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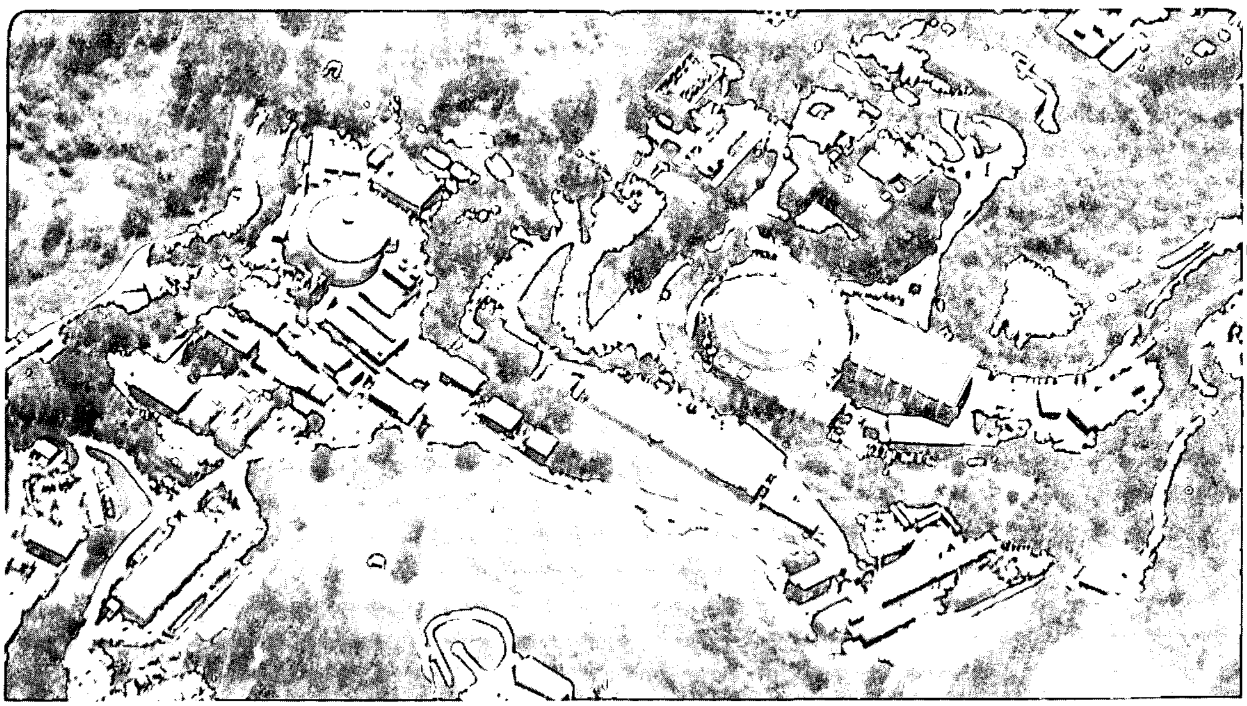

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PHYSICS AT THE PLANCK SCALE*

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Abstract

Effective supergravity theories suggested by superstrings can be explored to determine their potential for successfully describing both observed physics at zero temperature and an inflationary cosmology. An important ingredient in this study is the dynamics of gaugino condensation, which has been the subject of recent activity.

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INTRODUCTION

We are accustomed to the well-established Standard Model of particle physics, namely the $SU(3)_c \times SU(2)_L \times U(1)$ gauge theory whose three (scale dependent) gauge coupling constants apparently converge, when appropriately extrapolated, to a common value at a momentum scale, Λ_{GUT}, that is a few orders of magnitude below the Planck scale. Grand unified theories assumed that above that scale the theory would be (at least up to the Planck scale) a conventional, renormalizable quantum field theory (RQFT). The more modern view is that the theory above the GUT scale is a superstring theory [1], possibly in as many as ten dimensions, and that just below it the effective theory is a supergravity theory that coincides with the standard model (or some RQFT extension of it) when the energy scale is low enough that nonrenormalizable couplings, that are suppressed by inverse powers of the Planck mass $m_{Pl} = (8\pi G_N)^{-\frac{1}{2}}$, can be ignored.

