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HIERARCHIES FROM MINIMAL SUPERGRAVITY

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A B S T R A C T

We discuss in detail the properties of minimal $N = 1$ supergravity theories, such as the possibility of soft supersymmetry breaking in the form of F-type mass splittings in the effective low energy theory. The dynamical generation of hierarchies among mass scales, which are not fixed at tree level, requires the existence of a light singlet. The conventional problems connected with light singlets do not exist. A standard model with M_P as the only input mass scale is presented; its properties suggest an embedding into a GUT with natural hierarchies.

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1. - INTRODUCTION

Recently an interesting scenario has been developed, in which the hierarchy between the weak interaction scale, M_W , and the Planck mass, M_P , can be explained naturally¹⁾⁻⁵⁾. The scenario is based on the Lagrangian of $N = 1$ supergravity⁶⁾ with non-canonical kinetic energies¹⁾⁻⁹⁾. These are chosen such that the scalar fields parametrize a Kähler manifold with a $U(n,1)/U(n) \times U(1)$ structure. This choice of the kinetic energies is minimal from the superspace point of view^{*)},⁷⁾⁻¹⁰⁾, and similar coset spaces appear in extended supergravity theories^{1),11)}.

All minimal theories have some common properties, which are independent of the choice of the superpotential: the supertrace, taken over the components of all chiral multiplets, vanishes identically (except for the gravitino contribution)^{1),4),5)}. This already indicates that the effective low energy theory will not contain supersymmetry breaking positive definite masses for the scalar fields, as is the case for canonical kinetic energies^{6),12)}.

Furthermore, it can be shown that the gravitino mass $m_{3/2}$ is never fixed at tree level if supersymmetry is spontaneously broken and the tree-level cosmological constant vanishes⁵⁾. In the case of canonical kinetic energies, this can only be achieved for particular choices of the superpotential¹³⁾.

This property of minimal theories can be used to fix $m_{3/2} \ll M_P$ with the help of radiative corrections²⁾⁻⁴⁾ without introducing scales different from M_P into the original theory.

However, one can still put in different scales by hand⁹⁾. Assuming $m_{3/2} \ll M_P$, the corresponding effective low energy theory has been worked out in Refs. 9) and 7) (2nd Ref.), but a mechanism which fixes the sliding gravitino mass in this case has not been proposed.

In the present approach, a possible source of explicit soft supersymmetry breaking in the effective low energy theory are gaugino masses m_g via non-minimal kinetic energies for the gauge superfields⁴⁾. However, in this case the contributions of the radiative corrections to the vacuum energy density will be in general of the order of $m_g^2 M_P^2$ or $m_g^2 M_{GUT}^2$. This is very undesirable, because

*) Hence we will subsequently call these theories minimal.

one has carefully arranged m_g to be undetermined at tree level, and one wants to fix it with the help of radiative corrections of the order of $m_g^4 (2)^{-4}$.

Another possibility for soft supersymmetry breaking in the context of theories with M_p as the only scale are F-type splittings of the order of $m_{3/2}^2$ among the masses of scalar fields⁵⁾. This is not a general phenomenon in this class of theories, but it depends on the particular form of the superpotential. The necessary ingredients will be discussed in more detail below.

With this form of soft supersymmetry breaking, an interesting result concerning the maximal contribution of the radiative corrections to the vacuum energy can be derived. Let us describe the F-type splitting with the help of a spurion field in superspace:

$$\int d^2\theta \, \eta \, \phi^2 + \text{h.c.} \quad (1.1)$$

Here ϕ denotes a chiral superfield, and the spurion field η is given by

$$\eta = m^2 \theta^2 \quad (1.2)$$

Now, to any number of loops, the quantum corrections to the effective action can be written in the form:

$$\int d^4\theta \, \phi^a \eta^b M^{2-a-b} \quad (1.3)$$

For simplicity we have not distinguished between ϕ and $\bar{\phi}$ or η and $\bar{\eta}$. M denotes the (supersymmetric) mass of the fields inside the loops. In (1.3) we have omitted powers of coupling constants, logarithms and numerical factors. Supercovariant derivatives, which may appear under the $d^4\theta$ integral, will not modify the subsequent considerations. Now we make the following assumptions concerning the VEVs of the scalar components A and the auxiliary components F of the chiral superfields:

$$\langle A \rangle \lesssim M \quad (1.4a)$$

$$\langle F \rangle \lesssim m^2 \quad (1.4b)$$

In practice these assumptions are fulfilled at least after shifting the superfields by a constant. Equation (1.4b) means that all supersymmetry breaking is induced by the soft explicit breaking, and no additional spontaneous breaking takes place. After inserting (1.2) and (1.4) into (1.3) it is easy to see that (1.3) is bounded by m^4 .

In order to obtain this result it has been crucial to observe that the spurion field in (1.1) and (1.2) carries mass dimension +1. Hence this bound will in general be violated in the presence of supersymmetry breaking involving dimensionless spurion fields (e.g., gaugino masses) unless all fields have masses $M \sim m$. Note that here no assumption on the magnitude of M had to be made.

In the next section we will discuss in more detail under which circumstances F-type splitting can appear in minimal $N = 1$ supergravity theories. A model which exhibits all the interesting features of this class of theories will be investigated in Section 3. In Section 4 we study the radiative corrections and the mechanism which fixes $m_{3/2} \ll M_P$ dynamically. This turns out to work differently compared with the scenarios investigated previously²⁾⁻⁴⁾. Conclusions and a promising outlook will be given in Section 5.

