Anomalous U(1) D-term Contribution in Type I String Models

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Abstract

We study the D-term contribution for anomalous U(1) symmetries in type I string models and derive general formula for the D-term contribution, assuming that the dominant source of SUSY breaking is given by 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 different features from the case in heterotic string models. The differences originate from the different forms of Kähler potential between twisted moduli fields in type I string models and the dilaton field in heterotic string models.

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1 Introduction

Superstring theory is a promising candidate for unified theory including gravity. One of important features is that 4-dimensional (4D) string models have several moduli fields including the dilaton field. Their vacuum expectation values (VEVs) determine couplings of 4D effective theory, e.g. gauge couplings, Yukawa couplings and Fayet-Iliopoulos (FI) coefficients. These moduli fields have perturbatively a flat potential. Non-perturbative effects are expected to stabilize these moduli. Such non-perturbative 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 s-spectra would be measured in 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 only due to $F$-terms. That leads to the universal relation, $M_{1/2}^I = -A_{IJK} = \sqrt{3}m_{3/2}$ and $m_I^2 = |m_{3/2}|^2$, where $m_{3/2}$ is the gravitino mass, $M_{1/2}^I$ is the gaugino mass, $A_{IJK}$ is the $A$-term and $m_I$ is soft scalar masses, while SUSY breaking due to other sources leads to non-universal 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 unbounded from below (UFB) direction [4]. 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 of 4D string models for both heterotic models and type I models have anomalous $U(1)$ symmetries [6, 7, 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 non-linearly. Such 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 ($D$-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.

\footnote{This problem would not be serious, if the age of the Universe is not long enough to reach the CCB minimum.}
called the “$D$-term contribution” after gauge symmetries are broken down [11, 12]. This contribution has a linear dependence on the VEV of $D$-component and it is proportional to the charge of broken symmetry. These features are different from those in the contribution from $F$-component, which has the quadratic form of the VEVs of $F$-component and it does not depend on the charge of broken symmetry explicitly. A magnitude of $D$-term condensation has been studied in grand unified theories [12, 13].

Since most of 4D string models have an anomalous $U(1)$ symmetry, its breaking, in general, induces a $D$-term contribution to scalar masses. For 4D heterotic models, the $D$-term contribution has been examined [15, 16, 17, 18]. In particular, in Ref. [16] it is taken into account that the FI term is dilaton-dependent. As a result, even in the dilaton-dominant SUSY breaking the $D$-term contribution induces non-universal scalar masses, and the additional terms are proportional to $U(1)$ charges. That has phenomenologically important implications. For example, the CCB and UFB constraints can be relaxed [19]. As another aspect, these $D$-term contributions 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 $U(1)$ charges leads to non-universal sfermion masses, which are dangerous from the viewpoint of FCNC constraints.

In this way, it is an important subject to study a magnitude of $D$-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 in 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. 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 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 section 3, we study the $D$-term contribution for anomalous $U(1)$ symmetries in the framework of type I models in section 4. In section 5, we discuss phenomenological implications of $D$-term contributions. Section 6 is devoted to conclusion.

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6See Ref. [20, 21, 22] for SUSY breaking scenario by the $F$-term of the twisted moduli fields.
2 D-term contribution

We explain the D-term contribution to soft SUSY breaking scalar masses based on supergravity theory (SUGRA) [13]. The matter sector in SUGRA is specified by two functions, the total Kähler potential $G(\phi^I, \bar{\phi}^I)$ and the gauge kinetic function $f_{\alpha\beta}(\phi^I)$ with $\alpha, \beta$ being indices of the adjoint representations of the gauge groups. The former is a sum of the Kähler potential $K(\phi^I, \bar{\phi}^I)$ and the logarithm of the superpotential $W(\phi^I)$

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

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 $\phi^I$ and their complex conjugate by $\bar{\phi}^I$. The real part of gauge kinetic function $\text{Re}f_{\alpha\beta}$ is related to the gauge coupling constants $g_\alpha$ as follows,

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

The scalar potential is given by

$$
V = M^2 e^{G/M^2} \left( G_I(G^{-1})^{IJ}G_J - 3M^2 \right) + \frac{1}{2} \left( \text{Re}f^{-1} \right)_{\alpha\beta} G_I(T^\alpha \phi)^I G_J(T^\beta \phi)^J,
$$

where $G_I = \partial G/\partial \phi^I$, $G_I = \partial G/\partial \bar{\phi}^I$ etc, $(\text{Re}f^{-1})_{\alpha\beta}$ and $(G^{-1})^{IJ}$ are the inverse matrices of $\text{Re}f_{\alpha\beta}$ and $G_{IJ}$, respectively, e.g., $(G^{-1})^{IJ}G_{J'K'} = \delta_{I'}^{J'}$, $T^\alpha$ are gauge transformation generators, and a summation over $\alpha,...$ and $I,...$ is understood. The $(T^\alpha \phi)^I$ are variations under gauge transformations up to infinitesimal parameters and they become constants for non-linear realizations. The $F$-auxiliary fields and the $D$-auxiliary fields are given by

$$
F^I = Me^{G/2M^2}(G^{-1})^{IJ}G_J, \quad D^\alpha = (\text{Re}f^{-1})_{\alpha\beta} G_I(T^\beta \phi)^I,
$$

