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MASSIVE HIDDEN MATTER AND GAUGINO CONDENSATION

B. de CARLOS, J.A. CASAS

Instituto de Estructura de la Materia (CSIC),
Serrano 119, 28006-Madrid, Spain

and

C.MUÑOZ

CERN, CH-1211 Geneva 23, Switzerland

Abstract

We formulate, in a duality-invariant way, the modification produced in the gaugino condensation mechanism by the presence of massive hidden matter (with masses higher than the condensation scale). This kind of scenario is, in fact, very common in string constructions. We observe that the same modification may also be obtained by renormalization group arguments when the effect due to the one-loop integration over heavy string modes is included. Comparison with the other possible scenario (masses smaller than the condensation scale) shows that a unified description for both cases is feasible. Then we explore the phenomenological capabilities of such a scenario. They turn out to be much more promising than the ordinary ones, opening a way for obtaining a correct value of the dilaton (and thus of the gauge coupling constant) dynamically. We have also examined the connection between the effective Lagrangian approach and the truncated superpotential approach for gaugino condensation, proving their equivalence in the physically relevant region.

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As has recently been pointed out, non-perturbative effects can be the crucial ingredient for understanding a number of still unsolved problems in string theory. More precisely, it is believed that gaugino condensation in the hidden sector [1–3] is very likely the determinant factor in breaking supersymmetry, thus fixing the values of the dilaton and moduli and lifting the vacuum degeneracy. Much work has recently been devoted to the study of gaugino condensation in the framework of string theories [4–12]. It has been learned that the assumption of duality, i.e. target space modular $SL(2, \mathbb{Z})$ symmetry in the effective four-dimensional supergravity action [13], modifies the usual form of the condensate [5–7] in essential agreement with previous one-loop calculations including string heavy modes [14]. This modification has proved to be extremely useful for fixing the vacuum expectation value (VEV) of T [5,6], i.e. the modulus that parametrizes the radius of the compactified space. Otherwise the inclusion of world-sheet non-perturbative effects in the superpotential is essential [15,9]. It has also been argued that, in order to fix the value of the dilaton S (and thus the gauge coupling constant, since at string tree level $\text{Re } S = g_{gut}^{-2}$) two or more condensates may be necessary [16,8–10]¹. Very likely the appearance of an intermediate scale after compactification, coming from the fact that hidden matter acquires mass in a dynamical way, is also required [9]. In fact, it is known that the existence of hidden matter is the most general situation and occurs in all promising string constructions [18–20].

In spite of these achievements there remain some pending questions. First, there are several approaches to a dynamic description of the condensation of gauginos (instanton calculations [21], effective Lagrangian approach [22], truncated superpotential approach [2,3], etc.). Although all these methods seem to yield similar answers, the connection between them has not yet been completely established (it is in fact to be proved what is the most correct one). Second, the case of an intermediate scale has not been considered yet in a completely consistent and dual-invariant way. Third, there are no known examples working satisfactorily. In particular, a reasonable value for the dilaton has proved to be elusive.

In this letter we examine in more depth the connection between the effective Lagrangian approach and the truncated superpotential approach for gaugino condensation, proving their equivalence in the region of interest. Then we study the intermediate scale scenario due to massive hidden matter from a duality-invariant point of view and argue that this kind of scenario is probably the most promising one in order to find a realistic value for the dilaton.

¹A different approach conjecturing the existence of duality associated with the S field can be found in ref. [17].

Let us consider the general case of a gauge group which contains several non-Abelian factors: $G = G_1 \otimes G_2 \otimes \dots \otimes G_p$. In the effective Lagrangian approach each gaugino composite bilinear field $(\lambda\lambda)_n$ corresponds to the scalar component of a composite superfield Y_n , such that $Y_n^3 = (\delta_{ab} W_\alpha^a \epsilon^{\alpha\beta} W_\beta^b)_n / S_o^3$ where S_o is the chiral compensator. The dual-invariant action is given by the Kähler potential