2. - MINIMAL $N = 1$ SUPERGRAVITY AND COSET TRANSFORMATIONS

The Lagrangian of $N = 1$ supergravity is characterized by a real function $G(z, z^*)$ of the complex scalar matter fields z , which is composed of a real function $K(z, z^*)$ (determining the matter kinetic energies) and the analytic superpotential $g(z)$ ^{6),*)}:

$$G(z, z^*) = -3 \ln K(z, z^*) + \ln |g(z)|^2 \quad (2.1)$$

In addition, an analytic function $f_{ab}(z)$ determining the gauge kinetic energies has to be specified⁶⁾. The class of theories under investigation is characterized by the minimal ansatz for $K(z, z^*)$ from the superspace point of view^{1),4)-10)}:

*) Here and in the following we use the convention $2.4 \times 10^{18} \text{ GeV} = (\bar{M}_P / \sqrt{8\pi}) = M_P = 1$.

$$K = 1 - \frac{|z|^2}{3} \quad (2.2)$$

This choice differs from the common ansatz leading to canonical kinetic energies⁶⁾⁻¹²⁾:

$$K = \exp \left\{ -\frac{|z|^2}{3} \right\} \quad (2.3)$$

It has already been pointed out that in the minimal case (2.2) the inner space is an Einstein space with non-vanishing curvature¹⁾⁻⁵⁾. The n scalar fields form a Kähler manifold with $U(n,1)/U(n) \times U(1)$ structure: the kinetic energies are not only invariant under trivial $U(n)$ -rotations, but also under non-linear (non-compact) transformations belonging to the corresponding coset. These transformations are characterized by a complex vector $\vec{a}$ in field space. Let us denote the field in the direction of $\vec{a}$ by y . Then the coset transformations take the form ($|\vec{a}| = a$)⁵⁾:

$$y \rightarrow y' = \frac{y + \sqrt{3} a}{1 + \frac{a^* y}{\sqrt{3}}} \quad z_i \rightarrow z'_i = z_i \frac{\sqrt{1 - |a|^2}}{1 + \frac{a^* y}{\sqrt{3}}} \quad (2.4)$$

Clearly it is possible to choose a different set of independent fields (or a different set of co-ordinates on the inner space). The authors of Ref. 4), for example, singled out a field z and worked with the fields ϕ_i , z related to the ones above through:

$$z = \frac{1}{2} \frac{\sqrt{3} - y}{\sqrt{3} + y} \quad \phi_i = \frac{z_i}{\sqrt{3} + y} \quad (2.5)$$

[up to another field redefinition $z \rightarrow f(z)$]. In these variables the superpotential g will of course have a different form, and K is now given by⁴⁾:

$$K = z + z^* - \frac{|\phi_i|^2}{3} \quad (2.6)$$

The transformations corresponding to (2.4) have been given in Ref. 3) (decomposed into imaginary translations, dilatations and conformal transformations) in the case of a single field z and $SU(1,1)$ symmetry. There the consequences of non-compact symmetries in minimal $N = 1$ supergravity theories have been discussed at length.

In the following we will continue to work with the variables corresponding to (2.2) and (2.4), since this choice exhibits more symmetry among the fields, and since there it might be easier to see which of the subsequent results are consequences of the $U(n,1)$ structure, and which not.

Inserting the transformation (2.4) into (2.2) and (2.1), its consequences can be written as follows⁵⁾:

$$K(z, z^*) \rightarrow K'(z', z'^*) = K(z', z'^*)$$

$$g(z) \rightarrow g'(z') = \frac{(1 + \frac{a^* y}{\sqrt{3}})^3}{(1 - |a|^2)^{3/2}} g(z'(z)) \quad (2.7)$$

Hence the kinetic energies remain invariant, whereas the form of the superpotential will, in general, change under such a transformation. Accordingly $U(n,1)$ is not an invariance of the theory (unless $g \equiv 0$), but the non-linear transformations (2.4) can still be useful, if one takes the corresponding modifications of the superpotential into account.

A number of properties of these modifications have already been given in Ref. 5): if $g(z)$ is not more than cubic in the fields z , the same will be true for $g'(z')$ in spite of the non-linear relation between z and z' . Furthermore the transformations (2.4) can always be used to shift the fields to $\langle z \rangle = 0$ at tree level. If the model under consideration exhibits supersymmetry breaking with vanishing cosmological constant, the modified superpotential will necessarily be in the "standard form":

$$g(y, z_i) = \mu \left(\frac{y}{\sqrt{3}} + 1 \right)^3 + O(z_i^2) \quad (2.8)$$

Such an ansatz for g has been present in all models based on minimal supergravity theories proposed so far (up to field redefinitions)^{2)-5), 9)}; according to the above statement it is not possible to avoid it. Here y denotes the field with $g'_y \neq 0$; after inserting (2.8) into the scalar potential of minimal theories^{5), 8), 9)} (with D-terms omitted)

$$V = K^{-2} \left\{ |g'_i|^2 - 3 \left| g - \frac{z_i g'_i}{3} \right|^2 \right\} \quad (2.9)$$

one finds that neither the VEV of y nor the gravitino mass

$$m_{3/2} = e^{\langle G \rangle / 2} \quad (2.10)$$

are fixed at tree level. This is a common property of all minimal theories, although the form (2.8) of the superpotential might not be manifest due to a different choice of independent variables¹⁾⁻⁵⁾.

In the following we will use (2.8) as a starting point, since this form turns out to be convenient for a discussion of the resulting effective low energy theory. Let us now have a closer look at the sliding y field. Its VEV is not fixed by the potential (2.9); what are the consequences for the superpotential, if we redefine the field y with the help of a (real) transformation (2.4)? These consequences are most easily discussed after introducing the variables

$$\gamma_{\pm} = \frac{y}{\sqrt{3}} \pm 1 \quad (2.11)$$

The most general cubic superpotential of the type (2.8) can be written as a sum of terms of the form

$$\mu \gamma_+^3 ; \quad \lambda z_i^3 ; \quad M z_i^2 \gamma_+ ; \quad N z_i^2 \gamma_- . \quad (2.12)$$

Now the modifications of g after a transformation (2.4), together with (2.7), are given by ($a = a^*$)

$$\begin{aligned} \mu &\rightarrow \mu' = \left(\frac{1+a}{1-a} \right)^{3/2} \mu & \lambda &= \lambda' \\ M &\rightarrow M' = \left(\frac{1+a}{1-a} \right)^{1/2} M & N &\rightarrow N' = \left(\frac{1-a}{1+a} \right)^{1/2} N \end{aligned} \quad (2.13)$$

In spite of the non-linear relation between the old and the new field variables, the modified superpotential has kept its form, with just rescaled coupling constants and masses. Hence the terms (2.12) are the natural building blocks of g , since each of them transforms separately in a simple way under the particular coset transformation considered here.

