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THE FLIPPED SU(5)×U(1) STRING MODEL REVAMPED

I. Antoniadis^{*)}, John Ellis, J.S. Hagelin⁺⁾ and D.V. Nanopoulos^{")}

CERN - Geneva

E R R A T A

In equation (1g), replace $\bar{\phi}^{-1}, \dots, 8$ by $\bar{\phi}^{-1}, \dots, 8$

On the second line of equation (7):

replace $(\underline{10}, \underline{1})_{(0, -\frac{1}{2}, -\frac{1}{2}, 0)}$ by $(\underline{10}, \underline{1})_{(0, \frac{1}{2}, -\frac{1}{2}, 0)}$

In equation (8), add:

$$b_1 + b_2 + b_4 \pm \alpha(+\zeta): (\underline{1}, \underline{4})^{5/4}_{(-1/4, 1/4, -1/4, -\frac{1}{2})} + (\underline{1}, \underline{4})^{-5/4}_{(1/4, -1/4, 1/4, -\frac{1}{2})}$$

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A B S T R A C T

We present a refined version of our three-generation flipped SU(5)×U(1) string model with the following properties.

- The complete massless spectrum is derived and shown to be free of all gauge and mixed anomalies apart from a single anomalous U(1).

- The imaginary part of the dilaton supermultiplet is eaten by the anomalous U(1) gauge boson, and the corresponding D-term is cancelled by large v.e.v.'s for singlet fields that break surplus U(1) gauge factors, leaving a supersymmetric vacuum with an SU(5)×U(1) visible gauge group and an SO(10)×SO(6) hidden gauge group.

- There are sufficient Higgs multiplets to break the visible gauge symmetry down to the Standard Model in an essentially unique way.

- All trilinear superpotential couplings have been calculated and there are in particular some giving $m_t, m_b, m_\tau \neq 0$.

- A renormalization group analysis shows that $m_t < 190$ GeV and $m_b = 3m_\tau$.

- Light Higgs doublets are split automatically from heavy Higgs triplets, leaving no residual dimension-five operators for baryon decay, and the baryon lifetime $\tau_B \sim 2 \times 10^{34 \pm 2} \text{ y}$.

- There are no tree-level flavour-changing neutral currents, but $\mu \rightarrow e\gamma$ may occur at a detectable level: $B(\mu \rightarrow e\gamma) \sim 10^{-11}$ to 10^{-14} .

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There is a regrettable schism in particle physics today. On the one hand, many theorists are convinced that string is the Theory of Everything (TOE), and strive hard to understand better its formalism without worrying about its physical signatures. On the other hand, experiment is confined dismally within the Standard Model, and many phenomenologists strive hard to constrain its parameters without worrying about its origin. Surprisingly little effort is directed towards deriving the Standard Model from some variant of string theory, or to seeking some other distinctive signature of the TOE. Many would doubtless consider any such effort premature and foolhardy until we understand string better. However, if we wait around for someone to tell us how non-perturbative string effects determine the correct vacuum state, we will have a long wait. Indeed, since we now have many tools - geometric [1] and algebraic [2], bosonic [3] and fermionic [4,5] - for constructing different string models, surely we should be using them to look for candidate vacua that are phenomenologically realistic.

It may well be that the Standard Model can be derived directly from string without invoking any field-theoretical intermediate scale of gauge symmetry breaking [6,7]. However, so far all phenomenological string models [7-9] have had four-dimensional groups larger than $SU(3) \times SU(2) \times U(1)$, though often only with extra $U(1)$ factors [8]. We think [9] that if one is to extend the Standard Model gauge group, it is interesting and desirable to embed it in a Grand Unified Theory (GUT), i.e., a simple non-Abelian group containing the $SU(3)$, $SU(2)$ and (maybe) $U(1)$ factors of the Standard Model. Such a framework would combine the physical advantages of GUTs (e.g., slow baryon decay, cosmological baryosynthesis, small neutrino masses, etc.) with the well-known benefits of the string. However, there is an obstacle to such a programme that we have emphasized previously [9]: in general, GUTs require Higgs fields in adjoint representations to break the gauge symmetry down to the Standard Model, and these do not exist in string theories with $N = 1$ supersymmetry and/or chiral fermions [10,11,4] that have level $K = 1$ Kac-Moody algebras.

The only viable GUT that does not require adjoint Higgses is flipped $SU(5) \times U(1)$ [12,9], which also possesses [9,13] other advantages such as natural Higgs doublet-triplet splitting, a seesaw neutrino mass matrix, no problematic fermion mass relations, and no troublesome $d = 5$ proton decay operators^{*)}. We have in the past constructed realistic flipped $SU(5) \times U(1)$ models in field theory and indicated how they could be derived from manifold compactifications [15] or the fermionic formulation of string theory [16,17]. However, although our efforts using free world-sheet fermions [4,5] encouraged us that we were on the right track, our

*) Other proposed flipped unification models [14] have not been shown to share all these desirable features.

previous models [16,17] suffered from various technical problems and lacunae, including the following. (1) The spectra of light states were not correctly stated [18]. (2) Certain desired Yukawa couplings, such as that responsible for the top quark mass, were absent [19]. (3) The hidden sector was not worked out in detail, and hence the absence of gauge anomalies in the effective low-energy field theory could not be checked. (4) The "eating" of the dilaton [20] was not discussed in detail. (5) Quark mass mixing and the weak interaction angles were not completely satisfactory. In this paper we remedy the above defects.

We present a new version of the three-generation flipped $SU(5) \times U(1)$ model derived from string theory with the following features. (a) There is an essentially unique pattern of gauge symmetry breaking down to the Standard Model. (b) We exhibit the complete anomaly-free spectrum. (c) The imaginary part of the dilaton supermultiplet^{*)} is "eaten" by an anomalous $U(1)$ gauge boson [20], and the dilaton contribution to the corresponding D-term is cancelled [21] by field v.e.v.'s that break at least two surplus $U(1)$ gauge group factors. (d) The resulting gauge group is just minimal flipped $SU(5) \times U(1)$ with a hidden sector gauge group $SO(10) \times SO(6)$, and no problematic states in the hidden sector. (e) We calculate Yukawa couplings that give tree-level masses to the t quark and to the b quark and τ lepton. We show that $m_t < 190$ GeV and $m_b = m_\tau$ at the GUT scale leading to the physical $m_b \approx 5$ GeV. (f) Proton decay and flavour-changing neutral currents are naturally suppressed. (g) The possibility exists of phenomenologically acceptable quark masses and mixing.

It should be stressed that in our model the chiral multiplets do not fill 27 representations of E_6 , but are Weyl spinors 16 of $O(10)$, and thus there are no extra low-energy states beyond the right-handed neutrinos. This is because our model is in the $(2,0)$ class of $N = 1$ supersymmetric string vacua, and not in the $(2,2)$ class [1,2]. We also emphasize that the existence of an anomalous $U(1)$, which is a generic feature of our construction, resolves problems associated with a massless dilaton that are as yet unresolved in $U(1)$ anomaly-free four-dimensional string models.

