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Phenomenology of 3-Family Grand Unified String Models

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Abstract

In the 3-family grand unified string models constructed so far, there is only one adjoint (and no higher dimensional representation) Higgs field in the grand unified gauge group. In this preliminary analysis, we address the proton-decay problem in the 3-family E_6 and related $SO(10)$ string models. In particular, we analyze the doublet-triplet splitting (within certain assumptions about non-perturbative dynamics). It appears that generically some fine-tuning is necessary to arrange for a pair of Higgs doublets to be light, while having color Higgs triplets superheavy. We also discuss charge-2/3 quark mass matrix that generically also seems to require some fine-tuning to have rank 1.

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One elegant way for superstring theory to include the standard model in its low energy effective theory [1] is via extensively studied supersymmetric (SUSY) grand unified theory (GUT) [2]. Spontaneous symmetry breaking of the grand unified gauge group to the standard model in the low energy effective field theory requires at least an adjoint (or other appropriate higher dimensional representation) Higgs field in the light mass spectrum. This requirement imposes a strong constraint on GUT string model-building. In addition, dynamical SUSY breaking via gaugino condensation requires an asymptotically-free hidden sector (SUSY breaking is then transmitted to the observable sector via interactions with gravity or other messenger/intermediate sector [3]). Recently, two of us have constructed grand unified string models that satisfy the above requirements and have three chiral families of fermions in the GUT gauge group [4]. In fact, a classification of such 3-family grand unified models was carried out within the framework of conformal (free) field theory and asymmetric orbifolds. The form of the superpotentials for these models was determined recently [5]. In this letter, we give a preliminary analysis of the phenomenological properties of some of these models. More specifically, we analyze the proton decay problem and the fermion mass matrix in the 3-family E_6 and the related $SO(10)$ string models.

The E_6 and $SO(10)$ GUTs [6] are well studied. Here, the string models are very specific realizations of these GUTs. Within the string model-building framework described above, we find only one 3-family E_6 model. In this paper, we make a preliminary study of the phenomenological properties of this model and the closely related $SO(10)$ string model. They have 5 right-moving and 2 left-moving fermion families. Upon generation–antigeneration pairings, one is left with three chiral families. Analysis shows that one pairing takes place at around the M_{GUT} scale, while the other pairing occurs at a (perhaps, much) lower scale.

The problem of color Higgs mediated proton decay [7] has a simple solution in the framework of $SO(10)$ models. This is achieved by giving appropriate *vev* to the adjoint (for example, via the Dimopoulos–Wilczek mechanism [8]). In the model we considered, the Higgs doublets and triplets generically become massive at the scale slightly below M_{GUT} (with the exception of a pair of doublets and a pair of triplets which may become massive at a much lower scale; the latter triplets are harmless for proton stability, however). As the string models have only one adjoint Higgs field, one can arrange for a pair of Higgs doublets to be light, while having color Higgs triplets heavy. This, however, requires certain fine-tuning since under spontaneous symmetry breaking of E_6 to the standard model, all Higgs doublets and triplets are heavy in a generic point in the moduli space. To obtain one light Higgs doublet, the moduli have to satisfy one constraint. In terms of the moduli entering the Higgs doublet mass matrix (which is deduced from the approximate superpotential used in this paper), this means a light Higgs doublet exists in a 10-dimensional sub-manifold in the original 11-dimensional moduli space. This is fine-tuning.

We should emphasize that our conclusions rely heavily on important assumptions about the non-perturbative string dynamics, an issue we shall only comment on briefly. Furthermore, some mixings of fields are ignored in this preliminary analysis.

To begin, let us look at the massless spectra of the E_6 and $SO(10)$ models [4] given in the following table. The gravity, dilaton and gauge supermultiplets are not shown. The E_6 model has gauge symmetry $SU(2)_1 \otimes (E_6)_3 \otimes U(1)^3$, while the $SO(10)$ model has gauge symmetry $SU(2)_1 \otimes SO(10)_3 \otimes U(1)^4$. The subscripts label the levels of the current algebra realizations of the gauge symmetries. In practice, this means that the fine structure constant of the level-

3 GUT gauge symmetry is 1/3 that of the level-1 $SU(2)$ gauge symmetry. The $U(1)$ charge of a particle is given by its integer value in the table multiplied by the normalization provided at the bottom of the table. Note that all massless particles with E_6 quantum numbers are singlets under $SU(2)$, *i.e.*, this $SU(2)$ plays the role of a hidden sector. The $U(1)$s play the role of intermediate/messenger symmetries.

$E1$ $SU(2)\otimes E_6\otimes U(1)^3$	Field	$T1(1,1)$ $SU(2)\otimes SO(10)\otimes U(1)^4$	Field
$(\mathbf{1},\mathbf{78})(0,0,0)_L$ $(\mathbf{1},\mathbf{1})(0,+6,0)_L$ $2(\mathbf{1},\mathbf{1})(0,-3,\pm 3)_L$	$\widehat{\Phi}$ U_0 $U_{+\pm},U_{-\pm}$	$(\mathbf{1},\mathbf{45})(0,0,0,0)_L$ $(\mathbf{1},\mathbf{1})(0,0,0,0)_L$ $(\mathbf{1},\mathbf{1})(0,+6,0,0)_L$ $2(\mathbf{1},\mathbf{1})(0,-3,\pm 3,0)_L$	Φ ϕ U_0 $U_{+\pm},U_{-\pm}$
$(\mathbf{1},\mathbf{27})(0,-2,0)_L$ $2(\mathbf{1},\mathbf{27})(0,+1,\pm 1)_L$	χ_0 $\chi_{+\pm},\chi_{-\pm}$	$(\mathbf{1},\mathbf{16})(0,-2,0,-1)_L$ $(\mathbf{1},\mathbf{10})(0,-2,0,+2)_L$ $(\mathbf{1},\mathbf{1})(0,-2,0,-4)_L$ $2(\mathbf{1},\mathbf{16})(0,+1,\pm 1,-1)_L$ $2(\mathbf{1},\mathbf{10})(0,+1,\pm 1,+2)_L$ $2(\mathbf{1},\mathbf{1})(0,+1,\pm 1,-4)_L$	$\mathbf{Q}_0$ $\mathbf{H}_0$ $\mathbf{S}_0$ $\mathbf{Q}_{+\pm},\mathbf{Q}_{-\pm}$ $\mathbf{H}_{+\pm},\mathbf{H}_{-\pm}$ $\mathbf{S}_{+\pm},\mathbf{S}_{-\pm}$
$(\mathbf{1},\overline{\mathbf{27}})(\pm 1,-1,0)_L$	$\widetilde{\chi}_\pm$	$(\mathbf{1},\overline{\mathbf{16}})(\pm 1,-1,0,+1)_L$ $(\mathbf{1},\mathbf{10})(\pm 1,-1,0,-2)_L$ $(\mathbf{1},\mathbf{1})(\pm 1,-1,0,+4)_L$	$\widetilde{\mathbf{Q}}_\pm$ $\widetilde{\mathbf{H}}_\pm$ $\widetilde{\mathbf{S}}_\pm$
$(\mathbf{2},\mathbf{1})(0,0,\pm 3)_L$ $(\mathbf{1},\mathbf{1})(\pm 3,+3,0)_L$	$D_\pm$ $\widetilde{U}_\pm$	$(\mathbf{2},\mathbf{1})(0,0,\pm 3,0)_L$ $(\mathbf{1},\mathbf{1})(\pm 3,+3,0,0)_L$	$D_\pm$ $\widetilde{U}_\pm$
$r_{U(1)}:\ (1/\sqrt{6},\ 1/3\sqrt{2},\ 1/\sqrt{6})$		$(1/\sqrt{6},\ 1/3\sqrt{2},\ 1/\sqrt{6},\ 1/6)$	

