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Fermion masses and proton decay in string-inspired $SU(4) \times SU(2)^2 \times U(1)_X$

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Abstract

We present a supersymmetric model of fermion masses with $SU(4) \times SU(2)^2 \times U(1)_X$ gauge group with matter in fundamental and antisymmetric tensor representations only. The up, down, charged lepton and neutrino Yukawa matrices are distinguished by different Clebsch–Gordan coefficients due to contracting over SU(4) and $SU(2)_R$ indices. We obtain a hierarchical light neutrino mass spectrum with bi-large mixing. The condition that anomalies be cancelled by a Green–Schwarz mechanism leads to fractional $U(1)_X$ charges which exclude B violation through dimension-4 and -5 operators.

The values of the adjustable parameters of the Standard Model (SM) Lagrangian may be an important clue to physics beyond it (BSM). Thus, the experimentally measured values of gauge coupling constants, fermion masses and quark mixing angles, and now neutrino mass-squared differences and mixing angles (strictly, already a BSM effect), can be compared to the predictions of various types of model with full or partial [1] gauge unification and/or flavour symmetries. Recent data on atmospheric and solar neutrinos [2], implying large 1–2 and 2–3 mixing angles, present challenges for any unified framework in which neutrinos form part of a multiplet with quarks [3].

In this Letter we revisit the string-inspired 4–2–2 models [4] (see also [5]), whose implications for fermion masses were previously investigated in [6,7] and which have several attractive features. Large Higgs representations (problematic to obtain in string models) are not required, the doublet–triplet splitting problem is absent, third generation fermion Yukawa couplings are unified [8] up to small corrections, and unification of gauge couplings is allowed and, if one assumes the model embedded in supersymmetric string, might be predicted [9].

Small effective Yukawa couplings arise from nonrenormalizable superpotential operators involving a singlet charged under $U(1)_X$ [10] and Higgses which receive vevs at the $SU(4) \times SU(2)_R$ breaking scale. A particular feature of the model is the presence of two a priori independent expansion parameters depending on the sign

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of the $U(1)_X$ charge. For nonrenormalizable operators involving products of $SU(4) \times SU(2)_R$ breaking Higgses, different contractions of gauge group indices lead to effective Yukawa couplings which differ between the up (u) and down (d) quarks and the charged leptons (e) and neutrinos (ν). This freedom is exploited to fit the different mass hierarchies in the u, d and e sectors and produce the CKM mixing.

Right-handed neutrinos (RHNs) are automatically present and obtain Majorana masses, again through non-renormalizable couplings to the $U(1)_X$-charged singlet and to Higgses. Their mass spectrum is hierarchical with the lightest RHN around 10^{11} GeV and the heaviest just below the Pati–Salam breaking scale. The RHN mass hierarchy cancels off against the hierarchical neutrino Dirac mass matrix to yield a seesaw mass matrix with large off-diagonal elements. The Green–Schwarz anomaly cancellation conditions can always be satisfied, and imply fractional $U(1)_X$ charges which disallow many B- and L-violating operators.

1. The model

The field content is summarized in Table 1, where i ranges from 1 to 3. The two D fields are introduced to give mass to colour triplet components of H and $\bar{H}$ once their sneutrino-like components obtain vevs as follows:

$$\langle H\rangle = \langle H_\nu\rangle = M_{\rm G}, \qquad \langle \bar{H}\rangle = \langle \bar{H}_\nu\rangle = M_{\rm G}.$$

In the stable SUSY vacuum the singlet ϕ obtains a vev to satisfy the anomalous D-flatness condition. This is a natural mechanism in the context of string models which results to the spontaneous breaking of the $U(1)_X$ about an order of magnitude below the string scale. In fact, the part of the effective Lagrangian cancelling the anomalies has a supersymmetric counterpart which is a sort of field-dependent Fayed–Iliopoulos term [11] and leads to $U(1)_X$ breaking. Thus, the string scenario provides a natural explanation why the extra $U(1)_X$ symmetry required to generate the fermion mass patterns does not survive down to low energies [11–13].

In general a string model may have more than one singlet ϕ_i and more than one set of Higgses H_i, $\bar{H}_i$, with different $U(1)_X$ charge; the Higgses may obtain masses through $H\bar{H}\phi$ couplings. In order to break the Pati–Salam group while preserving SUSY we require that one H–$\bar{H}$ pair be massless at this level. This "symmetry-breaking" Higgs pair may be a linear combination of many H_i and $\bar{H}_i$ in the charge basis, which will in general complicate the expressions for fermion masses (as will the presence of many ϕ_i fields).

