Gauge and Matter Condensates in Realistic String Models

S. KALARA, JORGEL. LOPEZ,* and D. V. NANOPOLIUS

Center for Theoretical Physics, Department of Physics, Texas A&M University
College Station, TX 77843-4242, USA
and
Astroparticle Physics Group, Houston Advanced Research Center (HARC)
The Woodlands, TX 77381, USA

ABSTRACT

We examine the inter-relationship of the superpotential containing hidden and observable matter fields and the ensuing condensates in free fermionic string models. These gauge and matter condensates of the strongly interacting hidden gauge groups play a crucial role in the determination of the physical parameters of the observable sector. Supplementing the above information with the requirement of modular invariance, we find that a generic model with only trilinear superpotential allows for a degenerate (and sometimes pathological) set of vacua. This degeneracy may be lifted by higher order terms in the superpotential. We also point out some other subtle points that may arise in calculations of this nature. We exemplify our observations by computing explicitly the modular invariant gaugino and matter condensates in the flipped $SU(5)$ string model with hidden gauge group $SO(10) \times SU(4)$.

September, 1991

* Supported by an ICSC–World Laboratory Scholarship.
1. Introduction and General Remarks

Nonperturbative dynamics of strongly interacting supersymmetric gauge theories has been explored using a wide variety of tools with great success [1, 2]. A better understanding of dynamics may lead to solutions to some deep-rooted problems like supersymmetry breaking [3], the gauge hierarchy problem [4], and even string compactification [5]. In the context of string theory, the task of incorporating strongly interacting gauge theory dynamics becomes even more compelling since many crucial properties of the string theory may not be discerned until a deeper grasp of the vacuum structure of this kind of gauge theories is obtained [6].

Typically a theory which aims to go beyond the Standard Model and explain some of its features, may contain many different mass scales beyond $M_Z$ and may involve many other gauge degrees of freedom, some of which may even be strongly interacting. Examples of such theories include all grand unified theories [7], technicolor theories [8], and supergravity theories with or without strongly interacting hidden sectors [9]. In any of these theories one may propose the existence of a strongly interacting gauge theory for no other reason, save the problem at hand, i.e., elimination of elementary Higgs fields [8], supersymmetry breaking [10], etc. However, in string theory one finds that the consistency of the theory forces upon us the existence of a completely hidden and/or semi-hidden gauge group which has to reckoned with. Furthermore, one finds that in a large class of models, the theory also contains hidden matter fields. The interaction of the hidden matter fields with themselves and with the observable matter fields is in principle completely calculable [11]. It is the interplay between nonperturbative dynamics of the hidden sector and its possible effect on the observable sector which is of great interest.

In this paper we examine the reciprocal influence of the nonperturbative dynamics of strongly interacting gauge theories and the interaction among the matter fields in a class of models. The type of models we consider are characterized by two mass scales: $\Lambda_G$ where the gauge theory becomes strong, and the mass scale $\mathcal{M}$ ($\mathcal{M} \gg \Lambda_G$). The key ingredient that allows us to probe the theory in great detail is the tamed ultraviolet behavior of the theory due to supersymmetry and the presence of genuinely stringy discrete symmetries.

In a typical supergravity theory, without the benefit of any string input, the dynamics of the gaugino condensate depends primarily on the gauge coupling constant $g$ and the masses of the matter fields charged under the strongly interacting gauge group. However, in string theory the gauge coupling $g$ is related to the vacuum expectation value of the dilaton.
field $S$ as $1/g^2 = 4\text{Re} \langle S \rangle$ and consequently becomes a dynamical variable. Furthermore, the interactions of the matter fields are only dictated by the string dynamics and, in a class of theories, are calculable [12]. Additionally, string theory imposes very strong discrete symmetries (the so-called target space modular symmetries [13]) on the theory, making the structure of the gaugino condensate very tightly constrained.

