ON THE LOW ENERGY $d = 4$, $N = 1$ SUPERGRAVITY THEORY
EXTRACTED FROM THE $d = 10$, $N = 1$ SUPERSTRING^{*)}

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A B S T R A C T

We extract some general features of the $N = 1$ supergravity, $d = 4$ theory which could be obtained as a low energy limit of the recently proposed anomaly-free superstring theories. One finds that after the breaking of the residual $d = 4$ supersymmetry the gravitino and gaugino masses are equal. Soft scalar couplings are also obtained, including trilinear terms as well as common scalar masses $m \sim m_{3/2}$. It is argued that a possible source of supersymmetry breaking may be the condensation of fermions in the Yang-Mills sector and we discuss this possibility in the context of the $E_6 \times E_6$ model. Finally, we discuss some implications for the low energy supersymmetric spectrum, particularly squark and gluino masses.

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Recently ⁽¹⁾ it has been shown by Green and Schwarz that one may construct an anomaly free $N = 1$, $d = 10$ supergravity theory coupled to Yang-Mills matter ^(2,3) as a limit of a $d = 10$, $N = 1$ superstring theory with gauge groups $SO(32)$ and $E_8 \times E_8$. This offers a possible unification of all known interactions incorporating chiral fermions in the low energy sector which could be identified with quarks and leptons ⁽²⁾. Candelas et al. have also shown recently ⁽⁴⁾ that if one further insists on the conservation of an unbroken supersymmetry in $d = 4$ after compactification of extra dimensions, then the internal 6-dimensional space must be a Ricci-flat Kähler manifold with $SU(3)$ holonomy group. The condition of the existence of an unbroken $d = 4$ supersymmetry strongly relates the values of $\langle F^{MN} \rangle$ and $\langle R_{MN}^{AB} \rangle$ in the extra dimensions in such a way that one is apparently forced to identify an $SU(3)$ subgroup inside $SO(32)$ or $E_8 \times E_8$ with the $SU(3)$ connection in the compact manifold. Identifying this $SU(3)$ with the one in the maximal E_8 subgroup $E_6 \times SU(3)$ one is led to a $N = 1$ supersymmetric E_6 theory with a number of families given by one half the Euler characteristic of the internal manifold. In this case the effective low energy theory would be a $N = 1$ supergravity theory in $d = 4$ coupled to an E_6 GUT.

It is an appealing fact that one may consistently conserve one supersymmetry in the process of compactification since it can be of help in solving the "gauge hierarchy" problem of GUT's. In particular, it has been shown in the last few years that spontaneously broken $N = 1$ ($d = 4$) supergravity coupled to GUT's offers an aesthetic scheme for facing that problem ⁽⁵⁾. In these theories the

Glashow-Salam-Weinberg $SU(2) \times U(1)$ symmetry is radiatively broken as a natural consequence of the breaking of supersymmetry. If it is true that the effective Lagrangian left over after compactification of the $d = 10$ superstring is an $N = 1$ supergravity Lagrangian then a natural question to ask is, how is this effective theory? Is it consistent with the low energy phenomenology?

Of course, a complete answer to these questions would require a detailed knowledge of the process of compactification, which seems a difficult task given the complexity of the 6-dimensional compact manifolds involved ⁽⁴⁾. Furthermore, there could also be purely "stringy" ways of compactification which could complicate the problem more. Our approach here will be less ambitious. We will assume that a good guess of the effective low energy theory after compactification may be obtained from examining the Lagrangian of $N = 1$ supergravity coupled to matter in 10 dimensions. Of course, the superstring theory gives rise to extra terms ^(1,2,4) (most notably, the Wess-Zumino terms cancelling the anomalies as well as some term ⁽⁴⁾ $R_{MNPQ} R^{MNPQ}$ to avoid no-go theorems for compactification ⁽⁶⁾) but we will concentrate on the Chapline-Manton action with the appropriate modified H_{MNP} three-form incorporating the Lorentz Chern-Simons symbol ^(7,8). We do not restrict ourselves to any specific compactification scheme but just assume that an unbroken supersymmetry in $d = 4$ is preserved. Most of our discussion applies to both the $SO(32)$ and $E_8 \times E_8$ gauge groups.