However, if we impose that all products $H_i\bar{H}_j$ have the same sign of $U(1)_X$ charge, and that all ϕ_i likewise have the same sign of charge (opposite to that of $H\bar{H}$), then the leading contributions to fermion masses from nonrenormalizable operators arise from the $H\bar{H}$ pair and ϕ field which have the smallest absolute value of $U(1)_X$ charge; other vevs will enter at higher order and will be small corrections. Hence, and for simplicity, we restrict ourselves to a single copy of H, $\bar{H}$ and ϕ.

Table 1
Field content and $U(1)_X$ charges

	SU(4)	$SU(2)_L$	$SU(2)_R$	$U(1)_X$
F_i	4	2	1	α_i
$\bar{F}_i$	$\bar{4}$	1	$\bar{2}$	$\bar{\alpha}_i$
H	4	1	2	x
$\bar{H}$	$\bar{4}$	1	$\bar{2}$	$\bar{x}$
ϕ	1	1	1	z
h	1	$\bar{2}$	2	$-\alpha_3 - \bar{\alpha}_3$
D_1	6	1	1	$-2x$
D_2	6	1	1	$-2\bar{x}$

Standard Model fermion mass terms arise after electroweak symmetry-breaking from the operators

$$W_D = y_0^{33} F_3 \bar{F}_3 h + \delta \sum_{m>0} y_m^{ij} F_i \bar{F}_j h \left(\frac{\phi}{M_S}\right)^m + \delta \sum_{n>0} y_n'^{ij} F_i \bar{F}_j h \left(\frac{H\bar{H}}{M_S^2}\right)^n + \cdots, \tag{1}$$

where M_S is the "string scale" which governs the suppression of nonrenormalizable terms in the effective theory, and only the 33 element is nonvanishing at renormalizable level. Nonrenormalizable terms may be suppressed by an overall factor δ of order 1. The couplings y_n^{ij} and $y_n'^{ij}$ are nonvanishing and generically of order 1 whenever the $\mathrm{U}(1)_X$ charge of the corresponding operator vanishes. Other higher-dimension operators may arise by multiplying any term by $(H\bar{H})^a \phi^b / M_S^{2a+b}$ for positive integer a and b such that $a(x+\bar{x}) + bz = 0$; however such terms are negligible unless the leading term vanishes.

Majorana masses arise from the operators

$$W_M = \frac{\bar{F}_i \bar{F}_j H H}{M_S}\left(\mu_0^{ij} + \sum_{p>0} \mu_p^{ij} \left(\frac{\phi}{M_S}\right)^p + \sum_{q>0} \mu_q'^{ij} \left(\frac{H\bar{H}}{M_S^2}\right)^q\right). \tag{2}$$

Again, the $\mathrm{U}(1)_X$ charges $\bar{\alpha}_i + \bar{\alpha}_j + 2x + pz$ or $\bar{\alpha}_i + \bar{\alpha}_j + 2x + q(x+\bar{x})$ must be zero if the couplings μ_p^{ij} and μ'^{ij}_q, respectively, are not to vanish.

For nonrenormalizable Dirac mass terms involving n products $H\bar{H}/M_S^2$ the gauge group indices may be contracted in different ways [7] leading to Clebsch factors $C^{ij}_{n(u,d,e,\nu)}$ multiplying the effective Yukawa coupling: these are generically numbers of order 1 and may be zero in some cases. Although the Clebsch coefficient for a particular operator O_n may vanish at order n, the coefficient for the operator $O_{(n+a);b}$ containing a additional factors $(H\bar{H})$ and b factors of ϕ is generically nonzero.

Dirac mass terms at the unification scale are then

$$\frac{m^{ij}}{m_3} = \delta_{i3}\delta_{j3} + \delta \frac{y_n^{ij}}{y_0^{33}} \left(\frac{\langle\phi\rangle}{M_S}\right)^m + \delta \frac{y'^{ij}_n}{y_0^{33}} C_n^{ij} \left(\frac{M_{\mathrm{G}}^2}{M_S^2}\right)^n \simeq \delta_{i3}\delta_{j3} + \delta\left(\epsilon^{|\hat{m}|} + C^{ij} \epsilon'^{|\hat{n}|}\right), \tag{3}$$

where $m_3 \equiv v_{u,d} y_0^{33}$ with v_u and v_d being the up-type and down-type Higgs vevs respectively, and we define

$$\epsilon \equiv \left(\frac{\langle\phi\rangle}{M_S}\right)^{|1/z|}, \qquad \epsilon' \equiv \left(\frac{M_{\mathrm{G}}^2}{M_S^2}\right)^{1/|x+\bar{x}|}.$$