In general the explicit introduction of the term μy_+^3 in (2.8) can be avoided by choosing a different set of independent fields. Starting with an arbitrary g , which is not in the standard form (2.8), only one coefficient has to be fine tuned in order to ensure $V_{\min} = 0$. Nevertheless at least one VEV will always be undetermined and $m_{3/2}$ will always slide [see the example in Ref. 5)], but it will not be straightforward to obtain a sensible low energy theory including soft supersymmetry breaking, as will be discussed below. Also the field redefinition (2.5) would get rid of the term μy_+^3 , but now this "fine tuning" is hidden in the particular asymmetric form of the function K (2.6).

Inserting (2.8) (with $\langle z \rangle = 0$, y real) into (2.10) we find for the gravitino mass:

$$m_{3/2} = \mu \left(\frac{\sqrt{3} + \gamma}{\sqrt{3} - \gamma} \right)^{3/2} \quad (2.14)$$

We will be particularly interested in the case $m_{3/2} \ll 1$ achieved through a VEV of y near $-\sqrt{3}$. Now we will shift y back to have a vanishing VEV using the transformation (2.4) with

$$\alpha = -1 + 2\varepsilon^2, \quad \varepsilon \ll 1 \quad (2.15)$$

Now the fields have regained correctly normalized kinetic energies, but we have to take into account the modifications (2.13) of the superpotential. To leading order in ε these are given by

$$\begin{aligned} \mu &\rightarrow \mu' = \varepsilon^3 \mu & \lambda &= \lambda' \\ \mathcal{M} &\rightarrow \mathcal{M}' = \varepsilon \mathcal{M} & N &\rightarrow N' = N/\varepsilon \end{aligned} \quad (2.16)$$

In the following we will always assume that we had to shift back y from a VEV near $-\sqrt{3}$. However, instead of treating the VEV of y as a parameter, which is undetermined at tree level and to be determined by radiative corrections, we will work with a superpotential modified through (2.16), where now ε is a parameter to be fixed by radiative corrections.

We will show in the next section how $\varepsilon \ll 1$ can be achieved naturally; let us first investigate some of its consequences.

Since we assume that the original mass parameters in the superpotential are all of the order of M_P , we find in the modified superpotential masses of the type $M' = \varepsilon M \ll M_P$ and $N' = N/\varepsilon \gg M_P$. Since we furthermore assume an ultra-violet cut-off of the quantum fluctuations of $O(M_P)$, quantum effects of mass eigenstates with a mass $\sim N'$ will be negligibly small, i.e., these fields will be frozen in at their tree level VEVs. This philosophy is also dictated by regarding the action as describing an effective theory below the Planck scale.

Couplings between fields with masses $\sim M'$ and y , stemming from terms in the superpotential of the form $M z_i^2 y_+$, will be of $O(\varepsilon)$; hence y will approximately decouple from all propagating fields.

In general, we expect renormalizable supersymmetry violating interactions in the scalar sector from the second term involving $|g - z_i g'_i / 3|^2$ in the scalar potential (2.9). However, it is easy to verify that this term cancels exactly $|g'_y|^2$, if terms of the form $N z_i^2 y_-$ are not present in the superpotential. Hence masses $\sim N'$ have to be present, if we wish to obtain supersymmetry breaking in the scalar sector. The precise form of the supersymmetry breaking will depend on the complete superpotential; an example, which leads to the desired F-type breaking, will be given in the next section.

The field-dependent mass matrix of chiral fermions in minimal theories is given by⁵⁾:

$$M_{ij}(z) = K^{-3/2} (g / |g|) \left\{ g''_{ij} - \frac{2}{3} g'_i g'_j / g \right\} \quad (2.17)$$

Since we have $\langle g'_i \rangle = 0$ except for $i = y$, the fermionic interactions of all fields except for y will deviate from their globally supersymmetric form only through non-renormalizable interactions suppressed by powers of M_P (which, of course, are also always present in the scalar sector). Inserting the standard form of g [(2.8) together with (2.12) and (2.16)] into (2.17) one discovers again that all renormalizable interactions involving propagating fields and y are suppressed by ϵ .

Having discussed general properties of minimal theories we will investigate the features of a particular model in the next section.

3. - TREE LEVEL ANALYSIS OF A STANDARD MODEL

In this section we will investigate a theory with $SU(3) \times SU(2) \times U(1)$ gauge invariance at tree level; the breakdown of $SU(2) \times U(1)$ will be achieved through radiative corrections, which will be the subject of the next section. The model contains the standard gauge superfields, the usual chiral superfields with the

quantum numbers of the known quarks and leptons and in addition the following chiral superfields [listed with their SU(2) quantum numbers]^{*)}:

$$\begin{aligned}\tilde{S}, \tilde{T}, U, R &= \underline{3} \\ H, H' &= \underline{2} \\ Y, Z &= \underline{1}\end{aligned}\quad (3.1)$$

Our starting point for the superpotential will be the standard form (2.8) together with terms of the type (2.12):

$$\begin{aligned}g = & \mu y_+^3 + M_1 \tilde{S} \tilde{T} y_+ + \frac{N}{2} \tilde{T}^2 y_- + \alpha \tilde{S} U^2 \\ & + M_2 U R y_+ + H' (\beta \tilde{S} + \gamma U + \lambda Z) H + \dots\end{aligned}\quad (3.2)$$

For simplicity we have omitted in (3.2) the standard Higgs-matter couplings involving H and H' as well as obvious contractions of the SU(2) indices. Here the $\tilde{S}$ - $\tilde{T}$ sector will be the source of soft supersymmetry breaking, the fields U and Z allow for a particular form of radiative corrections (see the next sections) and R ensures $\langle S \rangle = \langle U \rangle = 0$ at tree level.

