Gaugino condensation in $\mathcal{N}=1$ supergravity models with multiple dilaton-like fields

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We study supersymmetry breaking by hidden-sector gaugino condensation in $\mathcal{N}=1$, $D=4$ supergravity models with multiple dilaton-like moduli fields. Our work is motivated by type I string theory, in which the low-energy effective Lagrangian can have different dilaton-like fields coupling to different sectors of the theory. We construct the effective Lagrangian for gaugino condensation and use it to compute the visible-sector gaugino masses. We find that the gaugino masses can be of order the gravitino mass, in stark contrast to heterotic string models with a single dilaton field. [S0556-2821(99)06818-6]

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I. INTRODUCTION

Before the recent string duality revolution, most string phenomenology centered on perturbative $\mathcal{N}=1$, $D=4$ heterotic string theories. The discovery of string duality and D-branes, however, opened a variety of new approaches to string phenomenology based on type I and type II string theories. For example, much recent work has focused on the intriguing possibility that our four-dimensional world lies at the intersection of a set of Dp-branes ($3\leqslant p\leqslant 9$) embedded in $(9+1)$-dimensional spacetime.[1]

A well-known problem with the usual perturbative heterotic string phenomenology is that hidden-sector gaugino condensation [4] gives rise to visible-sector gaugino masses that are much smaller than the scale of supersymmetry breaking [5,6]. This is a consequence of the fact that a single dilaton couples to all gauge and matter fields. Gaugino condensation induces a small F term for the dilaton field, and the dilaton couplings then give small masses to the gauginos.

In type I models, however, the situation can be very different. In these models, the hidden and visible sectors can reside on different D-branes. Each sector has its own dilaton-like fields [2]. The hidden-sector dilatons receive small F terms from gaugino condensation. However, these F terms are not responsible for the visible-sector gaugino masses, and therefore the gaugino masses are not forced to be small.[2]

Inspired by this possibility, in this paper we study the question of supersymmetry breaking by hidden-sector gaugino condensation in $\mathcal{N}=1$, $D=4$ supergravity models with multiple dilatonlike fields. We compute the visible-sector gaugino masses and find that they can indeed be of the order of the supersymmetry breaking scale. We see that type I models offer an appealing solution to the gaugino mass problem associated with heterotic string theories.

We approach this problem in the spirit of effective field theory. We take our visible sector to be composed of multiple pure $\mathcal{N}=1$ super Yang-Mills theories, each coupled to gravity, and each to its own dilatonlike field. We take our hidden sector to be composed of gaugino condensate fields, one on each set of branes. The condensate fields are also coupled to gravity and to their associated dilatons.[3] By construction, this theory describes the low-energy limit of a type I string theory, where each super Yang-Mills theory resides on its own set of D-branes. Of course, our analysis also applies to other string or M theory vacua with multiple dilaton-like fields.

Various $\mathcal{N}=1$, $D=4$ type I models have been proposed in the literature [7–19]. Particularly simple examples can be constructed from type I models with D9-branes and D5-branes compactified on $T^2\times T^2\times T^2$, where R_i is the radius of the ith two-torus, $i=1,2,3$. (D5-branes that wrap on the ith two-torus are denoted as $D5_i$-branes.) These models have one dilaton field S, three untwisted moduli fields T_i ($i=1,2,3$) and additional twisted moduli fields from the closed string sector. Other D-brane configurations can be obtained from these by T duality.

One such example is a type I model with two sectors, built from D9-branes and $D5_1$-branes. The gauge bosons arise from open strings ending on the D9- and $D5_1$-branes. The low-energy effective Lagrangian is as follows[4] [1,15]:

$$\begin{aligned}\mathcal{L}=&\frac{1}{8}\int \mathrm{d}^4\theta\frac{E}{R}S(\mathcal{W}^{\alpha}\mathcal{W}_{\alpha})_{D9}+\frac{1}{8}\int \mathrm{d}^4\theta\frac{E}{R^{\dagger}}\overline{S}(\mathcal{W}_{\dot\alpha}\mathcal{W}^{\dot\alpha})_{D9}\\&+\frac{1}{8}\int \mathrm{d}^4\theta\frac{E}{R}T_1(\mathcal{W}^{\alpha}\mathcal{W}_{\alpha})_{D5_1}\\&+\frac{1}{8}\int \mathrm{d}^4\theta\frac{E}{R^{\dagger}}\overline{T}_1(\mathcal{W}_{\dot\alpha}\mathcal{W}^{\dot\alpha})_{D5_1}+\cdots\end{aligned}\tag{1.1}$$

where

$$K=-\ln(S+\overline{S})-\sum_{i=1}^{3}\ln(T_i+\overline{T}_i)+\cdots\tag{1.2}$$

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[1]See [1,2,3] for recent reviews.

[2]This possibility was also observed in [1].

[3]To simplify our presentation, we ignore all charged chiral superfields; this restriction does not affect the results of our analysis.

[4]For convenience, we use the Kähler superspace formulation throughout this paper [20].

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is the Kähler function. The ellipsis denotes possible contributions from other neutral chiral superfields, and

$$\langle S+\bar{S}\rangle=\frac{R_1^2R_2^2R_3^2}{\pi\lambda_I\alpha'^3},\quad \langle T_i+\bar{T}_i\rangle=\frac{R_i^2}{\pi\lambda_I\alpha'},\quad i=1,2,3. \tag{1.3}$$

In these expressions, λ_I is the string coupling, $\alpha'=M_I^{-2}$, and M_I is the string scale. The field strengths $(\mathcal{W}^\alpha\mathcal{W}_\alpha)_{D9}$ and $(\mathcal{W}^\alpha\mathcal{W}_\alpha)_{D5_1}$ contain gauge fields residing on D9- and $D5_1$-branes, respectively. Note that the moduli S, T_1 are dilaton-like fields, while T_2, T_3 are not.[5]

