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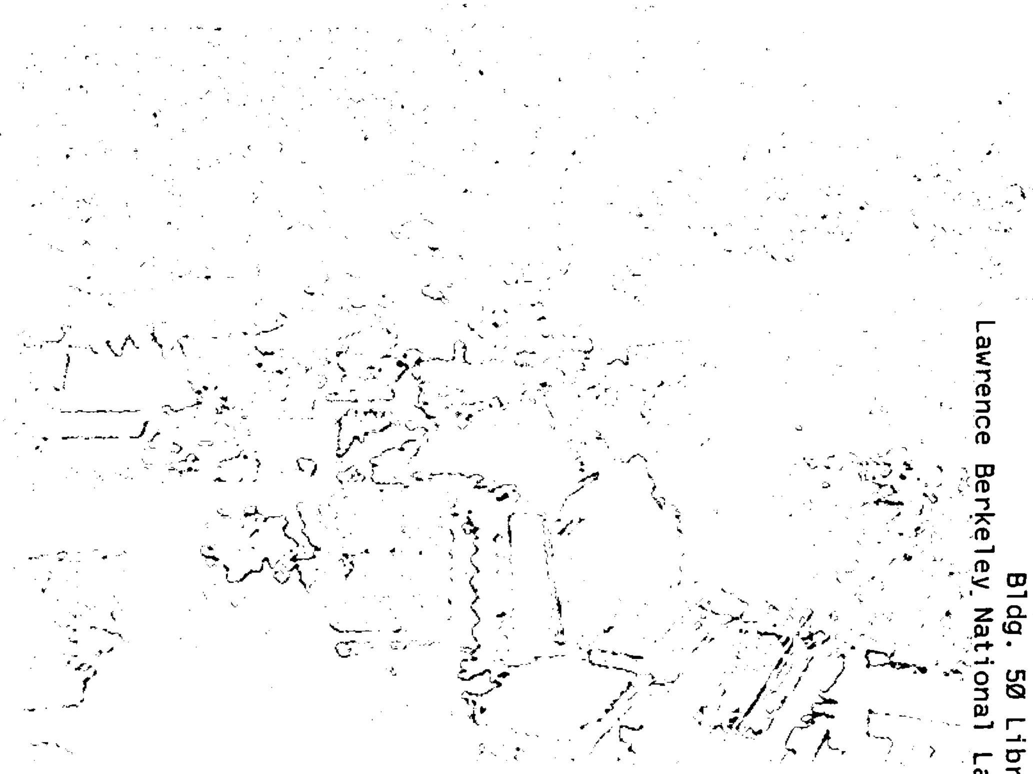

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GUT and SUSY Breaking by the Same Field [1]

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Abstract

We present a model in which the same modulus field breaks both SUSY and a simple GUT gauge group down to the SM gauge group. The modulus is stabilized by the inverted hierarchy mechanism in a perturbative region so that the model is calculable. This is the first example of this kind in the literature. All mass scales (other than the Planck scale) are generated dynamically. In one of the models doublet-triplet splitting is achieved naturally by the sliding singlet mechanism while another model requires fine tuning. The gauge mediation contribution to the right handed slepton $(\text{mass})^2$ is negative. But, for the modulus vacuum expectation value close to the GUT scale, the supergravity contribution to the slepton $(\text{mass})^2$ is comparable to the gauge mediation contribution and thus a realistic spectrum can be attained.

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1 Introduction

One of the central issues in studying supersymmetric extensions of the Standard Model (SM) is how to break supersymmetry (SUSY) and mediate SUSY breaking to the *s*particles. In models of dynamical SUSY breaking, SUSY is broken by the non-perturbative effects of a gauge group. Thus, the SUSY breaking scale is related to the energy scale at which some gauge group becomes strong and, in turn, to the Planck scale by dimensional transmutation. For the mediation of SUSY breaking to the sparticles, two mechanisms have been discussed in the literature – gravity and SM gauge interactions.

The measurements of $\sin^2\theta_W$ are in very good agreement with the predictions of SUSY grand unified theories (GUT's). This has led to a lot of interest in SUSY GUT's. One of the important issues in SUSY GUT's is the origin of the energy scale $\sim 2\times 10^{16}$ GeV at which the GUT symmetry breaks down to the SM.

There have been efforts to generate the GUT scale dynamically. In the models of Cheng [1] and Graesser [2], the GUT scale is related to the dynamical scale of a gauge group, but SUSY breaking is unrelated to GUT symmetry breaking, *i.e.*, there is a separate dynamical scale for SUSY breaking and GUT symmetry breaking.