$$K_{eL} = -\log(S + \bar{S}) - 3 \log[(T + \bar{T}) - (S + \bar{S})^{-1/3} \sum_n |Y_n|^2] \quad (1)$$

and the modular weight -3 superpotential

$$W_{eL} = \frac{1}{6} \sum_n \beta_n Y_n^3 \log[Y_n^3 f_n(S) h(T)] \quad (2)$$

where β_n is the one-loop coefficient of the beta function $\beta_n(g) = -\beta_n g^3$, $f_n(S) = \exp(3S/2\beta_n)$ and $h(T) \propto \eta^6(T)$ (for more details see e.g. refs.[6,10]). Eq.(2) is slightly modified if the level of the Kac-Moody algebra that generates G_n is different from one. The scalar potential is given by $V_{eL} = e^{G_{eL}} ((G_{eL})_l (G_{eL}^{-1})_m^l (G_{eL})^m - 3)$ [23], with $G_{eL} = K_{eL} + \log |W_{eL}|^2$. The explicit form of V_{eL} was obtained in ref.[10]

$$\begin{aligned} V_{eL} = & [6(1 - \sum_n |Z_n|^2)]^{-2} \left\{ \left| \sum_n \beta_n Z_n^3 [1 + 2S_R f'_n/f_n] \right|^2 \right. \\ & + 3 \sum_n \beta_n^2 |Z_n|^4 [1 + \log[(2S_R)^{1/2} (2T_R)^{3/2} f_n h Z_n^3]]^2 \\ & \left. + \frac{1}{3} \left| \sum_n \beta_n Z_n^3 [3 + 2T_R h'/h] \right|^2 - 3 \left| \sum_n \beta_n Z_n^3 \right|^2 \right\} \quad (3) \end{aligned}$$

where $S_R = \text{Re } S$, $T_R = \text{Re } T$ and $Z_n \equiv (2S_R)^{-1/6} (2T_R)^{-1/2} Y_n$. In the truncated superpotential approach each condensate is substituted by a function of S and T (its presumed VEV). This gives an S and T -dependent superpotential whose functional form can be essentially inferred by symmetry considerations [2,3]. More precisely²

$$W_{trunc} = \frac{\Omega(S)}{\eta^6(T)} \quad (4)$$

where $\Omega(S) = \sum \Omega_n(S)$ and $\Omega_n(S) \propto \exp(\frac{-3S}{2\beta_n})$. The corresponding scalar potential takes the form [5]

$$V_{trunc} = \frac{1}{2S_R (2T_R)^3 |\eta|^{12}} \left\{ |2S_R \Omega_S - \Omega|^2 + 3|\Omega|^2 \left(\left| \frac{4\eta' T_R}{\eta} + 1 \right|^2 - 1 \right) \right\} \quad (5)$$

²As has been pointed out in ref.[12], dual invariance just imposes $W_{trunc} = \frac{\Omega(S)}{\eta^6(T)} H(T)$ with $H(T)$ some modular invariant form, even though loop calculations seem to indicate $H(T) = 1$. Analogously the function $h(T)$ in eq.(2) can also be multiplied by $H(T)$, so all the results obtained in this paper are straightforwardly extended for a non-trivial $H(T)$.

where $\Omega_S = \partial\Omega/\partial S$.

It is easy to check that minimization of V_{eL} (eq.(3)) essentially yields the same results as that of V_{trunc} , provided that relation

$$Y_n^3 = e^{-1}\eta^{-6}(T)e^{-\frac{3S}{2\beta_n}} \quad (6)$$

is fulfilled for the extrema (this implies the vanishing of the F term associated with the condensate). Along this direction V_{eL} takes the form