The vacuum configuration that defines the effective renormalizable quantum field theory is determined by the vacuum expectation values (*vevs*) of scalar fields φ of the effective supergravity theory that are generally of order $< \varphi >\sim m_{Pl}$. Obviously, truncation to the effective renormalizable theory cannot be made without first determining these scalar *vevs*. These in turn determine:

- the value of the cosmological constant;
- the scales of supersymmetry (SUSY) and hence of gauge symmetry breaking–more specifically whether or not the possibility exists of generating a sufficiently large hierarchy between the observed scale of electroweak symmetry breaking and the fundamental scale m_{Pl};
- the shape of the effective scalar potential at temperatures not far below the Planck temperature, which determines whether or not a successful inflationary scenario is possible in the context of the theories considered.

In this talk I will focus on work done mainly in collaboration with Pierre Binétruy that draws on features common to a broad class of effective supergravity theories obtained as classical vacua of the heterotic superstring theory [2]. The idea is to use constraints from the presumed underlying superstring theory to determine the effective "low energy" supergravity theory, much as one uses constraints from the QCD theory of quarks and gluons to determine an effective low energy theory of hadrons.

THE HETEROTIC STRING

According to the presently most popular hope for a fully unified theory, the Standard Model is an effective theory that is the low energy limit of the heterotic string theory [2]. 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 [3], 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 [4], 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 (denoted by λ or $\tilde{g}$) of the hidden sector:

$$< \bar{\lambda}\lambda >_{hid} \sim \Lambda_c^3 \neq 0. \tag{1}$$

The condensate (1) breaks SUSY [5], and by itself would generate a positive cosmological constant. If a single gaugino condensate 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:

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

The *vev* (2) can arise if H-flux is trapped around a 3-dimensional space-hole in the compact 6-dimensional manifold, in a manner analogous to the Hosotani mechanism for breaking the gauge symmetry. When (1) and (2) are both present, λ and H_{LMN} couple in such a way [6] that the overall contribution to the classical cosmological constant vanishes. There are other potential sources of SUSY breaking, such as a gravitino condensate [7], that might play a similar role.

The particle spectrum of the effective four dimensional field theory includes the gauge supermultiplets $W^a = (\lambda_L^a, F_{\mu\nu}^a - i\tilde{F}_{\mu\nu}^a)$ (that is, gauginos $\tilde{g}^a = \lambda^a$ and gauge bosons g^a) and the chiral supermultiplets $\Phi^i = (\varphi^i, \chi_L^i)$ 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:

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

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 is the supergravity multiplet containing the graviton G and the gravitino $\tilde{G}$, and a number of gauge singlet chiral supermultiplets. One of these, $S = (s, \chi^S)$, includes the "dilaton" Res, where s is the (complex) scalar component of S, whose *vev* determines the inverse squared gauge coupling constant at the GUT scale:

$$g^2(\Lambda_{GUT}) = < (\mathrm{Re}s)^{-1} > . \tag{4}$$

Among the other gauge singlet chiral multiplets that may be present are the so-called moduli $T_i = (t_i, \chi_i^T)$, whose scalar components t_i correspond to quantum fluctuations in the topology of the compact manifold. One of these corresponds to fluctuations in the overall size of the manifold and its *vev* determines the compactification, or GUT, scale, that is, the inverse radius of compactification R:

$$\Lambda^2_{GUT} = R^{-2} = m^2_{St} < (\mathrm{Re}t)^{-1} > = m^2_{Pl} < (\mathrm{Re}s\mathrm{Re}t)^{-1} >, \tag{5}$$

where m^2_{St} is the string tension. The total number of moduli is equal to the number of matter generations ($\#27$'s $- \#\overline{27}$'s) and is determined by the detailed topology of the compact manifold.

These gauge nonsinglet scalar fields and their axion superpartners Ims, Imt_i, are important because classically their values are flat directions in the effective potential. The quantum effects that lift these degeneracies determine the overall vacuum configuration and therefore the scales of the elementary particle theory. For the same reason these fields are prime candidates for inflatons.