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

$$
V = V_F + V_D \equiv \left( F^J K_{JI} F^I - 3M^4 e^{G/M^2} \right) + \frac{1}{2} \text{Re}f_{\alpha\beta} D^\alpha D^\beta.
$$

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

$$
V_{\text{soft}} = (m_F^2)_{IJ} \phi^I \phi^J + (m_D^2)_{IJ} \phi^I \bar{\phi}^J + \cdots,
$$

\begin{align}
(m_F^2)_{IJ} & \equiv \left| m_{3/2} \right|^2 + \left( \frac{V_F}{M^2} \right) \langle K_{JI} \rangle \\
& \quad + \langle F^I \rangle \langle \bar{F}^J \rangle \left\langle \partial_{\nu_1} K_{IJ} (K^{-1})^{J\nu_1 \nu_2} \partial_{\nu_2} K_{\nu_2 J} - \partial_{\nu_1} \partial_{\nu_2} K_{IJ} \right\rangle,
\end{align}

\begin{align}
(m_D^2)_{IJ} & \equiv \langle D^\alpha \rangle \left\langle \frac{\partial}{\partial \phi^I} \phi^J \left( G_V(T^\alpha \phi)^I \right) \right\rangle,
\end{align}

where $V_F$ is the gravitational potential, $m_{3/2}$ is the gravitino mass, and $K$ is the total Kahler metric.
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 the phenomenological ground. 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 scalar masses \([11, 12]\). It is proportional to the charge of broken symmetry and appears when the rank of gauge group lowers on the breakdown of 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 \( \langle D^\alpha \rangle \),

\[
\left[ (M^2_V)^{\alpha\beta} + \left( \frac{V_F}{M^2} + 2|m_{3/2}|^2 \right) \langle \text{Re} f_{\alpha\beta} \rangle \right] \langle D^\beta \rangle \\
= \langle F^I \rangle \langle \bar{F}^J \rangle \left\{ \frac{\partial}{\partial \phi^I} \left( G_{T^\alpha \phi}^*(T^\alpha \phi)^I \right) \right\} + \frac{1}{2} \left\{ \frac{\partial}{\partial \phi^I} \text{Re} f_{\alpha\beta} \right\} \langle \langle (T^\alpha \phi)^I \rangle \langle D^\beta \rangle \langle D^\gamma \rangle \rangle , \tag{9}
\]

where \((M^2_V)^{\alpha\beta} = \langle \langle \bar{\phi} T^3 \rangle \rangle K_{IJ}(T^\alpha \phi)^I \rangle \) is the mass matrix of the gauge bosons, up to the normalization factor including the gauge coupling constants.

We require that the SUSY is broken down by non-vanishing \( F \)-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

\[
\langle D^\beta \rangle = \langle F^I \rangle \langle \bar{F}^J \rangle \left\{ \frac{\partial}{\partial \phi^I} \left( G_{T^\alpha \phi}^*(T^\alpha \phi)^I \right) \right\} \left( M^{-2}_V \right)^{\alpha\beta} , \tag{10}
\]

where \((M^{-2}_V)^{\alpha\beta} \) is the inverse matrix of \((M^2_V)^{\alpha\beta} \). The formula (10) is the master equation in our analysis.

For a later convenience, we write down the formula of gaugino masses \( M_{1/2}^\alpha \):

