CERN-TH.5883/90

NUB-3005

HIDDEN SECTORS WITH HIDDEN MATTER

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Abstract

Using previous results on non-perturbative effects in massive supersymmetric Yang-Mills theories, we extend recent considerations on duality invariant gaugino condensation, within the string effective action approach, to the formation of hidden matter field condensates. The masses of the condensing matter fields are determined by some string background parameters corresponding to vacuum expectation values of additional gauge singlet hidden matter fields. The vacuum degeneracy of these fields is shown to be lifted by non-perturbative condensation effects.

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CERN-TH.5883/90

NUB-3005

October 1990

The so-called “hidden” gauge sectors occur in practically all “realistic” compactifications of heterotic superstring theory. The “observable” matter representations containing quarks and leptons are singlets with respect to these hidden gauge groups. The couplings between both sectors have a purely gravitational character.

The existence of hidden sectors may have some interesting phenomenological consequences. In particular, non-perturbative condensation of hidden gauginos provides a promising mechanism for supersymmetry breaking [1]. Not much attention has been paid so far¹ to the hidden matter fields which are present in many interesting models, including the orbifold compactifications [3,4] and the flipped $SU(5)$ model [5]. In this letter, we shall discuss some dynamical aspects of such non-minimal hidden sectors.

Non-perturbative effects in globally supersymmetric gauge theories, including supersymmetry breaking, have been extensively studied in the past – see [6] for a review. In particular, Witten’s index theorem [7] provides a powerful criterion for dynamical supersymmetry breaking. According to this theorem, global supersymmetry is *not* broken in pure SYM and in SYM coupled to real matter representations, at least in the case of massive matter fields. Non-perturbative formation of the gaugino and matter condensates can be understood in terms of the effective theory of gauge singlet composite supermultiplets. The effective Lagrangians for pure SYM and massive SQCD have been constructed in [8] and [9], respectively.

The effective Lagrangian analysis of pure SYM was later extended in [10,11,12,13] to *locally* supersymmetric non-Abelian gauge sectors of compactified heterotic superstring theory [14]; in this case, a non-zero gaugino condensate breaks local supersymmetry [1]. More recently, the target space modular invariance [15,16], which reflects the underlying string symmetry under duality transformations [17], has been incorporated into the effective Lagrangian description of gaugino condensation in pure SYM [18,19,20]. Specifically, in [19], the gaugino condensation has been described by a duality-invariant effective action for the gauge-singlet gaugino bound states coupled to the fundamental

¹For a notable exception, see [2].

fields like the dilaton and moduli, whereas in [18] the gaugino condensate has been replaced by its vacuum expectation value to yield a duality-invariant action that depends on the fundamental fields only. We shall comment on the relation between these two approaches later in the text.

In the heterotic superstring theory, all fundamental parameters are dynamically determined by the vacuum expectation values (VEVs) of some scalar fields. The tree-level gauge coupling constant can be expressed in terms of the VEV of the scalar component of the dilaton superfield: $g^2 = \frac{1}{4}(\text{Re}S)^{-1}$ [21]. This important VEV determines also the Planck mass: $M_P^2 = 32\pi(\text{Re}S)/\alpha'$. The string analogue of the one-loop, running coupling constant may depend on the moduli of the compact space [18,19,22,23]. These moduli correspond to scalar components of massless chiral superfields T_k , hence $g = (g(S, T_k))$.

In the context of heterotic superstring compactification, the existence of massive matter fields in the hidden sector with masses possibly much below the Planck mass is related to the so-called stringy Higgs effect [24]. Specifically in orbifold compactification there exists, in addition to the moduli T_k of the internal space, another type of moduli – the Wilson line background fields [4,25]. These correspond to some chiral matter fields A with vanishing potential in the low-energy effective action following from string perturbation theory. Non-vanishing VEVs $\langle A \rangle$ break the hidden gauge group G to a subgroup H , giving masses to the hidden matter fields Q via the cubic superpotential of the form $w_m \propto AQQ$. To be specific, we shall consider in detail one example of a non-minimal hidden sector: $H = SU(N)$ SQCD with M massive flavours $Q \oplus \bar{Q}$ in the representations $N \oplus \bar{N}$; our discussion can be readily extended to any SYM with massive matter. In this case, the mass term is equal to $w_m = -\sum_{i,j=1}^M A^i Q^i \bar{Q}^j$.² Here, for the sake of simplicity, we assume that the fields A , Q and $\bar{Q}$ belong to the untwisted sector of superstring compactification; otherwise, some additional, moduli-dependent factors [16] would be present in the mass term, as required by the duality symmetry.