EXAMPLES OF INFLATIONARY SCENARIOS

Because the S field couples to all the Yang-Mills fields and their superpartners, condensation of the hidden gauginos induces a potential for s [6]. For illustrative purposes, Fig. 1 shows the potential in the effective theory obtained by Dine *et al.* [6] for the field variable $\phi = -\ln\mathrm{Re}(s)/\sqrt{2}$, which is the canonically normalized scalar field that rolls according the classical equations of motion. The shape of the potential depends on the ratio $\hat{c} = |\tilde{c}/\tilde{h}|$, where

$$\tilde{c} \propto \int dV^{lmn} < H_{lmn} >, \quad \tilde{h} \propto < \bar{\lambda}\lambda >_{hid}, \tag{6}$$

are parameters of the effective theory that reflect the *vevs* of hidden dynamical variables. The potential has vanishing vacuum energy, and for $\hat{c} > 1.21$ the only

global minimum is the SUSY conserving one: $m_{\tilde{G}} = 0$. For $\hat{c} < 1.21$ the vacuum is degenerate in the t direction and local SUSY breaking is possible. It turns out [8,9] that slow rollover inflation with a suitable number of e-foldings can occur only for the potential of Fig. 1(e), with the parameter fine-tuned to the value $\hat{c} = .937 \pm (3 \times 10^{-4})$. For $\hat{c}$ just above this value the field ϕ would have to tunnel to the true vacuum, entailing the problems of the old inflationary scenario, and for values just below it there is no plateau from which a slow rollover can start. In addition one must find a mechanism for stabilizing ϕ at the required initial value; temperature corrections cannot do this [8]. One suggested mechanism [9] is to take into account that condensation does not occur instantaneously at a critical temperature $T_c \sim \Lambda_c$. Ellis *et al.* [9] parameterized the turn-on of condensation by $\tilde{h} = \tilde{h}(T) = \tilde{h}_0(1 - T/T_c)^n$ for temperatures $T < T_c$. This effectively turns Fig. 1 into a moving picture with the potential evolving from 1(a) to 1(e); it then becomes plausible that the field ends up in the second minimum in 1(b) and 1(c) and evolves to the plateau in 1(e) from which it rolls to the true vacuum. Aside from the fine tuning problem, one needs to understand better the dynamics of gaugino condensation to see whether such a "moving picture" is plausible. Gaugino condensation has been the subject of much recent work [10–14] which I will describe below.

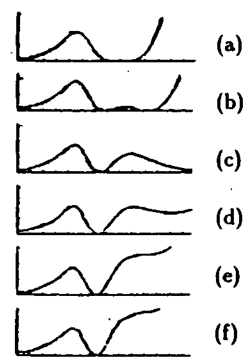

Figure 1. The shape of the potential of Dine *et al.* [6] for the variable $\phi = -\ln(\mathrm{Re}s)/\sqrt{2}$, plotted in the region $-2.3 < \phi < 3.7$, and for (a) $\hat{c} = 1.21$, (b) $1 < \hat{c} < 1.21$, (c) $\hat{c} = 1$, (d) $.937 < \hat{c} < 1$, (e) $\hat{c} = .937$, (f) $\hat{c} < .937$. In all cases the potential vanishes at the SUSY conserving ($m_{\tilde{G}} = 0$) configuration Re$s \to \infty$ ($\phi \to -\infty$); this is the only zero of the potential for $\hat{c} > 1.21$.

The fine tuning problem could be evaded in a chaotic inflationary picture [15] as I will point out later. In addition to the scalar field ϕ or Res there is its axion (pseudoscalar) superpartner. We found [8] that its potential as determined by the effective theory of Dine *et al.* does not allow for sufficient inflation. A recent reexamination [16] of axion-induced inflation suggests that axions with characteristic scale parameter f of order of the Planck scale can lead to sufficient inflation; however the specific values (including the height of the potential, which here is proportional to the gravitino mass) required in that analysis are not satisfied by the potential for Ims in this effective theory.