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

### 3 Anomalous \( U(1) \) \( D \)-term in heterotic string models

Effective SUGRA is derived from 4D string models taking a 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 \( K(\phi^I, \bar{\phi}^I) \) and the gauge kinetic function \( f_{\alpha\beta}(\phi^I) \) are given by

\[
K(\phi^I, \bar{\phi}^I) = -\ln(S + \bar{S} - 2\delta_{GS} V_A) - \sum_a \ln(T^a + \bar{T}^a) \\
+ \sum_\kappa \prod_\alpha (T^a + \bar{T}^a)^{a_\kappa} \bar{\phi}^\kappa e^{2\delta_{A\kappa} V_A \phi^\kappa} + \ldots , \tag{12}
\]

\[
f_{\alpha\beta}(\phi^I) = k_\alpha S \delta_{\alpha\beta} + \varepsilon_\alpha T^\alpha \delta_{\alpha\beta} , \tag{13}
\]
where $S$ is the dilaton field, $T^a$ are the moduli fields, $\phi^\kappa$ are matter fields with modular weights $n^\kappa_a$ and $U(1)_A$ charges $q^A_a$, and $V_A$ is the $U(1)_A$ vector superfield. Also in the above, $k_a$ is a Kac-Moody level (hereafter we set $k_a = 1$, for simplicity), $\varepsilon^\alpha_a$ is a model-dependent parameter coming from 1-loop correction and $\delta^A_{GS}$ is the GS coefficient of $U(1)_A$ given by

$$\delta^A_{GS} = \frac{1}{192\pi^2} \sum_\kappa q^A_\kappa .$$

(14)

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

$$D^A = (\text{Re} f^{-1})_A \left( \frac{\delta^A_{GS}}{S + \bar{S}} + \sum_\kappa \prod_a (T^a + \bar{T}^a)^{n^\kappa_a} q^A_\kappa |\phi^\kappa|^2 \right) ,$$

where we neglect terms from higher order terms in $K(\phi^I, \bar{\phi}^J)$. 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 non-linearly as $S \rightarrow S' = S + i\delta^A_{GS}\theta(x)$ under $U(1)_A$. Then the formula (10) for $\langle D^A \rangle$ reads

$$\langle D^A \rangle = \frac{1}{(M^2)^A} \left( 2\delta^A_{GS} \frac{|\langle F^S \rangle|^2}{\langle S + \bar{S} \rangle^3} + \sum_\kappa \prod_a (T^a + \bar{T}^a)^{n^\kappa_a} q^A_\kappa |\phi^\kappa|^2 \right) + \ldots \right) .$$

(16)

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

$$(M^2)^A = \frac{(\delta^A_{GS})^2}{\langle S + \bar{S} \rangle^2} + \left( \sum_\kappa \prod_a (T^a + \bar{T}^a)^{n^\kappa_a} (q^A_\kappa)^2 |\phi^\kappa|^2 \right) ,$$

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

$$(M^2)^A = \frac{\delta^A_{GS}}{\langle S + \bar{S} \rangle} \left( \frac{\delta^A_{GS}}{\langle S + \bar{S} \rangle} - \left( \frac{\sum_\kappa \prod_a (T^a + \bar{T}^a)^{n^\kappa_a} (q^A_\kappa)^2 |\phi^\kappa|^2}{\sum_\kappa \prod_a (T^a + \bar{T}^a)^{n^\kappa_a} q^A_\kappa |\phi^\kappa|^2} \right) \right) .$$

(17)

(18)

In explicit models, we find that $\delta^A_{GS} = \mathcal{O}(10^{-1}) - \mathcal{O}(10^{-2})$. Hence we will neglect terms with a higher order of $\delta^A_{GS}$. 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, $\langle D^A \rangle$ is given by

$$\langle D^A \rangle = \frac{1}{(M^2)^A} \left( 2\delta^A_{GS} \frac{|\langle F^S \rangle|^2}{\langle S + \bar{S} \rangle^3} + \sum_\kappa (T + \bar{T})^{n^\kappa_a} q^A_\kappa |F^\kappa|^2 \right) + \left( \frac{|\langle F^T \rangle|^2}{\langle T + \bar{T} \rangle} \sum_\kappa (T + \bar{T})^{n^\kappa_a} n^\kappa_a (n^\kappa_a - 1) q^A_\kappa |\phi^\kappa|^2 \right) + \left( \frac{\langle F^\bar{T} \rangle}{\langle T + \bar{T} \rangle} \sum_\kappa (T + \bar{T})^{n^\kappa_a} q^A_\kappa F^\kappa \bar{\phi}^\kappa \right) + \text{h.c.} .$$

(19)

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Footnote: In (12) and (13), all fields stand for superfields with a same notation for chiral superfields and anti-chiral superfields as its scalar components.
Here we consider the case that the dilaton and the overall moduli fields are dominant sources to 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

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

(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

\[
\langle V_F \rangle = \frac{|\langle F^S \rangle|^2}{(S + S)^2} + 3 \frac{|\langle F^T \rangle|^2}{(T + T)^2} - 3 |m_{3/2}|^2 .
\]

(21)

