The Fayet-Iliopoulos term in Type-I string theory and M-theory

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Abstract

The magnitude of the Fayet-Iliopoulos term is calculated for compactifications of Type-I string theory and Horava-Witten M-theory in which there exists a pseudo-anomalous $U(1)_X$. Contrary to various conjectures, it is found that in leading order in the perturbative expansion around the weakly-coupled M-theory or Type-I limits, a result identical to that of the weakly-coupled $E_8 \times E_8$ heterotic string is obtained. The result is independent of the values chosen for the Type-I string scale or the size of the M-theory 11th dimension, only depending upon Newton’s constant and the unified gauge coupling.

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1 Introduction

One of the most phenomenologically useful features of many string compactifications is the existence of a U(1) symmetry with apparent field theoretic anomalies \[1\]. These anomalies are cancelled by a four-dimensional version of the Green-Schwarz mechanism \[2\], which involves shifts of the model independent axion of string theory under gauge and general coordinate transformations. Such shifts are able to compensate the field-theoretic anomalies of the pseudo-anomalous U(1)\(_X\) symmetry if and only if it has equal U(1)\(_X\)G\(_2\) anomalies for all gauge groups G as well as with gravity (here and in the following we omit for simplicity the affine level factors \(k_G\) which can easily be reinstated if required).

As originally noticed by Dine, Seiberg, and Witten \[1\], when this mechanism is combined with supersymmetry, a Fayet-Iliopoulos (FI) term is induced in the D-term for U(1)\(_X\). For the weakly-coupled E\(_8\)×E\(_8\) heterotic string the magnitude of the induced FI term depends on the field-theoretic anomaly coefficient, the string coupling, \(g_{str}\), and the reduced Planck mass \(M_s = M_{\text{Planck}}/\sqrt{8\pi}\). A detailed calculation gives \[1, 3, 4\]

\[
\xi^2 = \frac{g^2_{str} \text{Tr}(Q_X)}{192\pi^2} M_s^2 ,
\]

where, following convention, this is expressed in terms of the total U(1)\(_X\)(gravity)\(^2\) anomaly proportional to \(\text{Tr}(Q_X)\).\[3\]

This FI term allows many phenomenologically interesting applications. These include, an alternate non-grand-unified theory explanation of the successful \(\sin^2 \theta_w = \frac{3}{8}\) relation at the unification scale \[4\], the possible flavor universal communication of supersymmetry breaking from a hidden sector to the minimal supersymmetric standard model (MSSM) \[3\], and an explanation of the texture of the fermion masses and mixings \[7\]. Many such applications require for their success the existence of a hierarchy between the contributions induced by the presence of the FI term and those contributions arising from other supergravity or string theory modes. Since the strength of supergravity corrections is controlled by the reduced Planck mass, the appropriate expansion parameter is expected to be

\[
\varepsilon \equiv \frac{\xi^2}{M_s^2} = \frac{g^2_{str} \text{Tr}(Q_X)}{192\pi^2} .
\]

For the value of \(g_{str}\) we may take the MSSM unified gauge coupling, however, a remaining uncertainty in \(\varepsilon\) is the size of the gravitational anomaly coefficient. For realistic string models, \(\text{Tr}(Q_X)\) is certainly quite large; for instance, in one typical semi-realistic example, (based on the free-fermionic construction) \(\text{Tr}(Q_X) = 72/\sqrt{3}\) \[3\]. However since the gravitational anomaly is sensitive to all chiral states in the model, both in the MSSM and hidden sectors, additional matter beyond the MSSM or large hidden

\[^{1}\text{An interesting aspect of this result is that it evades the “proof” that a non-zero FI term is inconsistent with local supersymmetry.}\]
sectors can increase $\text{Tr}(Q_X)$ dramatically. Thus in the context of the weakly coupled $E_8 \times E_8$ string the expansion parameter is unlikely to be any smaller than $\varepsilon \sim 1/100$.

