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SUPERSYMMETRIC FLIPPED SU(5) REVITALIZED

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A B S T R A C T

We describe a simple $N = 1$ supersymmetric GUT based on the group $SU(5) \times U(1)$ which has the following virtues: the gauge group is broken down to the $SU(3)_C \times SU(2)_L \times U(1)_Y$ of the Standard Model using just $\underline{10}, \underline{10}$ Higgs representations, and the doublet-triplet mass splitting problem is solved naturally by a very simple missing-partner mechanism. The successful supersymmetric GUT prediction for $\sin^2 \theta_w$ can be maintained, whilst there are no fermion mass relations. The gauge group and representation structure of the model may be obtainable from the superstring.

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Although supersymmetric GUTs have many attractive features, most of them have serious defects. Generally¹⁾, one has had to fine-tune parameters in order to keep Higgs doublets light while giving large masses to their colour-triplet partners, so as to inhibit proton decay. Although it is possible²⁾ to avoid this doublet-triplet mass splitting problem in a natural way, in practice this has required the introduction of large Higgs representations. Moreover, it was not always possible to find discrete symmetries that forbade all the unwanted superpotential couplings³⁾. Recently, creative activity in supersymmetric GUT model-building has subsided as authors have tried to fit into the phenomenological strait-jacket suggested by the superstring. Manifold compactifications of the ten-dimensional heterotic superstring suggested subgroups of E_6 as gauge groups⁴⁾, with matter particles assigned to subsets of $\underline{27}$ and $\overline{27}$ representations of E_6 . Many of the possible subgroups included Abelian $U(1)$ factors as well as non-Abelian factors, with equal couplings at a grand unification scale fixed by the Hosotani⁵⁾ flux-breaking mechanism. More recently, this phenomenological strait-jacket has been somewhat loosened as orbifold compactifications of the ten-dimensional heterotic superstring have been explored⁶⁾, as well as formulations of string theories directly in four dimensions⁷⁾. Such models could have gauge groups of rank ≤ 22 , with a correspondingly wide choice of matter representations. Thus the time is ripe for a resurgence of model-building, with the rules somewhat different from those applied previously. Models with additional Abelian $U(1)$ factors whose couplings are equal to the non-Abelian coupling(s) at the grand unification scale can now be tolerated, and "small" matter representations are favoured: totally antisymmetric representations of $SU(N)$ groups, and vector and spinor representations of $SO(N)$ groups. In this paper we report on one of the minimal supersymmetric GUTs that can be constructed according to these rules.

It is based on the group $SU(5) \times U(1)$ ^{8),9)} with just $\underline{10}$, $\overline{10}$, $\underline{5}$, $\overline{5}$ and $\underline{1}$ representations of $SU(5)$. The group is broken down to the Standard Model $SU(3)_C \times SU(2)_L \times U(1)_Y$ by fields in the $\underline{10} + \overline{10}$ representations, and then to $SU(3)_C \times U(1)_{em}$ by doublets from the $\underline{5} + \overline{5}$ representations as usual. The triplet partners of the Higgs doublets match with states from the $\underline{10} + \overline{10}$ to acquire naturally large masses at the first stage of gauge symmetry breaking. The successful supersymmetric $SU(5)$ prediction¹⁰⁾ for $\sin^2 \theta_w$ is retained, while the fermion mass relations¹¹⁾ are jettisoned. There is a see-saw mechanism¹²⁾ which guarantees small neutrino masses. The model has a natural embedding in $SO(10)$ ^{*)},

*) If this is not done, the model has no monopoles to cause cosmological problems⁹⁾.

and may well be obtainable from manifold or orbifold compactifications of the ten-dimensional heterotic superstring, or from four-dimensional superstring theories, as we comment at the end of the paper.

First we construct the model, starting with an enumeration of its chiral superfields. There are N_g generations of matter fields with the following $SU(5) \times U(1)$ transformation properties:

$$F_i = (10, 1), \quad \bar{F}_i = (\bar{5}, -3), \quad l_i^c = (1, 5) \quad (1)$$

The particle assignments in the F_i and $\bar{F}_i$ are^{8),9)}

$$F = \begin{pmatrix} 0 & d^c & -d^c & d & u \\ -d^c & 0 & d^c & d & u \\ d^c & -d^c & 0 & d & u \\ -d & -d & -d & 0 & \nu^c \\ -u & -u & -u & -\nu^c & 0 \end{pmatrix}; \quad \bar{F} = \begin{pmatrix} u^c \\ u^c \\ u^c \\ \nu \\ e \end{pmatrix} \quad (2)$$

and similarly for the other generations. There are two conjugate pairs of Higgs representations:

$$H = (10, 1), \quad \bar{H} = (\bar{10}, -1), \quad h = (5, -2), \quad \bar{h} = (\bar{5}, 2) \quad (3)$$

