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DUALITY AND SUPERSYMMETRY BREAKING IN STRING THEORY

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Abstract

Target-space duality is incorporated in previously proposed effective actions describing non-perturbative supersymmetry breaking in string theory via gaugino condensation. Duality-preserving vacua with broken supersymmetry and fixed unified coupling constant do generically occur. The question of the vanishing of the cosmological constant is also briefly addressed.

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Despite many interesting attempts to explain it, the origin of supersymmetry breaking in superstring theory still remains an unsolved puzzle. The problem is of crucial importance, not only because supersymmetry is not observed in Nature, but also because the breaking of supersymmetry should eliminate some embarrassing, gravitationally coupled, particles such as the dilaton and the scalar moduli. The removal of the corresponding flat directions should also provide “unique” vacuum expectation values for the scalar fields that determine the fundamental constants of the theory as well as the geometry and topology of the compactified space. Among these constants, one appears to stand out: once supersymmetry is broken, no satisfactory mechanism is known which protects the cosmological constant from acquiring unacceptably large values. Most probably, the resolution of the cosmological constant problem lies deep in the realm of quantum gravity; since our considerations treat gravity classically, we shall not regard here the vanishing of the cosmological constant as a necessary constraint. We shall return to the cosmological constant problem at the end of this paper.

In this note we shall consider a mechanism for supersymmetry breaking based on the following two main theoretical ingredients:

- 1) *Target-space duality* [1]. This symmetry, which appears to be valid at any string-loop order [2], acts on the moduli of the target-space (in its simplest form it transforms the radius R into α'/R). It has been already invoked for providing a minimal compactification scale [3] and, more generally, a minimal observable length in string theory [4]. In the context of $N = 1$ supergravity with a single modulus superfield $T = R^2/\alpha' + iB/\alpha' + \dots$, the $R \leftrightarrow \alpha'/R$ symmetry extends to the full group $SL(2, \mathbb{Z})$ of modular transformations acting on T [5]. The implications of this symmetry for the Kähler potential and the superpotential of $N = 1$ supergravity have been discussed at length in [5, 6].

- 2) *Non-perturbative gluino condensation* [7]. Strongly-coupled “hidden” gauge sectors occur in practically all “realistic” compactifications of heterotic superstring theory. In this context, gaugino condensation has been considered as a promising mechanism for observable supersymmetry breaking.

Gaugino condensation has been extensively analysed within the context of globally supersymmetric Yang-Mills (SYM) theory without matter. The occurrence of a gaugino condensate can be argued on the basis of Witten’s index [8] and of effective Lagrangians [9]. It can also be explicitly computed by instanton techniques (see [10] for a review) with the result:

$$\langle \lambda\lambda \rangle = \gamma (\Lambda_{\text{SYM}})^3 \quad (1)$$

where γ is a calculable constant, and Λ_{SYM} is the SYM strong interaction scale. In the 2-loop approximation, it is given by:

$$\Lambda_{\text{SYM}}^{2\text{-loop}} = \mu g^{\beta_1/2\beta_2^2} \exp\left(\frac{-8\pi^2}{\beta_0 g^2}\right), \quad (2)$$

where μ is the scale at which the coupling constant g is defined, and β_0, β_1 are the beta-function coefficients: $\beta(g) = -\frac{\beta_0}{(4\pi)^2} g^3 - \frac{\beta_1}{(4\pi)^4} g^5 + \dots$. Global supersymmetry remains unbroken in the presence of a non-zero condensate, in agreement with Witten’s index theorem [8].

Non-Abelian gauge sectors of compactified heterotic superstring theory have some very distinct features. They are coupled to supergravity, with supersymmetry promoted to a local symmetry. As a long-distance phenomenon, the formation of a gaugino condensate is expected to remain unaffected by gravitational interactions. The most important difference, though, between the global and local supersymmetry is that, in the latter case, a non-zero gaugino condensate breaks local supersymmetry, providing a seed for observable supersymmetry breaking [7]. Moreover, the gauge coupling g as well as the

scale μ in eq.(2) correspond to vacuum expectations of some scalar fields.