The E_6 model has 5 left-handed $\mathbf{27}$s: χ_0, χ_{++}, χ_{--}, χ_{+-} and χ_{-+}, and 2 left-handed $\overline{\mathbf{27}}$s: $\widetilde{\chi}_+$ and $\widetilde{\chi}_-$. Their lowest order non-vanishing couplings in the superpotential are [5]

$$
\begin{aligned}
W &= \lambda_0\, U_0\,(U_{++}U_{--}+U_{+-}U_{-+}) + \lambda_1\, \chi_0\,(\chi_{++}\chi_{--}+\chi_{+-}\chi_{-+})\,\widehat{\Phi} \\
&+ \lambda_2\, \widetilde{\chi}_+\widetilde{\chi}_-\,(\chi_{++}\chi_{--}+\chi_{+-}\chi_{-+}) + \lambda_3\left(\chi^3_{++}U_{--}+\chi^3_{+-}U_{-+}+\chi^3_{-+}U_{+-}+\chi^3_{--}U_{++}\right)\widehat{\Phi}^2 \\
&+ \lambda_4\, \chi_0^3\, U_0 D_+ D_- \widehat{\Phi}^2 + \lambda_5\left(\widetilde{\chi}^3_+\widetilde{U}_- + \widetilde{\chi}^3_-\widetilde{U}_+\right) D_+ D_- \widehat{\Phi} \ +\ \ldots
\end{aligned}
\tag{1}
$$

where the traces over the irreps of the gauge group are implicit here. The λ_k are polynomials of $\widehat{\Phi}^3$, with $\lambda_k[0] \neq 0$. In principle, one can extract the $\lambda_k[\widehat{\Phi}^3]$ polinomials from tree level string scatterings; they may also be determined using string modular invariant properties in a much simpler way. If all the *vevs* are below the string scale it is a reasonable approximation to ignore the higher order terms in the superpotential. It is clear that the gauge discrete symmetries in superstring theory impose powerful selection rules on the couplings. *A priori*, the superpotential may contain other gauge invariant terms such as the three-point couplings $\chi_0\chi_{\alpha+}\chi_{\beta-}$ where $\alpha, \beta = \pm$. However, these terms are forbidden by the stringy discrete symmetries of the model [5]. Similarly, couplings such as $\chi_0\chi_{++}\chi_{+-}\widehat{\Phi}$ are not allowed even though they contain gauge singlets.

The spectra of the $SO(10)$ and E_6 models are very similar. In particular, the spectrum of the $SO(10)$ model is the same as that of the E_6 model with a non-zero *vev* of the adjoint $\widehat{\Phi}$ of $(E_6)_3$ such that $(E_6)_3$ is broken down to $SO(10)_3 \otimes U(1)$ (the last $U(1)$ in the $SO(10)$ part of the table). Thus, the $\mathbf{27}$ of $(E_6)_3$ branches into $\mathbf{16}(-1)+\mathbf{10}(+2)+\mathbf{1}(-4)$ of $SO(10)_3 \otimes U(1)$ (denoted as $\mathbf{Q}$, $\mathbf{H}$ and $\mathbf{S}$ respectively). The adjoint $\mathbf{78}$ branches into $\mathbf{45}(0)+\mathbf{1}(0)+\mathbf{16}(+3)+\overline{\mathbf{16}}(-3)$. Note that while the $\mathbf{45}(0)$ and $\mathbf{1}(0)$ are present in the $SO(10)$ model, the $\mathbf{16}(+3)$ and $\overline{\mathbf{16}}(-3)$ are missing because of the super-Higgs mechanism.

From the above discussion, it is clear that to get the superpotential for the $SO(10)$ model from that of the E_6 model, we can simply replace χs by $(\mathbf{Q}+\mathbf{H}+\mathbf{S})$s (and similarly for $\widetilde{\chi}$s), and $\widehat{\Phi}$ by $\Phi+\phi$ in (1). The allowed 4-point couplings in the $SO(10)$ superpotential are:

$$
\begin{aligned}
W = {} & \lambda_{11}\left\{\mathbf{S}_0(\mathbf{H}\cdot\mathbf{H}) + \mathbf{H}_0(\mathbf{S}\cdot\mathbf{H})\right\}(\Phi+\phi) + \lambda_{12}\left\{\mathbf{H}_0(\mathbf{Q}\cdot\mathbf{Q}) + \mathbf{Q}_0(\mathbf{H}\cdot\mathbf{Q})\right\}(\Phi+\phi) \qquad (2) \\
& + \lambda_{21}\,\tilde{\mathbf{S}}_+\tilde{\mathbf{S}}_-(\mathbf{S}\cdot\mathbf{S}) + \lambda_{22}\,\widetilde{\mathbf{Q}}_+\widetilde{\mathbf{Q}}_-(\mathbf{Q}\cdot\mathbf{Q}) + \lambda_{23}\left(\tilde{\mathbf{S}}_+\widetilde{\mathbf{Q}}_- + \tilde{\mathbf{S}}_-\widetilde{\mathbf{Q}}_+\right)(\mathbf{S}\cdot\mathbf{Q}) + \\
& + \lambda_{24}\left(\tilde{\mathbf{S}}_+\widetilde{\mathbf{H}}_- + \tilde{\mathbf{S}}_-\widetilde{\mathbf{H}}_+\right)(\mathbf{S}\cdot\mathbf{H}) + \lambda_{25}\left(\widetilde{\mathbf{H}}_+\widetilde{\mathbf{Q}}_- + \widetilde{\mathbf{H}}_-\widetilde{\mathbf{Q}}_+\right)(\mathbf{H}\cdot\mathbf{Q}) \\
& + \lambda_{26}\,\widetilde{\mathbf{H}}_+\widetilde{\mathbf{H}}_-(\mathbf{H}\cdot\mathbf{H}) + \lambda_{27}\left\{\widetilde{\mathbf{Q}}_+\widetilde{\mathbf{Q}}_-(\mathbf{S}\cdot\mathbf{H}) + (\tilde{\mathbf{S}}_+\widetilde{\mathbf{H}}_- + \tilde{\mathbf{S}}_-\widetilde{\mathbf{H}}_+)(\mathbf{Q}\cdot\mathbf{Q})\right\} + \ldots
\end{aligned}
$$

Here, we introduce the following notation:

$$
(\mathbf{A}\cdot\mathbf{B}) \equiv \mathbf{A}_{++}\mathbf{B}_{--} + \mathbf{B}_{++}\mathbf{A}_{--} + \mathbf{A}_{+-}\mathbf{B}_{-+} + \mathbf{B}_{+-}\mathbf{A}_{-+} \qquad (3)
$$

where A and B could be any of the fields $\mathbf{S}, \mathbf{H}, \mathbf{Q}$. Only the terms originating from the λ_1- and λ_2–terms in the E_6 superpotential (1) are shown. The couplings λ_{1i} stem from the λ_1–term in the E_6 superpotential and are *a priori* different because of the Clebsch–Gordan coefficients which arise upon breaking E_6 to $SO(10) \otimes U(1)$. The same is true for the λ_{2i} couplings.

Classically, the adjoint Higgs Φ is a flat modulus, so it can develop an arbitrary *vev*. However, it is likely that gaugino condensate will lift this flatness. Let us consider the following rather standard scenario [3]. With only one pair of chiral doublets $D_\pm$, the gauge group $SU(2)$ is asymptotically-free. Assuming that $\alpha_{GUT} \approx 1/24$, we have $\alpha_{st} = \alpha_{SU(2)} = 3\,\alpha_{GUT} \approx 1/8$. (This factor of 3 difference reflects the level–3 current algebra realization of the GUT symmetry). In string theory, the coupling α_{st} is not a free parameter, but is determined by the expectation value of the dilaton. Gaugino condensation in the $SU(2)$ hidden sector generates an effective scalar potential which, hopefully, stabilizes the dilaton expectation value in a self-consistent way. (Since globally $N=1$ supersymmetric $SU(2)$ with one flavor does not have a quantum vacuum, one might hope that including non-perturbative corrections to the Kähler potential [9] would stabilize the moduli and the dilaton). Here we shall simply make this assumption and consider the consequences. The scale of gaugino condensation is given by

$$\mu^2 = M_{st}^2 \exp\left(\frac{-4\pi}{b_0\alpha_{st}}\right) \approx \left(10^{13} GeV\right)^2 \tag{4}$$

where $b_0 = 6-1 = 5$ and the string scale $M_{st} \approx 10^{18}\ GeV$. This yields the supersymmetry breaking scale in the visible sector [10]

$$m_S \approx \frac{\mu^3}{M_{st}^2} \approx (10 \sim 100)\ TeV \tag{5}$$

It is worth noticing that m_S is very sensitive to α_{st}. The above estimate is tightly related to the fact that the GUT symmetry is realized at level 3.

Due to the presence of the Kähler potential $K(z,\bar{z})$ in the F-term of the effective scalar potential V_F:

$$V_F(z,\bar{z}) = e^{K/M_{pl}^2}\left(D_z W K_{z\bar{z}}^{-1}\overline{D_z W} - 3\frac{|W|^2}{M_{pl}^2}\right) \tag{6}$$

(where $D_zW \equiv W_z + WK_z/M_{pl}^2$, with M_{pl} being the Planck mass) the gaugino condensate will generate an effective potential for the adjoint Higgs field. The height of this potential term is rather small, of the order of m_S^4. Dimensional arguments suggest that the $SO(10)$ adjoint Higgs may develop string (or GUT) scale *vev*; at the same time, the adjoint Higgs fields of the remaining gauge symmetry (such as QCD) will have masses of the order m_S. These relatively light Higgs fields are a generic prediction of string GUT models. (There may exist model-dependent mechanisms for generating intermediate masses for the adjoints, but we will not address this issue here). Because of the difficulties in the analysis of dynamical symmetry breaking, we shall simply treat the adjoint Higgs as flat moduli.

Let us consider the possible scenarios for breaking the $SO(10)$ down to that of the standard model, *i.e.*, $SU(3)_c \otimes SU(2)_w \otimes U(1)_Y$. There is one adjoint Higgs field in the massless spectrum of the model. The adjoint must break $SO(10)$ down to its subgroups that contain $SU(3)_c \otimes SU(2)_w \otimes U(1)_Y$. Here we are interested in the breaking $SO(10) \supset SU(3)_c \otimes SU(2)_w \otimes U(1)_Y \otimes U(1)$ by turning on the *vev* of the adjoint $\Phi = \epsilon \otimes \mathrm{diag}(a,a,a,b,b)$, where ϵ is a 2×2 antisymmetric matrix and $a, b \neq 0$, $a \neq b$.

