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SUPERSYMMETRY BREAKING FROM SUPERSTRINGS*

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Abstract

The gauge hierarchy problem is briefly reviewed and a class of effective field theories obtained from superstrings is described. These are characterized by a classical symmetry, related to the space-time duality of string theory, that is responsible for the suppression of observable supersymmetry breaking effects. At the quantum level, the symmetry is broken by anomalies that provide the seed of observable supersymmetry breaking, and an acceptably large gauge hierarchy may be generated.

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INTRODUCTION

In this talk I will report on some recent work [1] with Pierre Binétruy on effective field theories obtained from superstrings. The physics motivation is the gauge hierarchy problem, which I will first briefly review. I will then review the theoretical framework in which we are working, namely effective supergravity theories obtained from the $E_8 \times E_8$ heterotic string.

A certain class of these theories is characterized by an invariance, at the classical level, under a group of global, nonlinear (like chiral symmetry in low energy pion physics) transformations among the fields of the effective theory. We have shown [1] that this symmetry protects the scalars and gauginos of the observed gauge group from acquiring masses when supersymmetry (SUSY) is broken in a "hidden" sector of the theory, that couples to our world with interactions of gravitational strength only.

This symmetry group includes chiral transformations on fermion fields, as well as scale transformations, and is therefore broken at the quantum level by the well known chiral and conformal anomalies. These anomalies, in collusion with nonperturbative effects in the strongly coupled gauge interactions of the hidden sector, provide the seed of SUSY breaking in the observable sector. We find [1] that a very mild hierarchy between the Planck scale and the scale (i.e., the gravitino mass) of SUSY breaking in the hidden sector is sufficient to generate an acceptably large (for phenomenology) hierarchy in the observed sector.

THE GAUGE HIERARCHY PROBLEM

The gauge hierarchy problem may be simply expressed in the context of the Standard Model by writing the renormalized Higgs mass m_H as

$$m_H^2 = \frac{\lambda}{8}(\mathrm{TeV})^2 = m_H^2(\mathrm{tree}) + \frac{g^2}{16\pi^2}\Lambda^2 + \cdots. \qquad (1)$$

Here g is the weak gauge coupling constant, and λ is the renormalized coupling constant for scalar self-couplings. The right hand side of (1) represents the classical value plus the sum of quantum corrections, which are quadratically divergent, as indicated by the appearance of the cut-off Λ. If perturbation theory makes sense, λ can be no larger than 1 (or at least 4π). Then the first equality suggests $m_H < (.35\text{–}1.2)$ TeV, and so we need $\Lambda < (8\text{–}30)$ TeV. I wish to emphasize that one cannot [2] evade the gauge hierarchy problem by a strongly interacting scalar sector, i.e., by letting $\lambda \gg 1$ in (1). Of course, purely within the context of the renormalizable standard model, there is not really a gauge hierarchy problem. The infinite quadratic divergences can be absorbed into a redefinition of the Higgs mass, whose value is simply fixed by measurement. However if the underlying theory includes Higgs couplings to heavier particles, such as GUT vector bosons, quantum corrections will include finite terms with Λ in (1) replaced by the masses of these particles.

There are three standard "solutions" to the gauge hierarchy problem, which I briefly recall. I will list them in what I view as increasing order of plausibility; many people would

disagree with my ordering.

Compositeness. In this scenario, the standard model is an effective theory, some or all of whose "elementary" particles are bound states of yet more elementary objects. The theory makes sense up to momentum scales of order of the inverse radius of compositeness r_c, so

$$\Lambda \to \Lambda_c \sim r_c^{-1} \qquad (2)$$

in (1). If quarks and leptons are composite, those with common constituents should couple to one another via four-fermion interactions with an effective Fermi constant $G \sim 4\pi r_c^2$. Existing experiments suggest $r_c < (\mathrm{TeV})^{-1}$; recent results from Tristan [3] give more stringent limits, with $\Lambda_c > 5$ TeV in one channel.