Our four-dimensional string model is constructed using the free fermionic formulation of Ref. [4]. It is generated by the following basis of eight vectors of boundary conditions for all the world-sheet fermions:

$$S = \left(\underbrace{1, \dots, 1}_{\chi^{\mu}, \chi^1, \dots, 6}, 0, \dots, 0; 0, \dots, 0 \right) \quad (1a)$$

*) This is identified with the two-index antisymmetric tensor field.

$$b_1 = (\underbrace{1, \dots, 1}_{\chi^\mu, \chi^{1,2}, y^{3, \dots, 6}, \bar{y}^{3, \dots, 6}}, \underbrace{0, \dots, 0}_{\bar{\chi}^{1, \dots, 5}, \bar{\eta}^1}; \underbrace{1, \dots, 1}_{\bar{\chi}^{1, \dots, 5}, \bar{\eta}^1}, \underbrace{0, \dots, 0}_{\bar{\chi}^{1, \dots, 5}, \bar{\eta}^1}) \quad (1b)$$

$$b_2 = (\underbrace{1, \dots, 1}_{\chi^\mu, \chi^{3,4}, y^{1,2}, \bar{y}^{1,2}, \omega^{5,6}, \bar{\omega}^{5,6}}, \underbrace{0, \dots, 0}_{\bar{\chi}^{1, \dots, 5}, \bar{\eta}^2}; \underbrace{1, \dots, 1}_{\bar{\chi}^{1, \dots, 5}, \bar{\eta}^2}, \underbrace{0, \dots, 0}_{\bar{\chi}^{1, \dots, 5}, \bar{\eta}^2}) \quad (1c)$$

$$b_3 = (\underbrace{1, \dots, 1}_{\chi^\mu, \chi^{5,6}, \omega^{1, \dots, 4}, \bar{\omega}^{1, \dots, 4}}, \underbrace{0, \dots, 0}_{\bar{\chi}^{1, \dots, 5}, \bar{\eta}^3}; \underbrace{1, \dots, 1}_{\bar{\chi}^{1, \dots, 5}, \bar{\eta}^3}, \underbrace{0, \dots, 0}_{\bar{\chi}^{1, \dots, 5}, \bar{\eta}^3}) \quad (1d)$$

$$b_4 = (\underbrace{1, \dots, 1}_{\chi^\mu, \chi^{1,2}, y^{3,6}, \bar{y}^{3,6}, \omega^{4,5}, \bar{\omega}^{4,5}}, \underbrace{0, \dots, 0}_{\bar{\chi}^{1, \dots, 5}, \bar{\eta}^1}; \underbrace{1, \dots, 1}_{\bar{\chi}^{1, \dots, 5}, \bar{\eta}^1}, \underbrace{0, \dots, 0}_{\bar{\chi}^{1, \dots, 5}, \bar{\eta}^1}) \quad (1e)$$

$$b_5 = (\underbrace{1, \dots, 1}_{\chi^\mu, \chi^{3,4}, \omega^{5,6}, \bar{\omega}^{5,6}, y^{1,5}, \bar{y}^{1,5}, y^{2,6}, \bar{y}^{2,6}}, \underbrace{0, \dots, 0}_{\bar{\chi}^{1, \dots, 5}, \bar{\eta}^2}; \underbrace{1, \dots, 1}_{\bar{\chi}^{1, \dots, 5}, \bar{\eta}^2}, \underbrace{0, \dots, 0}_{\bar{\chi}^{1, \dots, 5}, \bar{\eta}^2}) \quad (1f)$$

$$\zeta = (0, \dots, 0, \underbrace{0, \dots, 0}_{\bar{\chi}^{1, \dots, 8}}, \underbrace{1, \dots, 1}_{\bar{\chi}^{1, \dots, 8}}) \quad (1g)$$

$$\alpha = (\underbrace{1, \dots, 1}_{y^{4,6}, \bar{y}^{4,6}, \omega^{4,6}, \bar{\omega}^{4,6}, \omega^{2,3,4,6}}, \underbrace{0, \dots, 0}_{\bar{\chi}^{1, \dots, 5}, \bar{\eta}^{1,2,3}}, \underbrace{\frac{1}{2}, \dots, \frac{1}{2}}_{\bar{\chi}^{1, \dots, 5}, \bar{\eta}^{1,2,3}}, \underline{A}) \quad (1h)$$

with

$$\underline{A} = (\underbrace{\frac{1}{2}, \dots, \frac{1}{2}}_{\bar{\phi}^{1, \dots, 4}}, \underbrace{\frac{1}{2}}_{\bar{\phi}^{5,6}}, \underbrace{0, 0}_{\bar{\phi}^{7,8}}) \quad (1i)$$

where we have used the notation of Ref. [4]: thus 1 stands for periodic fermions, 0 for antiperiodic, and $\frac{1}{2}$ for those twisted by a phase $-i$. The semicolon separates left- from right-movers: we have chosen a basis in which all left-movers ($\chi^\mu, \chi^i, y^i, \omega^i$; $i = 1, 2, \dots, 6$) are real, among which supersymmetry is realized non-linearly, 12 right-movers ($\bar{y}^i, \bar{\omega}^i$) are real, and 16 ($\bar{\phi}^{1,2, \dots, 5}, \bar{\eta}^{1,2,3}, \bar{\phi}^{1,2, \dots, 8}$) are complex. We make the following choice of generalized GSO projection [22] coefficients:

$$c\left(\begin{smallmatrix} b_i \\ S \end{smallmatrix}\right) = c\left(\begin{smallmatrix} b_i \\ b_j \end{smallmatrix}\right) = c\left(\begin{smallmatrix} \alpha \\ b_i \end{smallmatrix}\right)_{i \neq 5} = c\left(\begin{smallmatrix} \alpha \\ \zeta \end{smallmatrix}\right) = -c\left(\begin{smallmatrix} \alpha \\ b_5 \end{smallmatrix}\right) = -1 \quad (2)$$

with the others specified by modular invariance and space-time supersymmetry.

The basis vectors $\{S, \zeta, 1 = b_1 + b_2 + b_3 + \zeta\}$ define an $N = 4$ space-time supersymmetric model with an $SO(28) \times E_8$ gauge group: S plays the rôle of the supersymmetry generator, since when added to a sector it gives its superpartner. The vectors b_1 and b_2 reduce to $N = 1$ supersymmetry, break $SO(28) \rightarrow SO(10) \times SO(6)^3$, and give six chiral families $(\underline{16} + \underline{4}) + (\underline{16} + \bar{\underline{4}})$, two from each of the sectors b_1, b_2 and b_3 . The vectors b_4, b_5 and 2α break $SO(6)^3 \rightarrow U(1)^6$, $E_8 \rightarrow SO(16)$ and lead to six chiral $SO(10)$ families. In addition, in the observable sector there are two extra

$\underline{16}+\overline{\underline{16}}$ pairs from b_4 and b_5 , and two $\underline{10}$'s from b_4+b_5 . Finally, the vector α breaks $SO(10) \rightarrow$ flipped $SU(5) \times U(1)$, $U(1)^6 \rightarrow U(1)^4$ and the "hidden" group $SO(16) \rightarrow SO(10) \times SO(6)$, and projects out half the chiral families. The four remaining $U(1)$'s correspond to the right-moving world-sheet currents $\eta_1 \eta^*_1$, $\eta_2 \eta^*_2$, $\eta_3 \eta^*_3$ and $\bar{\omega}_2 \bar{\omega}_3$.

