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SHARPENING THE FLIPPED $SU(5)$ STRING MODEL

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ABSTRACT

We present an improved version of the flipped $SU(5)$ string model which accommodates numerous existing and expected features of the low-energy physics world, such as: (i) two Higgs doublets with adequate mixing between them; (ii) hierarchy of quark and lepton masses, including relations such as $m_s/m_b \sim m_\mu/m_\tau$ and $m_c/m_t \sim (m_e/m_\tau)^{1/2}$; (iii) sufficiently suppressed proton decay operators, with gauge-boson mediated decays favoring $p \rightarrow \bar{\nu}_\tau \pi^+$ and $n \rightarrow \bar{\nu}_\tau \pi^0$; (iv) nearly massless ν_μ and ν_τ , and ν_e in the eV range; (v) heavy top quark ($m_t \lesssim 170 \text{ GeV}$) and large ratio of vacuum expectation values ($\tan \beta \lesssim 33$).

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For many years our quest for a Theory of Everything (TOE), that is a theory that unifies electroweak, strong, and gravitational interactions in a self-consistent and falsifiable framework, has been based on intuition, imagination, and grand vision. This fact has sharply divided the physics community into “traditionalists” and “adventurers” raiding for the treasured TOE. Recently things have changed dramatically: a sharp eye’s view of the LEP ($\sin^2 \theta_w, \alpha_s$) data^{1,2} makes it apparent that the grand unification (GUT) program is working and in its best form, namely supersymmetric grand unification (SUSY GUTs).³ While no explicit supersymmetric particle has yet been seen, the LEP data implicitly encompass them all through the running of the gauge couplings via the renormalization group equations (RGEs). This is the second instance in which SUSY GUTs have shown their strength. The big LEP discovery of the existence of only three light neutrino species had been predicted by SUSY GUTs through the running of the Yukawa coupling RGEs and the fitting to the observed m_b/m_τ ratio.⁴ We should remark that while ordinary GUTs had already been ruled out based on unsuccessful searches for proton decay, the recent LEP data make this conclusion unescapable.^{1,2} We find these recent developments in SUSY GUTs to provide very strong hints to the real value of this general framework (and its natural extension to superstring unification). We are therefore strongly encouraged to continue in our endeavor of developing a realistic string-derived model with singularly unique properties, namely the flipped $SU(5)$ model.⁵

Even though there is to date not a shred of experimental evidence against the standard model, many extensions of it have been considered in an effort to understand its several arbitrary parameters. Ordinary GUT models³ have been quite successful in this respect, but are plagued by the notorious gauge hierarchy problem which renders them incalculable. It is well-known that (globally) supersymmetric models solve this problem most elegantly. However, supersymmetry must be softly broken at low energies;⁶ this phenomenon finds its origin in locally supersymmetric, *i.e.*, supergravity models,⁷ preferably of the no-scale type.⁸ The latter are to be understood as effective unified theories including gravity, although their predictability is limited to energies below the Planck mass. String theory⁹ then takes over providing the ultimate unification of all particle interactions. Clearly, a string-derived model automatically includes all the steps in the above conceptual chain.

In this letter we present an improved version of the flipped $SU(5)$ string model which ac-

commodates several existing and expected features of low-energy physics, such as: hierarchical fermion masses, two Higgs doublets with adequate mixing between them, neutrino masses, heavy top quark, etc. We cannot emphasize enough that all the features that this model possesses are *derived* from string theory in the so-called free fermionic formulation of the heterotic string in four dimensions.¹⁰ This computationally powerful¹¹ approach to string model-building constitutes one of many possible alternatives,¹⁰ which arise due to the very large perturbative vacuum degeneracy of string theory; nonperturbative string effects are expected to select the realized vacuum.