Technicolor. In this case only the Higgs sector is composite. The theory [4] mimics the observed properties of QCD. New asymptotically free gauge interactions are assumed, which break the electroweak symmetry via a technifermion condensate

$$< \bar{f}^T f^T > \simeq (f_{\pi T})^3 \equiv (\frac{1}{4}\mathrm{TeV})^3. \qquad (3)$$

Here $f_{\pi T}$ is the strength of the coupling to the axial current of the technipion π^T, analogous to the pion decay constant, f_π. This number is fixed at 250 GeV, so as to correctly reproduce the observed W, Z masses. The scale at which the effective "low energy" theory ceases to be valid is determined by the scale Λ_{TQCD} at which the technigauge interactions become strong:

$$\Lambda \to \Lambda_{TQCD} \sim f_\pi^T \simeq 250\ \mathrm{GeV}. \qquad (4)$$

As yet, no one has succeeded in constructing an experimentally viable, nor a grand unifiable, model that incorporates this idea.

SUSY. In this case [5] the quantum corrections on the right hand side of (1) are damped by cancellations between boson and fermion loops, which are complete if SUSY is unbroken. Since observation tells us that SUSY is certainly broken, the effective cut-off is provided by the fermion-boson mass splitting:

$$\Lambda \to \Lambda_{SUSY} = |m_{fermion} - m_{boson}|. \qquad (5)$$

It is possible to construct viable SUSY extensions of the standard model, but the scale parameter (5) is simply put in by hand, so we have not really solved the gauge hierarchy problem in this way.

THE HETEROTIC STRING

According to the presently most popular hope for a fully unified theory, the Standard Model is an effective theory that is a low energy limit of the heterotic [6] string [7] theory. Starting from a string theory in 10 dimensions with an $E_8 \times E_8$ gauge group, one ends up, at energies sufficiently below the Planck scale, with a supersymmetric field theory in 4 dimensions [8], with a generally smaller gauge group $\mathcal{H} \times \mathcal{G}$. $\mathcal{H}$ describes a "hidden sector", that has interactions with observed matter of only gravitational strength, and $\mathcal{G} \supset SU(3)_c \times SU(2)_L \times U(1)$ is the gauge group of observed matter. Part of the gauge symmetry may be broken (or additional gauge symmetries may be generated) by the $10 \to 4$ dimensional compactification process itself, and part of it may be broken by the Hosotani mechanism [9], in which gauge flux is trapped around space-tubes in the compact manifold. There are now many more examples of effective theories from superstrings than one once thought could emerge. For illustrative purposes, I will stick to the original "conventional" scenario, in which the "observed" E_8 is broken to E_6, long known to be the largest phenomenologically viable GUT, by the compactification process. Then the observed sector is a supersymmetric Yang-Mills theory, with gauge bosons and gauginos in the adjoint representation of $\mathcal{G} \subset E_6$, coupled to matter, i.e., to quarks, squarks, leptons, sleptons, Higgs, Higgsinos, ...

The hidden sector is assumed to be described by a pure SUSY Yang-Mills theory, $\mathcal{H} \subset E_8$, which is asymptotically free, and therefore infrared enslaved. At some energy scale Λ_c, below the compactification scale Λ_{GUT} at which all the gauge couplings are equal, the hidden gauge multiplets become confined and chiral symmetry is broken, as in QCD, by a fermion condensate. In this case the fermions are the gauginos of the hidden sector:

$$< \lambda\lambda >_{hid} \sim \Lambda_c^3 \neq 0. \tag{6}$$

The condensate (6) breaks SUSY [10], and by itself would generate a positive cosmological constant. If this were the only source of SUSY breaking, and of a cosmological constant, the condensate would be forced dynamically to vanish, due to the condition that the vacuum energy be minimized.

