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LOCAL SUPERSYMMETRY AND THE PROBLEM OF THE MASS SCALES*

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ABSTRACT

Spontaneously broken supergravity might help us to understand the puzzle of the mass scales in grand unified models. We describe the general mechanism and point out the remaining problems. Some new results on local supercolor are presented.

To begin let us remind you of the specific motivation that led to the application of supersymmetry[1] to grand unified models. It was the problem of three mass scales that we did not understand in terms of each others. These were the Planck mass $M_P \sim 10^{19}$ GeV which we know from the gravitational interactions, the speculative grand unification scale $M_X \sim 10^{14}$-10^{16} GeV related to proton decay and $M_W \sim 100$ GeV, the breakdown scale of the weak interactions. A first inspection of these three scales shows that M_W is tiny compared to the other scales: $M_W \approx 0$. One would like to understand why this is the case. An explanation could be a symmetry that keeps M_W small. Since the breakdown of the weak interactions is realized through vacuum expectation values of fundamental scalar fields we have only one choice for such a symmetry: Supersymmetry. In order to have $M_W \neq 0$ supersymmetry should be (spontaneously) broken at a scale M_S that is related to M_W. Once such a relation is established there remains the question whether one has gained anything. One has replaced M_W by M_S and one has now to face the problem to explain M_S in terms of M_X and M_P. There have however been some improvements. The first is a technical improvement due to the special behavior of supersymmetric field theories. It gives us the possibility that mass relations established at the tree graph level remain stable in perturbation theory. Secondly it is possible that M_S is much larger than M_W. The ratio can be as big as $M_S/M_W \sim 10^9$ as shown in the most promising globally supersymmetric models.[2] Such a large scale, however, leads necessarily to the introduction of supergravity since the gravitino mass $m_{3/2}$ is given by M_S^2/M_P which is now as big as the

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weak scale. Spontaneously broken local supersymmetry has additional desired properties. It admits vanishing vacuum energy in the case of broken supersymmetry[3] (which was impossible in the global case). Through the superHiggs effect the goldstino is absorbed; and finally it admits the Planck scale M_P to appear explicitly in the model. Maybe the small scale M_W can only be understood in terms of the two scales M_X and M_P.[4]

One thus is tempted to apply N = 1 supergravity to the models of high energy physics.[5,6] N = 1 supergravity is a nonrenormalizable theory and we can at best use it as an effective Lagrangian approach with the central assumption that there exists a satisfactory theory of gravity (which we do not know yet) for which N = 1 supergravity is a (good) approximation. This final theory should provide us with a cutoff for our nonrenormalizable approximation.

In such a model we assume that the gravitino mass is of the order of the weak scale and there are now two questions to be answered:

(1) How does one obtain a small scale $m_{3/2}$?

(2) How does the breakdown of local supersymmetry induce the breakdown of the weak interactions ?

We restrict ourselves to models where $m_{3/2}$ is the only small scale and where the weak interactions are restored in the limit $m_{3/2} \to 0$. Before we start our discussion we give the necessary formulae following the work of Cremmer et al.[5] for a single chiral superfield (z,x,h) coupled to supergravity. The model is characterized by a Kähler potential ($M = M_P$)

$$G(z,z^*) = -\frac{k}{M^2} - \log\left(\frac{|g|^2}{M^6}\right) \tag{1}$$

where g(z) is the superpotential and

$$k = -3M^2 \log(-\phi/3) \tag{2}$$

represents the choice of kinetic terms. Globally "normal" kinetic terms correspond to[7]

$$\phi = -3\left(1 - \frac{zz^*}{3M^2}\right) . \tag{3}$$

To discuss the model in supergravity one often makes the simplifying assumption of minimal kinetic terms

$$\phi = -3 \exp\left(-\frac{zz^*}{3M^2}\right) \tag{4}$$

which we will also use throughout this paper. In the presence of supergravity the scalar kinetic terms become

$$M^2 \frac{\partial G}{\partial z \partial z^*} |\partial_\mu z|^2 = M^2 G_{1zz^*} |\partial_\mu z|^2 . \tag{5}$$

