Anomalous $U(1)$ D-term contribution in type I string models

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We study the D-term contribution for anomalous $U(1)$ symmetries in type I string models and derive a general formula for the D-term contribution, assuming that the dominant source of supersymmetry breaking is given by the F terms of the dilaton, (overall) moduli, or twisted moduli fields. On the basis of the formula, we also point out that there are several features different from the case of heterotic string models. The differences originate from the different forms of the Kähler potential between twisted moduli fields in type I string models and the dilaton field in heterotic string models.

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I. INTRODUCTION

Superstring theory is a promising candidate for a unified theory including gravity. One of the important features is that four-dimensional (4D) string models have several moduli fields including the dilaton field. Their vacuum expectation values (VEVs) determine the couplings of 4D effective theory, e.g., gauge couplings, Yukawa couplings, and Fayet-Iliopoulos (FI) coefficients. These moduli fields have perturbatively a flat potential. Nonperturbative effects are expected to stabilize these moduli. Such nonperturbative effects may also break supersymmetry (SUSY) at the same time. If SUSY is broken, SUSY breaking terms, e.g., gaugino masses and soft scalar masses, are induced. The pattern of SUSY breaking terms depends on couplings of gauge and matter fields to moduli fields. These spectra of superpartners should be measured in the near future. Thus, it is very important to study SUSY breaking terms in 4D string models.

Actually, such analyses have been done extensively both in heterotic models [1,2] and in type I models [3]. For example, the dilaton-dominant SUSY breaking in 4D heterotic models has high predictability, when we consider the scalar potential due only to F terms. That leads to the universal relations $M^{\alpha}_{1/2}=-A_{IJK}=\sqrt{3}m_{3/2}$ and $m_I^2=|m_{3/2}|^2$, where $m_{3/2}$ is the gravitino mass, $M^{\alpha}_{1/2}$ is the gaugino mass, A_{IJK} is the A term, and m_I is the soft scalar mass, while SUSY breaking due to other sources leads to nonuniversal relations. The universal spectrum of sfermion masses is favorable from the viewpoint of flavor changing neutral current (FCNC) constraints. On the other hand, the high predictability may face problems. For example, this pattern of SUSY breaking terms easily leads to color and/or charge breaking (CCB) or the direction unbounded from below (UFB) [4].[1] Similarly, SUSY breaking terms have been studied in type I models when we consider the scalar potential only due to F terms [3,5].

Most 4D string models for both heterotic models and type I models have anomalous $U(1)$ symmetries [6–8]. Many 4D type I models have been built, e.g., through the type IIB orientifold construction. The anomaly is cancelled by the Green-Schwarz mechanism, where certain fields transform nonlinearly. This role is played by the dilaton field in heterotic models and twisted moduli fields in type I models, respectively. Then these fields generate FI terms, whose magnitudes are determined by VEVs of the dilaton field and twisted moduli fields. Other chiral matter fields develop their VEVs along the almost D-flat direction (the D-flat direction in the SUSY limit) and $U(1)$ symmetries are broken. As a phenomenological application of anomalous $U(1)$ symmetry, it can be used as a flavor symmetry for the Froggatt-Nielsen mechanism [9,10]. If one can assign $U(1)$ charges suitably to quarks and leptons, realistic Yukawa matrices can be derived.

In general, there appears an additional contribution to soft SUSY breaking scalar masses called the "D-term contribution" after gauge symmetries are broken down [11–13]. This contribution has a linear dependence on the VEV of the D component and it is proportional to the charge of the broken symmetry. These features are different from those in the contribution from the F component, which has the quadratic form of the VEVs of the F component and it does not depend explicitly on the charge of the broken symmetry. The magnitude of the D-term condensation has been studied in grand unified theories [12,14].

Since most 4D string models have an anomalous $U(1)$ symmetry, its breaking, in general, induces a D-term contribution to the scalar masses. For 4D heterotic models, the D-term contribution has been examined [15–18]. In particular, in Ref. [16] it is taken into account that the FI term is dilaton dependent. As a result, even in dilaton-dominant SUSY breaking the D-term contribution induces nonuniversal scalar masses, and the additional terms are proportional to the $U(1)$ charges. That has phenomenologically important implications. For example, the CCB and UFB constraints can be relaxed [19]. As another aspect, these D-term contribu-

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[1]This problem is not serious, if the age of the Universe is not long enough to reach the CCB minimum.

tions have an important implication for the Froggatt-Nielsen mechanism. In order to derive realistically hierarchical Yukawa matrices, one has to assign different $U(1)$ charges for different families. In this case, the D-term contribution proportional to the $U(1)$ charges leads to nonuniversal sfermion masses, which are dangerous from the viewpoint of FCNC constraints.

Thus, it is important to study the magnitude of the D-term condensation for each model. In this paper, we study the D-term contribution for anomalous $U(1)$ symmetries in 4D type I models and point out that there are several different features from the case of heterotic models. Such difference comes from the fact that in type I models the twisted moduli fields play a role in the Green-Schwarz (GS) anomaly cancellation mechanism, that is, the FI term depends on the twisted moduli fields. Their F components can contribute to SUSY breaking.[2] Their Kähler potential is expected to be different from that of the dilaton field. Furthermore, unlike the dilaton VEV in heterotic models, the VEVs of twisted moduli fields can be taken freely.

This paper is organized as follows. In the next section, we explain the D-term contribution to the soft SUSY breaking scalar masses and the general formula for the VEV of the D auxiliary fields. After reviewing the D-term contribution for the anomalous $U(1)$ symmetry based on heterotic models in Sec. III, we study the D-term contribution for anomalous $U(1)$ symmetries in the framework of type I models in Sec. IV. In Sec. V, we discuss phenomenological implications of the D-term contributions. Section VI is devoted to the conclusions.

II. D-TERM CONTRIBUTION

We explain the D-term contribution to soft SUSY breaking scalar masses based on supergravity (SUGRA) theory [14]. The matter sector in SUGRA is specified by two functions, the total Kähler potential $G(\phi^I,\overline{\phi^I})$ and the gauge kinetic function $f_{\alpha\beta}(\phi^I)$ with α,β being indices of the adjoint representations of the gauge groups. The former is a sum of the Kähler potential $K(\phi^I,\overline{\phi^I})$ and the logarithm of the superpotential $W(\phi^I)$

$$G(\phi^I,\overline{\phi^I})=K(\phi^I,\overline{\phi^I})+M^2\ln\frac{|W(\phi^I)|^2}{M^6}, \tag{1}$$

where M is the gravitational scale defined by use of the Planck mass M_{Pl} such as $M\equiv M_{Pl}/\sqrt{8\pi}$. We have denoted scalar fields in the chiral multiplets by ϕ^I and their complex conjugates by $\overline{\phi^I}$. The real part of the gauge kinetic function $\mathrm{Re}\,f_{\alpha\beta}$ is related to the gauge coupling constants g_α as follows:

$$\langle \mathrm{Re}\,f_{\alpha\beta}\rangle=\frac{1}{g_\alpha^2}\,\delta_{\alpha\beta}. \tag{2}$$

