LBNL-39422
UC-414
Preprint

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October 1996
Submitted to
Physics Letters B

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LBNL-39422
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This work was supported in part by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Division of High Energy Physics, of the U.S. Department of Energy under Contract No. DE-AC03-76SF00098; the National Science Foundation under Grant PHY-95-14797; and the Alfred P. Sloan Foundation.

October 16, 1996

LBNL-39422
UCB-PTH-96/40

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Abstract

Horizontal flavor symmetries can drastically suppress Dirac neutrino masses well below those of the corresponding charged leptons. We show that models can be constructed where the light neutrino mass eigenvalues are small enough to give the MSW solution to the solar neutrino problem, with a right-handed neutrino scale no larger than a TeV. We present a model of this type where the right-handed neutrino scale is generated by the spontaneous breakdown of gauged $U(1)_{B-L}$, in a radiative breaking scenario driven by the right-handed neutrino Yukawa couplings. The model allows for a solution to the μ problem, and predicts the existence of a Z' boson within the reach of the LHC or the Tevatron.

*This work was supported in part by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U.S. Department of Energy under Contract DE-AC03-76SF00098. HM was also supported in part by the National Science Foundation under grant PHY-95-14797 and by the Alfred P. Sloan Foundation.

1 Introduction

The seesaw mechanism [1] has been proposed as a natural explanation for the lightness of the three known neutrino species. In models with right-handed neutrinos, the 'seesaw' refers to the widely disparate eigenvalues of the neutrino mass matrix

$$\begin{pmatrix} \sim 0 & M_{LR} \\ (M_{LR})^T & M_{RR} \end{pmatrix} , \tag{1.1}$$

where M_{LR} are the entries generated through ordinary electroweak Higgs couplings, and M_{RR} are the Majorana masses for the right-handed states. Since the the right-handed neutrinos are singlets under the standard model gauge group, they can develop masses that are much larger than the electroweak scale. The result is immediate in the case of one neutrino flavor: the eigenvalues of the two-by-two matrix above are of order M_{RR} and M_{LR}^2/M_{RR}, assuming $M_{RR} \gg M_{LR}$.

This result is significant in light of the experimental indications that the light neutrino mass eigenvalues are smaller than 10 eV[†]. The observed solar neutrino deficit may be explained by MSW or vacuum ν_e-ν_x oscillation, with $\Delta m^2 \sim 10^{-5}$ eV^2 or $\sim 10^{-10}$ eV^2, respectively, where Δm^2 is the difference in squared masses of the two relevant neutrino flavors. The atmospheric neutrino anomaly may result from ν_μ-ν_τ oscillation with $\Delta m^2 \sim 10^{-2}$ eV^2. If neutrinos are the hot dark matter in the Universe, then $m \sim 4$ eV is preferred. Furthermore, if neutrino masses are less than approximately 10 eV, it is possible to avoid a number of other cosmological constraints, including those from big bang nucleosynthesis, overclosure of the Universe, and the distortion of the cosmic microwave background radiation. Finally, the recent LSND results suggest ν_e-ν_μ oscillation with $\Delta m^2 \sim 10$ eV^2. While it is unlikely that all the current neutrino anomalies will turn out to be real, it is still reasonable to conclude that the mass range below 10 eV is the most interesting one for neutrino physics[‡].

In implementing the seesaw mechanism, many have assumed that M_{LR} should be comparable to the mass of the corresponding charged lepton or

[†]The mass ranges given in this paragraph can be found in Ref. [2], and in references therein.

[‡]See Ref. [3] for counterexamples.

up-type quark, as a reasonable ansatz. This is natural in many models, such as SO(10) grand unified theories where the Yukawa matrices of the charged leptons and/or the up-type quarks are related to those of the neutrinos by a gauge symmetry. In models of this type, one needs $M_{RR} \simeq 10^{12}$ GeV to keep the tau neutrino mass below 10 eV, assuming $M_{LR}^{33} \sim m_{top}$. While it is quite interesting that the the seesaw mechanism allows the light neutrino masses to serve as a probe of physics at very high energy scales, it is unfortunate that the mechanism cannot be proven directly in experiment. In particular, there is no hope of producing such heavy right-handed neutrinos or studying their interactions at collider experiments in the imaginable future.

