

MULTIPLE SEESAW AT LOW ENERGY

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*(Received June 3, 2002)**Dedicated to Stefan Pokorski on his 60th birthday*

A new mechanism for generating neutrino masses without a high-energy mass scale is proposed. The mechanism needs a fundamental mass scale M in the 100–1000 TeV region and a minimal field content beyond the Standard Model one containing a pair of fermion singlets and a pair of weak doublet fermions for each neutrino mass, all of them with a mass of order M . The neutrino mass appears by a multiple seesaw-type tree-level diagram. We provide an explicit model based on supersymmetry and an Abelian symmetry which provides the required fermion mass matrix. The mechanism is natural in the context of string theories with a low fundamental scale. Within an explicit example where the Abelian symmetry is also responsible for the generation of fermion masses and mixings, we give a hint relating the fermion mass matrices and the weak mixing angle. By assuming the weak–strong couplings unification, one naturally finds $\sin^2 \theta_w = 1/4$ at the fundamental scale.

PACS numbers: 14.60.Pq, 14.60.St, 12.60.Jv

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1. The minimal model

Let us consider one active neutrino in a theory whose fundamental mass scale M is low, of the order of 100–1000 TeV [3, 4]. In such a context, the smallness of the neutrino masses cannot be explained by the standard seesaw mechanism [1], which would give much too large values. A successful generation of neutrino masses asks in this case for two conditions. The first is to forbid the operator responsible for the seesaw mass $(1/M) HH\nu_L\nu_L$. The second is to generate a small neutrino mass by some other mechanism.

Our proposal involves a minimal set of two Standard Model singlet fermions Ψ_1, Ψ_2 and two $SU(2)_L$ doublets Ψ_4, Ψ_3 of hypercharges $Y = \pm 1/2$, respectively. The symmetries of the model, whose discussion (in a particular realization) is postponed to the next section, give mass terms in the Lagrangian of the form

$$\frac{1}{2} V \mathcal{M} V^T, \quad (1)$$

where the vector V denotes the collection of fermionic fields

$$V = (\nu_L, \Psi_1, \Psi_2, \Psi_3, \Psi_4).$$

The mass matrix \mathcal{M} in (1) is given by

$$\mathcal{M} = \begin{pmatrix} 0 & m_1 & 0 & 0 & 0 \\ m_1 & M_1 & M_2 & m_2 & m_3 \\ 0 & M_2 & 0 & m_4 & m_5 \\ 0 & m_2 & m_4 & 0 & M_3 \\ 0 & m_3 & m_5 & M_3 & 0 \end{pmatrix}. \quad (2)$$

The entries m_i denote electroweak-type mass terms, given by the vacuum expectation value(s) of the Higgs field(s), whereas the capital letters M_a denote Majorana mass terms (independent of the electroweak scale), generically of the order of the fundamental scale, in our case assumed to be of the order 100–1000 TeV. The eigenvalues/eigenvectors of the mass matrix (2) are obtained rather easily in the relevant approximation $M_a \gg m_i$. In this limit, the mass matrix has a block-diagonal form of a 3×3 matrix times a 2×2 mass matrix. There are four heavy eigenstates of mass approximately given by

$$\begin{aligned} \lambda_{2,3} &\simeq \frac{1}{2} \left(M_1 \pm \sqrt{M_1^2 + 4M_2^2} \right), \\ \lambda_{4,5} &\simeq \pm M_3. \end{aligned} \quad (3)$$

