LOCALLY SUPERSYMMETRIC SU(5) GRAND UNIFICATION

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We consider the coupling of the SU(5) GUT to $N = 1$ supergravity. A general strategy to obtain a reasonable low-energy phenomenology is discussed. Very simple superpotentials naturally lead to large ($\sim m_\mathrm{p}$) vacuum expectation values for the Higgses breaking SU(5) but no large expectation values for the scalar quarks and leptons. Various schemes for obtaining naturally massless Higgs doublets are discussed. Local supersymmetry is broken at a scale $\sim 10^{11}$ GeV and the weak interaction scale is generated from the soft terms breaking global supersymmetry which are left after the breakdown of supergravity. The 24-Higgs scalar potential is very flat and may have cosmological relevance.

Supersymmetric GUTs [1–3] seem to provide the only known viable solution to the gauge hierarchy problem. Recently it has been argued [4] that the global supersymmetry in these models may be broken at a large scale $\gtrsim 10^{11}$ GeV. If the SUSY-breaking scale is so large, it is in principle not clear whether one can still neglect the influence of gravity on the supersymmetric GUTs. In this context, Weinberg [5] has recently pointed out that the coupling of $N = 1$ supergravity to the SU(5) scalar potential may split the degeneracy of the vacuum energy which is typical of SUSY GUTs. More generally, Cremmer et al. [6] have considered the coupling of $N = 1$ supergravity to Yang–Mills theories including also an arbitrary number of chiral multiplets. They obtained (for the case of vanishing cosmological constant) a tree-level mass formula relating the masses of the particles in the theory to the gravitino mass. Ellis and Nanopoulos [7] have remarked that in fact the mass formula implies a lower bound on the masses of the scalars in the theory $m_\mathrm{s} \gtrsim m_{3/2}$, where $m_{3/2}$ refers to the gravitino mass. Since the Higgs doublet mass should not exceed M_W, one finds that $N = 1$ supergravity has to be broken at most at $\lesssim (M_\mathrm{W} m_\mathrm{p})^{1/2} \sim 10^{10-11}$ GeV. This is certainly a strong constraint on models with large supersymmetry breaking [4].

There are more indications that one cannot ignore gravity in grand unified models. It has recently been shown [8] that radiatively generated quark and lepton masses in SUSY GUTs are negligibly small. This implies that in order to explain the failure of the relation $m_\mathrm{e} = m_\mathrm{d}$ in SU(5) one cannot make use of radiative corrections. Gravity-induced interactions of the type proposed by Ellis and Gaillard [9] are the natural source of the masses of (at least) the first generation of fermions.

Since gravity seems to play an important rôle in various aspects of grand unification, it seems logical to combine both ingredients. In this paper we consider the coupling of the SU(5) model to $N = 1$ supergravity [‡1]. We will show how one can obtain a phenomenologically attractive model with very simple superpotentials for the chiral superfields.

Let us consider the coupling of chiral superfields to $N = 1$ supergravity. It is well known that the scalar potential which, in the global case is

$$V_\mathrm{G} = \sum_i |W_{,i}|^2 + \sum_a \tfrac{1}{2} |D^a|^2 , \qquad (1)$$

is modified [11,12,5,6] when supersymmetry is local and takes the form (in Planck units):

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‡1 For a review on supergravity, see ref. [10].

$$V = \exp\left(\sum_i |Z_i|^2\right)\left(\sum_i |W_{,i} + WZ_i^*|^2 - 3|W|^2\right)$$

$$+ \sum_a \tfrac{1}{2}|D^a|^2 , \tag{2}$$

where W is the $\theta = 0$ component of the superpotential, D^a are the usual gauge auxiliary fields and Z_i denote the scalar fields involved. Here we consider for the moment the "minimal" coupling to $N = 1$ supergravity with "canonical" kinetic terms.