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

\[
\langle D^A \rangle = \frac{1}{\langle M^2_F \rangle^4} \left[ \frac{A^A}{S + S} \left( 2 \frac{|\langle F^S \rangle|^2}{(S + S)^2} - |m_{3/2}|^2 \right) \right.
\]

\[
- \frac{|\langle F^T \rangle|^2}{(T + T)^2} \left( \sum_\kappa (T + T)^{n_\kappa} n_\kappa q^A_\kappa |\phi^\kappa|^2 \right) \right] .
\]

(22)

By use of the parametrization

\[
\frac{\langle F^S \rangle}{(S + S)} \equiv \sqrt{5} C |m_{3/2}| e^{i\alpha_S \sin \theta} , \quad \frac{\langle F^T \rangle}{(T + T)} \equiv C |m_{3/2}| e^{i\alpha_T \cos \theta} ,
\]

(23)

the VEVs \( \langle V_F \rangle \) and \( \langle D^A \rangle \) can be expressed as

\[
\langle V_F \rangle = 3 \left( C^2 - 1 \right) |m_{3/2}|^2 ,
\]

(24)

\[
\langle D^A \rangle = |m_{3/2}|^2 \left[ \left( 1 - 6C^2 \sin^2 \theta \right) \frac{\langle \sum_\kappa (T + T)^{n_\kappa} q^A_\kappa |\phi^\kappa|^2 \rangle}{\langle \sum_\kappa (T + T)^{n_\kappa} q^A_\kappa |\phi^\kappa|^2 \rangle} \right.
\]

\[
- C^2 \cos^2 \theta \left( \frac{\langle \sum_\kappa (T + T)^{n_\kappa} n_\kappa q^A_\kappa |\phi^\kappa|^2 \rangle}{\langle \sum_\kappa (T + T)^{n_\kappa} q^A_\kappa |\phi^\kappa|^2 \rangle} \right) \right] .
\]

(25)

Here \( C \) is a constant and \( \theta \) is a parameter called the “goldstino angle”.

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

\[
M_{1/2}^\alpha = \frac{1}{2 \langle \text{Ref}_\alpha \rangle} \left( \langle F^S \rangle + \varepsilon_\alpha \langle F^T \rangle \right) .
\]

(26)

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, \( U(1)_A \) \( D \)-term
contribution to scalar masses appears, \( (m^2_2)_I = q^A_I \langle D^A \rangle \), and its magnitude is rather large as \( \langle D^A \rangle = \mathcal{O}(m^2_{3/2}) \). If \( U(1)_A \) charges are different between the first and second families, the non-universality among sfermion masses would be dangerous from the viewpoint of FCNC constraints. On the other hand, with \( U(1)_A \) \( D \)-term contribution we can relax the CCB and UFB bounds [19].

In a certain case [24], the \( F \)-components of matter fields can also contribute to the breakdown of SUSY. For instance, \( \langle F^\kappa \rangle \) contributes to the SUSY breaking if \( \langle \partial W/\partial \phi^\kappa \rangle = \mathcal{O}(m^3_3/2M) \). Then the dominant part of the \( D^A \) condensation comes from the second term in the r.h.s. of Eq. (19),

\[
\langle D^A \rangle = \frac{\langle \sum_{\kappa}(T + \bar{T})^n_{\kappa} q^A_I |F^\kappa|^2 \rangle}{\langle \sum_{\kappa}(T + \bar{T})^{n_{\kappa}}(q^A_I)^2 |\phi^\kappa|^2 \rangle}. \tag{27}
\]

The magnitude is estimated [18] as \( \langle D^A \rangle = \mathcal{O}(m^2_{3/2}/\delta^A_{GS}) \).

### 4 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 \( U(1)_i \) symmetries, i.e. \( \prod_i U(1)_i \). We denote the \( U(1)_i \) vector multiplet by \( V_i \). The Kähler potential \( K(\phi^I, \bar{\phi}^I) \) is given by

\[
K(\phi^I, \bar{\phi}^I) = \hat{K}\left(M_\ell + \bar{M}_\ell - 2 \sum_i (\delta_{GS})^I_{\ell} V_i \right) - \ln(S + \bar{S}) + \sum_a \ln(T^a + \bar{T}^a) + \sum_{\kappa} \prod(S + \bar{S})^{n^S_{\kappa}} (T^a + \bar{T}^a)^{n^a_{\kappa}} \bar{\phi}^\kappa e^{2q^A_I V_i \phi^\kappa} + \cdots , \tag{28}
\]

where chiral matter fields \( \phi^\kappa \) have the “modular weights” \( n^S_{\kappa} \) and \( n^a_{\kappa} \) with respect of \( S \) and \( T^a \), and \( (\delta_{GS})^I_{\ell} \) are model-dependent GS coefficients. Here, \( M_\ell \) is a twisted moduli field associated with the \( \ell \)-th fixed point. For simplicity, we use the notation \( m_\ell \) defined by \( m_\ell \equiv M_\ell + \bar{M}_\ell - 2 \sum_i (\delta_{GS})^I_{\ell} V_i \) 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]

\[
\hat{K}(m_\ell) = \frac{1}{2} m^2_\ell . \tag{29}
\]