The adjoint Higgs field of $SO(10)$ does not carry any other gauge quantum numbers, so there is no danger of mixing the $SO(10)$ quantum numbers with other gauge quantum numbers in the adjoint breaking. Since the adjoint by itself cannot break the $SO(10)$ gauge symmetry down to that of the standard model, some other fields must acquire *vev*s to achieve further breaking of the gauge symmetry. These fields must carry $SO(10)$ quantum numbers, and contain $SU(3)_c \otimes SU(2)_w$ singlets, *i.e.*, they must come from $\mathbf{16} + \overline{\mathbf{16}}$ of $SO(10)$.

Let us denote the values of the *vev*s of $\mathbf{S}_0, \mathbf{S}_i, \mathbf{Q}_0, \mathbf{Q}_i, \widetilde{\mathbf{S}}_\pm$ and $\widetilde{\mathbf{Q}}_\pm$ as $s_0, s_i, q_0, q_i, \widetilde{s}_\pm$ and $\widetilde{q}_\pm$ correspondingly. From the D-term of the effective scalar potential corresponding to a $U(1)$ gauge group:

$$V_D = \alpha_{st}\pi(\sum_A X_A\langle\phi_A\rangle^2)^2 \tag{7}$$

(where X_A is the charge of field ϕ_A), the D-flatness condition follows:

$$q_0^2 + q^2 = \tilde{q}^2 \ , \qquad \text{where} \quad q \equiv \sqrt{\sum_i q_i^2} \ , \ \ \tilde{q} \equiv \sqrt{\tilde{q}_+^2 + \tilde{q}_-^2} \ , \tag{8}$$

This implies at least one of the $\widetilde{\mathbf{Q}}_\pm$ *vevs* must be non-zero. Thus, the $\mathbf{16} + \overline{\mathbf{16}}$ breaking scale is given by $\tilde{q}$.

To analyze the $SU(3)_3 \otimes SU(2)_3 \otimes U(1)_Y \otimes U(1)$ spectrum of the model, let us consider the following branchings of the $SO(10)$ irreps **45**, **16** and **10** under the breaking $SO(10)_3 \supset SU(3)_3 \otimes SU(2)_3 \otimes U(1)_Y \otimes U(1)$:

$$\begin{aligned}
\mathbf{45} = \ & (\mathbf{1},\mathbf{1})(0,0) + [(\mathbf{1},\mathbf{1})(0,0) + (\mathbf{8},\mathbf{1})(0,0) + (\mathbf{1},\mathbf{3})(0,0) + (\mathbf{3},\mathbf{2})(-5,0) + (\overline{\mathbf{3}},\mathbf{2})(+5,0)] \\
& + [(\mathbf{1},\mathbf{1})(+6,+4) + (\overline{\mathbf{3}},\mathbf{1})(-4,+4) + (\mathbf{3},\mathbf{2})(+1,+4)] \\
& + [(\mathbf{1},\mathbf{1})(-6,-4) + (\mathbf{3},\mathbf{1})(+4,-4) + (\overline{\mathbf{3}},\mathbf{2})(-1,-4)] \ ,
\end{aligned} \tag{9}$$

$$\begin{aligned}
\mathbf{16} = \ & \overset{q}{(\mathbf{1},\mathbf{1})(0,-5)} + [\overset{D}{(\overline{\mathbf{3}},\mathbf{1})(+2,+3)} + \overset{L}{(\mathbf{1},\mathbf{2})(-3,+3)}] \\
& + [\overset{E}{(\mathbf{1},\mathbf{1})(+6,-1)} + \overset{U}{(\overline{\mathbf{3}},\mathbf{1})(-4,-1)} + \overset{Q}{(\mathbf{3},\mathbf{2})(+1,-1)}] \ ,
\end{aligned} \tag{10}$$

$$\mathbf{10} = [\overset{H}{(\mathbf{3},\mathbf{1})(-2,+2)} + \overset{h}{(\mathbf{1},\mathbf{2})(+3,+2)}] + [\overset{H'}{(\overline{\mathbf{3}},\mathbf{1})(+2,-2)} + \overset{h'}{(\mathbf{1},\mathbf{2})(-3,-2)}] \ . \tag{11}$$

and the corresponding fields from $\tilde{\chi}$ are denoted with a tilde.

The spontaneous symmetry breaking by the adjoint Higgs field discussed above correspond to giving appropriate *vevs* to the two neutral singlets in **45**. The **16** acquires *vev* in the direction of the singlet $(\mathbf{1},\mathbf{1})(0,-5)$, leaving $U(1)_Y$ unbroken. Standard notation is introduced for the quark (D,U,Q), lepton (L,E), Higgs doublet (h,h') and color Higgs triplet (H,H') superfields. Let's rewrite the superpotential (2) in terms of these superfields and $\phi, \Phi, s_i, \tilde{s}_\pm, q_i$ and $\tilde{q}_\pm$ (*vevs* of $SU(3)_c \otimes SU(2)_w \otimes U(1)_Y$ singlets). Generation indices and explicit coupling constants are suppressed for brevity:

$$\begin{aligned}
W = & \left\{(\phi + \Phi)h + \tilde{s}\tilde{h} + \tilde{q}\tilde{L}\right\} Q\,U + \left\{(\phi + \Phi)h' + \tilde{s}\tilde{h}'\right\}(Q\,D + E\,L) \\
& + (\phi + \Phi)(q\,h\,L + s\,h\,h') + (\tilde{s}\,s + \tilde{q}\,q)\left(\tilde{h}\,h' + \tilde{h}'\,h + \tilde{L}\,L\right) + s\tilde{q}h'\tilde{L} + \tilde{s}q\tilde{h}L \\
& + (\phi + \Phi)(q\,HD + s\,HH') + (\tilde{s}\,s + \tilde{q}\,q)\left(\widetilde{H}H' + \widetilde{H}'H + \widetilde{D}\,D\right) + s\tilde{q}H'\widetilde{D} + \tilde{s}q\widetilde{H}D \\
& + \left\{(\phi + \Phi)H + \tilde{s}\widetilde{H} + \tilde{q}\widetilde{D}\right\}(Q\,Q + E\,U) + \left\{(\phi + \Phi)H' + \tilde{s}\widetilde{H}'\right\}(Q\,L + U\,D) + \ldots
\end{aligned} \tag{12}$$

Only the terms relevant for the chiral fermion mass matrices and Higgs doublet and triplet mass matrices are shown. In particular, the first line of (12) shows terms contributing to the top, bottom and lepton mass matrices, terms in the second line are responsible for Higgs doublet masses, terms in the next line give masses to color Higgs triplets. The last line contains terms that are responsible for unacceptably rapid proton decay if the color Higgs triplets are not sufficiently massive. As can be seen from the first line of (12), mixing of h with $\widetilde{L}$, h' with L, H with $\widetilde{D}$ and H' with D affects charge-2/3 quark mass matrix differently than lepton and charge-$-1/3$ quark mass matrices. The last two matrices are affected by mixing in an identical way. So the leptons and the charge-$-1/3$ quarks are expected to have the same mass matrix at the GUT scale. It is worth noticing that, due to the underlying gauge and discrete symmetries of the string model, the lepton-number-violating terms LLE and LQD and the baryon-number-violating term UDD are absent.