In the models of Goldberg [3], Kolda and Polonsky [4] and Chacko, Luty and Ponton [5], there is a connection between GUT and SUSY breaking. However, there are two different sectors (and potentials) for GUT and SUSY breaking, but with related parameters (and one dynamical scale). Once SUSY is broken in one sector, a potential is generated for a field in another sector determining the GUT scale. In other words, in these models, the field breaking the GUT symmetry/determining the GUT scale is different from the field breaking SUSY.

In the model of Hirayama, Ishimura and Maekawa [6], the field breaking SUSY and GUT symmetry is the same. However, the GUT gauge group is

$SU(5) \times SU(3) \times SU(2) \times U(1)$, which is *not* a simple gauge group.[4] Also an assumption about a non-calculable Kähler potential is required for the model to work.

In this paper, we present a model in which not only are SUSY breaking and GUT symmetry breaking related, but the *same* field breaks both SUSY and a GUT gauge group down to the SM gauge group. However, unlike the model of reference [6], the GUT gauge group is *simple.* The well known inverted hierarchy mechanism is used to generate a local minimum for the modulus field in a perturbative region, thus making the model calculable, unlike the model of reference [6]. There are no dimensionful parameters in the model other than the Planck scale. The mediation of SUSY breaking to the sparticles is by a combination of gravity and SM gauge interactions.

2 General Structure

The gauge group of the model is:[5]

$$SU(6)_{GUT} \times SU(6)_S \tag{1}$$

and the particle content is

$$\begin{aligned} \Sigma &\sim (\mathbf{35}, \mathbf{1}) \\ Q &\sim (\mathbf{6}, \mathbf{6}) \\ \bar{Q} &\sim (\bar{\mathbf{6}}, \bar{\mathbf{6}}). \end{aligned} \tag{2}$$

The superpotential is

$$W_1 = \lambda_Q \Sigma Q \bar{Q} + \frac{\lambda_\Sigma}{3} \Sigma^3. \tag{3}$$

[4]Thus, the unification of the SM gauge couplings at $\sim 2 \times 10^{16}$ GeV is not an automatic consequence of the model.

[5]This model was used in [7] as a model of gauge mediation. However, in [7], the SM was an additional gauge group, *i.e.*, it was not embedded in the $SU(6)$ gauge symmetry.

Σ^3 lifts all flat directions in Σ except tr Σ^2 [7] along which the vacuum expectation value (vev) of Σ, upto $SU(6)_{GUT}$ rotations, is[6]

$$\langle\Sigma\rangle = \frac{v}{\sqrt{12}} \,\mathrm{diag}[1,1,1,-1,-1,-1]. \tag{4}$$

This can be seen as follows. The vev of Σ breaks $SU(6)_{GUT}$ to $SU(3) \times SU(3) \times U(1)$. The resulting Nambu-Goldstone fields, with their $SU(3) \times SU(3)$ quantum numbers, are:

$$(\mathbf{3},\bar{\mathbf{3}}) + (\bar{\mathbf{3}},\mathbf{3}). \tag{5}$$

Σ decomposes as:

$$(\mathbf{3},\bar{\mathbf{3}}) + (\bar{\mathbf{3}},\mathbf{3}) + (\mathbf{8},\mathbf{1}) + (\mathbf{1},\mathbf{8}) + (\mathbf{1},\mathbf{1}). \tag{6}$$

Thus, the $(\mathbf{3},\bar{\mathbf{3}}) + (\bar{\mathbf{3}},\mathbf{3})$ components of Σ are eaten in the gauge symmetry breaking. The $(\mathbf{1},\mathbf{8}) + (\mathbf{8},\mathbf{1})$ components get a mass from the Σ^3 term and the $(\mathbf{1},\mathbf{1})$ component is the flat direction. Thus, far out along this flat direction, $Q, \bar{Q}$ and all components of Σ other than the flat direction are heavy. The only light fields are the $SU(6)_S$ gauge field and the flat direction parametrized by tr Σ^2. We will denote the flat direction (both the chiral superfield and the vev of it's scalar component) by v. The dynamical scale, Λ_L, of the pure $SU(6)_S$ gauge theory is related to the dynamical scale, Λ, of the high energy $SU(6)_S$ by the matching relation at the mass of $Q, \bar{Q}$ (we assume $v \gg \Lambda$):

$$\left(\frac{\Lambda_L}{\lambda_Q v/\sqrt{12}}\right)^{18} = \left(\frac{\Lambda}{\lambda_Q v/\sqrt{12}}\right)^{12}. \tag{7}$$

Gaugino condensation in the low energy $SU(6)_S$ generates the superpotential:

$$W = 6\Lambda_L^3 = \sqrt{3}\lambda_Q \Lambda^2 v. \tag{8}$$

[6]We use the normalization tr $T_a T_b = 1/2\, \delta_{ab}$, where the T's are the generators for the fundamental representation of a gauge group.