In the preparatory examples treated in the literature [14,15,16,17], i.e., examples with only one modulus field and the case where all matter fields of the hidden gauge group acquire masses through trilinear superpotential terms, most of the stringy information has been incorporated. Specifically, for the case of $SU(N)$ with $M$ flavors in the fundamental representation, when the mass $\mathcal{M}$ of the matter fields is $\mathcal{M} \ll \Lambda_G$, one finds [15]

$$\frac{Y^3}{32\pi^2} = \left(32\pi^2 e\right)^{M/N-1}[c\eta(T)]^{2M/N-6} [\det \mathcal{M}]^{1/N} e^{-32\pi^2 S/N}, \quad (1.1a)$$

$$\Pi_{ij} = \frac{Y^3}{32\pi^2} \mathcal{M}_{ij}^{-1}, \quad (1.1b)$$

where $T$ is the modulus field and $\eta$ the Dedekind function, $c$ is an unknown constant, the composite superfield $Y$ is defined as $Y^3 = W_\alpha W^\alpha / S_0^3$ where $S_0$ is the chiral compensator, and $\Pi_{ij} = \langle C_i C_j \rangle$ are the matter condensates. Using one-loop renormalization group equations, the complementary case of $\mathcal{M} \gg \Lambda_G$ can also be incorporated [16]. For an alternate approach using an effective Lagrangian see Ref. [17]. However, a generic full-fledged string theory example introduces its own set of complications. It is the purpose of this paper to examine some of these involutions.

In a realistic example, to obtain a coherent picture of the vacuum structure of the theory requires a multipronged approach. Generically one finds that some of the matter fields which carry nontrivial gauge quantum numbers of the strongly interacting theory do not acquire mass at the trilinear level of the superpotential. As it can be easily shown that a strongly interacting supersymmetric gauge theory with massless matter fields is beset with the problem of pathological vacuum structure [1], this necessitates that the mass calculations for the matter fields be carried up to quartic order and beyond.

The presence of the modular invariance further restricts the type of terms that can arise [3]. In a class of models higher order nonrenormalizable terms can be calculated [12] and their modular invariance properties inferred [18]. The key point to note here is that the requirement of a stable vacuum mandates that the superpotential be probed up to the level at which the determinant of the mass matrix is nonzero. The presence of a zero
eigenvalue in the mass matrix invariably leads to both theoretical (unstable vacua \[1\]) and experimental (breakdown of the equivalence principle \[19\]) difficulties.

Furthermore, strongly interacting supersymmetric gauge theories are notorious for containing a large class of degenerate vacua. Unless the effects of the superpotential are taken into account, the degeneracy remains unbroken. Additionally, many of the expositions in strongly interacting gauge theory dynamics are based on the global and/or local symmetries of the theory \[1\], thus in the absence of complete knowledge of the superpotential an unambiguous identification of the vacuum structure may not be possible.

The free-fermionic formulation of string theory \[20\] is specially suited to explore these questions, since the nonrenormalizable terms can be explicitly calculated \[12\] and modular properties of these terms can also be determined. Succinctly put, a nonrenormalizable term $\Phi_1 \ldots \Phi_N$ can be calculated by evaluating the correlator $\langle \Phi_1 \ldots \Phi_N \rangle$. After taking into account the massless exchanges, a nonzero piece at low momenta signifies a presence of such a term in the superpotential. A priori, the existence of such a term in the superpotential would be inconsistent with modular invariance. However, it can be shown \[18\] that the nonzero value of the correlator $\langle \Phi_1 \ldots \Phi_N \rangle$ implies a nonzero value of the following series of correlators $\langle T^p \Phi_1 \ldots \Phi_N \rangle$, where $p \geq N - 3$ and $T^p$ is a definite product of $p$ moduli fields of the theory. Supplementing this observation with the requirement of modular invariance, we find that the coefficient of a given nonrenormalizable term can be summed up into a modular covariant function made up of products of Dedekind $\eta$ functions of the moduli fields involved. A modular invariant gaugino condensate will also have to be reconciled with the nontrivial modular properties of nonrenormalizable terms.