$$V_{eL}(S, T) = \frac{1}{(1 - \sum_n |Z_n|^2)^2} V_{trunc}(S, T) \quad (7)$$

where the role of $\Omega(S)$ is played by the expression

$$\Omega(S) = \frac{1}{6e} \sum_n \beta_n f_n^{-1}(S) \quad (8)$$

Then for $S_R > \beta_n$, which includes any realistic (weak coupling) case, $\sum |Z_n|^2 \ll 1$, so $V_{eL} = V_{trunc}$.³ The key point is, therefore, to prove eq.(6). It has been noticed by Lüst and Taylor [11] that eq.(6) is obtained from $\partial W_{eL}/\partial Y_n = 0$, which amounts to the fact that, in the limit $S_R \rightarrow \infty$ ($M_p \rightarrow \infty$), global supersymmetry is unbroken. This has been argued to be in fact the case by the authors of ref.[11] on grounds of general arguments (e.g. the Witten index). It is however desirable to prove (6) explicitly, i.e. directly from the extremum conditions. This, in turn, will give the actual range of S_R for which eq.(6) is correct (note that $S_R \rightarrow \infty$ is too vague, and hardly contains the realistic case, i.e. $S_R = O(1)$). It is a tremendous task to find the general solutions of the extremum conditions for V_{eL}

$$\frac{\partial V_{eL}}{\partial Y_n} = 0, \quad \frac{\partial V_{eL}}{\partial S} = 0, \quad \frac{\partial V_{eL}}{\partial T} = 0 \quad (9)$$

It is more sensible to look for general relations between Y_n , S and T using eq.(9) and duality arguments. In particular, after some work with eq.(9) one arrives at the following equation

$$\begin{aligned} & \sum_n \beta_n^2 |Z_n|^4 \left[\frac{3S_R}{\beta_n} - 4T_R \frac{\eta'}{\eta} \right] \left(1 + \log \left[\eta^6 Y_n^3 e^{\frac{3S}{2\beta_n}} \right] \right)^* \\ & + S_R \sum_n Z_n^3 \left(\sum_m \beta_m Z_m^3 \left[1 + \frac{3S_R}{\beta_m} \right] \right)^* + S_R \left(\sum_n \beta_n Z_n^3 \left[1 + \frac{3S_R}{\beta_n} \right] \right) \left(\sum_m Z_m^3 \right)^* \\ & - \frac{4}{9} T_R \left[\frac{\eta'}{\eta} + 2T_R \left(\frac{\eta''}{\eta} - \left(\frac{\eta'}{\eta} \right)^2 \right) \right] \left(\sum_n \beta_n Z_n^3 \right) \left(\sum_m \beta_m Z_m^3 \left[3 + 12T_R \frac{\eta'}{\eta} \right] \right)^* \\ & - \frac{4}{9} T_R \left(\frac{\eta'}{\eta} \right) \left(\sum_m \beta_m Z_m^3 \left[3 + 12T_R \frac{\eta'}{\eta} \right] \right) \left(\sum_n \beta_n Z_n^3 \right)^* = 0 \end{aligned} \quad (10)$$

³This result was first obtained in ref.[9], but with the assumption of eq.(6).

As stated above, $|Z_n|^2 \ll 1$ for a wide region of S_R . Keeping only the leading terms in $|Z_n|^2$ the previous expression simplifies to

$$\sum_n \beta_n^2 |Z_n|^4 \left[\frac{3S_R}{\beta_n} - 4T_R \frac{\eta'}{\eta} \right] \left(1 + \log \left[\eta^6 Y_n^3 e^{\frac{3S}{2\beta_n}} \right] \right)^* = 0 \quad (11)$$

This condition can be fulfilled in different ways. However, not all of them will also be a solution of the complete extremum equations, eq.(9). More precisely, only the solutions of (11) which are invariant under modular $SL(2, \mathbb{Z})$ transformations, i.e.

$$T \rightarrow \frac{aT - ib}{icT + d}, \quad a, b, c, d \in \mathbb{Z}, \quad ad - bc = 1 \quad (12)$$

may be solutions of (9). It is not difficult to check by inspection that the only modular invariant way to fulfil (11) is to demand the vanishing of the quantity between parenthesis for each n , thus leading to eq.(6), which is $SL(2, \mathbb{Z})$ modular invariant. This confirms the equivalence between the two approaches in a wide range of values of S_R , which includes the realistic case $S_R = O(1)$.