Below the scale Λ_L, we have only the field v with the above superpotential with $F_v = \sqrt{3}\lambda_Q\Lambda^2$. Thus, SUSY is broken and with a canonical Kähler potential, $v^\dagger v$, the vacuum energy is $3\lambda_Q^2\Lambda^4$. The vev v is undetermined at this level. To determine v, we need to include the corrections to the Kähler potential of v. The dominant corrections, for $v \gg \Lambda$, are due to the wavefunction renormalization Z of Σ.[7] Thus, the potential for v is:

$$V = \frac{3\lambda_Q^2\Lambda^4}{Z(v)}. \tag{9}$$

Since $v \gg \Lambda$, we can compute Z in perturbation theory. The one loop Renormalization Group Equation (RGE) for Z is:

$$\frac{dZ(v)}{d(\ln v)} = \frac{2Z(v)}{16\pi^2}\left(12g_6^2(v) - 6\lambda_Q^2(v) - \frac{16}{3}\lambda_\Sigma^2(v)\right), \tag{10}$$

where g_6 is the $SU(6)_{GUT}$ gauge coupling. The potential can develop a minimum by the inverted hierarchy mechanism [8] as follows. We can choose the gauge and Yukawa couplings so that, for large v, λ dominates in the above RGE so that $Z(v)$ decreases with increasing v, whereas, for small v, g_6 dominates so that $Z(v)$ increases with v. Thus, there is a minimum of V at v such that $\lambda(v) \sim g_6(v)$ so that $dZ(v)/d(\ln v) = 0$.[8] Due to the logarithmic dependence of Z, λ and g_6 on v, it is possible that at the minimum $v \gg \Lambda$ which is required for the perturbative calculation to be valid.

To get the SM gauge group from the unbroken gauge group, $SU(3) \times SU(3) \times U(1)$, we identify one $SU(3)$ with $SU(3)_c$ and we need to break the (other) $SU(3) \times U(1)$ to $SU(2)_L \times U(1)_Y$. For achieving this, we use the model in [5] with a slight modification. We next discuss the model.

---

[7]There are corrections to the Kähler potential from higher dimensional operators. But, for $v \gg \Lambda$, these are smaller than the corrections due the wavefunction renormalization [7].

[8]This is a local minimum only since there is a supersymmetric minimum near the origin with $\langle\Sigma\rangle \sim \Lambda$ diag$[2, 2, -1, -1, -1, -1]$ and $\langle Q\bar{Q}\rangle \sim \Lambda^2$ diag$[1, 1, -2, -2, -2, -2]$. However, since $v \gg \Lambda$, the tunneling rate from the "false" vacuum to this global minimum is very small [7].

3 Specific Models

Add the following particle content and superpotential:[9]

$$\begin{array}{rcl} S & \sim & (\mathbf{1},\mathbf{1}) \\ H & \sim & (\mathbf{6},\mathbf{1}) \\ \bar{H} & \sim & (\bar{\mathbf{6}},\mathbf{1}) \end{array} \tag{11}$$

$$W_2 = S(H\bar{H} - \Sigma^2). \tag{12}$$

The F-flatness condition for S forces H and $\bar{H}$ to acquire vev.[10] We look for a minimum with the vev's of $H, \bar{H}$ in the form:

$$\langle H \rangle = \langle \bar{H} \rangle \sim v\,(1,0,0,0,0,0). \tag{13}$$

This breaks $SU(3) \times U(1)$ to $SU(2) \times U(1)$.

We now discuss the mass spectrum. The superpotential has a separate $SU(6)$ symmetry acting on Σ and $H, \bar{H}$. The $SU(6)_H$ is broken to $SU(5)$ resulting in the Nambu-Goldstone fields (with $SU(3)_c \times SU(2)_L$ quantum numbers):

$$(\mathbf{3},\mathbf{1}) + (\bar{\mathbf{3}},\mathbf{1}) + (\mathbf{1},\mathbf{2}) + (\mathbf{1},\mathbf{2}) + (\mathbf{1},\mathbf{1}). \tag{14}$$