A second example contains D9-branes as well as $D5_i$-branes compactified on all three tori $(i=1,2,3)$. It has four sectors, and four dilaton-like fields, S, T_1, T_2, T_3, as defined in Eq. (1.3). These fields couple to the gauge fields on the D9- and $D5_i$-branes and give rise to the effective Lagrangian

$$\begin{aligned}\mathcal{L}=&\frac{1}{8}\int d^4\theta\frac{E}{R}S(\mathcal{W}^\alpha\mathcal{W}_\alpha)_{D9}+\frac{1}{8}\int d^4\theta\frac{E}{R^\dagger}\bar{S}(\mathcal{W}_{\dot\alpha}\mathcal{W}^{\dot\alpha})_{D9}\\&+\sum_{i=1}^{3}\frac{1}{8}\int d^4\theta\frac{E}{R}T_i(\mathcal{W}^\alpha\mathcal{W}_\alpha)_{D5_i}\\&+\sum_{i=1}^{3}\frac{1}{8}\int d^4\theta\frac{E}{R^\dagger}\bar{T}_i(\mathcal{W}_{\dot\alpha}\mathcal{W}^{\dot\alpha})_{D5_i}+\cdots.\end{aligned} \tag{1.4}$$

The plan of this paper is as follows. In Sec. II, we define our $\mathcal{N}=1$, $D=4$ supergravity model with multiple dilaton-like fields. In Sec. III, we argue that the extra dilaton-like fields allow gaugino masses to be as large as the gravitino mass. We also present an explicit example of this scenario.

In Appendix A, we exhibit the linear-chiral duality for $\mathcal{N}=1$, $D=4$ supergravity models with multiple linear supermultiplets. In Appendix B we extract the relevant pieces of the component supergravity Lagrangian. Finally, in Appendix C we check our results by comparing with the chiral supermultiplet formulation.

II. SUPERGRAVITY WITH MULTIPLE DILATON-LIKE FIELDS

In this section we define the $\mathcal{N}=1, D=4$ supergravity model that we will study. We start by considering a system with N different types of D-branes. Of these, we take $\tilde{N}$ to be $Dp_{\tilde{A}}$-branes $(\tilde{A}=1,\ldots,\tilde{N})$ with weakly coupled visible-sector fields on their world volumes. We take the remaining $N-\tilde{N}$ to be Dp_A-branes $(A=1,\ldots,N-\tilde{N})$ with strongly coupled hidden-sector fields in the condensation phase.

Let us first construct the effective theory of the visible sector. For each value of $\tilde{A}$, let $\mathcal{W}_{\tilde{A}}$ denote the super Yang-Mills field strength on the $\tilde{A}$th set of branes, and let $\tilde{V}_{\tilde{A}}$ be the real superfield which contains the $\tilde{A}$th dilaton-like field [21]. The field $\tilde{V}_{\tilde{A}}$ obeys the constraint

$$-(\bar{\mathcal{D}}^2-8R)\tilde{V}_{\tilde{A}}=\mathcal{W}_{\tilde{A}}\mathcal{W}_{\tilde{A}}. \tag{2.1}$$

This constraint couples the dilatons to the super Yang-Mills fields.

In Kähler superspace [20], the supergravity kinetic terms are given by a Lagrangian $\mathcal{L}$ and a Kähler function K. For the visible-sector fields, we take them to be

$$\begin{aligned}\mathcal{L}&=\int d^4\theta E\left\{-3+\tilde{N}+\sum_{\tilde{A}=1}^{\tilde{N}}\tilde{f}_{\tilde{A}}(\tilde{V}_{\tilde{A}})\right\},\\K&=\sum_{\tilde{A}=1}^{\tilde{N}}\{\ln\tilde{V}_{\tilde{A}}+\tilde{g}_{\tilde{A}}(\tilde{V}_{\tilde{A}})\},\end{aligned} \tag{2.2}$$

where

$$\tilde{V}_{\tilde{A}}\frac{d\tilde{g}_{\tilde{A}}}{d\tilde{V}_{\tilde{A}}}=\tilde{f}_{\tilde{A}}-\tilde{V}_{\tilde{A}}\frac{d\tilde{f}_{\tilde{A}}}{d\tilde{V}_{\tilde{A}}}. \tag{2.3}$$

In these expressions, the leading terms describe the tree-level couplings of the gauge and dilaton-like fields; the functions $\tilde{g}_{\tilde{A}}(\tilde{V}_{\tilde{A}})$ and $\tilde{f}_{\tilde{A}}(\tilde{V}_{\tilde{A}})$ contain corrections beyond tree level. The condition (2.3) guarantees that the Einstein gravity term is canonically normalized. (For simplicity, we do not include mixings between the different $\tilde{V}_{\tilde{A}}$'s. Such mixings arise at the loop level, but they do not change our conclusions.[6])

In the hidden sector, the effective Lagrangian is different because the super Yang-Mills fields are in a strongly coupled condensation phase. This leads us to replace the Yang-Mills fields by chiral condensate superfields U_A [22]. These fields are contained within the real superfields V_A [23,24], where

$$-(\bar{\mathcal{D}}^2-8R)V_A=U_A. \tag{2.4}$$

The hidden-sector dilaton-like fields are also contained in the fields V_A.

The effective Lagrangian for this sector contains kinetic and superpotential terms. The kinetic terms are as above,

$$\begin{aligned}\mathcal{L}&=\int d^4\theta E\left\{-3+N-\tilde{N}+\sum_{A=1}^{N-\tilde{N}}f_A(V_A)\right\},\\K&=\sum_{A=1}^{N-\tilde{N}}\{\ln V_A+g_A(V_A)\},\end{aligned} \tag{2.5}$$

[5] Our definition of the S and T_i moduli agrees with [20], but differs from [1].

