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HIERARCHICAL SUPERGRAVITY-INDUCED $SU(2) \times U(1)$
BREAKING IN $SU(5)$ GUTS

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ABSTRACT

In contrast with naïve extensions of the standard supersymmetric $SU(5)$ model to local supersymmetry, we show that different choices of representations in the heavy sector of the theory involving the 75, 50 and $\overline{50}$ representations of $SU(5)$ produce a low-energy effective $SU(3) \times SU(2) \times U(1)$ theory, where the $SU(2) \times U(1)$ breaking is entirely induced by the gravitino mass scale. We discuss how in these models one can avoid that the tree-level hierarchy $M_W \ll M_X$ become unstable in perturbation theory.

It has recently been recognized¹⁾ that the hierarchy problem of Grand Unified Theories (GUTs) can be solved by supersymmetric Yang-Mills theories²⁾. In fact, non-renormalization theorems³⁾ avoid certain types of radiative corrections and allow a natural separation of the weak and grand unification scales which would be problematic to understand on the basis of non-supersymmetric GUTs.

However, the simplest extension of the Georgi-Glashow SU(5) model faces other theoretical problems -- the triplet-doublet problem and the origin of supersymmetry breaking. An elegant way to solve the former is to have a group-theoretical reason for the Weinberg-Salam doublets being massless. This can be achieved by using a different set of representations in the heavy sector of the theory⁴⁾. In fact, the adjoint representation of Higgs can be replaced by the 75 dimensional real representation, and additional 50 and $\bar{50}$ dimensional representations must be introduced to allow the Higgs triplets of the 5 and $\bar{5}$ representations to have a large mass, while the Higgs doublets remain massless⁴⁾.

A suitable choice of the superpotential for the heavy sector of the theory, which has a supersymmetric solution with the desired properties of gauge-breaking of SU(5) into SU(3) $\times$ SU(2) $\times$ U(1) is, for example,

$$\lambda_1 \Sigma^3 + M \Sigma^2 + \lambda_2 \theta \Sigma \bar{\theta} + \lambda_3 \theta \Sigma H + \lambda_4 \bar{\theta} \Sigma \bar{H} + M' \theta \bar{\theta} \quad (1)$$

Σ , θ , $\bar{\theta}$, H , $\bar{H}$ denote 75, 50, $\bar{50}$, 5, and $\bar{5}$ representations of SU(5), respectively. We usually denote the chiral multiplets and their first scalar component by the same symbol to simplify notation. Also, obvious group contractions must be understood.

A supersymmetric solution of (1) exists which breaks SU(5) down to SU(3) $\times$ SU(2) $\times$ U(1) with

$$\langle \Sigma \rangle = O(M), \quad \langle \theta \rangle = \langle \bar{\theta} \rangle = \langle H \rangle = \langle \bar{H} \rangle = 0 \quad (2)$$

M being the GUT scale. We also assume that all mass scales in (1) are $O(M)$.

Of course, there are other supersymmetric minima⁵⁾ than (2), but the one given by (2) can be selected in a GUT coupled to $N = 1$ supergravity by demanding the vanishing of the cosmological constant to the order of M^6/M_P^2 where M_P is the Planck mass⁶⁾. We also note that in (1) we could replace the $\theta \bar{\theta}$ mass term with a coupling to a singlet X with further interactions

$$\lambda_5 X \theta \bar{\theta} + \lambda_6 X^3 + \lambda_7 X \Sigma^2$$

but this is not relevant for what follows.

The second problem to be solved is to obtain the $SU(3)_c \times SU(2) \times U(1)$ breaking down to $SU(3)_c \times U_{e.m.}(1)$ and the supersymmetry breaking. Efforts with global supersymmetry, spontaneously broken at a scale $M_S \simeq (10^2-10^3)$ GeV, were proved unsuccessful. Realistic models, seemed to need an intermediate supersymmetry (SUSY) breaking scale $M_S = O(\sqrt{M_W M_P})$. However, for these models the gravitino mass is⁷⁾ $m_{3/2} = O(M_W)$ and supergravity corrections can no longer be neglected^{7,8)}. Using the supergravity induced supersymmetry breaking discussed in Ref. (8) one can show that the correct tree-level pattern of $SU(2) \times U(1)$ breaking, as well as the correct mass spectrum for low-energy particles, can be achieved by coupling the $SU(5)$ GUT to $N = 1$ supergravity with the following additional Higgs superpotential terms⁹⁾

$$\alpha Y H \bar{H} + \frac{\beta}{3} y^3 \quad (3)$$

where Y is an $SU(3)$ singlet.