The tree-level gauge coupling constant can be expressed in terms of the vacuum expectation value of the scalar component of the dilaton superfield S : $g^2 = \frac{1}{4}(\text{Re}S)^{-1}$, which is also Newton's constant in string units [11]. The scale μ corresponds to the scale below which the "running" of the coupling constant takes place, and can be evaluated order by order in the loop expansion either by direct string computation [12] or by using the effective field theory approximation [3]. The latter method, based on the fact that "running" is entirely an infrared, low-energy effect [13], gives a particularly clear physical picture of μ as a geometrical quantity. The scale of "running" in string theory depends therefore on the moduli of the compact space. These moduli correspond to scalar components of massless chiral superfields T_i , hence $\mu = \langle \mu(T_i) \rangle$.

The one-loop corrections to the coupling constant at momentum scale Q are dominated by contributions of particles with masses smaller than the string mass $M_s = (\alpha')^{-1/2}$. These contributions include Kaluza-Klein modes and winding modes, whose spectra are determined by the geometry of compactification. One typically obtains [3]:

$$g^{-2}(Q) = g^{-2} + \beta_0(4\pi)^{-2} \log(Q/\mu)^2 + \text{const} \cdot (M_s/\mu)^2 + \dots \quad (3)$$

with the parameter μ depending on the moduli of the compact manifold. The simplest case of a toroidal compactification characterized by one radius R is particularly illuminating. For $R \gg (\alpha')^{1/2}$, the light spectrum consists of Kaluza-Klein modes only, and the summation over all excitations gives $\mu = 1/R$. For $R \ll (\alpha')^{1/2}$, the light spectrum consists entirely of "dual" winding modes, and the corresponding $\mu = R/\alpha'$. This is how duality manifests itself in the loop corrections to the gauge coupling [3].

We shall now show how to incorporate target-space modular invariance [5,6] in the framework of effective Lagrangians for gaugino condensation [9,14,15,16]. Assuming that compactification preserves $N = 1$ supersymmetry, the effective theory at scales lower

than μ should be of the general type of $N = 1$ supergravity coupled to gauge and matter fields [17]. The most general theory of this type is described, in the notation of [17,18], by the Lagrangian:

$$\mathcal{L} = -\frac{1}{2} \left[e^{-K/3} S_0 \bar{S}_0 \right]_D + [S_0^3 w]_F + [f_{ab} W_a^\alpha \epsilon^{\beta\gamma} W_b^\beta]_F, \quad (4)$$

where K is the Kähler potential, w is the superpotential and f_{ab} is the gauge function. All these functions depend upon the fundamental chiral matter superfields which have conformal weight 0, while the gauge superfield W has conformal weight $3/2$. Finally, S_0 is the so-called chiral compensator superfield of conformal weight 1 which ensures the correct weights of each of the three terms in eq.(4).

In the following, we shall limit our considerations to the simplest models of compactification parametrized by a single modulus $T = R^2/\alpha' + iB/\alpha' + \dots$ [11]. In this case, as shown in [11], the tree-level Kähler potential is of the no-scale type [19], and is given by:

$$K = -\log(S + \bar{S}) - 3\log(T + \bar{T} - |A_i|^2) \quad (5)$$

where A_i represent untwisted matter fields (for generalizations to twisted sectors, see [6]). The gauge function is $f_{ab} = \delta_{ab} S$.

In order to construct the effective theory of gaugino condensation we follow [9,14,15,16] and introduce the composite superfield:

$$U = \delta_{ab} W_a^\alpha \epsilon^{\beta\gamma} W_b^\beta \quad (6)$$

The superfield U has conformal weight 3 and is thus not adequate for the general treatment of [17]. This problem is easily coped with by introducing the superfield:¹

$$\tilde{U} = U/S_0^3 \quad (7)$$

¹An equivalent procedure, using the formalism of [20], has been considered in [16].

of canonical weight 0. Note that, in terms of $\tilde{U}$, the gauge kinetic term of eq.(4) amounts to the following contribution to the tree-level superpotential:

$$\tilde{w}^{tree} = S\tilde{U} \quad (8)$$

The full superpotential w will be an analytic function of S , T , $\tilde{U}$ and of the scalar matter fields.

