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Particle Physics and Superstrings

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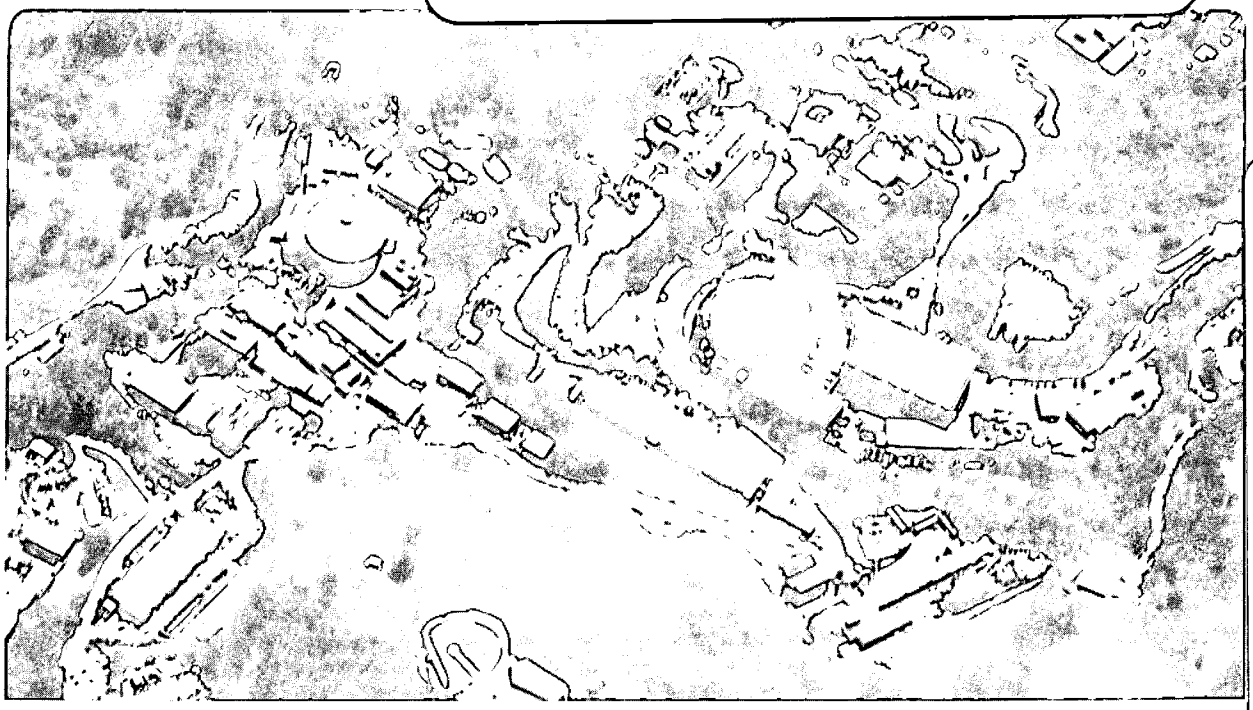

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PARTICLE PHYSICS AND SUPERSTRINGS*

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PARTICLE PHYSICS AND SUPERSTRINGS

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Abstract

Implications of recent precision measurements of the standard model gauge coupling constants are discussed in the context of superstring theory.

1. Introduction

This talk is intended to serve as a reminder that, as particle physicists, our task is to explain observed phenomena and make predictions for new phenomena. The new element in information from experiment is the high precision measurement at LEP of the three coupling constants, α_1, α_2 α_3, of the standard model $SU(3)_c \otimes SU(2)_L \otimes U(1)$ gauge group, as determined at the Z mass scale: $\mu = m_Z$. The two-loop analysis of the renormalization group equations (RGE) has been performed by many groups,[1,2] and recent trends in the data have suggested for some time that the minimal supersymmetric extension of the standard model (MSSM) is favored over the standard model (SM) if unification of the three coupling constants at a single scale is assumed. What the newest LEP data provide is higher precision and therefore a firmer basis for this conclusion.