We label the states in the H multiplet by the same symbols as in the F multiplet, with an additional subscript: d_H^c, d_H etc... States in the $\bar{H}$ multiplet are labelled by bars: $\bar{d}_H^c, \bar{d}_H$, etc... States in the h are denoted by (D, D, D, h^-, h^+) and in the $\bar{h}$ by $(\bar{D}, \bar{D}, \bar{D}, h^+, h^-)^T$. In addition to the non-singlet representations $(1, 3)$, the model employs $(N_g + 1)$ singlets $\phi_m = (1, 0)$. Note that the $U(1)$ assignments of the $F, \bar{F}, l^c, H$ (and $\bar{H}$) representations are those appearing naturally when spinorial $\underline{16}$ (and $\overline{16}$) representations of $SO(10)$ are decomposed with respect to $SU(5) \times U(1)$, whilst the h and $\bar{h}$ are naturally embedded in a vector $\underline{10}$ of $SO(10)$.

The superpotential of our model is

$$W = \left. \begin{aligned} & \lambda_1^i F_i F_j h + \lambda_2^i F_i \bar{F}_j \bar{h} + \lambda_3^i \bar{F}_i l_j^c h \\ & + \lambda_4 H H h + \lambda_5 \bar{H} \bar{H} \bar{h} \\ & + \lambda_6^{im} F_i \bar{H} \phi_m + \lambda_7^m h \bar{h} \phi_m + \lambda_8^{mnp} \phi_m \phi_n \phi_p \end{aligned} \right\} \quad (4)$$

which is the most general set of trilinear interactions invariant under the Z_2 symmetry

$$H \rightarrow -H \quad (5)$$

Note also that the couplings are just those which would be allowed by $SO(10)$.

Next we discuss gauge symmetry breaking. There is just one F-flat and D-flat direction: $\langle 0|H|0\rangle = \langle 0|\bar{H}|0\rangle \neq 0$, which is suitable for generating a large vacuum expectation value (v.e.v.) of order the grand unification scale M_{GUT}^* . These v.e.v.'s can always be rotated into the v_H^c and $\bar{v}_H^c$ directions, showing that the Standard Model group is unbroken:

$$SU(5) \times U(1) \longrightarrow SU(3)_C \times SU(2)_L \times U(1)_Y \quad (6)$$

$$\langle 0|H, \bar{H}|0\rangle = M_{GUT}$$

Note that unlike conventional GUTs where the full symmetry is broken by an adjoint Higgs v.e.v., and there is a considerable ambiguity in the choice of low-energy group, here the desired pattern (6) is unique. When this occurs, all spin-zero components of the H and $\bar{H}$ representations are eaten or acquire large masses via the supersymmetric Higgs mechanism, except for the d_H^c and $\bar{d}_H^c$, and one neutral chiral superfield. The second stage of gauge symmetry breaking: $SU(3)_C \times SU(2)_L \times U(1)_Y$ occurs via small v.e.v.'s for the doublet components of the h and $\bar{h}$ in the usual way. Later we will discuss a scenario for generating these v.e.v.'s starting from a seed of soft supersymmetry breaking.

Now we discuss the normal doublet-triplet mass splitting mechanism of this model. The λ_4 and λ_5 couplings in the superpotential (4) combine the uneaten d_H^c and $\bar{d}_H^c$ with the D and $\bar{D}$ respectively to form massive Dirac eigenstates with masses

$$m_{d_H^c D} = \lambda_4 \langle 0|v_H^c|0\rangle, \quad m_{\bar{d}_H^c \bar{D}} = \lambda_5 \langle 0|\bar{v}_H^c|0\rangle \quad (7)$$

The economy and elegance of this mechanism will be appreciated by all those who sweated previously over $\underline{75} + \underline{50} + \underline{50}$ representations in $SU(5)$ models^{2),3)}. More simply, we are applying the missing partner mechanism using the lowest-dimensional suitable representation of $SU(5)$, namely the $\underline{10}$.

*) The λ_8 terms in the superpotential (4) prevent large v.e.v.'s for the singlet fields ϕ_m .

