry - the Peccei-Quinn symmetry - which is spontaneously broken by the complex scalar field σ at a large scale, χ. The associated pseudo-Goldstone boson is then the so-called "invisible axion", $A \simeq Im\,\sigma$, [29]. In the singular USM considered in this paper, due to the structure of the mass matrix M_N, equ.(14), the axion couples at tree level only to the heavy singlets N_1 and N_3 (Fig.3). An effective diagonal coupling of the axion to the ordinary charged fermions of the standard type: $\frac{m_f}{\chi}\bar{f}\gamma_5 f A$ [33] is generated by the two diagrams in Fig.1. Assuming that the mass of the two pseudoscalar Higgs fields $\tilde{H}_1^0$ and $\tilde{H}_2^0$ is of the order of v_R [7], and using for the coupling of the cubic interaction $\phi_L^0\phi_L^0 A$ the scale obtained in the Appendix: $\gamma \simeq v_R^2/\chi$, these diagrams give:

$$Y^2\,\frac{v_L v_R \gamma}{M_F M_{\tilde{H}}^2}\,\bar{f}\gamma_5 f\,A \simeq \frac{m_f}{\chi}\,\bar{f}\gamma_5 f\,A\,. \tag{26}$$

We recall that according to the results given in Section 3, for the first two generations the mass of the heavy singlets is $M_F \simeq Y\chi$ and the mass of the ordinary fermions is $m_f \simeq Y v_L v_R/\chi$, while for the third generation $M_F \simeq Y\,v_R$ and $m_f \simeq Y v_L$. In addition to the diagrams of Fig.1, however, there are additional contributions to the coupling of the axion to the second generation from the type of diagrams shown in Fig.3, which give an overall diagonal coupling to the charged fermions of the more general form: $(m_f/\chi)\,\bar{f}_2(\alpha+\beta\gamma_5)f_2\,A$, where α and β are constants.

The decay of the left-handed tau neutrino into a left-handed electron neutrino and the axion proceeds through the graph of Fig.2. Basically, what happens is that ν_{τ_L}, which, up to a small mixing with ν_{e_L}, is a linear combination of the mass degenerate states ν_3 and ν_4, decays through its orthogonal mass-degenerate eigenstate $\tilde{N}_2$. The axion couples to $\tilde{N}_2$ and ν_{eL} via its mixing with the pseudoscalar Higgs $\tilde{H}_1^0$, as discussed in the Appendix. The resulting lifetime of the Simpson neutrino, for $\chi \simeq 10^{10}$ GeV, is:

$$\tau(\nu_{\tau L} \to \nu_{eL} A) = \frac{32\pi}{m_{\nu_\tau}}\left(\frac{\sqrt{2}M_{\tilde{H}}^2}{Y v_L \gamma}\right)^2 \simeq \frac{32\pi}{m_{\nu_\tau}}\left(\frac{\chi}{m_{\nu_\tau}}\right)^2 \simeq 1.3\cdot 10^{12}\,\text{sec}., \tag{27}$$

where we have used $M_{\tilde{H}} = v_R$, $\gamma = v_R^2/\chi$, and set $m_{\nu_\tau} = Y v_L/\sqrt{2}$. This value of the ν_τ lifetime is still compatible with the cosmological limit, especially if one takes into account the uncertainties in the knowledge of the parameters used above. Moreover, we notice that with such a lifetime, the ν_τ's from the SN87A could not have produced any delayed ν_e signal in the Kamiokande detector [10].

[7] These pseudoscalars are the left-right symmetric generalization of the single pseudoscalar present in the usual two-Higgs doublet models. A preliminary study of the Higgs sector, which will be discussed elsewhere, suggests that our choice of the pseudoscalar masses is consistent with the Higgs potential and the hierarchy of the VEV's.