Thus we see that the question of gaugino condensate cannot be answered independently of the superpotential. The effect of the superpotential is felt in determining the vacuum structure of the theory lifting the degeneracy of the vacuum, and in the specific value of the gaugino condensate.

2. Application to Flipped $SU(5)$

We now present an explicit example of the calculation of hidden massive matter condensates and the various subtle points that arise in realistic string models. To this end we study the flipped $SU(5)$ string model \[21\] which has a rich hidden sector spectrum composed of two gauge groups, namely $SO(10)$ and $SU(4)$. Since $SO(10)$ is expected to become strongly interacting at a scale ($\Lambda_{10} \approx 10^{14−16}$ GeV \[22\]) much higher than the
respective $SU(4)$ scale ($\Lambda_4 \approx 10^{10-12} \text{GeV}$ [22]), in the following we deal exclusively with $SO(10)$ condensates. Besides $\Lambda_{10}$, two other mass scales come into play: the $SU(5) \times U(1)$ beaking scale, $V, \overline{V} \sim 10^{15-16} \text{GeV}$; and the scale of singlet vevs needed to cancel the anomalous D-term, $\langle \phi \rangle \sim 10^{17} \text{GeV}$ [21, 23, 24]. All these scales are to be compared with the string unification scale $M_{SU} \approx 1.24 \times g \times 10^{18} \text{GeV}$ [23, 24] and the scale of nonrenormalizable terms in the superpotential $M \approx 10^{18} \text{GeV}$ [22].

We first need to determine the mass matrix for the $SO(10)$ fields. These belong to the $\mathbf{10}$ of $SO(10)$ and are denoted by $T_i, i = 1 \rightarrow 5$. There are generally three sources of mass terms for these fields:

(i) $T_i T_j \phi^{N-2}$: at $N=3$ we have the following terms [6, 25]

$$\Phi_{23} T_i^2, \Phi_{31} T_i^2, \Phi_{23} T_4^2, \Phi_{31} T_5^2, \phi_3 T_4 T_5,$$

where the various $\phi$'s are singlet fields which will generally get vevs $\langle \phi \rangle \sim 10^{17} \text{GeV}$, but could vanish for consistency or phenomenological reasons. Following the methods in [24], it can be shown that no new terms arise at any order $N \geq 4$, except for small corrections to the above terms (e.g., $T_1^2 \Phi_{31} \bar{\phi}_1^2 / M^2$).

(ii) $T_i T_j F_1, 3 \bar{F}_5 \phi^{N-4}$: here $F_{1,3}$ and $\bar{F}_5$ are the $SU(5)$ fields which get vevs $V, \overline{V}$ respectively. No such terms occur at $N=4, 6, 8$. At $N=5$ we get $T_1 T_4 F_1 \bar{F}_5 \phi_3 / M^2$ and at $N=7$

$$T_3 T_4 F_3 \bar{F}_5 \{\phi_2 \phi_3 \bar{\phi}_-^-, \phi_2 \bar{\phi}_3 \phi^+, \phi_3 \bar{\phi}_2 \phi^+\} / M^4,$$

$$(2.2a)$$

$$T_3 T_5 F_3 \bar{F}_5 \{\phi_3 \bar{\phi}_-^-, \phi_3 \bar{\phi}_3 \phi^+ \Phi_{31}\} / M^4.$$

$$(2.2b)$$

Yet higher order terms ($N \geq 9$) are further suppressed by powers of $\langle \phi \rangle / M \sim 10^{-2}$.