The breaking of $SU(6)_\Sigma$ to $SU(3) \times SU(3) \times U(1)$ generates the Nambu-Goldstone fields:

$$(\mathbf{3},\mathbf{2}) + (\bar{\mathbf{3}},\mathbf{2}) + (\mathbf{3},\mathbf{1}) + (\bar{\mathbf{3}},\mathbf{1}), \tag{15}$$

which is the same as Eqn.(5) but with quantum numbers under $SU(3) \times SU(2)$ shown. The following fields are eaten in the breaking of the $SU(6)$ gauge symmetry to the SM gauge group:

$$(\mathbf{3},\mathbf{2}) + (\bar{\mathbf{3}},\mathbf{2}) + (\mathbf{3},\mathbf{1}) + (\bar{\mathbf{3}},\mathbf{1}) + (\mathbf{1},\mathbf{2}) + (\mathbf{1},\mathbf{2}) + (\mathbf{1},\mathbf{1}). \tag{16}$$

---

[9] Henceforth, we will suppress the Yukawa couplings in the superpotential.

[10] In [5], the terms $S(\mathrm{tr}\Sigma^2 - \Phi^2)$ and $T(H\bar{H} - \Phi^2)$ (where S, T, Φ are singlets) were used instead to relate the $H, \bar{H}$ and Σ vev's to the vev of the GUT modulus Φ.

The various fields decompose as:

$$\begin{aligned}
\Sigma &\sim (\mathbf{3},\mathbf{2})+(\bar{\mathbf{3}},\mathbf{2})+(\mathbf{3},\mathbf{1})+(\bar{\mathbf{3}},\mathbf{1})+(\mathbf{8},\mathbf{1})+ \\
&\quad (\mathbf{1},\mathbf{2})+(\mathbf{1},\mathbf{2})+(\mathbf{1},\mathbf{3})+(\mathbf{1},\mathbf{1})+(\mathbf{1},\mathbf{1}) \\
H &\sim (\mathbf{3},\mathbf{1})+(\mathbf{1},\mathbf{2})+(\mathbf{1},\mathbf{1}) \\
\bar{H} &\sim (\bar{\mathbf{3}},\mathbf{1})+(\mathbf{1},\mathbf{2})+(\mathbf{1},\mathbf{1}).
\end{aligned} \tag{17}$$

As mentioned before, the $(\mathbf{8},\mathbf{1})+(\mathbf{1},\mathbf{2})+(\mathbf{1},\mathbf{2})+(\mathbf{1},\mathbf{3})+(\mathbf{1},\mathbf{1})$ components of Σ (which transform as $(\mathbf{8},\mathbf{1})+(\mathbf{1},\mathbf{8})$ under $SU(3)\times SU(3)$: see Eqn.(6)) get a mass from the Σ^3 term. The $(\mathbf{3},\mathbf{2})+(\bar{\mathbf{3}},\mathbf{2})$ components of Σ and the $(\mathbf{1},\mathbf{2})+(\mathbf{1},\mathbf{2})$ components of $H,\bar{H}$ are eaten by the broken gauge symmetry (see Eqn.(16)). From Eqns.(14) and (15) there are two pairs of Nambu-Goldstone triplets in Σ and $H,\bar{H}$. From Eqn.(16) only one combination of these two pairs is eaten.[11] The other combination is massless. The remaining SM singlet in Σ is the flat direction tr Σ^2. One combination of the SM singlets in $H,\bar{H}$ is eaten by the broken symmetry (see Eqn.(16)) or in other words is constrained by the D-flatness condition. The other combination is parametrized by $H\bar{H}$. The singlet S marries one combination of tr Σ^2 and $H\bar{H}$ due to the superpotential W_2. The orthogonal combination of Σ^2 and $H\bar{H}$ is massless. Thus, the massless fields are this flat direction and a pair of triplets in Σ, H and $\bar{H}$.

To make these triplets heavy[12] , we can use the sliding singlet mechanism [9, 5]. Add the following to the superpotential:

$$W_3 = H(\Sigma + X)\bar{h} + \bar{H}(\Sigma + \bar{X})h, \tag{18}$$

where

$$X \sim (\mathbf{1},\mathbf{1})$$

[11]Without the $H,\bar{H}$, the $(\mathbf{3},\mathbf{1})+(\bar{\mathbf{3}},\mathbf{1})$ components of Σ, which along with the $(\mathbf{3},\mathbf{2})+(\bar{\mathbf{3}},\mathbf{2})$ components form $(\mathbf{3},\bar{\mathbf{3}})+(\bar{\mathbf{3}},\mathbf{3})$ under $SU(3)\times SU(3)$, are eaten as mentioned before (see Eqns.(5) and (6)).

