Generalised Fayet–Iliopoulos Terms in Supergravity

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Abstract—The U(1) vector multiplet theory with the Fayet–Iliopoulos (FI) term is one of the oldest and simplest models for spontaneously broken rigid supersymmetry. Lifting the FI term to supergravity requires gauged R-symmetry, as was first demonstrated in 1977 by Freedman within N = 1 supergravity. There exists an alternative to the standard FI mechanism, which is reviewed in this conference paper. It is obtained by replacing the FI model with a manifestly gauge-invariant action such that its functional form is determined by two arbitrary real functions of a single complex variable. One of these functions generates a superconformal kinetic term for the vector multiplet, while the other yields a generalised FI term. Coupling such a vector multiplet model to supergravity does not require gauging of the R-symmetry. These generalised FI terms are consistently defined for any off-shell formulation for N = 1 supergravity, and are compatible with a supersymmetric cosmological term.

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1. INTRODUCTION

As is well-known, the observed accelerating expansion of the universe can be accounted for by a small positive cosmological constant. It is therefore desirable to look for theoretical mechanisms that naturally lead to the required features of the cosmological constant, which are: (i) its positivity; and (ii) its small value. In the last five years it has been widely appreciated that such a mechanism is potentially provided by spontaneously broken local supersymmetry. In particular, this is one of the ideas which have been advanced within the framework of pure de Sitter N = 1 supergravity models, see e.g. [1–7]. Actually the observation of a connection between spontaneously broken local supersymmetry and a positive cosmological constant is not new and, in fact, goes back to two seminal works published in 1977 [8, 9]. In one of them, Freedman [8] derived the locally supersymmetric extension of the FI term by gauging the R-symmetry. Eliminating the auxiliary field \( \tilde{N} \) yields a cosmological term with the cosmological constant being positive and proportional to \( \xi^2 \), where \( \xi \) is the parameter appearing in the rigid supersymmetric FI action [10]

\[
S = \frac{1}{2} \int d^4 x d^2 \theta W^a W_a - 2 \xi \int d^4 x d^2 \theta d^2 \Phi V,
\]

with \( W_a \) the chiral gauge-invariant field strength [11]. In the other work, Deser and Zumino [9] elaborated on the super-Higgs effect [12] and demonstrated that the spontaneous breaking of local supersymmetry is accompanied by the appearance of a positive contribution to the cosmological constant proportional to the square of the scale of supersymmetry breaking. The complete cosmological constant in [9] also includes a negative contribution coming from the supersymmetric cosmological term discovered by Townsend [13]. If a U(1) vector multiplet theory with the FI term is consistently coupled to supergravity, no supersymmetric cosmological term is permissible [14, 15], as in the case of [8].

It should be mentioned that both 1977 papers [8, 9] discussed above made use of N = 1 supergravity without auxiliary fields—off-shell supergravity simply did not exist at the time. However, within a year the main ideas and results of [8, 9] were extended to off-shell supergravity. In the framework of the old minimal formulation for N = 1 supergravity [16–18], the modern description of the FI term was given by Stelle and West [14]. The supersymmetric cosmological term in old minimal supergravity was constructed in [21, 22]. A complete off-shell extension of the construction attempted by Deser and Zumino [9] was given by Lindström and Roček [23] who proposed the first off-shell model for pure de Sitter supergravity in four dimensions. They coupled a nilpotent covariantly chi-

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1 For the new minimal formulation [19] the FI term was described in [20].

2 Since Deser and Zumino [9] made use of supergravity without auxiliary fields, it was next to impossible to construct a complete supergravity-Goldstino action in their setting.
r al scalar to old minimal supergravity, with a super-symmetric cosmological term included. They reduced the theory to components and computed the same cosmological term as in [9] and also in [3–7].

We see that some of the results obtained in recent publications [1–7] were derived for the first time in the late 1970s. Unfortunately, back in the 1970s nobody was interested in generating a positive cosmological constant. Everyone wanted it to vanish.