An additional chiral $SU(5)$ singlet z is also needed to cancel the cosmological constant with $\langle z \rangle = O(M_P)$, z being responsible for the spontaneous supersymmetry breaking¹⁰⁾. The fermion partner of the scalars z is the goldstino and it is eaten by the massive gravitino through the super-Higgs effect.

The full $SU(5)$ theory is defined by adding to (1) and (3) and to the z superpotential the Yukawa coupling of normal matter, i.e. additional superpotential terms for the 10 and $\bar{5}$ representations of $SU(5)$:

$$\alpha_{ij} 10_i 10_j H + \beta_{ij} 10_i \bar{5}_j \bar{H} + \gamma_{ij} 10_i 10_j \bar{\Theta} \quad (4)$$

To solve the hierarchy problem, the coupling of the singlet z to matter must be such that the non-vanishing gravitino mass $m_{3/2}$ sets up the scale of supersymmetry $M_S = O(\sqrt{m_{3/2} M_P})$ as well as the scale of $SU(2) \times U(1)$ breaking $O(m_{3/2})$.

At the tree level it turns out that the complete scalar potential of the Yang-Mills supergravity coupled system with superpotential given by the sum of (1), (3), and (4) and the z -superpotential to be defined below, has $SU(3) \times U_{e.m.}(1)$ invariant minima such that

$$\langle H_2 \rangle, \langle \bar{H}_2 \rangle, \langle y \rangle = O(m_{3/2}), \quad \langle z \rangle = O(M_P) \quad (5)$$

all the other VEVs but (2) being zero.

We note that, in contrast with other proposals the three-level hierarchy is entirely induced by the gravitino mass. Previously proposed models for locally supersymmetric GUTs did not solve the hierarchy problem. In Ref. (11) $SU(5)$ GUTs

were proposed. However, the form of the scalar potential, as well as its minimization, was incorrect^{*)}. In Ref. (11) the weak interaction scale was explicitly introduced by assuming a definite unnatural relation between the couplings M , M' of the adjoint Higgs Δ (24-dimensional representation) and the 5 and $\bar{5}$ $SU(5)$ representations:

$$\frac{M}{2} \Delta^2 + M' H \bar{H} \quad (6)$$

If one makes this *ad hoc* assumption, then one merely reproduces the low-energy electro-weak theory proposed in Ref. (13), where the condition $M = M'$ is equivalent to a relation between m and $m_{3/2}$ in the superpotential of Ref. (13).

Other proposed models suffer from the problem of a non-vanishing cosmological constant¹⁴⁾.

In addition, as has recently been pointed out by Nilles, Srednicki and Wyler¹⁵⁾, and Lahanas¹⁶⁾, all these models do not maintain the hierarchy of scales in perturbation theory. This is due to the presence of light singlets in these models, which are coupled to the heavy sector of the theory¹⁷⁾. As we will see below, and as was already shown in Ref. (13), the supergravity couplings, in the limit $M_P \rightarrow \infty$ but with $m_{3/2}$ fixed, introduce soft-breaking terms in the scalar potential which induce radiative effects for light singlets, similar to those discussed by Polshinski and Susskind¹⁷⁾ in the context of spontaneously broken globally supersymmetric models with a geometric hierarchy.

For example, the model discussed in Refs. (11) and (12) has an unstable hierarchy because the Higgs doublets receive a huge one-loop radiative mass $O(m_{3/2} M_X)$. Possible viable alternatives have been proposed when the $SU(2) \times U(1)$ breaking is driven by radiative corrections in an effective low-energy theory with or without non-vanishing masses for scalar quarks and leptons at the tree-level^{18,19)}.

Remarkably enough we will show that the use of the 75 and 50, $\bar{50}$ representations of $SU(5)$ provide the first examples of GUTs in which this problem can be avoided. This can be achieved by a technically natural fine tuning in the superpotential which describes the heavy sector of the theory. We demonstrate that the minimal model, which has a stable one-loop hierarchy, must contain at least 24, 75, 50, and $\bar{50}$ representations in the heavy sector of the theory. An alternative possibility, which does not require any fine-tuning at the one-loop level, needs a 75 representation and a doubling of the 50, $\bar{50}$, 5, and $\bar{5}$ representations, respectively.