Thus the following massless matter particles are produced by the sectors $b_{1,2,3,4,5}$, $S+b_4+b_5$, 0 and their superpartners in the observable $SU(5) \times U(1) \times U(1)^4$ sector.

(a) The $b_{1,2,3}$ sectors produce three $SO(10)$ chiral families $M_\alpha = F_\alpha + \bar{f}_\alpha + \lambda_\alpha^c$ ($\alpha=1,2,3$) with $F \equiv (\underline{10}, 1/2)$, $\bar{f} \equiv (\bar{5}, -3/2)$ and $\lambda^c \equiv (\underline{1}, 5/2)$ of $SU(5) \times U(1)$ and the following extra $U(1)$ charges:

$$\begin{pmatrix} M_1 \end{pmatrix}_{(-\frac{1}{2}, 0, 0, 0)} ; \begin{pmatrix} M_2 \end{pmatrix}_{(0, -\frac{1}{2}, 0, 0)} ; \begin{pmatrix} F_3 \end{pmatrix}_{(0, 0, \frac{1}{2}, -\frac{1}{2})} + \begin{pmatrix} \bar{f}_3 + \lambda_3^c \end{pmatrix}_{(0, 0, \frac{1}{2}, \frac{1}{2})} \quad (3)$$

(b) The $b_{4,5}$ sectors give

$$\begin{pmatrix} F_4 \end{pmatrix}_{(-\frac{1}{2}, 0, 0, 0)} ; \begin{pmatrix} f_4 \end{pmatrix}_{(\frac{1}{2}, 0, 0, 0)} ; \begin{pmatrix} \bar{\lambda}_4^c \end{pmatrix}_{(\frac{1}{2}, 0, 0, 0)} \text{ and } \begin{pmatrix} \bar{F}_5 \end{pmatrix}_{(0, \frac{1}{2}, 0, 0)} ; \begin{pmatrix} \bar{f}_5 \end{pmatrix}_{(0, -\frac{1}{2}, 0, 0)} ; \begin{pmatrix} \lambda_5^c \end{pmatrix}_{(0, \frac{1}{2}, 0, 0)} \quad (4)$$

respectively, where $f \equiv (\underline{5}, 3/2)$, $\bar{\lambda}^c \equiv (\underline{1}, -5/2)$ and $\bar{F} \equiv (\overline{\underline{10}}, -1/2)$.

(c) The $S+b_4+b_5$ sector gives

$$h_{45} \equiv (\underline{5}, -1)_{(-\frac{1}{2}, -\frac{1}{2}, 0, 0)} ; \phi_{45} \equiv (\underline{1}, 0)_{(\frac{1}{2}, \frac{1}{2}, 1, 0)} ; \phi_i \equiv (\underline{1}, 0)_{(\frac{1}{2}, -\frac{1}{2}, 0, 0)} \quad (i=1 \text{ to } 4) ;$$

$$\phi^+ \equiv (\underline{1}, 0)_{(\frac{1}{2}, -\frac{1}{2}, 0, 1)} ; \phi^- \equiv (\underline{1}, 0)_{(\frac{1}{2}, -\frac{1}{2}, 0, -1)} \quad (5)$$

(and their conjugates $\bar{h}_{45}$, etc.) obtained by acting on its vacuum with the $\bar{\psi}^{a*}$, $\bar{\eta}^3$, $\bar{\omega}^2 + i\bar{\omega}^3$ fermionic oscillators for the first four states respectively, and by those of $\bar{y}^{5,6}$, $\bar{\omega}^{5,6}$ for the ϕ_i : $i = 1$ to 4 (and their complex conjugates for $\bar{h}_{45}$, etc.).

(d) Finally, the Neveu-Schwarz 0 sector gives

$$h_1 \equiv (\underline{5}, -1)_{(1, 0, 0, 0)} ; h_2 \equiv (\underline{5}, -1)_{(0, 1, 0, 0)} ; h_3 \equiv (\underline{5}, -1)_{(0, 0, 1, 0)} ;$$

$$\Phi_{23} \equiv (\underline{1}, 0)_{(0, -1, 1, 0)} ; \Phi_{31} \equiv (\underline{1}, 0)_{(1, 0, -1, 0)} ; \Phi_{12} \equiv (\underline{1}, 0)_{(-1, 1, 0, 0)} ;$$

$$\text{their complex conjugates; } \Phi_I \equiv (\underline{1}, 0)_{(0, 0, 0, 0)} \quad (I=1 \text{ to } 5) \quad (6)$$

*) We thank A. Faraggi and K. Yuan who have each independently verified the massless spectrum of our model.

obtained by acting on the vacuum with the fermionic oscillators of $\chi^1+i\chi^2$ with $\bar{\psi}^{a*}\bar{\eta}^1$, $\bar{\eta}^2\bar{\eta}^3$, $\bar{y}^1\bar{\omega}^1$ (for h_1, Φ_{23}, Φ_1), $\chi^3+i\chi^4$ with $\bar{\psi}^{a*}\bar{\eta}^2, \bar{\eta}^1\bar{\eta}^3$, $\bar{y}^4\bar{\omega}^4$ (for h_2, Φ_{31}, Φ_2), and $\chi^5+i\chi^6$ with $\bar{\psi}^{a*}\bar{\eta}^3$, $\bar{\eta}^1\bar{\eta}^2$, $\bar{y}^1\bar{\omega}^1$, $\bar{y}^2\bar{\omega}^2$, $\bar{y}^3\bar{\omega}^3$ and $\bar{y}^4\bar{\omega}^4$ (for $h_3, \Phi_{12}, \Phi_{3,4,5}$).

Now we also list the massless matter fields transforming under the hidden $SO(10) \times SO(6)$ gauge group.