In this block-diagonal limit, the lightest neutrino stays massless, despite the presence of the Majorana mass term $M_1\Psi_1\Psi_1$ in (2). This is easily explained

by considering the 3×3 upper block in (2), which has zero determinant. Indeed, by inverting the 2×2 matrix block for the singlets Ψ_1, Ψ_2

$$\mathcal{M}_2 = \begin{pmatrix} M_1 & M_2 \\ M_2 & 0 \end{pmatrix}, \quad (4)$$

we get in the seesaw diagram entries in the singlet propagators of the type $S_{12}(p=0) = (1/M_2)$ and $S_{22}(p=0) = (M_1/M_2)^2$, where p is the momentum flowing into the fermion propagator. The main point here is that the $S_{11}(p=0)$ entry *is zero*. Consequently, the usual seesaw diagram obtained by integrating out the heavy Ψ_1 field is forbidden. The neutrino mass appears only by taking into account the small 3×2 block off-diagonal terms in (2). In this case, the light eigenvalue is most easily obtained by computing the determinant of the matrix (2)

$$\det \mathcal{M} = -2 m_4 m_5 m_1^2 M_3. \quad (5)$$

By combining (3) and (5), we obtain the value of the light neutrino eigenstate, which in our simplified model is mostly given by the electron neutrino

$$m_\nu = \lambda_1 \simeq -\frac{2m_1^2 m_4 m_5}{M_2^2 M_3}. \quad (6)$$

As expected, the corresponding eigenvector $|\hat{\nu}_L\rangle$ is mainly composed of the lightest neutrino state. The explicit expression is given by

$$|\hat{\nu}_L\rangle \simeq |\nu_L\rangle - \frac{2m_1 m_4 m_5}{M_2^2 M_3} |\Psi_1\rangle - \frac{m_1}{M_2} |\Psi_2\rangle + \frac{m_1 m_5}{M_2 M_3} |\Psi_3\rangle + \frac{m_1 m_4}{M_2 M_3} |\Psi_4\rangle. \quad (7)$$

The final result (6) can be easily understood as a result of a multiple seesaw-type diagram obtained by integrating out the heavy fields $\Psi_1, \Psi_2, \Psi_3, \Psi_4$ and inserting the appropriate Higgs v.e.v.(s). The first step is to integrate out the weak fermion doublets ψ_3, ψ_4 . The Majorana mass matrix (4) is then modified to

$$\mathcal{M}_2 = \begin{pmatrix} M_1 & M_2 \\ M_2 & \frac{2m_4 m_5}{M_3} \end{pmatrix}. \quad (8)$$

The second step involves integrating out the singlet ψ_1, ψ_2 fermion fields in the remaining 3×3 mass matrix

$$\mathcal{M}_3 = \begin{pmatrix} 0 & m_1 & 0 \\ m_1 & M_1 & M_2 \\ 0 & M_2 & \frac{2m_4 m_5}{M_3} \end{pmatrix}, \quad (9)$$

leading to the eigenvalues $\lambda_{2,3}$ in (3) and $\lambda_1 = m_\nu$ in (6).

The Lagrangian leading to the mass matrix (2) breaks the U(1) lepton number which we define to act on the various fields as

$$\begin{aligned}\nu_L &\rightarrow e^{i\alpha}\nu_L, & \psi_1 &\rightarrow e^{-i\alpha}\psi_1, & \psi_2 &\rightarrow e^{i\alpha}\psi_2, \\ \psi_3 &\rightarrow e^{-i\alpha}\psi_3, & \psi_4 &\rightarrow e^{-i\alpha}\psi_4.\end{aligned}\quad (10)$$

The breaking of the lepton number (by $\Delta L = 2$) is due to the Majorana mass parameters M_1 and M_3 . However, only M_3 is relevant for the neutrino mass generation. The smallness of the neutrino mass (6) is then understood by the fact ψ_3 and ψ_4 do not directly couple to the lightest neutrino. The transmission of the lepton number breaking goes therefore through the heavy fermions ψ_1, ψ_2 and generates a cubic suppression in the heavy masses (6).