What we will demonstrate in this letter is that horizontal flavor symmetries often naturally lead to a suppression of the matrix M_{LR}, so that its eigenvalues are significantly *smaller* than the corresponding charged lepton or up-type quark masses. As is well known, models with horizontal flavor symmetries have a separate motivation in that they provide a natural framework for understanding the hierarchical form of other fermion mass matrices. In theories where the hierarchy between the electroweak scale and any fundamental high-energy scale (such as M_{Planck}) is stabilized by supersymmetry, horizontal symmetries have an additional virtue: they also restrict the form of the soft supersymmetry-breaking scalar mass matrices and thereby may suppress the large flavor changing neutral current processes expected when the superparticle spectrum is generic [4].

What is significant about the suppression of M_{LR} is that it allows us to construct models where the right-handed neutrino scale is at or slightly above the electroweak scale and where the light neutrino masses fall in a desirable range. A right-handed neutrino scale near 1 TeV is very natural since we can imagine this situation arising in a radiative breaking scenario: If the right-handed neutrinos ν transform under an additional U(1) gauge symmetry, then the mass squared of an exotic Higgs field ρ that is also charged under this U(1) can be driven negative through renormalization group running, as a consequence of the Yukawa coupling $\rho\nu\nu$. This is completely analogous to the situation in the minimal supersymmetric standard model (MSSM), where the large top quark Yukawa coupling drives the up-type Higgs mass negative, triggering electroweak symmetry breaking at a scale comparable to the superparticle masses. In the present case, the additional U(1) gauge symmetry is also spontaneously broken near the electroweak scale and the right-handed neutrinos develop Majorana masses. Indeed, such additional U(1) gauge factors appear in many superstring compactifications [5]. The model we will present below demonstrates that the new Z' boson is likely to be within the reach of the LHC, or the Tevatron after the main injector upgrade.

Below we will elaborate on these points by presenting a specific model based on the flavor group $(S_3)^3$. Aside from bringing new physics down to the TeV scale, the neutrino physics of this model is interesting in its own right. Thus, this discussion complements the phenomenology presented in Ref. [6, 7] §.

2 A Model

The flavor group that we will assume is the discrete, non-Abelian gauge symmetry $(S_3)^3$. While an Abelian horizontal symmetry may be as effective in suppressing the Dirac mass matrix M_{LR}, non-Abelian symmetries provide a natural means of suppressing flavor-changing neutral current effects originating from superparticle loops, as we describe below. This particular flavor symmetry has been discussed extensively in Refs. [6, 7, 9], so we will provide only a brief review.

The group S_3 has three representations, $\mathbf{2}$, $\mathbf{1}_S$ and $\mathbf{1}_A$, where the latter two are trivial and nontrivial singlet representations, respectively. The three generations of the quark chiral superfields Q, U, and D, are assigned to $\mathbf{2}+\mathbf{1_A}$ representations of S_3^Q, S_3^U and S_3^D, respectively. The first two generation fields are embedded in a doublet to maintain the degeneracy of the corresponding squarks in the flavor symmetry limit. After the flavor symmetry is spontaneously broken, the remaining approximate squark degeneracy is sufficient to suppress flavor changing neutral current effects from superparticle exchange, like those contributing to K^0-$\overline{K}^0$ mixing. The Higgs fields both transform as $(\mathbf{1}_A, \mathbf{1}_A, \mathbf{1}_S)$'s, so that the top quark Yukawa coupling is invariant under the flavor symmetry group; this provides a natural explana-

§The neutrino mass spectrum has been considered in the context of other flavor groups [8]. However, we know of no reference in which a drastic suppression of the neutrino Dirac mass matrix was either discussed or achieved.

tion for the heaviness of the top quark relative to the other fermions. The remaining Yukawa couplings can then be treated as small flavor symmetry breaking spurions, with the transformation properties

$$Y_U \sim \left(\begin{array}{c|c} (\tilde{\mathbf{2}},\tilde{\mathbf{2}},\mathbf{1}_S) & (\tilde{\mathbf{2}},\mathbf{1}_S,\mathbf{1}_S) \\ \hline (\mathbf{1}_S,\tilde{\mathbf{2}},\mathbf{1}_S) & (\mathbf{1}_S,\mathbf{1}_S,\mathbf{1}_S) \end{array}\right) , \tag{2.2}$$

$$Y_D \sim \left(\begin{array}{c|c} (\tilde{\mathbf{2}},\mathbf{1}_A,\mathbf{2}) & (\tilde{\mathbf{2}},\mathbf{1}_A,\mathbf{1}_A) \\ \hline (\mathbf{1}_S,\mathbf{1}_A,\mathbf{2}) & (\mathbf{1}_S,\mathbf{1}_A,\mathbf{1}_A) \end{array}\right) , \tag{2.3}$$