*) Revised versions of both papers, quoted as Ref. (11), have recently come out¹²⁾.

In order to describe the soft-breaking terms and the solution to the instability of the hierarchy, we consider in some detail the coupling of the goldstino-multiplet z to supergravity and to the matter fields generally denoted by y .

Let us introduce the following linear function of the singlet z ¹⁰⁾

$$g(z) = A(z + \delta) \quad (7)$$

In Ref. (13) it was shown that when $g(z)$ is coupled to matter one can still solve the extremum condition for z , at zero potential

$$V = 0 \quad \partial V / \partial z = 0 \quad (8)$$

provided the matter superpotential $h(y_i)$ does not contain parameters $O(M_P)$. In the low-energy limit, neglecting terms $O(1/M_P)$ in the final Lagrangian the z -field decouples from matter and the resulting potential is ^{13,9)}

$$V(y, y^*) = |\tilde{h}^i|^2 + (A-3)m_{3/2}(h + h^*) + \frac{1}{2}(D^\alpha)^2 \quad (9)$$

where $h(y)$ is the matter superpotential, and

$$\tilde{h}^i = \frac{\partial h}{\partial y_i} + m_{3/2} y^{*i} \quad D^\alpha = e_\alpha y^{*i} T^{\alpha j}_i y_j \quad (10)$$

e_α is the gauge-coupling constant and $T^{\alpha j}_i$ are the group generators in the representation of the y 's. The parameter A depends on the particular way the z -singlet couples to the matter multiplets y_i ⁹⁾.

The parametrization (9) covers, for example, the cases ²⁰⁾

$$g(z, y_i) = g(z) \left[1 + \frac{g(z)^{q-1}}{\langle g(z) \rangle^{q-1}} h(y_i) \right] \quad (11)$$

where $q = 1, 2, \dots$. For $q = 1$ we recover the original form of Ref. (13) with $A = 3 - \sqrt{3}$, while for $q = 2$ we recover the multiplicative structure proposed in Ref. (21) with $A = 3$. For $q > 2$, $A - 3 = \sqrt{3}(q - 2) > 0$, and we fulfil the requirement of Ref. (9) for the breaking of $SU(2) \times U(1)$ purely induced by the gravitino mass $m_{3/2}$. We note that more general expressions than (11) and (9) can also be obtained ²⁰⁾.

It is crucial to notice for the discussion which follows that the effective theory described by the potential given by Eq. (9) can be rewritten in a superfield formalism by adding to the usually supersymmetric Lagrangian the following soft terms²²⁾

$$[\phi^{*a}(\eta+\eta^*)\phi_a]_D \quad (A-3)[\eta h(\phi)]_F + h.c. \quad (12)$$

where ϕ_a are matter chiral superfields whose first component is y_a and η is a spurion chiral superfield whose only non-vanishing component is the last one, $\eta_F = m_{3/2}$. After performing the θ -integrations and eliminating the auxiliary non-propagating fields F_a :

$$F_a = h_a^*(y^*) + m_{3/2} y_a \quad (13)$$

one gets back the potential described by Eq. (9). We just note that the situation described in Ref. (21) corresponds to the absence of the second invariant in Eq. (12).

We now turn to a detailed discussion of how the tree-level hierarchy could be destroyed in locally supersymmetric GUTs by radiative corrections as well as the general strategy we propose to avoid such an instability. In the models proposed so far, the $SU(2) \times U(1)$ breaking at low energy is provided by the coupling of the H and $\bar{H}$ superfields to a singlet superfield (Y). At the tree-level the VEVs of the Y -components come from the soft SUSY breaking terms. One then has to make sure that loop corrections involving large VEV insertions of $O(M_X)$ cannot generate VEVs of $O(m_{3/2} M_X)$ for the Y -components which would give rise to masses of $O(m_{3/2} M_X)$ for the $H_2, \bar{H}_2$ doublets^{15,16}.