The potential is given by

$$V = -M^4 \exp(-G)\left[3 + \frac{|G_{1z}|^2}{(G_{1zz^*})}\right] . \tag{6}$$

The transformation law for the chiral fermion is

$$\delta\chi = 2M^2 \exp(-G/2)\, \frac{G_{1z^*}}{\sqrt{-2G_{1zz^*}}}\, \varepsilon + \dots \tag{7}$$

where the dots denote field dependent terms. Formula (7) is the relevant expression to decide whether supersymmetry is broken or not.

We can now compare the corresponding formulae in the global and local case (we use minimal kinetic terms). The potential

$$V_L = \exp\left(\frac{zz^*}{M^2}\right)\left[|F_L|^2 - \frac{3|g|^2}{M^2}\right] \quad ; \quad V_G = |F_G|^2 . \tag{8}$$

The order parameter

$$F_L = \frac{\partial g}{\partial z} + \frac{z^*}{M^2} g \quad ; \quad F_G = \frac{\partial g}{\partial z} . \tag{9}$$

The supersymmetry by breaking scale is given by

$$M_S^2 = |F_L| \exp\left(\frac{zz^*}{2M^2}\right) \quad ; \quad M_S^2 = |F_G| = E_{vac}^2 . \tag{10}$$

Observe that the relation $M_S = E_{vac}$ is no longer valid in the local case. The massless Goldstino is absorbed by the gravitino which has a mass

$$m_{3/2} = \frac{M_S^2}{\sqrt{3}\, M} .$$

This concludes our presentation of the notation.

The most exciting property of $N = 1$ supergravity is the possibility to induce large mass gaps. A large hierarchy of mass scales can appear in which two large masses induce a small mass. This behavior has first been observed in a pure Yang Mills gauge theory coupled to supergravity.[4,8] A condensation of the gauginos λ at a scale $\langle\lambda\lambda\rangle = \mu^3$ was shown to break local supersymmetry at a scale $M_S^2 \sim \mu^3/M$ resulting in a gravitino mass $m_{3/2} \sim \mu^3/M^2$. A scale μ as large as the grand unification scale M_X and the scale M induce a small gravitino mass of order of a few TeV. A closer look at the situation showed that in this case nothing else could have happened since we know that the $\lambda\lambda$-condensation does not break global supersymmetry.[9] The breakdown scale M_S^2 has to be suppressed by $1/M$ to disappear in the global limit $M \to \infty$.

Meanwhile the general coupling of Yang Mills interactions to N = 1 supergravity has been worked out.[6] Consider a pure gauge theory and a chiral superfield in supergravity.[10] The transformation of the chiral fermion now reads as follows

$$\delta\chi = \left(\exp\,(-G/2)\,G' + \frac{1}{4}\,f'(z)\,\bar{\lambda}\lambda\right)\varepsilon \quad . \tag{11}$$

Suppose that G' vanishes at the minimum. A condensation of the gauge fermion breaks supersymmetry provided that $f'(z) = \partial f(z)/\partial z$ is nonzero. The function f(z) denotes nonminimal kinetic terms for the gauge interactions. Such nonminimal kinetic terms are known to exist in extended supergravities, whereas in N = 1 supergravity this function is a free parameter. Suppose we choose

$$f(z) = 1 + \sigma\,\frac{z}{M} \quad . \tag{12}$$

The condensate $\langle\lambda\lambda\rangle = \mu^3$ then leads to a breakdown of supersymmetry at the scale

$$M_S^2 = \frac{1}{4}\,\sigma\,\frac{\mu^3}{M} \tag{13}$$

and a gravitino mass

$$m_{3/2} = \frac{\sigma\mu^3}{4\sqrt{3}\,M^2} \quad . \tag{14}$$

This is true for general f as long as the vacuum expectation values of the scalar fields are not much larger than the Planck scale.