We suppress higher-order terms involving products $\epsilon\epsilon'$. Thus $\hat{m} = -(\alpha_i - \alpha_3 + \bar{\alpha}_j - \bar{\alpha}_3)$ and $\hat{n} = -(\alpha_i - \alpha_3 + \bar{\alpha}_j - \bar{\alpha}_3)$. The signs of $\hat{m}$ and $\hat{n}$ must be the same as z and $x+\bar{x}$, respectively, for the mass term to exist. Since the integers m and n are always positive, we have $\epsilon^{|\hat{m}|} \equiv (\langle\phi\rangle/M_S)^m$. Majorana mass terms are

$$M_M^{ij} = M_R\left(\mu_p^{ij} \epsilon^{|\hat{p}|} + \mu_q'^{ij} \epsilon'^{|\hat{q}|}\right) \simeq M_R\left(\epsilon^{|\hat{p}|} + \epsilon'^{|\hat{q}|}\right), \tag{4}$$

where $M_R \equiv M_{\mathrm{G}}^2/M_S \equiv \epsilon' M_S$, the term with $p = q = 0$ is understood, and $\hat{p} = -(\bar{\alpha}_i + \bar{\alpha}_j + 2x)$ and $\hat{q} = -(\bar{\alpha}_i + \bar{\alpha}_j + 2x)$. The full neutrino mass matrix in the basis (ν, ν^c) is of the "seesaw" form and the resulting light neutrino mass matrix is

$$m_\nu = -m_{D\nu} M_M^{-1} m_{D\nu}^T. \tag{5}$$

2. Parameter choices and mass matrices

So far the discussion has been independent of the choice of $\mathrm{U}(1)_X$ charges. The fermion mass terms are invariant under the family-independent shifts of $\mathrm{U}(1)_X$ charge

$$\alpha_i \to \alpha_i + \zeta, \qquad \bar{\alpha}_i \to \bar{\alpha}_i + \bar{\zeta},$$

$$x \to x - \bar{\zeta}, \qquad \bar{x} \to \bar{x} + \bar{\zeta} \tag{6}$$

(as are the HHD_1 and $\bar{H}\bar{H}D_2$ couplings). Thus for the purpose of investigating fermion masses, we are free to assign $\alpha_3 = \bar{\alpha}_3 = 0$ and apply ζ and $\bar{\zeta}$ shifts at the end of the calculation. The charges are then $x + \bar{x} = 1$, $z = -1$, $\alpha_1 = -4$, $\bar{\alpha}_1 = -2$, $\alpha_2 = -3$, $\bar{\alpha}_2 = 1$, which generate the charge matrices

$$Q_X[M_D] = \begin{pmatrix} -6 & -3 & -4 \\ -5 & -2 & -3 \\ -2 & 1 & 0 \end{pmatrix}, \qquad Q_X[M_M] = 2x + \begin{pmatrix} -4 & -1 & -2 \\ -1 & 2 & 1 \\ -2 & 1 & 0 \end{pmatrix}. \tag{7}$$

We set $\delta = 1$ and substitute ϵ and ϵ' by a single expansion parameter $\eta \simeq 0.06$ (or $\sqrt{\eta} \simeq 0.24$) via $\epsilon = \eta$, $\epsilon' = \sqrt{\eta}$. Thus the neutral gauge singlet $H\bar{H}\phi/M_S^3$ has a vev of $\eta^{3/2} \simeq 0.015$.

We obtain a hierarchical common Dirac mass matrix and (x-dependent) Majorana RHN mass matrix in powers of η. We then have to specify Clebsch coefficients for operators involving one or more powers of $H\bar{H}/M_S^2$. By taking linear combinations of operators with different contractions over $\mathrm{SU}(4) \times \mathrm{SU}(2)_R$ indices one can obtain any vector in the space (u, d, ν, e), since the operators constitute a complete set over this space. Thus at first sight the model has little prodictivity.