The potential for the modulus field t is flat at the tree level of the effective theory of Dine *et al.*, and the degeneracy in Ret is lifted at the one loop level. The earliest attempts [17–19] to evaluate the one-loop corrections resulted in a potential with a negative (and even unbounded [17,19]) cosmological constant. Maeda *et al.* [20] studied the case [18] of a bounded potential by simply adding a constant energy density to fix the global minimum at zero, and found that sufficient inflation could occur. Their key observation was that the t field is driven to a value away from its true vacuum when (as in a chaotic early universe, for example) the other fields are not at their vacuum values. A problem with this analysis is that one cannot simply simply add a constant energy density to the potential; in superstring theory everything is determined by the *vevs* of scalar fields. Another problem is that their ansatz violates general results that show that the SUSY conserving vacuum is always a (possibly degenerate) global minimum of the potential if it is bounded from below.

With Dawson and Hinchliffe we [21] subsequently showed that additional quantum corrections from loop momenta between the condensation and compactification scales could restore the boundedness of the effective one-loop potential, and that local SUSY breaking could occur for reasonable values of the relevant parameters, in which case the potential vanishes at its minimum which remains degenerate in one direction in parameter space. The resulting potential, shown in Fig. 2 in the ($\bar{c}$,Ret) plane with the other *vevs* fixed at their ground state values, does not yield sufficient inflation when the t field is considered in isolation. On the other hand all observable SUSY breaking effects are found to vanish along the minimum, which is good for phenomenology, since nonvanishing effects at this order would be much too large. The origin of this suppression of observable SUSY breaking has been identified [10]; it is related to a classical

symmetry of the effective theory that I will display below.

These last results rely on the assumption that the true vacuum is a minimum with respect to all *vevs*, including the parameter $\tilde{h}$ that reflects the size of gaugino condensation. To determine whether such an assumption is justified requires a better understanding of the dynamics of gaugino condensation, as does a full exploration of possible inflationary scenarios. Gaugino condensation, will be the subject of the remainder of the talk, with the end result that, when constraints from string theory are fully implemented, the picture sketched above may be considerably altered.

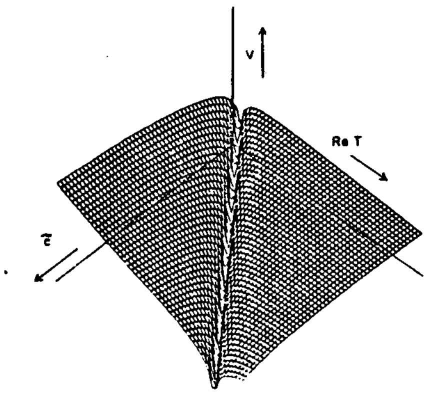

Figure 2. The one loop potential [21] in the (Ret, $\bar{c}$) plane, with the remaining parameters fixed at their ground state values, for the effective theory of Dine *et al.* [6]. The continuous degeneracy is reduced to a discrete one when the quantization condition (2,6) for $\bar{c}$ is taken into account.

CONSTRUCTION OF THE EFFECTIVE LAGRANGIAN

The general supergravity lagrangian was first given in its complete form by Cremmer *et al.* [22]. Its symmetry structure is most readily manifest in the superfield formulation later developed by Binétruy *et al.* [23], in which the Lagrangian can be expressed by just three terms:

$$\mathcal{L} = \mathcal{L}_E + \mathcal{L}_{pot} + \mathcal{L}_{YM}. \tag{7}$$

The first term

$$\mathcal{L}_E = -3\int d^4\Theta \mathcal{E}\mathcal{R} + h.c. \tag{8}$$

is the generalized Einstein term. It contains the pure supergravity part as well as the (noncanonical, i.e. including derivative couplings) kinetic energy terms for the chiral supermultiplets. The second term:

$$\mathcal{L}_{pot} = \int d^2\Theta e^{K(\Phi,\bar{\Phi})/2} W(\Phi) + h.c., \tag{9}$$

contains the Yukawa couplings and the scalar potential, and the third term

$$\mathcal{L}_{YM} = \frac{1}{4}\int d^2\Theta f(\Phi) W^a_\alpha W^\alpha_a + h.c. \tag{10}$$