Let us briefly recall the field content of the $N = 1$, $d = 10$ theory ^(7,8). It contains in the supergravity multiplet the fields

e_M^A (Zehnbein), ψ_M (Rarita-Schwinger-Majorana-Weyl spinor), λ (right-handed Majorana-Weyl spinor), B_{MN} (antisymmetric two-form potential) and φ (scalar, "dilaton"). This multiplet may be coupled to a $N = 1, d = 10$ Yang-Mills matter multiplet consisting of A_M (gauge field in the adjoint of an arbitrary group G) and χ (left-handed Majorana-Weyl spinor, also in the adjoint). Cancellation of anomalies may be obtained for $G = SO(32)$ or $E_8 \times E_8$ ⁽¹⁾. The Lagrangian density up to four-fermion interactions is :

$$\mathcal{L} = \mathcal{L}_B + \mathcal{L}_{FK} + \mathcal{L}_{FB} \quad (1a)$$

$$\mathcal{L}_B = -\frac{1}{2k^2} e R - \frac{e}{(k\varphi)^2} (\partial_M \varphi)^2 - \frac{e}{4(g^2\varphi)} F_a^{MN} F_{MN}^a - \frac{3ek^2}{2(g^2\varphi)^2} H^{MNP} H_{MNP} \quad (1b)$$

$$\mathcal{L}_{FK} = -\frac{e}{2} \bar{\psi}_M \Gamma^{MNP} D_N(\omega) \psi_P - \frac{e}{2} \bar{\lambda} \Gamma^M D_M(\omega) \lambda - \frac{e}{2} \text{tr}(\bar{\chi} \not{D}(\omega) \chi) - \frac{e}{12} \bar{\psi}_M \frac{(\not{\partial}\varphi)}{\varphi} \Gamma^M \lambda \quad (1c)$$

$$\mathcal{L}_{FB} = -\frac{2ke}{8(g^2\varphi)^{1/2}} \text{Tr}(\bar{\chi} \Gamma^M \Gamma^{NP} F_{NP})(\psi_M + \frac{\sqrt{2}}{12} \Gamma_M \lambda) + \frac{k^2 e}{8(g^2\varphi)} H_{MNP} (\bar{\psi}_Q \Gamma^{QMNP} \psi_R + 6\bar{\psi}^M \Gamma^N \psi^P - \sqrt{2} \bar{\psi}_Q \Gamma^{MNP} \Gamma^Q \lambda + \text{tr}(\bar{\chi} \Gamma^{MNP} \chi)) \quad (1d)$$

where the Γ 's are generalized antisymmetric gamma matrices and the block latin indices $M, N, P, \dots$ run from 1 to 10. The Yang-Mills strenghts are defined as usual :

$$\frac{1}{2} F_{MN}^a = \partial_{[M} A_{N]}^a + f^{abc} A_M^b A_N^c \quad (2)$$

and the H three-form is defined in differential form notation :

$$H = dB - \omega_{3Y} + \omega_{3L} \quad (3)$$

with B being the two index field in the supergravity multiplet and ω_{3Y} and ω_{3L} Yang-Mills and Lorentz Chern-Simons symbols :

$$\omega_{3Y} = \frac{1}{30} \text{Tr} (A \wedge F - \frac{1}{3} A \wedge A \wedge A) \quad (4)$$

$$\omega_{3L} = \text{tr} (\omega \wedge R - \frac{1}{3} \omega \wedge \omega \wedge \omega) \quad (5)$$

with Tr meaning trace in the adjoint representation and ω is the Lorentz connection. Notice that the only field with non-polynomial interactions (apart from pure gravity) is the scalar ϕ . This is easily understood if one obtains the $d = 10$ supergravity Lagrangian upon dimensional reduction of the $d = 11$ Lagrangian ⁽⁷⁾ since then ϕ is contained in the elf-bein and in $d = 10$ still conserves its non-polynomial couplings. Notice also that g^2 always appears in the combination $(g^2 \phi)$ and may be absorbed ⁽⁹⁾ in the definition of ϕ and thus it is not a free parameter of the model ^{*)}.