The scalar potential is given by

$$V=M^2e^{G/M^2}[G_I(G^{-1})^{I\bar{J}}G_{\bar{J}}-3M^2] + \frac{1}{2}(\mathrm{Re}\,f^{-1})_{\alpha\beta}G_I(T^\alpha\phi)^IG_J(T^\beta\phi)^J, \tag{3}$$

where $G_I=\partial G/\partial\phi^I$, $G_{\bar{I}}=\partial G/\partial\overline{\phi^I}$, etc., $(\mathrm{Re}\,f^{-1})_{\alpha\beta}$ and $(G^{-1})^{I\bar{J}}$ are the inverse matrices of $\mathrm{Re}\,f_{\alpha\beta}$ and $G_{I\bar{J}}$, respectively, and $(T^\alpha\phi)^I$ are gauge variations up to infinitesimal parameters. The F auxiliary fields and the D auxiliary fields are given by

$$F^I=Me^{G/2M^2}(G^{-1})^{I\bar{J}}G_{\bar{J}},\quad \mathrm{D}^\alpha=(\mathrm{Re}\,f^{-1})_{\alpha\beta}G_I(T^\beta\phi)^I, \tag{4}$$

respectively. In terms of F^I and D^α, the scalar potential takes the form

$$V=V_F+V_D\equiv(\overline{F^J}K_{\bar{J}I}F^I-3M^4e^{G/M^2})+\frac{1}{2}\mathrm{Re}\,f_{\alpha\beta}\mathrm{D}^\alpha\mathrm{D}^\beta. \tag{5}$$

By taking the flat limit of V, we obtain the soft SUSY breaking terms for scalar fields. Here we are interested in the scalar mass terms

$$V_{\mathrm{soft}}=(m_F^2)_{I\bar{J}}\phi^I\overline{\phi^J}+(m_D^2)_{I\bar{J}}\phi^I\overline{\phi^J}+\cdots, \tag{6}$$

$$(m_F^2)_{I\bar{J}}\equiv\left(|m_{3/2}|^2+\frac{\langle V_F\rangle}{M^2}\right)\langle K_{I\bar{J}}\rangle+\langle F^{I'}\rangle\langle\overline{F^{J'}}\rangle \times\langle\partial_{I'}K_{I\bar{J}''}(K^{-1})^{\bar{J}''I''}\partial_{\bar{J}'}K_{I''\bar{J}}-\partial_{I'}\partial_{\bar{J}'}K_{I\bar{J}}\rangle, \tag{7}$$

$$(m_D^2)_{I\bar{J}}\equiv\langle\mathrm{D}^\alpha\rangle\left\langle\frac{\partial}{\partial\phi^I\overline{\phi^J}}[G_{I'}(T^\alpha\phi)^{I'}]\right\rangle, \tag{8}$$

where $m_{3/2}=\langle e^{K/2M^2}W/M^2\rangle$ is the gravitino mass. The magnitude of $m_{3/2}$ is expected to be $\mathcal{O}(1)$ TeV on phenomenological grounds. The first term in Eq. (6) originates from the F-term scalar potential V_F and so we will refer to it as the F-term scalar mass. On the other hand, the second term, Eq. (8), is the D-term contribution to the scalar masses [11,12]. It is proportional to the charge of the broken symmetry and appears when the rank of the gauge group is lowered on the breakdown of the gauge symmetry.

By taking the VEV of $(\partial V/\partial\phi^I)(T^\alpha\phi)^I$ and using the stationary condition, we derive the useful formula for $\langle\mathrm{D}^\alpha\rangle$,

[2]See Refs. [20–22] for the scenario of SUSY breaking by the F term of the twisted moduli fields.

$$\left[(M_V^2)^{\alpha\beta} + \left(\frac{\langle V_F \rangle}{M^2} + 2|m_{3/2}|^2 \right) \langle \mathrm{Re}\, f_{\alpha\beta} \rangle \right] \langle \mathrm{D}^\beta \rangle = \langle F^I \rangle \langle \overline{F^J} \rangle \left\langle \frac{\partial}{\partial \phi^I \partial \overline{\phi^J}} [G_{I'} (T^\alpha \phi)^{I'}] \right\rangle + \frac{1}{2} \left\langle \frac{\partial}{\partial \phi^I} \mathrm{Re}\, f_{\beta\gamma} \right\rangle \langle (T^\alpha \phi)^I \rangle \langle \mathrm{D}^\beta \rangle \langle \mathrm{D}^\gamma \rangle, \tag{9}$$

where $(M_V^2)^{\alpha\beta} = \langle (\overline{\phi} T^\beta)^{\overline{J}} K_{I\overline{J}} (T^\alpha \phi)^I \rangle$ is the mass matrix of the gauge bosons, up to a normalization factor including the gauge coupling constants.

We require that the SUSY is broken down by nonvanishing F-component VEVs of $\mathcal{O}(m_{3/2} M)$ and its effect is mediated through the gravitational interaction. When the extra gauge boson mass is much larger than $m_{3/2}$, the last term is negligibly small compared with other terms in Eq. (9). Then the formula is simplified as

$$\langle \mathrm{D}^\beta \rangle = \langle F^I \rangle \langle \overline{F^J} \rangle \left\langle \frac{\partial}{\partial \phi^I \partial \overline{\phi^J}} [(G_{I'} (T^\alpha \phi))^{I'}] \right\rangle (M_V^{-2})^{\alpha\beta}, \tag{10}$$

where $(M_V^{-2})^{\alpha\beta}$ is the inverse matrix of $(M_V^2)^{\alpha\beta}$. The formula (10) is the master equation in our analysis; when we apply this formula to an anomalous $U(1)$, a certain field transforms nonlinearly and its gauge transformation $(T^\alpha \phi)^I$ becomes a field-independent constant.

For later convenience, we write down the formula for the gaugino masses $M_{1/2}^\alpha$:

$$M_{1/2}^\alpha \delta_{\alpha\beta} = \frac{1}{2 \langle \mathrm{Re}\, f_{\alpha\beta} \rangle} \langle F^I \rangle \left\langle \frac{\partial}{\partial \phi^I} f_{\alpha\beta} \right\rangle. \tag{11}$$