[12]Giving mass to these Nambu-Goldstone triplets is equivalent to getting the orientation of the Σ and $H,\bar{H}$ vev's in Eqns.(4) and (13).

$$
\begin{aligned}
\bar{X} &\sim (\mathbf{1}, \mathbf{1}) \\
h &\sim (\mathbf{6}, \mathbf{1}) \\
\bar{h} &\sim (\bar{\mathbf{6}}, \mathbf{1}).
\end{aligned} \tag{19}
$$

$F_X = F_{\bar{X}} = 0$ forces $h = \bar{h} = 0$. $F_h = F_{\bar{h}} = 0$ along with the form of the $H, \bar{H}$ vev's makes the singlets slide so that $X = \bar{X} = -v/\sqrt{12}$. Thus, the form of the $(\Sigma + X)$ vev is such that the triplets in $H, \bar{H}$ get a mass with the triplets in $h, \bar{h}$. There is no mass term for the doublets in $h, \bar{h}$ with those in $H, \bar{H}$. However, the $H, \bar{H}$ vev's with the above superpotential give a mass term for the doublets (and also the triplets) in Σ with those in $h, \bar{h}$ (there is also a mass term for the doublets in Σ from the Σ^3 term). Also, the $H, \bar{H}$ vev's give mass to the first (SM singlet) components of $h, \bar{h}$ with combinations of tr Σ^2 and $X, \bar{X}$. Thus, the only massless field is the flat direction which is now a combination of tr Σ^2, $H\bar{H}, X$ and $\bar{X}$. Along this flat direction, both SUSY and the GUT symmetry are broken.

To get the usual pair of light Higgs doublets, we duplicate the above structure of $S, H, \bar{H}, h, \bar{h}, X$ and $\bar{X}$ [9, 5]. The superpotential is:

$$
W_2 + W_3 = \sum_{i=1}^{2} S_i (H_i \bar{H}_i - \Sigma^2) + \sum_{i=1}^{2} H_i (\Sigma + X_i) \bar{h}_i + \sum_{i=1}^{2} \bar{H}_i (\Sigma + \bar{X}_i) h_i. \tag{20}
$$

$F_{S_2} = 0$ forces $H_2 \bar{H}_2 = \Sigma^2$. We look for a minimum with the vev's of $H_2, \bar{H}_2$ aligned with $H_1, \bar{H}_1$, *i.e.*, $H_2 = \bar{H}_2 \sim v(1,0,0,0,0,0)$. Then, as before, the sliding singlet mechanism gives mass for the triplets in $H_2, \bar{H}_2$ with those in $h_2, \bar{h}_2$. As before, due to the vev's of $H_2, \bar{H}_2$, the SM singlets in $h_2, \bar{h}_2$ get a mass with two combinations of tr Σ^2 and $X_2, \bar{X}_2$. Thus, the flat direction is now a combination of tr Σ^2, $H_i \bar{H}_i$, X_i and $\bar{X}_i$ with $i = 1, 2$. Only one combination of the doublets in h_1, h_2 marries the doublet in Σ due to the $\bar{H}$ vev's (similarly for the doublets in $\bar{h}_{1,2}$). This leaves one pair of massless doublets in the $h, \bar{h}$'s which can be the usual Higgs doublets. There is also a pair of massless doublets in the H's since only one pair is eaten in the gauge symmetry breaking (see Eqn.(16)). Also, there is a massless SM singlet in the H's which can be seen as follows. The $H, \bar{H}$'s have four SM singlets. The $F_S = 0$ conditions relate two combinations of these, namely $H_1\bar{H}_1$ and $H_1\bar{H}_2$, to Σ^2. One combination is eaten by the broken gauge symmetry (see Eqn.(16)); in other words, one combination of the vev's is constrained by the D-flatness condition. This leaves one combination of the vev's unconstrained, *i.e.*, one massless SM singlet in $H, \bar{H}$'s. We discuss two ways to give mass to the extra pair of doublets and the SM singlet in $H, \bar{H}$.[13]

In the first model we add the superpotential $W_4 + W_5$ where:

$$W_4 = \frac{1}{M}\left(\left(H_1\bar{H}_1\right)\left(H_2\bar{H}_2\right) - \left(H_1\bar{H}_2\right)\left(H_2\bar{H}_1\right)\right), \tag{21}$$