The other particle masses in this model appear as follows. The $\lambda_1 \bar{F} F h$ couplings give masses to the charge $-1/3$ quarks d, s, b , etc..., which are unrelated to the charged lepton masses. Thus we discard the previous, successful mass relation for m_b/m_τ ¹¹⁾, but also the problematic mass relations for $m_{d,s,e,\mu}$. The $\lambda_2 \bar{F} \bar{F} h$ couplings give masses to the charge $2/3$ quarks u, c, t , etc., and to the neutrinos:

$$m_u = m_{\nu_e}, m_c = m_{\nu_\mu}, m_t = m_{\nu_\tau}, \text{ etc.} \quad (8)$$

This is not disastrous, because the $\lambda_6 \bar{F} H \phi$ couplings provide a see-saw mechanism¹²⁾, since they give large Dirac masses

$$\lambda_6^{im} \langle 0 | \bar{\nu}_H^c | 0 \rangle \nu_i^c \phi_m \quad (9)$$

coupling the right-handed neutrinos ν_i^c to the singlets ϕ_m . The general form of the neutrino mass matrix is

$$(\nu_i, \nu_i^c, \phi_m) \begin{pmatrix} 0 & m_u^{ij} & 0 \\ m_u^{ji} & 0 & \lambda_6^{im} \langle 0 | \bar{\nu}_H^c | 0 \rangle \\ 0 & \lambda_6^{mi} \langle 0 | \bar{\nu}_H^c | 0 \rangle & 0 \end{pmatrix} \begin{pmatrix} \nu_i \\ \nu_j^c \\ \phi_m \end{pmatrix} \quad (10)$$

which contains N_g massless eigenstates^{*} $\nu_i + O(m_u/M_{\text{GUT}}) (\phi_i)$ and Dirac eigenstates $(\nu_i^c \pm \phi_i)/\sqrt{2} + O(m_u/M_{\text{GUT}}) \nu_i$ with masses $O(M_{\text{GUT}})$. The $\lambda_3 \bar{f} l^c h$ give masses to the charged leptons, which are unrelated to any quark masses. Finally, we note that the couplings $\lambda_7 h \bar{h} \phi_m$ provide the $h \bar{h}$ mixing which is essential to the establishment of an electroweak symmetry-breaking vacuum with both $\langle 0 | h^0 | 0 \rangle$, $\langle 0 | \bar{h}^0 | 0 \rangle \neq 0$, and to avoid an unacceptable electroweak axion. Thus we see that all of the Yukawa couplings in the superpotential (4) have important phenomenological roles to play.

The low-energy spectrum of this model contains the same non-singlet fields as the minimal supersymmetric Standard Model. Therefore the conventional renormalization group estimates¹⁰⁾ of M_{GUT} ($\approx 10^{16}$ GeV) and $\sin^2 \theta_w$ (≈ 0.230 to 0.240) are also applicable to this model if the $U(1)$ coupling is equal to the $SU(5)$ coupling at M_{GUT} . This would be the case if the model were embedded in $SO(10)$, or if the Hosotani⁵⁾ mechanism were used to break a larger group down to $SU(5) \times U(1)$ directly since the gauge kinetic function $f_{ab} \propto \delta_{ab}$ in superstring

^{*}) Thus there is clearly no problem with cosmological nucleosynthesis, and no neutrino dark matter. One combination of the ϕ_m acquires a negative supersymmetry breaking mass squared and a mass $O(m_W)$ from the λ_7, λ_8 terms in the superpotential (4), and a v.e.v. $\langle 0 | \phi_m | 0 \rangle = O(m_W)$.

models. We note that over the course of time the experimentally preferred value of $\sin^2\theta_w$ has drifted up to ≈ 0.230 ¹³⁾, and now favours the supersymmetric GUT prediction¹⁰⁾ over the original non-supersymmetric GUT prediction¹⁴⁾ ≈ 0.210 to 0.220 .

Next we discuss the generation of the two required gauge symmetry-breaking scales M_{GUT} and m_W in the context of renormalization group corrections to bare soft supersymmetry-breaking terms¹⁵⁾. We take the point of view that the seed of soft supersymmetry-breaking in the observable sector is likely to be a universal gaugino mass $m_{\frac{1}{2}}$, that the squared scalar masses $\tilde{m}_F^2$ etc. are likely to be $\ll m_{\frac{1}{2}}$ initially, and that the coefficients A_k ($k = 1, 2, \dots, 8$) of trilinear soft supersymmetry-breaking terms are also negligible initially^{16), 17)} (this point of view is not, however, crucial). The following sequence of events occurs as one decreases the renormalization scale μ from m_P to m_W .

(1) Since the singlets ϕ_m have no gauge interactions, the renormalization group equations first drive $\tilde{m}_{\phi_m}^2 < 0$. However, the ϕ_m are prevented from acquiring large v.e.v.'s by the λ_8 couplings, which mean there are no F-flat directions in the singlet sector.