The condensate will in general lead to a cosmological constant as can be seen from the general form of the potential

$$V = -3\exp\,(-G) + \left|\exp\,(-G/2)\,G' + \frac{1}{4}\,f'\,\lambda\lambda\right|^2 \quad . \tag{15}$$

This cosmological constant can however be cancelled by the scalar sector at the cost of one fine tuning of parameters. This will be seen later in special examples.

The most important result is Eq. (14), $m_{3/2} \sim \sigma\mu^3/M^2$. In principle one could imagine such a relation to occur in models with only chiral superfields. Suppose one has a superpotential g(z) with an intrinsic scale μ. In general one would then expect to have at the minimum $g_o = \lambda\mu^3$ where λ is a coupling constant. If global supersymmetry would be broken this would lead to $M_S^2 \sim \lambda\mu^2$ not much different from the scale μ. Suppose now however that we have broken supergravity with vanishing vacuum energy. Formulae (8) and (10) then lead to

$$M_S = \frac{\sqrt{3}\,|g|}{M} \sim \frac{\lambda\mu^3}{M} \tag{16}$$

comparable to (13). Unfortunately nobody was able up to now to find a superpotential that leads to (16) where vanishing vacuum energy is achieved for the absolute minimum of the potential. The only known superpotential that has broken supersymmetry and $E_{vac} = 0$ at the absolute minimum is[5]

$$g(z) = m^2(z+\beta) \tag{17}$$

with $\beta = \pm(2-\sqrt{3})M$. The supersymmetry breakdown is of order m and one has to choose $m \sim 10^{11}$ GeV to obtain a gravitino mass in the TeV range. In the dynamical case [Eq. (14)] the scale μ could, however, be as large as the grand unification scale to obtain $m_{3/2} \sim O(M_W)$.

Given now a scale $m_{3/2} \sim O(M_W)$ we are confronted with the question how to apply the supersymmetry breaking[11-20] to models of quarks, leptons and Higgs particles. The supersymmetry breakdown has to appear in general in a distant (hidden) sector in order to keep the splittings in the observable particle spectrum as small as $m_{3/2}$. This can be achieved of the hidden sector couples only gravitationally to the observable sector. The superpotential is split into two pieces, e.g.

$$g(z,L_i) = g(z) + g(L_i) \tag{18}$$

where z denotes the hidden fields and L_i the observable fields. We now want to discuss the question under which circumstances the breakdown of supergravity can induce the breakdown of the weak interactions and thus provide a link between $m_{3/2}$ and M_W. We will split this discussion into two parts and first discuss the $SU(3)\times SU(2)\times U(1)$ model in the TeV-region and later on include grand unified models.

In the first case the L_i in (18) denote light fields; quarks, leptons, etc. According to our motivation we demand the naturalness condition[15] that $g(L_i)$ be scale invariant. Superpotentials that violate this condition necessarily contain explicit small mass parameters (~100 GeV) for which we do not understand why they are small. This condition gives

$$\sum_i \frac{\partial g}{\partial L_i} L_i = 3g \quad . \tag{19}$$

The low energy theory contains besides the quark and lepton superfield Higgses in the $2+\bar{2}$ representation of SU(2). This is however not enough if the naturalness condition is imposed since $g(L_i)$ would then consist only of the Yukawa couplings and have a Peccei Quinn symmetry under which H and $\bar{H}$ have the same charge. The simplest extension is the inclusion of a singlet Y with $YH\bar{H}+Y^3$ couplings.[15]

To discuss the low energy potential it is convenient to go to the flat limit $M\to\infty$, $m_{3/2}$ fixed.[13] One then arrives at a softly broken globally supersymmetric theory. The soft breaking term include gaugino masses, scalar masses and scalar trilinear couplings. If one restricts oneself to a pure scalar hidden sector with minimal kinetic terms and imposes (19) the low energy potential turns out to be[15]

$$V_{LE} = \left|\frac{\partial g}{\partial L_i}\right|^2 + Am_{3/2}(g+g^*) + m_{3/2}^2 |L_i|^2 \quad . \tag{20}$$