The \( M_\ell \)-dependence of Kähler metric of \( \phi^\kappa \) is also unclear. In the orbifold limit, the Kähler metric \( K_{\kappa,\bar{\kappa}} \) does not depend on the twisted moduli \( M_\ell \) as in Eq. (28). For a large value of \( M_\ell \), however, it would receive a correction \( \Delta K_{\kappa,\bar{\kappa}}(M, \bar{M}) \).

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8See also Ref. [25].

9See for effective low-energy Lagrangian of type I models Ref. [3] and references therein.
The gauge kinetic function \( f_{\alpha \beta}(\phi^I) \) is given by
\[
f_{\alpha \beta}(\phi^I) = \hat{f}(S, T^a)\delta_{\alpha \beta} + \sum_{\ell} s_{\ell}^a M_{\ell \alpha \beta},
\] (30)
where \( s_{\ell}^a \) 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 \( M_{\ell} \to M_{\ell}' = M_{\ell} + i(\delta_{GS})_i^f \theta(x) \) under \( U(1)_i \).

The \( U(1)_i \) \( D \)-components are given by
\[
D^i = (\text{Ref}^{-1})_i \left( - (\delta_{GS})_i^\ell \frac{\partial \bar{K}}{\partial \bar{m}_\ell} + \sum_\kappa \Pi_a (S + \bar{S})^{n_\kappa} (T^a + \bar{T}^a) n_\kappa q_i^a |\phi^\kappa|^2 \right),
\] (31)
where we assume that \( U(1) \) kinetic mixing is absent for simplicity. According to the formula (10) for \( D \)-term condensation, we obtain
\[
\langle D^i \rangle = \frac{1}{(M_{V}^2)^i} \left( - (\delta_{GS})_i^\ell \left( \frac{\partial^3 \bar{K}}{\partial \bar{m}_\ell \partial \bar{m}_\ell' \partial \bar{m}_{\ell''}} F^{\bar{M}_{\ell'}} \bar{F}^{\bar{M}_{\ell''}} \right) \right.
\]
\[
\left. + \langle \sum_\kappa \Pi_a (S + \bar{S})^{n_\kappa} (T^a + \bar{T}^a) n_\kappa q_i^a |\phi^\kappa|^2 \rangle \right) + \ldots
\] (32)
Here \( (M_{V}^2)^i \) is (essentially) the \( U(1)_i \) gauge boson mass given by
\[
(M_{V}^2)^i \equiv (\delta_{GS})_i^\ell (\delta_{GS})_i^\ell \left( \frac{\partial^3 \bar{K}}{\partial \bar{m}_\ell \partial \bar{m}_\ell' \partial \bar{m}_{\ell''}} \right) + \langle \sum_\kappa \Pi_a (S + \bar{S})^{n_\kappa} (T^a + \bar{T}^a) n_\kappa q_i^a |\phi^\kappa|^2 \rangle
\]
\[
\left. + (\delta_{GS})_i^\ell \left( \frac{\partial \bar{K}}{\partial \bar{m}_\ell} \right) \frac{\langle \sum_\kappa \Pi_a (S + \bar{S})^{n_\kappa} (T^a + \bar{T}^a) n_\kappa q_i^a |\phi^\kappa|^2 \rangle}{\langle \sum_\kappa \Pi_a (S + \bar{S})^{n_\kappa} (T^a + \bar{T}^a) n_\kappa q_i^a |\phi^\kappa|^2 \rangle} \right) .
\] (33)
where we have used the almost \( D \)-flatness conditions of \( U(1)_i \). In the case with the canonical Kähler potential (29), the \( (M_{V}^2)^i \) is reduced to
\[
(M_{V}^2)^i \equiv \left( (\delta_{GS})_i^\ell \right)^2 + (\delta_{GS})_i^\ell \langle m_\ell \rangle \left( \frac{\langle \sum_\kappa \Pi_a (S + \bar{S})^{n_\kappa} (T^a + \bar{T}^a) n_\kappa q_i^a |\phi^\kappa|^2 \rangle}{\langle \sum_\kappa \Pi_a (S + \bar{S})^{n_\kappa} (T^a + \bar{T}^a) n_\kappa q_i^a |\phi^\kappa|^2 \rangle} \right) .
\] (34)
If \( \langle m_\ell \rangle \ll \mathcal{O}(\delta_{GS}) \), the first term is dominant in Eq. (34), 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 \( n_\kappa = \sum_a n_\kappa^a \). Then the \( \langle D^i \rangle \) in Eq. (32) can explicitly be written down as
\[
\langle D^i \rangle = \frac{1}{(M_{V}^2)^i} \left( - (\delta_{GS})_i^\ell \left( \frac{\partial^3 \bar{K}}{\partial \bar{m}_\ell \partial \bar{m}_\ell' \partial \bar{m}_{\ell''}} F^{\bar{M}_{\ell'}} \bar{F}^{\bar{M}_{\ell''}} \right) \right.
\]
\[
\left. + \langle \sum_\kappa \Pi_a (S + \bar{S})^{n_\kappa} (T^a + \bar{T}^a) n_\kappa q_i^a |\phi^\kappa|^2 \rangle \right),
\]
The expressions (35) or (38) are our main results for the D-term condensation in type I models. Note that $\langle \phi^\kappa \rangle \ll \mathcal{O}(M)$ and $\langle \partial W/\partial \phi^\kappa \rangle \ll \mathcal{O}(m_{3/2}M)$. In this case, since $\langle F^\kappa \rangle = \mathcal{O}\left(m_{3/2} \langle \phi^\kappa \rangle \right)$, the VEV $\langle V_F \rangle$ is simplified as