(2) The couplings $\lambda_{4,5,6}$ drive $\tilde{m}_{H,\bar{H}}^2 < 0$ at a large scale μ_1 , and then v.e.v.'s $\langle 0 | H, \bar{H} | 0 \rangle = M_{\text{GUT}} \sim \mu_1$ form, since there is an F- and D-flat direction in the potential. The value of μ_1 and hence M_{GUT} may be determined either by the relative strengths of the gauge and Yukawa couplings, or by the scale at which soft supersymmetry-breaking first appears. In a scenario of spontaneous dynamical supersymmetry-breaking by a hidden-sector gaugino condensate¹⁸⁾, this scale can be identified with the scale Λ_c at which the hidden-sector gauge interactions become strong. Thus we might expect $M_{\text{GUT}} \sim \mu_1 \sim \Lambda_c$ may be $\sim (10^{15}$ to $10^{16})\text{GeV}$, as previously suggested in order to obtain an appropriate value of $m_{\frac{1}{2}} \sim \Lambda_c^5/m_P^4 \sim 1\text{ TeV}$ ¹⁷⁾, and to obtain interesting cosmic strings from hidden-sector flux tubes¹⁹⁾. Now we may understand why the value of M_{GUT} where $SU(5) \times U(1) \rightarrow SU(3)_C \times SU(2)_L \times U(1)_Y$ which was inferred from the low-energy renormalization group analysis is so close to the value of the condensate scale advocated for totally independent reasons.

(3) The couplings $\lambda_{1,2,3,4,5,7}$ drive $\tilde{m}_{h,\bar{h}}^2 < 0$, probably also at a large scale close to μ_1 . This is because the gauge interactions of the $\underline{5}$ (h) and $\overline{\underline{5}}$ ($\bar{h}$) are weaker than those of the $\underline{10}$ (H) and $\overline{\underline{10}}$ ($\bar{H}$), and because the combinatorial factors in the $\lambda_{4,5}$ contributions to the renormalization group equations are larger for the h and $\bar{h}$. However, the Yukawa coupling λ_7 prevents h and $\bar{h}$ from acquiring large v.e.v.'s.

(4) There is no more excitement until μ decreases to the same order of magnitude as the soft supersymmetry-breaking coefficients $|\tilde{m}_1^2|$. At this scale v.e.v.'s of order $|\tilde{m}| \sim m_W$ can form in directions which are not F- or D-flat. Thus we expect a v.e.v. $\langle 0|\phi_m|0\rangle \sim 1$ TeV which gives the required $h\bar{h}$ mixing, and also $\langle 0|h,\bar{h}|0\rangle \sim m_W$ so that $SU(3)_C \times SU(2)_L \times U(1)_Y \rightarrow SU(3)_C \times U(1)_{em}$.

Having described our flipped $SU(5)$ model and its pattern of symmetry-breaking, we next make some comments on some of its cosmological aspects. Because $SU(5) \times U(1)$ is not a simple non-Abelian group, there are no monopoles to worry about⁹⁾, unless one embeds the model in $SO(10)$ in four dimensions. Because the magnitudes of the soft supersymmetry-breaking coefficients are $O(m_W)$, there is no gauge symmetry-breaking before the $SU(5)$ coupling becomes strong or until the temperature T falls below $\sim m_W$. In the latter case there is the danger that too much entropy will be generated during the $SU(5) \times U(1) \rightarrow SU(3)_C \times SU(2)_L \times U(1)_Y$ phase transition²⁰⁾. This problem is less severe²¹⁾, but still worrisome if baryogenesis occurs in a cold Universe at low temperatures. It may be avoided by $SU(5)$ strong-coupling effects. Vacua with $\langle 0|\phi_m|0\rangle \neq 0$ have a Z_3 symmetry which could in principle lead to a problem with domain walls. However, this Z_3 symmetry could be broken²²⁾ by higher-order terms in the superpotential such as ϕ_m^4 .

Finally, we discuss the possibility of obtaining our flipped $SU(5)$ model from an underlying superstring theory. One idea would be to obtain it from a manifold compactification, perhaps with a gauge bundle extended to $SU(4)$ ²³⁾. Generic models of this type have problems with the non-perturbative instability of the compactification manifold^{24),22)}, but models may exist which satisfy a finite set of consistency conditions²⁵⁾. If such a model exists, Hosotani⁵⁾ flux-breaking from $SO(10) \rightarrow SU(5) \times U(1)$ is possible in principle. Moreover, the spectrum of "light" four-dimensional particles may include relics from split $SO(10)$ supermultiplets, such as the H and $\bar{H}$ in our model. Alternative possibilities using orbifold compactifications⁶⁾ can be extracted from four-dimensional formulations of superstrings⁷⁾. The most general formulation which is now available⁷⁾ offers appetizing possibilities²⁶⁾. It is easy to construct models whose matter fields come from vectors, spinors and singlets of orthogonal groups. This is the case for the matter fields of our $SU(5) \times U(1)$ model, but is not the case for the Higgses of conventional $SU(5)$ and $SO(10)$ models.

We will present the superstring inspiration for our supersymmetric flipped SU(5) model in a future publication. We think the model is sufficiently cute that it warrants more detailed investigation.

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