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Sliding singlet mechanism in N=1 supergravity GUT

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ABSTRACT

We show how the sliding singlet mechanism may be successfully incorporated in SU(6) grand unified theories coupled to N=1 supergravity, to keep the weak doublet higgs light, while its color triplet partner acquires a large mass.

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Grand unified theories based on N=1 supergravity, where supersymmetry is broken at 10^{12}-10^{13} GeV in a hidden sector, has become popular recently[1]. The simplest description of the model is obtained if we take the limit $M_{Planck} \to \infty$, keeping the gravitino mass m_g fixed. In this limit, if $W(\phi)$ denotes the superpotential involving the observable sector superfields ϕ_i, the effective action of the theory is given by the action for a globally supersymmetric theory with superpotential W, together with explicit soft supersymmetry breaking terms of the form,

$$\{ m_g \sum_i z_i \frac{\partial W}{\partial z_i} + m_g (A-3)\, W(z) \} + h.c. - m_g^2 \sum_i |z_i|^2 \qquad (1)$$

where z_i's are the scalar components of the superfields.

Although supersymmetry, if unbroken, maintains any mass hierarchy that is present at the tree level[2], it does not tell us why some mass scales are smaller than others at the tree level. For example, it does not explain why the weak doublet higgs is so light compared to its color triplet partner. One explanation is provided by the missing partner mechanism[3], which needs the introduction of higggs fields belonging to the large representations (50 and $\overline{50}$) of SU(5). Another explanation is based on the sliding singlet mechanism[4], which is the topic of discussion of the present paper. This mechanism, when incorporated in the SU(5) gauge theories, has been shown to be unstable under radiative corrections[5-8]. We shall show that there exist models based on the SU(6) gauge group, with soft supersymmetry breaking terms induced by N=1 supergravity, where the sliding singlet mechanism may be successfully incorporated. A different type of globally supersymmetric SU(6) model with arbitrary soft supersymmetry terms of dimension two has been proposed by Dimopoulos and Georgi[9], which also uses the sliding singlet mechanism.

Before writing down the SU(6) model, we shall illustrate the problem of using the sliding singlet mechanism in the SU(5) model. We consider an SU(5) grand unified theory with fields $\Phi(24)$, $S(1)$, $H(5)$ and $\widetilde{H}(\bar{5})$. The superpotential W is taken as,

$$W = \alpha S H \widetilde{H} + \beta \Phi H \widetilde{H} + \lambda (\Phi^3 + M_1 \Phi^2) + M_2 H \widetilde{H} \qquad (2)$$

where, for simplicity, we have dropped the SU(5) indices. If we ignore the supersymmetry breaking terms, the tree level potential is given by,

$$V = \sum_i |\partial W / \partial z_i|^2 + \sum_a |z^\dagger T_a z|^2$$

$$= |(3\lambda\Phi^2 + 2\lambda M_1 \Phi - \beta H \widetilde{H})_{24}|^2 + |(\alpha S + \beta \Phi + M_2) H|^2$$

$$+ |\widetilde{H}(\alpha S + \beta \Phi + M_2)|^2 + |(\alpha H \widetilde{H})_1|^2 + \sum_a |z^\dagger T_a z|^2 \qquad (3)$$

where T_a's denote the generators of the group, and z the set of all scalar fields. The potential has a local minimum at,

$$\langle \Phi \rangle = \frac{4}{3} M_1 \begin{pmatrix} 1 & & & & \\ & 1 & & & \\ & & 1 & & \\ & & & -3/2 & \\ & & & & -3/2 \end{pmatrix}$$

$$\langle H \rangle = \langle \tilde{H} \rangle = 0$$

$$\langle S \rangle \text{ arbitrary} \tag{4}$$

If S takes the vacuum expectation value (vev),

$$\langle S \rangle = (-\beta \langle \Phi_{55} \rangle - M_2)/\alpha = (2\beta M_1 - M_2)/\alpha \tag{5}$$

then the weak doublet higgs is massless, whereas its color triplet partner has mass of order $M_1 \sim M$. For a general vev of S, both, the weak doublet and the color triplet higgses have mass of the order of the grand unification mass M. It is argued that whatever mechanism makes the weak doublet higgs to acquire a vev of order 100 GeV (which may be achieved, for example, by adding explicit negative mass^2 term in the Lagrangian for H and $\tilde{H}$), makes it energetically favorable for S to take the vev given in (5), so as to minimize the second and third term of (3).