One of the motivations for studying flipped $SU(5)$ models¹² is evident at the field theoretical level.¹³ Since the particles in the $\bar{\mathbf{5}}, \mathbf{10}, \mathbf{1}$ representations are “flipped” (*i.e.*, $u \leftrightarrow d$, $e \leftrightarrow \nu$, $u^c \leftrightarrow d^c$, $e^c \leftrightarrow \nu^c$) relative to their traditional assignments, the $\mathbf{10}_h, \bar{\mathbf{10}}_h$ representations of Higgs bosons can break the gauge symmetry $SU(5) \times U(1)_{\tilde{Y}} \rightarrow SU(3)_C \times SU(2)_L \times U(1)_Y$ most economically, and they do this *uniquely*, as favored by cosmology.¹⁴ The presence of these Higgs representations also allows for a very nice solution to the doublet-triplet splitting problem of the Higgs pentaplets, through the coupling $\mathbf{10}_h \cdot \mathbf{10}_h \cdot \mathbf{5}_h \rightarrow \langle \mathbf{10}_h \rangle d_h^c D$, where $\mathbf{5}_h = \{D, H\}$, and H is one of the electroweak Higgs doublets. Furthermore, the coupling $\mathbf{10}_f \cdot \bar{\mathbf{10}}_h \cdot \phi \rightarrow \langle \bar{\mathbf{10}}_h \rangle \nu^c \phi$ provides a natural extended see-saw mechanism involving singlet fields ϕ . Finally, gauge coupling unification in flipped $SU(5)$ is fully compatible with the recent LEP data.¹ A more compelling fact in the string context is that string models whose gauge groups are represented by level-one Kac-Moody algebras do not have adjoint Higgs representations in their spectrum¹⁵ to break the unified gauge symmetry. This fact makes flipped $SU(5)$ string model-building possible, while it disallows the traditional grand unified groups such as $SU(5)$, $SO(10)$, and E_6 .¹⁶

The flipped $SU(5)$ string model⁵ has the following gauge symmetry group at the Planck scale: $G = G_{obs} \times G_{hid} \times U(1)^4$, where $G_{obs} = SU(5) \times U(1)_{\tilde{Y}}$ and $G_{hid} = SU(4)_h \times SO(10)_h$. The massless spectrum consists of 70 fields: (i) Three $\mathbf{16}$ $SO(10) \supset SU(5) \times U(1)$ chiral multiplets, $F_i(\mathbf{10}, \frac{1}{2})$, $\bar{f}_i(\bar{\mathbf{5}}, -\frac{3}{2})$, $l_i^c(\mathbf{1}, \frac{5}{2})$, $i = 1, 2, 3$; extra matter fields, $F_4(\mathbf{10}, \frac{1}{2})$, $f_4(\mathbf{5}, \frac{3}{2})$, $l_4(\mathbf{1}, -\frac{5}{2})$, and $\bar{F}_5(\bar{\mathbf{10}}, -\frac{1}{2})$, $\bar{f}_5(\bar{\mathbf{5}}, -\frac{3}{2})$, $l_5^c(\mathbf{1}, \frac{5}{2})$; and four Higgs-like fields in the $\mathbf{10}$ of $SO(10) \supset SU(5) \times U(1)$, $h_i(\mathbf{5}, -1)$, $\bar{h}_i(\bar{\mathbf{5}}, 1)$, $i = 1, 2, 3, 4, 5$. (ii) 20 gauge singlet fields $\phi_{45}, \phi^+, \phi^-, \phi_i$, ($i = 1, 2, 3, 4$), $\Phi_{12}, \Phi_{23}, \Phi_{31}$, and their “barred” counterparts; and five extra fields Φ_I , $I = 1 \rightarrow 5$.

(iii) 22 hidden sector fields in the following representations of $SO(10)_h \times SU(4)_h$, $T_i(\mathbf{10}, \mathbf{1})$, $D_i(\mathbf{1}, \mathbf{6})$, $i = 1 \rightarrow 5$; $\tilde{F}_i(\mathbf{1}, \mathbf{4})$, $\tilde{\bar{F}}_i(\mathbf{1}, \bar{\mathbf{4}})$, $i = 1 \rightarrow 6$.