More general forms[13,20] can occur if one does not impose (19). A is a pure number and depends on the hidden sector. The sppearance of this terms breaks all R-symmetries that might have been present in the model. In the case under consideration the coefficient of the last term is universal. We will see later on that this need not be the case in general. Gaugino masses are generally expected[6] of order $m_{3/2}$, but could be smaller. Radiative corrections in connection with this masses break the universality of the $m_{3/2}^2|L_i|^2$ terms, such that even negative mass^2 for the Higgses might occur.[4,14] But let us first discuss the tree graph level situation. Since $g(L_i)$ is scale invariant we know that $V_{LE}=0$ at the point where all fields have vanishing vacuum expectation values. To induce a breakdown of $SU(2)\times U(1)$ V_{LE} must admit solutions at negative energies. The trilinear term is the only term that can give negative contributions. For $A=3$ one obtains from (20) using (19)

$$V_{LE} = |g_{1i} + m_{3/2}\, L_i^*|^2 \geq 0 \quad . \tag{21}$$

This shows that $A \geq 3$ is a necessary condition to induce an $SU(2)\times U(1)$ breakdown at the <u>tree graph level</u>, and this result is still independent of the special form of $g(L_i)$. A depends on the hidden sector. For $g = m^2(z+\beta)$ one obtains $A = 3-\sqrt{3}$. Even of one allows more scalars in the hidden sector one usually obtains this value as long as one uses (18). A model with $A=3$ has been found[16] by replacing (18) by

$$g(z,L_i) = m^2(z+\beta)\exp\left(\frac{g(L_i)}{m^2 M}\right) \quad . \tag{22}$$

This model has the nice property that A does not receive quadratic divergencies in one loop gravity corrections.[21] Generalizations of (22) have led to models with $A > 3$.

If one includes a strongly interacting Yang Mills theory in the hidden sector the situation becomes more general. The low energy potential reads[10]

$$V = |g_{1i}|^2 + Am_{3/2}(g+g^*) + Bm_{3/2}^2|L_i|^2 \tag{23}$$

with

$$A = -\sqrt{3}\,(z_o/M) \tag{24}$$

and

$$B = -2 + M^2|G_z'|^2 + \frac{1}{2}\left(G'\,\frac{f'\lambda\lambda}{4m_{3/2}} + h.c.\right) \quad . \tag{25}$$

The main difference is now that B can vary. Weak interaction breakdown at the tree level requires now either $A \geq 3\sqrt{B}$ or $B \leq 0$.

The cosmological constant induced by the condensate has to be cancelled by the scalar sector. Let us use the superpotential $g = m^2(z+\beta)$ to do this. We have explicitly computed two cases $\beta = 0$ and $\beta = 2M$. In both cases one is able to cancel the cosmological constant by adjusting m^2 to special values.[10] In the case $\beta = 0$ we obtain

$$A = -1.33 \quad ; \quad B = 1.59 \tag{26}$$

for $\beta = 2M$

$$A = \sqrt{3} \quad ; \quad B = -2 \tag{27}$$

It is evident that we can reach all values $1.59 \geq B \geq -2$ by varying β in the range $0 \leq \beta \leq 2M$. One thus can obtain models where B is close to zero. B remained still universal and this might cause problems as has been pointed out by Frère, Jones and Raby.[18] They showed that in models with $A \geq 3\sqrt{B}$ the absolute minimum usually corresponds to broken electric charge, due to the small Yukawa coupling that is responsible for the electron mass. This is a serious problem but it could perhaps be cured with cosmological arguments — the absolute minimum is separated by a high barier.[22]

This problem exists strictly because of the universality of B. This universality, however, is broken by radiative corrections which enhances the B values for quarks and leptons and reduces those for the Higgses of the model has a top quark mass larger than 20 GeV.[23] A model with A close to 3 and B_0 close to one would not suffer from this problem. It might even be that B is changed to negative values by radiative corrections. Starting with $B = 1$ this however requires a large lower bound $m_t \sim 60$ GeV on the top quark mass.[24] It might therefore be useful to consider models with potential like (23) which have a small B.[10,19]

It is thus possible to construct models in which the gravitino mass induces the breakdown of the weak interactions, and which contain no small mass parameters. The breakdown of $SU(2) \times U(1)$ is solely induced by supergravity which can be read off from (23); in the limit $m_{3/2} \to 0$ $SU(2) \times U(1)$ is restored.

