CONSTANT VERSUS FIELD DEPENDENT GAUGE COUPLINGS IN SUPERSYMMETRIC THEORIES

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We briefly discuss the differences between considering the gauge coupling as a constant or as a field, the dilaton, in $N = 1$ supersymmetric theories. We emphasize the differences regarding supersymmetry breaking. Recent developments on the nonperturbative dynamics of these theories provide new ideas on the induced dilaton potential and its stabilization.

Thanks to the work of Seiberg and many others, the understanding of supersymmetric theories has improved considerably during the past three years\textsuperscript{1-5}. In particular, the nonperturbative dynamics determining the possible phases of these theories has been very well understood. This is important for understanding issues such as chiral symmetry breaking, supersymmetry breaking, the vacuum structure etc. In the context of string theory, having $N = 1$ supersymmetric theories as their low energy effective theories, this progress should be reflected on the possibility to address the most important obstacles for the theory to make contact with low energy physics, namely, lifting the vacuum degeneracy and breaking supersymmetry.

Superstring theories include always in their spectrum a massless field called the dilaton $S$ which provides the bare gauge coupling. It is not only massless but has an exactly flat potential in perturbation theory, and therefore it is one of the many ‘moduli’ of the theory. Nonperturbative effects generically lift this potential and the dilaton will get a mass. Depending on the nature of these effects, the mass of the dilaton will be determined by the supersymmetry breaking scale and then is expected to be small, or is fixed at the Planck scale and therefore the dilaton does not appear in the low-energy spectrum of the theory. In the first scenario the nonperturbative effect responsible for breaking supersymmetry is the same that fixes the dilaton whereas the second scenario is in two steps: a Planck scale effect fixes the dilaton and a low energy effect breaks supersymmetry. These two different scenarios will gener-

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ally be differentiated in the low energy action by having the gauge coupling either constant or field dependent. The two-steps scenario fits with the gauge mediated supersymmetry breaking scenario recently revived \[6-8\] whereas the one-step scenario fits with the more standard gravity mediated supersymmetry breaking scenario.

The recent studies in supersymmetric gauge theories consider the gauge coupling as a constant and therefore it applies to the two-steps scenario directly. The search for models that break supersymmetry therefore has a direct application to string theory, only on this scenario. It is then interesting to ask what are the implications of the new understanding of supersymmetric gauge theories for the more standard scenario where the dilaton survives at low energies.

Previous attempts to understand these issues were based on a limited knowledge of supersymmetric theories, and the typical models considered included a string hidden sector with several gauge group factors and matter charged under one of the factors. Concentrating on the dilaton field, the standard superpotentials that emerge in this case are of the form

\[W = \sum_i A_i e^{-a_i S}\]  

These models generally do not break supersymmetry in the $S$ sector, they have a supersymmetric minimum at finite values of $S$, but also have a runaway solution to zero coupling ($S \to \infty$). This behaviour at infinity has been argued on very general grounds by Dine and Seiberg some time ago \[14\]. They argue that at zero coupling the theory should be free and therefore the potential should vanish there. This is a source of a cosmological problem pointed out by Brustein and Steinhardt \[16\]. Taking the superpotential above, being so steep, if the dilaton field starts at any value, it may never end up at the local minimum with non-vanishing coupling but will roll all the way to the runaway vacuum.

On the other hand, besides the field $S$, there is usually another modulus, the field $T$ measuring the size of the compact space. This field has also a flat potential to all orders in perturbation theory that gets lifted by nonperturbative effects. The properties of $T$ and $S$ are very similar and this similarity was actually at the origin of the proposal of $S$ duality, given that there existed a better established $T$ duality. There are even some models that have the symmetry $S \leftrightarrow T$. In the same way that $S$ represents the string coupling, $T$ represents the coupling of the underlying $2D$ sigma model. Curiously the potential for the $T$ field found in simple examples, blows up for large values of $T$. Unlike what it was naively expected, that it should runaway to the weak coupling limit $T \to \infty$. This may be understood in the following way: the gauge coupling in $10D$, $g_{10}$ is related to the gauge coupling in $4D$, $g_4$ by
\(1/g_4^2 = R^6/g_{10}^2\) where \(R\) is the size of the compact \(6D\) space and so it is the real part of the field \(T\). A large value of \(T\) is a large value of \(R\), combined with a relatively small value of \(g_4\) implies a very strong string coupling in \(10D\), therefore the blowing-up of the potential is a strong string coupling effect, not controlled in string perturbation theory where the calculation was performed.

