Title
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Publication Date
1988-05-01

LBL-25278

Lawrence Berkeley Laboratory
UNIVERSITY OF CALIFORNIA

Physics Division

Presented at the XXIII Recontres de Moriond on Electroweak
Interactions and Unified Theories, Les Arcs, France,
March 6–13, 1988

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May 1988

Prepared for the U.S. Department of Energy under Contract Number DE-AC03-76SF00098.

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May 1988

UCB/PTH/88/12
LBL-25278

Low-energy structure of four-dimensional superstrings*

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Abstract

The $N = 1$, $d = 4$ supergravity theories derived as the low-energy limit of four-dimensional superstrings are discussed, focusing on the properties of their effective potentials. Gauge symmetry breaking is possible along several flat directions. A class of superpotential modifications is introduced, which describes supersymmetry breaking with vanishing cosmological constant and $Str\ M^2 = 0$ at any minimum of the tree level potential. Under more restrictive assumptions, there are minima with broken supersymmetry at which also $Str\ f(M^2) = 0$ for any function f, so that the whole one-loop cosmological constant vanishes. This result is interpreted in terms of a new discrete boson-fermion symmetry, relating particles whose helicities differ by 3/2, e.g. the graviton and the 'dilatino'.

Talk given at the XXIII Rencontres de Moriond - Electroweak Interactions and Unified Theories - Les Arcs, France - 6-13 March 1988.

*This work was supported in part by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U. S. Department of Energy under Contract DE-AC03-76SF00098 and in part by NSF under grant PHY-8515857.

†Supported by a Fellowship of the Istituto Nazionale di Fisica Nucleare, Sezione di Padova, Italy.

It was already emphasized in the previous talks [1),2)] that, even if the heterotic string provides the framework for a mathematically consistent quantum theory unifying gravity with the other particle interactions, there are many ways of constructing acceptable string models in four space-time dimensions: Calabi-Yau compactifications, covariant lattice constructions, orbifolds, fermionic constructions, etc.. In particular, $d = 4$ string vacua with $N = 1$ space-time supersymmetry seem a desirable starting point to tackle the problems of the cosmological constant and of the gauge hierarchy, and in contrast with vacua with $N > 1$ extended supersymmetry they allow for the existence of chiral fermions. To describe physics at energies $E << M_{Planck}$, one can then consider the corresponding effective theories, which are particular versions of $N = 1$, $d = 4$ supergravity with gauge and matter fields [3)]. At present, only for a restricted class of string models the resulting low-energy theory has been fully worked out (barring possible string-loop and non-perturbative effects): the simplest examples [4),5)], to be considered in the following, are $N = 1$, $d = 4$ heterotic superstrings which can be constructed in terms of free world-sheet fermions with periodic-antiperiodic boundary conditions [2)]. This has to be contrasted with the case of Calabi-Yau manifolds and orbifolds, where one has only some approximate results [6),7)].

For the string models under consideration, the building blocks of the corresponding $N = 1$, $d = 4$ supergravities are the following:

- The gravitational supermultiplet, whose physical degrees of freedom are the spin-2 graviton $g_{\mu\nu}$ and the spin-$\frac{3}{2}$ gravitino ψ^μ.
- The vector supermultiplets, whose physical degrees of freedom are the spin-1 gauge bosons A^a_μ and the spin-$\frac{1}{2}$ gauginos λ^a.
- The chiral supermultiplets $(\Phi^i, \tilde{\Phi}^i)$, whose physical degrees of freedom are complex spin-0 and two-components spin-$\frac{1}{2}$ fields. Some of them are always present in any model of this kind and can be obtained by a $Z_2 \times Z_2'$ projection from the corresponding theory with $N = 4$ supersymmetry ('untwisted' sector). One, the 'dilaton' supermultiplet $(S, \tilde{S})$, is singlet under the gauge group, and its spin-0 component contains the degrees of freedom of the dilaton and of the antisymmetric tensor. The other supermultiplets $(y^i, \tilde{y}^i)$ of the untwisted sector transform under real representations of the gauge group. In addition, if certain conditions are met [2)], there can be additional matter supermultiplets $(z^\alpha, \tilde{z}^\alpha)$, transforming in chiral representations of the gauge group ('twisted' sector): their origin is a purely string phenomenon, connected with the requirement of modular invariance.