Candelas et al. have studied under what conditions one is left

^{*)} The canonical dimensions of the fields in $d = 10$ are :

$$\begin{aligned} d[R] &= 2; \quad d[g^2 \phi] = -6; \quad d[F] = 2; \quad d[K] = -4; \quad d[H] = 3; \quad d[B] = 2; \\ d[A] &= 1; \quad d[\psi] = d[\chi] = d[\lambda] = \frac{9}{2}. \end{aligned}$$

with an unbroken $d = 4$ supergravity after compactification of the six extra dimensions. They find that the internal manifold K must be Kähler and Ricci flat and have $SU(3)$ holonomy group ⁽⁴⁾. Furthermore it must be

$$H_{mnp} = 0 \quad (6a)$$

$$\Gamma^{mn} F_{mn}^a \eta = 0 \quad (6b)$$

where small latin letters refer to internal dimensions and η is a covariantly constant spinor ($\nabla_m \eta = 0$). Recalling the definitions (3), (4) and (5) one finds that $H_{mnp} = 0$ requires as an integrability condition

$$\frac{1}{30} \text{Tr}(F \wedge F) = \text{tr}(R \wedge R) \quad (7)$$

which is a strong restriction on compactifications ⁽⁴⁾, an obvious solution being to identify some $SU(3)$ gauge fields with the $SU(3)$ connections ω .

In order to obtain a realistic theory at low energies we eventually have to break slightly the remnant $d = 4$ supersymmetry. Let us then look to our Lagrangian in Eq. (1), assume that somehow the 6 internal dimensions are compactified and see how one could give a mass to the four-dimensional gravitino. Examining the second term in Eq. (1d) one notices that a v.e.v. for H_{mnp} in internal dimensions would give a mass to the four dimensional gravitino and in general would break supersymmetry. Notice that the required

mixing of the gravitino with would be goldstinos (λ or some χ if $F_{mn}\Gamma^{mn} \neq 0$) is also present in $\mathcal{L}_{FB}$. Thus all the ingredients for a 4-dimensional super-Higgs effect are present giving rise to a gravitino mass ^{*)}

$$m_{3/2} = \frac{k^2}{4(g^2\varphi)} \langle H_{mnp} \Gamma^{mnp} \rangle \quad (8)$$

Moreover the last term of Eq. (1d) shows us a remarkable fact. Once supersymmetry is broken the four-dimensional χ fields in real reps. acquire also masses equal to $m_{3/2}$. Since the gauginos should be part of the χ zero modes one expects masses M for them

$$M = m_{3/2} \quad (9)$$

If there are further real rep. fermionic modes (e.g. like higgsinos $H_1 + H_2$) they may also obtain similar masses $\mu_0 = m_{3/2}$ but this is more model dependent because the Higgsinos may be both contained in a complex rep. (e.g. a 27 of E_6) in which case no such Higgsino mass should appear. Furthermore they could obtain some extra contribution from the third term of Eq. (1c) (this is not possible for gauginos). Recall also that in order to preserve the softness of the global susy breaking this possible Higgsino mass should come along with an equal mass for their bosonic partners.

^{*)} The canonical dimensions of the fields in $d = 4$ are as in $d = 10$ except for $d[g^2\varphi] = 0$; $d[K] = -1$; $d[\psi] = d[\lambda] = d[\chi] = \frac{3}{2}$.

Notice that Eq. (9) implies that if we want to use the unbroken supersymmetry to solve the gauge hierarchy problem one should have $M \approx M_W \approx 10^{12} \text{ GeV}$ which implies :

$$\langle H \rangle \sim M_W M_P^2 \sim (10^{13} \text{ GeV})^3 \quad (10)$$

Let us also mention that in order to break the $d = 4$ supersymmetry we have seen that one needs $\langle H \rangle \neq 0$, $\langle F_{mn} F^{mn} \rangle \neq 0$ in agreement with the results of Candelas et al.