On the other hand the potentials for the \(S\) field seem to behave very different from those for the \(T\) field. It is then valid to question the general assumption that the potential for the \(S\) field runs away to \(\infty\) and study different alternatives to the sum of negative exponentials of the equation above.

In this talk we will present several models illustrating the difference between constant and dilaton dependent gauge couplings as well as different examples where the dilaton potential does not runaway to infinity. We also argue that the inclusion of field-dependent gauge couplings can qualitatively change whether or not a given model spontaneously breaks supersymmetry. The main difference is due to the additional requirement of extremizing the superpotential with respect to the coupling-constant field. For instance, it can happen that a supersymmetry-breaking ground state for fixed gauge coupling becomes supersymmetric once the coupling constant is allowed to relax to minimize the energy. In particular we show that most of the models with dynamical supersymmetry breaking, when the gauge coupling is field dependent, do not break supersymmetry. Furthermore, we find that the opposite of this is also possible, supersymmetry can be unbroken for fixed gauge coupling, but breaks down once the gauge coupling is considered as a field.

An example on the difference between having field dependent or independent gauge couplings is the simplest case of gaugino condensation for a pure gauge theory having a simple gauge group and no matter multiplets. In this case, for constant gauge couplings, gauginos condense without breaking supersymmetry. The reason is that the gaugino condensate is given as the lowest component of a chiral superfield \(U = \langle \lambda \lambda \rangle\) and a non-vanishing value for the lowest component does not break supersymmetry. On the other hand, once a field dependent coupling constant is introduced via a chiral field \(S\), whose real part gives the coupling constant \(\text{Re} \ S = 1/g^2\), a non-vanishing gaugino condensate will break supersymmetry because it will be proportional to the \(F\) term of the \(S\) field, and a nonvanishing \(F\) term breaks supersymmetry. However, the dynamics of the dilaton field for a single gaugino condensate in pure Yang-Mills theory has a runaway behaviour \(S \to \infty\) and the gaugino condensate vanishes \(U = \text{const.} e^{-aS} \to 0\). Therefore in both cases, field dependent and field independent coupling constant, supersymmetry is not broken. But in the first case gauginos condense whereas in the second case they do not condense.

We will now study potentials including matter fields and we will consider
$N = 1$ supersymmetric models with gauge group $SU(N_c)$. We represent the matter multiplets with chiral superfields, $Q^i_\alpha \in R$ (and $\tilde{Q}^i_\alpha \in \tilde{R}$), where ‘$i$’ is the flavour index, and ‘$\alpha$’ is the gauge index. The kinetic microscopic action for the model is given by

$$L_{\text{kin}} = \frac{1}{4} f \text{Tr} W_\alpha W_\alpha,$$

where $f$ is the gauge kinetic function and $W_\alpha$ the chiral gauge superfield and we take standard kinetic terms for the matter supermultiplets. At tree level in string theory on has $f = kS$ with $k$ the Kac-Moody level. The microscopic superpotential relating the matter supermultiplets is taken to vanish identically, $W(Q, \tilde{Q}) = 0$.