(iii) $T_i T_j T_k T_l \phi^{N-4}$: no such terms occur at $N=4, 5, 7$. At $N=6$ we get

$$T_3 T_3 T_4 T_4 \phi_{45} \phi^+/M^3,$$

$$(2.3a)$$

$$T_1 T_1 T_4 T_5 \{\phi_1 \bar{\phi}_2, \phi_2 \bar{\phi}_1\} / M^3,$$

$$(2.3b)$$

$$T_1 T_1 T_4 T_5 \phi_1 \Phi_{31} / M^3.$$

$$(2.3c)$$

With yet higher orders further suppressed.

We should point out that a very useful constraint in searching for non-vanishing higher-order nonrenormalizable terms in the superpotential is given by the modular invariance of the all-orders superpotential, as discussed above.
The resulting $T_i$ mass matrix can be written as follows

$$
M_{ij} = \begin{pmatrix}
\bar{\Phi}_{23} & 0 & 0 & \eta & 0 \\
0 & \Phi_{31} & 0 & 0 & 0 \\
0 & 0 & \delta_{44} & \delta_{34} + \epsilon_4 & \epsilon_5 \\
\eta & 0 & \delta_{34} + \epsilon_4 & \delta_{33} + \Phi_{23} & \phi_2 \\
0 & 0 & \epsilon_5 & \phi_2 & \Phi_{31}
\end{pmatrix}
$$

where

$$
\delta_{ij} = \langle T_i T_j \rangle \phi_{45} \phi^+ \frac{1}{M^3} \equiv \frac{\delta}{M} \Pi_{ij},
$$

$$
\epsilon_{4,5} = F_3 F_5 \frac{\phi^3_{4,5}}{M^4},
$$

$$
\eta = F_1 F_5 \phi_3 \frac{1}{M^2},
$$

with the relevant $\phi^3_{4,5}$ given in (2.2). It is important to estimate the sizes of the various entries in $M_{ij}$. With the above mentioned scales we obtain: $\langle \phi \rangle \sim 10^{-1} M$, $\eta \sim 10^{-5} M$, $\epsilon_{4,5} \sim 10^{-7} M$, and $\delta \sim 10^{-2}$.

Since $\Pi_{ij} = \frac{Y^3}{32\pi^2} M_{ij}^{-1}$, and $M_{ij}$ depends on $\Pi_{ij}$ also, the solution of the resulting equations can be rather nontrivial. We now invert the matrix $M$ in various levels of approximation to exhibit the subtleties that can arise. As a first approximation let us drop all nonrenormalizable contributions to $M$. In this limit the matrix breaks up into three blocks with a zero eigenvalue for $T_3$. The solution is $\Pi_{11,22,44,45,55} \sim \frac{Y^3}{32\pi^2} \frac{1}{\langle \phi \rangle}$ and $\Pi_{33} \to \infty$, with all the other condensates vanishing. Clearly this is an unphysical situation and we are forced to consider nonrenormalizable terms.

Next we keep the $\delta_{ij}$ terms and neglect the $\epsilon_{4,5}$ and $\eta$ terms. Schematically the matrix becomes

$$
\begin{pmatrix}
\phi & 0 & 0 & 0 & 0 \\
0 & \phi & 0 & 0 & 0 \\
0 & 0 & \delta_{44} & \delta_{34} & 0 \\
0 & 0 & \delta_{34} & \delta_{33} + \phi & \phi_2 \\
0 & 0 & 0 & \phi_2 & \phi
\end{pmatrix}
$$

The $2 \times 2$ submatrix involving $T_{1,2}$ remains unchanged, i.e., $\Pi_{11} \sim \Pi_{22} \sim \frac{Y^3}{32\pi^2} \frac{1}{\langle \phi \rangle}$. To simplify the calculation without loss of generality, let us set $\langle \phi_2 \rangle = 0$, thus decoupling the $T_5$ field, i.e., $\Pi_{45} = 0$, $\Pi_{55} \sim \frac{Y^3}{32\pi^2} \frac{1}{\langle \phi \rangle}$. The remaining $2 \times 2$ matrix can be easily inverted and the following three equations result

$$
\Pi_{44} \left\{ 1 - \frac{Y^3}{32\pi^2} \frac{1}{D M} \delta \right\} = 0,
$$