where we use the notation $\tilde{\mathbf{2}} \equiv \mathbf{2} \otimes \mathbf{1}_A$,¶. Notice that the Yukawa matrices above involve at most 7 irreducible multiplets of $(S_3)^3$. In Ref. [9], $(S_3)^3$ was spontaneously broken by a set of 'flavon' fields ϕ, representing only four of these multiplets,

$$\begin{array}{ll} \phi_1(\tilde{\mathbf{2}},\tilde{\mathbf{2}},\mathbf{1}_S) & \phi_2(\tilde{\mathbf{2}},\mathbf{1}_S,\mathbf{1}_S) \\ \phi_3(\tilde{\mathbf{2}},\mathbf{1}_A,\mathbf{2}) & \phi_4(\mathbf{1}_S,\mathbf{1}_A,\mathbf{1}_A) \end{array}$$

This was the minimal number needed to obtain realistic quark masses and mixing angles [10], assuming the Yukawa textures

$$Y_U = \left(\begin{array}{cc|c} h_u & h_c\lambda & -h_t V_{ub} \\ O(h_u) & h_c & -h_t V_{cb} \\ \hline 0 & 0 & h_t \end{array}\right), \tag{2.4}$$

$$Y_D = \left(\begin{array}{cc|c} h_d & h_s\lambda & 0 \\ O(h_d) & h_s & 0 \\ \hline 0 & 0 & h_b \end{array}\right). \tag{2.5}$$

Here the h_q are Yukawa couplings, and $\lambda \simeq 0.22$ is the Cabbibo angle. The hierarchical form of the Yukawa matrices presented above can be understood as a consequence of a sequential breaking of the flavor symmetry. Since different components of a single multiplet will generally obtain comparable vevs at a given stage of symmetry breaking, we must assume that the 2×2 blocks of Y_U and Y_D are each generated by two flavon fields that acquire vevs at somewhat different scales, namely $\phi_1 = \phi_1' + \phi_1''$ and $\phi_2 = \phi_2' + \phi_2''$, where

$$\phi_1' = \left(\begin{array}{cc} 0 & a' h_c\lambda \\ 0 & h_c \end{array}\right) , \quad \phi_1'' = \left(\begin{array}{cc} h_u & O(h_u) \\ O(h_u) & O(h_u) \end{array}\right), \tag{2.6}$$

¶ $\tilde{\mathbf{2}} = (a,b)$ is equivalent to $\mathbf{2} = (b,-a)$.

and

$$\phi_2' = \begin{pmatrix} 0 & ah_s\lambda \\ 0 & h_s \end{pmatrix} \; , \; \phi_2'' = \begin{pmatrix} h_d & \mathcal{O}(h_d) \\ \mathcal{O}(h_d) & \mathcal{O}(h_d) \end{pmatrix} . \tag{2.7}$$

Here a and a' are order one constants, with $a - a' = 1$. Note that the entries labelled $\mathcal{O}(h_u)$ and $\mathcal{O}(h_d)$ above can be set to zero without noticeably affecting the quark masses and mixing angles. However, in any estimates where these entries are significant, we will assume they are nonvanishing with the magnitudes given above.

The model is extended to the charged lepton sector by assigning the L and E chiral superfields to $\mathbf{2} + \mathbf{1_A}$ representations of S_3^D and S_3^Q respectively. This is the only choice that leads to a qualitative similarity between the down quark and charged lepton Yukawa matrices, while simultaneously forbidding dangerous baryon-number-violating, Planck-suppressed operators (e.g. $QQQL/M_{Pl}$) in the flavor symmetry limit. Differences between the charged lepton and down quark mass eigenvalues can originate from fluctuations in the order 1 coefficients that multiply the symmetry breaking operators; thus we assume that the electron-muon Yukawa matrix is given by $3\phi_2' + \frac{1}{3}\phi_2''$.

If right-handed neutrinos ν are to be included in the model, then we must decide on their transformation properties under $(S_3)^3$. A natural choice is to repeat the $\mathbf{2} + \mathbf{1_A}$ representation structure of the other matter fields. Aesthetics also suggests that we consider the possibility that the ν transform under S_3^U, the only S_3 factor that we haven't utilized in the lepton sector[‖]. As before, we can then determine flavor structure of M_{LR} and M_{RR}:

$$M_{RR} \sim \left(\begin{array}{c|c} (\mathbf{1}_S, \mathbf{2}+\mathbf{1}_S, \mathbf{1}_S) & (\mathbf{1}_S, \tilde{\mathbf{2}}, \mathbf{1}_S) \\ \hline (\mathbf{1}_S, \tilde{\mathbf{2}}, \mathbf{1}_S) & (\mathbf{1}_S, \mathbf{1}_S, \mathbf{1}_S) \end{array} \right) \tag{2.8}$$

$$M_{LR} \sim \left(\begin{array}{c|c} (\mathbf{1}_A, \tilde{\mathbf{2}}, \mathbf{2}) & (\mathbf{1}_A, \mathbf{1}_S, \mathbf{2}) \\ \hline (\mathbf{1}_A, \tilde{\mathbf{2}}, \mathbf{1}_A) & (\mathbf{1}_A, \mathbf{1}_S, \mathbf{1}_A) \end{array} \right) \tag{2.9}$$