To determine the superpotential for the quantum ‘effective action’ which generates the irreducible correlation functions of the theory (as opposed, say, to the theory’s Wilson action) we study the operators whose correlations we wish to explore. Of particular interest, however, are those fields which can describe the very light scalar degrees of freedom of the model, since these describe the system’s vacuum moduli and symmetries. In the absence of a microscopic superpotential for the matter fields $Q$ and $\tilde{Q}$, these light degrees of freedom are described classically (and hence also to all orders of perturbation theory) by the $D$-flat directions, which parametrize the zeroes of the classical scalar potential. It is well known that these $D$-flat directions can be parametrized in terms of a suitably chosen set of gauge-invariant holomorphic polynomials. We take the arguments of the superpotential to be $W(U, M^i_j)$, where $M^i_j = \langle Q^i_\alpha \tilde{Q}^\alpha_j \rangle$, $U = \langle \text{Tr} W_\alpha W_\alpha \rangle$. Although the gaugino condensate field, $U$, does not similarly describe a $D$-flat direction, it is nonetheless convenient to keep it as an argument of the effective action. The superpotential is completely determined by the twin conditions of linearity and symmetry under the model’s global flavour symmetries. As was demonstrated in [18], the fact that $S$ only couples to the microscopic theory via the kinetic term implies, as an exact result, that the effective superpotential necessarily has the form

$$W = \frac{1}{4} US + f(U, M^i_j).$$

That is, $S$ can only appear linearly, and moreover only in the term $\frac{1}{4}US$. Second, the function $f(U, M^i_j)$ is determined by the various global chiral symmetries of the underlying supersymmetric gauge theory. In the absence of a superpotential for the matter fields, $Q^i_\alpha$ and $\tilde{Q}^i_\alpha$, the underlying gauge theory admits the classical global symmetry $SU(N_f)_L \times SU(N_f)_R \times U(1)_A \times U(1)_B \times U(1)_R$, of which the factors $U(1)_A \times U(1)_R$ are anomalous. Invariance of the effective superpotential under the anomaly-free symmetries implies the fields $M^i_j$ can appear only through the invariant combination $\det M$. (For $N_c < N_f$ we imagine the expectation value of the baryon operator, $B^{i_1 \cdots i_{N_c}} = \epsilon^{\alpha_1 \cdots \alpha_{N_c}} Q^{i_1 \alpha_1} \cdots Q^{i_{N_c} \alpha_{N_c}}$ to be minimized by zero). The two
anomalous symmetries, $U(1)_A$ and $U(1)_R$, then fix the form of the unknown function $f(U, \det M)$.

From these considerations, it is clear that $W$ has the general structure

$$ W = \frac{1}{4} U S + \frac{U}{32\pi^2} \left[ (N_c - N_f) \log \left( \frac{U}{\mu^2} \right) + \log \left( \frac{\det M}{\mu^{2N_f}} \right) + C_0 \right]. \quad (3) $$

Symmetry arguments cannot determine the constants $\mu$ and $C_0$. Indeed $C_0$ may be chosen to vanish through an appropriate choice for $\mu$.

Since $W$ is the superpotential for the effective action — as opposed to the Wilson action — the correct procedure for ‘integrating out’ fields is to remove them by solving their extremal equations for $W$, rather than by performing their path integral. Furthermore, for supersymmetric theories this should be done using the effective superpotential, $W$, rather than the effective scalar potential $V$. Performing this operation for the gaugino condensate $U$ one obtains

$$ W = c \left( \frac{\mu^{3N_c-N_f} e^{-8\pi^2 S}}{\det M} \right)^{1/(N_c-N_f)} = c' \left( \frac{\Lambda^{3N_c-N_f}}{\det M} \right)^{1/(N_c-N_f)} \quad (4) $$

where $c = -\frac{a}{12\pi^2} \exp \left( \frac{C_0 + a}{a} \right)$, $c' = c \exp \left( -\frac{C_0}{a} \right)$ and the second equality defines the RG-invariant scale, $\Lambda = \mu^{3N_c-N_f} e^{-8\pi^2 S/(3N_c-N_f)}$.

It is convenient to distinguish four different cases depending on the matter content: i) $N_f < N_c$, ii) $N_f = N_c$, iii) $N_f > N_c$ and iv) $N_f > 3N_c$.