This scenario, however, breaks down under radiative corrections[5-8]. This effect is most suitably described in the superfield formulation[9]. The presence of soft breaking terms in the Lagrangian induce radiative corrections of the form,

$$-(a\, m_g F_s^* M + b\, m_g^2 S^* M + h.c.) \tag{6}$$

where a and b are two constants, and F_S is the auxiliary field corresponding to the singlet field S. If for the sake of simplicity, we include the terms of (6) in the effective action, but ignore the tree level soft breaking terms in the potential,[F1] we get the effective potential as,

$$V = |(3\lambda\Phi^2 + 2\lambda M_1\Phi - \beta H\tilde{H})_{24}|^2 + |(\alpha S + \beta\Phi + M_2)H|^2$$
$$+ |\tilde{H}(\alpha S + \beta\Phi + M_2)|^2 + |(\alpha H\tilde{H} - a^* m_g M)|^2$$
$$+ (b\, m_g^2 S^* M + h.c.) + \sum_a |z^\dagger T_a z|^2 \qquad (7)$$

which is obtained by eliminating the auxiliary field F_S from the effective action. The fourth term generates a mass term of H, H of order am_gM, and thus destroys the mass hierarchy, even if S adjusts itself to have the vev given in Eq. (5).

Even if a vanishes, the $bm_g^2S^*M$ term destroys the mass hierarchy as follows. If S takes the value given in (5), and H, H acquire vev of order $m_g \sim 100$ GeV, the potential as a function of S has the form,

$$\{|\langle H\rangle|^2 + |\langle\tilde{H}\rangle|^2\}\,|(\alpha S - 2\beta M_1 + M_2)|^2 + (b\, m_g^2 S^* M + h.c.) \qquad (8)$$

Minimizing the above potential with respect to S we find that $(\langle S\rangle - (2\beta/\alpha)M_1 + M_2/\alpha) \sim bM$, which is inconsistent with the original assumption that the minimum of the potential lies at $S=(2\beta/\alpha)M_1-(M_2/\alpha)$.

We shall now demonstrate how the SU(6) model solves these problems. The minimal model of this kind consists of the fields $\Phi(35)$, $\Phi_0(1)$, $S^{(i)}(1)$, $H^{(i)}(6)$ and $\tilde{H}^{(i)}(\bar{6})$ $(i=1,2)$. The superpotential is given by,

$$W=(\lambda_1\Phi^3+M_1\Phi^2+\lambda_2\Phi_0\Phi^2+M_2{}^2\Phi_0)$$
$$+\sum_{i=1}^{2}\{\alpha_1^{(i)}\Phi H^{(i)}\tilde{H}^{(i)}+\alpha_2^{(i)}\Phi_0 H^{(i)}\tilde{H}^{(i)}+\alpha_3^{(i)}S^{(i)}H^{(i)}\tilde{H}^{(i)}\} \tag{9}$$

For simplicity, we have ignored the quark-lepton fields. The detailed discussion of a realistic model has been given elsewhere[10]. Including the soft supersymmetry breaking terms induced by N=1 supergravity, we get the total potential as,

$$V=|(3\lambda_1\Phi^2+2M_1\Phi+2\lambda_2\Phi_0\Phi+\sum_{i=1}^{2}\alpha_1^{(i)}H^{(i)}\tilde{H}^{(i)}-m_g\Phi^*)_{35}|^2$$
$$+|(\lambda_2\Phi^2+M_2{}^2+\sum_{i=1}^{2}\alpha_2^{(i)}H^{(i)}\tilde{H}^{(i)}-m_g\Phi_0^*)_1|^2$$
$$+\sum_{i=1}^{2}\{|\tilde{H}^{(i)}(\alpha_1^{(i)}\Phi+\alpha_2^{(i)}\Phi_0+\alpha_3^{(i)}S^{(i)})-m_g H^{(i)*}|^2$$
$$+|(\alpha_1^{(i)}\Phi+\alpha_2^{(i)}\Phi_0+\alpha_3^{(i)}S^{(i)})H^{(i)}-m_g\tilde{H}^{(i)*}|^2$$
$$+|(\alpha_3^{(i)}\tilde{H}^{(i)}H^{(i)}-m_g S^{(i)*})_1|^2\}-\{m_g(A-3)W+h.c.\}$$
$$+\sum_a|z^\dagger T_a z|^2 \tag{10}$$