However, in specific string models these coefficients are calculable in terms of cocycle factors [14]; in the absence of a specific string construction we impose that the C_n^{ij} should be either small integers or simple rational numbers. We can take all C^{ij} equal to unity apart from the following, where we quote coefficients up to a possible complex phase: $C_d^{12} = C_d^{22} = 1/3$, $C_u^{23} = 3$, $C_u^{11} = C_u^{12} = C_u^{21} = C_u^{22} = C_u^{31} = 0$, all multiplying the leading nonvanishing entry. The ratio $C_e^{22}/C_d^{22} = 3$ is the usual Georgi–Jarlskog choice to fit the different ratios m_s/m_b and m_μ/m_τ. As explained above, when C_n^{ij} vanishes for the leading term, the next-to leading term is smaller by a factor $\eta^{3/2}$. Hence the Dirac mass matrices at the GUT scale take the following form, up to complex phases and factors of order 1:

$$\frac{m_u}{m_{t0}} = \begin{pmatrix} \eta^{9/2} & \eta^3 & \eta^2 \\ \eta^4 & \eta^{5/2} & 3\eta^{3/2} \\ \eta^{5/2} & \eta & 1 \end{pmatrix}, \qquad \frac{m_d}{m_{b0}} = \begin{pmatrix} \eta^3 & \frac{\eta^{3/2}}{3} & \eta^2 \\ \eta^{5/2} & \frac{\eta}{3} & \eta^{3/2} \\ \eta & \eta & 1 \end{pmatrix}, \qquad \frac{m_e}{m_{\tau 0}} = \begin{pmatrix} \eta & \eta^{3/2} & \eta^2 \\ \eta^{5/2} & \eta & \eta^{3/2} \\ \eta & \eta & 1 \end{pmatrix}, \tag{8}$$

where m_{t0} is the top mass at the GUT scale and $m_{b0} = m_{\tau 0} = m_{t0}/\tan\beta$. For this simple choice, the resulting eigenvalues and quark mixings can be RG evolved to observable energies, where they yield acceptable fits. For example, the CKM mixing angle θ_{12} is $\sqrt{\eta} \simeq 0.24$, while the angle θ_{13} is $\eta^2 \simeq 0.0035$; the ratio $(m_u/m_t)_{|M_Z}$ is $\eta^{9/2}/\zeta^3 \simeq 6 \times 10^{-6}$, where $\zeta \simeq 0.83$ accounts for the RG evolution of $y_{u,c,t}$ in the large $\tan\beta$ (fixed point) regime. No fine-tuned cancellations between unknown order 1 coefficients are needed. Such coefficients arise from the couplings y_m^{ij}, etc., which are SU(4) symmetric, and thus affect the Dirac mass matrix in the same way in each sector. This constraint may have consequences for a more detailed comparison with data.

The light neutrino mass matrix depends on x through the Majorana RHN matrix: we find two possible values[1]

$$x = 1 \quad \text{and} \quad x = \frac{3}{2}. \tag{9}$$

The seesaw mass matrix is then

$$m_\nu = \frac{(m_\tau \tan\beta)^2}{\eta^{2x} M_R} \begin{pmatrix} \eta & \sqrt{\eta} & \sqrt{\eta} \\ \sqrt{\eta} & 1 & 1 \\ \sqrt{\eta} & 1 & 1 \end{pmatrix} \tag{10}$$

1 For $x < 1$, m_ν no longer has a form consistent with bi-large mixing; if $x > 3/2$ then the RHNs are too light and the light ν masses too large.

up to order 1 factors (different from those in the Dirac matrices), where we display leading terms in η. With a mildly fine-tuned (at the 20% level) choice of order 1 coefficients one can obtain a neutrino mass spectrum with normal hierarchy and bi-large mixing. We did not find any charge assignments consistent with an inverted hierarchy or degenerate light neutrino spectrum.

If we take $x = 1$, $\tan\beta = 40$ and $M_R \equiv \eta^{1/4} M_{\rm G}$ to be 1.2×10^{16} GeV, then the largest entries in m_ν are of order 0.12 eV. The spectrum of RHNs comprises two superheavy states of mass $(1 \pm \eta/2)M_R$ and one light state of mass $\eta^4 M_R \simeq 1.5 \times 10^{11}$ GeV.