In the first case, $N_f < N_c$, the only invariant are meson fields and a non-vanishing superpotential $W = c \left( \frac{\Lambda^{3N_c-N_f}}{\det M} \right)^{1/(N_c-N_f)}$ is obtained. Since the scale $\Lambda$ in terms of the coupling constant is given by $\Lambda = \mu^{3N_c-N_f} e^{-8\pi^2 S/(3N_c-N_f)}$ minimizing the superpotential with respect to $S$ gives a runaway behaviour $S \to \infty$ and $W_S \propto e^{-8\pi^2 S/(N_c-N_f)} \to 0$. As in the pure Yang-Mills case, a superpotential is dynamically generated but its minimum is at vanishing potential and a supersymmetric vacuum is obtained. However, for a field independent gauge coupling we do not extremize the superpotential with respect to $S$ and we will get a non-vanishing vacuum for finite value of $\det M$. So we have a runaway potential in the $M$ direction. Adding tree level terms as a function of $M$ cannot avoid the runaway potential for $S$ in the field dependent case but may avoid the runaway potential for $M$ in the constant case, fixing $M$ at a finite value.

Another interesting case is when the matter content is $N_f = N_c$. In this case the second term in eq.(3) vanishes and extremizing the superpotential
with respect to $U$ gives the quantum constraint $\det M = \Lambda^{2N_c}$, where we have taken the baryons v.e.v. to vanish. If $< M_{ij} > = 0$ then the quantum constraint will be satisfied only for a runaway dilaton field (i.e. vanishing $\Lambda$). This is always possible for a field dependent gauge coupling but it will not be satisfied for a finite value of the gauge coupling or a constant gauge coupling and, therefore, supersymmetry will be broken. Furthermore, we can add to the superpotential eq. (8) tree level terms like $M^a + M^b$ for the mesons, that do not destroy the symmetries yielding a finite value of $M$ and thus stabilizing the dilaton field through the quantum constraint.

For $N_f > N_c + 1$ the exponent in eq. (4) is positive. In this case, a runaway behaviour for the dilaton is no longer favoured since $W \rightarrow \infty$. However, there is always a solution with $M \rightarrow 0$ and a runaway of the superpotential ($W \rightarrow 0$) in the plane $S - M$ is again not avoided. This includes the selfdual region $\frac{3N_f}{2} < N_f < 3N_c$.

Finally consider $3N_f > N_c$ with all baryons minimized by zero. There are a number of criticisms which might be raised against using non-asymptotically free gauge theories and against the generation of a non-perturbative superpotential\footnote{Notice that the way this constraint is realized is different from the one assumed in reference 25 for instance, but this does not change any of the results of that paper.}. However, the weakness in the arguments lie, in general, in its making an insufficient distinction between the effective action and the Wilson action\footnote{The Wilson action, $S_w$, describes the dynamics of the low-energy degrees of freedom of a given system, and is used in the path integral over these degrees of freedom in precisely the same way as is the classical action. The Wilson action for SQCD at scales for which quarks and gluons are the relevant degrees of freedom would therefore depend on the fields $W_\alpha$, $Q_\alpha$ and $\bar{Q}_\alpha$. As a result, the vanishing of $\det(Q\bar{Q})$ would indeed preclude the generation of a superpotential of the type $e^{-8\pi^2 S/\det(Q\bar{Q})}$ within the Wilson action. By contrast, it is the effective action, $\Gamma$, which is of interest when computing the v.e.v.s of various fields. And it is $M_{ij} = < Q_\alpha \bar{Q}_\beta >$ which serves as an argument of $\Gamma$. Since the expectation of a product of operators is not equal to the product of the expectations of each operator, it need not follow that $\det M = 0$ when $N_c < N_f$. Extremizing with respect to $M_{ij}$, and substituting the result back into $W$ gives the super-}. Let us introduce a mass term $Tr(\mu M)$ for the quark fields in eq. (4) and make the mass $\mu$ dynamical, as it is always the case in string theory, by adding a trilinear term for the field $\mu$. Eq. (4) becomes then $W(M, \mu, S) = Tr(\mu M) + \frac{a}{2} Tr(\mu^2) + k \left( e^{-8\pi^2 S/\det M} \right)^{1/(N_c - N_f)}$, where $k = N_c - N_f$. Extremizing with respect to $M_{ij}$, and substituting the result back into $W$ gives the super-
potential \( W(\mu, S) = \frac{k}{N_f} Tr(\mu^3) + k' \left(e^{-8\pi^2 S/N_c} \det \mu\right)^{1/N_c} \), where \( k' = N_c \). If \( \mu^i_j \) were a constant mass matrix this last equation would give the superpotential for \( S \) in SQCD. It is noteworthy that so long as \( k' \neq 0 \) the result has runaway behaviour to \( S \to \infty \) regardless of the values of \( N_c \) and \( N_f \). We extremize, now, \( W \) with respect to the field \( \mu^i_j \), to obtain the overall superpotential for \( S \). The extremum is obtained for \( \mu^i_j = \left(-he^{-8\pi^2 S/N_c}\right)^{N_c/(N_f-3N_c)} \delta^i_j \), and the superpotential is then given by