The effects of supersymmetry breaking in the scalar sector are not so transparent. The scalars corresponding to the "observable" sector will in general be originated from the A_m^a fields. From the $\mathcal{L}_B$ Lagrangian one notices that their kinetic terms will be originated in the $F_a^{\mu m} F_{\mu m}^a$ terms and that in general they will have non-minimal kinetic couplings of the general form $(f(z, z^*) \delta_b^a) \cdot (\partial^\mu A_a) \cdot (\partial_\mu A^b)$ where z are singlet fields from the supergravity sector. The usual supersymmetric scalar couplings will in general come from the $F_a^{mn} F_{mn}^a$ terms. More scalar couplings will be contained in the $H^{mnp} H_{mnp}$ terms. Thus for example after supersymmetry breaking through an $\langle H_{mnp} \rangle \neq 0$ one has a term

$$- \frac{3k^2 e}{(g^2 \varphi)^2} \langle H^{mnp} \rangle \left(\partial_{[m} B_{np]} - \frac{1}{30} \text{Tr} (A_{[m} F_{np]} - \frac{1}{3} A_{[m} A_n A_p]) \right) \quad (11)$$

Since $\langle H \rangle \sim m_{3/2} \frac{(g^2 \varphi)}{K^2}$ one gets e.g. trilinear scalar terms in four dimensions $\sim m_{3/2} (\text{scalar})^3$ as expected in a spontaneously broken $N = 1$, $d = 4$ supergravity theory ⁽⁵⁾. In fact the term $\text{tr}(A)^3$ is connected to the "superpotential" W in $d = 4$ and

e.g. the term $H_{mnp} H^{mnp}$ will contain a piece of the form $|W|^2$ as expected also in a $N = 1, d = 4$ theory. Going on with the analogy, one also expects soft masses for the scalars from the terms $\sim H^{mnp} A^a_{[m} F_{np]a}$, in analogy with the couplings $W\phi^a W_a$ one has in the $N = 1, d = 4$ supersymmetric scalar potential (10). However we must remark that one does not know the precise connection between the physical scalar zero modes and the scalar fields appearing in the $d = 10$ theory since that depends on the details of the compactification. Thus this means that we cannot make definite statements on things like e.g. the value of the "A" parameter (11) of low energy supergravity models. As we stated above, the "observable" scalars seem to have in general non-minimal kinetic terms of the form $f(z, z^*) (\partial^\mu A^a)^2$, where the z 's are singlets from the supergravity "hidden sector". The function $f(z, z^*)$ will in general appear as a consequence of rescaling after compactification. Under these conditions one expects that the breaking of supergravity will induce common masses for all the scalars

$$m = m_{3/2} (1 + \delta) \quad (12)$$

where δ will be a constant depending on f derivatives (12).

In order to explore the consequences of supersymmetry breaking we have assumed that somehow H_{mnp} (and possibly $F_{mn} \Gamma^{mn}$) acquire a v.e.v. but have not specified how that could occur. We now discuss a possible dynamical way in which it could take place. To be more specific we will concentrate on the $E_8 \times E'_8$ case. In this

model the extra E'_8 remains unbroken in $d = 4$ and the particles coupling to it (χ' and A') interact only gravitationally with the "observable" fields. If this E'_8 is asymptotically free its coupling will get strong at a "confining" scale Λ . Now looking to Eq. (1d) one finds a term

$$\frac{k^2}{8(g^2\varphi)} H_{mnp} \text{tr}(\bar{\chi}' \Gamma^{mnp} \chi') \quad (13)$$