For completeness, we give also the lifetime of the right-handed states, though we think that they might be allowed to be stable by cosmology, since they decoupled at a very high temperature, $T \sim 1$ TeV. The decay of the right-handed tau neutrino proceeds through a diagram similar to Fig.2, where $\tilde{N}_2$ is replaced by N_2, while ν_{eR} decays through the diagram shown in Fig.4. The corresponding lifetimes are:

$$\tau(\nu_{\tau R} \to \nu_{eL} A) \simeq \tau(\nu_{\tau L} \to \nu_{eL} A) \left(\frac{m_{\nu_{\tau L}}}{m_{\nu_{\tau R}}}\right) \simeq 3.8 \cdot 10^5 \,\text{sec}\,, \tag{28}$$

$$\tau(\nu_{eR} \to \nu_{eL} A) \simeq \tau(\nu_{\tau L} \to \nu_{eL} A) \left(\frac{m_{\nu_{\tau L}}}{m_{\nu_{eR}}}\right) \simeq 9.2 \cdot 10^6 \,\text{sec}\,, \tag{29}$$

where we have used $m_{\nu_{\tau R}} = Y v_R/\sqrt{2} = 60$ GeV and $m_{\nu_{eR}} = Y v_R^2/2\chi = 2.5$ GeV . The right-handed muon neutrino gets equal contributions from the type of diagrams shown in Figs.3 and 4, resulting in a lifetime:

$$\tau(\nu_{\mu R} \to \nu_{eL} A) \simeq \frac{\pi}{m_{\nu_{\mu R}}} \left(\frac{\chi}{v_R}\right)^2 \left(\frac{\chi}{m_{\nu_{\tau R}}}\right)^2 \simeq 7.7 \cdot 10^7 \,\text{sec}\,, \tag{30}$$

where we have used the same value for the mass of the electron and the muon right-handed neutrinos.

Finally, we would like to point out that, since the dominant decay mode of the neutrinos involves the pseudoscalar Higgs fields and their mixing with the axion, a detailed study of the Higgs spectrum and its possible phenomenological implications for future experiments would be interesting and instructive.

7 Conclusions

In this paper we have studied the lepton masses and their hierarchy, in the context of a realistic three-generation Universal Seesaw scenario. In the charged sector, we have found that the observed hierarchy can be understood in terms of the symmetry breaking scales of the model. On the other hand, the neutral lepton sector, at least in the most general case, does not show any mass hierarchy between the three standard neutrinos. We have however noticed the existence of an interesting model, which gives a hierarchical spectrum, in the case where an additional assumption is made; namely, that the Majorana- and Dirac-type coupling constants, generally assumed to be independent to each other, are proportional. This model is particularly interesting, since its spectrum contains an intermediate-mass Dirac tau neutrino and two very light Majorana states, which are essentially the electron and the muon neutrinos. Identifying the 17 keV neutrino with our Dirac state, and make use of the bounds on the Peccei-Quinn scale, we find that the mass of the two light states turns out to be in the right range, required by the MSW solution of the solar neutrino problem. A cosmologically consistent lifetime for the Simpson neutrino is obtained through its decay into an axion and the electron neutrino.

8 Appendix

The most general structure for a left-right symmetric Higgs potential consisting of one left and one right doublet field $\phi_{L,R}$, and a singlet field σ which under $L \leftrightarrow R$ transforms as $\sigma \leftrightarrow -\sigma$, is [20]:

$$\begin{aligned} V &= \Lambda \left[(\phi_L^\dagger \phi_L)^2 + (\phi_R^\dagger \phi_R)^2 \right] - 2(M^2 - \delta\sigma^2)(\phi_L^\dagger \phi_L + \phi_L^\dagger \phi_L) \\ &+ 2\Gamma (\phi_L^\dagger \phi_L)(\phi_R^\dagger \phi_R) + \lambda\sigma^4 - 2m^2\sigma^2 + 2\gamma\sigma (\phi_L^\dagger \phi_L - \phi_R^\dagger \phi_R), \end{aligned} \tag{31}$$