[‖]In fact, the choice S_3^Q would not give us the desired suppression of M_{LR}, which we would then expect to be of the same order as the charged lepton masses. The choice S_3^D does give us a high degree of suppression, but has other phenomenological difficulties, as we discuss in the Appendix. It is worth noting that two of the three natural charge assignments lead to a drastic suppression of M_{LR}.

Given our assumption that flavor symmetry breaking originates only from the vevs of the four flavons ϕ_i, it is possible to construct the $(S_3)^3$ representations shown above from products of the flavons. Thus, we will obtain the entries of the neutrino mass matrices in terms of products of the quark Yukawa couplings. A useful way to display our result is to express the Yukawa couplings and the third generation mixing angles in terms of powers of the Cabibbo angle: $h_u \sim \lambda^8$, $h_c \sim \lambda^4$, $h_t \sim 1$, $h_d \sim \lambda^7$, $h_s \sim \lambda^5$, $h_b \sim \lambda^3$, $V_{cb} \sim \lambda^2$, and $V_{ub} \sim \lambda^3$. We then obtain:

$$M_{RR} \sim c_0\langle\rho\rangle \begin{pmatrix} 1+c_1\lambda^6 & c_1\lambda^{10} & c_2\lambda^{10} \\ c_1\lambda^{10} & 1-c_1\lambda^6 & c_2\lambda^6 \\ c_2\lambda^{10} & c_2\lambda^6 & c_3 \end{pmatrix} \tag{2.10}$$

$$M_{LR} \sim \langle H_u\rangle\lambda^{10} \begin{pmatrix} d_1\lambda & d_1\lambda^5 & d_3\lambda \\ -d_1 & d_1\lambda^3 & -d_3\lambda^2 \\ d_2 & -d_2\lambda^3 & d_4\lambda^2 \end{pmatrix} \tag{2.11}$$

where ρ is the field whose vacuum expectation value determines the right-handed neutrino mass scale, and the c_i and d_i are order 1 coefficients. Notice that M_{LR} is suppressed by an overall factor of λ^{10}. This suppression has a simple interpretation. Had we chosen ν to transform as a $\mathbf{2}+\mathbf{1}_A$ under S_3^Q, we would expect flavor symmetry breaking to occur at the same order in the symmetry breaking as the down quark Yukawa matrix. However, with ν transforming under S_3^U, the matrix M_{LR} has a flavor symmetry structure that is completely different from all the other Yukawa matrices. Since the quark Yukawa couplings are the origin of flavor symmetry breaking in this model, we can only construct M_{LR} by going to *one higher order* in the symmetry breaking parameters, which, roughly speaking, is of the order of a typical Yukawa coupling squared. The precise value of the suppression is specific to the group theory of the given flavor model.

We may now compute the quantity of interest, the Majorana mass matrix for the light neutrino states. In the case of three flavors, the seesaw mechanism described earlier can be generalized:

$$M_{LL} = M_{LR}M_{RR}^{-1}(M_{LR})^T \ . \tag{2.12}$$

From Eqs. (2.10) and (2.11) above, we obtain

$$M_{LL} = \frac{\langle H_u\rangle^2 \lambda^{20}}{c_0\langle\rho\rangle}\times$$
$$\begin{pmatrix} (d_1^2 + d_3^2/c_3)\lambda^2 & -d_1^2\lambda - (d_3^2/c_3)\lambda^3 & d_1 d_2\lambda + (d_3 d_4/c_3)\lambda^3 \\ -d_1^2\lambda - (d_3^2/c_3)\lambda^3 & d_1^2 + (d_3^2/c_3)\lambda^4 & -d_1 d_2 - (d_3 d_4/c_3)\lambda^4 \\ d_1 d_2\lambda + (d_3 d_4/c_3)\lambda^3 & -d_1 d_2 - (d_3 d_4/c_3)\lambda^4 & d_2^2 + (d_4^2/c_3)\lambda^4 \end{pmatrix} \tag{2.13}$$

where we have retained higher order terms that lift a zero eigenvalue present at leading order. Notice that if $\langle\rho\rangle \approx 1$ TeV, $\lambda \approx 0.22$ and $\langle H_u\rangle \approx 175$ GeV, the overall scale of this matrix is of order 2×10^{-3} eV. This is approximately the correct magnitude to obtain the MSW solution to the solar neutrino problem, as we will see below. Given the very high power in λ, the predicted mass range can easily vary within an order of magnitude.