$$
\langle V_F \rangle = \frac{\langle F^S \rangle^2}{\langle S + S \rangle^2} + 3 \frac{\langle F^T \rangle^2}{\langle T + T \rangle^2} + \sum_{\ell \nu} \left( \frac{\partial^2 \hat{K}}{\partial m_{\ell} \partial m_{\nu}} F^{M_\ell} F^{M_{\nu}} \right) - 3 \left| m_{3/2} \right|^2.
$$

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

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

Then a careful calculation leads to

$$
\langle D^i \rangle = \frac{1}{(M_V^2)^\ell} \left( - \delta_{GS} \right)^i \left( \frac{\partial^3 \hat{K}}{\partial m_{\epsilon \ell} \partial m_{\nu} \partial m_{\nu}} F^{M_\epsilon} F^{M_{\nu}} \right)
$$

$$
+ \left| m_{3/2} \right|^2 \left( \sum_{\kappa} \langle S + \bar{S} \rangle \bar{n}_{\kappa} (T + \bar{T}) n_{\kappa} q_{\kappa}^i \langle \phi^\kappa \rangle^2 \right)
$$

$$
- \frac{\langle F^S \rangle^2}{\langle S + S \rangle^2} \left( \sum_{\kappa} \langle S + \bar{S} \rangle \bar{n}_{\kappa} (T + \bar{T}) n_{\kappa} q_{\kappa}^i \langle \phi^\kappa \rangle^2 \right)
$$

$$
- \frac{\langle F^T \rangle^2}{\langle T + T \rangle^2} \left( \sum_{\kappa} \langle S + \bar{S} \rangle \bar{n}_{\kappa} (T + \bar{T}) n_{\kappa} q_{\kappa}^i \langle \phi^\kappa \rangle^2 \right).
$$

The expressions (35) or (38) are our main results for the $D$-term condensation in type I models. Note that $\langle D^i \rangle$ becomes independent of $\langle F^{M_\ell} \rangle$ if the third derivative of $\hat{K}$ vanishes; $\langle \partial^3 \hat{K}/\partial m_{\ell} \partial m_{\nu} \partial m_{\nu} \rangle \ll \mathcal{O}\left((\delta_{GS})^\ell\right)$. 

9
The soft terms can be calculated by using the parametrization similar to Eq. (23),

\[
\langle F^S \rangle = \frac{\langle F^T \rangle}{\langle T + \bar{T} \rangle} \equiv C |m_{3/2}| e^{i\alpha T} \cos \theta \sin \phi ,
\]

\[
\langle F^{\lambda_\ell} \rangle \equiv \sqrt{3} C |m_{3/2}| e^{i\alpha_\ell \Phi_\ell} \cos \theta \cos \phi , \quad \sum_\ell \Phi_\ell^2 = 1 ,
\]

where we assumed that the Kähler metric of the twisted moduli is diagonal, and \( \theta, \phi \) and \( \Phi_\ell \) are “goldstino angles”. Then \( \langle V_F \rangle \) is the same as in Eq. (24), and the \( \langle D^i \rangle \) becomes