We now add the usual mass for the electron and a new mass mixing ψ_3^+ and the right-handed electron in the effective theory. The charged lepton mass matrix in the theory is of the form

$$(e_L^c \Psi_3^+) \begin{pmatrix} m_6 & 0 \\ m_7 & M_3 \end{pmatrix} \begin{pmatrix} e_R \\ \Psi_4^- \end{pmatrix}, \quad (11)$$

and the physical electron mass is approximately given by $m_e \sim m_6$, while the charged new leptons have a mass of order M_3 . The physical charged states mix slightly the light and the heavy states. The physical electron state is for example given by

$$|\hat{e}_R\rangle \simeq |e_R\rangle - \frac{m_7}{M_3} |\Psi_4^-\rangle, \quad |\hat{e}_L\rangle \simeq |e_L\rangle - \frac{m_7 m_6}{M_3^2} |\Psi_3^-\rangle. \quad (12)$$

With the minimal field content (2) only one neutrino linear combination acquires a mass at tree level. In order to give a tree-level mass of the type (6) for a second neutrino, we need a second pair of singlet fields ψ_5, ψ_6 and a second pair of weak fermion doublets ψ_7, ψ_8 . The second light neutrino should couple to ψ_1 and ψ_5 , but not to ψ_2 and ψ_6 . The need of doubling the exotic fermion spectrum can be easily shown by following the integrating out procedure outlined above. Indeed, first of all we need to double the singlet fermion content in order to provide two (linear combinations of) light neutrinos to couple to two different singlets. On the other hand, by integrating out the weak doublets we generate the Majorana mass matrix for the singlet fields. With only one pair of weak doublets, the Majorana mass matrix for the singlets will have zero eigenvalues and as a result one zero mass eigenstate will survive in the physical spectrum. Adding a second pair of weak doublets will solve this problem by generating large Majorana masses for all singlets $\psi_1, \psi_2, \psi_5, \psi_6$. The charge lepton mass matrix in this

case is a straightforward generalization of (11). The light charged states get mixed with heavy charged leptons, analogously to (12). By promoting the mass entries in (2) and (11) to 3×3 matrices in the flavor space, all the three light neutrinos get masses.

The mixing between the light and the heavy states (7) and (12) generates contributions to the anomalous magnetic moment of the electron and muon, the electric dipole moment and $\mu \rightarrow e\gamma$. For values of the fundamental scale $100 \text{ TeV} < M < 1000 \text{ TeV}$, the mixings are, however, very small and the results of such processes are experimentally unobservable, independently of the mixing angles between the light neutrino flavors.

There already exist in the literature other mechanisms [2] for getting small neutrino masses in the presence of large extra dimensions in models with a low string scale [3,4]. The mechanism we put forward in our paper does not rely crucially on large extra dimensions. Its phenomenology is different in nature from the previous ones and future experiments could distinguish it from other mechanisms.

The model just presented was based on a minimal field content and a specific mass matrix (2), which contained a certain number of important zeros. Their presence must be assured by some symmetry of the underlying theory. We were unable to find such a symmetry for an appropriate non-supersymmetric extension of the Standard Model. The simplest example we found has low-energy supersymmetry and an additional Abelian symmetry, to be explained in the next section.

2. A supersymmetric example based on an Abelian symmetry

The goal of this section is to present an explicit example of a symmetry which produces a mass matrix of the form (2). It will also forbid radiative corrections to generate the usual seesaw neutrino mass by symmetry arguments. The model we consider is the minimal SuperSymmetric Standard Model (MSSM), enlarged with two additional superfield singlets ϕ_1, ϕ_2 and two superfields ϕ_3, ϕ_4 of $SU(2)_L \times U(1)_Y$ quantum numbers $(2, -1/2)$, $(2, 1/2)$. Their fermionic components are the heavy fermion fields present in the previous section. The model is, therefore, the minimal supersymmetrization of the previous one. The symmetry which will guarantee the existence of the mass matrix (2) is an Abelian $U(1)_X$ symmetry. For the purpose of the present section, its nature (global or local, family dependent or independent) is irrelevant. We will come back later on its interesting implications if the symmetry is local. There are several possible charge assignments which fulfill our requirements and we just present here one particularly interesting example. We denote by small letters $h_1, h_2, l \dots$ the $U(1)_X$ charges of the MSSM fields $H_1, H_2, L \dots$ and by $x_1 \dots x_4$ the $U(1)_X$ charges of the heavy

superfields $\phi_1 \dots \phi_4$. As usual, we assume the existence of the (super)field Φ singlet under the Standard Model of $U(1)_X$ charge -1 , whose scalar v.e.v. spontaneously breaks the Abelian symmetry. We also assume, for reasons related to proton stability and lepton flavor violation, the existence of R -parity under which the heavy superfields $\phi_1 \dots \phi_4$ have a matter-type parity.