Perhaps the easiest way to study the physical implications of Eq. (2.13) is to rotate to a new basis $\nu_L' = U\nu_L$, where

$$U = \begin{pmatrix} 1 & 0 & 0 \\ 0 & -d_1/n & d_2/n \\ 0 & d_2/n & d_1/n \end{pmatrix} \tag{2.14}$$

with $n = (d_1^2 + d_2^2)^{1/2}$. Then M_{LL} becomes

$$M'_{LL} = \frac{\langle H_u\rangle^2 \lambda^{20}}{c_0\langle\rho\rangle}\times$$
$$\begin{pmatrix} (d_1^2 + d_3^2/c_3)\lambda^2 & d_1 n\lambda & (d_3/c_3)(d_1 d_4 - d_2 d_3)\lambda^3/n \\ d_1 n\lambda & n^2 & \mathcal{O}(\lambda^4) \\ (d_3/c_3)(d_1 d_4 - d_2 d_3)\lambda^3/n & \mathcal{O}(\lambda^4) & \mathcal{O}(\lambda^4) \end{pmatrix}. \tag{2.15}$$

We have retained order one coefficients only where they are relevant to our estimates below. Note that we principally will be interested in the 1-2 mixing, since Δm_{12}^2 and $\sin^2 2\theta_{12}$ are in the appropriate range for the the MSW solution. However, we will check that the 1-3 mixing (which corresponds to a much smaller Δm_{13}^2) does not lead to an unacceptable depletion of electron neutrinos that would be observable in terrestrial solar neutrino experiments**.

**The ν_μ-ν_τ mixing is not observable experimentally, given the small value of Δm^2.

Let us consider the neutrino mixing quantitatively. The 1-2 mixing angle from Eq. (2.15) is given by

$$\sin^2 2\theta_{12} = \frac{4d_1^2(d_1^2+d_2^2)\lambda^2}{(d_1^2+d_2^2-(d_1^2+d_3^2/c_3)\lambda^2)^2+4d_1^2(d_1^2+d_2^2)\lambda^2}. \tag{2.16}$$

If we choose the order 1 coefficients $d_1 = 0.4$, $d_2 = 2$, $d_3 = 1$, $c_0 = 2$, $c_3 = 1$, and we set $\langle\rho\rangle = 2$ TeV, we find

$$\sin^2 2\theta_{12} = 7.6\times 10^{-3} , \tag{2.17}$$
$$\Delta m_{12}^2 = 5.1\times 10^{-6}\ \text{eV}^2 . \tag{2.18}$$

This is consistent with the preferred range of the small angle MSW solution to the solar neutrino problem [11]

$$\sin^2 2\theta = 3\times 10^{-3}\text{–}1.1\times 10^{-2} , \tag{2.19}$$
$$\Delta m^2 = 3\times 10^{-6}\text{–}1\times 10^{-5}\ \text{eV}^2 . \tag{2.20}$$

Note that all the order 1 coefficients in this example are within the range 0.4–2, and the solution involved no fine-tuning. For this parameter set, $\sin^2 2\theta_{13} \sim 0.1$ and $\Delta m_{13}^2 = 7\times 10^{-10}$ eV^2. With Δm_{13}^2 this small, the 1-3 mixing does not lead to a significant depletion of the electron neutrino flux, and can be ignored.

It is also possible to achieve the large angle MSW solution, though in this case some fine-tuning is involved. If we choose $d_1 = 1.1$, $d_2 = 0.4$, $d_3 = 2.5$, $c_0 = 2/3$, $c_3 = 0.4$ and $\langle\rho\rangle = 1$ TeV, we obtain

$$\sin^2 2\theta_{12} = 0.5 , \tag{2.21}$$
$$\Delta m_{12}^2 = 1.8\times 10^{-5}\ \text{eV}^2 , \tag{2.22}$$

which is consistent with the preferred range for the large angle MSW solution [11]

$$\sin^2 2\theta = 0.5\text{–}0.9 , \tag{2.23}$$
$$\Delta m^2 = (1\times 10^{-5}\text{–}1\times 10^{-4})\ \text{eV}^2 . \tag{2.24}$$

In this parameter set, the order 1 coefficients fall within the range 0.4–2.5. Unfortunately, $\Delta m_{13}^2 = 5\times 10^{-6}$ implies that there would be significant depletion of electron neutrinos observed in the ^{71}Ga experiments, unless $\sin^2 2\theta_{13} \lesssim 10^{-4}$ [2]. This can be achieved in our model, providing we tolerate a 7% fine-tuning, $d_1 d_4 - d_2 d_3 \sim 0.07$.