If a condensate $\langle \bar{\chi}' \chi' \rangle \simeq \Lambda^3$ appears in four dimensions (13) then upon application of the equations of motion, H will in general develop a v.e.v.. In the basis for the extra dimensions introduced in Candelas et al. one expects :

$$\langle H_{ijk} \rangle = \frac{g^2\varphi}{24} \Lambda^3 \epsilon_{ijk} \quad (14)$$

where the i, j, k indices transform under the $SU(3)$ holonomy group. Then one obtains for the gravitino mass :

$$m_{3/2} \sim \frac{k^2 \Lambda^3}{16} \quad (15)$$

which implies $\Lambda \sim 5 \times 10^{13} \text{ GeV}$ if $m_{3/2} \sim M_W$. The actual confining scale of E'_8 will depend on the low energy E'_8 non-singlet content of the theory as well as on the value of g'_8 at the Planck scale. In one possible model (4) one only has a vector E'_8 multiplet in the low energy sector. In this case one can calculate that the β -function coefficient is $b = -90$ and in order to get $\Lambda \sim 10^{13} \text{ GeV}$ one needs $\alpha'_8 \sim 1/180$ ($g'_8 \sim 0.26$) for the $d = 4$ coupling. In the ten-dimensional theory one expects $\alpha'_8 = \alpha_8$ but

in general this does not necessarily imply that $\alpha_6 = \alpha'_8$ (in obvious notation) for the $d = 4$ theory since a rescaling of fields and couplings will in general be required after compactification. Notice that if the 4-dimensional equality $\alpha_6 = \alpha'_8$ was correct this would imply $\alpha'_8 \geq 1/24$ since one expects for an E_6 GUT $\alpha_6 \geq 1/24$ (in general $\sim 1/4 - 1/8$, see Ref. (14)). This would lead to a confining scale Λ too close to the Planck mass. Observe also, that the compactification scale need not coincide with the Planck scale but could be lower.

Let us comment now on the low energy effective theory below the Planck mass and the low energy susy spectrum. At the Planck mass we will have the boundary conditions $m \sim M = m_{3/2}$. Let us assume for a moment that $\delta = 0$ in Eq. (12) and $m = m_{3/2}$, which could occur if the z degrees of freedom decouple after compactification. For a wide range of parameters one may get radiatively the breaking of the $SU(2) \times U(1)$ symmetry. One may calculate the evolution of squark ($\tilde{q}$), slepton ($\tilde{\ell}$) and gluino ($\tilde{g}$) masses from their renormalization group equations. In the case of minimal low energy content (three families + 2 Higgses) one obtains the numerical results (15,16)

$$m_{\tilde{q}}^2(M_w) = m^2 + 7.4 M^2 = 8.4 M^2 \quad (16a)$$

$$m_{\tilde{\ell}_R}^2(M_w) = m^2 + 0.15 M^2 = 1.15 M^2 \quad (16b)$$

$$M_{\tilde{g}}(M_w) = 3 M \quad (16c)$$

so that one has $m_{\tilde{q}} \approx 0.97 M_{\tilde{g}}$ and squarks and gluinos will be practically degenerate. If we further impose the experimental constraint $m_{\tilde{\ell}_R} \geq 20\text{GeV}$ one gets $M \geq 18.5\text{GeV}$ implying

$$m_{\tilde{q}} \gtrsim 54\text{GeV} ; \quad M_{\tilde{g}} \gtrsim 56\text{GeV} \quad (17)$$

which may still be phenomenologically interesting at present collider energies. A recent analysis ⁽¹⁶⁾ also shows that in the minimal case of three generations and $m_t \approx 40\text{GeV}$ one has (for $B = A - 1$) that only values $A \approx 3$ or $-3 \lesssim A \lesssim -1$ are allowed by the appropriate breaking of the $SU(2) \times U(1)$ symmetry. Also the bare Higgsino mass is bounded above by $\mu_0 \lesssim 1.5 m_{3/2}$ and for the wide region $0.3 m_{3/2} \lesssim \mu_0 \lesssim 1.5 m_{3/2}$ the values of A are confined to be $A \sim 3$ or $A \sim -1$. Thus this shows us how the specific boundary conditions at the Planck scale very strongly affect the low energy predictions. In the case of four light families (as is the case in a specific $E_8 \times E_8'$ example) one gets $m_{\tilde{q}} \approx 2.5M$ and $M_{\tilde{g}} \approx 1.6M$ so that in this case the gluinos are always lighter than the squarks ($m_{\tilde{q}} \approx 1.56 M_{\tilde{g}}$). Requiring $m_{\tilde{e}_R} \geq 20\text{GeV}$ one gets that in this case one should have $M_{\tilde{g}} \gtrsim 30\text{GeV}$ and $m_{\tilde{q}} \gtrsim 46\text{GeV}$. Let us however notice that in both the 3 and 4 family cases the unification mass is $\sim 3 \times 10^{16}\text{GeV}$ instead of the desired value $\sim M_p / \sqrt{8\pi} \sim 2.4 \times 10^{18}\text{GeV}$. A third possibility for low energy content is the existence of three complete E_6 families (3 27's) with all the particles light. One must assume in this case that the Higgs triplet couplings to quarks and leptons are hierarchically suppressed in order not to give too fast proton decay. With this low energy content one finds $M_x \sim 10^{18}\text{GeV}$ as de-