Neglecting higher-derivative terms, which are not relevant for the present considerations, the effective theory is completely determined by two functions, the gauge kinetic function $f_{ab}(\Phi^i)$ and the Kähler potential $\mathcal{G}(\Phi^i, \overline{\Phi}_i)$: f_{ab} is an analytic function of the scalar fields, transforming as a symmetric product of adjoint representations of the gauge group; $\mathcal{G}$ is a real scalar function of the scalar fields and their conjugates, and is conventionally written as the sum of a 'kinetic' and a 'superpotential' part

$$\mathcal{G}(\Phi^i, \overline{\Phi}_i) = J(\Phi^i, \overline{\Phi}_i) + \log |g(\Phi^i)|^2. \tag{1}$$

All the different terms in the supergravity lagrangian (kinetic terms, masses, interactions) can be written [3)] in terms of f_{ab} and $\mathcal{G}$.

To illustrate the structure of the low-energy theory, let us consider a toy example [8),9)], corresponding to string models with no massless states in the twisted sector and gauge group $H \equiv SO(m_0) \times SO(m_1 - 2) \times SO(m_2 - 2) \times SO(m_3 - 2)$ [$m_0 = 2 + 4k$, $m_0 + m_I = 0$ $(mod.8)$, $m_0 + m_I \neq 8, 16$ $(I = 1,2,3)$, $m_0 + m_1 + m_2 + m_3 = 44$]. Their entire massless spectrum and interactions can be obtained from the underlying $N = 4$ model based on the gauge group $G = SO(38) \times U(1)^3$. Exploiting the fact that in $N = 4$ supergravity the lagrangian is completely determined by the gauge group, one can easily derive [4)] the functions f_{ab} and $\mathcal{G}$ of the corresponding $N = 1$ theory. The f_{ab} function has the universal form [6),7),4),5)]

$$f_{ab} = \delta_{ab} S, \tag{2}$$

and the scalar fields parametrize the Kähler manifold

$$\mathcal{M} = \frac{SU(1,1)}{U(1)} \times \prod_{A=1}^{3} \frac{SO(2,n_A)}{SO(2)\times SO(n_A)}. \tag{3}$$

The first submanifold is associated to the S field, while the remaining three submanifolds are parametrized by the $n_A = m_0(m_A-2)+(m_B-2)(m_C-2)+1$ $[(ABC)=P(123)]$ fields

$$y^{i_A} \equiv (y^1_A, y^{\hat{\imath}_A}) \quad (i_A = 1,\ldots,n_A;\ \hat{\imath}_A = 2,\ldots,n_A). \tag{4}$$

The y^1_A $(A=1,2,3)$ are singlets under the gauge group H, while the $y^{\hat{\imath}_A}$ $(A=1,2,3;\ \hat{\imath}_A = 2,\ldots,n_A)$ transform as elements of the coset $SO(38)/H$. In a convenient parametrization, the kinetic function for the scalar fields is

$$\begin{array}{ll} J = J_0 + \sum_{A=1}^{3} J_A, & J_0 = -\log(S+\overline{S}), \\ J_A = -\log Y_A, & Y_A = (1 - y^{i_A}\overline{y}_{i_A} + \frac{1}{4}|y^{i_A}y^{i_A}|^2), \end{array} \tag{5}$$

where bars denote complex conjugation and repeated indices are summed, and the superpotential is given by

$$g = g_{susy} = c_{\hat{\imath}_1\hat{\imath}_2\hat{\imath}_3} y^{\hat{\imath}_1} y^{\hat{\imath}_2} y^{\hat{\imath}_3}, \tag{6}$$

with the coefficients $c_{\hat{\imath}_1\hat{\imath}_2\hat{\imath}_3}$ proportional to the structure constants of $SO(38)$.

