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SCALES IN SUPERSTRING MODELS

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A B S T R A C T

We argue that supersymmetry breaking in the observable low energy sector of superstring models is triggered by gaugino masses $m_{\frac{1}{2}}$ which are proportional to the product of the gravitino mass and a non-vanishing F-term. The scales of global supersymmetry breaking $m_{\frac{1}{2}}$, weak gauge symmetry breaking m_W , grand unification m_X , compactification m_C and the gravitino mass $m_{3/2}$ are then determined by radiative corrections in a variant of the no-scale models proposed previously: $m_{\frac{1}{2}}, m_W \propto m_P \exp(-O(1)/\alpha_t)$, $m_X, m_C, m_{3/2} \propto m_P$. Key ingredients are an additional $SU(3)_C \times SU(2)_L \times U(1)_Y$ singlet field and an extra low energy $U(1)$ gauge symmetry expected in superstring models.

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Superstring models¹⁾ offer enticing solutions to the problems of quantum gravity²⁾, of unifying all the fundamental particle interactions³⁾, of flavour^{4),5)}, and of the cosmological constant⁶⁾. Of course, one or two minor technical details remain to be worked out, and some problems remain. In particular, how is the hierarchy of different physical mass scales generated and maintained? These scales include the scales of weak gauge symmetry breaking m_W , of global supersymmetry breaking $m_{0, \frac{1}{2}}$, of local supersymmetry breaking $m_{3/2}$, of grand unification m_X , of compactification m_C , and the Planck mass m_P . The raw material for understanding these scales is provided by the set of zero-mass particles in the superstring theory, namely the particles of $d = 10$, $N = 1$ chiral supergravity⁷⁾ left over when one takes the field theory or zero-slope limit. Among these is a scalar dilaton field ϕ whose a priori undetermined vacuum expectation value is related to that of the gauge coupling in ten dimensions. Compactification⁴⁾ of the ten-dimensional field theory to a four-dimensional supergravity theory introduces a new scalar field σ : $m_C = e^{-\sigma/2} m_P$ whose value is also undetermined at the tree level⁸⁾. When one compactifies maintaining a low energy $N = 1$ supersymmetry by using a Calabi-Yau manifold, one obtains a zero cosmological constant in four dimensions⁴⁾ and hence a degenerate vacuum state for different values of ϕ , σ and the observable gauge non-singlet field ϕ^1 . The resulting low energy theory has a no-scale⁹⁾⁻¹²⁾ $SU(1,1)/U(1) \times (SU(n,1)/SU(n) \times U(1))$ structure based on the Kähler potential⁸⁾

$$G = -\ln(S+S^+) - 3\ln(T+T^+ - 2\phi^i \phi_i^+) + \ln|F|^2 \quad (1)$$

where $\text{Re}S \equiv S_R = e^{3\sigma} \phi^{-3/4}$, $\text{Re}T \equiv T_R = e^{\sigma} \phi^{3/4} + \phi_1^+ \phi_1^1$ and F is the superpotential, and a non-trivial chiral kinetic function f for the gauge fields W^α

$$f_{\alpha\beta} W^\alpha W^\beta : f_{\alpha\beta} = \delta_{\alpha\beta} S \quad (2)$$

Such no-scale supergravity models have been advocated¹⁰⁾⁻¹²⁾ in the past for the purposes of generating a hierarchy $m_{W, 0, \frac{1}{2}} \ll m_P$ with $m_{3/2} = 0 (m_W^2/m_P \text{ to } m_P)$ through non-gravitational radiative corrections. The problem of scales in superstring theories becomes that of determining S_R and T_R dynamically in some variant of this no-scale scenario. Non-perturbative gauge dynamics can¹³⁾ generate a non-trivial S -dependence in the superpotential $F = C + h \exp(-3S/2b_0)$ and hence provide a mechanism¹⁴⁾ for determining S and thus the four-dimensional gauge coupling constant $g^2 = 1/S_R$. It therefore remains to determine T_R by radiative corrections, for which we need a seed of global supersymmetry breaking in the observable sector. Certainly the gaugino

condensation¹³⁾⁻¹⁵⁾ which fixes S_R permits local supersymmetry breaking with a non-zero gravitino mass

$$m_{3/2} = e^{G/2} = \left(\frac{1}{16 S_R T_R^3} \right)^{1/2} |W| = \frac{\Theta(c)}{S_R^{1/2} T_R^{3/2}} \quad (3)$$

in natural Planck units: $m_P/\sqrt{8\pi} = 1$, but how is this fed through to the observable sector?