Let us now turn to the massive hidden matter issue. The case of "quark" ($N + \bar{N}$) representations of $SU(N)$ acquiring masses below the condensation scale was recently analyzed in ref.[11] following the lines of ref.[24]. Here we wish to analyze the case of quark masses *higher* than the condensation scale (Λ), but *smaller* than the unification scale (M_{gut}). We follow a different, and quite simple, approach⁴. For the sake of simplicity we work with a single non-Abelian gauge group G whose matter representations acquire all the same mass M_I ($\Lambda < M_I < M_{gut}$) through a Higgs mechanism. The decoupling theorem ensures that at the condensation scale the theory should behave as a pure Yang-Mills (SYM) theory, described by a superpotential similar to that of eq.(2). There is, however, an important difference. In eq.(2), the renormalization group (RG) invariant scale, $\mu \exp\{-\int^g dg'/\beta(g')\}$, calculated at the one-loop level at the unification point ($g^{-2}(M_{gut}) = S_R$), has been implicitly substituted (for more details see e.g. ref.[21]). The effective SYM theory, which we are dealing with, has a different gauge coupling at the unification point, related to S_R by

$$S_R \rightarrow S_R + (\tilde{\beta}_o - \beta_o) \log \left(\frac{M_I}{M_{gut}} \right)^2 \quad (13)$$

where $\tilde{\beta}_o$ and β_o are the coefficients of the beta function between M_{gut} and M_I , and between M_I and Λ respectively. Then eq.(2) is modified to

$$W_{eL} \rightarrow \frac{1}{6} \beta_o Y^3 \log \left(Y^3 \eta^6(T) \left(\frac{M_I}{M_{gut}} \right)^{\frac{3}{\tilde{\beta}_o}(\tilde{\beta}_o - \beta_o)} e^{\frac{3S}{2\beta_o}} \right) \quad (14)$$

⁴A preliminary study of this problem was performed in ref.[9]

(This modification is equivalent to calculating the above-mentioned RG invariant scale at the unification point taking into account the contribution of the heavy matter between M_I and M_{gut} .) Eq.(14) is not yet the final form of our effective superpotential⁵. Suppose that ξ_1, ξ_2 are two matter representations acquiring mass through the term $A\xi_1\xi_2$ in the superpotential when the G -singlet field, A , acquires dynamically a non-vanishing VEV (we are assuming here that A, ξ_1, ξ_2 are untwisted fields; the modification for the twisted case is straightforward). This type of couplings naturally appear in string constructions [18–20]. Using the corresponding Kähler potential

$$K = -\log(S + \bar{S}) - 3\log[(T + \bar{T}) - (S + \bar{S})^{-1/3}|Y|^2 - 2\sum_i |\xi_i|^2 - 2|A|^2] \quad (15)$$

and the usual expression for the unification mass, it is straightforward to write the kinetic and mass terms for ξ_1, ξ_2 [23], and obtain M_I/M_{gut} in terms of A . After some algebra, the result turns out to be simply $M_I/M_{gut} = |A|$. Expressing this in an analytical way, we obtain from (14) a correctly defined superpotential. Finally we should note that A has modular weight -1 . Hence, in order to keep modular invariance in the Lagrangian, we have to substitute $A \rightarrow \eta^2(T)A$. Therefore the final form of the effective superpotential is

$$W_{eL} = \frac{1}{6}\beta_o Y^3 \log\left(Y^3 \eta^6(T) (\eta^2 A)^{\frac{3}{\beta_o}(\tilde{\beta}_o - \beta_o)} e^{\frac{3S}{2\tilde{\beta}_o}}\right) \quad (16)$$

If we choose to work with the truncated superpotential approach, then we have to write

$$W_{trunc} = \frac{\Omega(S, T, A)}{\eta^6(T)} \quad (17)$$

with

$$\Omega(S, T, A) \propto (A\eta^2)^{3(1-\tilde{\beta}_o/\beta_o)} e^{\frac{-3S}{2\tilde{\beta}_n}} \quad (18)$$