As we described in the previous section, we will work in the following with a modified superpotential, where we assume that we had to shift back y from a VEV near $-\sqrt{3}$. As in (2.16) the dimensionful parameters in g will be modified according to

$$\begin{aligned}\mu &\rightarrow \varepsilon^3 \mu & M_1, M_2 &\rightarrow \varepsilon M_1, \varepsilon M_2 \\ N &\rightarrow N/\varepsilon\end{aligned}\quad (3.3)$$

Now the scalar potential (apart from the D terms) is given by

$$V = K^{-2} \left\{ \sum_{i \neq Y} |g'_i|^2 + \frac{N}{M_P^2} \left(\varepsilon^2 \mu y_+^2 + \frac{M_1}{3} \tilde{S} \tilde{T} + \frac{M_2}{3} U R \right) \tilde{T}^{*2} + h.c. \right\} \quad (3.4)$$

The second term in (3.4) is due to an incomplete cancellation between $|g'_Y|^2$ and $-3|g_{-Z_1} g'_1/3|^2$. Clearly this term would not be affected by contributions from the trilinear Higgs-matter couplings neglected here. In (3.4) we have reintroduced the correct powers of M_P ; y_+ and y_- are now given by

^{*)} Unless it is necessary, we will not distinguish between chiral superfields and their scalar components. Subsequently, we will regard H' as a row vector, H as a column vector and $\tilde{S}$, $\tilde{T}$, U and R as 2x2 matrices.

$$y_{\pm} = \frac{y}{\sqrt{3}M_p} \pm 1 \quad (3.5)$$

As mentioned before, we assume

$$M_1 \sim M_2 \sim N \sim \mu^{1/3} \sim M_p \quad (3.6)$$

It is already obvious from the superpotential that the mass matrix in the $\tilde{S}$ - $\tilde{T}$ sector is not diagonal. After (3.3) it is diagonalized to leading order in ϵ with the help of the redefinitions

$$\begin{aligned} \tilde{T} &= T + \epsilon^2 \frac{M_1}{N} S \\ \tilde{S} &= S - \epsilon^2 \frac{M_1}{N} T \end{aligned} \quad (3.7)$$

After inserting (3.7) into (3.4) one obtains the following results. For a large range of parameters fulfilling (3.6), the absolute minimum of V with $V_{\min} = 0$ is assumed for vanishing VEVs of all fields except for y , which is undetermined and assumed to be zero after a shift. The appearance of y_+ and y_- (3.5), which has been motivated in the previous section, ensures a vanishing of the cosmological constant. The mass squared of T is $\sim N^2/\epsilon^2$, and the mixing of T with S is negligibly small. Hence we will put $T = 0$ subsequently. The second (supersymmetry violating) term of the potential (3.4) reads after (3.7) with $T = 0$:

$$\frac{\epsilon^4 M_1^2}{N M_p^2} \left(\mu \epsilon^2 y_+^2 + \frac{\epsilon^2 M_1^2}{3N} S^2 + \frac{M_2}{3} UR \right) S^{*2} + h.c. \quad (3.8)$$

Let us now have a look at the couplings involving the field y . Trilinear scalar couplings to matter fields z with masses M always behave like $\epsilon^n M z^2 y$ with $n > 0$ and can be ignored for $\epsilon \ll 1$. In addition we will neglect dimensionless couplings with positive powers of ϵ and non-renormalizable interactions. As a result, y decouples completely. Furthermore, quartic scalar couplings appearing in (3.8) as well as in (3.4) after expanding K^{-2} (these would break supersymmetry strongly) vanish with ϵ .

The remaining effective theory looks like a globally supersymmetric one with a superpotential

$$g_{\text{eff}} = \frac{1}{2} M_S S^2 + \alpha S U^2 + M_u UR + H'(\beta S + \gamma U + \lambda Z) H \quad (3.9)$$

and explicit soft supersymmetry breaking of the form (here S denotes the scalar component field)

$$m^2 \text{Re } S^2 \quad (3.10)$$

with

$$M_S = \varepsilon^3 M_1^2 / N \quad M_u = \varepsilon M_2$$

$$m^2 = 2 \varepsilon^6 M_1^2 / N M_P^2.$$

It is interesting to compare these masses with the gravitino mass. Therefore we step back to (3.2) together with (3.3) and find (μ has dimension M^3)

$$m_{3/2} = \varepsilon^3 \mu / M_P^2 \quad (3.11)$$

The most important part of this section is the evolution from the original superpotential (3.2) to the effective theory described by (3.9) and (3.10): the term $\sim \tilde{N} \tilde{T}^2 y_-$ in (3.2) is the original source of supersymmetry breaking; however, since the field $\tilde{T}$ tends to vanish identically due to its huge mass, we had to mix it with a light field $\tilde{S}$ in order to have some supersymmetry breaking left over in the light sector. In the effective theory (3.9) and (3.10), the field S has a supersymmetric mass M_S and F-type splitting $\sim m^2$; both M_S and m are of the order of the gravitino mass (3.11), or, in terms of ε and the Planck scale $\sim \varepsilon^3 M_P$. (Of course we chose the initial parameters such that $M_S > m$.) Other supersymmetric masses in the effective theory, as M_u in our case, turn out to be $\sim \varepsilon M_P$ or $m_{3/2}^{1/3} M_P^{2/3}$. The appearance of such mass scales has already been recognized in Ref. 3). Some speculations about its usability will be made in the last section. From this section we want to keep in mind the effective theory (3.9) and (3.10), and in particular the dependence of all mass scales on the as yet undetermined parameter ε as given by (3.10).

4. - RADIATIVE CORRECTIONS AND THE GENERATION OF HIERARCHIES

In this section we will investigate the radiative corrections in a theory described by (3.9) and (3.10)^{*)}, and we will show how the as yet undetermined parameter ϵ can be fixed naturally to be $\ll 1$. Let us first have a general look at the logarithmic divergences.

Of course, there will be the usual wave function renormalizations. They contribute to coupling constant renormalizations which have the following simple properties. They never introduce a new coupling constant (or mass) which has not been there from the beginning; they never change the sign of a coupling constant (at least after renormalization group improvement); finally, assuming an ultraviolet cut-off of $O(M_P)$ and a renormalization point near $m_{3/2}$, these renormalizations numerically never exceed $O(1)$ unless one allows for huge coupling constants. Hence these contributions to the effective scalar potential can be taken into account through a minor change in the parameters of the tree level potential. In the following we will assume that these modifications of the tree level parameters have been performed (corresponding to the use of running parameters with a renormalization point near $m_{3/2}$, where only the contributions to β and γ functions due to wave function renormalization have been taken into account), and subsequently we will not worry about wave function renormalization any more.