Another source of SUSY breaking is the (quantized) vacuum expectation value of an antisymmetric tensor field H_{LMN}, that is present in 10-dimensional supergravity:

$$H_{LMN} = \nabla_L B_{MN}, \quad L, M, N = 0, \ldots, 9,$$
$$\int dV^{lmn} < H_{lmn} > = 2\pi n \neq 0, \quad l, m, n = 4, \ldots, 9 \tag{7}$$

The vev (7) can arise if H-flux is trapped around a 3-dimensional space-hole in the compact 6-dimensional manifold, in a manner analagous to the Hosotani mechanism for breaking the gauge symmetry. When (6) and (7) are both present, λ and H_{LMN} couple in such a way [11] that the overall contribution to the classical cosmological constant vanishes. There are other potential sources of SUSY breaking, such as a gravitino condensate [12], that might play a similar role.

The particle spectrum of the effective four dimensional field theory includes the gauge supermultiplets $W^a = (\lambda^a, F^a_{\mu\nu} - i\tilde{F}^a_{\mu\nu})$ (gauginos and gauge bosons) and chiral supermultiplets $\Phi^i = (\varphi^i, \chi^i)$ that contain the matter fields (φ^i = squarks, sleptons, Higgs particles,

..., χ^i = quarks, ...). In the "conventional" scenario these are all remnants of the gauge supermultiplets in ten dimensions:

$$A_M \to A_\mu + \varphi_m, \quad \mu = 0, \ldots, 3, \quad m = 4, \ldots, 9. \tag{8}$$

Thus for each gauge boson A_M in ten dimensions, there are potentially one gauge boson A_μ and six scalars φ_m (and their superpartners) in four dimensions. However not all of these are massless. In the "conventional" picture ($E_8 \to E_6$ in the observed sector) the massless 4-vectors are in the adjoint of E_6, while the massless scalars are in $(27 + \overline{27})$'s that make up the difference: $(\text{adjoint})_{E_8} - (\text{adjoint})_{E_6}$. In addition there are gauge singlet chiral supermultiplets associated with the structure of the compact manifold. Two of these, $S = (s, \chi^S)$ and $T = (t, \chi^T)$ are of special interest. Their scalar components are [13]

$$s = e^{3\sigma}\phi^{-\frac{3}{4}} + 3i\sqrt{2}D,$$
$$t = e^{\sigma}\phi^{\frac{3}{4}} - i\sqrt{2}a + \frac{1}{2}\sum_i |\varphi^i|^2. \tag{9}$$

In (9) ϕ is the dilaton of ten-dimensional supergravity, D and a are two axions that are remnants of the antisymmetric tensor (7):

$$a \propto \epsilon^{lm} B_{lm}, \quad \partial_\mu D \propto \epsilon_{\mu\nu\rho\tau}\phi^{-\frac{3}{2}} e^{6\sigma} H^{\nu\rho\tau}, \tag{10}$$

and σ is the "breathing mode" or "compacton" whose vev determines the size of the compact manifold with metric $g_{lm} = g^{(0)}_{lm} e^{\sigma}$. Thus the GUT—or compactification—scale, which is the inverse of the radius R of compactification, is determined by the vev (in Planck mass units)

$$\Lambda^2_{GUT} = R^{-2} = < e^{-4\sigma} > = < (\text{ResRet})^{-1} > . \tag{11}$$

The total number of gauge singlet chiral multiplets, as well as the number of matter generations (#27's − #$\overline{27}$'s) is determined by the detailed topology of the compact manifold.

THE LAGRANGIAN OF THE EFFECTIVE FIELD THEORY

The classical lagrangian for a general supergravity theory in four dimensions is determined [14,15] by three functions of the chiral superfields:

$$\Phi^a = \Phi^i, S, T, \ldots \tag{12}$$

These are

i) A gauge field normalization function $f(\Phi) = \bar{f}(\bar{\Phi})^\dagger$. In the superfield formulation [16] the Yang-Mills part of the lagrangian is given by

$$\mathcal{L}_{YM} = \frac{1}{4}\int d^2\Theta f(\Phi) W^a_\alpha W^\alpha_a + h.c. = -\frac{1}{4}\{\text{Re}f(\varphi) F^a_{\mu\nu} F^{\mu\nu}_a + \text{Im}f(\varphi)\tilde{F}^a_{\mu\nu} F^{\mu\nu}_a\} + \cdots . \tag{13}$$