The FI terms can consistently be defined only for a restricted class of supergravity–matter theories. It is appropriate here to include a quote from the work [15] which identified such dynamical systems: “in order for a U(1) gauge theory with a Fayet–Iliopoulos term to be consistently coupled to supergravity, preserving gauge invariance, the superpotential must be R invariant. A supersymmetric cosmological term and therefore an explicit mass-like term for the gravitino is forbidden by gauge invariance”. It can be shown [24] that all such theories can be realised within the new minimal auxiliary field formulation (and hence they possess dual realisations in old minimal supergravity). Since these supergravity-matter systems are not extremely attractive for phenomenological applications, one might ask the following question: Could there exist a locally supersymmetric generalisation of, or an alternative to, the standard FI term that would be free of the limitations of the latter? This important question was posed for the first time shortly before Christmas 2017 by Cribiori, Farakos, Tournoy and Van Proeyen [25] who also provided a positive answer below) and

The constructions presented in [26, 27] are based on the assumption that local supersymmetry is in a spontaneously broken phase ab initio. In terms of a massless vector multiplet described by a gauge prepotential $V$, this assumption is simply the requirement that the real gauge-invariant descendant $\bar{D} W \equiv D^a W_a = \bar{T}_a \bar{W}^a = \bar{D} \bar{W}$ is nowhere vanishing, i.e. $(\bar{D} W)^{-1}$ exists. As usual, $W_a$ denotes the covariantly chiral field strength [16]

$$ W_a : = - \frac{1}{4} (\bar{D}^2 - 4R) \bar{T}_a V , \quad \bar{D}_b W_a = 0 , \quad (2.1) $$

which is invariant under gauge transformations of the form

$$ \delta_b V = \lambda + \bar{\lambda} , \quad \bar{D}_a \lambda = 0 . \quad (2.2) $$

The gauge prepotential is chosen to be super-Weyl inert, $\delta_b V = 0^4$.

In the phase with unbroken supersymmetry, there exists a unique gauge-invariant action for the vector multiplet coupled to conformal supergravity [16]

$$ S_{\text{Maxwell}}[V] = \frac{1}{8} \int d^4x d^2 \theta \delta E V D_a \alpha (2.3) $$

$$ \times (\bar{D}^2 - 4R) \bar{T}_a V = \frac{1}{2} \int d^4x d^2 \theta \bar{E} W^2 , \quad (2.4) $$

$$ \times \mathcal{H} \left( \frac{\bar{D}^2 W^2}{(\bar{D} W)^2} - \frac{\bar{D}^2 \bar{W}^2}{(\bar{D} \bar{W})^2} \right) , $$

where $W^2 : = W^a \bar{W}_a$ and $\bar{E}$ is the chiral integration measure [22]. In the spontaneously broken phase, however, a more general action can be introduced to describe the dynamics of a self-interacting vector multiplet coupled to conformal supergravity [27]. It has the form

$$ S[V] = \frac{1}{2} \int d^4x d^2 \theta \bar{E} W^2 + \int d^4x d^2 \theta \bar{E} \bar{W} \bar{W}^2 (\bar{D} W)^2 \quad (2.5) $$

$$ \times \mathcal{H} \left( \frac{\bar{D}^2 W^2}{(\bar{D} W)^2} - \frac{\bar{D}^2 \bar{W}^2}{(\bar{D} \bar{W})^2} \right) , $$

where $\mathcal{H}(z, \bar{z})$ is a real function of one complex variable.

We now turn to the generalised FI term introduced in [27]. It is given by

$$ \mathcal{J}[V; \gamma] = \int d^4x d^2 \theta \bar{E} \gamma V^\delta , \quad (2.6) $$

$$ \gamma(\delta) : = \gamma \left( \frac{\bar{D}^2 W^2}{(\bar{D} W)^2} - \frac{\bar{D}^2 \bar{W}^2}{(\bar{D} \bar{W})^2} \right) , $$

where $V$ denotes the following composite gauge-invariant scalar [28]

$$ V : = -4 W^2 \bar{W}^2 (\bar{D} W)^2 , \quad (2.7) $$

and $\gamma(\delta, \bar{z})$ is a real function of one complex variable, which is subject only to the condition $\gamma(1, 1) \neq 0$ (see below) and may also depend on super-Weyl invariant matter superfields. In the right-hand side of (2.5) $\gamma$ is a real scalar with super-Weyl transformation

$$ \delta_b \gamma = (\sigma + \bar{\sigma}) \gamma . \quad (2.8) $$

It is the nowhere vanishing scalar $\gamma$ which encodes information about a specific off-shell supergravity theory. Within the new minimal formulation for $N = 1$ supergravity [19, 20], $\gamma$ can be identified with the corresponding linear compensator $L$.

$$ (\bar{D}^2 - 4R) L = 0 , \quad L = L . \quad (2.9) $$

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3 References [25] and [26] appeared in the arXiv within less than a month of each other.