We argue in this paper that the dominant source of global supersymmetry breaking in the observable sector must be the gaugino mass $m_{1/2}$, with non-zero scalar masses m_0 for the gauge non-singlet fields ϕ^i in (1) being generated by radiative corrections¹²⁾. A non-zero gaugino mass arises only when one takes into account (i) renormalization of the chiral kinetic function f (2) by loops in the observable sector, and (ii) non-zero F-terms $F_i \equiv \partial F / \partial \phi^i$;

$$m_{1/2} = \frac{m_{3/2}}{4 \text{Re} f} G^i (G^{-1})^j_i \frac{\partial f}{\partial \phi^j} \quad (4a)$$

$$= - \frac{b g^2}{4} \left(\frac{1}{16 S_R T_R^3} \right)^{1/2} \frac{F_j^+ U_j}{U} \quad (4b)$$

$$= \Theta\left(\frac{\alpha}{4\pi}\right) \left(\frac{1}{16 S_R T_R^3} \right)^{1/2} \langle 0 | H | 0 \rangle \quad (4c)$$

where $G_i \equiv \partial G / \partial \phi^i$, etc., and

$$f_{\alpha\beta} = \delta_{\alpha\beta} (S + b \ln U(\phi^i)) \quad (5)$$

with b related to the observable sector gauge β -function and U a gauge singlet product of observable chiral fields, and $\langle 0 | H | 0 \rangle$ is a generic Higgs vacuum expectation value. The F-term can be non-zero if one augments the minimal supersymmetric standard model chiral fields by adding an $SU(3) \times SU(2) \times U(1)$ gauge singlet N which in fact appears naturally in superstring models where E_6 is broken down to $SU(3) \times SU(2) \times U(1)^2$ as a low energy group¹⁶⁾:

$$F \supset \lambda H \bar{H} N \quad (6)$$

where H and $\bar{H}$ are the usual two light Higgs doublets. This introduction of the

superpotential term (6) also solves the Higgs doublet mixing problem¹⁷⁾. Radiative corrections dominated by the top quark Yukawa coupling $\alpha_t \equiv \lambda_t^2/4\pi$ can then generate weak gauge symmetry breaking at a scale given by dimensional transmutation¹⁸⁾:

$$\langle 0|H|0\rangle, \langle 0|N|0\rangle = \mathcal{O}(\mu_0): \mu_0 \sim \mu_C \exp\left(-\frac{\mathcal{O}(1)}{\alpha_t}\right) \quad (7)$$

with the global supersymmetry breaking scale $m_{\frac{1}{2}}$ also $\mathcal{O}(\mu_0)$ ¹⁰⁾. We see from Eq. (4) that T_R is therefore determined dynamically to have a value which is not hierarchically different from unity¹⁹⁾:

$$T_R = \mathcal{O}(1) \quad (8)$$

This together with $S_R = 1/g^2$ means that the dilaton field ϕ , the compactification scale $m_C = e^{-\sigma/2} m_P$ and the grand unification scale $m_X = \mathcal{O}(m_C^4/m_P^3)$ are all of order the Planck scale, thereby avoiding the problems²⁰⁾ of strong and/or weak coupling which would have arisen if $T_R \gg 1$ as in some no-scale models¹⁰⁾⁻¹²⁾. The result (8) also means that the gravitino mass $m_{3/2}$ (3) is determined dynamically to be close to m_P ¹⁹⁾.