Here Θ is a complex two-component fermionic variable in superspace: $x \to x, \Theta, \bar{\Theta}$, and we have indicated some of the terms that appear after Θ integration when the superfields are expanded in terms of their component fields. The first term in this expansion implies that the gauge coupling constant is proportional to the *vev* $< \mathrm{Re} f(\varphi) >$.

ii) The Kähler potential $K(\Phi, \bar{\Phi}) = K(\Phi, \bar{\Phi})^{\dagger}$, which determines, for example, chiral multiplet kinetic energy terms:

$$\mathcal{L}_{KE}(\Phi) = K_{a\bar{b}} \partial_{\mu} \varphi^{a} \partial^{\mu} \bar{\varphi}^{\bar{b}} + \cdots, \quad K_{a\bar{b}} = \frac{\partial^2 K}{\partial \varphi^a \partial \bar{\varphi}^{\bar{b}}}. \tag{14}$$

iii) The superpotential $W(\Phi) = \overline{W}(\bar{\Phi})^{\dagger}$, which determines the Yukawa couplings and the scalar potential:

$$\mathcal{L}_{pot} = \int d^2\Theta e^{K/2} W(\Phi) + h.c. = -e^{\mathcal{G}} [\mathcal{G}_a (\mathcal{G}^{-1})^{a\bar{b}} \mathcal{G}_{\bar{b}} - 3) + \cdots, \tag{15}$$

where on the right hand side I have introduced the generalized Kähler potential

$$\mathcal{G} = K + \ln |W|^2 \tag{16}$$

of Cremmer it et al. [15]. In fact, the theory defined above is classically invariant [15,16] under a Kähler transformation that redefines both the Kähler potential and the superpotential in terms of a holomorphic function $F(\Phi) = \bar{F}(\bar{\Phi})^{\dagger}$:

$$K \to K' = K + F + \bar{F}, \quad W \to W' = e^{-F} W, \tag{17}$$

provided one also transforms the fermions by a chiral rotation; for example

$$W^{a} \to e^{-\mathrm{Im} F/2} W^{a}, \quad \lambda^{a} \to e^{-\mathrm{Im} F/2} \lambda^{a}. \tag{18}$$

This last transformation is anomalous at the quantum level, a point that will be important in the discussion below. One can fix the "Kähler gauge" by a specific choice of the function F. In particular, choosing $F = -\ln W$ casts the lagrangian in a form [15] that depends on only two functions of the scalar fields, f and $\mathcal{G}$.

Here I will describe a prototype [13] supergravity model from superstrings, with non-perturbative SUSY breaking included [11]. The functions (13)—(15) are given in terms of the superfields (12) by

$$f = S, \tag{19a}$$

$$K = -\ln(S + \bar{S}) - 3\ln(T + \bar{T} - |\Phi|^2), \quad |\Phi|^2 = \sum_{i=1}^{N} \Phi^i \bar{\Phi}^i, \tag{19b}$$

$$W(\Phi) = c_{ijk} \Phi^i \Phi^j \Phi^k + c + h e^{-3 b_0 S/2}. \tag{19c}$$

The last two terms in the superpotential W are parametrizations of nonperturbative SUSY breaking effects [11]. The parameter c is proportional to the *vev* of the antisymmetric tensor

field strength (7), and the last term represents the gaugino condensate (6), where b_0 is a group theory number that determines the β-function of the hidden sector Yang-Mills theory. The form of this term can be understood in terms of the standard R.G.E. result,

$$\Lambda_c = e^{-b_0/2g^2} \Lambda_{GUT} = \left\langle \frac{e^{-b_0 \mathrm{Re} s/2}}{\sqrt{\mathrm{Re} s \mathrm{Re} t}} \right\rangle, \tag{20}$$