Alternatively, with $x = 3/2$, $\tan\beta = 45$ and M_R being $\sqrt{\eta}$ times the reduced Planck mass 2.4×10^{18} GeV, i.e., $M_R \simeq 6 \times 10^{17}$ GeV, the largest entries in m_ν are of order 0.05 eV. The RHN masses are now of order $\sqrt{\eta} M_R$, $\eta^{3/2} M_R$ and $\eta^5 M_R \simeq 5 \times 10^{11}$ GeV. Thus the correct scale of light ν masses follows, with RHN mass terms derived from either the SUSY-GUT or heterotic string scale via nonrenormalizable operators. The lightest RHN masses are rather large for standard thermal leptogenesis [15] if one takes into account gravitino production (given $m_{3/2} = 1$–10 TeV), but might be considered for nonthermal leptogenesis [16] or in case the gravitino is light or very heavy.

3. Anomalies and *B* and *L* violation

For gauge coupling unification at the string scale (up to threshold corrections) the non-Abelian gauge groups are required to have equal Kač–Moody levels $k_4 = k_{2L} = k_{2R} = 1$. The $U(1)_X$ mixed anomalies can only be cancelled by a Green–Schwarz mechanism if anomaly coefficients obey the relation $A = \text{const} \times k$, hence we require the 4–4–1_X, 2_L–2_L–1_X and 2_R–2_R–1_X coefficients to be equal: $A_4 = A_{2L} = A_{2R}$. We have (up to an overall factor)

$$A_4 = 2\sum_i(\alpha_i + \bar\alpha_i) + 2(x + \bar x) + 2(-2x - 2\bar x), \qquad A_{2L} = 4\sum_i \alpha_i + 2(-\alpha_3 - \bar\alpha_3),$$
$$A_{2R} = 4\sum_i \bar\alpha_i + 4(x + \bar x) + 2(-\alpha_3 - \bar\alpha_3), \tag{11}$$

from which we obtain the requirements

$$\alpha_3 + \bar\alpha_3 - 2(x + \bar x) = 0, \tag{12}$$
$$\sum_i (\alpha_i - \bar\alpha_i) - (x + \bar x) = 0. \tag{13}$$

The generation-independent shifts of Eq. (6) produce shifts in the L.H.S. of Eqs. (12) and (13), of $\zeta + \bar\zeta$ and $3(\zeta - \bar\zeta)$, respectively. Thus, given an initial $U(1)_X$ charge assignment, it is always possible to satisfy the anomaly conditions, without altering the fermion mass matrices.

For our chosen set of charges we have $\zeta = 13/6$ and $\bar\zeta = -1/6$. The shifts make many operators fractionally charged: the $FFFF$ operator which would lead to $D = 5$ proton decay has charge $\sum\alpha + 26/3$, which cannot be cancelled by any singlet combination of H, $\bar H$ and ϕ fields. The only surviving operators coupling matter to Higgs triplets are FFD_1 and $FF\bar H\bar H$, whose charges are both shifted by $2(\zeta + \bar\zeta) = 4$. For $D = 5$ proton decay one would also require appropriate mass terms for intermediate states, either D_1D_1 or $D_1\bar H\bar H$; but the charges of these operators are both shifted by $4\bar\zeta = -2/3$, hence one cannot construct the mass terms.

We must also verify that all Higgs triplet states are massive enough to evade bounds from $D = 6$ operators. Down squark-like states 3_H and $\bar 3_{\bar H}$ in H and $\bar H$ respectively survive after breaking to the SM group. Mass terms for the H–$\bar H$–D_1–D_2 system follow from the superpotential operators

$$W_{\rm h} = HHD_1 + \bar H\bar H D_2 + D_1 D_2 \frac{H\bar H H\bar H}{M_S^3}, \tag{14}$$

inserting vevs we obtain up to factors of order 1

$$W_h = (\bar{3}_1 \quad \bar{3}_{\bar{H}} \quad \bar{3}_2) \begin{pmatrix} M_G & \mu_D & 0 \\ 0 & M_G & 0 \\ 0 & 0 & \mu_D \end{pmatrix} \begin{pmatrix} 3_H \\ 3_2 \\ 3_1 \end{pmatrix}, \tag{15}$$

where $\mu_D = M_G^4/M_S^3 \equiv \eta^{3/4} M_G$ and we decompose each sextet $D_{1,2}$ into a **3** and $\bar{\mathbf{3}}$ of SU(3). There are two mass eigenvalues of approximately M_G and one eigenvalue μ_D from the combination $\bar{3}_2$–3_1. Clearly there is no transition between 3_1 and either $\bar{3}_1$ or $\bar{3}_H$, at any order.