To get the flavour of what is going on, let us consider an extremely simple coupling of $N = 1$ supergravity to SU(5). Take as our chiral superfields just an adjoint ϕ and the usual fermion families $3(\bar{\mathbf{5}} + \mathbf{10})$. Take also a constant superpotential $W = \Delta$. This will give rise essentially to an "M term" [13] coupling directly to the M auxiliary field of minimal off-shell supergravity. The potential for the scalars is now:

$$V_\Delta = \exp\left(\sum_i |Z_i|^2\right) \Delta^2\left(\sum_i |Z_i|^2 - 3\right)$$

$$+ \sum_a \tfrac{1}{2}|D^a|^2 . \tag{3}$$

Let us consider the minimization of this potential. The D^2 term is positive definite and hence will prefer to be small. This automatically selects a null vacuum expectation value (vev) for the scalar quarks and leptons because they are contained in complex representations and would give rise to a non-vanishing D^2 term [14]. However, the adjoint ϕ can acquire a vev without creating a D^2 term. This is a nice feature of all the SUSY models. Concerning the first term in eq. (3), one could naively think that it is minimized by setting $\Sigma_i |Z_i|^2 = 0$. However, the exponential factor in eq. (3) ensures that in fact the minimum is obtained for

$$\sum_i |Z_i|^2 = 2 \to \phi^2 = 2 , \tag{4}$$

which means that ϕ acquires a vev of order m_p. This result is remarkable, given the extreme simplicity of the superpotential. Notice, however that there is a negative cosmological constant in this "order-zero" model. One needs Weinberg–Salam doublets to break SU(2) × U(1), etc. We now construct realistic models in which the Weinberg–Salam (WS) doublets remain light. In order to do this, we recall that, according to Ellis and Nanopoulos, the scalars get masses $\gtrsim m_{3/2} \sim (|W|/m_\mathrm{p}^2)$ once local-SUSY is broken. This means that, at the minimum, $|W| \lesssim M_\mathrm{W} m_\mathrm{p}^2 \ll m_\mathrm{p}^3$, i.e., *our superpotential must be such that* $|W| \sim M_\mathrm{W}$ *in Planck units.* So it cannot have terms of the form λ tr ϕ^3 or M tr ϕ^2 which imply $|W| \sim M_\mathrm{X}^3$. On the other hand, we do not need them since, as we have shown above, ϕ acquires a vev $\neq 0$ without the need of extra couplings. However, we know that in usual SUSY guts there is always a term in the superpotential W which, at the desired minimum, is at most M_W^3. This is the term which gives a large mass to the colour triplets in the $\mathbf{5} + \bar{\mathbf{5}}$ of Higgses:

$$W_\mathrm{A} = \lambda \mathrm{H}_1 \phi \mathrm{H}_2 + m\mathrm{H}_1\mathrm{H}_2 . \tag{5}$$

The mass parameter m is tuned to yield ~massless WS doublets. A possible alternative, as suggested in ref. [2], is to use instead an SU(5) singlet χ:

$$W_\mathrm{B} = \lambda \mathrm{H}_1 \phi \mathrm{H}_2 + \beta\chi\mathrm{H}_1\mathrm{H}_2 , \tag{6}$$

which, after minimization [2], implies a massless doublet at the tree level. Either of the superpotentials (5) or (6) can be present in W without giving a large mass to the scalars since, at the desired minimum they are of order M_W. For the sake of definiteness, let us consider the superpotential $W = \Delta + W_\mathrm{B}$, with $\Delta \sim M_\mathrm{W} m_\mathrm{p}^2$, as desired in order not to give too large a mass to the scalars once local SUSY is broken. The scalar potential now has the form:

$$V = \exp\left(\sum_i |Z_i|^2\right)\left[\sum_i |W_{,i}|^2 + \sum_i (W^* Z_i W_{,i} + \mathrm{h.c.})\right.$$

$$\left. + |W|^2\left(\sum_i |Z_i|^2 - 3\right)\right] + \sum_a \tfrac{1}{2}|D^a|^2 , \tag{7a}$$

where

$$\sum_i |W_{,i}|^2 = |(\lambda\phi_j^i + \beta\chi\delta_j^i)\,\mathrm{H}_2^j|^2 + |\mathrm{H}_1^j(\lambda\phi_j^i + \beta\chi\delta_j^i)|^2$$

$$+ |\lambda \mathrm{H}_{1_i}\mathrm{H}_2^j|^2 + |\beta\mathrm{H}_1\mathrm{H}_2|^2 , \tag{7b}$$

$$\sum_i (W^* Z_i W_{,i} + \mathrm{h.c.})$$

$$= 3W^*(\lambda\mathrm{H}_1\phi\mathrm{H}_2 + \beta\chi\mathrm{H}_1\mathrm{H}_2) + \mathrm{h.c.} , \tag{7c}$$