together with the relation implied by (13) and (19a) (there are no free parameters in the string theory!) between the *vev* of s and the gauge coupling constant g at the GUT scale (11), where all gauge couplings are equal:

$$g^2(\Lambda_{GUT}) = < (\mathrm{Re} s)^{-1} >. \tag{21}$$

The structure of the condensate term in W is further justified by symmetry considerations [11,17]. For $c = h = 0$, the theory is formally invariant under the Kähler transformation (17) with

$$F = i\alpha, \quad K \to K, \quad W \to e^{i\alpha} W, \quad \lambda \to e^{i\alpha} \lambda, \quad \alpha \text{ real}. \tag{22}$$

This symmetry, which is just the "R-symmetry" of renormalizable SUSY models, is broken at the quantum level (which cannot be ignored for the strongly interacting hidden Yang-Mills sector) due to the chiral anomaly; under (22)

$$\delta \mathcal{L} = -\frac{i\alpha}{6 b_0} (F\tilde{F})_{hid}. \tag{23}$$

However, because of the coupling (19a), (13) of the Yang-Mills supermultiplet to the S-supermultiplet, the variation (23) can be cancelled by a shift in S:

$$S \to S - \frac{2i\alpha}{3 b_0}. \tag{24}$$

The combined transformations (22) and (24) are an exact (neglecting the c-term and quantum corrections in the observed gauge sector) invariance of the theory; this is reflected by the transformation property

$$W(S) = e^{-3 b_0/2} h \to e^{i\alpha} W(S) \tag{25}$$

of the superpotential for S in (19c).

The general features of the theory defined by (19), first obtained by Witten [13] for the case of a simple torus compactification, are common to a broad class of more realistic models [18]. These possess the following properties at the classical level. The gravitino mass $m_{\tilde{G}}$ can be nonvanishing, so that local supersymmetry is broken. The cosmological constant vanishes, as do the observable gaugino masses m_{λ}, the gauge nonsinglet scalar masses m_{φ}, and "A-terms", which are trilinear gauge nonsinglet scalar self-couplings that, if present, would also break SUSY. Thus there is no manifestation of SUSY breaking in the observable sector.

One loop corrections have been evaluated [19] in this effective (nonrenormalizable) theory, cut off at the scale of gaugino condensation (20), with the result that the classical features described above are unchanged at the one loop level.

CLASSICAL SYMMETRIES OF THE THEORY

The class of 4 d theories considered possesses [20] a classical nonlinear, noncompact global symmetry. They are in fact nonlinear σ-models, much like the effective pion theory of low energy QCD, where chiral $SU(2)$ symmetry is realized via nonlinear transformations among the pion fields. The difference here is that the global symmetry group is the noncompact group $SU(1,1)$:

$$T \to T' = \frac{aT - ib}{icT + d}, \quad \Phi^i \to \Phi^{i\prime} = \frac{\Phi^i}{icT + d}, \quad S \to S' = S, \qquad ad - bc = 1, \quad a,b,c,d \text{ real}. \tag{26}$$

For $c = h = 0$, Eqs.(26) in fact represent a Kähler transformation (17), with

$$F = 3\ln(icT + d), \tag{27}$$

under which the full lagrangian is invariant provided the fermion fields undergo a chiral transformation (18).

The group of transformations (26) includes a subset, with $a = d = 0$, $bc = -1$, under which

$$t \to b^2/t, \tag{28}$$

where b^2 is a finite, continuous, positive, real parameter. The string scale M_S is related to the Planck scale M_P by

$$M_P = <(\mathrm{Re}s)^{\frac{1}{2}}> M_S, \tag{29}$$

so when the theory is expressed in string mass units, (28) corresponds to an inversion of the radius of compactification (11):

$$R^2 = \Lambda_{GUT}^{-2} = < \mathrm{Re}s\mathrm{Re}t > /M_P^2 = < \mathrm{Re}t > /M_S^2 \to b^2/R^2. \tag{30}$$

For the special case of integer b, this is the well known "duality" transformation, which leaves the string spectrum invariant.