The above potential has a local minimum at,[F2]

$$\langle \Phi \rangle = \frac{M_2}{\sqrt{6}\lambda_2} \begin{pmatrix} 1 & & & & & \\ & 1 & & & & \\ & & 1 & & & \\ & & & -1 & & \\ & & & & -1 & \\ & & & & & -1 \end{pmatrix} + O(m_g) \tag{11}$$

$$\langle \Phi_0 \rangle = -M_1/\lambda_2 + O(m_g) \tag{12}$$

$$\langle S^{(i)} \rangle = -\{\alpha_1^{(i)} \langle \Phi_{66} \rangle + \alpha_2^{(i)} \langle \Phi_0 \rangle\}/\alpha_3^{(i)} + O(m_g) \tag{13}$$

$$\langle H_6^{(i)} \rangle = \langle \tilde{H}_6^{(i)} \rangle = \{m_g \langle S^{(i)} \rangle^* / \alpha_3^{(i)}\}^{1/2} + O(m_g{}^2/\sqrt{m_g M}) \tag{14}$$

$$\langle H_\ell^{(i)} \rangle = \langle \tilde{H}_\ell^{(i)} \rangle = 0 \qquad 1 \leq \ell \leq 5 \tag{15}$$

Note that the vev of $S^{(i)}$ is no longer arbitrary after we include the soft supersymmetry breaking terms in the potential. There is a nearly degenerate minimum where $S^{(i)}$, $H^{(i)}$ and $\tilde{H}^{(i)}$ acquire zero vev, but the degeneracy is lifted by the radiative correcions[10]. If we ignore the contribution to the potential from the D terms (the last term of Eq.(10)), then at this minimum the $SU(2)^{weak}$ doublet parts of $H^{(i)}$ and $\tilde{H}^{(i)}$ acquire masses $\lesssim m_g$. One linear combination of these fields turn out to be exactly massless, and is absorbed by the gauge bosons through higgs mechanism. Another linear combination gets mass of order $\sqrt{m_g M}$ from the D-term of the potential,and becomes degenerate with the gauge bosons. But two other linear combinations of $H^{(i)}$, $H^{(i)\dagger}$ $(i=1,2)$ remain light$(\sim m_g)$, which may serve as the higgs bosons of the $SU(2)^{weak}\times U(1)$ symmetry breaking. In this scenario, $SU(6)$ group is broken to $SU(3)\times SU(3)\times U(1)$ at a scale of order $M(10^{16}-10^{17}\text{GeV})$, which is broken to $SU(3)\times SU(2)\times U(1)$ at a scale of order $\sqrt{m_g M}\sim 10^{10}\text{GeV}$. The $SU(3)\times SU(2)\times U(1)$ symmetry may then be broken to $SU(3)$ $U(1)$ at a scale of order $m_g(\sim 100\text{GeV})$ through radiative corrections, as has been discussed by several authors[11]. There are other nearly degenerate minima of the potential with different symmetries, but for suitable choices of the parameters of the theory, the minimum discussed above may be shown to be of lowest energy[10].

We may now see the effect of including radiatively induced terms of the form,

$$-\sum_{i=1}^{2}\left(m_g\, a^{(i)} F_{S^{(i)}}^{*} M + m_g^2\, b^{(i)} S^{(i)*} M + h.c.\right) \qquad (16)$$

in the effective action. In the corresponding potential, obtained after eliminating the auxiliary fields, the $\Sigma|\alpha_3^{(i)}H^{(i)}H^{(i)}-m_g S^{(i)*}|^2$ term is replaced by,

$$\sum_{i=1}^{2}\{|(\alpha_3^{(i)}H^{(i)}\tilde{H}^{(i)} - m_g S^{(i)*} + a^{(i)*}m_g M)_1|^2$$

$$+ (b^{(i)}m_g^2 M S^{(i)*} + a^{(i)}m_g^2 M S^{(i)*} + h.c.) \qquad (17)$$

The values of $<\Phi>$, $<\Phi_0>$ and $<S^{(i)}>$ remain unchanged except for terms of order m_g, whereas we now have,

$$<H_6^{(i)}> = <\tilde{H}_6^{(i)}> = \sqrt{\{m_g S^{(i)*} - a^{(i)*}m_g M\}/\alpha_3^{(i)}} + O(m_g^2/\sqrt{m_g M})$$

$$(18)$$

Thus the presence of the $a^{(i)}$ term changes the vev of $H_6^{(i)}$ and $\tilde{H}_6^{(i)}$, but does not produce large mass of $H_{4,5}^{(i)}$ or $\tilde{H}_{4,5}^{(i)}$, and we still have a pair of low mass weak doublet higgs.