The simplest class of one-loop super diagrams that, in our example, will potentially contribute to the generation of a $H_2 \bar{H}_2$ mass term proportional to $m_{3/2} M_X$ are indicated in Fig. 1. The soft breaking terms in Eq. (12) are taken into account by inserting spurions $(\eta+\eta^*)$ in each propagator and $(A-3)\eta$ or its conjugate in the corresponding vertices. An additional even number of mass insertions will give further diagrams in this class. The mass insertions appearing in the loop are those resulting from the shift of the fields by their respective VEVs. For the $(3,1,-1/3)$ components contributing to the tadpole in Fig. 1 the resulting mass terms are as follows

$$M_{\theta\theta} \theta_3 \bar{\theta}_3 + M_{\theta H} \theta_3 H_3 + M_{\theta H} \bar{\theta}_3 \bar{H}_3 \quad (14)$$

Since $M_{\theta H}$ and $M_{\bar{\theta} \bar{H}}$ are essential in order to avoid light $\Delta B \neq 0$ fields the only possible way to cancel the 1-loop diagrams in Fig. 1 is to put $M_{\theta \bar{\theta}} = 0$. In this way there is no bilinear term in the Lagrangian that connects the (θ_3, H_3) sector to the $(\bar{\theta}_3, \bar{H}_3)$ sector. This can be achieved if there are at least two different representations of superfields with non-vanishing VEVs contributing to $M_{\theta \bar{\theta}}$. Then $M_{\theta \bar{\theta}} = 0$ is obtained by fine-tuning their corresponding $\theta \bar{\theta}$ couplings. The lowest representations coupling 50 to a $\bar{50}$ behave like a singlet, a 24-plet, or a 75-plet^{*}). If one would choose a singlet (or, equivalently, an invariant mass term) together with either a 24 or a 75, then if $M_{\theta \bar{\theta}}$ is fine-tuned to zero, other components of the 50 and $\bar{50}$ supermultiplets will remain massless as well. This should be avoided^{**}). On the contrary, the choice of 24 and 75 representations allows for a cancellation of just the $\theta_3 \bar{\theta}_3$ mass term, and therefore should be preferred. It goes without saying that the above mechanism can be applied to any model employing light singlets for the $SU(2) \times U(1)$ breaking.

Once these one-loop tadpoles have been avoided through the fine-tuning described above, one still has to take care of higher loop corrections. Indeed, the diagrams illustrated in Fig. 2 could spoil the hierarchy[†]). We can contemplate different ways out of this problem:

- i) The coupling constants may be chosen to be small enough, typically of $O(10^{-3})$, so that the final contribution of these two-loop diagrams becomes harmless. In this case, one has to check if this value is consistent with the present limit on proton decay, since the masses of the scalar triplets that are proportional to these coupling constants fix the strength of the dimensional five operators responsible for proton decay in SUSY GUTs. One then finds values around 10^{13} - 10^{14} GeV for these masses which are not in contradiction with the experimental bounds²³).
- ii) One can avoid any intermediate line in Fig. 2 by demanding that $H\theta$ and $\bar{H}\bar{\theta}$ couple to different superfields than those coupled to $\theta\bar{\theta}$, and then by requiring that the two sets of supermultiplets do not communicate through

*) The 50-representation decomposes under $SU(3) \times SU(2) \times U(1)$ into: $(\bar{6}, 1, -4/3) + (8, 2, -1/2) + (1, 1, 2) + (3, 2, 7/6) + (6, 3, 1/3) + (\bar{3}, 1, 1/3)$. The $SU(3) \times SU(2) \times U(1)$ invariant masses coming from the $(50 \times 75 \times 50)$ coupling are in the ratios, 1:0:3:1:-1:1, respectively. Those arising from a $(50 \times 24 \times \bar{50})$ coupling are in the ratios, -8:-3:12:7:2:2, respectively.

**) These components would be given masses of $O(\lambda \sqrt{m_{3/2}} M)$ by radiative corrections, but they could introduce dangerous $H_2 \bar{H}_2$ mass terms at the two-loop level discussed below.

†) We acknowledge an enlightening discussion with M. Sredincki on this problem.

mass terms. This is achieved by forbidding any term in the Lagrangian which mixes supermultiplets in the two sets. For instance, one could couple $\theta\bar{\theta}$ to a 24 and a 75 (denoted by Δ and Σ , respectively) and then couple θH and $\bar{\theta}\bar{H}$ to a different 75 (denoted by Σ'), while excluding any coupling of Σ' to Δ and Σ . This leads to a superpotential for the heavy sector which, for example, may take the form,

$$\lambda_2 \theta \Sigma \bar{\theta} + \lambda_3 \theta \Sigma H + \lambda_4 \bar{\theta} \Sigma' \bar{H} + \lambda_5 \theta \Delta \bar{\theta} + f(\Sigma, \Sigma', \Delta), \quad (15)$$

with, for example,

$$f(\Sigma, \Sigma', \Delta) = \alpha_1 \Sigma^3 + \alpha_2 \Delta^3 + \alpha_3 \Delta^2 \Sigma + \alpha_4 \Sigma^2 \Delta + \\ + M_\Sigma \Sigma^2 + \alpha_5 \Sigma'^3 + M_{\Sigma'} \Sigma'^2.$$