²Hereafter, the $SU(N)$ gauge indices will be suppressed.

The effective action of superstring SQCD depends on the light degrees of freedom represented by the gauge-singlet supermultiplets. We introduce the composite superfields $Y^3 = W_\alpha W^\alpha / S_0$ representing the gaugino bound states [8,12,13,19], and $\Pi_j^i = Q^i \bar{Q}_j$ ($i, j = 1, \dots, M$) representing the matter bound states [9]. The chiral compensator S_0 [26,27] has been included in the definition of the superfield Y^3 in such a way as to fix its conformal weight equal to the canonical weight 0 [19]. The action depends also on the dilaton S and the two types of moduli: the masses A_i^j ($i, j = 1, \dots, M$) and the internal background fields T_k . Here, we shall limit our considerations to the simplest models of compactification parametrized by a single modulus T related to the size of the internal space. In this case, the effective Lagrangian must be invariant under the $SL(2, Z)$ modular group, generated by the duality transformation $T \rightarrow 1/T$ and the shift $T \rightarrow T + i$ [15]. The superfields A, Π and Y^3 transform under the modular transformations $T \rightarrow \frac{aT+b}{cT+d}$ with modular weights $-1, -2$ and -3 , respectively, whereas S is inert under the modular transformations.

The Lagrangian of the effective theory of composite supermultiplets has the usual supergravity [26] form:

$$\mathcal{L} = -\frac{1}{2} [e^{-K/3} S_0 \bar{S}_0]_D + [S_0^3 w]_F + \text{c.c.}, \quad (1)$$

where K is the Kähler potential and w is the superpotential.³ The effective superpotential, as dictated by Ward identities [9] and by the modular invariance [19], is:

$$w = \frac{1}{32\pi^2} Y^3 \log \{ e^{32\pi^2 S} [c \eta(T)]^{6N-2M} Y^{3N-3M} \det \Pi \} - \text{tr } 4\Pi, \quad (2)$$

where c is an unknown constant and η is the Dedekind eta function.

There exists a simple method of analysing the weak coupling limit of the effective theory. In the limit $g \rightarrow 0$ ($\langle \text{Re} S \rangle \rightarrow \infty$), the Planck mass $M_p \rightarrow \infty$, therefore gravity decouples and the model behaves like the familiar *globally* supersymmetric QCD, with

³The $SU(N)$ gauge kinetic term is contained in the effective superpotential of the composite theory [8,12,13,19], cf. eq.(2).

the superpotential given by eq.(2) [9]. As expected from Witten's index theorem [7], and as confirmed by the effective Lagrangian analysis of [9], global supersymmetry is *not* broken. Therefore the minimum of the potential occurs at $\partial w / \partial Y = h_Y = 0$ and $\partial w / \partial \Pi = h_\Pi = 0$, where Y and Π are the scalar components, and h_Y and h_Π are the auxiliary components, of the corresponding superfields. These conditions can be rewritten in the superfield form:

$$\frac{\partial w}{\partial Y} = \frac{\partial w}{\partial \Pi} = 0. \quad (3)$$