where the parameters λ, Λ and $\Lambda + \Gamma$ are positive. The spontaneous breaking of the left-right symmetry may be induced à la Chang-Mohapatra-Parida [23] when the sigma field acquires a non-vanishing vacuum expectation value χ. The minimization of this potential leads to the following relations for the three vacuum expectation values v_L, v_R and χ (all chosen to be real) [20]:

$$\begin{aligned} v_L^2 &= \hat{M}^2 - \hat{\gamma}\chi, \\ v_R^2 &= \hat{M}^2 + \hat{\gamma}\chi, \\ \chi^2 &= \hat{m}^2 + \alpha\hat{\gamma}^2, \end{aligned} \tag{32}$$

where

$$\begin{aligned} \hat{M} &= \frac{M}{\sqrt{\Lambda + \Gamma}}, \qquad \hat{\gamma} = \frac{\gamma}{\sqrt{\Lambda - \Gamma}}, \\ \hat{m} &= \frac{m}{\sqrt{\lambda}}, \qquad \alpha = \frac{\Lambda - \Gamma}{\lambda}. \end{aligned}$$

In order to obtain the desired hierarchy $v_L \ll v_R$, the parameter $\hat{\gamma}$ must be approximately: $\hat{\gamma} \simeq \hat{M}^2/\chi$, in which case $v_R^2 \simeq 2\hat{M}^2$. Assuming, for simplicity, that the parameter $\Lambda - \Gamma$ is of order one, the scale of the cubic-Higgs coupling is:

$$\gamma \simeq \frac{v_R^2}{\chi}. \tag{33}$$

As was shown in Ref.[20], the absolute minimum of the potential corresponds to the desired hierarchy $v_L \ll v_R \ll \chi$ for a wide range of the parameters.

When one introduces the extra two Higgs doublets with opposite Peccei-Quinn numbers to ϕ_L and ϕ_R, needed for the multigeneration case, the physical Higgs spectrum resembles the one of a left-right symmetric two-doublet model, apart from a small mixing between the left and the right sector, and the mixing generated by the coupling of the σ field with the doublets, which will be discussed next.

In order to identify σ with the field which breaks the Peccei-Quinn symmetry à la Dine-Fishler-Srednicki (DFS) [29], the cubic interaction term in the Higgs potential must be chosen to be of the following form [21]:

$$2\gamma\sigma\left(\phi_{L2}^{\dagger}\phi_{L1} - \phi_{R1}^{\dagger}\phi_{R2}\right) + \text{h.c.}. \tag{34}$$

After spontaneous symmetry breaking, the Higgs fields can be written as:

$$\phi_{L(R)i} = \begin{pmatrix} \phi_{L(R)i}^{+} \\ \phi_{L(R)i}^{0} + v_{L(R)i} \end{pmatrix}, \qquad \sigma = \frac{(Re\,\sigma + iA)}{\sqrt{2}} + \chi, \tag{35}$$

where $i = 1, 2$, and we have used the fact that approximately, up to terms of order $\sim v_R/\chi$, the axion field, which is the Goldstone boson associated with the spontaneous breaking of the Peccei-Quinn symmetry, is just the imaginary part of the σ field. Of course, our axion has to be an "invisible" one, which means that the breaking scale $<\sigma>= \chi \equiv \sqrt{2} f_a$ lies in the range allowed by astrophysics and cosmology [8][28]:

$$10^{10}\,\text{GeV} \leq \chi \leq 1.4 \cdot 10^{12}\,\text{GeV}. \tag{36}$$

As one can easily check, by substituting equ.(35) in equ.(34), there is a mixing term between the axion and the pseudoscalar Higgs fields $\tilde{H}_1^0$ and $\tilde{H}_2^0$:

$$\mathcal{L}_{A\tilde{H}} = \sqrt{2}\,\gamma\, A\,(v_L \tilde{H}_1^0 + v_R \tilde{H}_2^0), \tag{37}$$

where $2\tilde{H}_1^0 = Im\phi_{L2}^0 - Im\phi_{L1}^0$, and $\tilde{H}_2^0 = \tilde{H}_1^0\,(L \leftrightarrow R)$, corresponding to our choice $v_{L(R)1} = v_{L(R)2}$. Since we do not expect any dramatic change in the value of the parameters characterizing the Higgs potential when one considers the multigeneration case, in the evaluation of the diagrams of Figs.1,2 and 4, we have assumed the same scale for the cubic coupling γ as given in equ.(33).

Acknowledgements

We would like to thank R. Phillips and S. Sarkar for useful discussions.

[8]We do not consider here the limit on the axion decay constant (f_a) from the SN87A, which is still controversial.

Figure Captions

Figure 1. Diagrams generating an effective diagonal coupling of the axion to the ordinary charged fermions of the standard type: $(m_f/\chi)\bar{f}\gamma_5 f A$. As usual, the " $\times$ " along a fermion-line denotes a mass insertion, while at the end of a scalar-line it denotes a vacuum expectation value. f and F represent ordinary and extra-singlet fermions, respectively. $\tilde{H}_1^0$ and $\tilde{H}_2^0$ are the two pseudoscalar Higgs fields present in the model.

Figure 2. Diagram generating the amplitude for the decay of the 17 keV Simpson neutrino. The decay proceeds through the interactions of the singlet field $\bar{N}_2$, with which $\nu_{\tau L}$ forms a single effective Dirac state.

Figure 3. Diagram which generates an extra contribution for the axion coupling to the second generation fermions.

Figure 4. Diagram which dominates the amplitude for the decay of the right-handed electron neutrino. N_2 is a heavy singlet field, degenerate with $\nu_{\tau R}$.

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Mass [keV]	**Neutrino Spectrum**
$L^2/K \simeq 10^{-7}$	$\nu_1 \simeq \cos\theta\,\nu_{eL} - \sin\theta\,\nu_{\tau L}$ $\nu_2 \simeq \nu_{\mu L}$
$L \simeq 17$	$\nu_{3,4} \simeq \frac{1}{\sqrt{2}}(\cos\theta\,\nu_{\tau L} + \sin\theta\,\nu_{eL}) \pm \frac{1}{2}(N_{2L} - N_{2R})$
$R^2/K \simeq 3\cdot 10^6$	$\nu_5 \simeq \cos\theta\,\nu_{eR} - \sin\theta\,\nu_{\tau R}$ $\nu_6 \simeq \nu_{\mu R}$
$R \simeq 6\cdot 10^7$	$\nu_{7,8} \simeq \frac{1}{\sqrt{2}}(cos\theta\,\nu_{\tau R} + \sin\theta\,\nu_{eR}) \pm \frac{1}{2}(N_{2R} + N_{2R})$
$K \simeq 10^9$	$\nu_{9,10} \simeq \frac{1}{2}\left[(N_{1L} - N_{1R}) \pm (N_{3L} - N_{3R})\right]$ $\nu_{11,12} \simeq \frac{1}{2}\left[(N_{1L} + N_{1R}) \pm (N_{3L} + N_{3R})\right]$

Table 1: Neutral lepton spectrum in the "Universal Seesaw Model" for the 17 keV neutrino.