We also need to impose the F-flatness conditions with respect to the fields $\mathbf{S}_0$, $\mathbf{S}_i$, $\widetilde{\mathbf{S}}_\pm$, $\mathbf{Q}_0$, $\mathbf{Q}_i$ and $\widetilde{\mathbf{Q}}_\pm$ (note that the F-flatness conditions with respect to the Higgs fields $\mathbf{H}_0$, $\mathbf{H}_i$ and $\widetilde{\mathbf{H}}_\pm$ are automatically satisfied). These conditions can be easily read from (2):

$$\partial W/\partial\widetilde{\mathbf{Q}}_\mp = 0 \quad \Rightarrow \quad \lambda_{22}\,\tilde{q}_\pm(q\cdot q) + \lambda_{23}\,\tilde{s}_\pm(s\cdot q) = 0, \tag{13}$$
$$\partial W/\partial\widetilde{\mathbf{S}}_\mp = 0 \quad \Rightarrow \quad \lambda_{23}\,\tilde{q}_\pm(s\cdot q) + \lambda_{21}\,\tilde{s}_\pm(s\cdot s) = 0, \tag{14}$$
$$\partial W/\partial\mathbf{Q}_{\bar{i}} = 0 \quad \Rightarrow \quad 2\lambda_{22}\,q_i\tilde{q}_+\tilde{q}_- + \lambda_{23}\,s_i\,(\tilde{s}_+\tilde{q}_- + \tilde{s}_-\tilde{q}_+) = 0, \tag{15}$$
$$\partial W/\partial\mathbf{S}_{\bar{i}} = 0 \quad \Rightarrow \quad \lambda_{23}\,q_i\,(\tilde{s}_+\tilde{q}_- + \tilde{s}_-\tilde{q}_+) + 2\lambda_{21}\,s_i\tilde{s}_+\tilde{s}_- = 0 \tag{16}$$

Two constraints need to be satisfied in order for a non-trivial solution for the F-flatness conditions to exist:

$$0 = \lambda_{21}\lambda_{22}\,(s\cdot s)(q\cdot q) - \lambda_{23}^2\,(s\cdot q)^2\ , \tag{17}$$
$$0 = 4\,\lambda_{21}\lambda_{22}\,\tilde{s}_+\tilde{s}_-\tilde{q}_+\tilde{q}_- - \lambda_{23}^2\,(\tilde{s}_+\tilde{q}_- + \tilde{s}_+\tilde{q}_-)^2 \tag{18}$$

The effective low energy couplings of the type $\lambda_{22}\,\mathbf{Q}\,\mathbf{Q}\,\widetilde{\mathbf{Q}}\widetilde{\mathbf{Q}}$ and $\lambda_{23}\mathbf{Q}\,\mathbf{S}\,\widetilde{\mathbf{Q}}\widetilde{\mathbf{S}}$ which are responsible for generation–antigeneration pairing can be summarized in the following 2×5 matrix:

$$\begin{array}{c|c|c|} & \widetilde{\mathbf{Q}}_+ & \widetilde{\mathbf{Q}}_- \\ \hline \mathbf{Q}_0 & 0 & 0 \\ \mathbf{Q}_{++} & \lambda_{22}\,q_{--}\tilde{q}_- + \lambda_{23}\,s_{--}\tilde{s}_- & \lambda_{22}\,q_{--}\tilde{q}_+ + \lambda_{23}\,s_{--}\tilde{s}_+ \\ \mathbf{Q}_{--} & \lambda_{22}\,q_{++}\tilde{q}_- + \lambda_{23}\,s_{++}\tilde{s}_- & \lambda_{22}\,q_{++}\tilde{q}_+ + \lambda_{23}\,s_{++}\tilde{s}_+ \\ \mathbf{Q}_{+-} & \lambda_{22}\,q_{-+}\tilde{q}_- + y\lambda_{23}\,s_{-+}\tilde{s}_- & \lambda_{22}\,q_{-+}\tilde{q}_+ + \lambda_{23}\,s_{-+}\tilde{s}_+ \\ \mathbf{Q}_{-+} & \lambda_{22}\,q_{+-}\tilde{q}_- + \lambda_{23}\,s_{+-}\tilde{s}_- & \lambda_{22}\,q_{+-}\tilde{q}_+ + \lambda_{23}\,s_{+-}\tilde{s}_+ \\ \hline \end{array} \tag{19}$$

When the fields $\mathbf{Q}_i, \mathbf{S}_i, \widetilde{\mathbf{Q}}_\pm$ and $\widetilde{\mathbf{S}}_\pm$ acquire *vevs* (denoted in the table above as $q_i, s_i, \tilde{q}_\pm$ and $\tilde{s}_\pm$, respectively), a generation–antigeneration pairing takes place. Generically, only one scenario of pairing is allowed by F-flatness constraints:

• $\tilde{q}_+ = \tilde{s}_+ = 0$, $\tilde{q}_- \neq 0$ and $\tilde{s}_- \neq 0$ (interchange of indices $+ \leftrightarrow -$ clearly provides an equivalent choice). This choice of *vevs* satisfies (15,18) and one is left to choose q_i and s_i so that (13,17) are satisfied. Only one generation–antigeneration pairing occurs: $\widetilde{\mathbf{Q}}_+$ pairs with a linear combination of $\mathbf{Q}_q$ and $\mathbf{Q}_s$ where

$$\mathbf{Q}_q = \frac{(q\cdot\mathbf{Q})}{\mathbf{q}} \equiv (q_{++}\mathbf{Q}_{--} + q_{--}\mathbf{Q}_{++} + q_{+-}\mathbf{Q}_{-+} + q_{-+}\mathbf{Q}_{+-})/\mathbf{q}\ ; \tag{20}$$
$$\mathbf{Q}_s = \frac{(s\cdot\mathbf{Q})}{\mathbf{s}} \equiv (s_{++}\mathbf{Q}_{--} + s_{--}\mathbf{Q}_{++} + s_{+-}\mathbf{Q}_{-+} + s_{-+}\mathbf{Q}_{+-})/\mathbf{s}\ ;$$
$$\mathbf{s} \equiv \sqrt{\sum_i s_i^2}\ , \qquad \mathbf{q} \equiv \sqrt{\sum_i q_i^2}$$