Thus, it seems that the small angle MSW solution arises more naturally in our model. This is encouraging given that the small angle solution provides a better fit to the current data than the large angle one [11]. It is worth pointing out that the superKamiokande and SNO experiments are likely to see a distortion in the electron energy spectrum if the small angle solution is correct. Such a distortion would be an unambiguous indication of neutrino oscillation since it does not rely on the normalization of the solar neutrino flux, which is extremely sensitive to the core temperature of the Sun, scaling as $\sim T^{18}$.

3 Right-handed Neutrino Scale

In the example above, we saw that a completely reasonable theory of neutrino masses and mixings could be obtained with neutrinos in the 10^{-3} eV range, even when the right-handed scale $\langle \rho \rangle \approx 1$ TeV. In this section, we show how $\langle \rho \rangle$ can naturally be slightly larger than the electroweak scale in a radiative breaking scenario. A new gauge boson becomes massive when the ρ field acquires a vev, and could easily lie just beyond the current limits set from direct searches at the Tevatron.

We will assume an additional U(1) gauge symmetry, under which the right-handed neutrinos have charge +1. The simplest choice for the purpose of illustration is $U(1)_{B-L}$, since only right-handed neutrinos are required to render this symmetry nonanomalous in the MSSM. In addition, this extension of the MSSM preserves unification of the ordinary gauge coupling constants††. We also assume the presence of a pair of exotic Higgs fields with equal and opposite B-L charges, ρ_{+2} and ρ_{-2}. The superpotential that we will consider is

$$W = \alpha S \rho_{+2} \rho_{-2} - \frac{\beta}{3} S^3 + \gamma S H_u H_d + \frac{1}{2} Y_{RR} \rho_{-2} \nu\nu + Y_{LR} L H_u \nu \qquad (3.25)$$

††The particular flavor symmetry group, $(S_3)^3$, is difficult to implement in a conventional grand unified model. However, string unification could be a viable option in this scenario.

where S is a gauge singlet, and α, β and γ are coupling constants. The matrices M_{RR} and M_{LR} described earlier correspond to $Y_{RR}\langle\rho_{-2}\rangle$ and $Y_{LR}\langle H_u\rangle$, respectively.

The potential of this theory is given by

$$\begin{aligned} V &= |\alpha\rho_{+2}\rho_{-2} - \beta S^2|^2 + |\alpha S\rho_{-2}|^2 + |\alpha S\rho_{+2}|^2 + 2g^2(|\rho_{+2}|^2 - |\rho_{-2}|^2)^2 \\ &\quad + m_+^2|\rho_{+2}|^2 + m_-^2|\rho_{-2}|^2 + m_s^2|S|^2 - (A_\alpha\alpha S\rho_{+2}\rho_{-2} - \frac{1}{3}A_\beta\beta S^3 + \text{h.c}) \end{aligned} \tag{3.26}$$

where g is the $U(1)_{B-L}$ gauge coupling, and (m_+^2, m_-^2, m_s^2, A_α, A_β) are soft supersymmetry breaking masses and trilinear couplings. In a radiative breaking scenario, the Yukawa couplings Y_{RR} drive the soft squared mass m_-^2 negative in the renormalization group running, so we will look for a minimum of this potential assuming that $m_-^2 < 0$, with the remaining mass squared parameters positive. For $\alpha = \beta = \gamma = 0.1$ (to be explained below), $g = 0.3$, and the following dimensionful input parameters (in units of GeV)