We consider the charge assignments:

$$h_1 = 4, \quad h_2 = 0, \quad l = -1, \quad e = 0, \quad x_1 = x_3 = 2, \quad x_2 = x_4 = -2. \quad (13)$$

The relevant terms in the superpotential of this model, compatible with holomorphicity and the charge assignment (13) are

$$\begin{aligned} W = & \lambda_1 \left(\frac{\Phi}{M} \right) L H_2 \phi_1 + \lambda_2 \left(\frac{\Phi}{M} \right)^4 H_2 \phi_1 \phi_3 + \lambda_3 \left(\frac{\Phi}{M} \right)^4 H_1 \phi_1 \phi_4 \\ & + \lambda_4 H_2 \phi_2 \phi_3 + \lambda_5 H_1 \phi_2 \phi_4 + M'_1 \left(\frac{\Phi}{M} \right)^4 \phi_1 \phi_1 + M'_2 \phi_1 \phi_2 \\ & + M'_3 \phi_3 \phi_4 + \mu' \left(\frac{\Phi}{M} \right)^4 H_1 H_2 + \lambda_6 \varepsilon^3 L E H_1 + \varepsilon^6 \lambda_7 \Phi_3 E H_1, \quad (14) \end{aligned}$$

where M is the fundamental mass scale of the model. We denote the various v.e.v.'s by $\langle \Phi \rangle / M \equiv \varepsilon < 1$, $\langle H_{1,2} \rangle \equiv v_{1,2}$. Due to their large supersymmetric masses in (14), assumed to be much larger than that the supersymmetry breaking scale, the scalar components of $\phi_1 \dots \phi_4$ get no v.e.v.'s. Then the fermionic mass matrix derived from (14) is precisely of the form (2), with the various mass parameters being equal to

$$\begin{aligned} m_1 &= \lambda_1 \varepsilon v_2, \quad m_2 = \lambda_2 \varepsilon^4 v_2, \quad m_3 = \lambda_3 \varepsilon^4 v_1, \quad m_4 = \lambda_2 v_2, \quad m_5 = \lambda_5 v_1, \\ M_1 &= M'_1 \varepsilon^4, \quad M_2 = M'_2, \quad M_3 = M'_3, \quad \mu = \mu' \varepsilon^4. \quad (15) \end{aligned}$$

The natural values for the “fundamental” couplings in (14) are $\lambda_i \sim 1$ and $M'_i, \mu' \sim M$. In this case, by applying the expression (6), we obtain the neutrino mass

$$m_\nu \sim \varepsilon^2 \frac{v_2^3 v_1}{M^3}. \quad (16)$$

For low values of $\tan \beta = v_2 / v_1$ and by taking as a particular example the value $\varepsilon \simeq 0.22$ motivated by considering the Abelian symmetry $U(1)_X$ as responsible for the generation of fermion masses and mixings (see next section), we find from (16) for the electron neutrino $M \sim 500 \text{ TeV}$. By a slight change of charge assignments we can also consistently find $100 \text{ TeV} < M < 1000 \text{ TeV}$. Notice that the charge assignment (13) generates an effective μ -term in (15) of the order $(1/400) M$, which is, therefore, in the electroweak energy range.

Consequently, the model automatically produces a successful μ -term in the low-energy theory.