There is one linear combination ($U_A(1)$) of the $U(1)^4$ which is anomalous, while the three orthogonal ones are not. This anomaly is cancelled by a string one-loop effect¹⁷ which effectively breaks $U_A(1)$ while inducing a Fayet-Iliopoulos contribution to the $U_A(1)$ D-term. D-flatness is restored by giving vacuum expectation values (vevs) to some singlet fields charged under $U_A(1)$, and this in turn breaks all superfluous $U(1)$'s.¹⁸ The choice of vevs is not unique and in effect parametrizes a discrete shift in moduli space from the original one-loop unstable vacuum to the new supersymmetric vacuum.¹⁹ The superpotential of the model (W) can be calculated explicitly and is fully known up to fifth order in the massless fields,^{20,21} but it is too extensive to be repeated here. The F-flatness conditions arising from W can be combined with the above D-flatness ones to yield a nontrivial set of constraints on the singlet vevs in the model.^{18,21} Furthermore, the *trilinear* F-flatness conditions are in fact all there are, since one can show that arbitrary products of four or more singlet fields do not contribute to the superpotential.²² This is an important property of the model that ensures the existence of flat directions of the potential to all orders in perturbation theory.

The model possesses several self-consistently determined scales ranging from the Planck scale down to the electroweak scale, which in effect map out all of the old “big desert” of grand unification. These scales are: (a) $M \approx 10^{18}$ GeV which sets the scale of the nonrenormalizable interactions;^{11,20} (b) $M_U = \mathcal{O}(10^{17})$ GeV which is where string gauge coupling unification is supposed to happen once nontrivial string threshold effects are taken into account;^{23,19} (c) $\langle \phi \rangle / M \sim 1/100 - 1/10$, as determined from the solution to the D- and F-flatness constraints;^{18,20} (d) $M_X = \mathcal{O}(10^{16})$ GeV, where $SU(5) \times U(1)_{\tilde{F}}$ breaks down to the standard model gauge group through vevs of the $\mathbf{10}_h$ and $\bar{\mathbf{10}}_h$ Higgs representations; (e) $\Lambda_{10(4)} = \mathcal{O}(10^{15(12)})$ GeV where the $SO(10)_h(SU(4)_h)$ hidden gauge group becomes strong;^{24,25} (f) various functions of the above scales, such as condensates of hidden matter fields, *e.g.*, $\langle TT \rangle \sim \Lambda_{10}^2$, $\langle \tilde{F}\tilde{F}, DD \rangle \sim \Lambda_4^2$.

We now discuss in turn the various low-energy features of the string model:

Higgs spectrum: One linear combination of the decaplets $F_{1,2,3,4}$, namely $F_\alpha \equiv \sum_{i=1}^4 \alpha_i F_i$, together with $\bar{F}_5$ will get $SU(5) \times U(1)_{\tilde{F}}$ breaking vevs, such that $V \equiv \langle F_\alpha \rangle = \langle \bar{F}_5 \rangle \equiv \bar{V}$,⁵ where

$\alpha_i = \langle F_i \rangle / V \equiv V_i / V$ and $V^2 = \sum_{i=1}^4 V_i^2$.¹⁸ The three orthogonal linear combinations to F_α , denoted by $F_\beta, F_\gamma, F_\delta$, constitute the three sets of light matter multiplets. The particular values of $V_i, \bar{V}$ will ultimately be determined dynamically, and here we take them as free parameters such that $V = \bar{V} = \mathcal{O}(10^{16})$ GeV. The trilinear superpotential term $F_4 \bar{F}_5 \phi_3$ indicates that F_4 should not get a vev, hence we take $V_4 \equiv 0$. The possibility of $V_{1,2,3} \neq 0$ is constrained by the Higgs triplet mass matrix which involves the following fields:²⁶ $\{D_{1,2,3,45}, \bar{d}_5^c\} \times \{\bar{D}_{1,2,3,45}, d_{\alpha,\beta,\gamma,\delta}^c\}$. The danger of having $V_{1,2,3} \neq 0$ is that not all of $d_\beta^c, d_\gamma^c, d_\delta^c$ may remain light at low energies. In previous studies (except in Ref. [18]) we had assumed $V_3 = 0$, but an examination of the superpotential reveals that $V_3 \neq 0$ does not affect the number of light d^c states; we thus take $V_3 \neq 0$. The choice $V_{1,2} \neq 0$ does affect the triplet mass matrix and one can show that only when $V_2 = 0$ the triplet mass matrix admits enough light d^c eigenstates.²⁷ In this case the light matter multiplets are $F_\beta = -\alpha_3 F_1 + \alpha_1 F_3$, $F_\gamma = F_2$, and $F_\delta = F_4$. The resulting Higgs triplet mass matrix is found to have all Higgs triplets heavy except for one pair that gets a mass of $\mathcal{O}(10^{10-11})$ GeV.^{20,22}