These equations can be used in order to eliminate the composite superfields Y and Π , giving rise to a “truncated” Lagrangian which depends on the “elementary” superfields A, S and T only. This truncated Lagrangian is analogous to the Lagrangians discussed in [11,18] in the case of pure SYM. Such a procedure of searching for a meta-stable vacuum can be completely self-consistent only if the “truncated” equations of motion admit solutions with large $\langle \text{Re} S \rangle$. Unfortunately, this is not the case, as we will see later on. In fact, the effective potential obtained in this way will not have a minimum stable with respect to the fluctuations of the dilaton field. Thus as far as the dilaton VEV is concerned, our analysis of SQCD will encounter the problems that have been discussed before [11,18,19] in the case of pure SYM.⁴ In such a situation, we shall adopt the strategy of [18,19], and assume the existence of an unspecified mechanism for fixing the dilaton VEV in the weak coupling regime; the region of strong coupling constant is not as interesting from the phenomenological point of view, since it does not fit into any perturbative unification scheme. We shall therefore consider $\langle \text{Re} S \rangle$ as a fixed, *large* parameter. This will justify eliminating the composite fields.

The equations of motion for the composite fields, obtained by applying eq.(3) to the superpotential of eq.(2), are:

$$\frac{1}{32\pi^2} Y^3 = (32\pi^2 e)^{M/N-1} [c \eta(T)]^{2M/N-5} [\det A]^{1/N} e^{-32\pi^2 S/N}, \quad (4)$$

⁴For some new ideas how to stabilise the dilaton VEV, see [28,29].

$$\Pi = \frac{1}{32\pi^2} Y^3 A^{-1}. \quad (5)$$

These equations provide the superstring generalization of the well-known [6,9] expressions for the gaugino and scalar condensates:

$$\frac{1}{32\pi^2} \langle \lambda\lambda \rangle = \Lambda^3 \left[\det \frac{m}{\Lambda} \right]^{1/N}, \quad (6)$$

$$\langle Q\bar{Q} \rangle = \frac{1}{32\pi^2} \langle \lambda\lambda \rangle m^{-1}, \quad (7)$$

where the strong interaction scale $\Lambda \propto \exp[-8\pi^2/(3N - M)g^2]$ (cf. with eq.(16)). The truncated superpotential, obtained by inserting eqs.(4) and (5) into the superpotential of eq.(2), can be written as

$$w = \Omega(S) h(T) [\det A]^{1/N}, \quad (8)$$

where

$$\Omega(S) = -N e^{-32\pi^2 S/N}, \quad (9)$$

$$h(T) = (32\pi^2 e)^{M/N-1} [c\eta(T)]^{2M/N-6}. \quad (10)$$

We shall discuss here the case of only one Wilson line background field A , assuming $A_j^i = A\delta_j^i$; the case of a larger number of Wilson lines can be studied in a similar way. The Kähler potential of the composite theory can be constructed by a straightforward extension of the results of [9,19]. It is easy to see that for the composite VEVs proportional to $e^{-32\pi^2 S/N}$, as in eqs.(4) and (5), the terms that depend on Y and Π are negligible for sufficiently large $\langle \text{Re}S \rangle$. Hence the Kähler potential of the truncated theory can be approximated by the tree-level form [21]:

$$K = -\log(S + \bar{S}) - 3\log(T + \bar{T} - \text{tr}A^\dagger A). \quad (11)$$

Now we are ready to investigate the vacuum structure of the low energy effective string theory with massive matter fields in the hidden sector, by studying the associated

scalar potential. Spontaneous breakdown of local supersymmetry occurs if at least one of the auxiliary fields $h_Z \propto \partial w / \partial Z + w \partial K / \partial Z$, $Z = S, T, A$, acquires a non-vanishing VEV. The superpotential of eq.(8) and the Kähler potential of eq.(11) yield the following scalar potential [26]:

$$V = \frac{|w|^2}{3S_R T_R^3 (1 - |\tilde{A}|^2)^2} \left\{ \left| T_R \frac{h'}{h} \right|^2 + 2T_R \left(\frac{M}{N} - 3 \right) \text{Re} \left(\frac{h'}{h} \right) + \left(\frac{M}{N} \right)^2 |\tilde{A}|^{-2} + \frac{3k(S, \bar{S})}{(1 - |\tilde{A}|^2)} \right\}, \quad (12)$$

where

$$k(S, \bar{S}) = |1 - S_R \Omega' / \Omega|^2. \quad (13)$$

Here we use the notation $S_R = 2\text{Re}S$, $T_R = 2\text{Re}T$ and introduce the variable $|\tilde{A}|^2 \equiv |A|^2/T_R$ which is inert under the duality transformations. The requirement of kinetic energy positivity restricts $|\tilde{A}|^2$ to the domain of $|\tilde{A}|^2 < 1$.