Notice that the $SU(3) \times SU(2) \times U(1)$ symmetries of the Σ , Δ VEVs and of the Σ' VEVs do not have to coincide at the tree-level minimization. However, radiative corrections could align these solutions by introducing couplings of Σ' to Σ and Δ of $O(m_{3/2}/M_X)$. But one could even enforce the alignment of the VEVs by introducing further supermultiplets. For instance, another 24, denoted by Δ' , could be coupled to the Σ' and to the Σ, Δ superfields, but not to the 50, $\bar{50}$, and 5, $\bar{5}$ ones. The two-loop diagrams would then be avoided through a new fine tuning at the level of these additional couplings.

- iii) One could just fine tune the coupling constants in Fig. 2 so that the total two-loop contributions are made to vanish.

Instabilities from higher loop diagrams are harmless for Yukawa couplings in the usual range (10^{-1} - 10^{-2}).

It is of some interest to discuss the low-energy spectrum of $SU(5)$ GUTs involving the components of the Y , H_2 , $\bar{H}_2$ multiplets and the $SU(3) \times SU(2) \times U(1)$ gauge multiplets. The low-energy spectrum will, in general, depend on the function $g(z, y_1)$ and for the class of functions given by (12) on the parameter A .

As noticed in Ref. (21) a great simplification occurs for $A = 3$. In any case $A \geq 3$ if the $SU(2) \times U(1)$ breaking is purely induced by the gravitino mass $m_{3/2}$. We just observe that if we have another energy scale no restriction on A occurs. This was just the case studied in Ref. (13) and it also contains, as a particular case, the model considered in Refs. (11) and (12). For $A = 3$ and with the low-energy potential given by Eq. (3) we get the following results. A first set of states are given by the charged and neutral vector supermultiplets associated with the $W^\pm$ and Z^0 vector bosons with mass splitting $O(m_{3/2})$. Their scalar

components are the opposite parity partners of the Goldstone bosons with masses (similar formulae also hold in the heavy sector)

$$M_{\pm}^2 = M_W^2 + 4m_{3/2}^2 \quad M_0^2 = M_Z^2 + 4m_{3/2}^2 \quad (16)$$

while their fermionic components, together with the gauginos associated with the $W^{\pm}$, Z bosons, give rise to two charged (Dirac) fermions and two neutral Majorana fermions with masses

$$m_{\pm(1,2)}^2 = M_W^2 + \frac{1}{2} m_{3/2}^2 \left(1 \pm \sqrt{1 + 4M_W^2/m_{3/2}^2} \right) \quad (17)$$

$$m_{0(1,2)}^2 = M_Z^2 + \frac{1}{2} m_{3/2}^2 \left(1 \pm \sqrt{1 + 4M_Z^2/m_{3/2}^2} \right)$$

There are four more neutral scalars with masses

$$M_{0(2,3)}^2 = \frac{1}{2} (1-x) m_{3/2}^2 \left[5-x \pm \sqrt{9-10x+x^2} \right] \quad (18)$$

$$M_{0(4,5)}^2 = \frac{1}{2} m_{3/2}^2 \left[9-2x+x^2 \pm (3+x) \sqrt{9-10x+x^2} \right]$$

and two neutral (Majorana) fermions with masses

$$m_{0(3)}^2 = m_{3/2}^2 \quad m_{0(4)}^2 = m_{3/2}^2 (2-x)^2$$

where $x = \beta/\alpha$. We remark that, in sharp contrast with Ref. (21), the scale of weak interactions is triggered by $m_{3/2}$ without any spontaneously broken residual global symmetry which would unavoidably lead to an unacceptable Goldstone boson.

Note that for $x = 0$ ($\beta = 0$) $M_{0(5)}^2$ vanishes, the corresponding state being the Goldstone boson of the spontaneously broken $U(1)$ symmetry alluded to before.