\[ W(S) = k'' \left(he^{24\pi^2 S/N_c}\right)^{(N_f-3N_c)/N_f} = k'' \Lambda^3, \]

with \( k'' = (-1)^{3N_c/(N_f-3N_c)} (N_c - N_f/3) \). Notice that eq. (5) takes the simple form \( W \propto \Lambda^3 \) when expressed in terms of the renormalization group invariant scale and it is valid for all values of \( N_f \) and \( N_c \). Eq. (5) gives a positive exponential of \( S \) if \( N_f > 3N_c \) where the theory is not asymptotically free. When this is combined with the potential for another, asymptotically-free gauge group we obtain a superpotential of the form of eq. (1) with positive and negative exponentials and a non-trivial minimum can be found for \( S \).

The extremal condition for the dilaton \( W_S = 0 \) gives a runaway behaviour \( S \to \infty \) for \( 3N_c > N_f \) but for non-asymptotically free gauge group the equation \( W_S = 0 \) is satisfied only if the mass field \( \mu \) has a vanishing v.e.v., i.e. \( < \mu > = 0 \). Minimizing the superpotential with respect to \( \mu \), \( W_\mu = 0 \), gives two solutions:

\( \mu = 0 \) and \( \mu = \left(-he^{-8\pi^2 S/N_c}\right)^{3N_c-N_f} \). For asymptotically free gauge group \( 3N_c > N_f \) both solutions are equivalent in the runaway limit \( S \to \infty \). In this case both minima are continuously connected in the \( S-\mu \) plane. On the other hand, if \( 3N_c < N_f \) then the solution \( \mu = 0 \) and \( \mu = \left(-he^{-8\pi^2 S/N_c}\right)^{3N_c-N_f} \) are driven apart by a large value of \( S \) and one cannot continuously go from one minimum to the other one. The barrier between both minima increases exponentially with increasing \( S \). This property can play an important role in the evolution of the dilaton field for cosmology.

Notice that since it is the effective action which we use, rather than the Wilson action, one might worry whether our analysis is invalidated by the appearance of nonlocal terms or holomorphy anomalies. We argue that this is not the case for the solution where \( \mu^i_j \neq 0 \), since in this case the matter multiplets have masses and for scales below their mass the theory is a pure gauge theory, which has a gap due to confinement. Since holomorphy anomalies arise due to massless states, they cannot occur if the theory has a gap. The same need not be true for the potentially runaway solution, for which \( \mu^i_j = 0 \),

\(^c\text{We thank G. Dvali for interesting discussions on this point.}\)
since in this phase there are massless matter and gauge multiplets which can produce such anomalies.

An example of dynamical global supersymmetry breaking with constant gauge couplings, where supersymmetry can be restored by the incorporation of the dilaton (ie by the field dependence of the gauge couplings) is the canonical example of dynamical global supersymmetry breaking, the so-called 3-2 model of Affleck et al. In this example the gauge group is $SU(3) \times SU(2)$. The fundamental matter spectrum is such that the $SU(2)$ factor is quantum constrained. The quantum constraint is of the form $YZ = \Lambda_4^2$. It is shown in [4] that, if we suppose the condensation scale for the $SU(2)$ factor to be much greater than that for the $SU(3)$ factor, and if we suppose a certain superpotential in the microscopic theory, then the effective superpotential can be written as

$$W = XY + \lambda(YZ - \Lambda_4^2).$$  \hspace{1cm} (6)