In this talk I will focus, for concreteness, on a recent analysis[3] of DELPHI data. Optimistically, one could interpret the LEP results as determining three fundamental scales of the string theory: the string tension $T \approx (1.3\times10^{18}GeV)^2$, the radius of compactification $R \approx 10^{-16}GeV^{-1}$, and the SUSY mass gap $m_S \approx 1TeV$. Needless to say, such an interpretation is fraught with caveats. The extraction of the parameters at the unification scale is highly model dependent,

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and there are significant uncertainties associated with threshold effects.

Nevertheless such precise data present new challenges to theorists. Is it now possible to rule out classes of string vacua? Can one fit the data for a specific compact topology and extract experimentally determined values for the string parameters within that class of string vacua? Conversely, given experimentally determined values for the scales T, R, m_S, is it possible to calculate them from the theory within a class of string vacua?

2. What Does the Data Really Tell Us?

Amaldi, de Boer and Fürstenau[3] have extrapolated the values of the α_i, as measured by the DELPHI Collaboration, using the two-loop RGE for both the SM and the MSSM with several assumptions on the Higgs content. They find that unification of the couplings in the minimal SM, with just one complex Higgs doublet, is ruled out by more than seven standard deviations. Increasing the Higgs content to six complex doublets, they find that unification is possible, but with a unification scale $\Lambda_{GUT} \approx 3 \times 10^{13}$ that is unacceptably low if the unified theory allows proton decay with the standard couplings of, for example, minimal[4] $SU(5)$.

In contrast, the MSSM with two chiral supermultiplets of Higgs doublets fits the data well, while increasing the Higgs content to four chiral supermultiplets requires a fit with both a low unification scale ($\Lambda_{GUT} \approx 2 \times 10^{14} GeV$) and a SUSY mass gap ($m_S \approx 2 \times 10^{10} GeV$) that is too high to be relevant to the solution of the gauge hierarchy problem. The MSSM fit assumes that the mass of one Higgs supermultiplet is equal to the Z-mass and the other is equal to m_S, an assumption that should be borne in mind when considering the implications of the data discussed below. One then has three output parameters, namely the unification scale Λ_{GUT}, the common value of the coupling α_G at that scale, and the SUSY mass gap, that are obtained from the three input parameters α_i using the RGE. Amaldi *et al.* find[3]

$$\Lambda_{GUT} = 10^{16\pm.3} GeV, \quad \alpha_G^{-1} = 25.7 \pm 1.7, \quad m_S = 10^{0\pm1} TeV. \tag{1}$$

They also allowed a spread in superpartner masses over the range $m_S/30$ to $30 \times m_S$; this spread tends to lower the value of m_S and slightly raise the value of Λ_{GUT}.

To get a qualitative idea of what this data implies more generally, assume that the standard model gauge couplings do indeed unify at a single scale Λ_{GUT}. Whether or not the theory is supersymmetric, there is a linear combination of the α_i^{-1}, namely

$$\alpha_1^{-1} + 2\alpha_3^{-1} - 3\alpha_2^{-1}$$

that is not renormalized at one loop if particles other than gauge bosons (and gauginos) form complete representations of $SU(5)$. (Whether or not there is an $SU(5)$ grand unified theory has no bearing on this statement; it is simply a reflection of the fact that $SU(5)$ is the minimal unifying group for the standard model, and, indeed, the "observed" fermion spectrum, including the top quark, fills out complete $SU(5)$ representations.) Then one gets the result

$$\sin^2\theta_w = \frac{1}{6} + \frac{5}{9}\frac{\alpha_{em}}{\alpha_3} + \Delta_I + \delta_{2-loop} + \text{threshold effects.} \tag{2}$$

The first two terms on the right hand side of (2) correspond to the one-loop contributions of the gauge (super)multiplets to the RGE, and Δ_I is the contribution of $SU(5)$-incomplete matter (super)multiplets. Inserting the DELPHI values[3] for $\sin^2\theta_w$, α_3 and α_{em}, Eq.(2) reads

$$0.2336 \pm 0.0018 = 0.2066^{+0.0019}_{-0.0018}{}^{+0.0061}_{-0.0039} + \delta_{2-\text{loop}} + \Delta_I + \text{threshold effects.} \tag{3}$$