A_α	A_β	m_+^2	m_-^2	m_S^2
200	100	$(256)^2$	$-(274)^2$	$(191)^2$

we find the vacuum expectation values

$\langle\rho_{+2}\rangle$	$\langle\rho_{-2}\rangle$	$\langle S\rangle$
1.3 TeV	1.4 TeV	0.7 TeV

and the mass squared eigenvalues

M_1^s	M_2^s	M_3^s
$(122\text{ GeV})^2$	$(322\text{ GeV})^2$	$(1655\text{ GeV})^2$
M_1^p	M_2^p	M_3^p
$(191\text{ GeV})^2$	$(332\text{ GeV})^2$	0

where the 0 eigenvalue corresponds to the degree of freedom that is 'eaten' by the U(1) gauge boson. Note that this solution has the desired feature that $m_-^2 < 0$ while m_+^2 and m_S^2 remain positive. The positivity of the scalar (s) and pseudoscalar (p) mass eigenvalues indicates that we have found a local minimum of the potential, which will suffice for our purposes. Note that $\langle\rho_{-2}\rangle$ is at 2 TeV, giving us the desired right-handed neutrino scale. Furthermore, with $\langle S\rangle$ at 1 TeV, and the auxiliary component $F_s = \alpha\rho_{+2}\rho_{-2} - \beta S^2$ at

(374 GeV)^2, we generate a μ parameter of order 100 GeV, and $B\mu$ of order (118 GeV)^2; this is in the proper range for electroweak symmetry breaking. Finally, we see that the Z' gauge boson develops a mass of 1.7 TeV, which is within the expected reach of the LHC, as we describe below.

The choice of couplings $\alpha = \beta = \gamma = 0.1$ was convenient for the purposes of illustration, since it allowed for a separation between the B-L and electroweak scales, so that the Higgs doublet vevs could be neglected in the minimization of Eq. (3.26). We could have chosen these parameters to be of order one if we had allowed smaller ρ vacuum expectation values, and analysed the full potential without approximation. This is beyond the scope of the present work. However, it is worth pointing out that smaller ρ vevs are not immediately excluded by the current bounds on the Z' boson mass. With our choice $g = 0.3$ (which in the given normalization is the approximate value one would expect from gauge coupling unification), the strongest constraint on the Z' mass comes from direct collider searches in the dilepton channel, yielding $m_{Z'} < 425$ GeV [12]; this implies $\langle\rho\rangle > 500$ GeV, assuming equal vevs for ρ_{+2} and ρ_{-2}. Thus, it seems there is no reason why the Z' in our model could not have a mass that is within the discovery reach of the Tevatron after the main injector upgrade, $m_{Z'} \lesssim 800$ GeV [13]. The corresponding reach at the LHC, $m_{Z'} \lesssim 4.3$ TeV [13], suggests that our model would also be testable in the limit $\langle\rho\rangle \gg \langle H\rangle$, that we obtained earlier by chooising $\alpha = \beta = \gamma = 0.1$. Such a parameter set can be justified by a symmetry argument, so it is not necessarily just a convenient limit. For example, we may imagine extending the flavor symmetry to the nonanomalous, discrete gauge group $(S_3)^3 \times Z_2$, where the Z_2 factor acts on the field S which is odd. In addition, we may assume there is a new flavon field ζ that is also Z_2 odd, and $\langle\zeta\rangle/M_{PL} \approx 1/10$. A spurion analysis similar to the one presented earlier would then suggest that all couplings involving an odd number of S fields should naturally be suppressed by 1/10. This would yield an effective theory in which α, β and γ are all of the same order, and are all relatively small.

In the scenario that we have described, there are two other issues that deserve comment. First, we have not justified the purely trilinear form of the superpotential in Eq. (3.25). Second, if B-L is gauged down to low energies, it is not immediately clear how we can generate and preserve a cosmic baryon asymmetry, taking into account the electroweak sphaleron effect[‡‡]. The answers to these questions are somewhat more speculative than the main points of this letter, but we include them for completeness. On the first issue, one might imagine the relevant fields are simply a part of the massless spectrum of string theory and the superpotential in Eq. (3.25) gives the complete set of renormalizable interactions. Alternatively, we could impose a global Z_3 or R-symmetry that restricts all renormalizable interactions to be trilinear in form. However, since global symmetries are thought to be violated by Planck-scale effects, this symmetry may not apply to the Planck-suppressed operators from which flavor symmetry breaking originates. On the second issue, we may imagine that baryogenesis proceeds via the Affleck-Dine mechanism [14] with large enough baryon to entropy ratio n_B/s so that the Bose condensate of the squarks and sleptons make the sphalerons heavy enough to preserve the B and L asymmetries even when $B - L = 0$ [15]. For instance, consider the operator

$$\frac{h_c}{M_*^2} Q_1 \overline{U}_1^* L_2 \overline{D}_2^* \tag{3.27}$$

where the charm Yukawa coupling originates from the usual spurion analysis, M_* is the reduced Planck mass, and the subscript indicates the generation. This operator respects all the unbroken symmetries in the low-energy effective theory, including $U(1)_{B-L}$. The estimate of the final baryon-to-entropy ratio that it generates is given by