The dependence of the scalar potential on the dilaton field is determined by the function $\Omega(S)$, eq.(9). As already anticipated before, the potential under consideration has no stable minimum with respect to the dilaton field S . In the absence of detailed knowledge of the dilaton dynamics, we shall consider the dilaton VEV, and the corresponding values of the functions Ω and k , as free parameters. In particular, we shall relax the constraint $k(S, \bar{S}) = |1 + 32\pi^2 S_R / N|^2 > 1$ for $S_R > 0$, which is due to the particular form of the function $\Omega(S)$, and consider arbitrary $k \geq 0$.

The potential has no singularities in the upper half-plane of complex iT , and it is bounded from below provided that $M < 3N$, i.e. if the model is asymptotically free; in this case $V \rightarrow \infty$ as $T_R \rightarrow 0, \infty$. The function $V(S, T, A)$ can be shown explicitly to be modular invariant and therefore to be stationary at the self-dual points $T = 1$ and $T = e^{i\pi/6}$. These points may correspond to either minima, maxima, or saddle points, depending on M/N and on the value of the parameter k . There can possibly exist some other, duality-breaking, stationary points. In the following we shall provide some examples for either of the two cases of unbroken versus broken duality symmetry.

The main interest of this note is the investigation of the vacuum structure of the theory with respect to the chiral matter field $\tilde{A}$, since its VEV determines the masses of the gauge non-singlet matter fields. For this discussion, it is convenient to determine the limiting behaviour of the potential at $\tilde{A} \rightarrow 0$ and $\tilde{A} \rightarrow 1$. The dominant term in the region of small $|\tilde{A}|$, up to an overall (multiplicative, positive) constant, is:

$$V \rightarrow \frac{M^2/N^2 \Omega(S)^2}{S_R T_R^{3-M/N}} |\eta(T)|^{4M/N-12} |\tilde{A}|^{2(M/N-1)} + \mathcal{O}(|\tilde{A}|^{2M/N}), \quad |\tilde{A}| \rightarrow 0. \quad (14)$$

At $\tilde{A} = 0$, the scalar potential diverges for $N > M$, however it goes to zero (to a constant) for $M > N$ ($M = N$). In the region of $|\tilde{A}| \approx 1$, we find:

$$V \rightarrow \frac{|\Omega(S)|^2}{S_R T_R^{3-M/N}} |\eta(T)|^{4M/N-12} \frac{k(S, \tilde{S})}{(1 - |\tilde{A}|^2)^3} + \mathcal{O}[(1 - |\tilde{A}|^2)^{-2}], \quad |\tilde{A}| \rightarrow 1. \quad (15)$$

The asymptotic behaviour of the potential at $\tilde{A} = 1$ depends therefore on the parameter k which is assumed to be determined by some unknown dilaton dynamics. As long as $k > 0$, the residue of the cubic pole is positive; hence there always exists a stable minimum with respect to $\tilde{A}$. However, if $k = 0$, the potential is dominated by the quadratic pole which may have a negative residue for T in the vicinity of the self-dual points; in such a case, the potential shows a runaway behaviour $V \rightarrow -\infty$ as $|\tilde{A}| \rightarrow 1$.

It is clear from the discussion of asymptotic limits that the dynamics depend on the parameters of the theory, in particular on the number of colours N and the number of flavours M . In the following, we shall discuss in detail four distinctly different cases.