is the Yang-Mills lagrangian. Here Θ is a complex two-component fermionic variable in superspace: $x \to x, \Theta, \bar{\Theta}$. Integration over the anticommuting variable Θ is equivalent to differentiation. The expansion of the above expressions in terms of component fields includes derivatives that are covariant with respect to general coordinate, gauge and Kähler transformations. A Kähler transformation is a redefinition of the Kähler potential $K(\Phi,\bar{\Phi}) = K(\Phi,\bar{\Phi})^\dagger$ and of the superpotential $W(\Phi) = \overline{W}(\bar{\Phi})^\dagger$ by a holomorphic function $F(\Phi) = \bar{F}(\bar{\Phi})^\dagger$ of the chiral supermultiplets $\Phi = (\varphi, \chi)$:

$$K \to K' = K + F + \overline{F}, \quad W \to W' = e^{-F}W, \tag{11}$$

Since this transformation changes $e^{K/2}W$ by a phase that can be compensated by a phase transformation of the integration variable Θ, the theory defined above is classically invariant [22,23] under Kähler transformations provided one transforms the superfields $\mathcal{R}$ and W^a_α by the same phase; for example

$$W^a_\alpha \to e^{-\mathrm{Im}F/2} W^a_\alpha. \tag{12}$$

This last transformation, which implies a chiral rotation on the left-handed gaugino field λ^a_L:

$$\lambda^a_\alpha \to e^{-\mathrm{Im}F/2} \lambda^a_\alpha, \tag{13}$$

is anomalous at the quantum level, a point that will be important in the discussion below. (Here a is a gauge index and α is a Dirac index.).

The theory is completely specified by the field content, the gauge group and the three functions K, W and f of the chiral superfields. 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 [22] that depends on only two functions of the chiral superfields, f and $\mathcal{G} = K + \ln|W|^2$. The gauge coupling constant is determined by the *vev*

$$< \mathrm{Re} f(\varphi) >= g^{-2}. \tag{14}$$

I will consider a prototype [24] supergravity model from superstrings, with just one modulus T and one matter generation. The generalization to three moduli and three matter generations is straightforward [14]. The functions K, W and f are given in terms of the superfields $\Phi = \Phi^i$, S and T by

$$f = S, \tag{15a}$$

$$K = -\ln(S+\bar{S}) - 3\ln(T+\bar{T}-|\Phi|^2), \quad |\Phi|^2 = \sum_i \Phi^i \bar{\Phi}^{\bar{\imath}}, \tag{15b}$$

$$W(\Phi) = c_{ijk}\Phi^i\Phi^j\Phi^k + \tilde{c}. \tag{15c}$$

where the last term in the superpotential W parameterizes the nonperturbative SUSY breaking induced by the *vev* (2,6) of the antisymmetric tensor field strength. Comparison of (15a) with (14) yields the relation (4), and using this with (5) gives for the *vev* of the Kähler potential ($< |\varphi^i|^2 > \ll <\mathrm{Re}t>$)

$$16 < e^K > \approx < (\mathrm{Re}s)^{-1}(\mathrm{Re}t)^{-3} >= g^{-4}\Lambda^6_{GUT}/m^6_{Pl}. \tag{16}$$

For $\tilde{c} = 0$, the supergravity theory defined above is classically invariant [25] under the nonlinear transformations

$$T \to T' = \frac{aT - ib}{icT + d}, \quad \Phi^i \to \Phi^{i'} = \frac{\Phi^i}{icT+d}, \quad S \to S' = S,$$
$$ad - bc = 1, \quad a,b,c,d \text{ real}, \tag{17}$$

that form the three-parameter group $SU(1,1)$ or $SL(2,\mathcal{R})$ which is the noncompact version of the rotation group $SU(2)$ or $O(3)$. Eq.(17) effects a Kähler transformation (11) with

$$F = 3\ln(icT + d), \tag{18}$$

under which the full lagrangian is invariant provided the gaugino fields undergo the chiral transformation (12,13). The group of transformations (17) includes a subset with $a = d = 0$, $bc = -1$, under which:

$$t \to b^2/t, \tag{19}$$

where b is a finite, continuous, real parameter. From (5) we see that when the theory is expressed in string mass units m_{St}, (19) corresponds to an inversion of the radius of compactification. For the special case of integer b, this is the well known "duality" transformation that leaves the string spectrum invariant.