\[
\langle D^i \rangle = \frac{|m_{3/2}|^2}{(M_T^2)^i} \left( -3C^2 (\delta_{GS})_i^\ell \left\langle \frac{\partial^3 \hat{K}}{\partial m_\ell \partial m_\ell \partial m_\ell} \right\rangle \Phi_\ell^* \Phi_\ell \cos^2 \theta \cos^2 \phi + (\delta_{GS})_i^\ell \left\langle \frac{\partial \hat{K}}{\partial m_\ell} \right\rangle \right. \\
- 3C^2 \sin^2 \theta \left\langle \sum_\kappa (S + \bar{S})^{n_\kappa} (T + \bar{T})^{n_\kappa} n_\kappa^i q_\kappa^i |\phi_\kappa|^2 \right\rangle \\
- C^2 \cos^2 \theta \sin^2 \phi \left\langle \sum_\kappa (S + \bar{S})^{n_\kappa} (T + \bar{T})^{n_\kappa} n_\kappa^i q_\kappa^i |\phi_\kappa|^2 \right\rangle \right) .
\]

The gaugino masses \( M_{1/2}^\alpha \) are calculated by use of Eqs. (11) and (30),

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

\[
M_{1/2}^\alpha = \frac{1}{2 (\text{Re} f_\alpha)} \left( \langle F^T \rangle + \sum_\ell s_\ell^\alpha \langle F^{\lambda_\ell} \rangle \right) \quad \text{for D5}_\alpha \text{-branes} .
\]

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\(_\alpha\)-branes. If the dilaton and/or an overall moduli dominant SUSY breaking occur, the magnitude of \( U(1) \), \( D \)-term can be sizable as \( \langle D^i \rangle = \mathcal{O}(m_{3/2}^2) \). On the other hand, the magnitude of \( \langle D^i \rangle \) can be small if the twisted moduli fields dominate the SUSY breaking; in this case, \( \langle D^i \rangle \) 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 \( \langle m_\ell \rangle \ll \mathcal{O}((\delta_{GS})_i^\ell) \).

Alternatively we can suppose, as in the heterotic case, that there exists dynamical superpotential \( W \) of chiral matter fields \( \phi_\kappa \) so that \( \partial W / \partial \phi_\kappa = \mathcal{O}(m_{3/2} M) \). In such case, the dominant part of the \( D^i \) condensation comes from the second term in the r.h.s. of Eq. (35),

\[
\langle D^i \rangle = \frac{\left\langle \sum_\kappa (S + \bar{S})^{n_\kappa} (T + \bar{T})^{n_\kappa} n_\kappa^i q_\kappa^i |F^\kappa|^2 \right\rangle}{(M_T^2)^i} .
\]

The magnitude is estimated as \( \langle D^i \rangle = \mathcal{O}(m_{3/2}^2 M^2 / (M_T^2)^i) \).
5 Phenomenological implications

In this section, we discuss phenomenological implications of $D$-term contributions. An important point is that the FI terms depend on the twisted moduli fields in type I models, while such role is played by the dilaton field in heterotic models. Here let us compare our result (38) in type I models with the $D$-term (22) in heterotic models.

The first term of the $D$-term condensation (38) 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 explicitly. Recall that the $D$-term condensation (22) in heterotic models depends explicitly on $F^S$. The difference is originated from the different forms of Kähler potential between $M_\ell$ and $S$. If $\hat{K}(M_\ell, \bar{M}_\ell)$ is the logarithmic form like $S$, this difference would disappear.

For the remaining terms in Eq. (38), we can estimate the order of magnitudes by using the fact that the $D$-term (31) almost vanishes. The second term is proportional to the FI terms as in the heterotic case. The last two terms can be estimated as

\[ \langle \sum_\kappa (S + \bar{S})^{n_\kappa} (T + \bar{T})^{n_\kappa} q_\kappa^\ell |\phi^\kappa|^2 \rangle = O\left( (\delta_{GS})_i^\ell \left< \frac{\partial \hat{K}}{\partial m_\ell} \right> \right), \tag{43} \]

\[ \langle \sum_\kappa (S + \bar{S})^{n_\kappa} (T + \bar{T})^{n_\kappa} q_\kappa^\ell |\phi^\kappa|^2 \rangle = O\left( (\delta_{GS})_i^\ell \left< \frac{\partial \hat{K}}{\partial m_\ell} \right> \right). \tag{44} \]

Therefore, we have

\[ \langle D^i \rangle = m_{3/2}^2 \times O\left( \frac{1}{(M_\ell^2)^i} (\delta_{GS})_i^\ell \left< \frac{\partial \hat{K}}{\partial m_\ell} \right> \right), \tag{45} \]

and its magnitude depends on $\langle m_\ell \rangle/(\delta_{GS})_i^\ell$ as is seen from Eq. (34). Notice that unlike the dilaton VEV in the heterotic case, the VEVs of twisted moduli fields $m_\ell$ can be taken as arbitrary value, depending on stabilization mechanism [20, 21, 22].