(i) $M = 0$, $\tilde{A} = 0$: this is the pure SYM case already discussed in [18,19]. For small values of k , the minimum of the potential with respect to T is away from the self-dual points which are a maximum and a saddle point at $T = e^{i\pi/6}$ and $T = 1$, respectively. For instance, if $k = 0$, there are an infinite number of degenerate minima, located at $T \approx 1.2$ and its modular transformed images. There is a phase transition at $k \approx 0.5$ where the duality-breaking extrema disappear, and for $0.5 \lesssim k \lesssim 2.0$ the minimum is at the self-dual point $T = 1$, whereas $T = e^{i\pi/6}$ becomes a saddle point. Finally, for $k \gtrsim 2.0$, $T = 1$ becomes a saddle point and the minimum is found at $T = e^{i\pi/6}$. In all these cases

local supersymmetry is spontaneously broken.

This is the right place to make some comments on the results of the two original works [18,19] that investigated the relation between target space modular invariance and non-perturbative gaugino condensation in pure SYM. The starting point of the analysis presented in [18] has been the “truncated” superpotential of the form written in eq.(8) with $M = 0$. In order to stabilise the potential with respect to the dilaton field, the function $\Omega(S)$ of eq.(9) has been modified to include an additive constant [11]. The dilaton VEV is then stabilised at a point that corresponds to the parameter $k = 0$. Hence the minima of the potential [18] with respect to the modulus T are located at the duality-breaking points discussed in case (i) for $k = 0$. On the other hand, the analysis of [19] employs the fully-fledged composite theory which, upon the truncation, yields the function $\Omega(S)$ given *exactly* by eq.(9). The dilaton VEV has been considered as a free parameter; however, as already mentioned, for the function $\Omega(S)$ given by eq.(9) the value of k is always bounded by $k > 1$. Hence the minima of the potential [19] with respect to the modulus T are located at the self-dual points discussed in case (i) for $k \gtrsim 0.5$. The fact that the duality-preserving minima and saddle points are interchanged for the dilaton VEV corresponding to $k \approx 2.0$ has also been pointed out in [19]. To summarise, one can interpolate between the duality-breaking vacua of [18] and the duality-preserving vacua of [19] by changing the unknown parameter k . We conclude that the issue of target space modular symmetry breaking is inseparably tied up with the dilaton dynamics.

(ii) $M < N$: since the scalar potential diverges for $\tilde{A} \rightarrow 0$, the model is dynamically “repelled” from its singular zero-mass limit [9]. The existence of a stable minimum with respect to $\tilde{A}$ depends on the dilaton dynamics. There is no minimum for $k = 0$, however, for $k > 0$, there always exists one non-trivial minimum with $\tilde{A} \neq 0$. Local supersymmetry is spontaneously broken at this minimum. For small values of k , $|\tilde{A}| \approx 1$, and the potential is negative. For large values of k , $|\tilde{A}| \ll 1$ and the potential is positive. Hence one can in principle fine tune the parameter k , i.e. the string coupling constant,

in such a way that the cosmological constant vanishes. It is also interesting to consider the limit $M/N \rightarrow 0$ keeping the string coupling constant fixed. The potential minimum is then at $\tilde{A} \approx 0$, and the case of pure SYM is rediscovered. This means that case (ii) contains the limiting case (i) in a consistent way.

To be more specific, we discuss the case $N = 10$, $M = 5$. The target space modular invariant generalization [18, 19, 20] of the strong interaction scale $\Lambda \propto \exp[-8\pi^2/(3N - M)g^2(S, T)]$ is:

$$\Lambda_{SQCD}^{1-loop} \approx \frac{\exp[-16\pi^2 S_R/(3N - M)]}{|\eta(T)|^2 T_R^{1/2}}. \quad (16)$$

Let us consider first $S_R = 10^{-3}$ which corresponds to $k \approx 1.06$, assuming the original form of the function $\Omega(S)$ given by eq.(9). The potential, being negative at this point, is then minimized at $T = 1$ and $|\tilde{A}| \approx 0.55$ (in Planck units). The corresponding strong interaction scale is $\Lambda \approx 1$. Next, take $S_R = 0.2$. A positive-energy minimum is now found at $|\tilde{A}| \approx 0.04$ and $T = e^{i\pi/6}$ with $\Lambda \approx 0.3$.⁵ The masses of matter fields turn out to be smaller than the strong interaction scale for these values of S_R . This justifies the effective Lagrangian description with M matter fields. The situation changes however for very large values of S_R , i.e. for smaller couplings.⁶ For example, if $S_R = 5$, then $|\tilde{A}| \approx 10^{-3}$ whereas $\Lambda \approx 10^{-14}$. This means that, for sufficiently small gauge couplings, the matter fields tend to decouple. The effective theory should then be described by pure SYM.