(e) The sectors $b_i+2\alpha$ and $b_i+2\alpha+\zeta$ for $i = 1$ to 5 give the following vector representations:

$$\begin{aligned} & \left(\begin{pmatrix} 1 \\ 6 \end{pmatrix} + \begin{pmatrix} 10 \\ 1 \end{pmatrix} \right)_{(0, \frac{1}{2}, \frac{1}{2}, 0)}; \left(\begin{pmatrix} 1 \\ 6 \end{pmatrix} + \begin{pmatrix} 10 \\ 1 \end{pmatrix} \right)_{(-\frac{1}{2}, 0, \frac{1}{2}, 0)}; \begin{pmatrix} 1 \\ 6 \end{pmatrix}_{(-\frac{1}{2}, -\frac{1}{2}, 0, \frac{1}{2})} + \begin{pmatrix} 10 \\ 1 \end{pmatrix}_{(-\frac{1}{2}, -\frac{1}{2}, 0, -\frac{1}{2})}; \\ & \begin{pmatrix} 1 \\ 6 \end{pmatrix}_{(0, \frac{1}{2}, \frac{1}{2}, 0)} + \begin{pmatrix} 10 \\ 1 \end{pmatrix}_{(0, -\frac{1}{2}, -\frac{1}{2}, 0)}; \begin{pmatrix} 1 \\ 6 \end{pmatrix}_{(\frac{1}{2}, 0, -\frac{1}{2}, 0)} + \begin{pmatrix} 10 \\ 1 \end{pmatrix}_{(-\frac{1}{2}, 0, \frac{1}{2}, 0)} \end{aligned} \quad (7)$$

which are $SU(5) \times U(1)$ singlets.

(f) The following sectors give spinorial representations of $SO(6)$:

$$\begin{aligned} & b_1 \pm \alpha (+\zeta): \begin{pmatrix} 1 \\ \bar{4} \end{pmatrix}_{(-\frac{1}{4}, \frac{1}{4}, \frac{1}{4}, \frac{1}{2})}^{-5/4} + \begin{pmatrix} 1 \\ \bar{4} \end{pmatrix}_{(-\frac{1}{4}, \frac{1}{4}, \frac{1}{4}, -\frac{1}{2})}^{-5/4} \\ & b_1 + b_4 + b_5 \pm \alpha (+\zeta): \begin{pmatrix} 1 \\ 4 \end{pmatrix}_{(-\frac{1}{4}, \frac{1}{4}, -\frac{1}{4}, \frac{1}{2})}^{5/4} + \begin{pmatrix} 1 \\ 4 \end{pmatrix}_{(-\frac{1}{4}, \frac{1}{4}, -\frac{1}{4}, -\frac{1}{2})}^{5/4} \\ & S + b_2 + b_4 \pm \alpha (+\zeta): \begin{pmatrix} 1 \\ 4 \end{pmatrix}_{(\frac{1}{4}, \frac{1}{4}, -\frac{1}{4}, \frac{1}{2})}^{-5/4} + \begin{pmatrix} 1 \\ \bar{4} \end{pmatrix}_{(-\frac{1}{4}, -\frac{1}{4}, \frac{1}{4}, -\frac{1}{2})}^{5/4} \\ & b_4 \pm \alpha (+\zeta): \begin{pmatrix} 1 \\ 4 \end{pmatrix}_{(\frac{1}{4}, -\frac{1}{4}, -\frac{1}{4}, -\frac{1}{2})}^{5/4} + \begin{pmatrix} 1 \\ \bar{4} \end{pmatrix}_{(-\frac{1}{4}, \frac{1}{4}, \frac{1}{4}, -\frac{1}{2})}^{-5/4} \\ & b_2 + b_3 + b_5 \pm \alpha (+\zeta): \begin{pmatrix} 1 \\ 4 \end{pmatrix}_{(-\frac{3}{4}, \frac{1}{4}, -\frac{1}{4}, 0)}^{-5/4} + \begin{pmatrix} 1 \\ \bar{4} \end{pmatrix}_{(-\frac{1}{4}, \frac{3}{4}, \frac{1}{4}, 0)}^{5/4} \end{aligned} \quad (8)$$

which are $SU(5)$ singlets, but carry the non-zero $U(1)$ charges indicated by the upper right index, corresponding to fractional electric charges $\pm \frac{1}{2}$.

We note some immediate consequences of the above spectrum. (i) There are only three generations, and all the fields needed to break $SU(5) \times U(1) \times U(1)^4$ down to the Standard Model ($\underline{10} + \overline{\underline{10}}, \underline{5} + \overline{\underline{5}}$, singlets) are available. (ii) Among the four surplus $U(1)$'s, there is one anomalous linear combination

$$U(1)_A \equiv -3U(1)_1 - U(1)_2 + 2U(1)_3 - U(1)_4 \quad (9a)$$

whilst the three orthogonal combinations

$$U(1)'_1 = U(1)_3 + 2 U(1)_4 \quad (9b)$$

$$U(1)'_2 = U(1)_1 - 3 U(1)_2 \quad (9c)$$

$$U(1)'_3 = 3 U(1)_1 + U(1)_2 + 4 U(1)_3 - 2 U(1)_4 \quad (9d)$$

are free of all gauge and mixed gravitational anomalies. This is a non-trivial consistency check of the massless spectrum, including the hidden sector. (iii) The appearance of fractionally-charged states in any sector involving $\pm\alpha$ is closely related to the breaking of $SO(10)$ to $SU(5) \times U(1)$.

We now turn to the Yukawa couplings, which can easily be calculated in the free fermionic formulation of four-dimensional string models [4]. They correspond to expectation values of the form

$$\langle \int d^2 z_1 \int d^2 z_2 \int d^2 z_3 V_1^F(z_1) V_2^F(z_2) V_3^B(z_3) \rangle \quad (10)$$

where the $V_i^{F(B)}$ are the vertex operators for fermionic (bosonic) states, and the z_i are the two-dimensional co-ordinates. Using $SL(2, \mathbb{C})$ invariance, one can choose, for instance, $z_1 = 1$, $z_2 = 0$, $z_3 = \infty$ so that (4) is identified with the corresponding operator product coefficient c_{123} :

$$V_1^F(z) V_2^F(w) \underset{z \rightarrow w}{\sim} c_{123} \frac{V_3^B(w)}{|z-w|^2} + \dots \quad (11a)$$

since the V_i are $(1,1)$ conformal fields normalized:

$$V_i(z) V_j^*(w) \underset{z \rightarrow w}{\sim} \frac{\delta_{ij}}{|z-w|^4} + \dots \quad (11b)$$

The three-point function (10) can be factorized as a product of several three-point functions corresponding to the different world-sheet fields: two-point functions do not enter, because all the fields are orthonormalized. The individual contributions can be summarized as follows. (i) For a complex two-dimensional fermion f one has

$$\langle e^{Q_1 \phi} e^{Q_2 \phi} e^{Q_3 \phi} \rangle = \delta_{Q_1 + Q_2 + Q_3, 0} \quad (12a)$$

where $e^{Q_i \phi}$ is the corresponding part in the vertex operator of a state having charge Q_i for the $U(1)$ current $f^* f = \partial \phi$: Q_i is related to the boundary condition α_i for the corresponding sector by $Q_i = (\alpha_i/2) + F_i$, where F_i is the two-dimensional fermion number. One also has

$$\langle e^{Q_1 \phi} e^{Q_2 \phi} \partial \phi \rangle = Q_1 \delta_{Q_1+Q_2,0} \quad (12b)$$

which usually only appears in a gauge coupling, since the current $\partial \phi$ corresponds to some adjoint quantum number, and there are no adjoint matter fields in chiral theories with a level $K = 1$ Kac-Moody algebra. (ii) For a real two-dimensional fermion f (left) and $\bar{f}$ (right) there are two twist fields $\sigma_{\pm}$ with conformal dimensions $(1/16, 1/16)$ corresponding to the two different fermion number projections $(-1)^{f\bar{f}} = \pm 1$ in the Ramond sector. Normalizing $\langle \sigma_+ \sigma_+ \rangle = \langle \sigma_- \sigma_- \rangle = 1$ and $\langle \sigma_+ \sigma_- \rangle = 0$, the non-vanishing three-point functions are [23]

$$\langle \sigma_+ \sigma_- f \rangle = \langle \sigma_+ \sigma_- \bar{f} \rangle = \frac{1}{\sqrt{2}} \quad (13)$$