It is worth noticing that the same result can also be obtained by using RG arguments as follows. It is known that the form of the gaugino condensate can be understood by dimensional arguments to correspond to the scale at which the gauge coupling becomes strong [2]. In our case the one-loop RG equations for the coupling constant read

$$\begin{aligned} \frac{1}{g^2(M_I)} &= \frac{1}{g^2(M_{gut})} + \tilde{\beta}_o \log\left(\frac{M_I^2}{M_{gut}^2}\right) \\ \frac{1}{g^2(\Lambda)} &= \frac{1}{g^2(M_I)} + \beta_o \log\left(\frac{\Lambda^2}{M_I^2}\right) \end{aligned} \quad (19)$$

⁵Note, for instance, that the presence of M_I/M_{gut} in (14) breaks the analyticity of the superpotential.

The coupling becomes strong at

$$\Lambda^2 = M_{gut}^2 \left(\frac{M_I^2}{M_{gut}^2} \right)^{1-\tilde{\beta}_o/\beta_o} e^{-\frac{1}{\beta_o g^2(M_{gut})}} \quad (20)$$

The gaugino condensate, $\langle \lambda\lambda \rangle = (S + \bar{S})\Lambda^3$ [6], is then given by

$$(S + \bar{S})M_{gut}^3 \left(\frac{M_I}{M_{gut}} \right)^{3(1-\tilde{\beta}_o/\beta_o)} e^{-\frac{3}{2\beta_o g^2(M_{gut})}} \quad (21)$$

After performing the usual Weyl rescaling of the gaugino fields [23] and writing the final expression in an analytical way, the superpotential turns out to be

$$W = A^{3(1-\tilde{\beta}_o/\beta_o)} \frac{e^{-\frac{3S}{2\tilde{\beta}_o}}}{\eta^{6\tilde{\beta}_o/\beta_o}} \quad (22)$$

which coincides with the result obtained in eq.(17). Note that in writing (22) we have introduced a factor $(\eta^{6\tilde{\beta}_o/\beta_o})^{-1}$ in order to keep the modular invariance of the action, which requires a modular weight -3 superpotential.

As it was mentioned above, in the pure Yang-Mills case the modification in the usual form of the condensate due to the assumption of modular invariance is known to be essentially in agreement with the modification coming from the one-loop contribution of the heavy string modes to the coupling constant [14,5-7]. More precisely, the latter contribution gives an additional term $\sim \log|\eta|^4$ in the first equation of (19) [14] (there is also a piece $\sim \log(T + \bar{T})$ which is absorbed into the usual definition of M_{gut} [5]). For the pure Yang-Mills case the equivalence is achieved only if the coefficient of this term coincides with $\tilde{\beta}_o$ (the beta function of the $N = 1$ theory)⁶. It is interesting to check that this is also the case in the scenario at hand, i.e. when massive hidden matter is present. To see this, note that the inclusion of the term $\tilde{\beta}_o \log|\eta|^4$ in the first equation of (19) corrects the second member of eq.(20) with a factor $|\eta|^{-4\tilde{\beta}_o/\beta_o}$. Then eq.(22) is obtained directly, without any additional assumption.

It is quite interesting to compare this result with that obtained in ref.[11] for quark masses smaller than Λ . There a superpotential involving Y (the gaugino condensate) and Π (the squark condensate), whose form is dictated by Ward identities and modular invariance, was used. After eliminating Y and Π from $\frac{\partial W}{\partial Y} = \frac{\partial W}{\partial \Pi} = 0$, which, as mentioned above, is claimed to be true for $S_R \rightarrow \infty$, a simplified form for the superpotential was obtained:

$$W \propto \eta^{2M/N-6} e^{-8\pi^2 S/N} A^{M/N} \quad (23)$$

⁶This is a rather polemic point in the current literature.