Why do we say that the seed of global supersymmetry breaking must be $m_{\frac{1}{2}}$? The no-scale mechanism¹⁰⁾ works if and only if the physical gauge singlet "Polonyi" field associated with variations in the scale of global supersymmetry breaking has a mass $\mathcal{O}(m_W^2/m_P)$. The $SU(n,1)/SU(n) \times U(1)$ Kähler manifold structure (1) describes the Goldstone bosons of the spontaneous breakdown of a global $SU(n,1)$ to an $SU(n) \times U(1)$ maximal compact subgroup, so that the Polonyi and gauge non-singlet scalar fields have zero mass¹²⁾. We presume that gravitational interactions respect order by order both the linearly and non-linearly realized symmetries, and therefore will not change this Kähler structure²¹⁾. This assumption is supported by the knowledge that counter terms for extended supergravities respect their non-compact global symmetries wherever this can be checked²²⁾. It may well be that the structure of gravitational loops in the superstring theory is not correctly reproduced by calculations in the effective $N = 1$ supergravity field theory. The Kähler structure (1) can only be broken by gauge and Yukawa couplings, which preserve a residual $SU(1,1)/U(1)$ for the Polonyi field while giving the gauge non-singlet scalars masses $\mathcal{O}(m_W)$ as soon as

global supersymmetry is broken elsewhere in the theory. Consider for example the effects of gravitational loops involving a massive gravitino. In one loop they produce quadratically divergent terms

$$\delta V = \frac{1}{32\pi^2} \text{Str } M^2 \Lambda^2 = -\frac{1}{8\pi^2} e^G \Lambda^2 \quad (9)$$

which are consistent with the assumed $SU(n,1)/SU(n) \times U(1)$ symmetry, and logarithmically divergent terms

$$\delta V = \frac{1}{64\pi^2} \text{Str } M^4 \ln M^2/\Lambda^2 = -\frac{1}{16\pi^2} e^{2G} \ln e^G/\Lambda^2 \quad (10)$$

which are not consistent, but may be cancelled in the superstring theory as we now assume.

Non-gravitational loop corrections can only generate non-zero scalar masses m_0 if $m_{\frac{1}{2}} \neq 0$, so we now see how this can arise. Equation (4a) tells us that we need $G^i \neq 0$ for some chiral field, and also $\partial f/\partial \phi^j \neq 0$. The Kähler potential (1) contains two sectors, one involving S alone and the other (T, ϕ^i) . Clearly $\partial f/\partial S \neq 0$ from Eq. (2), while gaugino condensation¹³⁾⁻¹⁵⁾ yields $G_S = 0$ as a first approximation¹⁴⁾ since

$$V = e^G \left(\frac{G_S G_{S^+}}{G_{SS^+}} + \dots \right) \quad (11)$$

Observable sector corrections to V (11) are $O(m_W^4/g^2)$ which alter the minimization condition for S so that

$$G_S = O\left(m_W^4/|c|\right) \quad (12)$$

so that this sector can only make a negligible contribution (4a) to $m_{\frac{1}{2}}$. Therefore we turn to the (T, ϕ^i) sector.

Clearly, $G_T = -3/(T+T^+-2\phi^i\phi_i^+) \neq 0$ in general, but f must be independent of T if the problem of the cosmological constant is to be solved. The imaginary part of T will then be the massless Goldstone boson advocated in Ref. 6). Actually, we know f is indeed⁸⁾ independent of T at the tree level (2), and radiative corrections are incapable of generating a T dependence, since T is a gauge singlet. The matrix $(G^{-1})_i^j$ contains off-diagonal elements, but as pointed out in Ref. 12, the combination $G^i(G^{-1})_i^j$ is a projector on the T -direction - in

the absence of a superpotential F . In its presence

$$G^i (G^{-1})^j_i = - \frac{\partial F^+ / \partial \phi^j}{F} \quad (13)$$

and we can obtain a non-zero gaugino mass (4a) iff $\partial f / \partial \phi^j \neq 0$.