The next step is to include grand unification. The superpotential $g(L_i)$ now contains also heavy fields. For these heavy fields we allow explicit mass parameters μ of the order of the grand unification scale. For the light fields we however still impose condition (19). This discussion is still relevant here since there is a hidden way[12,20] to break this condition in grand unified models, which usually leads to a breakdown of $SU(2) \times U(1)$ through a fine tuning. Let us take a simple example to explain this

$$g = \lambda A_{24} H_5 \bar{H}_{\bar{5}} + m H_5 H_{\bar{5}} \quad . \tag{28}$$

In this model one has to fine tune to keep the Higgs doublets massless. One solves the equations $|g_{1i}| = 0$ and determines the vacuum expectation value of A_{24}. One then adjusts m in such a way that the Higgs doublets remain massless. This is the right way to fine tune in a globally supersymmetric model. In local supersymmetry one has to fine tune differently. The vev of A_{24} is determined from $|g_{1i} + (L_i^*/M^2)g| = 0$ and differs from the one in the global theory slightly by an amount of order $m_{3/2}$. If one now fine tunes in the global limit one induces a small $m\bar{H}H$ mass parameter in the local theory. This leads sometimes to a breakdown of $SU(2) \times U(1)$, but this is not a breakdown induced by supergravity. It can be removed by a slightly different choice of the fine tuning procedure. Models that avoid the fine tuning through group theoretical reasons[25] (like models with 75, 50 and $\overline{50}$ representations) immediately rule out this possibility. As a result, the low energy effective potential remains unchanged.

To discuss grand unified models we use a toy model with one light (L) and one heavy field (B) and a superpotential[26]

$$g = \mu B^2 + \lambda_1 B^3 + \lambda_2 B^2 L + \lambda_3 L^3 + \lambda_4 BL^2 \quad . \tag{29}$$

In the limit $M \to \infty$ and μ, $m_{3/2}$ fixed this leads to the following potential

$$\begin{aligned} V = & \left|2\mu B + 3\lambda_1 H^2 + 2\lambda_2 BL + \lambda_4 L^2\right|^2 \\ & + \left|\lambda_2 B^2 + 3\lambda_3 L^2 + 2\lambda_4 BL\right|^2 + Bm_{3/2}^2(|B|^2 + |L|^2) \\ & + Am_{3/2}(\lambda_1 B^3 + \lambda_2 B^2 L + \lambda_3 L^3 + \lambda_4 BL^2 + h.c.) \\ & + (A-1)m_{3/2}\,(\mu B^2 + h.c.) \end{aligned} \tag{30}$$

Observe that all splittings are of order $m_{3/2}$ except for the last term which is of order $\mu m_{3/2}$. Thus the heavy fields are split by a large amount, and one has to worry whether this might induce large splittings in the low energy sector through radiative corrections. The light fields are coupled to the heavy ones through the couplings λ_2 and λ_4. Let us discuss the two terms separately.

The effect of λ_4 can be seen in the tadpole of Fig. 1. It gives a contribution $\mu m_{3/2}(L^2 + L^{*2})$ in leading order. This is however cancelled by the graph in Fig. 2. What remains are contributions of order $m_{3/2}^2$. Graphs where a vertex $\mu\lambda_1$ (in Fig. 1) is replaced by $Am_{3/2}\lambda_1$ also give contributions of order $m_{3/2}^2(L^2 + L^{*2})$. Thus the hierarchy remains stable. The exercise shows however that $m_{3/2}$ cannot be much larger than the weak interaction scale M_W, even of the parameters A and B are very small.