Once the Kähler potential $\mathcal{G}$ is known, one can examine the structure of the classical potential, given by the well-known supergravity formula [3)]

$$V = e^{\mathcal{G}}(\mathcal{G}^i \mathcal{G}^{-1}{}_i{}^j \mathcal{G}_j - 3) + V_{D-terms}, \tag{7}$$

where we follow the standard notation $\mathcal{G}_i \equiv \partial\mathcal{G}/\partial\Phi^i$, $\mathcal{G}^i \equiv \partial\mathcal{G}/\partial\overline{\phi}_i$, etc.. In terms of a generic superpotential g, one obtains, after some algebra [8),9)]

$$V = V_0 + \sum_{A=1}^{3}(V_A^I + V_A^{II}) + V_{D-terms}, \tag{8}$$

$$V_0 = e^J |g - (S+\overline{S})g_S|^2, \tag{9}$$

$$V_A^I = e^J Y_A \left(|g_{i_A}|^2 - |y^{i_A} g_{i_A} - g|^2\right), \tag{10}$$

$$V_A^{II} = e^J \left|\overline{y}_{i_A}\left[g_{i_A} - \frac{1}{2}\overline{y}_{i_A}(y^{j_A}g_{j_A} - g)\right]\right|^2, \tag{11}$$

$$V_{D-terms} = e^{J_0}\left[\sum_{A=1}^{3} \frac{\overline{y}_{i_A} T^{a i_A}{}_{j_A} y^{j_A}}{Y_A}\right]^2. \tag{12}$$

One important thing to note is that, apart from V_A^I, all the addenda contributing to V are manifestly positive semi-definite. If we now take $g = g_{susy}$, the remarkable homogeneity properties of the superpotential (6), $y^{i_A}g_{i_A} - g = 0$ $(A=1,2,3)$, imply that all minima of the potential have zero vacuum energy and correspond to $<g_{i_A}> = 0$ $(\Rightarrow <g> = 0)$. Since the gravitino mass is given by the general formula [3)]

$$m^2_{3/2} = \left\langle e^{\mathcal{G}}\right\rangle = \left\langle \frac{|g|^2}{(S+\overline{S})\prod_{A=1}^{3} Y_A}\right\rangle, \tag{13}$$

this in turn implies that supersymmetry is unbroken. However, there is the possibility of gauge symmetry breaking along the many flat directions: for instance, purely real directions along which at least one of the $SO(m)$ factors of the gauge group remains unbroken.

On the way towards the formulation of realistic string models, one has to face a very severe problem, the obvious experimental fact that supersymmetry is broken in Nature. The ultimate goal would be to find a four-dimensional string model with broken supersymmetry, vanishing (or acceptably small) cosmological constant and the possibility of a stable hierarchy $m_W << M_{Planck}$. Several suggestions for the breaking of supersymmetry have been made, but there are not yet completely satisfactory results. In string models with supersymmetry broken at the Planck scale [1),2)], the flat background is in general unstable, and a huge cosmological constant is generated at one loop. Attempts to break supersymmetry perturbatively on strings with a small gravitino mass have been frustrated by the appearance of a series of no-go theorems [10)]. An alternative possibility is offered by non-perturbative effects like gaugino [11)] or gravitino [12)] condensation. In all these scenarios one would like to find some new symmetry, different from supersymmetry [see, e.g., ref.[13]], protecting the cosmological constant and the gauge hierarchy.

Given the present lack of knowledge, we adopt here a more modest approach. We assume as a working hypothesis that in the effective theory supersymmetry breaking is spontaneous (the pairing of bosonic and fermionic degrees of freedom is preserved by the integration over the massive modes) and parametrized, with respect to the case of unbroken supersymmetry, by a superpotential modification, with no changes in the geometrical structure (Kähler metric) of the theory. This is the case, for example, in the effective theory [14] of string models where supersymmetry is spontaneously broken by a coordinate-dependent compactification [15]. The requirement of vanishing vacuum energy at the minima allows only superpotential modifications which preserve the positivity of the scalar potential, and this severely restricts the possible choices. In the following we shall consider the following ansatz [8),9)] (inspired by the 'gaugings' of $N = 4$ supergravity [16]), which might have a superstring generalization)