The precise value of the expansion parameter $\varepsilon$ is quite important phenomenologically. For example, in the case that the soft masses of the MSSM squarks and sleptons are communicated by $U(1)_X$, the MSSM gauginos, being necessarily uncharged under $U(1)_X$, only receive soft masses via supergravity. Thus the gaugino masses are suppressed by the factor $\varepsilon$ relative to the soft masses of the MSSM squarks and sleptons (by assumption $U(1)_X$ charged – at least for the first and second generations) $[6]$. This leads to either unacceptably light MSSM gauginos, or so heavy squarks and sleptons that naturalness, or stability of the $U(1)_{\text{em}}$ and $SU(3)$ color preserving minimum is threatened. Although this problem is ameliorated when a non-zero F-component for the dilaton superfield of string theory is included $[1]$, it is presently unclear if a fully consistent model is possible (especially the question of whether $F$-terms for the other, non-universal, moduli are induced). In any case it is certainly interesting to ask how the $U(1)_X$ scenarios become modified outside of the weakly-coupled $E_8 \times E_8$ case.

Another of the proposed uses of the string-induced FI term is D-term inflation $[10]$. In these models the inflationary vacuum energy density is due to a non-zero D-term rather than the more usual F-term contributions. This has the advantage, over the F-term dominated models of supersymmetric inflation of naturally possessing sufficiently flat inflaton potentials for successful slow-roll inflation, even after the effects of spontaneous supersymmetry breaking due to $V_{\text{vac}} \neq 0$ are taken into account. Since the magnitude of the FI term sets the size of the Hubble constant during inflation, as well as its rate of change along the slow-roll direction, it sets in turn the magnitude of the fluctuations in the cosmic microwave background radiation (CMBR). An analysis $[10]$ shows that the observed amplitude of fluctuations in the CMBR requires $\xi = 6.6 \times 10^{15}$ GeV. This is to be compared with the larger value, $\xi_{\text{string}} \gtrsim 2 \times 10^{17}$ GeV, obtained by substituting the deduced value of the MSSM unified coupling into the weakly-coupled heterotic string theory prediction Eq.(1). The most attractive proposal for the solution of this mismatch is that the weak-coupling prediction for $\xi$ is modified in the strong-coupling (Horava-Witten M-theory, or Type-I) limit. This seems not unreasonable since it is well known that in such a limit many $E_8 \times E_8$ weak-coupling predictions are amended, most notably the prediction for the scale of string unification $[11]$.

Given these various motivations for considering the FI term in more general contexts, we show in the following sections that a simple calculation of the induced FI term in Type-I string theory and M-theory is possible. The calculation is quite general, being independent of almost all details of the compactification down to four dimensions. However it is important to keep in mind two assumptions under which we work: First, we assume that the standard model gauge group arises from within the perturbative excitations of the theory (rather than from D-brane dynamics), and second, that there are not large corrections to the normalisation of the kinetic terms of various fields arising from corrections to the minimal Kahler potential. Although
such corrections to the Kahler potential are certainly possible as we move away from
the weakly coupled M-theory or Type-I domain into the domain of intermediate cou-
pling, it seems very likely that the FI term itself is protected from such corrections
by the surviving $N = 1$ spacetime supersymmetry.

2 The FI term

The calculation of the magnitude of the FI term induced in a string theory with an
anomalous $U(1)_X$ symmetry is greatly simplified by a judicious use of supersymmetry
together with the anomaly constraints. Since we are concerned with the phenomeno-
logical applications of the FI-term, we are interested in compactifications of either
Type-I string theory or M-theory down to four dimensions, in which the low-energy
limit is an $N=1$ supersymmetric field theory. In this situation the dilaton $\phi$ and the
model-independent axion $a$ combine to form the lowest complex scalar component of
the dilaton chiral superfield, $S = \exp(-2\phi) + ia + \ldots$. Expanded out in terms of
components, the coupling of the dilaton superfield to a gauge field kinetic term of a
gauge group $G$ reads

$$
\frac{1}{4} \int d^2 \theta S W^a W^a + \text{h.c.} = -\frac{1}{4e^{2\phi}} (F^a)^2 + \frac{1}{2e^{2\phi}} (D^a)^2 - \frac{a}{4} F^a F^a + \ldots ,
$$

(3)

where $F^a_{mn} = \frac{1}{2} \epsilon_{mnlp} F^{lp}$. Here four-dimensional Lorentz indices are denoted $m, n, \ldots$
and $a, b, \ldots$ are adjoint indices of $G$. The expectation value of the dilaton sets the
value of the four-dimensional gauge couplings at an appropriate scale, here always
taken to be the string unification scale, so $\langle e^{2\phi} \rangle = g_{\text{unif}}^2$.