In this model the absence of the seesaw operator LLH_2H_2 after integrating out the vector-like heavy states is not an accident. This operator has $U(1)_X$ charge -2 and is forbidden by holomorphicity in the superpotential¹. On the other hand, the neutrino mass (16) obtained by integrating out the heavy states is described by the higher-dimensional operator

$$\frac{\varepsilon^2}{M^3} LLH_2H_2 (H_1H_2) , \quad (17)$$

of $U(1)_X$ charge 2 that can appear in the superpotential. As we just proved, this operator is indeed produced at tree-level by a diagram containing three propagators of massive states. The charge of the operator (17) explains also the ε^2 factor in (16). Notice also that the imposition of R -parity and assignment of matter parity for the heavy fields $\phi_1 \cdots \phi_4$ forbids in the superpotential operators of the type

$$M_5 \left(\frac{\Phi}{M} \right)^2 H_2 \phi_3 + M_6 \left(\frac{\Phi}{M} \right)^2 H_1 \phi_4 . \quad (18)$$

Operators of the type (18), if present, would generate mixings of the heavy leptons with Higgsinos which would ask for a more careful treatment of the resulting (7×7) mass matrix. It is not obvious that such terms would destroy the prediction for the neutrino mass (16) or create other serious problems, but we choose here the simple option of eliminating them by imposing R -parity.

The lepton flavor violating processes in our model are of two types. The first are typical of supersymmetric models and constrains as usual the sparticle spectrum. The second type are generated by the new lepton violating sector $\Phi_1 \dots \Phi_4$ in the superpotential (14). They put a lower bound on the fundamental scale, which is however largely satisfied in our class of models by imposing phenomenologically relevant values for neutrino masses.

3. Electroweak angle and Froggatt–Nielsen mechanism at low energy

Up to now the details of the Abelian flavor symmetry were irrelevant for producing the multiple seesaw mechanism we are proposing. We are now

¹ Actually, non-analyticity in ε of the type $1/\varepsilon^n$ can be found in the superpotential if chiral-type, with respect of $U(1)_X$, heavy states are integrated out. A sufficient condition to obtain analytic superpotential in ε is to integrate out vector-like heavy states, which is indeed the case under consideration.

tempted to go one step further and try to use the $U(1)_X$ symmetry as a horizontal symmetry responsible for the generation of fermion masses and mixings [5]. The quark and charged lepton mass matrices in this framework can be related to mixed anomalies of $U(1)_X$ with the Standard Model gauge group [6]. Let us denote by A_1, A_2, A_3 the mixed anomalies $[U(1)_Y]^2 U(1)_X$, $[SU(2)_L]^2 U(1)_X$ and $[SU(3)_c]^2 U(1)_X$ respectively. A straightforward generalization of the results of [6] give the relations between the determinants of the up quark Y_U , down quark Y_D and charged leptons Y_L mass matrices

$$\begin{aligned}\det(Y_D^{-2} Y_L^2) &= \varepsilon^{A_1 + A_2 - \frac{8}{3} A_3 - 2(h_1 + h_2) - 2(x_3 + x_4)}, \\ \det(Y_U Y_D^{-2} Y_L^3) &= \varepsilon^{\frac{3}{2}(A_1 + A_2 - 2A_3) - 3(x_3 + x_4)}.\end{aligned}\quad (19)$$

Let us define the number $x = h_1 + q_3 + d_3 = d_1 + l_3 + e_3$ related to the angle $\tan \beta = (v_2/v_1)$ approximately by the relation $\tan \beta = (m_t/m_b) \varepsilon^x$. Phenomenologically relevant values of x are then $x = 0, 1, 2$. Considering the most viable mass matrices obtained in the references [6, 7] we find the relations

$$\begin{aligned}A_1 + A_2 - 2A_3 &= 12 + 2x + 2(x_3 + x_4), \\ A_1 + A_2 - \frac{8}{3}A_3 &= 2(h_1 + h_2 + x_3 + x_4).\end{aligned}\quad (20)$$