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

On the other hand, it is possible that the VEVs of twisted moduli fields $m_\ell$ are suppressed, i.e., $O\left( \langle m_\ell \rangle/(\delta_{GS})_i^\ell \right) \ll 1$. In this case, $D$-term contribution can be suppressed. This is a sharp contrast to the heterotic case where $D$-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 of whose end is on the D9 branes. In this case, the gaugino mass is obtained

\[ M_{1/2} = \sqrt{3} m_{3/2} , \]

where we have taken \( |s_\ell^\alpha \langle m_\ell \rangle| \ll \text{ReS} \). Since \( n^K = 0 \), the F-term scalar masses are universal, i.e.,

\[ (m^2_F)_I = |m_{3/2}|^2 . \]

This spectrum is the same as the dilaton-dominant SUSY breaking in heterotic models. In addition, we have to add the D-term contribution, \( (m^2_D)_I = q_i^I \langle D^i \rangle \) with \( \langle D^i \rangle \) given by Eq. (40). When \( \left| \langle m_\ell \rangle / (\delta_{GS})^I \right| \ll 1 \), the D-term contribution becomes simplified to

\[ (m^2_D)_I = q_i^I \left| m_{3/2} \right|^2 \frac{\langle m_\ell \rangle}{(\delta_{GS})^I} . \]

Thus, if \( \left| \langle m_\ell \rangle / (\delta_{GS})^I \right| \ll 1 \), the D-term contribution is small and the total soft scalar masses become almost universal. This has important implications on FCNC constraints as well as CCB and UFB bounds. That is, if this \( U(1) \) symmetry is relevant to the flavor symmetry, the suppressed D-term contribution would be favorable to avoid FCNC constraints.\(^{10}\)

For this purpose, we need to realize a suppression like \( \left| \langle m_\ell \rangle / (\delta_{GS})^I \right| \ll O(10^{-2}) \) for \( m_{3/2} = O(100) \) GeV. Whether it is possible or not, depends on the stabilization mechanism of \( M_\ell \).

Next let us consider the case that the single twisted moduli field \( M_\ell \) gives a dominant source in the SUSY breaking. In this case, the gaugino mass is written as

\[ M_{1/2}^\alpha = \frac{\sqrt{3}}{2} s_\ell^\alpha g_\ell^2 m_{3/2} . \]

It is interesting to note that if \( s_\ell^\alpha \) are proportional to the coefficients of 1-loop beta function of gauge couplings, like the case of ‘mirage gauge coupling unification’ in Ref. [30], this spectrum of gaugino masses resembles that in the anomaly mediation scenario [31]. Since Kähler potential of matter fields does not depend on \( M_\ell \) for small \( \langle M_\ell \rangle \), the F-term scalar masses are universal, i.e.,

\[ (m^2_F)_I = |m_{3/2}|^2 . \]

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

\[ (m^2_D)_I = q_i^I \left| m_{3/2} \right|^2 \left( -3C^2 \frac{\delta^3 \hat{K}}{(\delta_{GS})^I} \left( \frac{\delta \hat{K}}{\delta m_\ell \partial m_\ell \partial m_{\ell'}} \right) \Phi_\ell \Phi_{\ell'} + \frac{\langle m_\ell \rangle}{(\delta_{GS})^I} \right) . \]

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

\(^{10}\)See e.g. Ref. [28], where such flavor \( U(1) \) symmetry is discussed in a type-I inspired model.
We note that even if flavor-dependent $D$-term contributions can be suppressed, radiative corrections due to the gaugino mass of (gauged) flavor $U(1)$ symmetry would be important. That might generate sizable effects on FCNC processes [29] when the gauged $U(1)$ is relevant to the flavor symmetry.

6 Conclusion

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

We have 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 can be taken as arbitrary values, although it depends on the stabilization mechanism. That is, whether $D$-term contributions are sizable or can be suppressed, depends on the VEVs of twisted moduli. This aspect will be important for CCB/UFB bounds and FCNC constraints.

The observed differences originate from the fact that in type I models, Kähler potential $\hat{K}$ of twisted moduli fields $M_\ell$ can take different forms from the dilaton Kähler potential. For instance, if the $\hat{K}$ takes the tree level form (29), the $M_\ell$-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. (48), are consistent with this property. Another remark concerns the sign of the FI term; The FI term in type I models can take both signs depending on the sign of the twisted moduli VEVs. Again this property is sharp contrast to the heterotic case. Further phenomenological impact of these properties will be discussed elsewhere.

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