It is known that all physical quantities depend on K and w only through the combination:

$$\mathcal{G} = K + \log(w\bar{w}) = K + (\log S\tilde{U} + h.c.) + \dots \quad (9)$$

Modular invariance of $\mathcal{G}$ guarantees modular invariance of the theory. Under the modular $SL(2, Z)$ transformations:

$$T \rightarrow \frac{aT - ib}{icT + d} \quad (ad - bc = 1), \quad (10)$$

whereas S is invariant and $A_i \rightarrow (icT + d)^{-1} A_i$. From eqs.(5), (9) and (10) we see that the invariance of $\mathcal{G}$ requires w to transform as:

$$w \rightarrow \frac{w}{(icT + d)^3} \quad (11)$$

up to a T -independent phase factor. In other words, the superpotential w should be a modular form of weight -3 , possibly with a non-trivial multiplier system [21]. This constrains $\tilde{U}$ to be a modular form of weight -3 .

The one-loop modification of the superpotential should be S independent, and of the form:

$$\tilde{w}^{1-loop} = \tilde{U} f(\tilde{U}, T) \quad (12)$$

with f a modular function of weight 0. General arguments based on the fulfilment of anomalous Ward identities require f to contain the term $\frac{\beta_0}{96\pi^2} \log \tilde{U}$ [9,14]. In order to

compensate for the lack of modular invariance, we have to multiply $\tilde{U}$ under the argument of this logarithm by a modular form of weight 3. There exist six classes of modular forms of weight 3, one for each multiplier system [21]. Moreover, in order to avoid unphysical singularities, we demand that the argument of the logarithm has neither zeroes nor poles in the upper half-plane of complex iT . This leaves us with the unique form:

$$\eta^8(T) = e^{-\pi T/2} \prod_{m=1}^{\infty} (1 - e^{-2m\pi T})^6 \approx e^{-\pi T/2} (1 - 6e^{-2\pi T} + \dots) \quad (13)$$

where $\eta(T)$ is the Dedekind eta function. At this point, it is convenient to introduce the chiral superfield Y defined by $Y^3 = \tilde{U}$, so that Y has canonical modular weight -1 . The effective Kähler potential and superpotential that incorporate modular-invariant one-loop corrections are then given by:

$$K = -\log(S + \bar{S}) - 3\log(T + \bar{T} - |Y|^2 - |A_i|^2) \quad (14)$$

$$w = 3bY^3 \log[c e^{S/36} Y \eta^2(T)], \quad (15)$$

where $b = \frac{\beta_0}{96\pi^2}$ and c is an arbitrary constant. The choice of the Y dependence of K is such as to reproduce the kinetic term of [9,14] for the gaugino condensate.²

We note that, owing to the asymptotic behaviour of the η function, see eq.(13), our effective theory contains, at large compactification radius (large T), a correction to the effective Lagrangian proportional to $[T\delta_{ab}W_a^\alpha e^{c\beta}W_\beta^\dagger]$. The existence of such a term is a consequence of the Green-Schwarz anomaly cancellation mechanism [22]. As far as the gauge coupling constant is concerned, this term corresponds to a correction proportional to $(M_s R)^2$, as anticipated in eq.(3) [3]. Finally, the logarithmic renormalization of the coupling constant implied by our superpotential, eq.(15), will be shown to agree with the explicit results of [23].

²Strictly speaking complete agreement would require an extra factor $(S + \bar{S})^{-1/3}$ in front of $|Y|^2$, which does not affect the results substantially, while considerably complicating the calculations.

We shall now proceed with the computation of the scalar potential in our theory and with the determination of possible supersymmetry-breaking candidate vacua. As seen from eq.(4), the scalar potential depends on the chiral compensator S_0 . On the other hand, the choice of S_0 determines the normalization of the gravitational action [17,18]: $\mathcal{L}_{\text{grav}} \propto [S_0 \bar{S}_0 e^{-K/3}]_{\theta=\bar{\theta}=0} R$. Since $(S + \bar{S}) = \alpha' M_P^2$, the conventional normalization corresponds to the choice:

$$[S_0 \bar{S}_0 e^{-K/3}]_{\theta=\bar{\theta}=0} = [S + \bar{S}]_{\theta=\bar{\theta}=0}. \quad (16)$$