the term

$$\exp\left(\sum_i |Z_i|^2\right)\left(\sum_i |Z_i|^2 - 3\right), \tag{7d}$$

and Z_i denotes any scalar in the model. We do not include the usual Yukawa couplings to quarks and leptons since, as we remarked above, the D^2 term gives zero vev to their scalar partners and hence its conclusion would not alter the minimization [14]. It is easy to convince oneself that the uncontrolled growth of $|W|$, i.e., $|W| \gg \Delta$, is energetically disfavoured, hence at the minimum $|W| \sim \Delta$. Looking at the potential (7a) we see that the term from (7d) is just like the potential in (3) and hence tends to give a $(\mathrm{vev})^2 \sim m_\mathrm{p}^2$ to the sum $\Sigma |Z_i|^2$. On the other hand, the term (7b) (which is the usual global SUSY potential but for the exponential factor) yields $\mathrm{H}_1 \sim \mathrm{H}_2 \sim 0$, ϕ and χ undetermined. However, the term (7c) is not positive definite and combined with the term (7b) may admit small vev's $\mathrm{H}_1 \sim \mathrm{H}_2 \sim M_\mathrm{W}$ and $\phi\lambda_5^5 + \beta\chi \sim M_\mathrm{W}$. The latter condition keeps the WS doublets light but the absolute magnitude of χ^2 (or ϕ^2) remains undetermined. However, the term (7d) gives a vev $\sim m_\mathrm{p}$ to χ and ϕ since they are the only fields in the sum $\Sigma |Z_i|^2$ which can acquire a vev $\neq 0$ without paying any energetical penalty. Thus the physical outcome of the minimization of the scalar potential is:

$$\mathrm{H}_1 \sim \mathrm{H}_2 \lesssim M_\mathrm{W},$$

$$\lambda\phi_5^5 + \beta\chi \lesssim M_\mathrm{W} \quad \text{(light W–S doublets)},$$

$$|\phi|^2 + |\chi|^2 \simeq 2m_\mathrm{p}^2 + \mathrm{O}(M_\mathrm{W}^2). \tag{8}$$

We do not think it is worth doing the exact [including $\mathrm{O}(M_\mathrm{W}^4)$ terms] minimization of the potential since, as we shall argue below, the radiative corrections may substantially affect the low-energy effective theory. Notice that there are two degenerate ground states, $\mathrm{SU}(4) \times \mathrm{U}(1)$ and $\mathrm{SU}(3) \times \mathrm{SU}(2) \times \mathrm{U}(1)$, so that the degeneracy problem still remains in this simple model. However, it may be a question to answer when the early times of the evolution of the Universe are better understood.

One may try to cancel the cosmological constant (c.c.) by adding an SU(5) singlet Y with the potential $M^2\mathrm{Y}$, $M^2 \sim \Delta/m_\mathrm{p}$. The complete superpotential is then

$$W = \Delta + M^2\mathrm{Y} + \lambda\mathrm{H}_1\phi\mathrm{H}_2 + \beta\chi\mathrm{H}_1\mathrm{H}_2 \tag{9}$$

and the scalar potential is changed by the addition of:

$$\exp\left(\sum_i |Z_i|^2\right)(M^4 + |W|^2|\mathrm{Y}|^2 - W^* 2M^2\mathrm{Y} + \mathrm{h.c.}). \tag{10}$$

The field Y gets a vev $\sim m_\mathrm{p}$ and M^2 can be tuned to cancel the c.c. However, it is easy to see that if one cancels the c.c., ϕ and χ do not get vev's $\sim m_\mathrm{p}$ but at most $\sim\Delta$. Hence *we will not tune M^2 to cancel the c.c.* in a supersymmetric invariant way. Till the c.c. problem is clarified [17] we will assume that the c.c. is cancelled by adding an explicit non-supersymmetric cosmological term [13]. However, the introduction of the singlet Y may still be useful, as we will show below. In the above models, local supersymmetry is broken since $|W| \sim \Delta \neq 0$. The massless gravitino combines with the Goldstino which in the case is a linear combination of the fermionic Y and the fermionic partners of the scalars ϕ and χ with vev $\neq 0$. The resulting massive gravitino has a mass $m_{3/2} \sim \Delta$. Notice that the only large mass term for the scalars in the adjoint ϕ comes from the D^2 term. This means that the scalars transforming like a $(8,1,0) + (1,3,0)$ of $\mathrm{SU}(3) \times \mathrm{SU}(2) \times \mathrm{U}(1)$ acquire masses of at most $\sim\Delta$, and that the same happens with their fermionic partners. This modifies the values given by the renormalization group for $\sin^2\theta_\mathrm{W}$ and M_X. One obtains $\sin^2\theta_\mathrm{W} \simeq 0.24$ and $M_\mathrm{X} \sim 1 \times 10^{18}$ GeV. This value for M_X is approximately consistent with the result obtained above that ϕ acquires a vev $\sim m_\mathrm{p}$.