Mass term has the form $(\lambda_{22}\,\mathbf{q}\,\tilde{q}_-\,\mathbf{Q}_q + \lambda_{23}\,\mathbf{s}\,\tilde{s}_-\,\mathbf{Q}_s)\,\widetilde{\mathbf{Q}}_+$.

To achieve second generation–antigeneration pairing one has to consider non-vanishing higher–point couplings that are present in the superpotential of this model, such as 8–point couplings $(\chi_{++}\chi_{--}+\chi_{+-}\chi_{-+})^2\,\tilde{\chi}_+^2\,\tilde{\chi}_-^2$ and $\chi_0(\chi_{++}\chi_{--}+\chi_{+-}\chi_{-+})^2\,\tilde{\chi}_+\,\tilde{\chi}_-\,\widehat{\Phi}$. These couplings will modify the F-flatness conditions and allow small non-vanishing $\tilde{s}_+$ and $\tilde{q}_+$. As a result, the second generation–antigeneration pairing can take place at a lower scale. (Here we note that F-flatness conditions will be modified once the supersymmetry is broken, so that the second generation–antigeneration pairing scale is very likely to be at least as high as m_S.)

Since the presence of light color Higgs triplets will mediate unacceptably rapid proton decay, one needs to arrange for them to be appropriately heavy. One needs also to arrange for a pair of Higgs doublets to remain light. To simplify the analysis, we assume that Higgs doublets and triplets come solely from **10** of $SO(10)$ (due to the mixing, they may also come from $\mathbf{16}+\overline{\mathbf{16}}$ s).

We give the adjoint Φ of $SO(10)$ a *vev* in the following direction: $\Phi = \epsilon\otimes\mathrm{diag}(a,a,a,b,b)$, where $a,b \neq 0,\ a\neq b$. Since $\tilde{s}_+ = \tilde{q}_+ = 0,\ \tilde{h}_-,\ \tilde{h}'_-,\ \widetilde{H}_-$ and $\widetilde{H}'_-$ decouple (they may become massive via the coupling $\tilde{\chi}_-^3\,\tilde{U}_+ D_+ D_-\,\hat{\Phi}^4$ at a much lower scale) and one is left with 6×6 mass matrix $\mathcal{M}_{ij}$ of Higgs doublet couplings $(\lambda'_{11}\phi + \lambda^a_{11}\,b)s\,h\,h'$ and $(\lambda'_{23}\tilde{s}\,s\ + +\tilde{h}'\,h)$ (mixing is disregarded, the coupling constants are not shown explicitly in the table):

$$
\begin{array}{cccccccc}
h_0 & h_{++} & h_{--} & h_{+-} & h_{-+} & \tilde{h}_+ & & \\
\left.\begin{array}{|cccccc|}
\hline
0 & (\phi+b)\,s_{--} & (\phi-b)\,s_{++} & (\phi+b)\,s_{-+} & (\phi-b)\,s_{+-} & 0 \\
(\phi-b)\,s_{--} & 0 & (\phi+b)\,s_0 & 0 & 0 & p_{--} \\
(\phi+b)\,s_{++} & (\phi-b)\,s_0 & 0 & 0 & 0 & p_{++} \\
(\phi-b)\,s_{-+} & 0 & 0 & 0 & (\phi+b)\,s_0 & p_{-+} \\
(\phi+b)\,s_{+-} & 0 & 0 & (\phi-b)\,s_0 & 0 & p_{+-} \\
0 & p_{--} & p_{++} & p_{-+} & p_{+-} & 0 \\
\hline
\end{array}\right. & & & & & & \begin{array}{c} h'_0 \\ h'_{++} \\ h'_{--} \\ h'_{+-} \\ h'_{-+} \\ \tilde{h}'_+ \end{array} &
\end{array}
\tag{21}
$$

where we denote $p_i = \lambda'_{23}\tilde{s}_-\,s_i + \lambda'_{22}\tilde{q}_-\,q_i$. The coupling constants λ'_i are corresponding λ_i from the (2) corrected by Clebsch-Gordan factors upon gauge symmetry breaking (11). Generically, $\mathrm{rank}(\mathcal{M}_{ij}) = 6$. Thus, to have $\mathrm{rank}(\mathcal{M}_{ij}) = 5$, *i.e.* one pair of light Higgs doublets, it is necessary to fine-tune the parameters entering the Higgs doublet matrix, *e.g.* adjoint *vev* b.

Let us begin with the case $b = 0$ (as in Dimopoulos–Wilczek mechanism). With some algebra, it can be shown that $\det(\mathcal{M}_{ij})$ vanishes only when $(s\cdot s) = (s\cdot q) = (q\cdot q) = 0$ (this is the only choice that satisfies the F-flatness conditions (13,14)). However, with these conditions imposed on the *vevs* s_i and q_i, the rank of $\mathcal{M}_{ij}$ invariably becomes four, *i.e.*, there are *two* pairs of light Higgs doublets. The extra pair of light Higgs doublets would affect the successful SUSY GUT prediction for $\sin^2\theta_W$. One thus has to consider fine–tuning b to some non–zero value in order to arrange for only one pair of doublets to exist. Indeed, one finds that if b is fine-tuned to satisfy

$$
0 = (p\cdot s)^2\,b^4\ +\ \phi^4\left((p\cdot s)^2\ -\ (p\cdot p)(s\cdot s)\right) \tag{22}
$$

$$- b^2 \phi^2 \left(2\,(p\cdot s)^2 - (p\cdot p)(s\cdot s) + 8(p_{++}s_{+-} - p_{+-}s_{++})(p_{--}s_{-+} - p_{-+}s_{--})\right)$$

then $\mathrm{rank}(\mathcal{M}_{ij}) = 5$.