It is easiest to calculate the FI term by starting with the field-theoretic expres-
sion for the $U(1)_X G^2$ anomaly. Such an anomaly implies that under a $U(1)_X$
gauge transformation $v_m \rightarrow v_m + \partial_m \omega$ the low-energy effective action is not invariant, but
transformes as

$$
\delta L_{\text{eff}} = \frac{\omega A}{8\pi^2} F \tilde{F} ,
$$

(4)

where $A$ is the anomaly coefficient. In a supersymmetric theory $\omega$ gets promoted to a
chiral superfield $\Omega$, and the $U(1)_X$ vector superfield $V$ transforms as $V \rightarrow V + i(\Omega - \Omega^*)/2$.

The apparent anomaly in the low-energy effective action Eq.(4) is cancelled by
assigning a shift in the axion under $U(1)_X$ gauge transformations. In terms of super-
fields,

$$
\delta S = -i \frac{A}{8\pi^2} \Omega,
$$

(5)

and then the shift in the gauge kinetic term Eq.(3) compensates the triangle anomaly
Eq.(4). However, the non-trivial transformation of $S$ implies that the conventional
Kahler potential term for the dilaton superfield $\ln(S + S^*)$ is no longer invariant. The
dilaton superfield kinetic energy term must be modified to

\[ C^2 \int d^4 \theta \ln \left( S + S^* + \frac{A}{4\pi^2} V \right), \quad \text{(6)} \]

then the transformation of the U(1) vector superfield cancels the variation. (In Eq.(6) \( C^2 \) is an all important normalization factor that will be calculated below from the underlying Type-I or M-theory.) The most important feature of Eq.(6) from our point view is that when this modified kinetic term is expanded out in component form, a term \textit{linear in the U(1) D-term} occurs:

\[ C^2 \left( \frac{1}{2} (\partial \phi)^2 + \frac{e^{4\phi}}{4} (\partial a)^2 + \frac{Ae^{2\phi}}{16\pi^2} D + \ldots \right). \quad \text{(7)} \]

This is nothing other than the FI term, \( \xi^2 \), (\( \xi \) has mass dimension 1 in the conventions being followed here):

\[ \xi^2 = \frac{C^2 A \langle e^{2\phi} \rangle}{16 \pi^2} = \frac{C^2 A g^2_{\text{unif}}}{16 \pi^2}. \quad \text{(8)} \]

Thus we see that supersymmetry allows us to calculate the FI term given the coefficient of the anomaly, \( A \), and the normalization, \( C^2 \), of the model-independent axion kinetic energy, in the basis that the axion itself is normalized so that it couples to the G gauge field strengths as in Eq.(3).

To fix \( C \) it is necessary to consider the origin of the axion \( a \) in string (or M-) theory. Within the framework of an effective 10-dimensional string-derived supergravity, the model independent axion arises from the 3-form field strength \( H_{\mu\nu\rho} \), which satisfies the Green-Schwarz modified Bianchi identity

\[ dH = - \text{tr} F \wedge F + \text{tr} R \wedge R. \quad \text{(9)} \]

Once we compactify down to four dimensions we may define the effective four-dimensional axion \( a \) in terms of \( H \) via

\[ H_{mnp} = M^2 \varepsilon_{mnpq} \partial^q a, \quad \text{(10)} \]

where \( M \) is a dimensionful normalization factor introduced so that \( a \) has mass dimension zero in four dimensions, and which can be fixed by demanding that \( a \) couples to \( F \tilde{F} \) correctly.