After substituting eqs.(14-16) into eq.(4), and eliminating the auxiliary fields, we obtain the following result for the scalar potential:

$$V(S, T, Y) = \frac{3(S + \bar{S})(Y\bar{Y})^2}{(T + \bar{T} - Y\bar{Y})^2} \left\{ 3b^2 \left| 1 + 3 \log[c e^{S/3b} Y \eta^2(T)] \right|^2 \right. \\ \left. + \frac{Y\bar{Y}}{(T + \bar{T} - Y\bar{Y})} \left| S + \bar{S} - 3b \log[c e^{S/3b} Y \eta^2(T)] \right|^2 \right. \\ \left. + 6b^2 Y\bar{Y} \left[2(T + \bar{T}) |\eta(T)/\eta(T)|^2 + \eta(T)/\eta(T) + \overline{\eta(T)/\eta(T)} \right] \right\}. \quad (17)$$

In eq.(17), for the sake of simplicity, we have set to zero all matter fields A_i .

The potential of eq.(17) can be shown to be explicitly modular invariant and therefore to be stationary at the self-dual points $T = 1$ and $T = e^{i\pi/6}$. It has a zero-value stationary point at $Y = 0$ and, for a wide range of the parameter c , one minimum at $Y \neq 0$ which will be discussed below. Notice that, in order to minimize the potential, the phase of the condensate Y has to be aligned with respect to T and S , in the sense that $c e^{S/3b} Y \eta^2(T)$ has to be real and positive.

A numerical analysis of the potential $V(S, T, Y)$ proves the existence of one non-trivial minimum at $Y \neq 0$, which corresponds to a non-vanishing gaugino condensate $(\lambda\lambda)$. We

have minimized the potential as a function of the magnitude of Y at fixed real S and c and complex T . The resulting value of V (up to an S -dependent factor) is then plotted as a 3D (and the corresponding contour) plot as a function of complex T for various values of $x = (S + \bar{S})/b$ and c , see Figs.1-3. For large enough c (e.g. $c = 1.7$ as in Figs.1-3) the self-dual point $T = 1$ turns from an absolute minimum to a saddle as x is increased. Conversely, the point $T = e^{i\pi/6}$ is a saddle at small x and becomes a minimum as x is increased. Only for small enough x , however, does the value of the potential at the minimum ($T = 1$) become negative, hence energetically favoured over the trivial minimum at $Y = 0$. The cross-over of the minima occurs around $x = 0.6$ with potential close to zero all along the boundary of the fundamental region (see Fig.2 and Fig.4). Finally, if c is too small (say less than 1.0), the potential has no local minimum in Y at fixed T other than the trivial minimum at $Y = 0$.

Postponing momentarily minimization with respect to S , we conclude that target-space duality remains unbroken ($T = 1$) by gaugino condensation. Supersymmetry, however, is dynamically broken, as can be shown by computing the derivatives of the function $\mathcal{G}$: $\mathcal{G}_i \neq 0$.³ The gravitino mass turns out proportional to $\langle \lambda\lambda \rangle / M_P^2$, assuming flat gravitational background; we shall return to the discussion of the gravitational field equations later on. Some phenomenological implications of such a supersymmetry breaking mechanism -along the lines of Refs. [20], [24]- will be discussed elsewhere.

We shall now compare our result for the gaugino condensate with the results of explicit instanton calculations in flat space-time [10]. The value of the quantity appearing under the logarithm in eq.(17): $c e^{S/3b} Y \eta^2(T)$ turns out to be S and T independent and equal

³It is easy to see that duality implies a relation of proportionality between $\mathcal{G}_T$ and $\mathcal{G}_Y$ at the fixed points, leaving open the possibility of duality-preserving, supersymmetry-breaking vacua. A similar relation for derivatives of V implies that stationary points in Y are automatically stationary in T at the fixed points.