When we allow $\tilde{c} \neq 0$, the $SU(1,1)$ symmetry can be formally maintained by allowing this parameter to transform like the superpotential in (11):

$$\tilde{c} \rightarrow \tilde{c}' = e^{-F}\tilde{c}. \tag{20}$$

It can be shown [10] that (20) corresponds to the correct transformation property of the antisymmetric tensor field $< H_{LMN} >$ in (2). The $SU(1,1)$ symmetry has recently been shown [10] to be at the origin of the cancellation of observable SUSY breaking effects, found [21] previously by explicit calculation in perturbation theory.

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 [10] $SU(1,1)$ and its anomalies, together with supersymmetry [26], to construct an effective lagrangian for the lightest hidden-sector chiral supermultiplet, denoted $H = (h, \chi^H)$, that is a bound state, with mass m_H, of the hidden gauge supermultiplet. Here the relevant classical symmetries are the chiral transformation (12,13) and the scale transformation

$$\Lambda_{GUT} = m_{Pl}(2g)^{\frac{1}{2}} < e^{K/6} > \rightarrow e^{\mathrm{Re}F/3}\Lambda_{GUT}, \tag{21}$$

which follows from (17), (18) and the identification (16) for the cut-off Λ_{GUT} of the effective theory. Scale and chiral invariance are broken at the quantum level by anomalies. The variation of the one-loop quantum corrected SUSY Yang-Mills lagrangian is well-known [26]:

$$\mathcal{L}_{YM}^{\text{Quantum}} \rightarrow \mathcal{L}_{YM}^{\text{Quantum}} - \frac{2b_0}{3}\int d^2\Theta F(T) W_\alpha^a W_a^\alpha + h.c., \tag{22}$$

where b_0 is the group theory number that determines the β-function for the Yang-Mills theory. We wish to construct an effective potential for a composite

chiral multiple H that represents the lightest bound state of the confined hidden Yang-Mills sector. Thus we identify [27,10]:

$$\begin{aligned}\mathcal{L}_{YM}^{\text{Quantum}} \Rightarrow \mathcal{L}_{pot}^{eff} &= \int d^2\Theta e^{K/2} 2b_0\lambda e^{-3S/2b_0} H^3 \ln(H/\mu) + h.c.\\ &= \int d^2\Theta[SU + U\ln(4Ug^2(S)/\Lambda_{GUT}^3(\Phi)\lambda\mu^3) + h.c.],\end{aligned} \tag{23}$$

where λ and μ are constants of order unity, U is the interpolating field for the Yang-Mills composite operator:

$$\frac{1}{4}W_\alpha^a W_a^\alpha \Rightarrow U = \lambda H^3 e^{K/2} e^{-3S/2b_0}, \tag{25}$$

and the factor

$$\begin{aligned}< 2b_0 \ln(H/\mu) > &= \frac{1}{g^2}\left[1 + \frac{2b_0 g^2}{3}\ln(4 < U > g^2/\Lambda_{GUT}^3\lambda\mu^3)\right]\\ &= \frac{1}{g^2}\left[1 + \frac{2b_0 g^2}{3}\ln(4e^{-1}\Lambda_c^3/\Lambda_{GUT}^3)\right]\end{aligned} \tag{26}$$

is the one-loop Yang-Mills field wave function renormalization from the compactification scale to the condensation scale. The composite field (25) and the effective lagrangian defined by (23) transform according to (12) and (22), respectively, provided we impose

$$H \rightarrow H' = e^{-F/3}H = \frac{H}{icT + d} \tag{27}$$

and take as Kähler potential (which determines the H-superfield kinetic energy) [10]:

$$K = -\ln(S + \bar{S}) - 3\ln(T + \bar{T} - |\Phi|^2 - |H|^2). \tag{28}$$

We can now solve the effective theory for the condensate *vev*:

$$< H > = h_0 = \mu e^{-1/3}, \quad \text{or} \quad < \bar{\lambda}\lambda >_{hid} = 4 < U > = \frac{\lambda h_0^3}{g^2}\Lambda_c^3. \tag{29}$$