The modified Bianchi identity, Eq.(9), implies that \( a \) satisfies the equation of motion

\[ M^2 \partial^2 a = \frac{1}{2} F^a \tilde{F}^a + (R - \text{dependent terms}), \quad \text{(11)} \]

which is equivalent to the statement that \( a \) has the required axion-like interaction term \( a F \tilde{F} \) in the effective action up to a normalization factor. Furthermore, after compactification, the low-energy four-dimensional kinetic energy for \( H_{mnp} \)

\[ \mathcal{L}_{\text{eff,4d}} = \ldots + NH_{mnp}H^{mnp} + \ldots, \quad \text{(12)} \]
induces the kinetic energy term for the axion. To proceed we must convert this to the normalization used in Eqs.(3), (7). Fixing the coupling of the axion to $F \tilde{F}$ gives the relation $M^2 = 1/32N$, while we find that $C^2 = e^{-4\phi}/32N$, and thus using Eq.(8)

$$\xi^2 = \frac{A}{512\pi^2 g_{\text{unif}}^2 N}.$$  \hfill(13)

Thus the calculation of the FI term is reduced to a computation of the coefficient, $N$, of the $H$ kinetic term in the four-dimensional low-energy effective action after compactification.

3 The low-energy effective action in M-theory

The Horava-Witten [12] 11-dimensional low-energy supergravity action is, in component form, given by (see also [13]):

$$\mathcal{L}_{11d} = \frac{1}{\kappa^2} \int_M d^{11}x \sqrt{g} \left( -\frac{1}{2} R - \frac{1}{48} G_{IJKL}G^{IJKL} - \frac{\sqrt{2}}{3456} \varepsilon^{I_1\ldots I_{11}} C_{I_1I_2I_3} G_{I_4\ldots I_7} G_{I_8\ldots I_{11}} \right)$$

$$+ \sum_{i=1,2} \frac{1}{8\pi (4\pi k^2)^{2/3}} \int_{\partial M} d^{10}x \left( -\text{tr} F_i^2 + \frac{1}{2} \text{tr} R^2 \right) + \text{fermi terms.} \hfill(14)$$

This describes a system in which the supergravity modes propagate in the 11-d bulk of $M$, while the $E_8 \times E_8$ super Yang-Mills (SYM) degrees of freedom are localized to the two 10-dimensional orbifold hyperplanes $\partial M_i$ ($i = 1, 2$) of $M$ at $x^{11} = 0, \pi \rho$ respectively. The normalization of the gravitational part of the action is appropriate for the 11th dimension being an $S^1$, of circumference $2\pi \rho$, with the fields satisfying $Z_2$ orbifold conditions (the $E_8$ generators obey $\text{tr}(T^a T^b) = \delta^{ab}$). The 4-form field strength $G_{IJKL}$ satisfies a modified Bianchi identity

$$(dG)_{11, IJKL} = -\frac{3\sqrt{2}}{2\pi} \left( \frac{\kappa}{4\pi} \right)^{2/3} \delta(x^{11}) (\text{tr} F_{I,J} F_{K,L}) - \frac{1}{2} R_{[I,J} R_{K,L]}$$

$$- \frac{3\sqrt{2}}{2\pi} \left( \frac{\kappa}{4\pi} \right)^{2/3} \delta(x^{11} - \pi \rho) (\text{tr} F_{[I,J} F_{K,L]} - \frac{1}{2} R_{[I,J} R_{K,L]}), \hfill(15)$$

and will provide, after compactification, the four-dimensional model-independent axion $a$.

Before we calculate the normalization of the axion kinetic term arising from this action, it is necessary to eliminate $\kappa$ and $\rho$ in terms of the 4d gauge and gravitational couplings. We can compactify the theory down to four dimensions, preserving $N = 1$ supersymmetry, by taking $M = S^1 \times R^4 \times K$, where $K$ is a (possibly deformed) Calabi-Yau manifold, and where, as usual, the spin connection is embedded in the $E_6$ gauge connection. This leaves an $E_6$ SYM theory which can be further broken by Wilson lines, etc. (We assume in the following analysis that the visible MSSM gauge
group arises from the $i = 1$ boundary theory.) Denote the volume of the Calabi-Yau measured at the $i = 1$ boundary by $V(0)$. As shown by Horava and Witten, in leading non-trivial order in the expansion in $\kappa^{2/3}$, the volume of $K$ decreases linearly with $x^{11}$. The maximum allowed value of $\rho$ is such that $V(x^{11} = \pi \rho)$ approximately equals the Planck volume. In any case, compactifying the bulk supergravity action on $S^1 \times K$ leads to

$$\frac{1}{2\kappa^2} \int d^{11}x \sqrt{g} R \to \frac{1}{2\kappa^2} \int d^4x \sqrt{g^{(4)}} V(0) \frac{1}{2} 2\pi \rho R^{(4)},$$

where the factor of $V(0)/2$ arises from averaging over the Calabi-Yau volume ($V(0) >> V(\pi \rho)$ in the scenarios where the hidden $E_8$ is strongly coupled). Comparing this to the standard Einstein-Hilbert action in four dimensions gives the usual result

$$G_N = \frac{\kappa^2}{8\pi^2 V(0) \rho}.$$  \hspace{1cm} (17)