(iii) $M = N$: for $k = 0$, there is no stable minimum with respect to $\tilde{A}$. For $0 < k \lesssim 1$, however, there exists one non-trivial minimum at $\tilde{A} \neq 0$, with either positive, negative, or zero energy. For $k \gtrsim 1$ there is only one minimum at $\tilde{A} = 0$. In all these cases local supersymmetry is spontaneously broken.

(iv) $N < M < 3N$: in the case of a larger number of flavors than the number of colors, one always finds a zero-energy supersymmetric minimum at $\tilde{A} = 0$. Thus the

⁵The zero-energy minimum is found at $S_R \approx 5 \times 10^{-3}$, $T = 1$, $|\tilde{A}| \approx 0.45$, $\Lambda \approx 1$.

⁶N.B. $V \rightarrow 0^+$ as $S_R \rightarrow \infty$.

massless limit of SQCD corresponds in this case to a stable supersymmetric vacuum. The condensates of eqs.(4) and (5) are equal zero at this minimum, and the VEV of the modulus T remains undetermined. In addition, for sufficiently small k , there exists another minimum at $\tilde{A} \neq 0$, with either positive, negative, or zero energy and broken supersymmetry.

Finally, let us briefly comment on the case $M \geq 3N$. For $M = 3N$, the one-loop SQCD beta function is zero, and the superpotential of eq.(8) is cubic in the Wilson line background field A . This is analogous to the cubic superpotential at string tree level among untwisted matter fields. The VEV of the modulus T remains undetermined. Nevertheless, the original degeneracy in the Wilson line background field is lifted due to non-perturbative effects. For $M > 3N$, asymptotic freedom is lost and one does not expect any condensations to occur at all. At least formally, the minimum of the potential with respect to A is at $A = 0$, where the two condensates are also zero and supersymmetry remains unbroken.

In conclusion, we have considered the duality invariant formation of gaugino and matter field condensates within the framework of effective supergravity action for the hidden sectors of heterotic superstring compactification. As a representative example, we have considered in detail $SU(N)$ SQCD with M matter fields. This model is described by three types of (complex) parameters, all related to the VEVs of some dynamical fields: the gauge coupling constant, the compactification radius and the masses of the matter fields. The vacuum degeneracy of these fields can be lifted due to the non-perturbative formation of condensates. For $M < N$, non-vanishing masses are dynamically generated for the matter fields and local supersymmetry is spontaneously broken. Simultaneously, the compactification scale is determined to be of the order of the Planck scale. For small values of the gauge coupling constant, the masses are small compared to the Planck mass, thus leaving the possibility of obtaining the desired hierarchy between the string scale and the scale of supersymmetry breaking. On the other hand, for $M > N$, there is no dynamical mass generation, the condensates are

vanishing and supersymmetry remains unbroken, while the compactification radius is left undetermined. All these features are in agreement with the previous investigations of globally supersymmetric Yang-Mills theories. The effective supergravity action does not share though the singular zero-mass behaviour, for $M < N$, with the global case. This difficulty is avoided since the masses correspond now to dynamical fields with the potential energy determined by the supergravity action. We have also clarified the relation between the effective Lagrangian formalism utilizing dynamical condensates and its truncated version with the composite fields replaced by their VEVs. We have repeatedly emphasized the fact that many phenomenologically interesting problems, including the issue of target space modular symmetry breaking, are very difficult to address without improving the present understanding of the dilaton dynamics.

Acknowledgements

We acknowledge interesting discussions with A. Font, L.E. Ibáñez, M. Peskin, F. Quevedo, S. Theisen and G. Veneziano.

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