III. ANOMALOUS $U(1)$ D-TERM IN HETEROTIC STRING MODELS

Effective SUGRA is derived from 4D string models by taking the field theory limit [2]. In this section, we review the D-term contribution for the anomalous $U(1)$ symmetry [$U(1)_A$] in 4D heterotic string models [16]. The Kähler potential $K(\phi^I, \overline{\phi^I})$ and the gauge kinetic function $f_{\alpha\beta}(\phi^I)$ are given by

$$K(\phi^I, \overline{\phi^I}) = -\ln(S + \overline{S} - 2\delta_{GS}^A V_A) - \sum_a \ln(T^a + \overline{T^a}) + \sum_\kappa \prod_a (T^a + \overline{T^a})^{n_\kappa^a} \overline{\phi^\kappa} e^{2 q_\kappa^A V_A} \phi^\kappa + \cdots, \tag{12}$$

$$f_{\alpha\beta}(\phi^I) = k_\alpha S \delta_{\alpha\beta} + \varepsilon_\alpha^a T^a \delta_{\alpha\beta}, \tag{13}$$

where S is the dilaton field, T^a are the moduli fields, ϕ^κ are matter fields with modular weights n_κ^a and $U(1)_A$ charges q_κ^A, and V_A is the $U(1)_A$ vector superfield.[3] Also, in the above, k_α is a Kac-Moody level (hereafter we set $k_\alpha = 1$, for simplicity), ε_α^a is a model-dependent parameter coming from the one-loop correction, and δ_{GS}^A is the GS coefficient of $U(1)_A$ given by

$$\delta_{GS}^A = \frac{1}{192\pi^2} \sum_\kappa q_\kappa^A. \tag{14}$$

The $U(1)_A$ D component is given by

$$\mathrm{D}^A = (\mathrm{Re}\, f^{-1})_A \left(\frac{\delta_{GS}^A}{S + \overline{S}} + \sum_\kappa \prod_a (T^a + \overline{T^a})^{n_\kappa^a} q_\kappa^A |\phi^\kappa|^2 \right), \tag{15}$$

where we neglect terms from higher order terms in $K(\phi^I, \overline{\phi^I})$. Following the custom in 4D SUGRA derived from string models, we take the $M = 1$ unit if no confusion is expected.

The $U(1)_A$ and its mixed anomalies due to matter fields are cancelled by the contribution from the dilaton field which transforms nonlinearly as $S \to S' = S + i \delta_{GS}^A \theta(x)$ under $U(1)_A$. Then the formula (10) for $\langle \mathrm{D}^A \rangle$ reads

$$\langle \mathrm{D}^A \rangle = \frac{1}{(M_V^2)^A} \left(2\delta_{GS}^A \frac{|\langle F^S \rangle|^2}{\langle S + \overline{S} \rangle^3} + \left\langle \sum_\kappa \prod_a (T^a + \overline{T^a})^{n_\kappa^a} q_\kappa^A |F^\kappa|^2 \right\rangle + \cdots \right). \tag{16}$$

Here $(M_V^2)^A$ is given by

$$(M_V^2)^A = \frac{(\delta_{GS}^A)^2}{\langle S + \overline{S} \rangle^2} + \left\langle \sum_\kappa \prod_a (T^a + \overline{T^a})^{n_\kappa^a} (q_\kappa^A)^2 |\phi^\kappa|^2 \right\rangle, \tag{17}$$

which can be rewritten with the help of the almost D-flatness condition of $U(1)_A$ into

$$(M_V^2)^A = \frac{\delta_{GS}^A}{\langle S + \overline{S} \rangle} \left(\frac{\delta_{GS}^A}{\langle S + \overline{S} \rangle} - \frac{\left\langle \sum_\kappa \prod_a (T^a + \overline{T^a})^{n_\kappa^a} (q_\kappa^A)^2 |\phi^\kappa|^2 \right\rangle}{\left\langle \sum_\kappa \prod_a (T^a + \overline{T^a})^{n_\kappa^a} q_\kappa^A |\phi^\kappa|^2 \right\rangle} \right). \tag{18}$$

[3]In Eqs. (12) and (13), all fields stand for superfields with the same notation for chiral superfields and antichiral superfields as for the scalar components.

In explicit models, we find that $\delta_{GS}^{A}=\mathcal{O}(10^{-1})$–$\mathcal{O}(10^{-2})$. Hence we will neglect terms with a higher order of δ_{GS}^{A}. With this assumption, the second term is dominant in Eq. (17).

For simplicity, we treat the case with the overall moduli, i.e., $T=T^1=T^2=T^3$. In this case, $\langle \mathrm{D}^A\rangle$ is given by

$$\langle \mathrm{D}^A\rangle=\frac{1}{(M_V^2)^A}\Bigg(2\,\delta_{GS}^{A}\frac{|\langle F^S\rangle|^2}{\langle S+\bar{S}\rangle^3}+\Bigg\langle\sum_\kappa (T+\bar{T})^{n_\kappa}q_\kappa^A|F^\kappa|^2\Bigg\rangle+\frac{|\langle F^T\rangle|^2}{\langle T+\bar{T}\rangle^2}\Bigg\langle\sum_\kappa (T+\bar{T})^{n_\kappa}n_\kappa(n_\kappa-1)q_\kappa^A|\phi^\kappa|^2\Bigg\rangle+\frac{\langle \bar{F}^{\bar{T}}\rangle}{\langle T+\bar{T}\rangle}\Bigg\langle\sum_\kappa (T+\bar{T})^{n_\kappa}n_\kappa q_\kappa^A F^\kappa\bar{\phi}^{\bar{\kappa}}\Bigg\rangle+\mathrm{H.c.}\Bigg). \tag{19}$$

Here we consider the case that the dilaton and the overall moduli fields are dominant sources for the SUSY breaking, e.g., $\langle F^S\rangle,\langle F^T\rangle=\mathcal{O}(m_{3/2}M)$. This situation is realized if $\langle\phi^\kappa\rangle\ll\mathcal{O}(M)$ and $\langle\partial W/\partial\phi^\kappa\rangle\ll\mathcal{O}(m_{3/2}M)$. In this case, from the expression (4) for $\langle F^I\rangle$, we find that the F terms $\langle F^S\rangle$ of the chiral matter fields are induced as

$$\langle F^\kappa\rangle=\left(m_{3/2}-n_\kappa\frac{\langle F^T\rangle}{\langle T+\bar{T}\rangle}\right)\langle\phi^\kappa\rangle. \tag{20}$$

Since the induced $\langle F^\kappa\rangle$ is much smaller than $\langle F^S\rangle$ and $\langle F^T\rangle$, the VEV $\langle V_F\rangle$ is simplified to

$$\langle V_F\rangle=\frac{|\langle F^S\rangle|^2}{\langle S+\bar{S}\rangle^2}+3\frac{|\langle F^T\rangle|^2}{\langle T+\bar{T}\rangle^2}-3|m_{3/2}|^2. \tag{21}$$

In Eq. (19), however, the terms including $\langle F^\kappa\rangle$ are comparable with the other terms, and we obtain

$$\langle \mathrm{D}^A\rangle=\frac{1}{(M_V^2)^A}\Bigg[\frac{\delta_{GS}^{A}}{S+\bar{S}}\left(2\frac{|\langle F^S\rangle|^2}{\langle S+\bar{S}\rangle^2}-|m_{3/2}|^2\right)-\frac{|\langle F^T\rangle|^2}{\langle T+\bar{T}\rangle^2}\Bigg\langle\sum_\kappa (T+\bar{T})^{n_\kappa}n_\kappa q_\kappa^A|\phi^\kappa|^2\Bigg\rangle\Bigg]. \tag{22}$$

Note that the D-term condensation depends on $\langle F^T\rangle$ as well as $\langle F^S\rangle$.