The singlet introduced in (9) may be used to give a direct mass to the gauginos [6] if the coupling of $N = 1$ supergravity to the Yang–Mills strength is not minimal, i.e.,

$$\int \mathrm{d}x^4 \mathrm{d}\theta^4 E(R^{-1} f_{\alpha\beta} W^\alpha W^\beta); \quad f_{\alpha\beta} = (1 + \mathrm{Y})\delta_{\alpha\beta}, \tag{11}$$

where W^α is the chiral Yang–Mills strength, E is the determinant of the vielbein and R is the scalar curvature. From eq. (11) one can see that there is a term in the lagrangian of the form [6]

$$f^{\alpha\beta}_{,\mathrm{Y}} F_\mathrm{Y}(\lambda_\alpha\lambda_\beta) + \mathrm{h.c.} \tag{12}$$

where $F_\mathrm{Y} = M^2$. This gives gaugino masses $\sim M^2/m_\mathrm{p} \sim \Delta/m_\mathrm{p}^2$.

It is difficult to say precisely what are the effective interactions at low energies. This is because radiative corrections may play an important role in generating the low-energy potential. This is what happens for example in the recently proposed SUSY GUTs with large supersymmetry breaking [4]. If, for example, the direct gaugino masses are large enough, they may induce radiatively [2,4] scalar masses which may be even bigger than the tree level ones coming from the coupling to $N = 1$ supergravity in the potential (7a). Let us see, however, what terms seem to be present at low energies apart from the globally supersymmetric terms in (7b).

(i) There are either direct [as in eq. (12)] or radiative (as suggested in ref. [15]) gaugino masses. Notice that, due to renormalization effects, the gluino mass will become much bigger than the other gaugino masses at low energies.

(ii) There are tree level scalar masses coming from the potential (7a), as argued by Ellis and Nanopoulos [7].

(iii) There are mass terms of the form $(\Delta/m_p)H_1H_2 +$ h.c. for the colour triplet of Higgses, coming from the term $W^*(\lambda H_1\phi H_2 + \beta H_1H_2\chi)$ + h.c. in the potential. This mass $(\Delta/m_p) \sim M_w m_p$ also generates radiative masses $\sim M_w$ for the rest of the scalars, as is well known from radiative SUSY GUTs [3,4].

(iv) Soft trilinear terms of the form

$$(\Delta/m_p^2)\,\lambda H_1\phi_3H_2 + \text{h.c.}\,, \quad (\Delta/m_p^2)\,h_t\phi_{t_L}\phi_{t_R}H_1 + \text{h.c.}\,,$$

where H_1 and H_2 are the WS doublets and ϕ_3 is the SU(2) triplet component of the superfield ϕ which, as we remarked above, is light. ϕ_{t_L} and ϕ_{t_R} are the scalar partners of the top quark, for example. Notice that all the SUSY breaking terms which are remnant at low energies are soft and hence the scalar masses are protected from quadratic divergences.

It is quite plausible that from all these terms one may obtain a sensible low-energy phenomenology and, in particular, that SU(2) × U(1) is broken to $U(1)_Q$. We know that, apart from positive (mass)^2 radiative terms which usually are present [2–4], the quarks and leptons are protected from acquiring vev's by the D^2 term in the potential. On the other hand, as shown in eq. (8), the WS doublets are not protected from acquiring a vev $\lesssim M_w$. Also, the vev's of H_1 and H_2 can arrange themselves to give a null D^2 term. Furthermore, as pointed out in ref. [2], if the Yukawa coupling of the top quark is large enough, a radiative negative (mass)^2 for the Higgs doublet may be generated (but not for the scalar leptons, for example). The fact that soft SUSY breaking terms of the type considered here may lead radiatively to the desired low-energy symmetry has also been recently studied in ref. [16].