The first set of errors on the right hand side of (3) reflect the experimental errors quoted[3] by Amaldi *et al.*, and the second set corresponds to a theoretical uncertainty of $\pm.13$ assigned by Langacker and Luo[2] to α_3 in order to roughly incorporate higher order QCD corrections. If we assume the MSSM value[5] $\delta_{2-\text{loop}} \approx 0.0029$ and neglect threshold effects, we find a discrepancy of about seven standard deviations from $\Delta_I = 0$. In the minimal SM the only contribution to Δ_I is a single complex doublet of Higgs scalars which is insufficient to account for the discrepancy, whereas two Higgs supermultiplets with the mass assignments mentioned above fit the data nicely. Adding more Higgs multiplets overcompensates the discrepancy unless their masses are pushed to much higher scales so as to reduce their contribution to the running of the couplings.

Quite generally, what the data is telling us is that either there are incomplete $SU(5)$ multiplets, and/or the gauge group must be modified at some intermediate scale Λ_{Int} before the unification scale is reached. For example, Langacker and Luo, who find fits similar to those of Amaldi *et al.* also found a good

fit[2] to the data using a nonsupersymmetric SO(10) model with the symmetry breaking pattern

$$SO(10) \xrightarrow{\Lambda_{GUT}\sim 10^{16}GeV} SU(3)_c \otimes SU(2)_L \otimes SU(2)_R$$
$$\xrightarrow{\Lambda_{Int}\sim 10^{10}GeV} SU(3)_c \otimes SU(2)_L \otimes U(1).$$

Both incomplete $SU(5)$ multiplets and intermediate scale gauge symmetry breaking can occur in models obtained from superstring theory.

3. Examples of Scenarios from Superstrings

In the three generation model of Greene *et al.*[6], the $E_8' \otimes E_8$ gauge group[7] of the heterotic string[8] is broken via Calabi-Yau compactification[9] to $E_8' \otimes E_6$, and E_6 is further broken[9] by the Hosotani mechanism[10] to $[SU(3)]^3$ which is in turn broken by the conventional Higgs mechanism, at several intermediate stages, to the Standard Model. The Hosotani mechanism entails noncontractable loops of gauge flux, such that the massless modes satisfy twisted boundary conditions that depend on their E_6 transformation properties, and incomplete 27 and $\overline{27}$ multiplets occur. Specifically the massless spectrum at the compactification scale consists of three complete matter generations (27's) and four complete pairs of $27 + \overline{27}$ (which contribute to the RGE at the two-loop level). In addition there are two incomplete multiplets, each consisting of two Higgs supermultiplets, one $SU(2)_L$ doublet lepton supermultiplet, and one charged plus one neutral $SU(2)_L$ singlet lepton supermultiplet. The neutral singlet does not contribute to the RGE, and the charged singlet supermultiplet contributes only through its $U(1)$ coupling. The lepton doublet has the same SM gauge couplings as the Higgs supermultiplets. Overall, the incomplete multiplets give a contribution to Δ_I in (2) that is slightly larger than that of six Higgs supermultiplets, which would much too large a contribution if there were a single unification scale. The question to be studied is whether or not there is a viable scenario for intermediate scale breaking that allows a fit of this model to the data.