where A is the generic field giving mass to M ($N + \bar{N}$) representations of a $SU(N)$ gauge group. Remarkably enough eq.(23) coincides with eq.(17), in spite of the fact that they have been obtained on completely different grounds. Consequently, we conclude that eq.(17) (or, equivalently eq.(22)) can be used in a safe way, independently of the final VEV of A . This simplifies the analysis notably. Let us finally note that eq.(17) represents a good example, and an extension, of the most general form of a modular invariant superpotential, formulated in ref.[12]. There it was argued that the superpotential for gaugino condensation should have a form $W = \frac{\Omega(S)}{\eta^6(T)} H(T)$ with $H(T)$ some modular invariant form. We see that in eq.(17) the role of $H(T)$ is played by $(A\eta^2(T))^{3(1-\tilde{\beta}_o/\beta_o)}$, i.e. a modular invariant function involving T and A . When multiple gaugino condensation is considered (see below) $\Omega(S, T, A)$ cannot even be factorized as $\frac{\Omega(S)}{\eta^6(T)} H(T, A)$.

Let us now explore the phenomenological capabilities of a massive hidden matter (MHM) scenario. The relevant Kähler potential is

$$K = -\log(S + \bar{S}) - 3\log[(T + \bar{T}) - 2|A|^2] \quad (24)$$

and the superpotential is of the form

$$W = \frac{\Omega(S, T, A)}{\eta^6(T)} \quad (25)$$

where, in general, $\Omega(S, T, A) = \sum_n \Omega_n(S, T, A)$. Each Ω_n will be of the type of eq.(18) with values of β_o and $\tilde{\beta}_o$ depending on how the breaking affects the corresponding gauge group. Terms of the type $A\xi_i\xi_j$ can be added to W , as well as $-2|\xi_i|^2$ terms inside the second log of eq.(24). However, when the D-terms associated with ξ_i are taken into account, it can be shown that $\langle \xi_i \rangle = 0$ corresponds to a stable minimum. Therefore we can ignore this additional complication. The scalar potential derived from (24), (25) is

$$V = \frac{1}{2S_R(2T_R - 2|A|^2)^3} \left\{ |2S_R W_S - W|^2 + \frac{1}{6}(2T_R - 2|A|^2) \left| \frac{\partial W}{\partial A} + 2\bar{A} \frac{\partial W}{\partial T} \right|^2 + \frac{1}{3}(2T_R - 2|A|^2)^2 \left| \frac{\partial W}{\partial T} - \frac{3W}{(2T_R - 2|A|^2)} \right|^2 - 3|W|^2 \right\} \quad (26)$$

For $|A|^2 = T_R$ there exists an infinite barrier which allows us to restrict the values of $|A|^2$ to the range $|A|^2 < T_R$. This is also the condition for kinetic energy positivity. The potential without MHM, i.e. eq.(5), is recovered from (26) if $\beta_o = \tilde{\beta}_o$ for all the condensates and $\langle A \rangle \rightarrow 0$.

The first consideration about (26) is that for a single condensate there is no stable minimum for any value of S_R , since V behaves as $V \sim \frac{|W|^2}{S_R}$. Of course this also applies

if no MHM is present. This leads us to consider the multiple gaugino condensation scenario[16,8–10]. In the absence of MHM, however, it does not help very much. For example, for two condensates the potential of eq.(5) can be shown to develop non-trivial stationary points for certain values of S_R , but they always fall far away from the realistic range [25]. In the presence of MHM things get more involved, but with much better prospects. To see this let us consider an important kind of minima: those for which $V \leq 0$. These minima are easier to manage, since a very useful bound on S_R can be derived for them. Namely, from (26) it is clear that $V \leq 0$ only if $|2S_R W_S - W|^2 \leq 3|W|^2$, which leads to the condition⁷

$$S_R \leq \frac{1}{2}(1 + \sqrt{3}) \left| \frac{W}{W_S} \right| \quad (27)$$

(Of course this bound also applies in the absence of MHM.) It is easy to check that for two condensates bound (27) can only be satisfied (for $S_R > (1 + \sqrt{3})\beta_o/3$) inside a very narrow band around $W_S = 0$.⁸ The width of the band is $(2\beta_o/3)^2(1 + \sqrt{3})S_o^{-1}$ approximately, where S_o is the point where the derivative vanishes. We will illustrate the implications of MHM by means of an explicit example. Consider a vacuum in which the hidden gauge group has two separate sectors with β_o and β'_o the respective one-loop beta functions. The contribution of gaugino condensation to the superpotential is