When we allow $c, h \neq 0$, the $S(1,1)$ symmetry can be formally maintained by allowing these parameters to transform like a superpotential, Eq.(17):

$$c \to c' = e^{-F}c, \quad h \to h' = e^{-F}h. \tag{31}$$

This makes sense when one recalls that c and h are actually the *vevs* of underlying dynamical variables; therefore their values will relax to those that minimize the total vacuum energy

density. This was precisely the attitude taken in [19], where it was found that observable SUSY breaking vanishes at the overall ground state of the one-loop corrected effective theory. It has recently been shown [1] that the classical $SU(1,1)$ symmetry is responsible for the cancellation of observable SUSY breaking effects found [19] by explicit calculation.

ANOMALIES AS THE SEED OF OBSERVABLE SYMMETRY BREAKING

The noncompact symmetry (26) and the R-symmetry (22) are broken at the quantum level by the chiral anomaly and also by the conformal anomaly. This latter anomaly arises because $SU(1,1)$ includes the scale transformation ($c = b = 0$, $ad = 1$ in (26)) $t \to a^2 t$, under which the cut-off for the theory (which at energies above the scale of hidden gaugino condensation is just Λ_{GUT}, Eq.(11)) scales as

$$\Lambda_{GUT}^2 \propto <(\mathrm{Re}t)^{-1}> \to a^{-2}\Lambda_{GUT}^2. \tag{32}$$

Then under $SU(1,1) \times U(1)_R$

$$\delta\mathcal{L} = \frac{2b_0}{3}\{\mathrm{Re}F(t)F^a_{\mu\nu}F^{\mu\nu}_a + \mathrm{Im}F(t)\tilde{F}^a_{\mu\nu}F^{\mu\nu}_a\} + \cdots = -\frac{2b_0}{3}\int d^2\Theta F(T)W^a_\alpha W^\alpha_a + h.c., \tag{33}$$

where $F(T)$ is the function defining the Kähler transformation (26) or (22).

The dominant observable effect of these anomalies is associated with the highest mass scale at which nonperturbative effects come into play. In the context of the effective 4-d field theory, these arise from instantons and gaugino condensation in the hidden Yang-Mills sector. Just as one can construct low energy effective Lagrangians for pseudoscalar mesons that are $q\bar{q}$ bound states using the symmetries of QCD and the chiral and conformal anomaly, one can use [1] $SU(1,1) \times U(1)_R$ and its anomalies, together with supersymmetry [21], to construct an effective lagrangian for the composite multiplet U:

$$\frac{1}{4}W^\alpha W_\alpha \to U = \lambda H^3 e^{K/2} e^{-3S/2b_0}, \tag{34}$$

or equivalently the chiral multiplet H, which is the lightest composite state, with mass m_H, of the (confined) hidden gauge sector. The Kähler transformation property of H

$$H \to H' = e^{-F/3}H \tag{35}$$

can be inferred from those of Φ^a, W^α. With this transformation property, the anomalies are correctly reproduced [21,22,1] by the following effective potential lagrangian for the composite chiral field:

$$\mathcal{L}^{eff}_{pot} = \int d^2\Theta e^{K/2} 2b_0\lambda e^{-3S/2b_0} H^3 \ln(H/\mu) \equiv \int d^2\Theta e^{K/2} W(H,S)$$
$$= \int d^2\Theta[SU + U\ln(4Ug^2/\Lambda_{GUT}^3\lambda\mu^3)], \tag{36}$$

which is also invariant [22] under the nonanomalous transformation (22) + (24). Aside from the numerical parameters (of order unity) μ and λ, the logarithmic term in (36) is precisely what is expected from the one-instanton contribution [23]. Note that Λ_{GUT} is the physical cut-off for the theory above the condensate scale, and that the gauge multiplets W^a are normalized with a factor $g^{-2} = < Res >$ relative to the canonical normalization. In addition, the ground state configuration is determined by the minimum with respect to H of the potential (36). This gives

$$< H >= h_0 = \mu e^{-1/3}, \quad < \bar{\lambda}\lambda >_{hid} = 4 < U > = \frac{\lambda h_0^3}{g^2}\Lambda_c^3. \tag{37}$$

Again (37) corresponds exactly to the one-instanton contribution [23].