We proceed to discuss the terms proportional to λ_2 (from $\lambda_2 B^2 L$). A graph like Fig. 3 which contains $(A-1)\mu m_{3/2}$ and $m_{3/2}\lambda_2$ terms explicitly is not cancelled by any other graph and contributes with $\mu m_{3/2}$ to the mass of the light particles. The reason is the light particle exchange in Fig. 3. In the previous case (λ_4) these

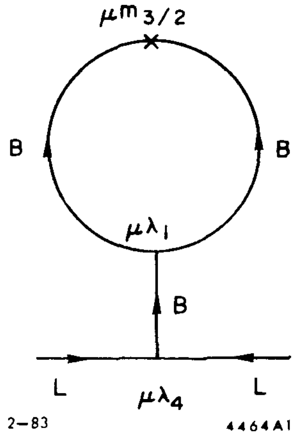

Fig. 1. Potentially dangerous contribution to light particle masses.

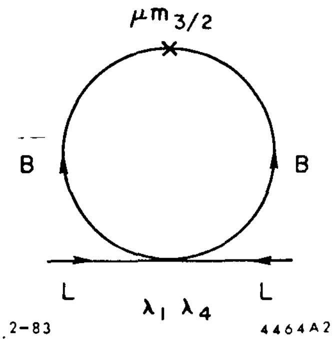

Fig. 2. Cancels contribution of Fig. 1 to leading order.

contributions were suppressed by the heavy propagator. As a result the tree graph level hierarchy in our toy model is spoiled.

In realistic models this problem occurs only if the model contains light singlets that couple via $\lambda_2 LB^2$ to heavy particles, but most of the locally supersymmetric grand unified models have this disease. It appears in graphs like Fig. 4 where a light singlet connects the light Higgs doublets with the heavy triplets. The problem usually persists even in more complicated models that contain light singlets.[27] The only natural way to solve the problem with a light singlet is to impose a discrete symmetry that does not allow the feeddown of $\mu m_{3/2}$ to the light particles. Such a symmetry, however, makes it difficult to generate masses for the Higgs triplets, which however can be solved by imagination and a fairly large amount of group theory.[28] The cleanest way, however, is the absence of light singlets.

It seems at the moment that the marriage of localsupersymmetry and grand unification leads to phenomenologically acceptable models. It might even be possible to understand the small mass scale M_W in form of a conspiracy of the two big scales M_X and M_P as described. Ultimately one might understand M_X from gauged extended supergravities where the gauge interactions become strong at M_X and breaks the last remaining supersymmetry. The small gravitino mass can induce the breakdown of $SU(2) \times U(1)$ in a natural way. This could be achieved in several ways depending on the parameters of the special model. Those parameters, however, can only be fixed by new experimental input.

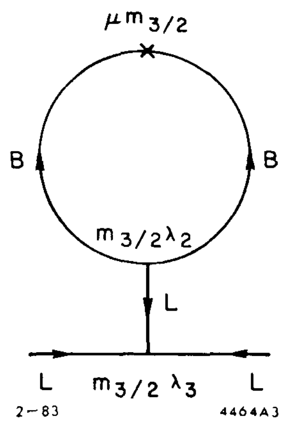

Fig. 3. Tadpole that spoils the hierarchy in the toy model.

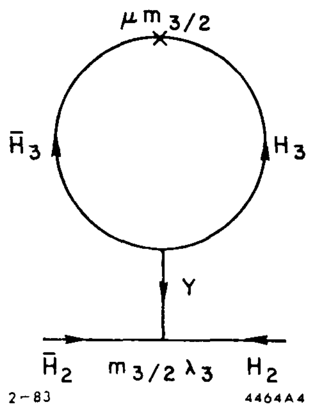

Fig. 4. The contribution that is relevant in grand unified models.

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