For the scalar quarks and leptons one has the general formula²¹⁾

$$M_{sq(sl)}^{(\pm)} = |m_{3/2} \pm m_{q(e)}| \quad (20)$$

From Eqs. (17)-(20) one can easily check the general mass formulae of $N = 1$ supergravity. The $W^{\pm}$ and the gravitino masses are related as follows

$$M_W^2 = \frac{g^2}{\alpha^2} (1-x)^2 m_{3/2}^2 \quad (21)$$

and may be rather different, depending on the ratio $x = \beta/\alpha$.

If one assumes the so-called minimal coupling^{a)} of matter to supergravity the gluinos and the photinos are massless at the tree level. They get one-loop radiative masses of $O(\alpha_s m_{3/2})$ and $O(\alpha_{em} m_{3/2})$, respectively, from the splitting of $O(m_{3/2})$ in the masses of the coloured and/or charged heavy particles [the charged particles whose spectrum is given in Eqs. (16) and (17) also contribute to the gaugino masses, while the quarks and lepton multiplets give a smaller contribution for $m_{q,\ell}/m_{3/2} \ll 1$].

It is also possible to get a direct gaugino mass from a non-minimal kinetic term for the gauge sector of the theory^{a)}. In this case the natural order of magnitude for gaugino masses would be $m_{3/2}^{18,21}$.

Although in the previous examples the stability of the hierarchy has been obtained through a (technically natural) fine-tuning, it is also possible to build another, more elegant, mechanism where the fine-tuning is replaced by an approximate symmetry as we will now discuss. This can be done by considering together with the set of $50 + \overline{50} + 5 + \overline{5}$ supermultiplets $(\theta, \overline{\theta}, H, \overline{H})$, an additional identical set, denoted by $(\theta', \overline{\theta}', H', \overline{H}')$. We then replace the superpotential in Eq. (1) by

$$\begin{aligned} & \lambda_1 \Sigma^3 + M \Sigma^2 + \lambda_2 \theta \Sigma \overline{\theta} + \lambda_3 \theta \Sigma H + \lambda_4 \overline{\theta} \Sigma \overline{H} + M' \theta \overline{\theta} \\ & + \lambda'_2 \theta' \Sigma \overline{\theta}' + \lambda'_3 \theta' \Sigma H' + \lambda'_4 \overline{\theta}' \Sigma \overline{H}' + M'' \theta' \overline{\theta}', \end{aligned} \quad (22)$$

which has a $U(1)$ symmetry (besides the usual $B - L$ symmetry). Next, we replace the terms appearing in Eq. (3) by

$$\alpha Y_H \overline{H}' + \alpha' Y_{H'} \overline{H} + \frac{\beta}{3} Y^3, \quad (23)$$

so that the $U(1)$ symmetry is broken by the Y -terms. Then, the one-loop diagrams in Fig. 1 are irrelevant since the only way to link the $(\theta \overline{\theta} H \overline{H})$ and the $(\theta' \overline{\theta}', H' \overline{H}')$ superfields is through a $\langle Y \rangle$ insertion. The diagrams in Fig. 2a are also unimportant for the same reason, but those in Fig. 2b remain since they involve the Y^3 -coupling that breaks the $U(1)$ symmetry. These two-loop diagrams can now be suppressed again either by involving small Yukawa couplings or by fine-tuning the effective $\theta \overline{\theta}$ or $\theta' \overline{\theta}'$ mass term to zero, or by the clever inclusion of new supermultiplets and approximate selection rules (exact symmetries are potentially dangerous since they lead to Goldstone particles when spontaneously broken).

To summarize: in the present paper we have investigated GUTs emerging from a spontaneously broken $N = 1$ local supersymmetry only in the mathematical limit $M_P \rightarrow \infty$ with M_X and $m_{3/2}$ fixed. Only in this limit is the theory renormalizable and the "règles du jeu" well defined. This is equivalent to having neglected all interactions containing M_P^{-1} explicitly. A deeper understanding of supersymmetry corrections, containing inverse powers of M_P , should however be undertaken so that we could trust the present approach as a consistent mathematical truncation of a more fundamental theory encompassing quantum gravity. We have proposed a mechanism which avoids at one loop, the instability of the tree-level hierarchy $m_W \ll M_X$ produced by the soft breaking terms of the supergravity-induced supersymmetry breaking.