It remains to specify the H-dependence of the Kähler potential. The symmetries of the theory dictate [1] the form

$$K = -\ln(S + \bar{S}) - 3\ln(T + \bar{T} - |\Phi|^2 - |H|^2). \tag{38}$$

The effective classical theory below the scale of condensation is determined by "integating out" the H-supermultiplet, that is, by the sum of tree diagrams with "light" particles ($m < \Lambda_c$) on external legs only. It turns out that there are no such diagrams with H-exchange, because vertices with a single H leg vanish at the H ground state.

On the other hand, retaining one-loop corrections from the H degrees of freedom, whose couplings explicitly include the anomalous symmetry beaking, one finds [1] that the effective low energy theory defined in this way is no longer totally $SU(1,1)$ invariant, although no observable SUSY breaking appears at the "classical" level of this effective theory. However, at the one-loop level of this effective theory, gaugino masses are generated in the observable sector that are of order

$$m_{\tilde{g}} \sim \frac{1}{(16\pi^2 m_P^2)^2} m_{\tilde{G}} m_H^2 \Lambda_c^2. \tag{39}$$

The factor $(4\pi)^{-4}$ appears in (39) because the effect arises first at two-loop order, the factor $m_{\tilde{G}}$ is the necessary signal of SUSY breaking, the factor m_H^2 is the signal of $SU(1,1) \times U(1)_R$ breaking, and Λ_c^2 is the effective cut-off. This last factor arises essentially for dimensional reasons: the couplings responsible for transmitting the knowledge of symmetry breaking to the observable sector are nonrenormalizable interactions with dimensionful coupling constants proportional to m_P^{-2}.

Solving [19] the minimization conditions for the effective theory at the one-loop level yields, for vacua with broken supersymmetry, the values

$$m_{\tilde{G}} \simeq \frac{1}{3} m_H \simeq \frac{1}{3}\Lambda_c \simeq (10^{-2}\text{–}10^{-1})\frac{m_P}{\sqrt{b_0 c}}, \tag{40}$$

where the parameter c is proportional to the vev (7) of H_{LMN}. The quantization condition (7) and dimensional analysis suggest [19] $c > 10^3 n$ if $c \neq 0$, or

$$m_{\tilde{g}} < 10^{-16} m_P \simeq 2\text{TeV}. \tag{41}$$

Once gauginos acquire masses, gauge nonsinglet scalars (in particular the Higgs particles) will acquire masses $m_\varphi \sim \frac{\alpha}{4\pi} m_{\tilde{g}}$ at the next loop order in the renormalizable gauge interactions.

The results reported here may be slightly, but not essentially, modified by the inclusion of a T-dependence in the superpotential $W(S,H)$ defined by (36): $\mu \to \mu(T)$. Such a modification is expected, so as to restore [24] the discrete form (a,b,c,d integers in (26)) of $SU(1,1)$, which is known [25] to be an exact symmetry of string theory. Such a term has recently been found [26] as a loop correction to the function f, Eq.(19a), from the heavy string modes. From the point of view of the four dimensional effective field theory, such a term is expected to arise [1] from T-couplings to the axial current, just as chiral anomalies in QCD induce a pion coupling to $(F\tilde{F})_{QED}$ via the pion coupling to the axial quark current.

The superstring context used here is not the most general one, but there is a broad class of models with similar features, so these results suggest that there is hope, after all, of extracting meaningful physics from superstrings.

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