[6] Type I models typically contain chiral superfields charged under the $Dp_{\tilde{A}}$ and $Dp_{\tilde{A}'}$ gauge groups $(\tilde{A}\neq\tilde{A}')$. These superfields mix the different dilaton-like fields. Our study ignores charged chiral superfields, so it is consistent to ignore mixings between different $\tilde{V}_{\tilde{A}}$'s.

where

$$V_A \frac{dg_A}{dV_A} = f_A - V_A \frac{df_A}{dV_A}. \tag{2.6}$$

The leading terms describe the tree-level couplings, while $g_A(V_A)$ and $f_A(V_A)$ contain corrections beyond the tree level.

The superpotential terms are generated by nonperturbative effects associated with gaugino condensation. In what follows, we take the superpotential to be given by

$$\begin{aligned}\mathcal{L}_A &= \int d^4\theta \frac{E}{R} \frac{1}{8} b_A U_A \ln(e^{-K/2} U_A) + \text{H.c.}\\ &= \int d^4\theta E b_A V_A \ln(e^{-K} \overline{U}_A U_A),\end{aligned} \tag{2.7}$$

where $b_A = 2b'_A/3$, and b'_A is the one-loop β-function coefficient of Dp_A-sector. The form of this term is dictated by the anomalies of the underlying super Yang-Mills theory [4,22,25,26,27].

In rigid supersymmetry, the superpotential is fixed by the chiral and conformal anomalies [22]. In local supersymmetry, the Kähler anomaly also comes into play [28,29]. In particular, it fixes the explicit K dependence of the superpotential. Under an arbitrary Kähler transformation, $K \to K + F + \overline{F}$, the fields $V_A \to V_A$ and $U_A \to U_A e^{(\overline{F}-F)/2}$. The superpotential then transforms as follows:

$$\begin{aligned}\mathcal{L}_A &\equiv \int d^4\theta E b_A V_A \ln(e^{-K} \overline{U}_A U_A)\\ &\to \mathcal{L}_A - \int d^4\theta \frac{E}{R} \frac{1}{8} b_A U_A F - \int d^4\theta \frac{E}{R^\dagger} \frac{1}{8} b_A \overline{U}_A \overline{F}.\end{aligned} \tag{2.8}$$

This is precisely the right transformation to match the Kähler anomaly of the underlying super Yang-Mills theory [4,25,26,27].

The full supergravity Lagrangian contains contributions from both of these sectors. It can be written as follows:[7]

$$\begin{aligned}\mathcal{L} = &\int d^4\theta E \left\{ -3 + N + \sum_{\tilde{A}=1}^{\tilde{N}} \tilde{f}_{\tilde{A}}(\tilde{V}_{\tilde{A}}) + \sum_{A=1}^{N-\tilde{N}} f_A(V_A) \right\}\\ &+ \sum_{A=1}^{N-\tilde{N}} \int d^4\theta E b_A V_A \ln(e^{-K} \overline{U}_A U_A),\end{aligned} \tag{2.9}$$

where the Kähler function is given by

$$\begin{aligned}K = &\sum_{\tilde{A}=1}^{\tilde{N}} \{\ln \tilde{V}_{\tilde{A}} + \tilde{g}_{\tilde{A}}(\tilde{V}_{\tilde{A}})\} + \sum_{A=1}^{N-\tilde{N}} \{\ln V_A + g_A(V_A)\}\\ &+ G(\Phi_1, \overline{\Phi}_1, \ldots, \Phi_n, \overline{\Phi}_n)\end{aligned} \tag{2.10}$$

and

$$\begin{aligned}&\tilde{V}_{\tilde{A}} \frac{d\tilde{g}_{\tilde{A}}}{d\tilde{V}_{\tilde{A}}} = \tilde{f}_{\tilde{A}} - \tilde{V}_{\tilde{A}} \frac{d\tilde{f}_{\tilde{A}}}{d\tilde{V}_{\tilde{A}}}, \quad \tilde{A} = 1, \ldots, \tilde{N},\\ &V_A \frac{dg_A}{dV_A} = f_A - V_A \frac{df_A}{dV_A}, \quad A = 1, \ldots, N - \tilde{N}.\end{aligned} \tag{2.11}$$

The hidden-sector superpotential couples all sectors through its explicit K dependence; it gives rise to gaugino masses in the visible sector.

III. GAUGINO MASSES

To find the gaugino masses, we need certain terms from the component-field Lagrangian. These terms are computed in Appendix B.

From Eq. (B4) we find that the scalar potential takes the following form:

$$\begin{aligned}\mathcal{V} = &\sum_{A=1}^{N-\tilde{N}} \frac{1}{16 l_A^2} (1 + f^A - l_A f_l^A) \overline{u}_A u_A + \frac{1}{16} \left\{ \left(\sum_{A=1}^{N-\tilde{N}} \frac{(1 + f^A - l_A f_l^A)}{l_A} u_A \right) \left(\sum_{B=1}^{N-\tilde{N}} b_B \overline{u}_B \right) + \text{H.c.} \right\}\\ &+ \frac{1}{16} \left\{ \begin{array}{l} -3 + \sum_{i,j} G_{\bar{j}} G^{-1}_{\bar{j}i} G_i \\ + \sum_{\tilde{A}=1}^{\tilde{N}} (1 + \tilde{f}^{\tilde{A}} - \tilde{l}_{\tilde{A}} \tilde{f}_{\tilde{l}}^{\tilde{A}}) + \sum_{A=1}^{N-\tilde{N}} (1 + f^A - l_A f_l^A) \end{array} \right\} \left| \sum_{A=1}^{N-\tilde{N}} b_A \overline{u}_A \right|^2.\end{aligned} \tag{3.1}$$

[7]Note also that these expressions contain contributions from extra moduli fields Φ_i, $i = 1, \ldots, n$. In the context of type I models, the Φ_i are twisted or nondilaton-like untwisted moduli fields.