Without explicit soft supersymmetry breaking, there would not exist any further logarithmic divergences to all orders in perturbation theory. However, as becomes clear from previous work on radiative determination of sliding mass scales²⁾, contributions to the effective potential involving large logarithms are crucial in order to obtain hierarchies naturally. The only explicit soft supersymmetry breaking in our effective theory consists of F-type breaking among some scalar fields. With the help of the analysis of Grisaru and Girardello¹⁴⁾, one finds immediately that the only logarithmic divergences generated in this case are proportional to the auxiliary components of gauge singlet fields. Hence light singlets (heavy ones would decouple) will play an important role in our class of theories; this is the reason why we introduced the field Z into our model (3.2). Light singlets have had a rather bad press recently^{12),15)}; however, as we will discuss near the end of this section, the conventional problems will not arise with the type of supersymmetry breaking considered here.

*) We will, however, not discuss quantum effects involving the graviton or the gravitino.

In spite of the fact that we will obtain logarithmically divergent contributions to the effective potential as mentioned before, we will not introduce counter terms to the effective potential or its renormalization group improvement. After introducing a cut-off of $O(M_p)$ we will simply keep the large logarithm. This way the basic mechanism for the generation of hierarchies, namely an interplay between the finite and the logarithmically divergent contributions to the effective potential, becomes most obvious. Of course, a summation of the leading logarithms with the help of a renormalization group improvement¹⁶⁾ can be performed in order to improve the results quantitatively, but a consideration of the finite contributions would require a complicated mass-dependent renormalization scheme¹⁷⁾, which would make it difficult to see what happens. Let us now discuss the radiative corrections loop after loop.

At one loop the scalar fields H and H' receive negative masses squared from the diagram given in Fig. 1; for $m^2 \ll M_S^2$ these masses are given by

$$- \frac{\beta^2 m^8}{96 \pi^2 M_S^6} (|H|^2 + |H'|^2). \quad (4.1)$$

A similar diagram contributes to the masses of the U field; however, due to the large supersymmetric mass of U this contribution will be negligibly small.

Also at one loop the gauginos of the $SU(2)$ gauge vector superfield receive masses from the diagram given in Fig. 2; including their q^2 dependence they are given by (for $m^2 \ll M_S^2$)

$$m_g(q^2) = \frac{g_2^2 m^2}{8 \pi^2 M_S} \quad \text{for } q^2 = 0$$

$$m_g(q^2) = \frac{g_2^2 m^2 M_S}{4 \pi^2 q^2} \ln(q^2/M_S^2) \quad \text{for } q^2 \gg M_S^2 \quad (4.2)$$

Here m_g denotes the corresponding gaugino mass and g_2 the $SU(2)$ gauge coupling constant.

At two loops, the graphs shown in Fig. 3 lead to positive masses squared of all scalars A transforming non-trivially under $SU(2)$ of the order of¹⁸⁾

$$\sim \frac{2 g_2^4}{(16 \pi^2)^2} \frac{m^4}{M_S^2} C_2(A) |A|^2 \quad (4.3)$$

In (4.3) we assume again $m^2 \ll M_S^2$, and $C_2(A)$ denotes the eigenvalue of the quadratic Casimir operator corresponding to the $SU(2)$ representation of the field A .

Subsequently we will assume the following inequality among coupling constants and masses:

$$\beta^2 < \frac{g_2^4 M_S^4}{4\pi^2 m^4} \quad (4.4)$$

In particular for $m^2 \ll M_S^2$ (4.4) will be easily fulfilled. As a result the two-loop contribution (4.3) (with $A = H, H'$) will dominate over the one-loop contribution (4.1), and the potential will still be stable and positive semi-definite so far. In fact, the only contributions to the scalar potential, which will not be overwhelmed by (4.3), will either be proportional to SU(2) singlet fields [where (4.3) does not exist] or else be logarithmically divergent.

Let us put aside the radiative masses of SU(2) singlets for the moment and concentrate on logarithmic divergent contributions, which will arise first at the three-loop level. As mentioned above, they will be proportional to the auxiliary component of a gauge singlet superfield. The only singlet in our effective theory is the field Z , and the corresponding diagram is shown in Fig. 4. After replacing F_Z by the scalars $H^\dagger H$ through its equation of motion, the logarithmically divergent contribution of this diagram to the effective potential is estimated to be

$$\frac{\lambda^2 \gamma^2 \alpha \beta m^2}{(16\pi^2)^3} H^\dagger H \ln(M_u^2/\Lambda^2) + h.c. \quad (4.5)$$

Without the field U it would not have been possible to draw a similar tadpole diagram to any number of loops due to the discrete symmetry $S, Z, H^\dagger H \rightarrow -S, -Z, -H^\dagger H$, which is only broken by the couplings to U . Since for the range of parameters we are interested in M_u is much larger than all other masses or VEVs (see below), it is only M_u^2 which appears in the argument of the logarithm in (4.5) apart from the cut-off Λ^2 . This has been verified by an explicit computation of a diagram of the type shown in Fig. 4. (There exist others with H and H' interchanged or with reversed arrows everywhere, leading to the complex conjugate result.)

In order to discuss the complete scalar potential we still have to give its tree-level part. For this we have to show the equations of motion of all the auxiliary fields; suppressing again obvious SU(2) indices, these are given by [$H(H')$ is a column (row) vector, matrix multiplication is understood throughout]

$$\begin{aligned}
F_S &= M_S S + \alpha (U^2 - \frac{1}{2} \text{Tr} U^2) + \beta (H \cdot H' - \frac{1}{2} H' \cdot H) \\
F_U &= \alpha (\{S, U\} - \text{Tr} S U) + M_U R + \gamma (H \cdot H' - \frac{1}{2} H' \cdot H) \\
F_R &= M_U U \\
F_H &= H' (\beta S + \gamma U + \lambda Z) \\
F_{H'} &= (\beta S + \gamma U + \lambda Z) H \\
F_Z &= \lambda H' \cdot H
\end{aligned} \tag{4.6}$$

In addition we have the D terms; omitting the contributions from quark and lepton superpartners the auxiliary field D_2 of the SU(2) vector supermultiplet is given by:

$$D_2 = [S, S^\dagger] + [U, U^\dagger] + [R, R^\dagger] + (H H^\dagger - \frac{1}{2} H^\dagger H) - (H' H'^\dagger - \frac{1}{2} H'^\dagger H') \tag{4.7}$$

Including the F-type mass splitting among the scalars S, the tree-level potential reads:

$$V = \sum_i |F_i|^2 + \sum_j \frac{g_j^2}{4} \text{Tr} D_j^2 + m^2 \text{Re Tr} S^2 \tag{4.8}$$

where j runs over the three gauge groups SU(3), SU(2) and U(1).