Alternatively, one could consider a model with only complete $SU(5)$ representations, but with large mass splittings within these representations, so that Δ_I vanishes only above some large threshold. The 27 of E_6 decomposes under $SU(5)$ as

$$27 = \bar{5} + 10 + 1 + (5 + \bar{5}) + 1. \tag{4}$$

The SM left-handed fermions and their superpartners are in $\bar{5} + 10$, and the $(5 + \bar{5})$ contain two Higgs supermultiplets plus two $SU(3)$ triplet supermultiplets. Mechanisms for large doublet-triplet mass splittings have been found, for example, in the context of the "flipped $SU(5)$" model.[11]

One could also try to fit a "minimal" Calabi-Yau scenario where the surviving gauge group after Hosotani symmetry breaking is $SU(3)_c \otimes SU(2)_L \otimes [U(1)]^n$. The massless spectra in these models also include incomplete representations due to twisted boundary conditions, although known realizations[12] seem to require at least four generations. In this case SM gauge unification can occur at a single scale but there may be additional corrections to the RGE due to mixing[13] of the additional Z's with the Z of the Standard Model.

The main point is that the precision of the data now available may allow more serious restrictions on the viability of models suggested by-or derived from-the heterotic superstring.

4. Superstring Interpretations of the Data

In effective theories from superstrings, the gauge coupling g at the unification scale is determined by the vacuum expectation value (*vev*) of the dilaton field Res; using the result[3] of Amaldi *et al.*:

$$\alpha_G^{-1} = \frac{4\pi}{g^2} \approx 26 \tag{5}$$

gives

$$< \mathrm{Re}\, s >= g^{-2} \approx 2. \tag{6}$$

This in turn determines the string tension T which is related to the reduced Planck mass: $m_{Pl} = (8\pi G_N)^{-\frac{1}{2}} \simeq 1.8 \times 10^{18} GeV$ by:

$$\sqrt{T} = g m_{Pl}, \tag{7}$$

giving, up to additional renormalization effects between the unification scale and the string scale, the result stated in the introduction. However, if there are complete $SU(5)$ matter supermultiplets in addition to those present in the MSSM, this will raise the value of α_G, with no other effect in one-loop order. For example, for three generations of chiral supermultiplets filling E_6 27 representations (note that $\Delta_I = 0$ in this case), the one-loop β-function vanishes for

$SU(3)_c$, so that the QCD coupling ceases to evolve once the SUSY threshold m_S is reached, giving

$$\alpha_G^{-1} = \alpha_3^{-1}(m_S) + \delta_{2-loop} + \cdots \approx 12 + \delta_{2-loop} + \cdots, \quad g^{-2} \approx 1.$$

What is the interpretation of the unification scale Λ_{GUT}? Several different scenarios are possible, with different interpretations. For "four-dimensional" string vacua,[14] where all the extra dimensions of the heterotic string are interpreted as internal dimensions, the "radius of compactification" $R \equiv T^{-\frac{1}{2}}$ is fixed by the string scale and there are no breathing modes, or moduli. One obtains[15] an enlarged gauge group of rank 22 or less, that breaks via the conventional Higgs mechanism to the SM gauge group, possibly via intermediate symmetry breaking steps. In "ten-dimensional" string vacua, the "observed" E_8 group might break via compactification to some group G that in turn breaks via the Higgs mechanism to the Standard Model. For example, the flipped $SU(5)$ model could emerge from the symmetry breaking scenario[16]

$$E_8 \xrightarrow{SU(4)_{Holonomy}} SO(10) \xrightarrow{Hosotani} SU(5) \otimes U(1)' \xrightarrow{Higgs} SM.$$

Then the value of $\Lambda_{GUT} < R^{-1}$, as determined by the extrapolation of the LEP data, is identified with the scale where Higgs breaking occurs. Since the $U(1)$ of the SM is not embedded in $SU(5)$ in this model, one generally expects $\alpha_2 = \alpha_3 \neq \alpha_1$ at this scale.[5]

In all the scenarios just described, Λ_{GUT} is interpreted as the scale where the gauge symmetry of an effective four dimensional field theory is broken by the Higgs mechanism to a smaller gauge group. In such cases the threshold effects are well understood, provided $\Lambda_{GUT} \ll \sqrt{T}$, and for "10-dimensional" strings $\Lambda_{GUT} \ll R^{-1}$. The LEP data provides no information on the radius R of compactification in such scenarios.