$$g = g_{new} = \omega(S) F(y_1^1, y_2^1, y_3^1) + g_{susy}, \tag{14}$$

where

$$F(y_1^1, y_2^1, y_3^1) = (1 + y_1^1)(1 + y_2^1)(1 + y_3^1), \tag{15}$$

and, for the time being, $\omega(S)$ is an arbitrary analytic function of S. More general superpotential modifications can be found in ref.[8),9),14)]. Using the modified superpotential of (14-15) and the general expression (8-12) for the scalar potential, it is a simple exercise to work out the minimization conditions. Minima with broken supersymmetry correspond to $\langle \omega F \rangle \neq 0$, in which case one must also have

$$< g_{susy\, \hat{\imath}_A} > = 0 \quad (\hat{\imath}_A = 2, \ldots, n_A;\ A = 1,2,3), \quad \Rightarrow \quad \langle g_{susy} \rangle = 0, \tag{16}$$

$$< \omega - (S + \overline{S}) \omega_S > = 0, \qquad < V_{D-terms} > = 0, \tag{17}$$

and

$$< y_A^1 (1 + \frac{1}{2} y_A^1) + \frac{1}{2} y^{\hat{\imath}_A} y^{\hat{\imath}_A} > \; = 0 \quad (A = 1,2,3). \tag{18}$$

Notice that there are infinite field configurations which satisfy (16-18), each of which corresponds to a different value for the gravitino mass (13): we have therefore a new realization of the so-called 'no-scale' models [17].

One major problem of all supergravity models, and in particular of no-scale models, is the stability of the classical potential against radiative corrections. To address this question, one can consider the one-loop effective potential

$$\begin{aligned} V_{1-loop} &= V \\ &+ \tfrac{1}{64\pi^2} Str \left[\Lambda^2 M^2 + \Lambda^4 \log\left(1 + \tfrac{M^2}{\Lambda^2}\right) - M^4 \log\left(1 + \tfrac{\Lambda^2}{M^2}\right) \right], \end{aligned} \tag{19}$$

where Λ is the cut-off of the effective theory, of the order of $\alpha'^{-1/2}$, α' being the string tension. In a general $N = 1$ supergravity, there are quadratic divergences proportional to

$$\begin{aligned} Str M^2 &\equiv \textstyle\sum_{J=0}^{3/2} (-1)^{2J} (2J+1) M_J^2 \\ &= 2 e^{\mathcal{G}} [N_{TOT} - 1 - \mathcal{G}^i (\mathcal{G}^{-1})_i{}^m R_m{}^n (\mathcal{G}^{-1})_n{}^j \mathcal{G}_j] + \ldots, \end{aligned} \tag{20}$$

where the dots stand for D-term contributions and N_{TOT} is the total number of chiral superfields. A very important geometrical quantity, which plays a crucial role [18] in the evaluation of $Str M^2$, is the Ricci tensor of the scalar manifold, $R_m{}^n \equiv (\partial/\partial\Phi^m)(\partial/\partial\overline{\Phi}_n) \log[\det \mathcal{G}_i{}^j]$. It must be stressed that $R_m{}^n$ involves the fourth derivatives of the function $\mathcal{G}$: this means that the simple knowledge of the quadratic part of the kinetic function J is not enough to discuss the quadratic divergences of the effective theory. Another important thing to note is that $R_m{}^n$ depends *only* on the kinetic function J and not on the superpotential g. In the simple model under consideration, the Kähler manifold (3) is a product of Einstein spaces, with

$$R_{i_A}{}^{j_A} = R_A \mathcal{G}_{i_A}{}^{j_A} = R_A J_{i_A}{}^{j_A} \quad (A = 0,1,2,3), \tag{21}$$

$$R_0 = 2 \quad \text{and} \quad R_A = n_A \quad (A = 1,2,3). \tag{22}$$

Therefore $Str M^2$ has a very simple expression in terms of the curvatures R_A

$$\begin{aligned} Str M^2 &= 2(N - 1 - \textstyle\sum_{A=0}^{3} a_A R_A) e^{\mathcal{G}} + \ldots, \\ a_A &= \mathcal{G}^{i_A} (\mathcal{G}^{-1})_{i_A}{}^{j_A} \mathcal{G}_{j_A} \quad (A = 0,1,2,3). \end{aligned} \tag{23}$$