The chiral function f is independent⁸⁾ of the ϕ^i at the tree level (2), but its real part is always related to g^2 : $\text{Re} f = 1/g^2$, so one might expect it to be subject to renormalization. In the spirit of the Coleman-Weinberg effective potential²³⁾, we replace the conventional renormalization scale μ by the magnitude of the field variables $|\phi^i|$. In leading order then

$$\text{Re} f = \frac{1}{g^2} \simeq S_R + b \sum_i \ln |\phi^i|^2 \quad (14)$$

where b is related to the conventional gauge β -function. To reconstruct from this a possible leading order form for the chiral function f (2), we note that for any product $U(\phi^i) = \prod_i \phi^i$

$$\text{Re} \ln U = \ln U(\phi^i) + \ln U^+(\phi^i) = \ln \prod_i |\phi^i|^2 = \sum_i \ln |\phi^i|^2 \quad (15)$$

to leading order. Therefore we make the ansatz (5), which, when substituted into Eq. (4a) together with the modified projector (13), gives us the expression (4b) for $m_{\frac{1}{2}}$.

For this expression to be non-zero, we need some F_i to be non-zero. However, all the $F_i = 0$ in the minimal supersymmetric standard model, and to get $F_i \neq 0$ one should add a trilinear term containing at least two fields with non-zero vacuum expectation values. Happily, superstring theories provide an observable sector low energy gauge group which is at least as large as $SU(3) \times SU(2) \times U(1)^2$, and an extra $SU(3) \times SU(2) \times U(1)_Y$ singlet field with non-zero $U(1)_E$ hypercharge for which (6) is the only additional superpotential coupling allowed¹⁶⁾. Under $SU(5) \times U(1)_E \times U(1)_F$ the $\underline{27}$ of E_6 decomposes to

$$\begin{aligned} \underline{27} = & (1, 1, 0) + (\underline{5}, -\frac{1}{2}, \frac{1}{2}) + (\bar{\underline{5}}, -\frac{1}{2}, -\frac{1}{2}) \\ & + (\underline{10}, \frac{1}{4}, -\frac{1}{4}) + (\bar{\underline{5}}, \frac{1}{4}, \frac{3}{4}) + (\underline{1}, \frac{1}{4}, -\frac{5}{4}) \end{aligned} \quad (16)$$

Therefore the singlet $(1, 1/4, -5/4)$ cannot couple to $H\bar{H}$, and it makes sense to restrict our attention to the additional superpotential term (6), where now

$N \sim (1,1,0)$. Since the kinetic function U (5) also contains $H\bar{H}N$, we find

$$m_{1/2} = -\frac{bg^2}{4} \left(\frac{1}{16 S_R T_R^3} \right)^{1/2} \lambda \left(\frac{v x}{\bar{v}} + \frac{\bar{v} x}{v} + \frac{v \bar{v}}{x} \right) \quad (17)$$

which has a finite non-zero value if $v \equiv \langle 0|H|0\rangle$, $\bar{v} \equiv \langle 0|\bar{H}|0\rangle$ and $x \equiv \langle 0|N|0\rangle$ are all non-zero. The low-energy effective potential is given by

$$V_{eff}(H, \bar{H}, N) = m^2 |H|^2 + \bar{m}^2 |\bar{H}|^2 + m_N^2 |N|^2 + \lambda A \text{Re } H\bar{H}N \\ + \lambda^2 (|H\bar{H}|^2 + |N|^2 (|H|^2 + |\bar{H}|^2)) + D\text{-terms} \quad (18)$$

The D term of interest to us is that corresponding to $U(1)_E$:

$$D_E = (|N|^2 - \frac{1}{2}|H|^2 - \frac{1}{2}|\bar{H}|^2) \quad (19)$$

The other hypercharge $U(1)_F$ may or may not be broken already at high energies. If it survives to low energies its D term would produce quartic terms in N but no novel effect, and therefore we neglect it for simplicity in the following. The boundary conditions for the soft SUSY breaking parameters in (18) are

$$m^2 = \bar{m}^2 = m_N^2 = A^2 = 0 \quad (20)$$

at the scale where the gaugino condensate is¹³⁾⁻¹⁵⁾ formed in the hidden sector of the theory, and $m_{3/2}$ is generated. This is the scale where the unbroken subgroup of the hidden E_8 becomes strong. Note that the singlet field N cannot destabilize the gauge hierarchy²⁴⁾ because in $SU(n,1)/SU(n) \times U(1)$ no-scale models there can be no supersymmetry breaking in the superheavy sector¹²⁾.