Similarly reducing the $i = 1$ SYM boundary theory leads to an expression for the unified gauge coupling $\alpha_{\text{unif}}$ of the four-dimensional theory

$$\alpha_{\text{unif}} = \frac{(4\pi \kappa^2)^{2/3}}{2V(0)}.$$ \hspace{1cm} (18)

It is important to recall that beyond leading order quantum corrections modify these classical relations; however, the above expressions for $G_N$ and $\alpha_{\text{unif}}$ suffice to determine the FI term.

As discussed by Witten [11], the additional dependence of $G_N$ on the size $\rho$ of the 11th dimension relative to $\alpha_{\text{unif}}$ (as compared to the standard weak-coupling $E_8 \times E_8$ result) allows one to somewhat adjust the string prediction of the unification scale. Roughly speaking the four-dimensional MSSM unification scale can be defined as the scale at which either string excitations, or the additional six dimensions in which the gauge degrees of freedom propagate, first become apparent. An arbitrarily low unification scale is not achievable within the Horava-Witten picture (at least if our world is situated at the weakly coupled boundary), since the second $E_8$ SYM theory moves into the strongly-coupled domain as $\rho$ increases. Indeed a calculation [11] indicates that it is just possible to take the length of the 11th dimension long enough so as to lower the string unification scale to match that inferred in the MSSM, $M_{\text{unif}} \simeq 2 \times 10^{16}$ GeV. In the following we will keep the volume of the Calabi-Yau space and the length of the 11th dimension as free parameters.

Having completed this preliminary discussion, the effective 3-form field strength $H$ is defined as $G_{11,mnp} = \eta H_{mnp}$. Here the normalization factor $\eta$ is necessary so that the gauge dependent piece of the modified Bianchi identity for $H$ restricted to the $i = 1$ boundary theory reads $dH = -\text{tr} F^1 \wedge F^1$, as in Eq.(9). It follows from the M-theory version of the Bianchi identity Eq.(15), that

$$\eta = \frac{3\sqrt{2}}{4\pi^2} \rho \left(\frac{\kappa}{4\pi}\right)^{2/3}.$$ \hspace{1cm} (19)
The axion kinetic term arises from the $G^2$ term in Eq.$(14)$, which expressed in terms of the correctly normalized 3-form field strength reads, after compactification to four dimensions:

$$
\frac{1}{12\kappa^2} \int_M d^{11}x \sqrt{g} G_{ABC,11}^2 \rightarrow \frac{1}{12\kappa^2} \pi \rho V(0) A^2 \int_M d^4x \sqrt{g} H_{mnp}^2.
$$

Comparing to the general expression for the $H$ kinetic energy leads gives $N$ which, using the relations Eqs.$(17)$ and $(18)$ between $G_N$, $\alpha_{\text{unif}}$, $\kappa$, $\rho$, and $V(0)$, may be written

$$
N = \frac{3}{32\pi^3 (4\pi)^{4/3} \rho \kappa^{2/3}} \frac{V(0)}{\rho \kappa V(0)} = \frac{3G_N}{16\pi \alpha_{\text{unif}}^2}.
$$

Finally, substituting this into the expression for the FI term, Eq.$(13)$, gives

$$
\xi_{\text{HW}}^2 = \frac{A \rho^2}{192\pi^2}.
$$

This is identical to the value obtained for the weakly-coupled $E_8 \times E_8$ heterotic string, independent of the size of the 11th dimension, or equivalently the unification scale in this strongly-coupled $E_8 \times E_8$ picture.