If the D-terms are vanishing, as it is the case for the models we consider,

$$a_0 = V_0 e^{-\mathcal{G}}, \qquad a_A = 1 + (V_A^I + V_A^{II})e^{-\mathcal{G}} \quad (A = 1,2,3), \tag{24}$$

so that at any minimum of the potential $a_0 = 0$, $a_A = 1$ $(A = 1,2,3)$, and thus $StrM^2 = 0$. It should be stressed that the vanishing of $StrM^2$ at any classical minimum does not depend on the details of the superpotential modification (14-15), but is guaranteed by the geometry of the manifold $\mathcal{M}$, by the positivity of the potential and by the fact that the Goldstino $\tilde{\eta}$ does not have a component along the dilatino direction, $\langle \mathcal{G}_S \rangle = 0$.

To investigate further the structure of the effective potential (19), one needs to evaluate $StrM^{2n}$ for $n > 1$, which requires the explicit computation of the fermion and boson mass matrices [9]. We restrict ourselves here, for simplicity, to the gauge-symmetric minimum

$$<y_A^1> = 0, \quad <y^{i_A}> = 0, \quad (i_A = 2, \ldots, n_A;\ A = 1,2,3). \tag{25}$$

One can easily check that at this minimum the non-singlet scalar fields y^{i_A} do not get any mass. However, supersymmetry is broken and the goldstino is $\tilde{\eta} = (\sum_{A=1}^{3} \tilde{y}_A^1)/\sqrt{3}$. Moreover, one finds

$$StrM^4 = 12 m_{\tilde{S}}^2 m_{3/2}^2 \propto <\omega_{SS}>, \tag{26}$$

which vanishes for $<\omega_{SS}> = 0$, corresponding to $m_{\tilde{S}} = 0$ and $m_{S_R}^2 = m_{S_I}^2 = m_{3/2}^2$. This is always the case for a linear $\omega(S) = \alpha + \beta S$, where α and β are complex parameters such that $Re(\overline{\alpha}\beta) > 0$. Indeed, one can show that not only $StrM^4 = 0$, but also $Str f(M^2) = 0$ for any function f. Therefore, the whole one-loop cosmological constant vanishes at the classical minimum (25). This result can be extended [19] to the gauge symmetry-breaking minima satisfying (16-18), with the only difference that supersymmetric masses for scalar and vector supermultiplets are generated from the superpotential couplings and from the gauge couplings.

It would be of main importance to understand the previous result in terms of some symmetry of the effective theory (and of the underlying string theory). With this in mind, it is useful to take a closer look

at the particle spectrum at the minimum (25). Neglecting the gauge non-singlet fields, which remain massless, the remaining fields can be paired in the following spectrum multiplets

	Multiplet	Mass
(1)	$(g_{\mu\nu}; \tilde{S})$	0
(2)	$(S; \psi_T^\mu)$	$m_{3/2}$
(3)	$(\eta; \psi_L^\mu \equiv \tilde{\eta}\gamma^\mu)$	$m_{3/2}$
(4)	$(\eta'; \tilde{\eta}')$	$m_{3/2}$
(5)	$(\eta''; \tilde{\eta}'')$	$m_{3/2}$

(27)