If we insert the charges (13) and consider the large $\tan \beta$ regime $x = 0$, we naturally find a solution to the Eqs. (20)

$$A_1 = 3A_2, \quad A_2 = A_3. \quad (21)$$

In the regime $\tan \beta \sim 1$, we also find the solution (21) provided that we change slightly the H_1 Higgs $U(1)_X$ charge to $h_1 = 5$. The relation (21) provides an interesting hint for the value of the weak mixing angle in such a model with low fundamental scale [3]. Indeed, in an effective string theory context, the values of the mixed gauge anomalies can be related to the value of the weak mixing angle at the energy scale, where the flavor symmetry is spontaneously broken. The relevant string theory under consideration is here Type I string theory with low fundamental mass scale $M_I \sim M$. The $SU(2)_L$ and $U(1)_Y$ gauge groups leave on a stack of “electroweak” branes. The tree-level gauge couplings in this case are given by (we consider for illustration D9 branes) [10]

$$\frac{4\pi^2}{g_a^2} = s + s_{ak} m_k, \quad (22)$$

where $s = ReS$ is the four-dimensional dilaton and $m_k = ReM_k$ are twisted moduli present in orbifold-type compactifications. Under the $U(1)_X$ gauge

transformation the twisted moduli get shifted according to

$$\begin{aligned} V_X &\rightarrow V_X + \frac{i}{2}(A - \bar{A}), \\ M_k &\rightarrow M_k + \frac{1}{2}\varepsilon_k A. \end{aligned} \quad (23)$$

Let us consider for simplicity the case of just one relevant twisted modulus M . In this case, the mixed anomalies are canceled provided the following condition holds

$$\frac{\varepsilon}{4\pi^2} = \frac{A_1}{s_1} = \frac{A_2}{s_2} = \frac{A_3}{s_3}. \quad (24)$$

If the twisted field has a rather large v.e.v. $\langle m \rangle \gg \langle s \rangle$, then the weak mixing angle at the scale of the symmetry breaking M is determined by the relation

$$\sin^2 \theta_W \simeq \frac{s_2}{s_1 + s_2} = \frac{A_2}{A_1 + A_2} = \frac{1}{4}. \quad (25)$$

This value is rather close to the experimentally measured one at the M_Z scale $\sin^2 \theta_W \simeq 0.231$, as pointed long ago by Weinberg and reemphasized recently in [9] in specific models. The model we put forward here is the MSSM with an additional Abelian symmetry. It is, therefore, natural to analyze the RGE for the gauge couplings and to explicitly find the energy scale at which (25) holds. A straightforward computation, with the one-loop β -functions of the SM between M_Z and an effective MSSM threshold, M_S , and with the MSSM β -functions above M_S gives for the scale μ_0 where (25) holds the expression

$$\mu_0 = M_Z \left(\frac{M_Z}{M_S} \right)^{1/4} e^{\frac{\pi}{4} \left[\frac{1}{\alpha_Y(M_Z)} - \frac{3}{\alpha_2(M_Z)} \right]} \simeq 90 - 150 \text{ TeV}, \quad (26)$$

a value within the energy range of the fundamental scale of our model.

We consider (25),(26) as an encouraging hint in order to embed our model containing the additional lepton-like multiplets into a fundamental theory. The relation $A_2 = A_3$ in (21) implies $s_2 = s_3$ and, therefore, unification of the $SU(3)_c$ strong coupling with the $SU(2)_L$ gauge coupling at the string scale. Even if this is an appealing feature, the logarithmic evolution of gauge couplings [11] cannot explain this strong-weak unification. There are several possibilities in order to overcome this difficulty. One possibility we mention here is the case where the $SU(3)_c$ strong interaction brane (but not the electroweak branes) experience a new large compact dimension in the TeV range [12]. The strong coupling will then start having a power-law evolution [4,13] and it will unify with the weak $SU(2)_L$ coupling at low energy. This is

similar with the unification mechanism proposed in [4], but uses the power-law evolution just for the strong gauge coupling. Another possibility is that the $SU(3)$ gauge field couples to additional moduli fields on the strong coupling brane. The v.e.v. of the additional (twisted or untwisted) fields can then modify the strong coupling in a phenomenologically successful, but unpredictable from our arguments, way.

We dedicate this work to our collaborator and friend Stefan Pokorski, on the occasion of his 60th birthday. Stefan greatly stimulated the interest for phenomenology of E.D. by completing, through uncountable discussions over the last years, his particle physics education. We would like to thank Christophe Grojean and Stephane Lavignac for helpful discussions. Work supported in part by the RTN European Program HPRN-CT-2000-00148.

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