4 All relevant information concerning the supergravity covariant derivatives $\bar{D}_a = (\bar{D}_a, \bar{D}_a \bar{\sigma} \bar{\delta} \bar{\delta})$ and the super-Weyl transformations (with the super-Weyl parameter $\sigma$ being chiral, $\bar{D}_a \bar{\sigma} = 0$) can be found in [27].
GENERALISED FAYET–ILIPOULOS TERMS IN SUPERGRAVITY

In pure old minimal supergravity [16–18], $\mathcal{Y}$ is given by $\mathcal{Y} = \mathcal{F}_0 S_0$, where $S_0$ is the chiral compensator, $\overline{D}_a S_0 = 0$, with super-Weyl transformation law $\delta_{\alpha} S_0 = \alpha S_0$. In the presence of chiral matter, however, $\mathcal{Y}$ must be deformed, see below. In principle, the super-Weyl transformation law (2.7) is the only condition on $\mathcal{Y}$. This means that the generalised FI term (2.5) is consistently defined for any off-shell formulation for $\mathcal{N} = 1$ supergravity, unlike the standard FI term in supergravity.

We denote by $\mathcal{J}_{\text{FI}}[\mathcal{V}; \mathcal{Y}]$ the generalised FI term (2.5) corresponding to the homogeneous function $\mathcal{G}(z, \overline{z}) = \eta z^n$, with $\eta$ a real parameter. This one-parameter family of generalised FI terms was introduced in [26]. Its special representative $\mathcal{J}_{\text{FI}}[\mathcal{V}; \mathcal{Y}]$ was proposed for the first time in [25].

The composite superfield $\mathcal{V}^{(6)}$ defined in (2.5) has three important properties. Firstly, it satisfies the nilpotency conditions

$$\mathcal{V}^{(6)} \mathcal{V}^{(6)} = 0, \quad \mathcal{V}^{(6)} \bar{D}_a \bar{D}_b \mathcal{V}^{(6)} = 0, \quad \mathcal{V}^{(6)} \bar{D}_a \bar{D}_b \bar{D}_c \mathcal{V}^{(6)} = 0,$$

which are characteristic of the Goldstino superfields studied in [28, 29]. Secondly, $\mathcal{V}^{(6)}$ is gauge invariant, $\delta_{\alpha} \mathcal{V}^{(6)} = 0$. Thirdly, $\mathcal{V}^{(6)}$ is super-Weyl inert, $\delta_{\sigma} \mathcal{V}^{(6)} = 0$.

The vector multiplet action (2.4) and the generalised FI term (2.5) are two sectors of the complete supergravity-matter action

$$S = S_{\text{SUGRA}} + S[\mathcal{V}] – 2\xi \mathcal{J}_{\text{FI}}[\mathcal{V}; \mathcal{Y}],$$

(2.10)

where $S_{\text{SUGRA}}$ denotes an action for supergravity coupled to all other supermultiplets, for instance

$$S_{\text{SUGRA}} = -3 \int d^4 x d^2 \theta \overline{\theta} E S_0 e^{-K(\phi, \overline{\phi})} S_0 + \left\{ \int d^4 x d^2 \theta \overline{\theta} S_0^3 W(\phi) + \text{c.c.} \right\}.$$

(2.11)

This corresponds to old minimal supergravity, with $\phi$ being matter chiral superfields taking their values in a Kähler manifold, with $K(\phi, \overline{\phi})$ its Kähler potential. In this case $\mathcal{Y}$ in the generalised FI term in (2.10) should be $\mathcal{Y} = \bar{S}_0 e^{-K(\phi, \overline{\phi})} S_0$, as was pointed out in [30, 31], and then our model can be written in a form similar to that with the standard FI term [14],

$$-3 \mathcal{Y} - 2\xi \mathcal{V}^{(6)} = -3 \bar{S}_0 e^{-K(\phi, \overline{\phi})+3\xi \mathcal{V}^{(6)}} S_0.$$

(2.12)

In general, the vector multiplet part of the action (2.10) is highly nonlinear. However its functional form drastically simplifies if the ordinary gauge field contained in $\mathcal{V}$ is a flat connection. Then the gauge freedom (2.2) allows one to choose a gauge in which $\mathcal{V}$ is a nilpotent superfield constrained by

$$\mathcal{V} V = 0, \quad \mathcal{V} \overline{\mathcal{D}}_a \mathcal{V} = 0.$$

(2.13)

In this gauge it can be shown (see also [28]) that

$$\mathcal{V} = -4 \mathcal{W} W^{2} \mathcal{W}^{-2} \Rightarrow \mathcal{V} \mathcal{D}^2 \mathcal{W}^2 = -\mathcal{V}(\mathcal{D} \mathcal{W})^2$$

$$\Rightarrow \mathcal{V}^{(6)} = \mathcal{G}(1,1) \mathcal{V}.$$

(2.14)