with, say, $M = M_{Pl}$ and

$$W_5 = S_3\left(H_1\bar{H}_2 - H_2\bar{H}_1\right), \tag{22}$$

where S_3 is a singlet. $W_2 + W_3 + W_5$ is invariant under $(H, \bar{H}, h, \bar{h}, S, X, \bar{X})_1 \leftrightarrow (H, \bar{H}, h, \bar{h}, S, X, \bar{X})_2$ and $S_3 \leftrightarrow -S_3$ and W_4 is invariant under two global $SU(2)$ symmetries – one with (H_1, H_2) as a doublet and the other with $\left(\bar{H}_1, \bar{H}_2\right)$ as a doublet. [14] We look for a minimum with $\langle S_3 \rangle = 0$. F_{S_3} gives an additional constraint between the $H, \bar{H}$ vev's giving a mass (with S_3) to the SM singlet mentioned above. The doublets in $H, \bar{H}$ have a mass matrix of the form [9]:

$$\begin{pmatrix} \langle H_2\bar{H}_2\rangle & -\langle H_1\bar{H}_2\rangle \\ -\langle H_2\bar{H}_1\rangle & \langle H_1\bar{H}_1\rangle \end{pmatrix}, \tag{23}$$

which has one zero eigenvalue corresponding to the eaten pair of doublets and one non-zero eigenvalue $\sim M_{GUT}^2/M$ which is the mass for the other pair of doublets. This shifts the prediction of $\sin^2\theta_W$ by about $+3 \times 10^{-3}$ if $\alpha_s(m_Z)$ and $\alpha_{em}(m_Z)$ are used as inputs.

[13] Giving mass to the extra pair of doublets and the SM singlet in $H, \bar{H}$ is equivalent to getting the alignment of the $H_2, \bar{H}_2$ vev's with the $H_1, \bar{H}_1$ vev's.

[14] Otherwise, we have to tolerate some fine tuning to get this form of the superpotential.

In the other method [5], we add the terms:

$$W_4' = (X_1 + X_2)\Delta^2 + \left(H_2 \Delta \bar{H}_1 - H_1 \Delta \bar{H}_2\right), \tag{24}$$

where Δ is a $\mathbf{35}$ of $SU(6)$. $W_2 + W_3 + W_4'$ is invariant under the symmetry $(H, \bar{H}, h, \bar{h}, S, X, \bar{X})_1 \leftrightarrow (H, \bar{H}, h, \bar{h}, S, X, \bar{X})_2$ and $\Delta \leftrightarrow -\Delta$. We look for a minimum with the vev of $\Delta = 0$ so that the F_X and F_H-flatness conditions are not affected. The vev's of $X_{1,2}$ give mass to Δ. $F_\Delta = 0$ gives a constraint between the vev's of the $H, \bar{H}$'s giving a mass (with a singlet in Δ) to the SM singlet mentioned above. Due to the $H, \bar{H}$ vev's, the massless pair of doublets in the H's gets a mass with those in Δ. Thus, the only massless field is the flat direction which breaks both SUSY and the GUT symmetry.

If we are willing to tolerate fine tuning to "solve" the usual doublet-triplet splitting problem to get a pair of light doublets, we can gauge only the $SU(5)$ subgroup of the $SU(6)$. Then, with only the Σ field and W_1, the generators of the global $SU(6)$ in Eqn.(15) are broken ($SU(6)_{global}$ is broken to $SU(3) \times SU(3) \times U(1)$). Of these generators, only $(\mathbf{3}, \mathbf{2}) + (\bar{\mathbf{3}}, \mathbf{2})$ are gauged. Thus, $SU(5)_{local}$ is broken down to the SM. We get a pair of massless triplets in Σ corresponding to the broken generators which are not gauged. These can be given a mass by adding:

$$H(\lambda_1 \Sigma_1 + \lambda_{24} \Sigma_{24}) \Sigma_{\bar{5}} + \bar{H}(\bar{\lambda}_1 \Sigma_1 + \bar{\lambda}_{24} \Sigma_{24}) \Sigma_5, \tag{25}$$

where $H, \bar{H}$ are fundamentals of $SU(5)$ and $\Sigma_5, \Sigma_{\bar{5}}, \Sigma_1$ and Σ_{24} denote components of Σ transforming as $\mathbf{5}, \bar{\mathbf{5}}, \mathbf{1}$ and $\mathbf{24}$, respectively, under $SU(5)$.[15] Since $\langle \Sigma_1 \rangle \sim \mathrm{diag}[1, 1, 1, 1, 1]$ and $\langle \Sigma_{24} \rangle \sim \mathrm{diag}[-3, -3, 2, 2, 2]$ (in $SU(5)$ space), we can fine tune the couplings $\lambda, \bar{\lambda}$ so that there is a mass term for the triplet in H ($\bar{H}$) with the triplet in $\Sigma_{\bar{5}}$ (Σ_5) but not for the doublets. Then, the doublets in $H, \bar{H}$ can be the usual Higgs doublets. [16]

[15] The superpotential in Eqn.(3) is invariant under the $SU(6)$ global symmetry whereas the one in Eqn.(25) is only $SU(5)_{local}$ invariant.