Thus, there are three constraints imposed on the 13 parameters that enter the Higgs doublet mass matrix: s_i, q_i, $\widetilde{s}_-$, $\widetilde{q}_-$, s_0, ϕ and b. These are the F-flatness conditions (13,14) and condition (22) for a pair of Higgs doublets to be light (constraints (15,16) were already satisfied by setting $\widetilde{s}_+ = \widetilde{q}_+ = 0$). The D-flatness constraint (8) can be satisfied for arbitrary $\widetilde{q}_-$, q_i by tuning q_0 which otherwise does not enter our analysis. The resulting moduli space generically has only one pair of light Higgs doublets. An extra pair(s) of light Higgs doublets occur in the subspaces specified by $b = 0$ or $s_0 = 0$.

The general expression for the light Higgs doublets is rather cumbersome and we choose to substitute constraint (22) with two simpler constraints from which (22) follows directly:

$$0 = p_{--}s_{-+} - p_{-+}s_{--}\;; \qquad \phi - b = 0 \tag{23}$$

In this subspace of the moduli space, the light Higgs doublet $\widehat{h}$ is a particular linear combination of the full set of fields h_0, h_i and $\widetilde{h}_-$. Its dependence on h_0 will be of importance later on.

The mass matrix of the Higgs triplet couplings $\mathcal{N}_{ij}$ has contributions from $(\lambda''_{11}\phi + \lambda^A_{11}\,a)s\,H\,H'$ and $(\lambda''_{23}\widetilde{s}\,s + \lambda''_{22}\widetilde{q}\,q)\left(\widetilde{H}\,H' + \widetilde{H}'\,H\right)$. It has a similar form as $\mathcal{M}_{ij}$ except b is changed to a and the coupling constants are different. Generically, it has rank 6 and thus all Higgs triplets acquire a mass which is slightly below M_{GUT} (except $\widetilde{H}_-$ and $\widetilde{H}'_-$ as discussed earlier).

So far we have specified only the first generation–antigeneration pairing (20). Let us denote that generation as $\mathbf{Q}_X$. The modified F-flatness conditions will allow for small non-vanishing $\widetilde{q}_+$ and $\widetilde{s}_+$ and the second generation (which we denote as $\mathbf{Q}_Y$) will pair as can be seen from (19). Let us for brevity write $\mathbf{Q}_X$ and $\mathbf{Q}_Y$ as

$$\begin{aligned}
\mathbf{Q}_X &= \frac{(l\cdot\mathbf{Q})}{\mathbf{l}} \equiv (l_{++}\mathbf{Q}_{--} + l_{--}\mathbf{Q}_{++} + l_{+-}\mathbf{Q}_{-+} + l_{-+}\mathbf{Q}_{+-})/\mathbf{l}\;; \qquad && \mathbf{l} \equiv \sqrt{\sum_i l_i^2} \\
\mathbf{Q}_Y &= \frac{(k\cdot\mathbf{Q})}{\mathbf{k}} \equiv (k_{++}\mathbf{Q}_{--} + k_{--}\mathbf{Q}_{++} + k_{+-}\mathbf{Q}_{-+} + k_{-+}\mathbf{Q}_{+-})/\mathbf{k}\;; \qquad && \mathbf{k} \equiv \sqrt{\sum_i k_i^2}
\end{aligned} \tag{24}$$

where l_i can be read from (20) and k_i have an analogous form with $\widetilde{q}_-$, $\widetilde{s}_- \to \widetilde{q}_+$, $\widetilde{s}_+$. It is important to note that neither $\mathbf{Q}_X$ nor $\mathbf{Q}_Y$ have a $\mathbf{Q}_0$ component. Let us now analyze the mass matrix for the remaining three generations. We have to choose a new basis $\mathbf{Q}_0$, $\mathbf{Q}_1$, $\mathbf{Q}_2$ such that there are no mass terms of the form $\mathbf{Q}_{1,2}\mathbf{Q}_X$ or $\mathbf{Q}_{1,2}\mathbf{Q}_Y$. Without loss of generality we choose (the overall normalization factors are not shown)

$$\begin{aligned}
\mathbf{Q}_1 &\sim (l_{+-}k_{++} - k_{+-}l_{++})\mathbf{Q}_{--} + (l_{--}k_{+-} - k_{--}l_{+-})\mathbf{Q}_{++} + (l_{++}k_{--} - k_{++}l_{--})\mathbf{Q}_{+-} \\
\mathbf{Q}_2 &\sim (l_{-+}k_{++} - k_{-+}l_{++})\mathbf{Q}_{--} + (l_{--}k_{-+} - k_{--}l_{-+})\mathbf{Q}_{++} + (l_{++}k_{--} - k_{++}l_{--})\mathbf{Q}_{-+}
\end{aligned} \tag{25}$$

The mass matrix for $\mathbf{Q}_0$, $\mathbf{Q}_1$ and $\mathbf{Q}_2$ is

$$
\begin{array}{c|ccc|}
 & \mathbf{Q}_0 & \mathbf{Q}_1 & \mathbf{Q}_2 \\
\hline
\mathbf{Q}_0 & 0 & a_{01} & a_{02} \\
\mathbf{Q}_1 & a_{01} & a_{11} & a_{12} \\
\mathbf{Q}_2 & a_{02} & a_{12} & a_{22} \\
\hline
\end{array}
\qquad (26)
$$

The coefficients a_{11}, a_{12} and a_{22} are proportional to h_0 and since the light Higgs generically depends on h_0 they do not vanish. The coefficients a_{01} and a_{02} depend on h_i and are also generically non-zero. Thus this matrix has rank three unless some fine-tuning is imposed on the s_i, q_i, $\tilde{s}_\pm$ and $\tilde{q}_\pm$. It remains to be shown whether one can arrange for the rank of the quark mass matrix to be one as suggested by phenomenology. Should the answer be positive it will require fine-tuning of the *vevs* in some rather specific manner.