If we fix H at its ground state value, we obtain an effective theory for Φ^i, S, T and the observable-sector Yang-Mills fields that is defined by ($15a, b$) and the superpotential

$$W(\Phi) = c_{ijk}\Phi^i\Phi^j\Phi^k + \tilde{c} + \tilde{h}e^{-3b_0 S/2}, \quad \tilde{h} = -\frac{2b_0}{3}\lambda\mu^3 e^{-1}, \tag{30}$$

which is precisely the theory of Dine *et al.* studied previously. This is a reasonable approximation because the (SUSY conserving) mass $m_H \sim \Lambda_c > m_{\tilde{G}}$ of the composite supermultiplet H is larger than the other masses generated at the classical level of the effective theory defined by (23). As mentioned earlier, the effective theory defined by (30) has the following properties at the classical level [6] and at the one-loop [21] level: the cosmological constant vanishes, the gravitino mass $m_{\tilde{G}}$ can be nonvanishing, so that local supersymmetry is broken, in which case the vacuum is degenerate, and there is no manifestation of SUSY breaking in the observable sector. Note that with $\tilde{h}$ now determined by (29,30), minimization of the effective potential with respect to the parameter $\tilde{h}$ is equivalent to minimization with respect to the parameter μ (or λ). The presence of an undetermined parameter reflects [10] an additional degree of freedom of the underlying theory, namely the gauge field strength $F_{\mu\nu}$, which does not appear explicitly as an independent propagating field in the composite supermultiplet formulation.

Including loop correction from the H-sector, one finds that masses are generated for the gauginos of the observable sector that are of order [10]

$$m_{\tilde{g}} \sim \frac{1}{(16\pi^2 m_{Pl}^2)^2} m_{\tilde{G}} m_H^2 \Lambda_c^2 < 10^{-15} m_{Pl} \simeq 2\ TeV.$$

$$\text{for } m_{\tilde{G}} < m_H \simeq \Lambda_c \sim 10^{-2} m_{Pl}. \qquad (31)$$

The factor $(4\pi)^{-4}$ appears in (31) because the effect arises first at two-loop order in the effective theory, the factor $m_{\tilde{G}}$ is the necessary signal of SUSY breaking, the factor m_H^2 is the signal of $SU(1,1)$ 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_{Pl}^{-2}. Once gauginos acquire masses, gauge nonsinglet scalars (in particular the Higgs particle, whose mass governs the scale of electroweak symmetry breaking and hence the gauge hierarchy) will acquire masses $m_\varphi \sim \frac{\alpha}{4\pi} m_{\tilde{g}}$ at the next loop order in the renormalizable gauge interactions.

This theory thus appears capable of generating a plausible gauge hierarchy, which in turn suggests that one should take seriously the study of its

possible inflationary scenarios. Specifically, in the context of chaotic inflation [15], we should not set the gauge nonsinglet scalars φ^i in (30) to zero at temperatures near the Planck temperature. We can parameterize their effect by letting $\hat{c} \rightarrow \hat{c}(\varphi) = \hat{c} + c_{ijk} < \varphi^i \varphi^j \varphi^k >$, which is temperature dependent. Then one recovers the "moving picture" scenario for Fig. 1 [28]. Moreover, one need not fine tune the zero-temperature parameter $\hat{c}$, since the the desired exponential expansion will occur if the field Res sits for a sufficiently long time in, for example, the second (false) minimum of Fig. 2d. No tunneling is required as the "graceful exit" is provided [29] by the potential changing its shape due to the evolution of the φ^i fields to their ground state values that are equal to zero (or nearly so with respect to the scales in question). This analysis has not been pursued further, because recent developments in understanding the symmetry of effective theories from superstrings appear to change the picture appreciably.