To see that the presence of the new radiatively induced terms does not shift $<S^{(i)}>$ from the value given in (13) by more than order m_g, we note that for $<H_6^{(i)}>$, $<H_6^{(i)}>$ given by (18), the dependence of the potential on $S^{(i)}$ is given by,

$$\sum_{i=1}^{2}[(|<H_6^{(i)}>|^2 + |<\tilde{H}_6^{(i)}>|^2)|\alpha_3^{(i)}S^{(i)} + \alpha_1^{(i)}<\Phi_{66}> + \alpha_2^{(i)}<\Phi_0> - m_g|^2$$

$$+ m_g^2(b^{(i)}S^{(i)*}M + a^{(i)}S^{(i)*}M + h.c.)$$

$$-\{m_g(A-3)<H_6^{(i)}><\tilde{H}_6^{(i)}>S^{(i)} + h.c.\}] \qquad (19)$$

Minimizing this term with respect to $S^{(i)}$, and using the fact that $|\langle H_6^{(i)}\rangle|^2 \sim m_g M$, we see that,

$$\langle S^{(i)}\rangle + \frac{\alpha_1^{(i)}}{\alpha_3^{(i)}} \langle \Phi_{66}\rangle + \frac{\alpha_2^{(i)}}{\alpha_3^{(i)}} \langle \Phi_0\rangle \sim m_g \qquad (20)$$

and hence the mass hierarchy is preserved. In the component field language, this may be explained by the fact that at the minimum where $\langle H_6^{(i)}\rangle$, $\langle \tilde{H}_6^{(i)}\rangle$ are of order $\sqrt{m_g M}$, $S^{(i)}$ acquires a mass of order $\sqrt{m_g M}$, and hence does not remain light any more. Also in this model, the $m_g(A-3)W$ term gives rise to a term of order $m_g^2(A-3)H\tilde{H}$ in the potential, due to the appearance of terms of order m_g on the right hand sides of Eqs.(13) and (20). (This term appears even at the tree level). Hence this class of models also provides a solution of the so called 'μ-problem'[12].

FOOTNOTES

[F1]In fact, if we include all the supersymmetry breaking terms given in Eq.(1), this model is unstable even at the tree level. This is due to the mass term of the singlet field S, and the cubic supersymmetry breaking terms. We may, however, consider the effect of arbitrary soft breaking terms which do not destroy the hierarchy at the tree level.

[F2]The potential may be minimized by first minimizing Eq.(10) without the $m_g(A-3)W+h.c.$ term. Taking this solution as the first approximation, we find the true minimum of the potential to any accuracy by iterative procedure.

REFERENCES

[1] H. P. Nilles, 'Supersymmetry, supergravity and particle physics', Universite de Geneve report No. UGVA-DPT 1983/12-412, and references therein.

[2] S. Dimopoulos and H. Georgi, Nucl. Phys. B193 (1981) 150; N. Sakai, Z. Phys. C11 (1981) 153; R. K. Kaul, Phys. Lett. 109B (1982) 19.

[3] H. Georgi, Phys. Lett. 108B (1982) 283; A. Masiero, D. V. Nanopoulos, K. Tamvakis and T. Yanagida, Phys. Lett. 115B (1982) 380; B. Grinstein, Nucl. Phys. B206 (1982) 387.

[4] E. Witten, Phys. Lett. 105B (1981) 267; D. V. Nanopoulos and K. Tamvakis, Phys. Lett. 113B (1982) 151.

[5] J. Polchinski and L. Susskind, Phys. Rev. D26 (1982) 3661.

[6] M. Dine, lectures at the Johns Hopkins workshop in

Florence, 1982.

[7] H. P. Nilles, M. Srednicki, and D. Wyler, Phys. Lett. 124B (1982) 337.

[8] A. B. Lahanas, Phys. Lett. 124B (1982) 341.

[9] S. Dimopoulos and H. Georgi, Phys. Lett. 117B (1982) 287; K. Tabata, I. Umemura and K. Yamamoto, Prog. Theor. Phys. 71 (1984) 615.

[10] A. Sen, 'A locally supersymmetric SU(6) grand unified theory without fine tuning and strong CP problems', Fermilab report No. Fermilab-Pub-83/106-Thy (Revised version).

[11] L. Alvarez-Gaume, J. Polchinski and M. B. Wise, Nucl. Phys. B221 (1983) 495; L. E. Ibanez, Nucl. Phys. B218 (1983) 514.

[12] J. E. Kim and H. P. Nilles, 'The μ-problem and the strong CP problem', Universite de Geneve report No. UGVA-DPT 1983/10-410.