It is now straightforward to write down all the trilinear superpotential terms for massless matter fields in the observable sectors (a) to (d) above:

$$\begin{aligned} W = g\sqrt{2} \Big[& F_1 F_1 h_1 + F_2 F_2 h_2 + F_4 F_4 h_1 + F_4 \bar{F}_5 h_{45} + F_3 \bar{F}_3 h_3 \\ & + \bar{F}_1 \bar{F}_1 h_1 + \bar{F}_2 \bar{F}_2 h_2 + \bar{F}_5 \bar{F}_5 h_2 + \frac{1}{\sqrt{2}} (F_4 \bar{F}_5 \phi_3 + \bar{F}_4 \bar{F}_5 \phi_2 + \bar{F}_4 \bar{F}_5 \phi_2) \\ & + \bar{F}_5 \bar{F}_5 h_2 + \bar{F}_4 \bar{F}_4 h_1 + \bar{F}_5 \bar{F}_4 h_{45} + h_1 \bar{h}_2 \phi_{12} + h_2 \bar{h}_3 \phi_{23} + h_3 \bar{h}_1 \phi_{31} + h_3 \bar{h}_{45} \phi_{45} + h.c. \\ & + \frac{1}{2} (\phi_{45} \bar{\phi}_{45} \Phi_3 + \phi^+ \bar{\phi}^+ \Phi_3 + \phi^- \bar{\phi}^- \Phi_3 + \phi_i \bar{\phi}_i \Phi_3) \\ & + (\phi_1 \bar{\phi}_2 + \bar{\phi}_1 \phi_2) \Phi_4 + (\Phi_{12} \Phi_{23} \Phi_{31} + \Phi_{12} \phi^+ \bar{\phi}^- + \Phi_{12} \phi_i \bar{\phi}_i + h.c.) \Big] \quad (14) \end{aligned}$$

The common overall normalization is fixed by calculating the non-Abelian fermion-fermion-gauge boson coupling g of $SU(5)$ [19].

We now discuss the pattern of gauge symmetry breaking. The anomalous $U(1)_A$ (9a) is broken by the Dine-Seiberg-Witten mechanism [20], in which a potentially large Fayet-Iliopoulos [24] D-term is generated by the v.e.v. of the dilaton field. Such a D-term would, in general, break supersymmetry and destabilize the string vacuum unless, as in our case, there is a direction in the scalar potential $\phi = \sum \alpha_i \phi_i$ which is F-flat and also D-flat with respect to the non-anomalous gauge symmetries and in which $\sum Q_i^A |\alpha_i|^2 < 0$. If such a direction exists, it will acquire a v.e.v., cancelling the anomalous D-term, restoring supersymmetry [21] and stabilizing the vacuum. Since the fields corresponding to such a flat direction typically also carry charges for the non-anomalous gauge symmetries, this mechanism for cancelling the anomalous D-term will in general break some of these symmetries spontaneously [16,17]. (The same mechanism was subsequently used in orbifold constructions [25].)

In our string model, at least three and typically all four of the surplus $U(1)$'s are broken by this mechanism. The general form of the anomalous D-term is [20,26]

$$D_A = \frac{g e^{\Phi_D}}{192 \pi^2} \text{Tr } Q_A + \sum_i Q_A^i |\phi_i|^2 \quad (15)$$

where Φ_D is the shifted dilaton field $\langle \Phi_D \rangle = 0$, and $\text{Tr } Q_A = 182$ in our model. As an example, if we give v.e.v.'s to $\bar{\Phi}_{23}$, Φ_{31} , ϕ_{45} and ϕ^+ (which will be needed later together with other v.e.v.'s), which is seen to be an F-flat direction of the superpotential (14), we have the following conditions for vanishing of the D-terms:

$$\left. \begin{aligned} U(1)'_1 : D'_1 &= -|\bar{\Phi}_{23}|^2 - |\Phi_{31}|^2 + |\phi_{45}|^2 + 2|\phi^+|^2 = 0 \\ U(1)'_2 : D'_2 &= -3|\bar{\Phi}_{23}|^2 + |\Phi_{31}|^2 - |\phi_{45}|^2 + 2|\phi^+|^2 = 0 \\ U(1)'_3 : D'_3 &= -3|\bar{\Phi}_{23}|^2 - |\Phi_{31}|^2 + 6|\phi_{45}|^2 - |\phi^+|^2 = 0 \\ U(1)_A : D_A &= -3|\bar{\Phi}_{23}|^2 - 5|\Phi_{31}|^2 - 2|\phi^+|^2 + \frac{182}{192\pi^2}g = 0 \end{aligned} \right\} \quad (16a)$$

which have the solution

$$|\bar{\Phi}_{23}|^2 = \frac{1}{2} |\Phi_{31}|^2 = |\phi_{45}|^2 = |\phi^+|^2 = \frac{1}{15} \frac{91}{192 \pi^2} g \quad (16b)$$

that breaks all four $U(1)$'s^{*}). Notice that these v.e.v.'s are two orders of magnitude smaller than the Planck scale, justifying a perturbative treatment. These v.e.v.'s imply that the superstring vacuum is not really $SU(5) \times U(1)^5$, but actually minimal flipped $SU(5) \times U(1)$. Nevertheless, these extra $U(1)$'s constrain the allowed Yukawa couplings, forbidding in particular trilinear couplings that could give masses to the first two generations of quarks and leptons. This and $\langle \Phi, \phi \rangle < 1$ offer a viable framework [17] for understanding the hierarchical pattern of quark and lepton masses, which would be difficult to explain if there were trilinear Higgs couplings of similar magnitude for all three generations.