$$W = \eta^{-6}(T) \left(e^{-3S/2\beta_o} + (A\eta^2)^{3(1-\tilde{\beta}'_o/\beta'_o)} e^{-3S/2\beta'_o} \right) \quad (28)$$

where we have allowed for an intermediate scale modifying the second condensate according to (17) ($\tilde{\beta}'_o$ is the beta function between M_I and M_{gut}). We find $W_S = 0$ for

$$\frac{3S_R}{2\beta_o} = \frac{1}{1 - \beta_o/\beta'_o} \log \left(\frac{\beta'_o}{\beta_o} \left(\frac{1}{|A\eta^2|} \right)^{3(1-\tilde{\beta}'_o/\beta'_o)} \right) \quad (29)$$

As has been shown above, this value of S_R is extremely close to the minimum of the potential. It is clear that for $\tilde{\beta}'_o = \beta'_o$, i.e. in the absence of massive hidden matter, S_R becomes too small. In order to get a reasonable value for S_R the condition $\beta'_o > \beta_o$ is needed. There are two possible cases: a) $\beta'_o > \beta_o > \tilde{\beta}'_o$ and b) $\beta'_o > \tilde{\beta}'_o > \beta_o$. It is not difficult to check that (a) leads to a scenario with quark masses higher than the condensation scale, which is the case studied in this letter, whereas (b) leads to the opposite case, which has been studied in ref.[11]. Let us also note that if $3(1 - \tilde{\beta}'_o/\beta'_o) < 1$ the scalar

⁷The only exception to this rule is when $W = W_S = 0$. Then V might also vanish, but with unbroken supersymmetry.

⁸This supports the claim that for $S_R \rightarrow \infty$ global supersymmetry is unbroken.

potential, given in eq.(26), behaves as $V \rightarrow \infty$ for $|A| \rightarrow 0, T_R$ and $|2S_R W_S - W| > 0$. This implies the existence of a non-trivial minimum for A , thus allowing the appearance of gaugino condensation. For a $SU(N)$ gauge group in which $M(N + \bar{N})$ representations acquire mass, the previous condition simply reads $M < N$, in agreement with previous results [21,11].

In order to get a quantitative flavour we take the following values for the parameters that appear in (28)

$$\beta_o = \frac{12}{16\pi^2}, \quad \beta'_o = \frac{15}{16\pi^2}, \quad \tilde{\beta}'_o = \frac{11}{16\pi^2} \quad (30)$$

This corresponds to a $SU(4) \times SU(5)$ scenario where four $(5 + \bar{5})$ representations acquire masses. Now, from (29), W_S vanishes for $S_R = 1.78$ if $|A\eta^2| = 2 \times 10^{-4}$, a perfectly reasonable value, thus leading to a realistic gauge coupling and a correct large mass hierarchy ($3S_R/2\beta_o = 35$). Of course the actual value of A must be obtained by a complete minimization [25]. Let us finally notice that if A possesses a non-trivial anomalous $U(1)$ charge, the corresponding Fayet-Iliopoulos term will also contribute to fixing the value of A [26,18–20].

To summarize, in this letter we have first examined the connection between the effective Lagrangian approach and the truncated superpotential approach for gaugino condensation, proving their equivalence in the physically relevant region ($S_R = O(1)$). Secondly we have formulated, in a duality-invariant way, the modification of the effective superpotential when massive hidden matter (with masses higher than the condensation scale) is present. This kind of scenario is, in fact, the usual one. We have observed that the same result may also be obtained by RG arguments when the effect due to the one-loop integration over heavy string modes is included. Comparison with the other possible scenario (masses smaller than the condensation scale) shows that a unified description for both cases is feasible. Finally, we have explored the phenomenological capabilities of such a scenario. They turn out to be much more promising than the ordinary ones, opening a way for obtaining a correct value of S_R (and thus of the gauge coupling constant) dynamically.

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