In this expression, $\tilde{l}_{\tilde{A}}=\tilde{V}_{\tilde{A}}|_{\theta=\bar{\theta}=0}$ and $l_A=V_A|_{\theta=\bar{\theta}=0}$ are the dilaton-like scalar fields which couple to the visible and hidden sectors, respectively. Moreover, $u_A=U_A|_{\theta=\bar{\theta}=0}$ is the gaugino condensate field on each of the Dp_A-branes; it depends on l_A according to Eq. (B7). Other notation is as follows:

$$g_l^A=\left.\frac{\mathrm{d}g_A(V_A)}{\mathrm{d}V_A}\right|_{\theta=\bar{\theta}=0},\quad \tilde{g}_{\tilde{l}}^{\tilde{A}}=\left.\frac{\mathrm{d}\tilde{g}_{\tilde{A}}(\tilde{V}_{\tilde{A}})}{\mathrm{d}\tilde{V}_{\tilde{A}}}\right|_{\theta=\bar{\theta}=0},$$

$$f_l^A=\left.\frac{\mathrm{d}f_A(V_A)}{\mathrm{d}V_A}\right|_{\theta=\bar{\theta}=0},\quad \tilde{f}_{\tilde{l}}^{\tilde{A}}=\left.\frac{\mathrm{d}\tilde{f}_{\tilde{A}}(\tilde{V}_{\tilde{A}})}{\mathrm{d}\tilde{V}_{\tilde{A}}}\right|_{\theta=\bar{\theta}=0}. \tag{3.2}$$

$$G_i=\left.\frac{\partial G}{\partial \Phi_i}\right|_{\theta=\bar{\theta}=0},\quad G_{\bar{j}}=\left.\frac{\partial G}{\partial \bar{\Phi}_j}\right|_{\theta=\bar{\theta}=0},$$

$$G_{i\bar{j}}=\left.\frac{\partial^2 G}{\partial \Phi_i\partial \bar{\Phi}_j}\right|_{\theta=\bar{\theta}=0}.$$

The gaugino mass has its origin in the explicit K dependence of the hidden-sector superpotential. We find Eq. (B5)

$$\begin{aligned} m_\lambda^{\tilde{A}}&=\left\langle \frac{1+\tilde{f}^{\tilde{A}}-\tilde{l}_{\tilde{A}}\tilde{f}_{\tilde{l}}^{\tilde{A}}}{1+\tilde{f}^{\tilde{A}}}\cdot\sum_{A=1}^{N-\tilde{N}}\frac{1}{4}b_A u_A\right\rangle,\\ &=\left\langle \frac{1+\tilde{f}^{\tilde{A}}-\tilde{l}_{\tilde{A}}\tilde{f}_{\tilde{l}}^{\tilde{A}}}{1+\tilde{f}^{\tilde{A}}}\right\rangle m_{\tilde{G}},\quad \tilde{A}=1,\ldots,\tilde{N}. \end{aligned}\tag{3.3}$$

The gravitino mass is simply

$$m_{\tilde{G}}=\left\langle \sum_{A=1}^{N-\tilde{N}}\frac{1}{4}b_A u_A\right\rangle. \tag{3.4}$$

To understand these formulas, let us first examine the gaugino mass in a theory with just one dilaton [6,30,31]. In this case the gaugino mass is as follows [6]:

$$m_\lambda=\left\langle\frac{1+bl}{bl}\right\rangle\left\langle\frac{1+f-lf_l}{1+f}\right\rangle m_{\tilde{G}}. \tag{3.5}$$

All scalar fields are evaluated at the minimum of the potential,

$$\mathcal{V}=\frac{1}{16l^2}\{(1+f-lf_l)(1+bl)^2-3b^2l^2\}\bar{u}u. \tag{3.6}$$

From this we see that any vacuum with zero cosmological constant satisfies $\langle 1+f-lf_l\rangle=3b^2\langle l^2\rangle+\mathcal{O}(b^3)$. Since $\langle 1+f\rangle$ and $\langle l\rangle$ are numbers of order 1, Eq. (3.5) implies that m_λ is smaller than $m_{\tilde{G}}$ by a factor of b.

Let us compare this to a model in which the visible and hidden sectors couple to different dilaton-like fields. The scalar potential of Eq. (3.1) receives contributions from both sectors, proportional to $\langle 1+\tilde{f}^{\tilde{A}}-\tilde{l}_{\tilde{A}}\tilde{f}_{\tilde{l}}^{\tilde{A}}\rangle$ and $\langle 1+f^A-l_Af_l^A\rangle$. However, to leading order, only $\langle 1+\tilde{f}^{\tilde{A}}-\tilde{l}_{\tilde{A}}\tilde{f}_{\tilde{l}}^{\tilde{A}}\rangle$ contributes to the gaugino masses. As above, minimizing the potential forces $\langle 1+f^A-l_Af_l^A\rangle$ to be $\mathcal{O}(b_A^2)$, but it does not constrain $\langle 1+\tilde{f}^{\tilde{A}}-\tilde{l}_{\tilde{A}}\tilde{f}_{\tilde{l}}^{\tilde{A}}\rangle$. This suggests that the gaugino masses can indeed be as large as the gravitino mass.

To see this explicitly, let us consider an example with just two types of D-branes ($\tilde{N}=1$ and $N=2$). Each has its own dilaton-like field, $\tilde{l}_{\tilde{A}}$ and l_A. The scalar potential follows from Eq. (3.1), with no summation over the indices A and $\tilde{A}$:

$$\begin{aligned}\mathcal{V}=&\frac{1}{16l_A^2}\{(1+f^A-l_Af_l^A)(1+b_Al_A)^2+b_A^2l_A^2\\ &\times[(1+\tilde{f}^{\tilde{A}}-\tilde{l}_{\tilde{A}}\tilde{f}_{\tilde{l}}^{\tilde{A}})-3]\}\bar{u}_Au_A.\end{aligned}\tag{3.7}$$