The doublet Higgs mass matrix can be similarly analyzed.^{20,22} One finds that one pair of Higgs doublets, namely H_1 and $\bar{H}_{45}$, remains massless and unmixed as long as $\langle \bar{\phi}_{45} \rangle = 0$. Furthermore, the fifth order superpotential term $h_1 \bar{h}_{45} \bar{\phi}_{45} \langle TT \rangle / M^2$ would give an effective mixing term $\mu H_1 \bar{H}_{45}$ with

$$\mu = \langle \bar{\phi}_{45} \rangle \langle TT \rangle / M^2 = \langle \bar{\phi}_{45} \rangle c (\Lambda_{10} / M)^2, \quad (1)$$

where c is a condensate-dependent constant of $\mathcal{O}(1/10)$. With $\Lambda_{10} / M \sim 10^{-4} - 10^{-3}$ we would need $\langle \bar{\phi}_{45} \rangle \sim 10^{10} - 10^{12}$ GeV to obtain $\mu \sim 1$ TeV. If $\bar{\phi}_{45}$ is massless and $\langle \bar{\phi}_{45} \rangle = 0$ at the Planck scale, then supersymmetry breaking in the hidden sector can very well trigger a vev of this size ($\sim \Lambda_4$). Fortunately, this is quite plausible since $\bar{\phi}_{45}$ appears in W as one of several fields combining with Φ_3 to get massive, and thus a light component exists. This solution to the so-called “ μ -problem” falls quite naturally among the set of solutions suggested by Kim and Nilles.²⁸

Hierarchical fermion masses: The list of massless states given above shows an excess over the standard model particles in the subset $\bar{f}_{1,2,3,5}, f_4; l_{1,2,3,5}^c, l_4$. However, there are superpotential

terms of the form²⁰ $\eta_i(\bar{f}_i f_4 + l_i^c l_4)$, $i = 1, 2, 3, 5$ which make f_4, l_4 and the linear combinations $\sum_i \eta_i \bar{f}_i$, $\sum_i \eta_i l_i^c$ heavy, leaving the three orthogonal linear combinations light. For suitable values of the η_i (calculable) parameters, this mechanism can be used to generate low-energy quark and lepton mixing from otherwise diagonal high-energy Yukawa matrices.²⁹ In this letter we pick the parameters η_i such that³⁰ $\eta_3 \gg \eta_{1,2,5}$, hence making $\bar{f}_{1,2,5}, l_{1,2,5}^c$ light and unmixed. This choice is in fact the original one proposed by Antoniadis, *et. al.*,^{5,18} and has been recently reconsidered in the context of proton decay.³¹ We now examine the fermion Yukawa matrices:

- (a) $Q = -1/3$: The relevant Yukawa couplings are $\lambda_b F_4 F_4 h_1$, $\lambda_s F_2 F_2 h_1$,²⁰ and $\lambda_d F_\beta F_\beta h_1$ (with $F_1 = \alpha_3 F_\beta$), with $\lambda_b = g\sqrt{2}$, $\lambda_d = \alpha_3^2 = (V_3/V)^2$, and

$$\lambda_s = g \left\{ \sum_{i=1}^4 c_{si} \langle \bar{\phi}_i \rangle^2 + c_{s\pm} \langle \bar{\phi}^+ \bar{\phi}^- \rangle \right\} / M^2. \quad (2)$$

- (b) $Q = 2/3$: The Yukawa couplings are²⁰ $\lambda_t F_4 \bar{f}_5 \bar{h}_{45}$ and $\lambda_c F_2 \bar{f}_2 \bar{h}_{45}$, with $\lambda_t = g\sqrt{2}$ and $\lambda_c = 3.04g \langle \bar{\phi}_4 \rangle / M$. There are also two other terms²⁰ $\lambda_u^{(1,5)} F_1 \bar{f}_{1,5} \bar{h}_{45} \rightarrow \lambda_u^{(1,5)} V_1 L_{1,5} \bar{H}_{45}$, but we must demand $\lambda_u^{(1,5)} \approx 0$ to keep $\bar{H}_{45}$ light. Thus, $\lambda_u = 0$.