We take v and $\bar{v}$ to be real. Then $\text{Im} x = 0$ and the potential (18) has in general an acceptable minimum for $\text{Re} x \neq 0$. As a pedagogical example to illustrate the formation of a suitable gauge symmetry breaking minimum, we consider the limit of small λ , where the similarities to minimal $SU(3) \times SU(2) \times U(1)$ no-scale models¹⁰⁾⁻¹²⁾ are most apparent. To zeroth order in λ , there is a minimum with v and $\bar{v}$ both $\neq 0$ if

$$\bar{m}^2 + m^2 \left(\frac{g^2 - g_E^2}{g^2 + g_E^2} \right) < 0 \quad (21)$$

where $g^2 \equiv g_2^2 + g'^2$. The trilinear term in (18) then generates $x \neq 0$ in first order in λ :

$$x = - \frac{\lambda A v \bar{v}}{2 (m_N^2 - \frac{1}{2} g_E^2 (v^2 + \bar{v}^2))} \quad (22)$$

The condition (21) reduces to that found in the minimal $SU(3) \times SU(2) \times U(1)$ supergravity model¹⁸⁾ when $g_E \rightarrow 0$. (A simple renormalization group analysis actually suggests $\alpha_E \equiv g_E^2/4\pi \sim 1/100 \ll \alpha_2$.) As in previous dimensional transmutation models¹⁸⁾, the condition (21) will be satisfied when the theory is renormalized at mass scales

$$\mu < \mu_0 = m_P \exp(-O(1)/\alpha_t) \quad (23)$$

When this occurs

$$V, \bar{V}, x \propto \mu_0 \quad (24)$$

as in previous dimensional transmutation models¹⁸⁾, and with the flat $SU(n,1)/SU(n) \times U(1)$ potential one can also determine

$$m_{1/2} \propto \mu_0 \quad (25)$$

as in previous no-scale models¹⁰⁾⁻¹²⁾. We have given here a simplified discussion, and in general the condition (21) should be replaced by a more general sufficient condition for gauge symmetry breaking²⁵⁾. However, the slow logarithmic variation of the couplings always guarantees an analogous expression (23) for μ_0 - though in general also depending on α_λ - and hence realization of the usual no-scale scenario¹⁰⁾⁻¹²⁾.

Equations (24), (25) and (4c) now tell us that T_R is dynamically determined (8). Since $S_R = 1/g^2 = O(1)$ also, we now deduce

$$\langle 0|\phi|0 \rangle \sim \langle 0|\sigma|0 \rangle \sim O(1) \quad (26)$$

so that the compactification scale $m_C = e^{-\sigma/2} m_P$ and the grand unification scale $m_X = m_C^4/m_P^3$ are both $O(m_P)$. Assuming that $C = O(1)$, this also means that the gravitino mass (3) is $O(m_P)$ ¹⁹⁾, as are the physical masses of the spin-zero particles associated with the real and imaginary parts of the S field. The latter is the only true axion in the superstring: the imaginary part of the T field is associated with a truly massless Goldstone field^{12),6)}. Thus there are no cosmological gravitino or axion problems.

After obtaining the above results we received two papers²⁰⁾ which discuss the problems of strong and weak coupling arising if the conditions (26) are not satisfied. Equation (3) suggests that to avoid them one needs $m_{3/2} = 0(m_P)$ as proposed previously¹⁹⁾. Closer examination of (24), (25) and (4c) suggests that

$$\phi = O(g^{5/3}) \quad , \quad e^\sigma = O(g^{-1/4}) \quad (27)$$

which avoids these strong/weak coupling problems if g is not $\ll 1$, and supports the validity of the original heterotic string $\rightarrow N = 1, d = 10$ supergravity $\rightarrow N = 1, d = 4$ supergravity scenario⁴⁾ of sequential symmetry breaking if g is a viable expansion parameter. The result (8) means that the observable non-singlet chiral fields ϕ^i need only an $O(1)$ finite renormalization for their kinetic terms [see, e.g., Refs. 11), 12) and 8)] to be normalized canonically, and therefore higher order $\prod_{i=1}^n \phi^i$ terms in the chiral superpotential have coefficients $O(m_P^{3-n})$ as usually expected: this would not be the case if $T_R \gg 1$.