RESTORATION OF MODULAR INVARIANCE

In the formalism presented above, the continuous classical symmetry $SU(1,1)$ or $SL(2,\mathcal{R})$ is broken by anomalies at the quantum level. However the discrete subgroup $SL(2,\mathcal{Z})$ [a,b,c,d integers in (17)] of $SL(2,\mathcal{R})$ is known [30] to be an exact symmetry to all orders in string perturbation theory. This so-called "modular invariance" is restored by adopting, instead of (23), the effective lagrangian [12]

$$\mathcal{L}_{YM}^{\text{Quantum}} \Rightarrow \mathcal{L}_{pot}^{eff} = \int d^2\Theta e^{K/2} 2b_0 \lambda e^{-3S/2b_0} H^3 \ln(H\eta^2(T)/\mu) + h.c., \qquad (32)$$

where

$$\eta(T) = e^{-\pi T/12} \prod_{m=1}^{\infty} (1 - e^{-2m\pi T}) \qquad (33)$$

is the Dedekind η-function. This is the unique function of the chiral superfields that has the required analyticity and $SL(2,\mathcal{Z})$ transformation properties [12]. This additional contribution to the Yang-Mills wave function renormalization was found independently [31] from the direct calculation of finite threshold corrections to the leading log approximation that arise from heavy string mode loops. The result (32) has been generalized to the cases of several gaugino condensates [13,14] and of several moduli [14].

The effective scalar potential [11,12] for the theory defined by (32) is unbounded from below. If one appeals to some unknown dynamics, such as wormholes, to restore positivity and erase the (infinitely negative) cosmological constant, one loses all predictive power because the same dynamics can affect other features; specifically one can make no statements about the gauge hierarchy (which involves expanding in quantum fluctuations about the true vacuum) nor about inflation (which involves tracing the evolution of the field configurations toward the true vacuum).

A cure for this disease has recently been found [14]. It suffices to reinterpret the above formalism, which is based on one-loop results for the Yang-Mills field renormalization in SUSY Yang-Mills theories. In the construction (23) or (32) the field renormalization is incorporated into a redefinition of the superpotential. We can instead incorporate it into a redefinition of the Kähler potential, which is in fact the interpretation that is supported by known renormalization effects in supersymmetry and in supergravity [32,10]. That is, instead of (32) and (28), we define the effective composite theory by

$$\mathcal{L}_{YM}^{\text{Quantum}} \Rightarrow \mathcal{L}_{pot}^{eff} = \int d^2\Theta S e^{K/2} \lambda e^{-3S/2b_0} H^3 + h.c., \tag{34a}$$

$$K = -\ln(S+\bar{S})$$
$$-3\ln(T+\bar{T}-|\Phi|^2-|H|^2[1-\tfrac{1}{8}b_0 g^2(S)\ln(H\eta^2(T)/\mu)+\text{h.c.}]). \tag{34b}$$

The two theories are identical (after appropriate redefinitions of field variables) to first order in the loop expansion parameter b_0; indeed it is a term of order b_0^2 that drives $< V > \to -\infty$ in the theory defined by (32). At its classical level, the theory defined by (34) has once again a vanishing cosmological constant and (for $\tilde{c} \neq 0$) a degenerate vacuum with local SUSY breaking ($m_{\tilde{G}} \neq 0$) possible, and again no SUSY breaking appears in the observable sector at the classical level of this effective theory. [14]

What remains to be done is to reexamine the issues of quantum corrections and inflationary potentials for this modified theory. The analysis is considerably more complicated in this case, and one can expect some qualitative differences from the scenarios described above. For example, one generally finds [33] a different hierarchy of masses, namely

$$m_S > m_H \sim m_{\tilde{G}}, \tag{35}$$

with SUSY-invariant masses for the S-supermultiplet up to terms of order $m_{\tilde{G}}$. This means, for example, that it makes no sense to fix the H-supermultiplet at its *vev* and to study the effective potential for s.

It may also be the case that the theory defined by (34) is sufficiently complex that a realistic picture will emerge only when all the moduli fields are included. For example, much of the degeneracy of the simple model defined by (15,23) is lifted by the appearance of the η-function (33). On the other hand, including three moduli (and three matter generations) will increase the degree of degeneracy. The degeneracy in the moduli directions plays an important role in suppressing observable SUSY breaking effects. Finally, in contrast with the simple model (15,23), which remains degenerate in the Imt direction to all orders in perturbation theory, the axion components of the moduli superfields may develop nontrivial potentials at the quantum level; this might provide additional interesting possibilities for inflation [16].

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