The gaugino masses $M_{1/2}^\alpha$ are calculated by use of Eq. (11) to be

$$M_{1/2}^\alpha=\frac{1}{2\langle\mathrm{Re}\,f_\alpha\rangle}(\langle F^S\rangle+\varepsilon_\alpha\langle F^T\rangle). \tag{23}$$

To obtain gaugino masses of $\mathcal{O}(m_{3/2})$, we need a dilaton-dominant SUSY breaking scenario in the weakly coupled region. In the strongly coupled region, the moduli F component can also lead to gaugino masses of $\mathcal{O}(m_{3/2})$ [23]. In any case, a $U(1)_A$ D-term contribution to the scalar masses appears, $(m_D^2)_I=q_I^A\langle \mathrm{D}^A\rangle$, and its magnitude is rather large as $\langle \mathrm{D}^A\rangle=\mathcal{O}(m_{3/2}^2)$. If $U(1)_A$ charges are different between the first and second families, the nonuniversality among sfermion masses would be dangerous from the viewpoint of FCNC constraints. On the other hand, with a $U(1)_A$ D-term contribution we can relax the CCB and UFB bounds [19].

In a particular case [24], the F components of matter fields can also contribute to the breakdown of SUSY.[4] For instance, $\langle F^\kappa\rangle$ contributes to the SUSY breaking if $\langle\partial W/\partial\phi^\kappa\rangle=\mathcal{O}(m_{3/2}M)$. Then the dominant part of the D^A condensation comes from the second term in the right-hand side (RHS) of Eq. (19),

$$\langle \mathrm{D}^A\rangle=\frac{\Big\langle\sum_\kappa (T+\bar{T})^{n_\kappa}q_\kappa^A|F^\kappa|^2\Big\rangle}{\Big\langle\sum_\kappa (T+\bar{T})^{n_\kappa}(q_\kappa^A)^2|\phi^\kappa|^2\Big\rangle}. \tag{24}$$

The magnitude is estimated [18] as $\langle \mathrm{D}^A\rangle=\mathcal{O}(m_{3/2}^2/\delta_{GS}^{A})$.

IV. ANOMALOUS $U(1)$ D-TERMS IN TYPE I MODELS

Next we turn to the type I case. In general, a 4D type I model has more than one anomalous $U(1)$ symmetry, i.e., $\Pi_i U(1)_i$. We denote the $U(1)_i$ vector multiplet by V_i. The Kähler potential $K(\phi^I,\bar{\phi}^{\bar{I}})$ is given by[5]

$$K(\phi^I,\bar{\phi}^{\bar{I}})=\hat{K}\left(M_\ell+\bar{M}_\ell-2\sum_i(\delta_{GS})_i^\ell V_i\right)-\ln(S+\bar{S})-\sum_a\ln(T^a+\bar{T}^a)+\sum_\kappa\prod_a(S+\bar{S})^{n_\kappa^s}\times(T^a+\bar{T}^a)^{n_\kappa^a}\bar{\phi}^{\bar{\kappa}}e^{2q_\kappa^i V_i}\phi^\kappa+\cdots, \tag{25}$$

where the chiral matter fields ϕ^κ have the "modular weights" n_κ^s and n_κ^a with respect to S and T^a, and $(\delta_{GS})_i^\ell$ are model-dependent GS coefficients. Here, M_ℓ is a twisted moduli field associated with the ℓth fixed point. For simplicity, we use the notation m_ℓ defined by $m_\ell\equiv M_\ell+\bar{M}_\ell-2\Sigma_i(\delta_{GS})_i^\ell V_i$ hereafter. The complete form of $\hat{K}$ is unknown, but in the orbifold limit $M_\ell\to 0$, it takes the tree level form [26]

$$\hat{K}(m_\ell)=\frac{1}{2}m_\ell^2. \tag{26}$$

The M_ℓ dependence of the Kähler metric of ϕ^κ is also unclear. In the orbifold limit, the Kähler metric $K_{\kappa\bar{\kappa}}$ does not

[4]See also Ref. [25].

[5]For models with effective low-energy Lagrangians of type I, see Ref. [3] and references therein.

depend on the twisted moduli M_ℓ as in Eq. (25). For a large value of M_ℓ, however, it would receive a correction $\Delta K_{\kappa\bar{\kappa}}(M,\bar{M})$.

The gauge kinetic function $f_{\alpha\beta}(\phi^I)$ is given by

$$f_{\alpha\beta}(\phi^I)=\hat{f}(S,T^a)\,\delta_{\alpha\beta}+\sum_\ell s_\ell^\alpha M_\ell \delta_{\alpha\beta}, \tag{27}$$

where s_ℓ^α is a model-dependent constant. The first term is D-brane dependent, e.g., $\hat{f}(S,T^a)=S$ for gauge groups from D9-branes and $\hat{f}(S,T^a)=T^a$ for gauge groups from D5$_a$-branes. The $U(1)_i$ and mixed anomalies due to matter fields are cancelled by the contribution from the twisted moduli fields, which transform as $M_\ell\to M'_\ell=M_\ell+i(\delta_{GS})_i^\ell\,\theta(x)$ under $U(1)_i$.

The $U(1)_i$ D components are given by

$$\mathrm{D}^i=(\mathrm{Re}\,f^{-1})_i\left(-(\delta_{GS})_i^\ell\frac{\partial\hat{K}}{\partial m_\ell}+\sum_\kappa\prod_a (S+\bar{S})^{n_\kappa^s}(T^a+\bar{T}^a)^{n_\kappa^a}q_\kappa^i|\phi^\kappa|^2\right), \tag{28}$$

where we assume that $U(1)$ kinetic mixing is absent for simplicity. According to the formula (10) for D-term condensation, we obtain

$$\langle\mathrm{D}^i\rangle=\frac{1}{(M_V^2)^i}\left(-(\delta_{GS})_i^\ell\left\langle\frac{\partial^3\hat{K}}{\partial m_\ell\partial m_{\ell'}\partial m_{\ell''}}F^{M_{\ell'}}\bar{F}^{\bar{M}_{\bar{\ell}''}}\right\rangle+\left\langle\sum_\kappa\prod_a (S+\bar{S})^{n_\kappa^s}(T^a+\bar{T}^a)^{n_\kappa^a}q_\kappa^i|F^\kappa|^2\right\rangle+\cdots\right). \tag{29}$$