A different interpretation applies if E_8 is broken at the compactification scale $R^{-1} \equiv \Lambda_{GUT}$ to the SM gauge group times some factor group, or to a larger group that subsequently breaks at an intermediate scale to the SM group. In the latter case the fit to an RGE analysis of the low energy data would include both Λ_{Int} and $\Lambda_{GUT} = R^{-1}$, as for the $SO(10)$ model[2] mentioned in Sect. 2. In these scenarios one directly extracts the value of R, for example, $R \approx 10^2 T^{-\frac{1}{2}}$ for the fit to the MSSM quoted in the Introduction. However, at the two-loop level, this

result is sensitive to the total number of complete $SU(5)$ matter representations that contribute to the RGE. For example, each additional generation of SM fermions was found[17] to increase the value of Λ_{GUT} by a factor 2; one would expect an even larger effect in the MSSM, as well as additional contributions from extra generations of $(5 + \bar{5})$ as in (4).

Since in this second class of scenarios, the GUT scale now characterizes the transition from a 10-dimensional field or string theory to a four dimensional field theory, threshold effects are much less well understood. Calculations[18,19] of the finite corrections[20] from integration over heavy string and Kaluza-Klein modes suggest that the α_i do not precisely unify at a single scale. The deviation from unification is in principle calculable for a given string vacuum, but is different for different vacua.[18,21]

Another problem is in the precise interpretation of the compactification scale, since the metric g_{ij} of the complex compact manifold (or orbifold) is not in general specified by a single scale. When one extracts[22-24] an effective 4-dimensional supergravity theory from a ten-dimensional one, the Kähler potential $K(Z, \bar{Z})$ that defines the low energy theory for the chiral multiplets Z, is related to the metric g_{ij}, $i, j = 1, 2, 3$, by

$$e^K = \frac{1}{16}(\mathrm{Re}s)^2 \det{}^{-4} g_{ij}. \tag{8}$$

A manifestly supersymmetric loop expansion of the effective (nonrenormalizable) 4-dimensional supergravity theory requires a field-dependent cut-off; consistent results have been found[25,26] with the cut-off

$$\frac{\Lambda_{GUT}}{m_{Pl}} = \left\langle \frac{16}{(\mathrm{Re}s)^2} e^K \right\rangle^{\frac{1}{4}} = < \det g_{ij} >^{-\frac{2}{3}}, \tag{9}$$

which for compactification on three tori of equal radius,[22] $g_{ij} = \delta_{ij}(Rm_{Pl})^{\frac{1}{2}}$, reduces to the intuitively expected value, $\Lambda_{GUT} = R^{-1}$. However even if (9) is the correct scale to be used in the RGE, there could be additional threshold effects due to a nonuniformity of the background metric.

Assuming that these ambiguities can be resolved in such a way that the parameters of the string vacuum can be reliably inferred from the data, one has then to produce predictions of the theory with which to compare them.

5. Can the Parameters of the String Vacuum be Calculated?

In this final section I will describe some attempts to answer the above question in the affirmative. There is a class of models based, for example, on orbifold compactification that have a "no-scale" structure,[27] with the Kähler potential (8) of the form[22,24]

$$K = -\ln(S+\bar{S}) - 3\ln\det(T_{ij}+T_{ij}-\sum_a \Phi_i^a\bar{\Phi}_j^a) + \text{twisted sector terms}, \quad (10)$$

where S is the chiral superfield that contains the dilaton Res, and $T_{ij}=\bar{T}_{ij}$ are the moduli superfields; the *vev* 's of their real scalar components $<\mathrm{Re}t_{ij}>$ determine the structure of the background metric and hence the compactification scale (9). The Φ_i^a are gauge nonsinglet matter superfields with scalar components φ_i^a. Neglecting possible massless modes in the twisted sector, the classical vacuum is supersymmetric and the *vev* 's $<s>$ and $<t>$ are undetermined classically in these models.