(2.7a)
\[ \Pi_{33} \left\{ 1 - \frac{Y^3}{32\pi^2} \frac{1}{D M} \delta \right\} = \frac{Y^3}{32\pi^2} \langle \phi \rangle D, \quad (2.7b) \]

\[ \Pi_{34} \left\{ 1 + \frac{Y^3}{32\pi^2} \frac{1}{D M} \delta \right\} = 0, \quad (2.7c) \]

where

\[ D = \left( \frac{\langle \phi \rangle + \delta}{M} \Pi_{33} \right) \frac{\delta}{M} \Pi_{44} - \left( \frac{\delta}{M} \Pi_{34} \right)^2. \quad (2.8) \]

Equation (2.7c) can be solved if (i) \( \Pi_{34} = 0 \) and/or (ii) \( \frac{Y^3}{32\pi^2} \frac{\delta}{D M} = -1 \). In case (ii) we get from (2.7b)

\[ \Pi_{44} = 0, \quad (2.9) \]

and from (2.7b)

\[ \Pi_{33} = \frac{1}{2} \frac{Y^3}{32\pi^2} \frac{\langle \phi \rangle}{D} = -\frac{1}{2} \frac{M}{\delta} \langle \phi \rangle \sim M^2. \quad (2.10) \]

Also, \( \Pi_{34} \) gets determined through \( D \),

\[ \Pi_{34} = \left( \frac{Y^3}{32\pi^2} \frac{M}{\delta} \right)^{1/2}. \quad (2.11) \]

In case (i), Eqn. (2.7c) gives \( \Pi_{44} = 0 \) and/or \( \frac{Y^3}{32\pi^2} \frac{\delta}{D M} = 1 \), both of which lead to unphysical or inconsistent solutions. Thus solution (ii) is preferred. Note that the field \( T_3 \) remains very light \( (\delta_{44} = 0) \), while \( \Pi_{33} \sim M^2 \) is finite.

For the case when all matter fields are the mass eigenstates, the usual clustering argument, which is based on the symmetries of the superpotential, would imply that \( \Pi_{34} = 0 \). However, we see that \( T_3, T_4 \) are not mass eigenstates due to a mixing term \( \epsilon_4 T_3 T_4 \sim 10^{-7} M T_3 T_4 \). The presence of such a term breaks the degeneracy of the vacuum: \( \Pi_{44} = 0, \Pi_{34} \neq 0 \) and \( \Pi_{44} \neq 0, \Pi_{34} = 0 \). One can repeat the above calculations taking \( \epsilon_4 \neq 0 \) (and \( \epsilon_5 = 0 \) for simplicity) and the only change in Eqs. (2.7) is in the right-hand side of Eq. (2.7d) which becomes \( -\frac{Y^3}{32\pi^2} \frac{\epsilon_4}{D} \). The solutions of the equations are more complicated, but they differ negligibly from the case \( \epsilon_4 = 0 \). That is, even though \( \epsilon_4 \) is negligible numerically, its presence determines the correct vacuum of the theory.

The above example shows the importance of knowing the higher-order superpotential to obtain physically acceptable results. We should point out that it may happen that some of the singlet vevs vanish, implying \( \delta = 0 \) and/or \( \epsilon_{4,5} = 0 \). As discussed above, this scenario may lead to unphysical solutions and thus these vev choices appear disfavored.
3. Calculation of the Gaugino Condensate