Here $(M_V^2)^i$ is (essentially) the $U(1)_i$ gauge boson mass given by

$$\begin{aligned}(M_V^2)^i&\equiv(\delta_{GS})_i^\ell(\delta_{GS})_i^{\ell'}\left\langle\frac{\partial^2\hat{K}}{\partial m_\ell\partial m_{\ell'}}\right\rangle+\left\langle\sum_\kappa\prod_a (S+\bar{S})^{n_\kappa^s}(T^a+\bar{T}^a)^{n_\kappa^a}(q_\kappa^i)^2|\phi^\kappa|^2\right\rangle\\&=(\delta_{GS})_i^\ell(\delta_{GS})_i^{\ell'}\left\langle\frac{\partial^2\hat{K}}{\partial m_\ell\partial m_{\ell'}}\right\rangle+(\delta_{GS})_i^\ell\left\langle\frac{\partial\hat{K}}{\partial m_\ell}\right\rangle\times\frac{\left\langle\sum_\kappa\prod_a (S+\bar{S})^{n_\kappa^s}(T^a+\bar{T}^a)^{n_\kappa^a}(q_\kappa^i)^2|\phi^\kappa|^2\right\rangle}{\left\langle\sum_\kappa\prod_a (S+\bar{S})^{n_\kappa^s}(T^a+\bar{T}^a)^{n_\kappa^a}q_\kappa^i|\phi^\kappa|^2\right\rangle},\end{aligned} \tag{30}$$

where we have used the almost D-flatness conditions of $U(1)_i$. In the case with the canonical Kähler potential (26), $(M_V^2)^i$ is reduced to

$$(M_V^2)^i\equiv[(\delta_{GS})_i^\ell]^2+(\delta_{GS})_i^\ell\langle m_\ell\rangle\times\frac{\left\langle\sum_\kappa\prod_a (S+\bar{S})^{n_\kappa^s}(T^a+\bar{T}^a)^{n_\kappa^a}(q_\kappa^i)^2|\phi^\kappa|^2\right\rangle}{\left\langle\sum_\kappa\prod_a (S+\bar{S})^{n_\kappa^s}(T^a+\bar{T}^a)^{n_\kappa^a}q_\kappa^i|\phi^\kappa|^2\right\rangle}. \tag{31}$$

If $\langle m_\ell\rangle\ll\mathcal{O}(\delta_{GS})$, the first term is dominant in Eq. (31), unlike the heterotic case (17).

Again we treat the case with the overall moduli, i.e., $T=T^1=T^2=T^3$, and denote $n_\kappa=\Sigma_a n_\kappa^a$. Then the $\langle\mathrm{D}^i\rangle$ in Eq. (29) can be written down explicitly as

$$\begin{aligned}\langle\mathrm{D}^i\rangle=\frac{1}{(M_V^2)^i}&\left(-(\delta_{GS})_i^\ell\left\langle\frac{\partial^3\hat{K}}{\partial m_\ell\partial m_{\ell'}\partial m_{\ell''}}F^{M_{\ell'}}\bar{F}^{\bar{M}_{\bar{\ell}''}}\right\rangle+\left\langle\sum_\kappa (S+\bar{S})^{n_\kappa^s}(T+\bar{T})^{n_\kappa}q_\kappa^i|F^\kappa|^2\right\rangle\right.\\&+\frac{|\langle F^S\rangle|^2}{\langle S+\bar{S}\rangle^2}\left\langle\sum_\kappa (S+\bar{S})^{n_\kappa^s}(T+\bar{T})^{n_\kappa}n_\kappa^s(n_\kappa^s-1)q_\kappa^i|\phi^\kappa|^2\right\rangle+\frac{|\langle F^T\rangle|^2}{\langle T+\bar{T}\rangle^2}\left\langle\sum_\kappa (S+\bar{S})^{n_\kappa^s}(T+\bar{T})^{n_\kappa}n_\kappa(n_\kappa-1)q_\kappa^i|\phi^\kappa|^2\right\rangle\\&+\frac{|\langle F^S\rangle|}{\langle S+\bar{S}\rangle}\frac{|\langle\bar{F}^{\bar{T}}\rangle|}{\langle T+\bar{T}\rangle}\left\langle\sum_\kappa (S+\bar{S})^{n_\kappa^s}(T+\bar{T})^{n_\kappa}n_\kappa^s n_\kappa q_\kappa^i|\phi^\kappa|^2\right\rangle+\mathrm{H.c.}+\frac{\langle\bar{F}^{\bar{S}}\rangle}{\langle S+\bar{S}\rangle}\left\langle\sum_\kappa (S+\bar{S})^{n_\kappa^s}(T+\bar{T})^{n_\kappa}n_\kappa^s q_\kappa^i F^\kappa\bar{\phi}^{\bar{\kappa}}\right\rangle\\&\left.+\mathrm{H.c.}+\frac{\langle\bar{F}^{\bar{T}}\rangle}{\langle T+\bar{T}\rangle}\left\langle\sum_\kappa (S+\bar{S})^{n_\kappa^s}(T+\bar{T})^{n_\kappa}n_\kappa q_\kappa^i F^\kappa\bar{\phi}^{\bar{\kappa}}\right\rangle+\mathrm{H.c.}\right).\end{aligned} \tag{32}$$

In the following, we mainly consider the case that SUSY is broken by the dilaton, the overall moduli, and/or the twisted moduli fields $\langle F^S\rangle$, $\langle F^T\rangle$, $\langle F^{M_\ell}\rangle=\mathcal{O}(m_{3/2}M)$. This situation is realized if $\langle\phi^\kappa\rangle\ll\mathcal{O}(M)$ and $\langle\partial W/\partial\phi^\kappa\rangle\ll\mathcal{O}(m_{3/2}M)$. In this case, since $\langle F^\kappa\rangle=\mathcal{O}(m_{3/2}\langle\phi^\kappa\rangle)$, the VEV $\langle V_F\rangle$ is simplified as

$$\langle V_F\rangle=\frac{|\langle F^S\rangle|^2}{\langle S+\overline{S}\rangle^2}+3\frac{|\langle F^T\rangle|^2}{\langle T+\overline{T}\rangle^2}+\sum_{\ell,\ell'}\left\langle\frac{\partial^2\hat{K}}{\partial m_\ell\partial m_{\ell'}}F^{M_\ell}\overline{F}^{\overline{M}_{\ell'}}\right\rangle-3|m_{3/2}|^2. \quad (33)$$

To calculate the $\langle \mathrm{D}^i\rangle$, however, it is important to note that $\langle F^\kappa\rangle$ is induced as

$$\langle F^\kappa\rangle=\left(m_{3/2}-n_\kappa^s\frac{\langle F^S\rangle}{\langle S+\overline{S}\rangle}-n_\kappa\frac{\langle F^T\rangle}{\langle T+\overline{T}\rangle}\right)\langle\phi^\kappa\rangle. \quad (34)$$