Let us now concentrate on VEVs which could lead to a negative vacuum energy. At first sight this could be achieved through a VEV of S such that $\text{Re Tr} S^2$ is negative, but it is easy to verify that for M_S sufficiently large compared with m (which we will assume subsequently), the sum $|F_S|^2 + |F_R|^2 + |F_Z|^2 + m^2 \text{Re Tr} S^2$ is positive semidefinite.

Taking the radiative corrections into account and assuming that (4.4) holds, it is only (4.5) which can allow the vacuum energy to become negative. Hence we will investigate the scalar potential for a non-vanishing VEV of $H' \cdot H$. Minimizing the tree level potential (4.8) for $\langle H' H \rangle \neq 0$, it follows that the VEVs of H and H' will be given by [up to SU(2) rotations]:

$$\langle H \rangle = \begin{pmatrix} 0 \\ h \end{pmatrix}, \quad \langle H' \rangle = \pm (0, h) \tag{4.9}$$

$\langle H \rangle$ and $\langle H' \rangle$ are real and their relative signs are such that the over-all sign of (4.5) is negative independently of the sign of the diagram Fig. 4.

In the presence of the VEVs (4.9), those of the other fields will arrange themselves such that the rest of the scalar potential is minimized. Due to the huge value of M_u , the vanishing of F_R will always ensure $\langle U \rangle = 0$. To a very good approximation the VEV of S will arrange itself such that F_S vanishes:

$$\langle S \rangle = \frac{\beta h^2}{2 M_S} \begin{pmatrix} 1 \\ -1 \end{pmatrix} \quad (4.10)$$

With m small compared to M_S , this VEV of S , inserted into (4.8) and (4.3) with $A = S$, contributes to the vacuum energy proportional to h^4 , but with a coefficient which is small compared with a similar contribution stemming from $|F_Z|^2$.

The VEV of R will ensure the vanishing of F_u :

$$\langle R \rangle = \frac{\gamma h^2}{2 M_u} \begin{pmatrix} 1 \\ -1 \end{pmatrix} \quad (4.11)$$

With M_u very large, the insertion of (4.11) into (4.3) with $A = U$ will result in a negligible contribution to the vacuum energy.

The VEV of Z which leads to $F_H = F_{H'} = 0$ is given by

$$\langle Z \rangle = \frac{\beta^2 h^2}{2 \lambda M_S} \quad (4.12)$$

With these VEVs of all the fields, the remaining non-negligible contributions to the vacuum energy are given as follows: the only surviving tree-level term is $|F_Z|^2$ or $\lambda^2 h^4$. Due to the assumption (4.4), we will neglect (4.1) with respect to (4.3) with $A = H$ or H' . Finally we have to consider the contribution (4.5). We did not work out factors of $O(1)$ in detail, since these will not affect the subsequent discussion. Before investigating the potential it is convenient to introduce the following definitions:

$$m^2 = \alpha M_S^2, \quad \alpha = 6 N_\mu / M_p^2 M_1^2 \quad [\text{see (3.10)}]$$

$$p \approx \alpha^2 g_2^4 / (16 \pi^2)^2, \quad q \approx \alpha \lambda^2 \gamma^2 \alpha \beta / (16 \pi^2)^3 \quad (4.13)$$

Now the potential as a function of h is given by

$$V(h^2) = \lambda^2 h^4 + M_S^2 (p + q \ln(M_u^2 / \Lambda^2)) h^2 \quad (4.14)$$

At this point we have to emphasize that M_S and M_u depend on the as yet undetermined parameter ε [see (3.10)]:

$$\begin{aligned} M_S &= r \varepsilon^3 M_P, \quad r = M_1^2 / N M_P \sim O(1) \\ M_u &= s \varepsilon M_P, \quad s = M_2 / M_P \sim O(1) \end{aligned} \quad (4.15)$$

The potential (4.14) will not only have to be minimized with respect to h^2 , but also with respect to ε , which will be fixed this way. One sees immediately that $V(h^2)$ can only be negative for

$$p + q \ln(M_u^2 / \Lambda^2) < 0 \quad (4.16)$$

Hence, taking the ε dependence of M_u into account, it is easy to achieve $\varepsilon \ll 1$ at the minimum of V . Since we can approximately identify [see (3.10)]

$$\ln(M_u^2 / \Lambda^2) \approx \frac{1}{3} \ln(M_S^2 / M_P^2) \quad (4.17)$$

we can express the result of the minimization of V in terms of h^2 and the mass scale M_S rather than ε :

$$\begin{aligned} h^2 &= q M_S^2 / 6 \lambda^2 \\ \ln(M_S^2 / M_P^2) &= -\frac{3p}{q} - 1 \end{aligned} \quad (4.18)$$

For $q < p$ a hierarchy of mass scales $M_S \ll M_P$ is easily obtained. Let us recall the basic ingredients for arriving at this result: the dominant radiative contributions to the masses of the scalars H and H' are a finite positive definite one $\sim p M_S^2$, involving only gauge interactions, and a logarithmically divergent one involving only Yukawa interactions, which contributes a negative amount to the vacuum energy. The argument of the logarithm depends on the as yet undetermined ratio of particle masses to the Planck scale. This second contribution dominates the first one only if the absolute value of the logarithm is larger than some ratio of coupling constants [see (4.16) and (4.18)]: this way an exponential hierarchy of mass scales appears naturally.