to $e^{-1/3}$ to a very good approximation. The corresponding condensate is given by

$$|\langle \lambda \lambda \rangle| = |\langle U \rangle| \propto |e^{-S/b_\eta} \eta^{-6} (T) S_0^3| = (S + \bar{S})(T + \bar{T})^{-3/2} |\eta|^{-6} |e^{-(S+\bar{S})/2b_\eta}|, \quad (18)$$

where we have used the definition $U = S_0^3 Y^3$ and eq.(16). On the other hand, instanton calculations give [10]:

$$|\langle \lambda \lambda \rangle| \propto (\Lambda_{SYM}^{2-loop})^3 = \mu^3 g^{-2} \exp\left(\frac{-24\pi^2}{\beta_0 g^2}\right), \quad (19)$$

where we have used the fact that $\beta_1/2\beta_0^2 = -\frac{2}{3}$ in SYM. Eqs.(18) and (19) agree⁴ provided that we identify

$$\mu = (T + \bar{T})^{-1/2}, \quad (20)$$

and take the coupling at scale μ as:

$$g^{-2}(\mu) = 2(S + \bar{S}) + \beta_0(2\pi)^{-2} \log |\eta(T)|. \quad (21)$$

At large T , eqs.(20) and (21) are consistent with [3] and [22], respectively. For arbitrary T , eq.(21) is consistent with the results of [23].

We finally come to two intimately related questions: minimization of the potential in S and the cosmological constant. With the normalization of our potential corresponding to the choice of the chiral compensator of eq.(16), minimization in the dilaton field S is coupled to the gravitational field equations obtained by varying the action $\Gamma \propto \int d^4x \sqrt{g} [(S + \bar{S})R + V]$ with respect to the metric tensor $g_{\mu\nu}$. If we eliminate the metric through its own equations of motion, we find that the potential to be minimized has an extra $(\text{Re}S)^{-2}$ factor in front and that, consequently, the potential runs away to $-\infty$ for $S \rightarrow 0$. Since our analysis neglects all sorts of quantum gravitational effects, it is not clear that the classical equations of motion should be used for the metric. If, instead, we put

⁴The fact that the resultant S -dependence of the condensate is consistent with the value of Λ_{SYM}^{2-loop} of eq.(2) was pointed out to us by P. Binétruy and M.K. Gaillard.

$R = 0$ by hand by invoking some unspecified S -independent mechanism for suppressing the cosmological constant, the minimization in S fixes $\text{Re}S$ to a small positive value which can be read from Fig.4 to be around 0.12b ($x=0.24$). At this value of x the vacuum with $T = 1$ is favoured and the potential is negative (the critical point for a negative minimum can be seen, from Fig.4, to lie around $x = 0.60$). To summarize, it is difficult to make definite statements about S without having a specific model for the cosmological constant. For instance, we did not take into account other effects that may occur in the fully-fledged theory. Possible sources of non-perturbative contributions to the superpotential include wormholes [25], gravitino condensation due to gravitational anomalies [26,27], and other condensates due to a hidden gauge group of the type $G_1 \times G_2 \times \dots$ [23]. Altogether, these effects may stabilize S and lead naturally to a vanishing cosmological constant.

In conclusion, we have found evidence that the incorporation of modular invariance in the non-perturbative dynamics of gluino condensation leads in general to a unique determination of couplings and compactification parameters in string theory through the removal of all tree-level flat directions. While the dynamically favoured values of the radii turn out to be the fixed points of the modular group, the coupling constant cannot be completely determined without a well-defined mechanism for suppressing the cosmological constant. Generally speaking, our results lend support to the claim [3] that compactifications of ten-dimensional superstrings to $D = 4$ with string-size compactification radii, and with large couplings at the unification scale, are energetically preferred.

If this picture should be further confirmed, superstring theory would be just a step short of fixing the fundamentals constants of Nature, the vanishing of the cosmological constant still remaining as the ultimate mystery.

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Figure Captions

Fig. 1: a) 3D plot of the minimum of the potential with respect to Y as a function of complex T at $x = 0.24$, $c = 1.7$

b) Contour plot of Fig.1a)

Fig. 2: a) 3D plot of the minimum of the potential with respect to Y as a function of complex T at $x = 0.6$, $c = 1.7$

b) Contour plot of Fig.2a)

Fig. 3: a) 3D plot of the minimum of the potential with respect to Y as a function of complex T at $x = 3.0$, $c = 1.7$

b) Contour plot of Fig.3a)

Fig. 4: x -dependence of the potential (17) at the points $T = 1$, $e^{i\pi/6}$, at the minimum in Y . Minima and saddles swap around $x = 0.6$, where the potential is very flat (see Fig.2) along the boundary of the fundamental region and close to 0.

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FIGURE 1a

FIGURE 1b

FIGURE 2b

FIGURE 2a

FIGURE 3b

FIGURE 4