Finally, the models considered here contain the independent scales M_P , M_X , and $m_{3/2}$, while supersymmetry is broken by pure supergravity effects at an intermediate scale $M_S = (m_{3/2} M_P)^{1/2}$. The corresponding value of $m_{3/2} = O(m_W)$ has been shown²⁴⁾ not to be in conflict with cosmological constraints²⁵⁾. It would be interesting to see whether M_X and $m_{3/2}$ could be generated dynamically or whether the GUT theory could be spontaneously broken at M_P as attempted in Ref. (14).

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REFERENCES

- 1) L. Maiani, in Proceedings of the Summer School of Gif-sur-Yvette (1979),
p. 3.
E. Witten, Nucl. Phys. B188 (1981) 513.
- 2) S. Ferrara and B. Zumino, Nucl. Phys. B79 (1974) 413.
A. Salam and J. Strathdee, Phys. Lett. 51B (1974) 353.
B. de Wit and D.Z. Freedman, Phys. Rev. D 12 (1975) 2286.
- 3) J. Wess and B. Zumino, Phys. Lett. 49B (1974) 52.
J. Iliopoulos and B. Zumino, Nucl. Phys. B76 (1974) 310.
S. Ferrara, J. Iliopoulos and B. Zumino, Nucl. Phys. B77 (1977) 413.
S. Ferrara and O. Piguet, Nucl. Phys. B93 (1975) 261.
M.T. Grisaru, M. Rocek and W. Siegel, Nucl. Phys. B159 (1979) 429.
- 4) A. Masiero, D.V. Nanopoulos, K. Tamvakis and T. Yanagida, Phys. Lett. 115B
(1982) 380.
- 5) F. Buccella, J.-P. Derendinger, S. Ferrara and C.A. Savoy, Phys. Lett. 115B
(1982) 375.
- 6) S. Weinberg, Phys. Rev. Lett. 48 (1982) 1776.
- 7) R. Barbieri, S. Ferrara, D.V. Nanopoulos and K.S. Stelle, Phys. Lett. 113B
(1982) 219.
- 8) E. Cremmer, S. Ferrara, L. Girardello and A. Van Proeyen, Phys. Lett. 116B
(1982) 231 and preprint CERN TH.3348 (1982).
- 9) H.P. Nilles, M. Srednicki and D. Wyler, preprint CERN TH.3432 (1982).
- 10) E. Cremmer et al., Phys. Lett. 79B (1978) 231; Nucl. Phys. B147 (1979) 105.
- 11) R. Arnowitt, A.H. Chamseddine and P. Nath, Northeastern University preprints
NUB Nos. 2559, 2565 (1982).
- 12) R. Arnowitt, A.H. Chamseddine and P. Nath, Phys. Rev. Lett. 49 (1982) 870,
NUB No. 2565 (1982), revised versions of Ref. 11).
- 13) R. Barbieri, S. Ferrara and C.A. Savoy, preprint CERN TH.3365 (1982).
- 14) L. Ibanez, preprint CERN TH.3374 (1982).
- 15) H.P. Nilles, M. Srednicki and D. Wyler, preprint CERN TH.2461 (1982).
- 16) A.B. Lahanas, preprint CERN TH.2467 (1982).
- 17) S. Polshinski and L. Susskind, SLAC preprint SLAC-PUB 2924 (1982).
- 18) J. Ellis, D.V. Nanopoulos and K. Tamvakis, preprint CERN TH.3418 (1982).
- 19) H.P. Nilles, preprints CERN TH.3330 and TH.3398 (1982).

- 20) E. Cremmer, S. Ferrara, L. Girardello and C.A. Savoy, in preparation.
- 21) E. Cremmer, P. Fayet and L. Girardello, Ecole Normale Supérieure preprint LPTENS 82/30 (1982).
- 22) L. Girardello and M.T. Grisaru, Nucl. Phys. B194 (1982) 65.
- 23) J. Ellis, D.V. Nanopoulos and S. Rudaz, Nucl. Phys. B202, 43 (1982).
- 24) J. Ellis, A.D. Linde and D.V. Nanopoulos, preprint CERN TH.3356 (1982).
- 25) S. Weinberg, Phys. Rev. Lett. 48 (1982) 1303.

FIGURE CAPTIONS

Fig. 1 : One-loop superfield diagrams contributing to a Higgs doublet $(\text{mass})^2$ of $O(m_{3/2}^2 M)$. The various contributions are obtained by inserting spurious chiral fields to each propagator and to each vertex (see text).

Fig. 2 : As in Fig. 1, for two-loop diagrams.

Fig. 1

Fig. 2a

Fig. 2b