The cancellation of the anomalous D-term (15) and the other conditions of D-flatness (16a) constrain the choice of v.e.v.'s. However, there are more F-flat directions than constraints, leaving some freedom to choose v.e.v.'s in a phenomenologically realistic way. Specifically, if we assume zero v.e.v.'s for $(\bar{\Phi}_{12})$ and $\Phi_{3,4}$, then differentiation of the superpotential (14) with respect to these fields yields four constraints:

^{*}) One surplus $U(1)$ factor $U(1)_1 - 3U(1)_2 + U(1)_3 + 2U(1)_4$ is retained if only Φ_{31} , ϕ_{45} , $\bar{\phi}^-$ and possibly their conjugates acquire v.e.v.'s, in which case D-flatness requires that $\Delta_{31} + 1/3 \Delta_{45} - 1/3 \Delta^- \propto g$, where $\Delta_{31} \equiv |\Phi_{31}|^2 - |\bar{\Phi}_{31}|^2$, etc.

$$\left. \begin{aligned} |\phi_{45} \bar{\phi}_{45} + \phi^+ \bar{\phi}^+ + \phi^- \bar{\phi}^- + \sum_i \phi_i \bar{\phi}_i| &= 0 \\ |\bar{\Phi}_{23} \bar{\Phi}_{23} + \phi^+ \phi^- + \sum_i (\phi_i)^2| &= |\bar{\Phi}_{23} \bar{\Phi}_{31} + \bar{\phi}^+ \bar{\phi}^- + \sum_i (\bar{\phi}_i)^2| = 0 \\ |\phi_1 \bar{\phi}_2 + \phi_2 \bar{\phi}_1| &= 0 \end{aligned} \right\} \quad (17)$$

on the remaining 21 $SU(5) \times U(1)$ singlets. Differentiation with respect to these fields leads to no further non-trivial constraints. The four constraints (17) combined with the four D-flatness constraints leave considerable freedom to choose among the remaining $21-8 = 13$ flat directions. Some of these flat directions, e.g., involving the Neveu-Schwarz sector fields $\Phi_{23}, \bar{\Phi}_{23}, \Phi_{31}, \bar{\Phi}_{31}, \Phi_{1,2,5}$, are known [27] to remain flat to all orders in string dynamics, since they correspond to lattice deformations [3] and the associated Higgs phenomenon can be realized directly at the string level [28]. A more careful analysis [29] is required for the remaining fields from the b_4+b_5 sector, because they contain some Z_2 twist fields. The degeneracy among these various vacua will be broken by radiative corrections in the presence of soft supersymmetry-breaking terms, which we plan to analyze in a later paper [30]^{*)}.

Our choice of v.e.v.'s continues as follows. V.e.v.'s $\langle \bar{F}_5 \rangle = \langle F \rangle$: $F \equiv \sum_{i=1}^3 \alpha_i F_i$, $|\alpha| = 1$ is obviously needed to break $SU(5) \times U(1)$ to the Standard Model. The equality $\langle \bar{F}_5 \rangle = \langle F \rangle$ is imposed by D-flatness for $SU(5)$, whilst $\langle F_4 \rangle = 0$ because of the superpotential term $\bar{F}_5 F_4 \phi_3$, which provides a superheavy Dirac mass for the v_4^c . This guarantees a light left-handed neutrino in $\bar{f}_5$, and also protects the $\bar{f}_5$ and $\bar{h}_{45}$ doublets from getting a superheavy mass via the $F_4 \bar{f}_5 \bar{h}_{45}$ coupling. The precise linear combination α_i of the F_i that gets a v.e.v. is not determined by trilinear couplings. We will first discuss the simplest case: $\langle F_1 \rangle \neq 0$, then consider the effects of introducing v.e.v.'s for $F_{2,3}$.

With $\langle F_1 \rangle \neq 0$, all components of F_1 are eaten or acquire superheavy masses through the supersymmetric Higgs mechanism except for the $d_{F_1}^c$. The latter combines with the D_{h_1} triplet through the superpotential coupling $F_1 F_1 h_1$, splitting the h_1 triplet from the h_1 doublet, which will be one of the supersymmetric Weinberg-Salam Higgs doublets. The other Higgses h_2, h_3 acquire large Dirac mass terms $h_2 \bar{h}_3$ and $h_3 \bar{h}_2$ from v.e.v.'s for the $SU(5) \times U(1)$ singlets $\Phi_{23}, \bar{\Phi}_{23}$ respectively. A large v.e.v. for $\bar{\phi}_{45}$ adds a Dirac mass term $h_3 \bar{h}_{45}$, with the result that one linear combination

*) Naïve consideration of radiative corrections might have led one to expect that $\langle \Phi_{12}, \bar{\Phi}_{12}, \Phi_{3,4} \rangle \neq 0$, since these fields have the most Yukawa couplings. However these v.e.v.'s would leave too few flat directions to maintain D-flatness and cancel the anomalous D-term (15). Thus the dilaton D-term precludes some phenomenologically unacceptable vacua.

$$\bar{h} \equiv \cos\bar{\theta} \bar{h}_{45} + \sin\bar{\theta} \bar{h}_2 : \tan\bar{\theta} = - \langle \bar{\phi}_{45} \rangle / \langle \bar{\phi}_{23} \rangle \quad (18)$$

remains light. The $\bar{h}$ triplets $\bar{D}_h$ receive superheavy masses from the doublet-triplet splitting term $\bar{F}_5 \bar{F}_5 \bar{h}_2$. We should stress that this mechanism [9] leaves no dangerous residual dimension-five baryon-number-violating operators [13], since it only generates Dirac mass terms $d_{F D_h}^C, \bar{d}_{F \bar{D}_h}^C$. The surviving trilinear superpotential couplings involving quarks and leptons are just:

$$g\sqrt{2} \cos\bar{\theta} F_4 \bar{F}_5 \bar{h} , \quad g\sqrt{2} F_4 F_4 h_1 , \quad g\sqrt{2} \bar{F}_1 \bar{L}_1 h_1 \quad (19)$$

where it is understood that only the Higgs doublets in $\bar{h}, h_1$ survive^{*}). In addition, we have the complete Higgs multiplets h_{12} and $\bar{h}_1$, but these do not couple to matter fields, and therefore do not mediate rapid baryon decay or flavour-changing neutral currents.

Let us summarize where we now stand. Without any recourse to non-renormalizable interactions, we have achieved natural doublet-triplet splitting with no residual dimension-five baryon-number-violating operators. We have found masses for the third-generation quarks and leptons: $(t_L, b_L, b_R) \in F_4, t_R \in \bar{F}_5, (v_\tau, \tau_L) \in \bar{F}_1$ and $\tau_R \in L_1^C$, with all other quarks and leptons massless at the level of trilinear superpotential interactions. The v_τ remains light thanks to a seesaw mass term provided by $\bar{F}_5 F_4 \phi_3$. We have recovered the successful GUT prediction $m_b = m_\tau$ at $Q^2 = m_{GUT}^2$, and there are no tree-level flavour-changing neutral currents, since only one pair of Higgs doublets couples to fermions.