A possible problem of the present model remains to be investigated: although in the presence of the VEVs (4.9) a positive definite mass for the scalar singlet Z arises from the tree-level terms $|F_H|^2$ and $|F_{H'}|^2$, these could be overwhelmed by radiative contributions [Z does not receive a positive definite mass from

radiative corrections of the type (4.3)] and the presently assumed vacuum with $\langle Z \rangle$ given by (4.12) could be unstable. With (4.4) the dominant contributions to the Z mass arise from a number of three-loop diagrams, where Z couples to H and H', which in turn have contact with the field S carrying the original supersymmetry breaking via gauge interactions. These contributions can be estimated to be of the order of

$$\sim \frac{\lambda^2 M_S^2 p}{16 \pi^2} \quad (4.19)$$

Inserting the value for h^2 from (4.18) into F_H and $F_{H'}$, in (4.6) the positive tree-level contribution to the mass² dominates (4.19) for

$$p/q \lesssim 16 \pi^2 / 3 \lambda^2 \quad (4.20)$$

One also has to worry about contributions to the potential which are linear in the scalar field Z. These arise from diagrams similar to the one in Fig. 4, but where two supersymmetry breaking masses (and one supersymmetric mass) are inserted into the S line. However, now the diagram is finite, and due to the large mass of the U field its contribution will be negligibly small. Hence, with (4.20), the potential will be stable in Z even without having to worry about the sign of the contributions (4.19).

A wide range of initial parameters will lead to an acceptable version of our model; let us discuss a particular choice numerically. For the dimensionful parameters (3.6) in our initial theory we assume

$$M_1 \sim 10 M_P, \quad N \sim M_P / 10, \quad \mu \sim M_P^3 / 60 \quad (4.21)$$

As a result the ratio a between the mass splitting among the scalar S fields, m^2 , and their common supersymmetric mass, M_S^2 [see (4.12)] becomes very small: $a \sim 10^{-4}$. For simplicity we assume a common value for the Yukawa couplings:

$$\alpha \sim \beta \sim \gamma \sim \lambda \sim .24 \quad (4.22)$$

Together with the known value of the SU(2) gauge coupling constant, this fixes the value of the parameters p and q in (4.13):

$$p \sim 7 \cdot 10^{-14} \quad q \sim 5 \cdot 10^{-15} \quad (4.23)$$

Their ratio determines the hierarchy of mass scales via (4.18)

$$\ln(M_S^2/M_P^2) \cong 42, \quad M_S/M_P \sim 6 \cdot 10^{-10} \quad (4.24)$$

and subsequently the VEV of $H = H'$:

$$h^2 \cong (120 \text{ GeV})^2 \quad (4.25)$$

Clearly the initial choice of parameters is natural although we had to arrange a to be small and the Yukawa couplings to be somewhat large in order to keep the hierarchy of mass scales from becoming too big. However, this way, previous assumptions as $m^2 \ll M_S^2$, (4.4) and (4.20) are also fulfilled. It is interesting to work out some further phenomenological details. For the mass M_S and the splitting m we find

$$M_S \sim 10^9 \text{ GeV}, \quad m \sim 10^7 \text{ GeV} \quad (4.26)$$

Hence the original supersymmetry breaking is quite large compared to the weak interaction scale, since it had to be fed through a number of loops. For the gravitino mass we obtain

$$m_{3/2} \sim 10^5 \text{ GeV} \quad (4.27)$$

which is comfortably above the bound obtained from cosmological considerations¹⁹⁾. The value of the VEV of the SU(2) triplet field S from (4.10) turns out to be small enough not to conflict with constraints originating from the known value of the ρ parameter. The mass M_U of the heavy fields U and R depends on the initial dimensionful parameter M_2 ; with $M_2 \sim M_P$ one finds

$$M_U \sim 10^{14-15} \text{ GeV} \quad (4.28)$$

Let us now have a look at the masses of the supersymmetric partners of the known gauge vector and matter fields. The masses of the SU(2) gauginos [see (4.2)] and scalar quarks and leptons with SU(2) quantum numbers [see (4.3)] are about the same $\sim 500 \text{ GeV}$. Most of the masses of the SU(3) and U(1) gauginos and the SU(2) singlet scalar quarks and leptons, however, are below the known experimental bounds. Hence the present model can only be called semirealistic. It is straightforward to see how this defect can and should be cured: the fields carrying the initial supersymmetry breaking F-type mass splitting (S in our case) should have SU(3) and U(1) quantum numbers as well. One might think of embedding

the SU(2) triplet field S into a 24 of a SU(5) GUT; let us postpone a discussion of this possibility until the next section.

Some more general properties of the radiative corrections in the present class of models should be discussed. Note that the radiative SU(2)×U(1) breaking is not driven by a large Yukawa coupling of a Higgs field to a (heavy) quark (e.g., the top quark), but that the coupling of the Higgs fields to a light singlet has been crucial. This coupling did not generate too large Higgs masses radiatively, since the F-type breaking is not as large as $m_{3/2}^{M_{\text{GUT}}}$, as it appears in Sugra-GUTs with canonical kinetic energies¹²⁾ (and since potentially dangerous diagrams involve at least three loops, see Fig. 4; in fact here they are highly welcome); furthermore, explicit soft supersymmetry breaking in the form of gaugino masses or positive definite scalar masses squared (corresponding to a dimensionless spurion field in superspace¹⁴⁾) do not exist. The finite gaugino and scalar masses generated radiatively do not trigger new divergences once they are inserted into higher loop diagrams, because they vanish rapidly at high momenta [see (4.2); the same qualitative q^2 -behaviour holds also for (4.3)].

Finally we want to comment on a contribution to the vacuum energy which has been neglected up to now. In general the one-loop corrections are given by¹⁶⁾:

$$(\Lambda^2/16\pi^2) \text{STr } M^2 + (1/64\pi^2) \text{STr } M^4 \ln(M^2/\Lambda^2) \quad (4.29)$$

Restricting ourselves to fields with spin ≤ 1 , the first term vanishes identically in the present class of models [which is not the case in Refs. 2)-4)]. After inserting the mass matrices of our effective theory (3.9) with (3.10) into (4.29), one discovers a term

$$\sim (m^4/64\pi^2) \ln(m^2/\Lambda^2) \quad (4.30)$$

which depends only on the scale m of the F-type mass splitting, not on any VEVs at first sight. However, m depends on the parameter ε , with respect to which we have to minimize the potential, and ε in turn depends on the VEV of the sliding field y before it was redefined to have a vanishing VEV [see (2.15)]. If we were to make the dependence of m on $\langle y \rangle$ explicit, we would have to take non-renormalizable interactions with powers of M_p in the denominator into account, which, however, would not be negligible since $\langle y \rangle \sim 0(M_p)$.