Then a careful calculation leads to

$$\begin{aligned}\langle \mathrm{D}^i\rangle=\frac{1}{(M_V^2)^i}\Bigg(&-(\delta_{GS})_i^\ell\left\langle\frac{\partial^3\hat{K}}{\partial m_\ell\partial m_{\ell'}\partial m_{\ell''}}F^{M_{\ell'}}\overline{F}^{\overline{M}_{\ell''}}\right\rangle\\&+|m_{3/2}|^2\left\langle\sum_\kappa (S+\overline{S})^{n_\kappa^s}(T+\overline{T})^{n_\kappa}q_\kappa^i|\phi^\kappa|^2\right\rangle\\&-\frac{|\langle F^S\rangle|^2}{\langle S+\overline{S}\rangle^2}\left\langle\sum_\kappa (S+\overline{S})^{n_\kappa^s}(T+\overline{T})^{n_\kappa}n_\kappa^s q_\kappa^i|\phi^\kappa|^2\right\rangle\\&-\frac{|\langle F^T\rangle|^2}{\langle T+\overline{T}\rangle^2}\left\langle\sum_\kappa (S+\overline{S})^{n_\kappa^s}(T+\overline{T})^{n_\kappa}n_\kappa q_\kappa^i|\phi^\kappa|^2\right\rangle\Bigg).\end{aligned} \quad (35)$$

The expressions (32) and (35) are our main results for the D-term condensation in type I models. Note that $\langle \mathrm{D}^i\rangle$ becomes independent of $\langle F^{M_\ell}\rangle$ if the third derivative of $\hat{K}$ vanishes: $\langle\partial^3\hat{K}/\partial m_\ell\partial m_{\ell'}\partial m_{\ell''}\rangle\ll\mathcal{O}((\delta_{GS})_i^\ell)$.

The soft terms can be calculated by using the parametrization

$$\frac{\langle F^S\rangle}{\langle S+\overline{S}\rangle}\equiv\sqrt{3}C|m_{3/2}|e^{i\alpha_S}\sin\theta,$$

$$\frac{\langle F^T\rangle}{\langle T+\overline{T}\rangle}\equiv C|m_{3/2}|e^{i\alpha_T}\cos\theta\sin\phi,$$

$$\langle F^{M_\ell}\rangle\equiv\sqrt{3}C|m_{3/2}|e^{i\alpha_\ell}\Phi_\ell\cos\theta\cos\phi,$$

$$\sum_\ell \Phi_\ell^2=1, \quad (36)$$

where C is a constant so that $\langle V_F\rangle=3(C^2-1)|m_{3/2}|^2$, and θ, ϕ, and Φ_ℓ are parameters called the "goldstino angles." We have assumed for simplicity that the Kähler metric of the twisted moduli is diagonal. Then the $\langle \mathrm{D}^i\rangle$ becomes

$$\begin{aligned}\langle \mathrm{D}^i\rangle=\frac{|m_{3/2}|^2}{(M_V^2)^i}\Bigg(&-3C^2(\delta_{GS})_i^\ell\left\langle\frac{\partial^3\hat{K}}{\partial m_\ell\partial m_{\ell'}\partial m_{\ell''}}\right\rangle\\&\times\Phi_{\ell'}\Phi_{\ell''}\cos^2\theta\cos^2\phi+(\delta_{GS})_i^\ell\left\langle\frac{\partial\hat{K}}{\partial m_\ell}\right\rangle\\&-3C^2\sin^2\theta\left\langle\sum_\kappa (S+\overline{S})^{n_\kappa^s}(T+\overline{T})^{n_\kappa}n_\kappa^s q_\kappa^i|\phi^\kappa|^2\right\rangle\\&-C^2\cos^2\theta\sin^2\phi\\&\times\left\langle\sum_\kappa (S+\overline{S})^{n_\kappa^s}(T+\overline{T})^{n_\kappa}n_\kappa q_\kappa^i|\phi^\kappa|^2\right\rangle\Bigg).\end{aligned} \quad (37)$$

By using this expression, we shall discuss the D-term contribution to the soft scalar masses, $(m_D^2)_I=q_I^i\langle \mathrm{D}^i\rangle$, in the next section.

The gaugino masses $M_{1/2}^\alpha$ are calculated by use of Eqs. (11) and (27),

$$M_{1/2}^\alpha=\frac{1}{2\langle\mathrm{Re}\,f_\alpha\rangle}\left(\langle F^S\rangle+\sum_\ell s_\ell^\alpha\langle F^{M_\ell}\rangle\right)\quad\text{for D9-branes,}$$

$$M_{1/2}^\alpha=\frac{1}{2\langle\mathrm{Re}\,f_\alpha\rangle}\left(\langle F^T\rangle+\sum_\ell s_\ell^\alpha\langle F^{M_\ell}\rangle\right)\quad\text{for D5}_a\text{-branes.} \quad (38)$$

To obtain sizable gaugino masses of $\mathcal{O}(m_{3/2})$, we need the dilaton and/or twisted moduli dominant SUSY breaking scenario on D9-branes, and the overall moduli and/or twisted moduli dominant SUSY breaking scenario on D5$_a$-branes. If the dilaton and/or an overall moduli dominant SUSY breaking occur, the magnitude of the $U(1)_i$ D-term can be sizable as $\langle \mathrm{D}^i\rangle=\mathcal{O}(m_{3/2}^2)$. On the other hand, the magnitude of $\langle \mathrm{D}^i\rangle$ can be small if the twisted moduli fields dominate the SUSY breaking; in this case, $\langle \mathrm{D}^i\rangle$ is negligibly small if $\langle\partial^3\hat{K}/\partial m_\ell\partial m_{\ell'}\partial m_{\ell''}\rangle\ll\mathcal{O}((\delta_{GS})_i^\ell)$ and $\langle m_\ell\rangle\ll\mathcal{O}((\delta_{GS})_i^\ell)$.

Up to here, we have assumed that SUSY is broken by the F terms of S, T, or M_ℓ. Alternatively we can suppose, as in the heterotic case, that there exists a dynamical superpotential W of the chiral matter fields ϕ^κ so that $\langle\partial W/\partial\phi^\kappa\rangle=\mathcal{O}(m_{3/2}M)$. In this case, the dominant part of the D^i condensation comes from the second term in the RHS of Eq. (32),

$$\langle \mathrm{D}^i\rangle=\frac{\left\langle\sum_\kappa (S+\overline{S})^{n_\kappa^s}(T+\overline{T})^{n_\kappa}q_\kappa^i|F^\kappa|^2\right\rangle}{(M_V^2)^i}. \quad (39)$$

The magnitude is estimated as $\langle \mathrm{D}^i\rangle=\mathcal{O}(m_{3/2}^2M^2/(M_V^2)^i)$.

V. PHENOMENOLOGICAL IMPLICATIONS

In this section, we discuss the phenomenological implications of the D-term contributions. An important point is that the FI terms depend on the twisted moduli fields in type I models, while this role is played by the dilaton field in heterotic models. Here let us compare our result (35) in type I models with the D term (22) in heterotic models.

The first term of the D-term condensation (35) is negligibly small, when the canonical term is dominant in $\hat{K}(M_\ell,\bar{M}_\ell)$. As a result, the D-term condensation does not depend on F^{M_ℓ} explicitly. Recall that the D-term condensation (22) in heterotic models depends explicitly on F^S. The difference originates from the different forms of Kähler potential between M_ℓ and S. If $\hat{K}(M_\ell,\bar{M}_\ell)$ has a logarithmic form like S, this difference will disappear.