We now calculate the only unknown left in the theory, namely the gaugino condensate \( Y^3 \). Equation (1.14) has been obtained by imposing the modular symmetry of \( SL(2, \mathbb{Z}) \) on the effective action. Generally the modular group is much larger. For the class of free-fermionic models under consideration one finds that the symmetry group is \( SL(2, \mathbb{Z})^3 \) accompanied by three moduli fields \cite{18}. Correspondingly the gaugino condensate is straightforwardly generalized, and in the case of \( SO(10) \) with \( M \) massive flavors we obtain

\[
\frac{Y^3}{32\pi^2} = (32\pi^2e)^{M/8-1}[c\eta(T_1)\eta(T_2)\eta(T_3)]^{-2}[\text{det} \tilde{M}]^{1/8}e^{-32\pi^2S/8},
\]

where \( \text{det} \tilde{M} \) is the determinant of the mass matrix \( M \) calculated with modular invariant mass terms. This way \( Y^3 \) has modular weight \((-1,-1,-1)\) \cite{17} has the correct modular properties in this generalized case.

Let us now compute \( \text{det} \tilde{M} \) for the case analyzed in Eqs. (2.9)–(2.11) above where

\[
\text{det} M = \overline{\Phi}_{23}\Phi_{31}\delta_{44}\delta_{33} - \delta_{34}^2 + \delta_{44}\Phi_{23}. \tag{3.2}
\]

The trilinear masses \( \overline{\Phi}_{23} \) and \( \Phi_{31} \) have modular weights \((-1, 0)\) and \((0, -1)\) respectively and thus we take instead \( \eta^2(T_1)\overline{\Phi}_{23} \) and \( \eta^2(T_2)\Phi_{31} \). The \( \delta_{ij} \) terms come from the following superpotential term

\[
a T_3T_3T_4\phi_{45}^+\eta^2(T_1)\eta^4(T_2), \tag{3.3}
\]

where \( a \) is an \( O(1) \) calculable constant and the \( \eta \) functions insure that this contribution to the superpotential has the appropriate modular weight (i.e., \((-1, -1)\)) \cite{18}. This implies that \( \delta/M \) defined in Eq. (2.5a) should be multiplied by \( \eta^2(T_1)\eta^4(T_2) \) as well. With this information and the modular weights of the fields involved, one can readily determine whether additional \( \eta \) functions are needed to make the mass terms in (3.2) modular invariant. The final result is

\[
\text{det} \tilde{M} = \overline{\Phi}_{23}\Phi_{31}\eta^2(T_1)\eta^4(T_2)[(\delta_{33}\delta_{44} - \delta_{34}^2)\eta^2(T_1) - \delta_{44}\Phi_{23}], \tag{3.4a}
\]

\[
= -\overline{\Phi}_{23}\Phi_{31}^2 a \frac{\phi_{45}\phi^+}{M^3} \frac{Y^3}{32\pi^2}\eta^6(T_1)\eta^8(T_2), \tag{3.4b}
\]

\[\text{1 The modular weights of fields in free fermionic models are given in} \cite{24}. \text{Here we also use} \ T_3, \phi_{45}, \phi^+: (-1/2, -1/2) \text{and} \ T_4: (0, -1/2). \text{Note that we only consider the modular properties under the} \ T_{1,2} \text{moduli fields. The modular transformations under} \ T_3 \text{are less obvious and are still under investigation.} \]
where Eqs. (2.9) and (2.11) have been used.

To make our results for \( Y_{32}^{3} \pi^{2} \) more transparent, we insert the missing units in Eq. (3.1) and set \( M = 5 \). The exponential factor then becomes \( M^{19/8} e^{-32\pi^{2}S/8} \) with \( M \approx 10^{18} \text{GeV} \). Using \( 4g^{2} = 1/S \) and introducing the \( SO(10) \) condensation scale \( \Lambda_{10} = M e^{8\pi^{2}/(3g^{2})} \), with \( \beta = -3 \times 8 + 5 = -19 \), this factor becomes \( \Lambda_{10}^{19/8} \), and thus we can write

\[
\frac{Y_{3}^{3}}{32\pi^{2}} = (32\pi^{2}e)^{-3/8} [c\eta(T_{1})\eta(T_{2})]^{-2} \Lambda_{10}^{3} \left( \frac{\det \tilde{M}}{\Lambda_{10}^{10}} \right)^{1/8},
\]