4 Type-I string theory

For Type-I string theory, the situation is quite similar. We start from the 10-dimensional supergravity Lagrangian derived from its' low-energy limit $[14]$:

$$
\mathcal{L}_{\text{Type-I}} = - \int d^{10}x \sqrt{g} \left( 8\pi^7 e^{-2\phi_I} \right) R + \frac{e^{-\phi_I}}{16(\alpha'_I)^3} F^a F^a + 3 \frac{1}{2048\pi^7 (\alpha'_I)^2} H^2 + \ldots.
$$

Here $H$ is the 3-form field strength, $\phi_I$ the 10d dilaton of Type-I string theory, and, as usual, $\alpha'_I$ has dimension of an inverse mass-squared. If we compactify this theory on a Calabi-Yau of volume $V$, the classical expressions for Newton’s constant and the 4d unified gauge coupling that follow from comparing to the standard 4d action are

$$
G_N = \frac{e^{2\phi_I}(\alpha'_I)^4}{128\pi^8 V},
$$

$$
\alpha_{\text{unif}} = \frac{e^{\phi_I}(\alpha'_I)^3}{\pi V}.
$$

7
As recently emphasised by Witten [11], it is now the different dilaton dependencies of the gauge and gravitational pieces of the action that allow one to adjust the string prediction of the unification scale to match the inferred MSSM unification scale. However unlike the case in either the weakly-coupled $E_8 \times E_8$ heterotic string or the standard Horava-Witten M-theory with $E_8^2$ strongly coupled, the Type-I unification scale can be made arbitrarily low. In the Type-I case of interest, the unification scale corresponds to leading order to $M_{\text{unif}} \simeq 2\pi (V^{-1/6})$, the mass of the first Kaluza-Klein excitations, and as this is lowered, the type-I string coupling is also reduced, moving the theory farther into the weakly coupled domain. (However one must be careful that the string worldsheet coupling is not becoming strong, see Ref. [15] for an extensive discussion of these issues.)

Independent of $V$, the expressions Eq.(24) satisfy the standard Type-I relation between the reduced Planck mass, $M_*$, and the gauge coupling, $\alpha_{\text{unif}} M_*^2 = 128 \pi^7 e^{-\phi_I}/\alpha'_I$. (Note that this relation differs by numerical factors from that stated in Ref. [11].) Solving for $\alpha'_I$ from Eqs.(24) gives

$$ (\alpha'_I)^2 = \frac{\alpha_{\text{unif}}^2 M_*^2 V}{16 \pi^5}. \quad (25) $$

As before, the quantity of interest for the calculation of the FI term is the coefficient of the compactified 3-form field strength kinetic term expressed in terms of $\alpha_{\text{unif}}, M_{\text{unif}}$ and $M_*$. Using Eq.(25) leads to

$$ N = \frac{3V}{2048 \pi^7 (\alpha'_I)^2} = \frac{3}{128 \pi^2 \alpha_{\text{unif}}^2 M_*^2}, \quad (26) $$

and thus from Eq.(13) we find that the FI term in Type-I string theory is given by

$$ \xi_I^2 = \frac{AM_*^2 g_{\text{unif}}^2}{192 \pi^2}, \quad (27) $$

again identical to the result in the weakly-coupled $E_8 \times E_8$ heterotic string. This is independent of the Type-I string scale, or equivalently the effective four-dimensional unification scale, at least as long as one stays in the domain of applicability of the above calculation.

5 Comments

We have shown that a quite simple calculation of the induced Fayet-Iliopoulos term is possible in the contexts of Type-I string theory and Horava-Witten M-theory. For fixed Newton’s constant and four-dimensional unified gauge coupling the value of the FI term is independent of all details of the underlying Type-I string theory or M-theory. This includes the moduli of the compactification down to four dimensions, the underlying string coupling, and in the case of M-theory the length of the 11th dimension.
One significant caveat to this work is that we have always assumed that the
gauge degrees of freedom of the MSSM appear in the perturbative spectrum of the
Type-I string theory or Horava-Witten M-theory. Clearly this is not a necessity,
for the MSSM degrees of freedom could in principle arise from D-brane excitations.
Interestingly, FI terms for theories containing anomalous U(1)’s do arise in the D-
brane construction of six- and four-dimensional gauge theories (see for example the
discussions of Refs. [16] and [17]). It would certainly be worthwhile to consider such
FI terms in the case of realistic (or semi-realistic) D-brane models of the unified
MSSM interactions.

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