The conditions for a nontrivial vacuum with vanishing cosmological constant are

$$\begin{aligned}&\left\langle\frac{\partial}{\partial\tilde{l}_{\tilde{A}}}(1+\tilde{f}^{\tilde{A}}-\tilde{l}_{\tilde{A}}\tilde{f}_{\tilde{l}}^{\tilde{A}})\right\rangle=0,\\ &\left\langle\frac{\partial}{\partial l_A}(1+f^A-l_Af_l^A)\right\rangle\\ &\quad=2b_A^2\langle l_A\rangle[3-\langle 1+\tilde{f}^{\tilde{A}}-\tilde{l}_{\tilde{A}}\tilde{f}_{\tilde{l}}^{\tilde{A}}\rangle]+\mathcal{O}(b_A^3),\\ &\langle 1+f^A-l_Af_l^A\rangle=b_A^2\langle l_A^2\rangle[3-\langle 1+\tilde{f}^{\tilde{A}}-\tilde{l}_{\tilde{A}}\tilde{f}_{\tilde{l}}^{\tilde{A}}\rangle]+\mathcal{O}(b_A^3),\end{aligned}\tag{3.8}$$

where all $\mathcal{O}(b_A^3)$ terms are suppressed. These equations have a consistent solution if $\langle 1+f^A-l_Af_l^A\rangle=\mathcal{O}(b_A^2)$ and $\langle 1+\tilde{f}^{\tilde{A}}-\tilde{l}_{\tilde{A}}\tilde{f}_{\tilde{l}}^{\tilde{A}}\rangle$ is $\mathcal{O}(1)$. Since this latter term fixes the visible-sector gaugino mass, we expect $m_\lambda^{\tilde{A}}$ to be of order $m_{\tilde{G}}$.

To actually compute the gaugino mass, we need to specify the functions f^A and $\tilde{f}^{\tilde{A}}$. [The functions g_A and $\tilde{g}_{\tilde{A}}$ are determined via Eq. (2.11)]. We choose them to stabilize the runaway vacuum typically associated with dilaton-like fields.[8] The general procedure is described in [30,34]. For now, we consider the following simple choice:[9]

$$f_A=\mathcal{P}\cdot e^{-\mathcal{Q}/l_A},\quad \tilde{f}_{\tilde{A}}=\tilde{\mathcal{P}}\cdot e^{-\tilde{\mathcal{Q}}/\tilde{l}_{\tilde{A}}},\quad \mathcal{P},\mathcal{Q},\tilde{\mathcal{P}},\tilde{\mathcal{Q}}>0. \tag{3.9}$$

Substituting Eqs. (3.9) into Eqs. (3.8), we find that the potential has an extremum, located at

[8]See [32,33] for recent reviews.

[9]This choice is motivated by the type I string theory discussed in [35].

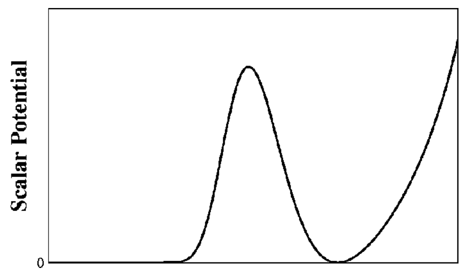

FIG. 1. The scalar potential, $\mathcal{V}$, plotted versus l_A for fixed $\tilde{l}_{\tilde{A}} = \langle \tilde{l}_{\tilde{A}} \rangle$.

$$\langle \tilde{l}_{\tilde{A}} \rangle = \frac{\tilde{\mathcal{Q}}}{2},$$

$$\langle l_A \rangle = \frac{\mathcal{Q}}{2} + \mathcal{O}(b_A^2), \tag{3.10}$$

$$\mathcal{P} = e^2 + \mathcal{O}(b_A^2),$$

for any value of $\tilde{\mathcal{P}}$. It is, in fact, a global minimum of the potential with vanishing cosmological constant, as illustrated in Figs. 1 and 2.

Let us now evaluate the gaugino mass in this vacuum. We find

$$m_\lambda^{\tilde{A}} = \frac{e^2 - \tilde{\mathcal{P}}}{e^2 + \tilde{\mathcal{P}}} m_{\tilde{G}}, \tag{3.11}$$

where

$$0.2 m_{\tilde{G}} < m_\lambda^{\tilde{A}} < 0.8 m_{\tilde{G}} \quad \text{for } 4.9 > \tilde{\mathcal{P}} > 0.8. \tag{3.12}$$

We see that, for reasonable values of $\tilde{\mathcal{P}}$, the gaugino mass is of the order of the gravitino mass.

IV. CONCLUSION

The presence of multiple dilaton-like moduli fields is a very important feature of $\mathcal{N}=1, D=4$ type I string models. The extra dilaton-like fields can change the resulting phenomenology in many ways. In this paper we examined their effect on supersymmetry breaking by hidden-sector gaugino condensation.

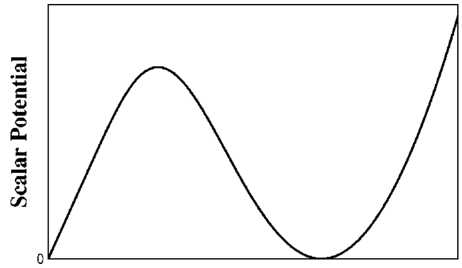

FIG. 2. The scalar potential, $\mathcal{V}$, plotted versus $\tilde{l}_{\tilde{A}}$ for fixed $l_A = \langle l_A \rangle$.

We studied a scenario in which different dilaton-like fields couple to the hidden and visible sectors. We assumed that supersymmetry is broken by gaugino condensation in the hidden sector. We found that the visible-sector gaugino masses can be as large as gravitino mass because of the extra dilaton-like fields. Our results stand in contrast to the usual heterotic string phenomenology, where the gaugino masses are typically much smaller.