which is a modular invariant generalization of the usual gaugino condensate expression [1]. Note that since \( \det \tilde{M} \) depends on \( \frac{Y_{3}^{3}}{32\pi^{2}} \) as well, the final expression for \( \frac{Y_{3}^{3}}{32\pi^{2}} \) is different, as follows

\[
\frac{Y_{3}^{3}}{32\pi^{2}} = (32\pi^{2}e)^{-3/7} [c\eta(T_{1})\eta(T_{2})]^{-2} \Lambda_{10}^{3} \left( \frac{\det \tilde{M}'}{\Lambda_{10}^{10}} \right)^{1/7},
\]

with

\[
\det \frac{\tilde{M}'}{\Lambda_{10}} = -\Phi_{23}\Phi_{31}^{2}a\phi_{45}\phi_{6}^{+}\tilde{A}_{10}^{2}M^{2}\eta^{4}(T_{1})\eta^{6}(T_{2}).
\]

Numerically, the \( \Lambda_{10}^{3}(\det \frac{\tilde{M}'}{\Lambda_{10}})^{1/7} \) terms determine the scale of \( \frac{Y_{3}^{3}}{32\pi^{2}} \). For typical values of the parameters we obtain \( \Lambda_{10} \sim 10^{15} \text{GeV} \), \( \det \frac{\tilde{M}'}{\Lambda_{10}} \sim 10 \), and \( \frac{Y_{3}^{3}}{32\pi^{2}} \sim (10^{15} \text{GeV})^{3} \). The matter condensates then become \( \Pi_{11} \sim \Pi_{22} \sim \Pi_{55} \sim (10^{14} \text{GeV})^{2} \), \( \Pi_{44} \sim 0 \), \( \Pi_{33} \sim (10^{18} \text{GeV})^{2} \), and \( \Pi_{34} \sim (10^{16} \text{GeV})^{2} \).

4. Conclusions

Gauge and matter condensates are very important in the analysis of realistic string-derived models, since these typically involve hidden gauge groups which communicate with the observable sector, although rather feebly. The calculations presented in this paper address some of the difficulties that are likely to arise in typical models, such as the degeneracy of the vacuum and the need to explore the superpotential to high orders. The latter requires an understanding of the modular invariant properties of the superpotential to all orders in nonrenormalizable terms. We have made the above points explicit in the particular case of the flipped \( SU(5) \) model with \( SO(10) \) hidden gauge group and matter fields in the fundamental representation. In this case the gaugino condensate has to be solved self-consistently together with the matter condensates, a feature that may only be appreciated in a calculation in an explicit string model.

Acknowledgments: This work has been supported in part by DOE grant DE-FG05-91-ER-40633.
References