The top quark mass term in (13) deserves more discussion. The relation $\lambda_t = g\sqrt{2} \cos\bar{\theta}$ is valid at the GUT scale, and the Yukawa coupling λ_t and the gauge coupling g are renormalized below this scale in the usual way [13]:

$$\frac{d\lambda_t}{dt} = \frac{\lambda_t}{8\pi^2} \left[3\lambda_t^2 + \frac{1}{2}\lambda_b^2 + \frac{1}{2}\lambda_h^2 - \frac{8}{3}g_3^2 - \frac{3}{2}g_2^2 - \frac{3}{5}\left(\frac{13}{18}\right)g_Y^2 \right], \text{ etc.} \quad (20)$$

where λ_h is an effective $h\bar{h}\Phi$ mixing coupling, and the U(1) coupling g_Y is normalized: $g_3 = g_2 = g_Y$ at $Q = m_{GUT}$. The right-hand side of (20) vanishes when

$$\lambda_t^2 = \left(\frac{1}{3} \text{ or } \frac{1}{4} \right) \left(\frac{8}{3}g_3^2 + \frac{3}{2}g_2^2 + \frac{13}{30}g_Y^2 \right) \equiv \lambda_{IR}^2 \quad (21)$$

^{*}) We will see later that the physical charge -1/3 quark and charged lepton mass terms are suppressed relative to those in (19) by a common mixing angle $\cos\bar{\theta}$ analogous to $\cos\theta$.

depending on whether we assume $\lambda_t \gg \lambda_{b,h}$ or $\lambda_t = \lambda_{b,h}$. The value (21) of λ_t is an effective infra-red fixed point [31], and if $\lambda_t < \lambda_{t_{IR}}$ at the GUT scale, it remains so at any lower renormalization scale. Thus, we have in general the upper bound

$$m_t = \lambda_t \bar{v} < \lambda_t \sqrt{v^2 + \bar{v}^2} = \frac{\lambda_t \sqrt{2}}{g_2} m_W < \frac{\sqrt{2} m_W}{g_2} \lambda_{t_{IR}} \quad (22a)$$

where $\langle \bar{h} \rangle \equiv \langle \bar{v} \rangle$. Using Eq. (21) at the renormalization scale $Q = m_t \approx m_W$ we obtain

$$m_t < m_W \sqrt{\frac{2}{3}} \left[\frac{8}{3} \left(\frac{g_3}{g_2} \right)^2 + \frac{3}{2} + \frac{13}{18} \tan^2 \theta_W \right]^{\frac{1}{2}} \approx 220 \text{ GeV} \quad (22b)$$

In fact, the Figure shows that $\lambda_t(m_W)$ is always significantly below $\lambda_{t_{IR}}(m_W)$, even if $\lambda_t > \lambda_{t_{IR}}$ at the GUT scale. This is because the renormalization group scaling of λ_t cannot keep pace with the growth of g_3 and hence $\lambda_{t_{IR}}$ at scales $Q \sim m_W$. Thus we find an actual upper bound of

$$m_t < 190 \text{ GeV} \quad (23)$$

which is independent of λ_t at m_{GUT} and is valid in any supersymmetric theory. If we now incorporate the extra string information (19), for $\cos \bar{\theta} < 1$ we find only a slight improvement of this bound, as seen in the Figure.

This upper bound is apparently a generic feature of four-dimensional string models, at least in their fermionic formulation. It is to be compared with the experimental lower bound $m_t \gtrsim 60 \text{ GeV}$ from $\bar{p}p$ collider experiments [32] and the upper bound $m_t \lesssim 200 \text{ GeV}$ from precision neutral current experiments [33].

Another interesting consequence of the fermion mass terms (19) is that the usual quark-lepton kinship rules are violated. The heaviest quarks and leptons belong to different $SU(5)$ representations $\bar{f}_5, \bar{f}_1$. One potentially dramatic consequence is that the baryon decay modes $p \rightarrow \mu^+ \pi^0$ or $\bar{\nu} \pi^0$, $n \rightarrow \mu^+ \pi^-$ or $\bar{\nu} \pi^0$ could dominate over $p \rightarrow e^+ \pi^0$, $n \rightarrow e^+ \pi^-$. Since the t^c is assigned to $\bar{f}_5$, the final-state u^c quark must belong to $\bar{f}_1$ or $\bar{f}_2$. (As we will see later, $\bar{f}_3$ acquires a large mass.) If $u^c \in \bar{f}_1$, then the dominant modes will be $p \rightarrow \bar{\nu}_\tau \pi^+$, $n \rightarrow \bar{\nu}_\tau \pi^0$ since $(\bar{\nu}_\tau, \tau_L) \in \bar{f}_1$. If $u^c \in \bar{f}_2$, the final-state leptons will be some linear combination of $(\bar{\nu}_e, e^+)$ and $(\bar{\nu}_\mu, \mu^+)$, creating the possibility of large $p \rightarrow \mu^+ \pi^0$, $n \rightarrow \mu^+ \pi^-$ branching ratios. The baryon lifetime should be similar to the estimate of Ref. [13], namely $2 \times 10^{34 \pm 2} \text{ y}$. In addition, the lepton flavour-changing decay $\mu \rightarrow e \gamma$ may be greatly enhanced relative to conventional supersymmetric GUTs [34]. In $SU(5)$, for example, $\mu \rightarrow e \gamma$

results from generation-changing slepton masses $\delta m_{\mu e}^2$ induced by renormalization effects between m_P and m_{GUT} . These have the form $\delta m_{\mu e}^2 \propto (K^*_{32} \lambda_t^2 K_{31}) \ln(m_{GUT}/m_P)$, where the combination $K^*_{32} K_{31}$ of Cabibbo-Kobayashi-Maskawa mixing angles is known experimentally to be $< 10^{-3}$, so that $B(\mu \rightarrow e\gamma) \sim 10^{-14}$ to 10^{-17} in conventional supersymmetric SU(5). However, in our flipped SU(5)×U(1) string model [35], since the usual quark-lepton kinship rules break down, one of $K_{32,31}$ is expected to be $O(1)$, whilst the other is suppressed by a single Cabibbo-like factor $\sin\Phi$. The branching ratio is therefore enhanced relative to SU(5) by the ratio $(\sin\Phi/10^{-3})^2 \sim 10^3$, giving $B(\mu \rightarrow e\gamma) \sim 10^{-11}$ to 10^{-14} in the range accessible to the MEGA experiment.

We will now briefly consider the possible effects of non-renormalizable superpotential couplings [17]. Since the selection rules governing such terms are incompletely understood, we present here a preliminary analysis of some relevant terms, leaving a more systematic study for elsewhere [29,30]. Interestingly enough, even the simplest non-renormalizable couplings consistent with gauge invariance and some discrete string symmetries are sufficient for phenomenological acceptability.