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APPENDIX A: LINEAR-CHIRAL DUALITY

In this appendix we prove that a supergravity model with N linear supermultiplets is dual to another supergravity model with N chiral supermultiplets.[10] We will show that both models can be obtained from the following Lagrangian:

$$\mathcal{L} = \int d^4\theta E \left\{ F(V_1, \dots, V_N) + \sum_{A=1}^{N} (S_A + \overline{S_A})(\Omega_A - V_A) \right\} \tag{A1}$$

with

$$K = K(V_1, \dots, V_N) \tag{A2}$$

and

$$-3 + \sum_{A=1}^{N} V_A \frac{\partial K}{\partial V_A} = F - \sum_{A=1}^{N} V_A \frac{\partial F}{\partial V_A}. \tag{A3}$$

In these expressions, the V_A's are *unconstrained* real superfields and the S_A's are ordinary chiral supermultiplets, for $A=1,\dots,N$. The field Ω_A is a Chern-Simons superform; it obeys

$$-(\overline{\mathcal{D}}^2 - 8R)\Omega_A = \mathcal{W}_A \mathcal{W}_A, \tag{A4}$$

for $A=1,\dots,N$.

Let us first integrate out the S_A's in $\mathcal{L}$. Their equations of motion are as follows:

$$-(\overline{\mathcal{D}}^2 - 8R)V_A = \mathcal{W}_A \mathcal{W}_A. \tag{A5}$$

This equation can be used to eliminate the second term in Eq. (A1), reducing $\mathcal{L}$ to a model with N linear supermultiplets. Note that the Einstein gravity term is canonically normalized because of Eq. (A3).

Let us now return to $\mathcal{L}$ and integrate out the V_A's. Their equations of motion are as follows:

$$S_A + \overline{S_A} = \frac{\partial F}{\partial V_A} - \frac{1}{3} \frac{\partial K}{\partial V_A} \left\{ F - \sum_{B=1}^{N} V_B (S_B + \overline{S_B}) \right\}. \tag{A6}$$

[10] This duality was briefly discussed in [36].

If we multiply Eq. (A6) by V_A and sum over $A=1,\ldots,N$, we find[11]

$$F(V_1,\ldots,V_N)-\sum_{A=1}^{N} V_A(S_A+\overline{S_A})=-3. \tag{A7}$$

It is now trivial to use Eq. (A7) to rewrite Eq. (A6) in a very simple form:

$$S_A+\overline{S_A}=\frac{\partial(K+F)}{\partial V_A}. \tag{A8}$$

Using these relations, we find

$$\mathcal{L}=-3\int d^4\theta E+\left\{\sum_{A=1}^{N}\frac{1}{8}\int d^4\theta\frac{E}{R}S_A(\mathcal{W}^{\alpha}\mathcal{W}_{\alpha})_A+\text{H.c.}\right\}, \tag{A9}$$

with

$$K=K(S_1+\overline{S_1},\ldots,S_N+\overline{S_N}). \tag{A10}$$

This completes the proof of linear-chiral duality.

APPENDIX B: COMPONENT-FIELD LAGRANGIAN

In this appendix, we compute the necessary elements of the component-field Lagrangian corresponding to the superfield Lagrangian (2.10)–(2.11). We use the chiral density method [30,37].

We start by enumerating the definitions of bosonic component fields. In the hidden sector, we have

$$\begin{aligned} l_A&=V_A|_{\theta=\bar\theta=0},\\ 2\sigma^m_{\alpha\dot\alpha}B^A_m-\frac{4}{3}l^A\sigma^a_{\alpha\dot\alpha}b_a&=[\mathcal{D}_\alpha,\mathcal{D}_{\dot\alpha}]V^A|_{\theta=\bar\theta=0},\\ u_A=U_A|_{\theta=\bar\theta=0}&\equiv-(\overline{\mathcal{D}}^2-8R)V_A|_{\theta=\bar\theta=0},\\ -4F_A&=-\mathcal{D}^2(\overline{\mathcal{D}}^2-8R)V_A|_{\theta=\bar\theta=0}. \end{aligned} \tag{B1}$$

In these expressions, the l_A are dilaton-like scalar fields, and the B^A_m are axionic degrees of freedom in the same supermultiplets. The fields u_A are the gaugino condensate fields of the hidden sector.

The visible-sector fields are defined in a similar way:[12]

$$\begin{aligned} &\tilde l_{\tilde A}=\tilde V_{\tilde A}|_{\theta=\bar\theta=0},\\ &2\sigma^m_{\alpha\dot\alpha}\tilde B^{\tilde A}_m-\frac{4}{3}\tilde l^{\tilde A}\sigma^a_{\alpha\dot\alpha}b_a+2\,\mathrm{Tr}(\lambda^{\tilde A}_\alpha\overline{\lambda}^{\tilde A}_{\dot\alpha})\\ &\qquad=[\mathcal{D}_\alpha,\mathcal{D}_{\dot\alpha}]\tilde V^{\tilde A}|_{\theta=\bar\theta=0},-\mathrm{Tr}(\lambda^{\tilde A}\lambda^{\tilde A})\\ &\qquad=-(\overline{\mathcal{D}}^2-8R)\tilde V_{\tilde A}|_{\theta=\bar\theta=0},\\ &8i\,\mathrm{Tr}(\lambda^{\tilde A}\sigma^m\mathcal{D}_m\overline{\lambda}^{\tilde A})+4\,\mathrm{Tr}(\lambda^{\tilde A}\lambda^{\tilde A})\overline{M}+2\,\mathrm{Tr}(\tilde{\mathcal{F}}^{mn}_{\tilde A}\tilde{\mathcal{F}}^{\tilde A}_{mn})\\ &\quad+i\epsilon^{mnpq}\,\mathrm{Tr}(\tilde{\mathcal{F}}^{\tilde A}_{mn}\tilde{\mathcal{F}}^{\tilde A}_{pq})-4\,\mathrm{Tr}(\tilde D^{\tilde A}\tilde D^{\tilde A})\\ &\qquad=-\mathcal{D}^2(\overline{\mathcal{D}}^2-8R)\tilde V_{\tilde A\theta=\bar\theta=0}, \end{aligned} \tag{B2}$$