Let us summarize the phenomenological properties of the E_6 and $SO(10)$ models analyzed in this preliminary analysis:

• (i) The mass of one generation–antigeneration pair is $\approx 10^{16}$ GeV while that of the other pair is (perhaps, much) lower. In our analysis, we assume that the generation and antigeneration come entirely from **16** and $\overline{\mathbf{16}}$ of $SO(10)$. *A priori*, after spontaneous symmetry breaking, the **10** of $SU(5)$ in **16** will mix with the **10** of $SU(5)$ in the adjoint **45** of $SO(10)$ (and similarly for the $\overline{\mathbf{16}}$). Clearly, a more careful analysis taking into account this mixing will be useful.

• (ii) In a generic solution to the flatness conditions, the Higgs triplet mass matrix has rank 6. One can arrange for a pair of Higgs doublets to remain light by fine-tuning the *vev* of the adjoint in its coupling to the doublets. We have chosen a scenario in which the Higgs fields come solely from the **10**s of $SO(10)$, that is, mixing with **16**s has been ignored. It is important to investigate possible mixing more carefully.

• (iii) Notice that only one generation-antigeneration pairing takes place at M_{GUT} and fine-tuning is necessary for the Higgs doublets to remain light. There is no light Higgs doublet arising naturally in this model, nor could one obtain a quark mass matrix with rank one in a natural way. These naturalness problems may result from the approximations we have made, for instance, the Kähler potential, the higher order terms in the superpotential and non-perturbative corrections were ignored. Perhaps, stringy symmetries such as target space modular invariance would improve our understanding of the superpotential and solve the fine-tuning problems in a natural way.

• (iv) The hidden sector gauge group $SU(2)$ is asymptotically-free. Assuming that stabilization of the dilaton expectation value at the correct value results from the dynamical SUSY breaking, and evaluating the scale of the gaugino condensation to be $\sim 10^{13} GeV$, we find that the $SUSY$ breaking scale in the visible sector is of order $10 \sim 100\ TeV$.

• (v) The $SU(3)_c$ and the $SU(2)_w$ adjoint Higgs fields are presumably quite light; they may be as light as the electroweak scale.

Of course, all the above properties depend crucially on the non-perturbative string dynamics that we have assumed. Clearly, better understanding of SUSY breaking in string theory is needed. On the other hand, we see that the spectra and their couplings are tightly restricted in the string models. This already allows non-trivial tests of the viability of such models. It is clearly important to calulate explicitly the perturbative couplings in the superpotential and the Kähler potential. This will provide further stringent tests on the viability of the 3-family E_6 and its related $SO(10)$ string models.

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REFERENCES

[1] For reviews, see *e.g.*, A.E. Faraggi, hep-ph/9405357;
F. Quevedo, hep-th/9511131; hep-th/9603074;
K. Dienes, hep-th/9602045;
L.E. Ibáñez, hep-th/9505098;
C. Kounnas, hep-th/9512034;
J. Lopez, Rept. Prog. Phys. **59** (1996) 819, hep-ph/9601208;
J.D. Lykken, hep-th/9607144.

[2] See *e.g.*, J. Pati, hep-ph/9611371; hep-ph/9607446;
J. Hisano, H. Murayama and T. Yanagida, Phys. Rev. **D49** (1994) 4966, Nucl. Phys. **B402** (1993) 46;
V. Lucas and S. Raby, hep-ph/9610293;
K. Dienes, hep-th/9606467;
G.Anderson, S. Raby, S. Dimopoulos, L. J. Hall and G. Starkman, Phys. Rev. **D49** (1994) 3660;
G. Aldazabal, A. Font, L. E. Ibanez and A. M. Uranga, Nucl. Phys. **B452** (1995) 3.

[3] See *e.g.*, T. Taylor, G. Veneziano and S. Yankielowicz, Nucl. Phys. **B218** (1983) 493;
J.-P. Derendinger, L.E. Ibáñez and H.P. Nilles, Phys. Lett. **155B** (1985) 65;
M. Dine, R. Rohm, N. Seiberg and E. Witten, Phys. Lett. **156B** (1985) 55;
T. Taylor, Phys. Lett. **164B** (1985) 43;
A. Font, L.E. Ibáñez, D. Lüst and F. Quevedo, Phys. Lett. **245B** (1990) 401;
S. Ferrara, N. Magnoli, T. Taylor and G. Veneziano, Phys. Lett. **245B** (1990) 409;
H.P. Nilles and M. Olechowski, Phys. Lett. **248B** (1990) 268;
P. Binétruy and M.K. Gaillard, Phys. Lett. **253B** (1991) 119;
L. Dixon, V. Kaplunovsky and J. Louis, Nucl. Phys. **B329** (1990) 27; **B355** (1991) 649.

[4] Z. Kakushadze and S.-H.H. Tye, Phys. Rev. Lett. **77** (1996) 2612, hep-th/9605221; Phys. Rev. **D54** (1996) 7520, hep-th/9607138; Phys. Lett. **392B** (1996) 335, hep-th/9609027; hep-th/9610106; hep-th/9701057.

[5] Z. Kakushadze, G. Shiu and S.-H.H. Tye, hep-th/9704113.

[6] H. Georgi, in Particles and Fields (ed. by C. Carlson), A.I.P. (1975);
H. Fritzsch and P. Minkowski, Ann. Phys. **93**, (1975) 193;
F. Gursey, P. Ramond and P. Sikivie, Phys. Lett. **60B** (1976) 177.

[7] S. Weinberg, Phys. Rev. **D26** (1982) 287;
N. Sakai and T. Yanagida, Nucl. Phys. **B197** (1982) 533;
S. Dimopoulos, S. Raby and F. Wilczek, Phys. Lett. **112B** (1982) 113.

[8] S. Dimopoulos and F. Wilczek, in Proc. Erice Summer School (ed. by A. Zichichi), I.T.P. preprint NSF-ITP-82-07;
K. S. Babu and S. M. Barr, Phys. Rev. **D48** (1993) 5354; ibid **D51** (1995) 2463;
J. Hisano, Phys. Rev. **D49** (1994) 4966.

[9] T. Banks and M. Dine, Phys. Rev. **D50** (1994) 7454;
P. Binétruy, M. K. Gaillard and Y.-Y. Wu, Nucl. Phys. **B481** (1996) 109, hep-th/9605170.

[10] P. Nath, R. Arnowitt and A.H. Chamseddine, Phys. Rev Lett. **49** (1982) 970; Phys. Lett **121B** (1983) 33.