Some further words about trilinear Yukawa couplings are in order. It could be that λ (6) plays a key rôle in determining the critical scale: $\mu_0 = m_P \exp(-O(1)/\alpha_\lambda)$: $\alpha_\lambda \equiv \lambda^2/4\pi$. This would not alter the essential logic of the no-scale scenario, though one would lose the direct connection to the observable top quark mass. Another remark is that in superstring models one may expect¹⁶⁾ complete 27's of chiral fields at low energies, and hence one pair of Higgs doublets $H, \bar{H}$ and one $SU(3) \times SU(2)_L \times U(1)_Y$ singlet N per conventional fermion generation. Thus the Yukawa coupling λ (6) is really a matrix, and there are more $QQ^C H$ and $LL^C H$ couplings. In restricting our attention to a single coupling λ (6) we are assuming that the highest (third?) generation coupling dominates as for quarks and leptons. Turning to the $QQ^C H$ and $LL^C H$ couplings, it is plausible to believe that the highest (third?) generation Higgses have the largest couplings to the highest (third?) generation quarks and leptons. Maybe the top quark Yukawa coupling only drives the highest Higgs mass squared negative, and the lower "Higgses" do not acquire vacuum expectation values. This might help explain why the first two generations of quarks and leptons are so light: their wave functions⁵⁾ may be almost orthogonal to that of the third generation Higgs doublets which acquire vacuum expectation values.

Some final comments concern neutral current phenomenology²⁵⁾ in the presence of the extra $U(1)_E$ gauge symmetry spontaneously broken at low energies. The entries in the (Z, Z_E) mass² matrix are

$$\begin{pmatrix} Z_1 & Z_E \end{pmatrix} \begin{pmatrix} M_W^2 / \cos^2 \theta_W & \frac{1}{2} \frac{g_2 g_E (V^2 - \bar{V}^2)}{\cos^2 \theta_W} \\ \frac{1}{2} \frac{g_2 g_E (V^2 - \bar{V}^2)}{\cos^2 \theta_W} & g_E^2 \left(X^2 + \frac{1}{4} (V^2 + \bar{V}^2) \right) \end{pmatrix} \begin{pmatrix} Z \\ Z_E \end{pmatrix} \quad (30)$$

In the limit $v^2 - \bar{v}^2 \ll v^2 + \bar{v}^2$ the $W^\pm$ and lighter neutral gauge boson masses give

$$\rho \equiv \frac{M_W^2}{m_{Z_1}^2 \cos^2 \theta_W} \simeq 1 + \frac{V^2}{X^2} \cos^2 2\phi \quad (31)$$

where $v \equiv V \sin \phi$, $\bar{v} \equiv V \cos \phi$. To satisfy the experimental limit²⁶⁾ $(\rho-1) \lesssim 0.03$, v and $\bar{v}$ should be roughly equal as expected in no-scale models¹⁰⁾⁻¹²⁾. Other neutral current information²⁶⁾ does not impose very restrictive bounds on the parameters of the model²⁵⁾ as long as $m_{Z_2} \gtrsim 0(m_Z)$ corresponding to $x \gtrsim 0(200)$ GeV for $g_E^2 \approx 1/8$. As the dynamical determination of $m_{\frac{1}{2}}$ typically yields $m_{\frac{1}{2}} = 0(1) \times V$: $V = 0(200)$ GeV, this requirement is quite reasonable. It is amusing to note that Eq. (31) implies that Z_1 is lighter than the standard model Z_0 , which does not disagree with $\bar{p}p$ collider results.

We will return to the issues of potential minimization and neutral current phenomenology in a forthcoming paper²⁵⁾.

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