These relations imply that the $\mathcal{V}$-dependent part of the action (2.10) becomes

$$S[\mathcal{V}] – 2\xi \mathcal{J}_{\text{FI}}[\mathcal{V}; \mathcal{Y}]$$

$$\Rightarrow \frac{h}{2} \int d^4 x d^2 \theta \overline{\theta} W^2 - 2\xi g \int d^4 x d^2 \theta d^2 \theta E Y V,$$

(2.15)

where we have denoted $h := 1 + \frac{1}{2} \mathcal{H}(1,1)$ and $g := \mathcal{G}(1,1)$. Under mild conditions

$$\mathcal{H}(1,1) > -2, \quad \mathcal{G}(1,1) \neq 0,$$

(2.16)

the functional (2.15) coincides (modulo an overall numerical factor) with the Goldstino multiplet action, $S_{\text{Goldstino}}$, which was proposed in [28] and is given by

$$S_{\text{Goldstino}} = \frac{1}{2} \int d^4 x d^2 \theta d^2 \theta E \mathcal{W},$$

(2.17)

where the Goldstino superfield $\mathcal{V}$ obeys the constraints (2.14) and is subject to the condition that $\mathcal{D} \mathcal{W}$ be nowhere vanishing. It should be noted that the condition $h > 0$ is equivalent to the requirement that the kinetic term of the Goldstino field has the correct sign. At the component level, the action (2.15) is highly nonlinear due to the nilpotency constraints (2.13). However, the functional form of (2.15) is universal, unlike the complete vector multiplet action in (2.10), which is a manifestation of the fact that the Volkov–Akulov action [32] is universal [33]5. We have thus shown that the supergravity-matter system (2.10) describes spontaneously broken local supersymmetry under the condition (2.16).

To arrive at (2.15), we have made use of the gauge fixing (2.13), which exists only if the component gauge field is a flat connection and which expresses the fact

5 The constraints (2.13) are invariant under local rescalings $V \rightarrow e^\rho V$, with the parameter $\rho$ being an arbitrary real scalar superfield. Requiring the action (2.15) to be stationary under such rescalings gives the constraint $f Y V = -\frac{1}{2} \mathcal{V} \mathcal{D}^2 \mathcal{W} = \frac{1}{8} \mathcal{V} \mathcal{D}^2 (\mathcal{D} \mathcal{W})^2 - 4 \mathcal{Y} \mathcal{D} \mathcal{W} V$, where $f = \xi g/h$. In conjunction with (2.13), this constraint defines the irreducible Goldstino multiplet introduced in [29].
that the gauge field is switched off. This is actually possible to avoid by: (i) making any assumption about the component gauge field; and (ii) imposing any gauge condition. In general, the unconstrained gauge prepotential \( V \) may be split in two superfields, one of which contains the \( U(1,1) \) gauge field and all purely gauge degrees of freedom, while the other contains the remaining physical component fields. The point is that the nilpotency conditions (2.9) allow us to interpret \( V^{(6)} \) as a Goldstino superfield of the type proposed in [28] provided its \( \{-\} \)-field is nowhere vanishing, which means that \( \bar{Q}^2 W^2 \) is nowhere vanishing, in addition to the condition \( \bar{Q} W \neq 0 \) imposed earlier. Then \( V^{(6)} \) contains only two independent component fields, the Goldstino and \( D \)-field. We then can introduce a new parametrisation for the gauge prepotential given by

\[
V = A^{(6)} + V^{(6)}. \tag{2.18}
\]

It is \( A^{(6)} \) which varies under the gauge transformation (2.2), \( \delta_{\chi} A^{(6)} = \lambda + \bar{\lambda} \), while \( V^{(6)} \) is gauge invariant by construction. Modulo purely gauge degrees of freedom, \( A^{(6)} \) contains a single independent physical field, the gauge one-form.

Our discussion shows that a flat-superspace limit of the vector multiplet action \( S[V] = 2 \bar{\epsilon}_{\beta} \bar{Q}^{(8)}[V; \Upsilon] \) in (2.10) gives a consistent model for spontaneously broken rigid supersymmetry under the conditions (2.16).