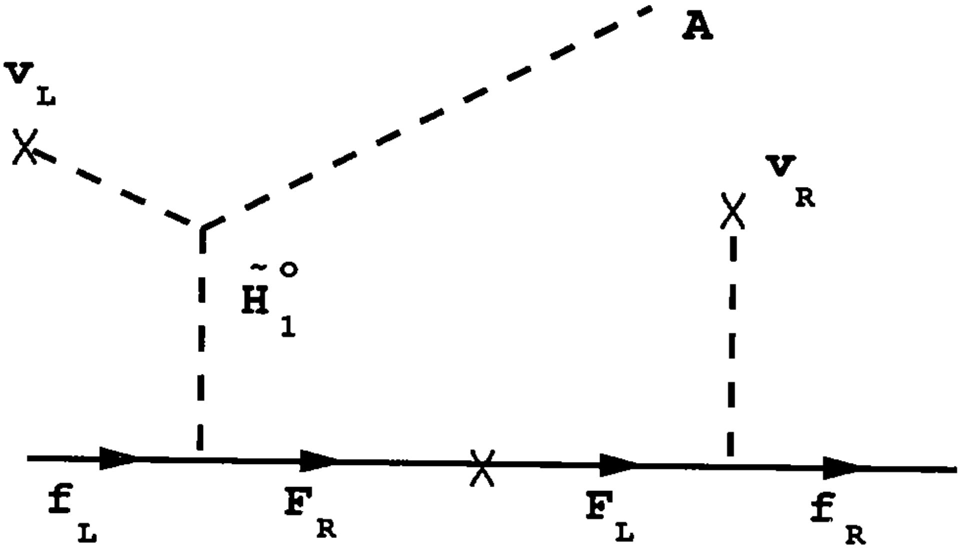


Fig. 1a

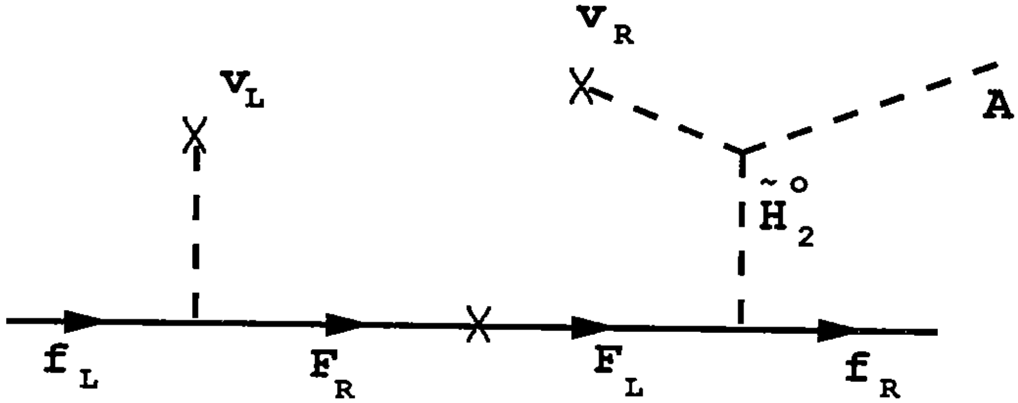
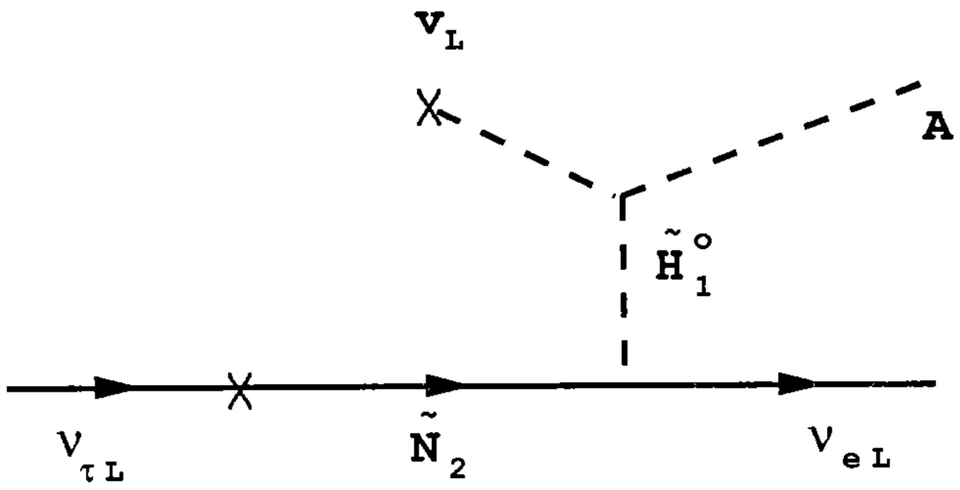
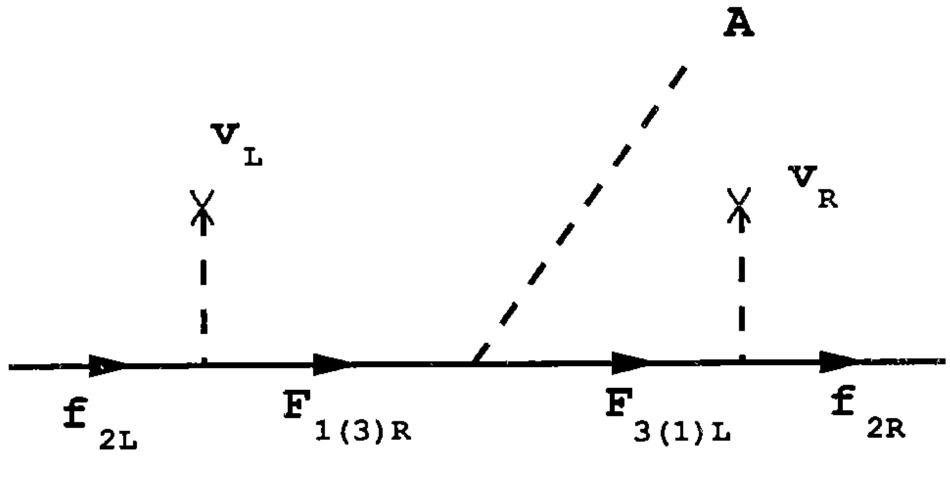