For the remaining terms in Eq. (35), we can estimate the order of magnitudes by using the fact that the D term (28) almost vanishes. The second term is proportional to the FI terms as in the heterotic case. The last two terms can be estimated as

$$\left\langle \sum_\kappa (S+\bar{S})^{n_\kappa^s}(T+\bar{T})^{n_\kappa} n_\kappa^s q_\kappa^i |\phi^\kappa|^2 \right\rangle = \mathcal{O}\left((\delta_{GS})_i^\ell \left\langle \frac{\partial \hat{K}}{\partial m_\ell} \right\rangle \right), \tag{40}$$

$$\left\langle \sum_\kappa (S+\bar{S})^{n_\kappa^s}(T+\bar{T})^{n_\kappa} n_\kappa q_\kappa^i |\phi^\kappa|^2 \right\rangle = \mathcal{O}\left((\delta_{GS})_i^\ell \left\langle \frac{\partial \hat{K}}{\partial m_\ell} \right\rangle \right). \tag{41}$$

Therefore, we have

$$\langle \mathrm{D}^i \rangle = m_{3/2}^2 \times \mathcal{O}\left(\frac{1}{(M_V^2)^i} (\delta_{GS})_i^\ell \left\langle \frac{\partial \hat{K}}{\partial m_\ell} \right\rangle \right), \tag{42}$$

and its magnitude depends on $\langle m_\ell \rangle / (\delta_{GS})_i^\ell$ as is seen from Eq. (31). Notice that, unlike the dilaton VEV in the heterotic case, the VEVs of twisted moduli fields m_ℓ can be taken as arbitrary values, depending on the stabilization mechanism [20–22].

If $\mathcal{O}(\langle m_\ell \rangle / (\delta_{GS})_i^\ell) \geq 1$, D-term condensations are sizable and their order is $\mathcal{O}(m_{3/2}^2)$. They significantly affect spectra of superpartners. This situation is the same as that in the heterotic case. For example, the CCB and UFB directions have been studied for type I models in Ref. [27] with the string scale varied. Hence, the D-term contributions have important effects as in the heterotic case if their magnitudes are $\mathcal{O}(m_{3/2}^2)$.

On the other hand, it is possible that the VEVs of twisted moduli fields m_ℓ are suppressed, i.e., $O(\langle m_\ell \rangle / (\delta_{GS})_i^\ell) \ll 1$. In this case, the D-term contribution can be suppressed. This is in sharp contrast to the heterotic case where the D-term contribution cannot be suppressed without fine-tuning.

To be concrete, let us first discuss the dilaton-dominant SUSY breaking with $\langle V_F \rangle = 0$ in type I models. For comparison with the heterotic models, we consider the case that the gauge multiplets originate from D9-branes and chiral matter fields originate from open strings, one end of which is on the D9-branes. In this case, the gaugino mass is obtained as

$$M_{1/2} = \sqrt{3} m_{3/2}, \tag{43}$$

where we have taken $|s_\ell^\alpha \langle m_\ell \rangle| \ll \mathrm{Re}\, S$. Since $n_\kappa^s = 0$, the F-term scalar masses are universal, i.e.,

$$(m_F^2)_I = |m_{3/2}|^2. \tag{44}$$

This spectrum is the same as for the dilaton-dominant SUSY breaking in heterotic models. In addition, we have to add the D-term contribution $(m_D^2)_I = q_I^i \langle \mathrm{D}^i \rangle$ with $\langle \mathrm{D}^i \rangle$ given by Eq. (37). When $|\langle m_\ell \rangle / (\delta_{GS})_i^\ell| \ll 1$, the D-term contribution becomes simplified to

$$(m_D^2)_I = q_I^i |m_{3/2}|^2 \frac{\langle m_\ell \rangle}{(\delta_{GS})_i^\ell}. \tag{45}$$

Thus, if $|\langle m_\ell \rangle / (\delta_{GS})_i^\ell| \ll 1$, the D-term contribution is small, and the total soft scalar masses become almost universal. This has important implications for FCNC constraints as well as CCB and UFB bounds. For instance, if this $U(1)$ symmetry is relevant to the flavor symmetry,[6] the suppressed D-term contribution is favorable to avoid FCNC constraints. For this purpose, we need only a suppression like $|\langle m_\ell \rangle / (\delta_{GS})_i^\ell| \leq \mathcal{O}(10^{-2})$ for $m_{3/2} = \mathcal{O}(100)$ GeV. Thus the FCNC can be parametrically suppressed in type I models. (Of course, it is necessary to find a proper mechanism for stabilizing the twisted moduli VEVs, but that is beyond the scope of the present paper.)

Next let us consider the case that the single twisted moduli field M_ℓ is a dominant source of the SUSY breaking. In this case, the gaugino mass is written as

$$M_{1/2}^\alpha = \frac{\sqrt{3}}{2} s_\ell^\alpha g_\alpha^2 m_{3/2}. \tag{46}$$

It is interesting to note that if s_ℓ^α are proportional to the coefficients of one-loop beta functions of the gauge couplings, as in the case of "mirage gauge coupling unification" in Ref. [29], this spectrum of gaugino masses resembles that in the anomaly mediation scenario [30]. Since the Kähler potential of the matter fields does not depend on M_ℓ for small $\langle M_\ell \rangle$, the F-term scalar masses are universal, i.e.,

$$(m_F^2)_I = |m_{3/2}|^2. \tag{47}$$

When $|\langle m_\ell \rangle / (\delta_{GS})_i^\ell| \ll 1$, the D-term contribution becomes

$$(m_D^2)_I = q_I^i |m_{3/2}|^2 \left(-\frac{3C^2}{(\delta_{GS})_i^\ell} \times \left\langle \frac{\partial^3 \hat{K}}{\partial m_\ell \partial m_{\ell'} \partial m_{\ell''}} \right\rangle \Phi_{\ell'} \Phi_{\ell''} + \frac{\langle m_\ell \rangle}{(\delta_{GS})_i^\ell} \right). \tag{48}$$

[6] See, e.g., Ref. [28], where this flavor $U(1)$ symmetry is discussed in a type I inspired model.

Thus, if $\langle \partial^3 \hat{K}/\partial m_\ell \partial m_{\ell'} \partial m_{\ell''} \rangle \ll \mathcal{O}((\delta_{GS})_i^\ell)$, the D-term contribution is small and the total soft scalar masses become almost universal.

Finally, we note that even if flavor-dependent D-term contributions can be suppressed, it is important to take account of radiative corrections due to the gaugino of flavor $U(1)_A$ symmetry [31]. To focus on it, let us assume that at the cutoff scale Λ (of order of the Planck scale), the soft terms are given by the universal values (43), (44), and that the D-term contributions are absent. According to Eq. (45), the latter requirement is satisfied if $\langle m_\ell \rangle = 0$ in the type I models we are considering. Recall that the absence of D-term contributions generally requires fine-tuning or complicated model building in the heterotic models.