sired as well as $\sin^2 \theta_w \approx 0.20$ at one loop. One gets $m_{\tilde{q}} \approx 2.3M$ and $M_{\tilde{g}} \approx M$ ($\beta_3 = 0$) so that $m_{\tilde{q}} \approx 2.3M_{\tilde{g}}$ so that also in this case the gluinos are lighter than squarks. The constraint $m_{\tilde{e}_R} \geq 20\text{GeV}$ implies $M_{\tilde{g}} \geq 19\text{GeV}$ and $m_{\tilde{q}} \geq 42\text{GeV}$. A general conclusion of all this is that for the minimal three family content (and assuming $m = m_{3/2}$) one has almost degenerate squark and gluinos ($m_{\tilde{q}}$ slightly smaller than $M_{\tilde{g}}$) but if one adds some extra $SU(3)$ non-singlet object which is light then the gluinos will be always lighter than squarks. However very light gluinos do not seem possible in any of the schemes.

Although it seems possible that a realistic $N = 1$ supergravity model may be obtainable from a superstring theory several difficult problems remain unsolved. One of them, as we mentioned above, is the problem of the quark-lepton spectrum. Naively one expects in this theories to have all the Yukawa couplings of order g and hence it is not clear from where the observable hierarchy of quark and lepton masses could come. Maybe the process of compactification may imply more complicated manifolds. A second problem is the old question of the doublet-triplet splitting of the Weinberg-Salam doublets and their $SU(3)_C$ triplet mates. In a general SUSY-GUT one has freedom to introduce fields in the theory as well as ad hoc couplings in order to give large masses to the unwanted triplets. We do not have this freedom in this theory and hence we have to believe in the existence of a compactification scheme in which the dangerous triplets do not appear as possible zero modes of the Dirac operator. This certainly would be a miracle. Another possibility is that this Higgs triplets may be as light as the

usual Weinberg-Salam doublets but they do not mediate fast proton decay because they couple extremely weakly to quarks and leptons due to some symmetry. This symmetry cannot commute with anything embedding $SU(3) \times SU(2) \times U(1)$ (like $SU(5)$, E_6 etc.) since we need some quantum number distinguishing between colour triplets and $W - S$ doublets. Further problems may arise in specific models (e.g. the E_6 theory coming from the $E_8 \times E_8'$ case) in order to get rid of unwanted particles like e.g. the 12 unseen states in each E_6 family ⁽¹⁷⁾. This includes some right-handed neutrinos as well as $\underline{5} + \bar{\underline{5}}$'s under the $SU(5)$ subgroup. However the theory seems to be rich enough so that one may get rid of all unwanted states ⁽¹⁴⁾. Another potential problem for these theories comes from cosmological bounds ⁽¹⁷⁾ on the gravitino mass which imply $m_{3/2} \gtrsim 10^4 \text{ GeV}$. Since in the present scheme $M = m_{3/2}$ the susy-gap would be too large for obtaining natural $SU(2) \times U(1)$ breaking at the M_W scale. Not with standing all these problems it seems really amazing if it is true that one can get all this low energy information from such intricate theories as superstrings are.

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