[16] The doublets in $\Sigma_{5,\bar{5}}$ get a mass from the Σ^3 term as before.

In all these models, the μ term has to be generated by some mechanism. Also, these models are only technically natural, *i.e.*, the superpotential is not the most general one allowed by symmetries. For example, in the model with the full $SU(6)$ symmetry gauged, we need the terms $SH\bar{H}$, tr Σ^3 and S trΣ^2 and so the term $H\Sigma\bar{H}$ is also allowed which is undesirable. So, these models should be viewed as existence proofs of models in which both a simple GUT gauge group and SUSY are broken by the same field.

4 MSSM Spectrum

4.1 Quarks and Leptons

The SM fermion Yukawa couplings can be generated using the method in [5] as follows. Add the following fields charged under $SU(6)_{GUT}$ and superpotential:

$$\begin{aligned} N_i &\sim \mathbf{15} \\ \bar{P}_{1i}, \bar{P}_{2i} &\sim \bar{\mathbf{6}} \\ Y &\sim \mathbf{15} \\ \bar{Y} &\sim \bar{\mathbf{15}} \end{aligned} \tag{26}$$

$$\begin{aligned} W_{Yukawa} = {} & N_i(\bar{P}_{1j}\bar{H}_1 + \bar{P}_{2j}\bar{H}_2) + N_i(\bar{P}_{1j}\bar{h}_1 + \bar{P}_{2j}\bar{h}_2) \\ & + N_i N_j Y + (X_1 + X_2)Y\bar{Y} + \bar{Y}(H_1 h_1 - H_2 h_2), \end{aligned} \tag{27}$$

where $i, j = 1, 2, 3$ are generation indices. This superpotential is invariant under the symmetry $(H, \bar{H}, h, \bar{h}, X)_1 \leftrightarrow (H, \bar{H}, h, \bar{h}, X)_2$ and $\bar{P}_1 \to i\bar{P}_2$, $\bar{P}_2 \to i\bar{P}_1$, $N \to -iN$, $Y \to -Y$ and $\bar{Y} \to -\bar{Y}$. For each generation, the $N\bar{P}\bar{H}$ terms make the $\mathbf{5}$ (under $SU(5)$) of the N and one combination of the $\bar{\mathbf{5}}$'s of $\bar{P}_{1,2}$ heavy, leaving the usual $\bar{\mathbf{5}} + \mathbf{10}$ massless. The $N\bar{P}\bar{h}$ terms give the down quark and lepton Yukawa couplings whereas the up quark Yukawa couplings arise from the terms NNY and $\bar{Y}Hh$ after integrating out the $Y, \bar{Y}$ fields.

4.2 Sparticle Spectrum

There is a gauge mediation (GM) contribution to the sparticle masses. The model has both "matter" messengers (the $Q, \bar{Q}$ fields and the heavy components of H, h's) and "gauge" messengers (the heavy gauge multiplets which have a non-supersymmetric spectrum since the field breaking the GUT symmetry has a non-zero F-component). We compute the sparticle spectrum using the method of [10]. In this method, the scalar (mass)^2, m_i^2, are computed from the RG scaling of the wavefunctions of the matter fields and the gaugino masses, M_A, are related to the RG scaling of the gauge couplings. The expressions for the masses are:

$$M_A(\mu) = \frac{\alpha_A(\mu)}{4\pi} \frac{F_v}{v} (b_A - b_6) , \tag{28}$$

where b_A's are the beta functions of the SM gauge couplings below the GUT scale and b_6 is the beta function of the $SU(6)_{GUT}$ above the GUT scale, and

$$\begin{aligned} m_i^2(\mu) &= \frac{1}{16\pi^2} \left(\frac{F_v}{v} \right)^2 \times \\ & \left(\sum_A \frac{2C_A^i}{b_A} \left(\alpha_A^2(\mu) \left(b_6 - b_A \right)^2 - b_6^2 \alpha_6^2 \right) + 2C_6^i b_6 \alpha_6^2 \right) , \end{aligned} \tag{29}$$

where C_A^i is the quadratic Casimir invariant for the scalar i under the gauge group A, *i.e.*, 4/3, 3/4 for fundamentals of $SU(3)_c$, $SU(2)_L$ respectively and 3/5 Y^2 for $U(1)_Y$. $C_6^i = 35/12$ for fields in $\bar{\mathbf{5}}$ of $SU(5)$ ($\bar{\mathbf{6}}$ of $SU(6)_{GUT}$) and 14/3 for fields in **10** of $SU(5)$ (**15** of $SU(6)_{GUT}$). The beta function for $SU(N_c)$ group is defined as $3N_c - N_{f,eff}$, where N_{eff} is the "effective" number of flavors. α_6 is the $SU(6)$ coupling at the GUT scale. The messengers do *not* form complete $SU(5)$ representations and thus the above mass spectrum is different from the models of gauge mediation with complete $SU(5)$ multiplets as messengers. For example, the gaugino masses are *not* unified at the GUT scale.