One can easily see that the equation of motion for $X$ implies $Y = 0$ and that the equation of motion for the Lagrange multiplier $\lambda$ implies $YZ = \Lambda_4^2$. For the case of constant gauge couplings ($\Lambda_i =$constant), the relations cannot be simultaneously satisfied and supersymmetry is said to be dynamically broken. However, for the case of field dependent gauge couplings ($\Lambda_i = \mu_i e^{-c_i S_i}$), the relations are satisfied by the runaway vacuum $S \to \infty$, for which $\Lambda_i = 0$. Therefore we learn that in this model, supersymmetry is restored by a runaway dilaton if the gauge couplings are conceived to be field-dependent.

The opposite can also happen. We can have broken supersymmetry for field dependent gauge coupling but unbroken supersymmetry for field independent gauge coupling. For instance, consider gaugino condensate for two gauge groups with gauge kinetic function $f = f(S,T)$, as in string models when one-loop corrections are included. Once T-Duality is imposed on the theory the superpotential becomes a function of $S$ and $T$, $W = A_1(T)e^{-a_1 S} + A_2(T)e^{-a_2 S}$ is given in eq.(1) where the $A_i$ coefficients are now $T$ dependent. It is well known that this superpotential in local supersymmetry has a non-supersymmetric vacuum. Supersymmetry is broken through the auxiliary field of the moduli $T$, i.e. $F_T \neq 0$, and the dilaton gets a finite v.e.v. However, for a field independent gauge coupling, in this example, supersymmetry is not broken. In global supersymmetry, there may be cases where this also happens because the field equations of the field $S$ may turn out to be inconsistent with the other field equations, but so far we have not found explicit examples showing this property.

Finally we can also write down asymptotically free models which can produce positive exponentials from product group models. As an example, let us consider the $SU(2) \times SU(2)$ model of Intriligator, Leigh and Seiberg with
invariants $X$ and $Y$ for which the nonperturbative superpotential is:

$$W_{np} = \frac{\Lambda_5}{XY - \Lambda_2^2}$$

If we add to this superpotential a tree level one of the form $W_p = \lambda_1 X + \lambda_2 (Y - A)$ where $\lambda_{1,2}$ are Lagrange multiplier fields and $A$ is a constant. We can see that integration of $\lambda_{1,2}$ implies $X = 0, Y = A$ and so the superpotential becomes $W_{np} = A \Lambda_5^4 / \Lambda_2^2$. In terms of the gauge couplings $k_1 S$ and $k_2 S$, where $k_{1,2}$ are the Kac-Moody levels of each of the two $SU(2)$ factors, this is proportional to $\exp(8\pi^2 (k_2 - k_1) S)$ and therefore, for $k_2 > k_1$ we have a positive exponential. Since $S$ is always positive this superpotential will always break supersymmetry, if we combine this model with a standard asymptotically free model we will have a sum of positive and negative exponentials and the situation will be just like the non-asymptotically free models for which $S$ can be fixed. In this case however, the limit $S \to \infty$ is never a minimum of the scalar potential, even though we can have $W \to 0$ in this limit, and so the Dine-Seiberg general argument still holds but in a very special way, because it would imply that the runaway minimum is not continuously connected to the finite dilaton minimum. This may also lead to very interesting cosmological features. Notice however that the tree-level terms were just chosen to do the job, just as an illustration that this is possible.

To summarize, we have illustrated with a few examples the difference of considering constant gauge couplings against field dependent gauge couplings. We have argued that the inclusion of field-dependent gauge couplings can qualitatively change whether or not a given model spontaneously breaks supersymmetry in the sense that a model with broken supersymmetry may turn out to be supersymmetric if $S$ is included. Furthermore, we have found that the opposite is also possible, i.e. supersymmetry can be unbroken for fixed gauge coupling, but breaks down once the gauge coupling is considered as a field. Finally the stabilization of the dilaton field may be achieved in models of product groups and non asymptotically free, in a way that may not lead to the standard runaway solution. Product group models have shown to provide a very rich structure and their study in more general cases than those considered here can lead to further surprises.

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