The previously light Higgs multiplets h_1, h_{45} can mix with the massive h_2 via v.e.v.'s for ϕ_{45} and $\bar{\Phi}_{31}$. The light Higgs eigenstates become the two linear combinations orthogonal to the massive combination $\langle \bar{\Phi}_{31} \rangle h_1 + \langle \Phi_{23} \rangle h_2 + \langle \phi_{45} \rangle h_{45}$. One combination of light Higgses can be removed by a coupling of the form

$$h_2 \bar{h}_1 < \phi^+ \phi^- > \quad (24)$$

where at least two of $\langle \Phi_{23}, \Phi_{31} \rangle$, $\langle \phi^+, \phi^- \rangle$, $\langle \phi_i \rangle$ must be non-zero to preserve F-flatness (17). The surviving Weinberg-Salam Higgs doublets are then $h \equiv \cos\theta h_1 + \sin\theta h_{45}$: $\tan\theta \equiv -\langle \bar{\Phi}_{31} \rangle / \langle \phi_{45} \rangle$ and $\bar{h}$ defined in Eq. (18). The b and τ Yukawa couplings $F_4 F_4 h$, $\bar{f}_1 l_1^c h$ are now suppressed relative to (19) by $\cos\theta$, and $\lambda_t > \lambda_{b,\tau}$ at the GUT scale if $\cos\bar{\theta} > \cos\theta$. This difference need not be great, since through renormalization they generate $\tilde{m}_h^2 < \tilde{m}_h^2$, therefore $\bar{v} > v$, and hence $m_t = \lambda_t \bar{v} \gg m_b = \lambda_b v$ [13]. Mixing between the light Higgs doublets $h, \bar{h}$ through an effective $h\bar{h}\phi$ interaction may result from the charge conjugate of (24). In this case only one, not both, of $\bar{\phi}^{+,-}$ should get a large v.e.v., to avoid a superheavy mass for $h\bar{h}$. Such an asymmetry between $\langle \bar{\phi}^{+,-} \rangle$ and $\langle \phi^{+,-} \rangle$ is in fact suggested by the cancellation [21] of the anomalous D-term (15) [20].

Finally, we consider the effects of v.e.v.'s for $F_{i=2,3}$. The precise linear combination α_i of the F_i that gets a v.e.v. is not determined by trilinear couplings, but may be fixed by non-renormalizable terms of the form $\bar{F}_5 F_i \phi < \dots >$, which would provide the remaining seesaw neutrino masses and stabilize the vacuum against v.e.v.'s for two linear combinations of the F_i . If $\langle F_2 \rangle \neq 0$, the trilinear coupling $F_2 F_2 h_2$ results in a doublet-triplet splitting term for $d_{F_2}^c D_{h_2}$. However, because of the non-renormalizable $h_2 \bar{h}_1$ mass term (24), one linear combination of

$d_{F_2}^c$ and $\bar{D}_{h_1}$ remains light. Since these states have identical Standard Model quantum numbers, and since $\bar{h}_1$ has no Yukawa couplings, this mixing has no immediate observable effects.

A v.e.v. for F_3 would permit a mass term of the form

$$(f_4 \bar{f}_3 \text{ or } \bar{l}_4^c l_3^c) \langle \bar{F}_5 F_3 \rangle \langle \bar{\phi}_{45} \rangle \quad (25)$$

which would eliminate all surplus matter fields $f_4, \bar{l}_4^c, \bar{f}_3$ and l_3^c from the light spectrum. This term also protects the light h doublet from being eaten by the mass term $\langle F_5 \rangle f_4 h_{45}$: the light h becomes a linear combination of h_1, h_{45} and $\bar{f}_3$. Since the $\bar{f}_3$ doublet has the same quantum numbers as the h , and has no couplings to light fields, this mixing has no immediate effect. Likewise, the mass term $\langle F_3 \rangle \bar{f}_3 \bar{h}_3$ plays no major rôle, since the fields $\bar{f}_3, h_3$ are already massive in this scenario.

It is easy to check that the large v.e.v.'s introduced above, namely

$$\langle \bar{\Phi}_{23}, \bar{\Phi}_{23}, \phi_{45}, \bar{\phi}_{45}, \bar{\Phi}_{31}, \phi^+ \phi^- \text{ with } \bar{\Phi}_{31} \text{ and/or } \phi_i, \bar{\phi}^+ \text{ or } \bar{\phi}^- \rangle \quad (26)$$

can cancel the anomalous D-term whilst preserving F-flatness and D-flatness for the remaining gauge symmetries. These v.e.v.'s combined with $\langle F_{1,2,3} \rangle$ are sufficient to provide masses and couplings for all three generations of quarks and leptons via non-renormalizable terms along the lines of Ref. [17]. Note that the smallness of the off-diagonal Cabibbo-Kobayashi-Maskawa matrix elements V_{ts}, V_{td} is guaranteed by the fact that the t_L and $b_{L,R}$ quarks belong to the same SU(5) multiplet F . A more complete analysis of non-renormalizable couplings is in progress [30].

Thus far, we have said little about the hidden sector matter fields. There are two classes of such states: those in class (e) which are SU(5)×U(1) invariant and couple to the observable sector only via U(1)⁴ gauge interactions and gravitational-strength couplings and are hence phenomenologically harmless, and those in class (f) that carry electric charges $\pm \frac{1}{2}$. It is easy to see that the v.e.v.'s (26) give large masses to all except two pairs of these charged states via trilinear couplings. The remainder would get large masses from dimension-five operators. Should these be absent, the remaining charge $\pm \frac{1}{2}$ states are in any case confined into integer-charged states by hidden SU(4) interactions that become strong at a scale $\Lambda \gg m_W$, so there are no light stable fractionally-charged states.

In this paper we have derived an improved version of our three-generation flipped SU(5)×U(1) string model [16,17] in which there are non-vanishing trilinear superpotential terms giving $m_t, m_b, m_\tau \neq 0$. We recover the successful GUT relation $m_b \approx 3m_\tau$ and find that $m_t < 190$ GeV after renormalization. The low-energy spectrum is completely free of gauge and mixed gravitational anomalies, and all surplus U(1) gauge factors are automatically broken by the Fayet-Iliopoulos [24] D-term genera-

ted by the Dine-Seiberg-Witten [20] mechanism. Baryon decay is automatically suppressed so that $\tau_B \approx 2 \times 10^{34 \pm 2} \text{y}$ [13], with the likelihood of unusual decay modes such as $\bar{\nu}\pi$ or $\mu^+\pi$ dominating. There are no tree-level FCNC's, but slepton mixing is likely to induce observable $\mu \rightarrow e\gamma$ decay at the one-loop level [35]. Plausible hypotheses about non-renormalizable superpotential terms yield the minimal flipped $SU(5) \times U(1)$ matter spectrum with non-zero light fermion masses and CKM flavour mixing. We believe we have identified a phenomenologically viable string vacuum.

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FIGURE CAPTION

The physical top quark Yukawa coupling $\lambda_t(m_W)$ as a function of $\lambda_t(m_{GUT})$ in the minimal supersymmetric Standard Model with light two Higgs doublets. Even for $\lambda_t(m_{GUT}) > \lambda_{t_{IR}}(m_{GUT})$, we find that $\lambda_t(m_W) < \lambda_{t_{IR}}(m_W)$, corresponding to $m_t < 190$ GeV. The upper bound on $\lambda_t(m_{GUT})$ from four-dimensional string models gives a slightly stronger upper bound on $\lambda_t(m_W)$.