Fig. 1b



v_L
A
$\tilde{H}_1^{\circ}$
$\nu_{\tau L}$
$\tilde{N}_2$
ν_{eL}

Fig. 2



A
v_L
v_R
f_{2L}
$F_{1(3)R}$
$F_{3(1)L}$
f_{2R}

Fig. 3

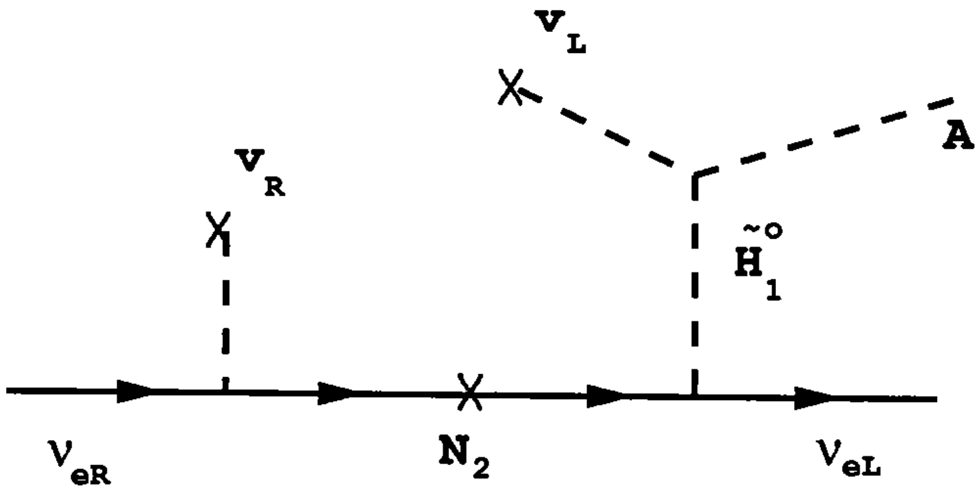


Fig. 4