- (c) $Q = -1$: The Yukawa couplings are²⁰ $\lambda_\tau \bar{f}_1 l_1^c h_1$, $\lambda_\mu \bar{f}_2 l_2^c h_1$, $\lambda_e \bar{f}_5 l_5^c h_1$ with²⁰ $\lambda_\tau = g\sqrt{2}$ and

$$\lambda_\mu = g \left\{ \sum_{i=1}^4 c_{\mu i} \langle \bar{\phi}_i \rangle^2 + c_{\mu\pm} \langle \bar{\phi}^+ \bar{\phi}^- \rangle \right\} / M^2, \quad (3)$$

$$\lambda_e = g \{ c_{e1} \langle \bar{\phi}_1 \rangle^2 + c_{e4} \langle \bar{\phi}_4 \rangle^2 + c_{e\pm} \langle \bar{\phi}^+ \bar{\phi}^- \rangle \} / M^2. \quad (4)$$

The above Yukawa couplings (at M_U) show several remarkable features: (i) $\lambda_b = \lambda_\tau$ which implies through the usual argument⁴ the correct low-energy m_b/m_τ ratio; (ii) the low-energy mass ratios m_s/m_b and m_d/m_b can be easily fit by natural values of the parameters, *i.e.*, $\langle \phi \rangle / M \sim 1/10$, $V_3/V \approx 1/25$; (iii) an explicit calculation shows that $c_{si} = c_{\mu i}$, $i = 1, 2, 3, 4$, and $c_{s\pm} = c_{\mu\pm}$. Thus, we obtain the successful relation $\lambda_s = \lambda_\mu$ (or equivalently $\lambda_s/\lambda_b = \lambda_\mu/\lambda_\tau$), independently of the values of the vevs involved; (iv) also, $c_{e4} = -c_{e1} = -c_{e\pm} \approx 3.32$, which implies the unexpected although realistic relation $\lambda_c/\lambda_t = (\lambda_e/\lambda_\tau)^{1/2} \mathcal{O}(1/2 - 2)$.

Another consequence of the Yukawa couplings is the *unambiguous* assignment of the light particles to the string representations, as follows

$$\begin{aligned}
t b \tau \nu_\tau : & \quad Q_4 d_4^c u_5^c L_1 l_1^c, \\
c s \mu \nu_\mu : & \quad Q_2 d_2^c u_2^c L_2 l_2^c, \\
u d e \nu_e : & \quad Q_\beta d_\beta^c u_1^c L_5 l_5^c.
\end{aligned} \tag{5}$$

Let us point out that $\theta_{QCD} = 0$ since $m_u = 0$. Also, note that to the present level of approximation the model has negligible KM mixing. There are a few effects that may perturb this result: (i) loop-induced low-energy effects after supersymmetry breaking, (ii) $V_2 \neq 0$ may still be allowable for special values of the parameters, (iii) non-perturbative string effects relaxing some of the string symmetries that hold in perturbation theory.

Proton decay: The various operators contributing to proton decay in this model have to be studied carefully since some of them may be potentially disastrous. Here we present an overview of the analysis which shows that in fact all these operators yield acceptable proton decay rates, and refer the reader to Ref. [32] for a more detailed treatment. The operators are:

- (a) Nonrenormalizable terms in the superpotential of the form^{33,31} $FFF\bar{f}\frac{1}{M}$ contain effective $d = 5$ $qqql$ operators which could be very large^{33,31} since the overall coefficient is expected²⁰ to be $\mathcal{O}(1)$. These terms appear first at fifth order^{20,31} (multiplied by $\langle\phi\rangle/M$) and always contain the $\bar{f}_3$ field, and thus are harmless here since $\bar{f}_3$ does not contain any light states.
- (b) Dimension-five Higgsino mediated operators³⁴ are constructed via the usual tree-level diagrams involving three superpotential couplings: FFh , $F\bar{f}\bar{h}$, and $h\bar{h}$. Since we are interested in taking $\bar{f}$ to be $\bar{f}_{1,2,5}$, the associated $\bar{h}$ is always $\bar{h}_{45}$. The needed mixing term $h\bar{h}_{45}$ is proportional to μ for $h_{1,2}$ and to $\langle\bar{\phi}_{45}\rangle$ for h_3 , and the $F\bar{f}\bar{h}$ vertex coupling is in effect the up-quark Yukawa matrix. A careful analysis³² shows that the coefficients of the resulting effective $d = 5$ operators are always suppressed enough, although the potential proton decay rates could be close to the current experimental limit.
- (c) Dimension-six Higgs boson mediated operators could be important because of the existence of $\mathcal{O}(10^{10-11})$ GeV Higgs triplet states, as mentioned above. There are two classes