where the entries in lines (2) and (3) combine to form a boson-fermion multiplet containing transverse and longitudinal degrees of freedom of the massive gravitino plus the S and η scalars. In (27), $\tilde{\eta}'$ and $\tilde{\eta}''$ are two linear combinations of the singlets $\tilde{y}_A^1$ $(A = 1,2,3)$ orthogonal to $\tilde{\eta}$, and η, η', η'' the corresponding linear combinations of spin-0 fields. Since all physical masses are either zero or equal to the gravitino mass, and the massive fermionic and bosonic degrees of freedom are equal in number, then $Str\ M^{2n} = (n_B^m - n_F^m) m_{3/2}^{2n} = 0$. Now it is clear that the vanishing of all the supertraces is due to a fermion-boson mass degeneracy in the spectrum. The spectrum classification (27) looks like an exact $N = 1$ supersymmetry in flat space-time background, however some of the spectrum multiplets in (27) are not multiplets of $N = 1$ supersymmetry. The degrees of freedom in the graviton-dilaton sector are classified by $N = 1$ supersymmetry in the multiplets $\{(g_{\mu\nu}, \psi_T^\mu), (S, \tilde{S})\}$, by the spectrum symmetry in the multiplets $\{(g_{\mu\nu}; \tilde{S}), (S; \psi_T^\mu)\}$, with masses $m = 0$ and $m = m_{3/2}$ respectively. It is important to note that the spectrum symmetry $\mathcal{S}$ connects states whose spins differ by 3/2: it can thus be interpreted as the combination of a discrete supersymmetry transformation $\mathcal{Q}$ with a permutation $\mathcal{P}$ of the physical degrees of freedom of $g_{\mu\nu}$ and S. Such a symmetry can be naturally formulated in a more 'stringy' language, recalling that the spectrum of any superstring model is given as a direct product of left- and right-moving oscillators. In the light-cone gauge, S and $g_{\mu\nu}$ appear in a symmetric form

$$(S, g_{\mu\nu}) \leftrightarrow |\mu>_L |\nu>_R \qquad (\mu, \nu = 3, 4). \tag{28}$$

The real and imaginary parts of the S field are associated with

$$(\Phi_D, \Phi_A) \equiv (|3>_L|3>_R + |4>_L|4>_R,\ |3>_L|4>_R - |4>_L|3>_R), \tag{29}$$

while the two physical degrees of freedom of the graviton $g_{\mu\nu}$ are associated with

$$(\Phi_1, \Phi_2) \equiv (|3>_L|3>_R - |4>_L|4>_R,\ |3>_L|4>_R + |4>_L|3>_R). \tag{30}$$

We want to introduce an operator $\mathcal{P}$ that interchanges the degrees of freedom $\{\Phi_D, \Phi_A\}$ and $\{\Phi_1, \Phi_2\}$. A natural solution, consistent with Lorentz invariance, is a left-permutation $\mathcal{P}:\ |3>_L \leftrightarrow |4>_L$. If $\mathcal{Q}$ is the operator associated to a discrete supersymmetry transformation, then the symmetry of the spectrum is simply generated by the operator $\mathcal{S} = \mathcal{QP}$. In the case of unbroken $N = 1$ supersymmetry, $\mathcal{S}$, $\mathcal{Q}$ and $\mathcal{P}$ are all symmetries of the spectrum. When supersymmetry is broken as in the previous example, $\mathcal{Q}$ and $\mathcal{P}$ are both broken, but $\mathcal{S}$ survives as an unbroken symmetry of the spectrum. It would be very interesting to examine whether the symmetries $\mathcal{P}$ and $\mathcal{S}$ are just accidental symmetries of the spectrum or they can be extended, in a Lorentz-covariant formulation, to the interactions of the effective theory and perhaps to the full string theory: the particularly symmetric way in which the graviton and the dilaton supermultiplets appear in the underlying two-dimensional sigma model give encouraging signals, and the problem is under scrutiny [20]. Unfortunately, string models where supersymmetry is spontaneously broken by a coordinate-dependent compactification [15] give superpotential modifications [14] which differ from (14-15), but other possibility for spontaneous supersymmetry breaking at the string level could exist. An indication in this sense is given by the existence of $d = 4$ one-loop partition functions with Atkin-Lehner symmetry and boson-fermion degeneracy at each mass level [21].

All the results presented here can be extended [8],[9] to more realistic models containing a massless twisted sector $(z^\alpha, \bar{z}^\alpha)$ with a net non-zero number of chiral families. In that case the analysis is complicated by the presence of additional superpotential terms of the form 'yzz' and

by the fact that the Kähler manifold for the scalar fields is no longer symmetric and mixes non-trivially twisted and untwisted sectors [5], but with some more effort one reaches equivalent conclusions.

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