To make contact with observed physics one must appeal to nonperturbative sources of SUSY breaking. A popular candidate mechanism is gaugino condensation[28] in the "hidden" E_8' SUSY Yang-Mills sector of the effective supergravity theory in four dimensions. This induces a superpotential[29] for the dilaton superfield S:

$$W = W(\Phi) + \tilde{h}e^{-3S/2b_0}, \quad (11)$$

where $W(\Phi)$ is the classical superpotential for matter and b_0 determines the hidden sector β-function. However, the vacuum energy density for the theory defined by (10) and (11) is positive semi-definite. The potential vanishes at its minimum, which occurs for $<W>=0$, that is, for $<\varphi>=0$ and for either $<s>\to\infty$ $(g\to 0)$ or $\tilde{h}=0$, so condensation does not occur and supersymmetry remains unbroken. Note that the parameter $\tilde{h}$ can be interpreted[30,31] as the *vev* of the scalar component h of the lightest chiral supermultiplet H that occurs as a bound state in the confined hidden Yang-Mills sector:

$$<\bar{\lambda}\lambda>_{hid} = 4<e^{-3s/2b_0}h^3e^{K/2}> = \frac{<h^3>}{g^2}\Lambda_c^3, \quad (12)$$

where

$$\Lambda_c = e^{-1/2b_0g^2}\Lambda_{GUT}$$

is the scale of gaugino condensation.

Another possibility is that E_8' is broken at compactification to a product gauge group[32] $\prod G_i$ with different β-functions b_i. This scenario has recently been reconsidered[33] by Dixon *et al.* who took into account finite threshold corrections[18-21] Δ_i that separate the values of the different coupling constants at the GUT scale. Then the effective superpotential takes the form[33]

$$W = W(\Phi) + \sum_i \tilde{h}_i e^{(-3S/2b_i+\Delta_i)}. \quad (13)$$

Now a vanishing energy density can occur for $\tilde{h}_i \neq 0$; for example a solution with $<\mathrm{Re}s>\approx 2$, $\alpha_G^{-1}\approx 25$ was estimated[33] in the context of "four dimensional" strings. However the superpotential vanishes at this minimum, and therefore so does the gravitino mass

$$m_{\tilde{G}} = <e^{K/2}W>, \quad (14)$$

so supersymmetry remains unbroken. Provided one does not consider an effective theory with an unbounded potential,[34-36] this result that gaugino condensation alone cannot break supersymmetry appears to be quite general.[32] The same result was found for vanishing vacuum energy in models[37,38] with several gaugino condensates using the composite bound state formalism and including threshold corrections.

An additional source of SUSY breaking could be[29] the (quantized) *vev* of the field strength H_{lmn} of ten-dimensional supergravity. Then the effective superpotential takes the form

$$W = W(\Phi) + \tilde{c} + \tilde{h}e^{-3S/2b_0}, \quad (15)$$

$$\tilde{c} \propto \int dV^{lmn} <H_{lmn}> = 2\pi n \neq 0, \quad l,m,n = 4,\ldots,9,$$
$$H_{LMN} = \nabla_L B_{MN}, \quad L,M,N = 0,\ldots,9. \quad (16)$$

For $\tilde{c}\neq 0$ this effective theory has the following properties both at the classical level[29] and at the one-loop[39] 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. Nonrenormalization theorems for supergravity, together with classical nonlinear, noncompact symmetries suggest[26] that these results will persist to all orders of the effective theory defined by (10)

and (15). The gravitino mass (14) is determined as (assuming $\Lambda_{GUT} = R^{-1}$)

$$m_{\tilde{G}} \approx \frac{\tilde{c}\Lambda_{GUT}^3}{4g^4 m_{Pl}^2} \overset{?}{\approx} 3\tilde{c} \times 10^{11} GeV \tag{17}$$

where the number on the right is the value suggested by the analysis[3] of Amaldi *et al.* if $\Lambda_{GUT} = R^{-1}$. At one loop the degeneracy in the moduli space is lifted; for the case of a single modulus,[22] $g_{ij} = \delta_{ij}(\mathrm{Res}\,\mathrm{Re}t)^{\frac{1}{4}}$, one obtains[39]

$$m_{\tilde{G}} \approx 0.1\Lambda_{GUT} \approx 0.1 m_{Pl}/b_0\tilde{c}$$

[which is not necessarily compatible with the last expression in (17)].