The above results depend on the beta functions of the SM gauge group below the GUT scale and the beta function of $SU(6)_{GUT}$ above the GUT scale. We assume that there are no particles with SM quantum numbers between the weak and the GUT scales so that $b_{1,2,3}$ are the usual MSSM beta functions. The $SU(6)$ beta function, b_6, depends on the particle content at the GUT scale and thus, in turn, on the method used to generate SM fermion Yukawa couplings and the method used to make the extra pair of doublets in $H, \bar{H}$ heavy. We consider the case where the above method is used to generate SM fermion Yukawa couplings and the higher dimensional operator (Eqn.(21)) is used to make the extra doublets heavy. In this case, the beta function b_6 (defined as $3N_c - N_{f,eff}$) is -11. We get $m^2_{\tilde{e}_R}(\mu \sim m_Z) \approx -8 \times 10^{-4} \left(F_v/v\right)^2$, whereas all other scalar (mass)^2 are positive. We have to add the supergravity (SUGRA) contribution to the (mass)^2 $\sim \left(F_v/M_{Pl}\right)^2$ where $M_{Pl} \sim 2\times 10^{18}$ GeV.[17] For $v \sim 6\times 10^{16}$ GeV, the two contributions to $m^2_{\tilde{e}_R}$ are comparable and thus we can get a phenomenologically acceptable spectrum.[18] However, since the supergravity contribution is comparable to the flavor blind GM contribution, we need to impose some flavor symmetries or alignment (of the SUGRA contribution with the Yukawa couplings) to avoid too large SUSY contributions to FCNC's. For the squarks, the GM contribution is larger so that less degeneracy is required in the SUGRA contribution.

5 Inverted Hierarchy

Since the flat direction is a combination of the fields tr $\Sigma^2, H_i\bar{H}_i, X_i$ and $\bar{X}_i$ $(i = 1, 2)$, the RGE analysis for the wavefunction of the flat direction involves too many Yukawa couplings. To simplify the analysis, we assume the the vev's of all the fields in the flat direction are of the same order and that among the Yukawa couplings, only the $\Sigma Q\bar{Q}$ coupling is large. The

[17]We assume that the SUGRA contribution to the (mass)^2 is positive.

[18]It might seem that this value of v is a bit larger than the "usual" GUT scale $\sim 2\times 10^{16}$ GeV. However, as mentioned earlier, the flat direction v is really a combination of ~ 7 fields. If we assume that all these fields have roughly the same vev, then the vev of each field, in particular, the Σ, H fields is $\sim v/\sqrt{7}$ which is closer to the usual GUT scale.

$SU(6)_{GUT}$ coupling at the Planck scale is fixed with the assumption of a desert between the weak and the GUT scales and the particle content at the GUT scale. We require $(F_v/v) \sim 10$ TeV to get the sparticle masses ~ 100 GeV to 1 TeV. With $v \sim 10^{16}$ GeV, this determines $F_v \sim \Lambda^2$ and hence the $SU(6)_S$ gauge coupling at the Planck scale. Then, we checked that for the $\Sigma Q\bar{Q}$ coupling ~ 2 at the Planck scale, we do get a minimum of the potential at around the GUT scale.

There is also a SUGRA contribution to the $(\text{mass})^2 \sim (F_v/M_{Pl})^2$ of the flat direction. For $v \sim 10^{16}$ GeV, we expect this to be comparable to the $(\text{mass})^2$ due to the inverted hierarchy which is $\sim -F_v^2/v^2\, d^2Z(v)/d(\ln v)^2$. It turns out that in this case the SUGRA contribution is smaller (by a factor of ~ 4) than the $(\text{mass})^2$ due to the inverted hierarchy. This results in a shift of the minimum of v by $\sim O(1/4)\, v$.

To summarize, we have presented a model in which the field breaking SUSY is the same as the field which breaks a simple GUT gauge group to the SM gauge group. The model is calculable – it uses the inverted hierarchy mechanism to generate a minimum for the field in a perturbative region. As far as we know, this is the first example of such a kind in the literature.

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