$$\frac{n_B}{s} \sim g_*^{-1/4} h_c \left(\frac{M_*}{m}\right)^{1/2} \left(\frac{\phi_0}{M_*}\right)^4, \tag{3.28}$$

where m is a typical supersymmetry-breaking scalar mass, g_* is the effective number of degrees of freedom in the Early Universe, and ϕ_0 is the initial amplitude of the scalar fields, when they begin to oscillate. In order to have sufficient suppression of the sphaleron effects, we need $n_B/s \gtrsim 10^{-2}$ which translates into the bound $\phi_0 \gtrsim 10^{16}$ GeV. Such a large initial amplitude may be generated if there is a scalar field with a negative mass squared during inflation, in a model with a non-minimal Kähler potential coupling to the

[‡‡] Note that we could have chosen a different gauged U(1) group, for example from $E_6 \to SO(10) \times U(1)$ breaking, so that the right-handed neutrinos are charged but $B - L$ is not a gauge symmetry. In this case, the constraints from baryogenesis are significantly weaker.

inflaton [16] or in no-scale supergravity [17]. Subsequent entropy production is then needed to dilute n_B/s down to the value required by nucleosynthesis; this possibly may be provided by the decay of a Polonyi-like field or by a late inflation with an e-folding of order 7, like that in Ref. [18].

4 Conclusions

We have shown that flavor symmetries allow for a class of models in which neutrino masses can be made extremely small, without requiring the mass scale of right-handed neutrinos to be larger than the electroweak scale. While our model was based on a discrete, non-Abelian family symmetry, the general idea should also be applicable to a large number of other flavor models, in particular those involving Abelian horizontal symmetries. The scenario we have outlined supports the notion that there are a wide class of models where right-handed neutrino Yukawa couplings may lead to a radiative breaking of additional U(1) gauge groups, leading naturally to Z' bosons with masses less than a few TeV. With the LHC expected to discover a Z' bosons with standard model couplings up to about 5 TeV [13], this scenario can be tested definitively in the future.

Acknowledgments

This work was supported in part by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U.S. Department of Energy under Contract DE-AC03-76SF00098. HM was also supported in part by the National Science Foundation under grant PHY-95-14797, and by the Alfred P. Sloan Foundation.

A Another Model

For completeness, we present the mass matrix M_{LL} that we obtain when the right-handed neutrinos ν transform as a $\mathbf{2} + \mathbf{1}_A$ under S_3^D rather than S_3^U. We obtain

$$M_{LL} \sim \frac{\langle H_u \rangle^2 \lambda^{16}}{c_0 \langle \rho \rangle} \begin{pmatrix} d_1^2 + d_3^2/c_3 & -d_3^2/c_3\lambda & -2d_1d_2\lambda \\ -d_3^2/c_3\lambda & d_1^2 & d_1d_2 \\ -2d_1d_2\lambda & d_1d_2 & d_2^2 \end{pmatrix}, \tag{A.1}$$

where the coefficients c_i and d_i multiply the same entries of M_{RR} and M_{LR} as in the example presented in the text. In this case, if we choose $\langle\rho\rangle \approx 1$ TeV, and $\langle H_u \rangle \approx 175$ GeV, the overall scale of this matrix is of order 1 eV. Again, the flavor symmetry structure of the model is responsible for a large suppression of the Dirac mass matrix, of order λ^8.

Notice from the determinant of Eq. (A.1) that two of the eigenvalues are of order 1 eV, while the remaining one is roughly λ^2 smaller. Thus, the corresponding Δm^2 values do not suggest a natural solution to either the solar, or atmospheric neutrino problems. The 1 eV mass scale is appropriate for the neutrinos to be hot dark matter candidates. However, with masses in this range the large mixing between the second and third generations is problematic in light of the bounds from disappearance experiments that search for $\nu_\mu \to \nu_x$. These searches completely exclude $\sin^2 2\theta_{23} > 0.1$ for $\Delta m^2 = 1$ eV^2, and place significantly stronger bounds for Δm^2 between 1–100 eV^2 [19]. Thus, the hot dark matter solution seems possible only for very small $d_1/d_2 < 0.15$ with d_2 of order 1, implying a rather unnatural value for d_1. For a somewhat smaller overall scale, (choosing $c_0 = 2$ and $\langle\rho\rangle = 3$ TeV, for example), we can obtain an acceptable theory, which explains no neutrino anomalies, but implies ν_μ-ν_τ mixing at a level that is likely to be measurable at the CHORUS or NOMAD experiments.

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Prepared for the U.S. Department of Energy under Contract No. DE-AC03-76SF00098