of diagrams³⁴ originating from superpotential couplings of the forms: $F\bar{f}\bar{h}/F^\dagger\bar{f}^\dagger\bar{h}^\dagger$ and $FFh/\bar{f}^\dagger l^{c\dagger}h^\dagger$. The first class yield effective operators suppressed by $1/\tan^2\beta$ relative to their counterpart in SUSY GUTs,³² while the second class give negligible contributions.³²

- (d) Dimension-six gauge boson exchange operators are sufficiently suppressed in flipped $SU(5)$ models,^{33,1} with a nucleon lifetime in the $10^{32}y - 10^{36}y$ range.¹ The expected dominant decay modes are easily determined^{34,18} from the particle assignment in Eq. (5): $p \rightarrow \bar{\nu}_\tau\pi^+$ and $n \rightarrow \bar{\nu}_\tau\pi^0$.

Neutrino masses: The see-saw neutrino mass matrix in flipped $SU(5)$ involves ν, ν^c , and singlet states ϕ ,^{13,35} through the superpotential couplings: $\lambda_u^{ij}F_i\bar{f}_j\bar{h}_{45} \ni \lambda_u^{ij}\nu_i^c\nu_j \langle \bar{H}_{45}^0 \rangle$; $w_{ij}F_i\bar{F}_5\phi_j \ni w_{ij}\bar{V}\nu_i^c\phi_j$; and $\mu_{ij}\phi_i\phi_j$. In the string model, the see-saw matrix turns out to be 14×14 $\{\nu_{1,2,5}, \nu_{\beta,2,4}^c, \phi_{1,2,3,4}, \bar{\phi}_{45}, \Phi_{3,4,5}\}$. An important point is that all the $\nu^c\nu^c$ entries vanish. However, there are two-loop radiatively induced contributions of the form³⁶ $m_{\nu^c\nu^c} \approx (\alpha_{GUT}/4\pi)^2 (3/16)\lambda_{d,s,b}M_X \approx 5 \times 10^8 m_{d,s,b}$, which contribute decisively to the structure of the see-saw matrix. If we choose $\langle \bar{\phi}_2 \rangle = 0$, then the matrix breaks up into three blocks with only three light eigenvalues: $\nu_1, m_1 = 0$; $\approx \nu_2, m_2 = 0$; and $\approx \nu_5, m_5 \approx m_t\sqrt{\lambda_d}\langle TT \rangle/M^2$. Hence (see Eq. (5)) ν_μ and ν_τ are massless to this level of approximation, and ν_e has a mass in the eV range, within the current experimental limit $m_{\nu_e} < 10 \text{ eV}$.³⁷

Top quark mass and $\tan\beta$: The one-loop RGEs for the top and bottom quark Yukawa couplings can be run from $M_X \approx 10^{16} \text{ GeV}$ down to M_Z with the initial conditions $\lambda_t = \lambda_b = g\sqrt{2} \approx 1$. At low energies, the relations $m_t = \lambda_tv_2$, $m_b = \lambda_bv_1$, and $\tan\beta = v_2/v_1$, give $m_t \lesssim 170 \text{ GeV}$ and $\tan\beta \lesssim 33$, where the inequality accounts for the supersymmetry breaking effects that have been neglected. This value of the top-quark mass falls within its expected interval.³⁸

In sum, we have presented a sharpened up version of the flipped $SU(5)$ string model which accommodates numerous features of the low-energy physics world in one consistent calculational framework. In the next step in the development of the model we would like to quantify the connection between the low- and high-energy regimes of the theory through RGE evolution and the computation of string threshold corrections and hidden sector matter condensates.

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