where

$$\begin{aligned} &-i\lambda^{\tilde A}_\alpha=\mathcal{W}^{\tilde A}_\alpha|_{\theta=\bar\theta=0},\qquad i\overline{\lambda}^{\tilde A}_{\dot\alpha}=\mathcal{W}^{\tilde A}_{\dot\alpha}|_{\theta=\bar\theta=0},\\ &-2\tilde D^{\tilde A}=\mathcal{D}^\alpha\mathcal{W}^{\tilde A}_\alpha|_{\theta=\bar\theta=0}=\mathcal{D}_{\dot\alpha}\mathcal{W}^{\dot\alpha}_{\tilde A\theta=\bar\theta=0},\\ &\tilde B^m_{\tilde A}=\frac{1}{2}\epsilon^{mnpq}\left\{\partial_q\tilde b^{\tilde A}_{pn}+\frac{1}{6}\mathrm{Tr}\left(\tilde a_{[q}\partial_p\tilde a_{n]}-\frac{2i}{3}\tilde a_{[q}\tilde a_p\tilde a_{n]}\right)^{\tilde A}\right\}. \end{aligned} \tag{B3}$$

In these expressions, the $\tilde l_{\tilde A}$ are dilaton-like scalar fields. The $\tilde B^{\tilde A}_m$ are dual field strengths of the antisymmetric tensors $\tilde b^{\tilde A}_{pq}$. The $\tilde{\mathcal{F}}^{\tilde A}_{mn}$ are the Yang-Mills field strengths, while the $\tilde a^{\tilde A}_m$ are the corresponding gauge fields. The fields M, $\overline{M}$, and b_a are the auxiliary fields of the supergravity multiplet [38]. The bosonic components of the Φ_i are $\phi_i=\Phi_i|_{\theta=\bar\theta=0}$ and $-4F_i=\mathcal{D}^2\Phi_i|_{\theta=\bar\theta=0}$.

Using these definitions, we find the following bosonic component-field Lagrangian:

[11]Equation (A7) is obtained by assuming $\Sigma_{A=1}^{N}V_A(\partial K/\partial V_A)\neq 3$, as is true for the models considered in this paper.

[12]We include the gaugino fields for the reader's convenience.

$$
\begin{aligned}
\frac{1}{\sqrt{-g}}\mathcal{L}_B = & -\frac{1}{4}\mathcal{R} - \sum_{\tilde{A}=1}^{\tilde{N}} \frac{1}{8\tilde{l}_{\tilde{A}}^2}(1+\tilde{l}_{\tilde{A}}\tilde{g}_{\tilde{l}}^{\tilde{A}})\nabla^m \tilde{l}_{\tilde{A}}\nabla_m \tilde{l}_{\tilde{A}} + \sum_{\tilde{A}=1}^{\tilde{N}} \frac{1}{8\tilde{l}_{\tilde{A}}^2}(1+\tilde{l}_{\tilde{A}}\tilde{g}_{\tilde{l}}^{\tilde{A}})\tilde{B}^m_{\tilde{A}}\tilde{B}^{\tilde{A}}_m - \sum_{\tilde{A}=1}^{\tilde{N}} \frac{1+\tilde{f}_{\tilde{A}}}{16\tilde{l}_{\tilde{A}}}\mathrm{Tr}(\tilde{\mathcal{F}}^{mn}_{\tilde{A}}\,\tilde{\mathcal{F}}^{\tilde{A}}_{mn}) \\
& - \sum_{A=1}^{N-\tilde{N}} \frac{1}{8l_A^2}(1+l_A g_l^A)\nabla^m l_A \nabla_m l_A + \sum_{A=1}^{N-\tilde{N}} \frac{1}{8l_A^2}(1+l_A g_l^A) B^m_A B^A_m + \sum_{A=1}^{N-\tilde{N}} \frac{i}{4} b_A B^m_A \nabla_m \ln\left(\frac{\bar{u}_A}{u_A}\right) - \frac{1}{2}\sum_{i,j} G_{i\bar{j}}\nabla^m\phi_i\nabla_m\bar{\phi}_j \\
& - \frac{1}{18}\left(N-3+\sum_{\tilde{A}=1}^{\tilde{N}}\tilde{l}_{\tilde{A}}\tilde{g}_{\tilde{l}}^{\tilde{A}} + \sum_{A=1}^{N-\tilde{N}} l_A g_l^A\right) b^a b_a + \sum_{\tilde{A}=1}^{\tilde{N}} \frac{1+\tilde{f}_{\tilde{A}}}{8\tilde{l}_{\tilde{A}}}\mathrm{Tr}(\tilde{D}^{\tilde{A}}\tilde{D}^{\tilde{A}}) + \frac{1}{2}\sum_{i,j} G_{i\bar{j}}F_i\bar{F}_j - \frac{1}{4}\left(\sum_{A=1}^{N-\tilde{N}} b_A u_A\right)\sum_i G_i F_i \\
& + \sum_{A=1}^{N-\tilde{N}} \frac{1}{8l_A}\{1+f_A+b_A l_A \ln(e^{-K}\bar{u}_A u_A) + 2b_A l_A\}F_A + \frac{1}{18}\left\{N-3+\sum_{\tilde{A}=1}^{\tilde{N}}(\tilde{f}^{\tilde{A}} - \tilde{l}_{\tilde{A}}\tilde{f}_{\tilde{l}}^{\tilde{A}}) + \sum_{A=1}^{N-\tilde{N}}(f^A - l_A f_l^A)\right\}\bar{M}M \\
& - \sum_{A=1}^{N-\tilde{N}} \frac{1}{8l_A}\{1+f_A+b_A l_A \ln(e^{-K}\bar{u}_A u_A)\}u_A\bar{M} - \frac{1}{12}\left(\sum_{B=1}^{N-\tilde{N}} b_B\bar{u}_B\right)\left\{N+\sum_{\tilde{A}=1}^{\tilde{N}}(\tilde{f}^{\tilde{A}} - \tilde{l}_{\tilde{A}}\tilde{f}_{\tilde{l}}^{\tilde{A}}) + \sum_{A=1}^{N-\tilde{N}}(f^A - l_A f_l^A)\right\}M \\
& - \sum_{A=1}^{N-\tilde{N}} \frac{1}{32l_A^2}(1+f^A - l_A f_l^A)\bar{u}_A u_A - \frac{1}{16}\left(\sum_{B=1}^{N-\tilde{N}} b_B\bar{u}_B\right)\left\{\sum_{A=1}^{N-\tilde{N}}\frac{(1+f^A-l_A f_l^A)}{l_A}u_A\right\} + \text{H.c.}
\end{aligned}
\tag{B4}
$$