The effective superpotential (15) may be obtained[31,26] by first constructing an effective theory for the composite superfield H, solving for the *vev* of its scalar component h and setting $H \equiv < H > = < h >$; for a suitably defined theory, this is equivalent to integrating out the superfield H at tree level. If instead the effective low energy theory is defined by integrating out the composite superfield at the one loop level, the chiral and conformal anomalies induce terms that break the nonlinear, noncompact symmetries in such a way that SUSY breaking is transmitted to the observable sector. Specifically a gaugino mass is generated that is of order[31]

$$m_{\tilde{g}} \sim \frac{m_{\tilde{G}} m_H^2 \Lambda_c^2}{(4\pi m_{Pl})^4} \overset{?}{\sim} \tilde{c} \times MeV, \tag{18}$$

where m_H is the mass of the H-supermultiplet, and the last figure is again the one suggested by (17) and the analysis[3] of DELPHI data for $\Lambda_{GUT} = R^{-1}$. The numbers in both (17) and (18)-which may be overkill-should probably be regarded as lower bounds in the general context of string models. They are very sensitive to the value of Λ_{GUT}; for $\Lambda_{GUT} = 10^{17} GeV$, they become, for example,

$$m_{\tilde{G}} \approx 3\tilde{c} \times 10^{14} GeV, \quad m_{\tilde{g}} \sim \tilde{c} \times 10 TeV.$$

The effective theory defined by (10) and (15) does not respect modular invariance, or target space duality, under which $R \to R^{-1}$. Restoration of modular invariance can be achieved[34,35] by replacing the constants

$$\tilde{h} \to \tilde{h}(T), \quad \tilde{c} \to \tilde{c}(T)$$

by appropriate functions of the moduli. At the level of the effective theory for the composite field H, the required T-dependence[35] of $\tilde{h}$ can be interpreted[19]

as threshold corrections[18] from integration over heavy Kaluza-Klein and string modes. This modification lifts the degeneracy in t at the classical level of the effective low energy theory, but the resulting potential[34-36] is unbounded in the direction $s \to 0$, $g \to \infty$. However, a reinterpretation[38] of the effective superpotential for H in terms of a wavefunction renormalization results in an effective theory that is once again positive semi-definite and of the no-scale form,[27] with a degenerate vacuum and no SUSY breaking in the observable sector at tree level. Local SUSY breaking, $m_{\tilde{G}} \neq 0$ can occur provided that modular invariance is broken by the classical nonperturbative *vev* (16): $\tilde{c} =$ constant $\neq 0$. For this theory the vacuum degeneracy is fully lifted at one loop. The results depend to some extent on uncertain threshold effects (that can be adjusted within reasonable limits to assure a vanishing cosmological constant at the one-loop level). A preliminary investigation[40] of the one-loop effective potential in the case of three moduli and a hidden E_8' gauge group shows a minimum of the potential with, for example, the parameters:

$$\alpha_G^{-1} \approx 25, \quad R \approx 2.8T^{-\frac{1}{2}} \approx 2.2 \times 10^{-18} GeV^{-1}, \quad m_{\tilde{G}} \approx 2 \times 10^{16} GeV,$$

where the gauge nonsinglet matter and twisted sector fields were set to zero. To determine whether or not observable SUSY breaking is sufficiently strongly suppressed in this model requires an understanding of how the threshold corrections depend on the gauge nonsinglet superfields Φ^i—note that for this model [see (8) and (10)] the radii of the three tori are actually given by $R_\alpha = < t_\alpha - \frac{1}{2}|\varphi_\alpha|^2 >$ with $\alpha = 1, 2, 3$.

The moral of this discussion is that, while the gap between superstring theory and phenomenology remains a wide one, attempts to bridge it may ultimately bear fruit.

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