On the other hand, taking into account (4.30) for the vacuum energy can certainly not be the end of the story. The vacuum energy in the presence of the VEVs of H , H' and the other scalar fields which have been considered so far is roughly of $O(m^4)$ and negative. Although this is already a big step forward compared with vacuum energies of $O(M_P^4)$, which are generally present in supergravity theories without fine tuning, it is nowhere near small enough to be compatible with the known upper bound on the cosmological constant. Clearly, at the present stage of our knowledge, a vanishing cosmological constant after radiative corrections could be achieved only, if at all, with the help of unaesthetic assumptions. First we note that (4.30) will not be able to cancel the negative contribution mentioned before, since it is negative as well. Hence some positive contribution to the vacuum energy of $O(m^4)$ has to appear from somewhere else. A possible source could be gravitational quantum effects (note that $m \sim m_{3/2}$) which are beyond the scope of the present paper.

This speculation makes it clear that presently we do not know the precise form of the contributions of $O(m^4)$ which achieve a vanishing of the vacuum energy. Hence it is consistent to neglect (4.30); it is important to notice that we are talking about effects which are not larger than the ones we have considered before. It is easy to see that functions of m or $m_{3/2}$ (which, in turn, depend on ϵ) exist which achieve a vanishing of the vacuum energy after having been added to V in (4.14), but which do not affect the mechanism leading to $\epsilon \ll 1$ at the minimum of the potential.

5. - CONCLUSIONS

In this paper we have discussed a number of properties of minimal $N = 1$ supergravity theories. Apart from the sliding of all mass scales at tree level we have shown the possibility of supersymmetry breaking in the form of F-type mass splittings in the effective low energy theory, which makes these theories similar to the ones with global supersymmetry and an O'Raifeartaigh-sector.

Some important differences between the radiative corrections in these theories and supergravity theories with canonical kinetic energies for the matter have been pointed out: unless one introduces supersymmetry breaking gaugino masses, the quantum contributions to the vacuum energy are bounded naturally by m^4 , where the scale m of the F-type mass splitting is of the order of the

gravitino mass. Light singlets no longer lead to any dangerous effects; they are even necessary in order to fix $m \ll M_P$ naturally without introducing mass scales different from M_P by hand.

This mechanism has been discussed in detail in the context of a semi-realistic standard model. The natural appearance of hierarchies is connected with a team-work of finite and logarithmically divergent radiative corrections. Most of the calculations have only been performed up to factors of $O(1)$ and the initial parameters of the model are not very restricted by the requirement of arriving at the correct value of M_W ; however, a value for $m_{3/2}$ of $\sim 10^5$ GeV or more seems to appear naturally in these theories, thus avoiding possible cosmological problems.

The number of fields we had to introduce in the model in addition to the supersymmetric partners of the known particles might appear somewhat unaesthetic, and the mass splitting among SU(2) singlet superpartners turns out to be too small. However, an enlargement of the model to a grand unified theory seems to be very promising: due to their quantum numbers, most of the fields introduced in addition to the standard ones would fit nicely into adjoint representations of SU(5), which would be needed anyhow. After such an extension, the initial supersymmetry breaking F-type mass splitting would be present among fields carrying SU(3) and U(1) quantum numbers as well, and radiative corrections similar to the ones discussed before could easily lead to a phenomenologically acceptable low energy particle spectrum. Furthermore, a mass scale (called M_U) close to the GUT scale (up to a power of a coupling constant) has already appeared automatically. Finally the singlet Z, which had to be introduced for different reasons, already couples exactly the way it is required to do for the sliding singlet mechanism²⁰⁾ (leading to naturally light Higgs doublets) to work.

The reader might have observed that we did not perform the conventional "flat limit" $M_P \rightarrow \infty$, $m_{3/2}$ fixed. This limit is a priori possible in the form $M_P \rightarrow \infty$, $\varepsilon \rightarrow 0$ with $\varepsilon^3 M_P$ fixed. However, it would lead to $M_U \rightarrow \infty$ if we assume that all initial mass parameters are of $O(M_P)$. Clearly, if we start with one mass scale (M_P) only, we can only be left with one mass scale ($m_{3/2}$) after a corresponding limit. Hence, in order not to "lose" the mass scale M_U , a detailed investigation of the effective low energy theory in terms of powers of ε (with ε , M_P finite) had to be carried out.

The present model has been investigated in order to study the properties of minimal $N = 1$ supergravity theories, and its features seem to lead straight on to grand unification. Sugra GUTs with M_p as the only mass scale, which have already been proposed in Refs. 4) and 21), can probably be formulated in a very elegant way within the class of models presented here. It remains to hope that open problems such as gravitational quantum corrections probably connected with the embedding of $N = 1$ supergravity into extended supergravities can be solved in the near future. From this direction we expect to obtain restrictions on sensible $N = 1$ supergravity theories; hopefully the ones discussed here will be among the survivors.

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FIGURE CAPTIONS

- Fig. 1 Supergraph contributing to the masses of the scalar components of H and H' (after interchanging H and H' in the picture). Here and in the following the superpropagator of S is meant to include the explicit supersymmetry breaking²²⁾.
- Fig. 2 Supergraph leading to gaugino masses of the SU(2) gauge vector superfield, which is denoted by V_2 .
- Fig. 3 Supergraphs contributing to the masses of all scalar components A of the superfields ϕ transforming non-trivially under SU(2).
- Fig. 4 Supergraph leading to a logarithmic divergence proportional to F_Z . In contrast to the previous figures, the superpropagator of S is supersymmetric here, and we have shown explicitly the insertion of a supersymmetry breaking mass splitting of the F-type, denoted by a cross.

Fig. 1

Fig. 2

Fig. 3

Fig. 4