We also find the following kinetic and mass terms for the gauginos and gravitino:

$$
\begin{aligned}
\frac{1}{\sqrt{-g}}\mathcal{L}_{m_\lambda} = & -\sum_{\tilde{A}=1}^{\tilde{N}} \frac{1+\tilde{f}_{\tilde{A}}}{4\tilde{l}_{\tilde{A}}}\mathrm{Tr}(i\lambda^{\tilde{A}}\sigma^m\nabla_m\bar{\lambda}^{\tilde{A}}) \\
& + \frac{1}{16}\left(\sum_{A=1}^{N-\tilde{N}} b_A\bar{u}_A\right) \\
& \times\left\{\sum_{\tilde{A}=1}^{\tilde{N}}\frac{(1+\tilde{f}^{\tilde{A}} - \tilde{l}_{\tilde{A}}\tilde{f}_{\tilde{l}}^{\tilde{A}})}{\tilde{l}_{\tilde{A}}}\mathrm{Tr}(\lambda^{\tilde{A}}\lambda^{\tilde{A}})\right\} + \text{H.c.}
\end{aligned}
$$

$$
\begin{aligned}
\frac{1}{\sqrt{-g}}\mathcal{L}_{m_{\tilde{G}}} = & \frac{1}{2}\epsilon^{mnpq}\bar{\psi}_m\bar{\sigma}_n\nabla_p\psi_q - \sum_{A=1}^{N-\tilde{N}}\frac{1}{8l_A}\{1+f_A \\
& + b_A l_A \ln(e^{-K}\bar{u}_A u_A)\}\bar{u}_A(\psi_m\sigma^{mn}\psi_n) + \text{H.c.}
\end{aligned}
\tag{B5}
$$

To extract the potential (3.1), we must eliminate the auxiliary fields. We will not do that here, except to note that the equations of motion for the auxiliary fields $(F_A+\bar{F}_A)$ are

$$1+f_A+b_A l_A \ln(e^{-K}\bar{u}_A u_A) + 2b_A l_A = 0. \tag{B6}$$

This fixes the modulus of the condensate field u_A to be

$$u_A\bar{u}_A = \exp\left[(K-2) - \left(\frac{1+f_A}{b_A l_A}\right)\right]. \tag{B7}$$

The modulus has the correct dependence on the gauge group and gauge coupling, as expected from the usual renormalization group arguments [6,31].

APPENDIX C: GAUGINO MASSES IN THE CHIRAL SUPERMULTIPLET FORMULATION

In this appendix we compute the gaugino masses using the chiral supermultiplet formulation. By linear-chiral duality, the result must be identical to the one obtained in Sec. III.

Following Appendix A, it is straightforward to write down the chiral supermultiplet formulation of the model defined by Eqs. (2.10)–(2.11). The real superfields $\tilde{V}_{\tilde{A}}$ and V_A dualize to chiral superfields[13] $\tilde{S}_{\tilde{A}}$ and S_A. The linear-chiral duality relations include Eq. (A5) and

$$\tilde{S}_{\tilde{A}} + \bar{\tilde{S}}_{\tilde{A}} = \frac{1+\tilde{f}_{\tilde{A}}(\tilde{V}_{\tilde{A}})}{\tilde{V}_{\tilde{A}}}. \tag{C1}$$

This gives rise to the following identities:

$$\tilde{s}_{\tilde{A}} + \bar{\tilde{s}}_{\tilde{A}} = \frac{1+\tilde{f}_{\tilde{A}}}{\tilde{l}_{\tilde{A}}}, \tag{C2}$$

[13]It can be shown that the linear-chiral duality established in Appendix A also applies to the effective theory description of Veneziano and Yankielowicz [22]. In particular, the linear-chiral duality relations remain the same.

$$\langle \overline{F}_{\tilde{\bar{S}}_{\tilde{A}}} \rangle = -\frac{\langle M \rangle}{3} \left\langle \frac{1+\tilde{f}^{\tilde{A}} - \tilde{l}_{\tilde{A}} \tilde{f}_{\tilde{l}}^{\tilde{A}}}{\tilde{l}_{\tilde{A}}} \right\rangle .$$

The first is the $\theta = \overline{\theta} = 0$ component of Eq. (C1). The second is obtained by acting on both sides of Eq. (C1) with the operator $\overline{\mathcal{D}}^2$, and then taking the vacuum expectation value of the lowest component.

The expression for the gaugino masses is standard,

$$m_{\lambda}^{\tilde{A}} = -\frac{\langle \overline{F}_{\tilde{\bar{S}}_{\tilde{A}}} \rangle}{\langle \tilde{s}_{\tilde{A}} + \tilde{\bar{s}}_{\tilde{A}} \rangle}, \tag{C3}$$

where $\tilde{s}_{\tilde{A}} = \tilde{S}_{\tilde{A}}|_{\theta=\overline{\theta}=0}$ and $F_{\tilde{S}_{\tilde{A}}} = -\frac{1}{4}\mathcal{D}^2 \tilde{S}_{\tilde{A}}|_{\theta=\overline{\theta}=0}$. Using Eq. (C2) and $\langle M \rangle = 3 m_{\tilde{G}}$, we find a final result for the gaugino masses that is identical to Eq. (3.3), obtained from the linear supermultiplet formalism.

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