[1] For a review see, e.g., D. Amati, et. al., Phys. Rep. 162 (1988) 169.
[2] K. Konishi, Phys. Lett. B 135 (1984) 439.
[3] L. Dixon, in Proceedings of The Rice Meeting, ed. by B. Bonner and H. Miettinen (World Scientific, 1990), p. 811, and references therein.
[4] I. Affleck, M. Dine, and N. Seiberg, Nucl. Phys. B 256 (1985) 557 and references therein.
[5] S. Ferrara, D. Lüst, A. Shapere, and S. Theisen, Phys. Lett. B 225 (1989) 363; S. Ferrara, D. Lüst, and S. Theisen, Phys. Lett. B 233 (1989) 147 and Phys. Lett. B 242 (1990) 39.
[6] J. L. Lopez and D. V. Nanopoulos, Phys. Lett. B 251 (1990) 73.
[7] See e.g., G. G. Ross, Grand Unified Theories, (Benjamin/Cummings, MA, 1983); C. Kounnas, A. Masiero, D. V. Nanopoulos and K. A. Olive, Grand Unification With and Without Supersymmetry and Cosmological Implications, (World Scient. Publ. Comp., Singapore, 1984).
[8] For a review see E. Farhi and L. Susskind, Phys. Rep. 74 (1981) 277.
[9] A. Lahanas and D. V. Nanopoulos, Phys. Rep. 145 (1987) 1.
[10] J. P. Derendinger, L. Ibáñez, and H. Nilles, Phys. Lett. B 155 (1985) 65; M. Dine, R. Rohm, N. Seiberg, and E. Witten, Phys. Lett. B 156 (1985) 55.
[11] J. Ellis, J. L. Lopez, and D. V. Nanopoulos, Phys. Lett. B 247 (1990) 257.
[12] S. Kalara, J. Lopez, and D. V. Nanopoulos, Phys. Lett. B 245 (1990) 421, Nucl. Phys. B 353 (1991) 650.
[13] See e.g., J. Schwarz, Caltech preprint CALT-68-1581 (1989).
[14] I. Antoniadis, J. Ellis, A. B. Lahanas, and D. V. Nanopoulos, Phys. Lett. B 241 (1990) 24; C. P. Burguess and F. Quevedo, Phys. Rev. Lett. 64 (1990) 2611; S. Ferrara, N. Magnoli, T. R. Taylor, and G. Veneziano, Phys. Lett. B 245 (1990) 409; H. P. Nilles and M. Olechowsky, Phys. Lett. B 248 (1990) 268; P. Binetruy and M. K. Gaillard, Phys. Lett. B 253 (1991) 119; M. Cvetic, et. al., Nucl. Phys. B 361 (1991) 194.
[15] D. Lüst and T. Taylor, Phys. Lett. B 253 (1991) 335.
[16] B. de Carlos, J. A. Casas, and C. Muñoz, Phys. Lett. B 263 (1991) 248; J. A. Casas and C. Muñoz, CERN preprint CERN-TH.6187/91.
[17] J. Louis, SLAC preprint SLAC-PUB-5645 (1991).
[18] S. Kalara, J. L. Lopez, and D. V. Nanopoulos, in preparation.
[19] J. Ellis, S. Kalara, K. Olive, and C. Wetterich, Phys. Lett. B 228 (1989) 264.
[20] I. Antoniadis, C. Bachas, and C. Kounnas, Nucl. Phys. B 289 (1987) 87; I. Antoniadis and C. Bachas, Nucl. Phys. B 298 (1988) 586; H. Kawai, D.C. Lewellen, and S.H.-H. Tye, Phys. Rev. Lett. 57 (1986) 1832; Phys. Rev. D 34 (1986) 3794; Nucl. Phys. B 288 (1987) 1; R. Bluhm, L. Dolan, and P. Goddard, Nucl. Phys. B 309 (1988) 330;
H. Dreiner, J. L. Lopez, D. V. Nanopoulos, and D. Reiss, Nucl. Phys. B \textbf{320} (1989) 401.

[21] I. Antoniadis, J. Ellis, J. Hagelin, and D. V. Nanopoulos, Phys. Lett. B \textbf{231} (1989) 65.

[22] G. Leontaris, J. Rizos, and K. Tamvakis, Phys. Lett. B \textbf{243} (1990) 220.

[23] V. Kaplunovsky, Nucl. Phys. B \textbf{307} (1988) 145; I. Antoniadis, J. Ellis, R. Lacaze, and D. V. Nanopoulos, CERN preprint CERN-TH.6136/91 (To appear in Phys. Lett. B).

[24] S. Kalara, J. L. Lopez, and D. V. Nanopoulos, Texas A & M University preprint CTP-TAMU-46/91 (To appear in Phys. Lett. B).

[25] J. Rizos and K. Tamvakis, Phys. Lett. B \textbf{251} (1990) 369.

[26] J. L. Lopez and D. V. Nanopoulos, Phys. Lett. B \textbf{256} (1991) 150.