Other B- and L-violating operators which might endanger the proton's longevity are

$$F_i \bar{H} h, \ \epsilon_{ab} F_{iA} \bar{F}_j^{Aa} F_{kB} \bar{H}^{Bb}, \ \bar{F}_i \bar{F}_j \bar{F}_k \bar{H}, \ F_i F_j F_k H h, \tag{16}$$

where we have indicated SU(4) and $SU(2)_R$ summations in the second term.[2] These give rise to superpotential terms LH_u; QD^cL and LE^cL; $U^cD^cD^c$; and $QQQH_d$, respectively. Out of these, $F\bar{H}h$ is not shifted, hence is allowed (suppressed by some powers of η) if $\bar{x}$ is integer; $F\bar{F}F\bar{H}$ is shifted by 4 and is allowed under the same condition; $\bar{F}^4$ is shifted by $-2/3$, hence is disallowed; and F^3Hh is shifted by $14/3$, hence also disallowed.

Regarding the R-parity violating LH_u term, this is allowed only in the case $x = 1$ and forbidden in the case $x = 3/2$. If it is allowed, R-parity violation in the lepton sector also has potentially dangerous effects. The bilinear operator $L_i H_2$ is known to contribute to neutrino masses [17] and there are experimental bounds on trilinear LLE couplings [18]. For $x = 1$, if there are no other discrete or continuous symmetries in the model, we expect the (mass dimension 1) coefficient of $L_i H_u$ to be a few orders of magnitude below the GUT scale, which is obviously unacceptable. Thus if no additional symmetries (such as R-parity) are present at the effective Lagrangian the case $x = 1$ seems to be excluded by these effects. However, in the case $x = 3/2$ potentially problematic L-violating terms are absent. In a more general model with many H–$\bar{H}$ pairs and singlets, with a range of $U(1)_X$ charges, the analysis of anomaly coefficients, proton decay operators and R-parity violation will be different. But at least in the minimal version, all dangerous B-violating operators vanish automatically.

The presence of R-parity violating operators at low energy implies an unstable LSP, which cannot be a dark matter particle (unless its lifetime is much longer than the age of the Universe). This may widen the available MSSM parameter space, since an unstable LSP cannot overclose the Universe, and a charged LSP is possible. The LSP lifetime and abundance should however satisfy other cosmological constraints [19]. The details of such a scenario likely depend on how the bilinear LH_u and mu-terms are suppressed, hence it is beyond the scope of this Letter to discuss the possibilities in detail. There are well-studied candidates in string models for dark matter beyond the LSP, including axions and superheavy hidden sector bound states [20].

The "mu-term" of the MSSM originates from the product hh whose $U(1)_X$ charge is -4: hence it receives only a mild suppression. Thus the "mu-problem" is not solved by the $U(1)_X$ symmetry alone. This is a rather generic problem in string model building [21]. Contrary to the fermions that are chiral and thus protected from direct mass terms the Higgs doublet mass is generically allowed. In realistic string models however, there are two additional mechanisms of mass term suppression. The first is additional Abelian symmetries that are generically present and the second is string related selection rules [22]. As an example for the latter one could refer to a configuration where one of the Higges arises from the untwisted sector while the singlet and the other Higgs come from the twisted sector. In this case the μ-term is forbidden in the effective superpotential. It can be however generated by the Kähler potential [23]. String selection rules and/or additional Abelian symmetries could also account for the elimination of the LH_u mixing term.

In conclusion, we have presented a string-inspired supersymmetric $SU(4) \times SU(2)^2 \times U(1)_X$ model with 5 discretely adjustable charges, 8 discretely adjusted Clebsch factors, 1 adjustable expansion parameter and 1 adjustable

[2] The conventional R-parity may be obtained from a $\mathbb{Z}_2$ symmetry acting either on F and $\bar{F}$, or on H and $\bar{H}$.

mass ratio (v/M_G), which is consistent with gauge unification and with all known elementary fermion masses and mixings, if the spectrum of light neutrinos is hierarchical. The simplest version of the model is also free from B violation through dimension-4 and -5 operators, but may allow L violation. The lightest RHN mass is a few times 10^{11} GeV, the lightest neutrino mass eigenstate is of order $\eta/2 \simeq 0.03$ times the heaviest, and the neutrino mixing angle θ_{13} is of order $\sqrt{\eta}/2 \simeq 0.12$. Other issues for further investigation include gauge unification, CP violation, supersymmetry-breaking and flavour-changing effects in both quark and lepton sectors, and cosmology including inflation, baryogenesis and dark matter.

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