Let M_A be the breaking scale of an anomalous $U(1)_A$ flavor symmetry. Above this scale, the $U(1)_A$ gaugino gives a radiative correction $\delta_A m_I^2$ to the sfermion mass squared. The radiative correction is proportional to the square of the soft mass parameter of the $U(1)_A$ gaugino; explicitly it is found to be

$$\delta_A m_I^2 = q_I^2 \frac{2}{b_A}\left(1 - \frac{\alpha_A^2(M_A)}{\alpha_A^2(\Lambda)}\right) M_{1/2}^2(\Lambda) \equiv q_I^2 \Delta_A, \quad (49)$$

where b_A is the one-loop beta function coefficient of the anomalous $U(1)_A$.

In addition, we have flavor-blind radiative corrections due to the remaining gauginos. For example, the radiative corrections from the $SU(3)\times SU(2)\times U(1)$ gauginos between the gauge unification scale M_X and the weak scale are evaluated as

$$\delta m_{\tilde{d}}^2 = 6.6 \times M_{1/2}^2(M_X), \quad (50)$$

$$\delta m_{\tilde{\ell}}^2 = 0.51 \times M_{1/2}^2(M_X), \quad (51)$$

$$\delta m_{\tilde{e}}^2 = 0.15 \times M_{1/2}^2(M_X) \quad (52)$$

for the right-handed squark, left-handed slepton, and right-handed slepton masses, respectively. Typically, these corrections dominate, for each type of sfermion $\tilde{f}$, over those from the $U(1)_A$ gaugino. Accordingly, we introduce the ratio

$$d_{\tilde{f}} \equiv \frac{\Delta_A}{m_0^2 + \delta m_{\tilde{f}}^2}, \quad (53)$$

where $m_0 = m_{3/2}$ is the initial value of the sfermion mass. Nondegeneracy of the sfermion mass squared may be estimated by this factor times the difference of the $U(1)$ charges squared q_I^2.

Figure 1 shows the degeneracy factors $d_{\tilde{d},\tilde{\ell},\tilde{e}}$ against b_X for $\alpha_A(\Lambda)\ln(\Lambda/M_X) = 0.06$, whereas Fig. 2 shows $d_{\tilde{d},\tilde{\ell},\tilde{e}}$ against $a = \alpha_A(\Lambda)\ln(\Lambda/M_X)$ in the case of $b_A = 50$. In each figure, the lower, middle, and upper curves correspond to $d_{\tilde{d}}$, $d_{\tilde{\ell}}$, and $d_{\tilde{e}}$, respectively. For the flavor-blind radiative corrections $\delta m_{\tilde{f}}^2$, we have kept only those from the $SU(3)\times SU(2)\times U(1)$ gauginos. For the initial condition on the gaugino masses, we used $M_{1/2}(M_X) \approx M_{1/2}(\Lambda) = \sqrt{3} m_{3/2}$. We see that the $U(1)$ gaugino of a gauged flavor symmetry can affect sfermion mass degeneracy and that this effect is important especially for sleptons. These radiative effects may be detected if the D-term contributions are suppressed.

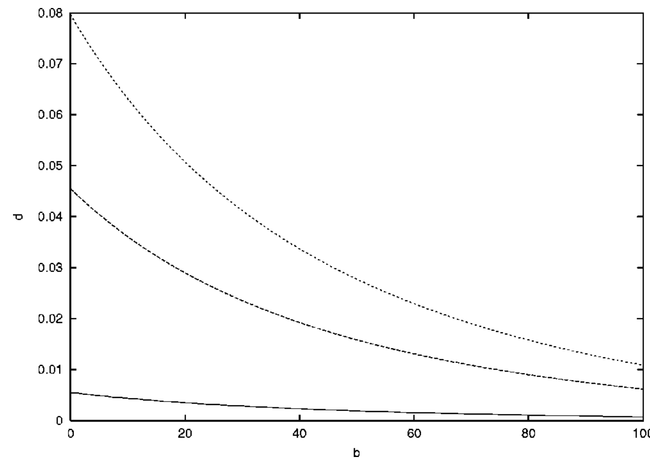

FIG. 1. The degeneracy of the sfermion masses against b_A for $\alpha_A(\Lambda)\ln(\Lambda/M_X) = 0.06$.

VI. CONCLUSION

We have studied the D-term contribution for anomalous $U(1)$ symmetries in type I models. Specifically, we derived a general formula for the D-term contribution, assuming that the dominant source of SUSY breaking is given by the F terms of the dilaton, (overall) moduli, or twisted moduli fields.

We also observed that there are several differences in the D-term contributions between the heterotic and type I models. One of the important differences is that the D-term contribution in type I models depends on the VEVs of twisted moduli fields, while that in heterotic models depends on the dilaton VEV. The former VEVs can be taken as arbitrary values, while the value of the dilaton VEV is known phenomenologically. Since the size of the D-term contribution

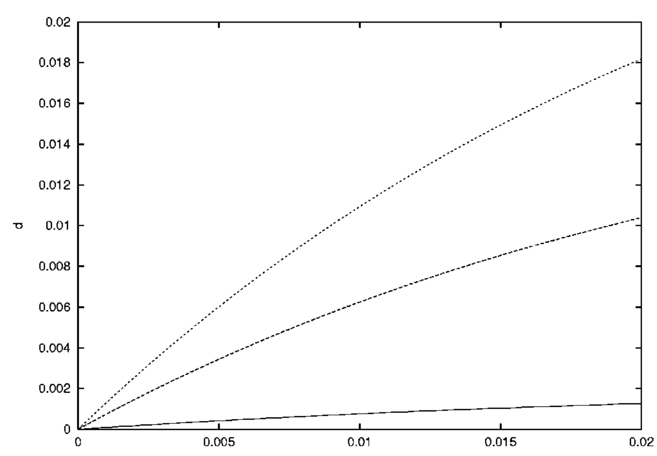

FIG. 2. The degeneracy of sfermion masses against $a = \alpha_A(\Lambda)\ln(\Lambda/M_X)$ for $b_A = 50$.

depends on the twisted moduli VEVs, it is important to find a proper mechanism for stabilizing their VEVs.

The observed differences originate mainly from the fact that the Kähler potential $\hat{K}$ of twisted moduli fields M_ℓ can take different forms from the dilaton Kähler potential. For instance, if $\hat{K}$ takes the tree level form (26), the M_ℓ-dependent FI term vanishes in the limit $\langle M_\ell \rangle \to 0$, and consequently the anomalous $U(1)$ D-term contribution also vanishes in that limit. Our results, e.g., Eq. (45), are consistent with this property. Another remark concerns the sign of the FI term; The FI term in type I models can take either sign depending on the sign of the twisted moduli VEVs. Again this property is in sharp contrast to the heterotic case. Further phenomenological impacts of these properties will be discussed elsewhere.

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