The introduction of the singlet χ in the superpotential W_B in order to obtain massless WS doublets certainly works at the tree level, but it is not clear that radiative corrections preserve the nearly masslessness of the doublets. In order to see if this mechanism works, one would have to calculate the radiative corrections to the potential in (7a), which is certainly a difficult task. If it does not work one would have to take the superpotential in eq. (5) and fine-tune the mass parameter m. Another mechanism to obtain naturally massless WS doublets has recently been suggested by the authors of ref. [18]. The part of their superpotential relevant to us is:

$$W_c = \lambda_1 H_1\varphi_1\phi + \lambda_2 H_2\varphi_2\phi + m\varphi_1\varphi_2\,, \tag{13}$$

where

$$H_1 \sim \mathbf{5}\,, \quad \varphi_1 \sim \mathbf{50}\,; \quad H_2 \sim \overline{\mathbf{5}}\,, \quad \varphi_2 \sim \overline{\mathbf{50}}\,; \quad \phi \sim \mathbf{75}\,,$$

and we will take $m \sim m_p$. The first two couplings give masses to the colour triplets in H_1 and H_2 once ϕ acquires a vev, but not to the WS doublets. The huge mass term for $\phi_1\phi_2$ forbids ϕ_1 and ϕ_2 to acquire a vev and hence, at the minimum, $W_c \sim 0$, which is the requisite we need in order not to have too large scalar masses. One can then consider a superpotential of the form

$$W = \Delta + M^2 Y + W_c\,,$$

and follow the same steps as we followed with the superpotentials W_A and W_B. However, one can see that in this case there is no term forbidding a large vev for H_1 and H_2 as well as for ϕ, so that there are many possible degenerate ground states at this level.

Let us remark that the strategy sketched in the above models is, *in principle, general* and that one can apply it to other superpotentials fulfilling certain requirements. The important point is that the value of the superpotential *W at the minimum* must verify $|W| \lesssim M_w m_p^2$. Then the superpotential, if possible, must be such that the WS doublets are forbidden from acquiring a large vev by the $\Sigma_i |W_{,i}|^2$ terms, but the value of the Higgses breaking SU(5) is not restricted by those terms. Then the term in (7d) does the rest of the job and gives a large vev to the ϕ Higgses. The weak interaction scale is generated through the geometric scale [4] Δ/m_p^2.

Another interesting point of these scalar potentials generated by $N = 1$ supergravity is that they are very flat. The barrier between different (in principle degenerate) ground states is of order $V \sim \Delta^2 \sim |M_W m_P|^2 \sim (10^{11})^4$ GeV^4, whereas the vev involved is of order $\sim m_P$. This type of scalar potential may be interesting for building an "inflationary universe" cosmological scenario [19].

Summing up the results presented in this letter, we have shown a general strategy to build specific SU(5) models coupled to $N = 1$ supergravity. Very simple superpotentials lead to the breaking of SU(5) close to the Planck mass. Local supersymmetry is broken at a scale $(M_W m_P)^{1/2} \sim 10^{11}$ GeV but there is an effective residual global supersymmetry only broken by soft terms which protects the scalars from quadratically divergent masses. The weak scale is generated both by tree level and radiative terms at a mass $\sim\Delta/m_P^2$, which is also the order of magnitude of the gravitino mass.

In this type of models, supersymmetry breaking certainly appears in a much more elegant way than in the models using O'Raifeartaigh potentials [3,4] or Fayet–Iliopoulos U(1) factors [20]. Moreover, local supersymmetric models include the effects of gravity which, as we remarked above, probably cannot be neglected. Hopefully, the SU(5) × local $N = 1$ supersymmetry is a remnant symmetry from a dynamical realization of $N = 8$ supergravity in which one of the supersymmetric generators is only broken at a scale $\sim 10^{11}$ GeV. Still, apart from the question of the origin of such an intermediate scale [‡2], there is the outstanding problem of naturally obtaining ~massless doublets. Maybe the remnant local supersymmetry does not commute with SU(5). Another question is the naturalness problem of what terms are present and what terms are not present in the superpotential W. Anyway, it seems worth exploring different possible W's which may lead to the scenario explained above.

I acknowledge conversations with John Ellis.

Note added. After this work was finished I became aware of the work in ref. [22]. These authors also consider the coupling of SU(5) to $N = 1$ supergravity and cancel the cosmological constant. However, they minimize the scalar potential expanding around $m_P \to \infty$, keeping only terms $\sim O(1/m_P)$. They also consider a grand unification mass different from the Planck mass.

[‡2] A scenario in which supergravity is broken dynamically has recently been suggested in ref. [21].

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