3. THE COMPONENT STRUCTURE

It is important to analyse the bosonic Lagrangian of the model (2.10) in the vector multiplet sector. For this purpose we introduce gauge-invariant component fields of the vector multiplet following [34]

\[
\bar{W}_{\alpha} = \chi_{\alpha}, \quad -\frac{1}{2} \bar{D}^\alpha W_{\alpha} = D, \tag{3.1}
\]

\[
D_{\alpha} \bar{W}_{\beta} = 2 i \hat{F}_{\alpha \beta} = i (\sigma^{ab})_{\alpha \beta} \hat{F}_{\alpha \beta},
\]

where the bar-projection \( U \) of a superfield \( U \) means switching off the superspace Grassmann variables, and

\[
\hat{F}_{\alpha \beta} = F_{\alpha \beta} - \frac{1}{2} (\Psi_{\alpha} \sigma_{\beta} \bar{\chi} + \chi_{\beta} \sigma_{\alpha} \bar{\Psi}_{\alpha}) + \frac{1}{2} (\Psi_{\beta} \sigma_{\alpha} \bar{\chi} + \chi_{\alpha} \sigma_{\beta} \bar{\Psi}_{\beta}),
\]

\[
F_{\alpha \beta} = \nabla_{\alpha} V_{\beta} - \nabla_{\beta} V_{\alpha} - \bar{T}_{\alpha \beta} V_{\gamma}, \tag{3.2}
\]

with \( V_{\alpha} = e_{a}^{\alpha}(x) V_{m}(x) \) the gauge one-form, and \( \Psi_{\beta}^a \) the gravitino. Here \( V_{\alpha} \) denotes a Lorentz-covariant derivative with torsion,

\[
[\nabla_{\alpha}, \nabla_{\beta}] = \bar{T}_{\alpha \beta} V_{\gamma} + \frac{1}{2} R_{\alpha \beta \gamma} M^{\alpha \beta}, \tag{3.3}
\]

where \( T_{\alpha \beta \gamma} \) and \( R_{\alpha \beta \gamma} \) are the torsion and curvature tensors, respectively. The former is determined by the gravitino via \( T_{\alpha \beta \gamma} = -i (\psi_{\alpha} \sigma_{\beta} \bar{\Psi}_{\gamma} - \psi_{\beta} \sigma_{\alpha} \bar{\Psi}_{\gamma}) \). Making use of the above relations leads to

\[
-\frac{1}{4} \bar{Q}^2 W^2 = D^2 - 2 F^2 + \text{fermionic terms}, \tag{3.4}
\]

\[
F^2 := F_{\alpha \beta} F_{\alpha \beta}.
\]

We conclude that the electromagnetic field should be weak enough to satisfy \( D^2 - 2 F^2 \neq 0 \), in addition to the condition \( D \neq 0 \) discussed above. Direct calculations yield the component bosonic Lagrangian

\[
\mathcal{L}(F_{\alpha \beta}, D) = -\frac{1}{2} (F^2 + F^2) + \frac{1}{2} D^2 \times \left[ 1 + \frac{1}{2} (1 - 2 F^2 D^2, 1 - 2 F^2 D^2) \right] - 2 F^2 - 2 F^2 \Upsilon.
\]

In order for the supergravity action in (2.10) to give the correct Einstein–Hilbert gravitational Lagrangian at the component level, the super-Weyl gauge \( \Upsilon = 1 \) should be imposed, see [34, 35] for the technical details.

Let us briefly discuss the structure of the Lagrangian (3.5). It is seen that the case [25]

\[
\mathcal{G}(z, \bar{z}) = \frac{g}{z \bar{z}}, \tag{3.6}
\]

is special since the last term in (3.5) becomes linear in \( D \) and independent of the field strength. Another special case corresponds to

\[
\mathcal{H}(z, \bar{z}) = \frac{2 h}{z \bar{z}}, \tag{3.7}
\]

since the second term in (3.5) becomes quadratic in \( D \) and independent of \( F \). The following table summarises the differences between the standard and generalised FI terms:

| FI term                              | \( J_{FI} \) |
|--------------------------------------|--------------|
| Standard FI term [10]                | \( J_{FI} = \frac{1}{2} D \) |
| Special FI-type term [25]            | \( J_{FI}^{(6)} = \frac{1}{2} D + O(\chi, F) \) |
| General FI-type term [26, 27]        | \( J_{FI}^{(6)} = \frac{1}{2} D + O(\chi, F) \) |
Making use of the local supersymmetry allows one to impose the unitary gauge \( \chi_a = 0 \), which is why little is lost in restricting our analysis to the bosonic sector (3.5). Imposing the gauge \( \chi_a = 0 \) red [10] uces \( J[^{-1}] \) to the standard FI term, \( \frac{1}{2} D \).

As follows from (3.5), the auxiliary field \( D \) may be integrated out using its equation of motion

\[
\frac{\partial}{\partial D} \mathcal{L}(F_{ab}, D) = 0,
\]

resulting in a model for nonlinear electrodynamics, \( \mathcal{L}(F_{ab}) \). The special feature of the choice (3.6) and (3.7) is that \( \mathcal{L}(F_{ab}) \) coincides with Maxwell’s Lagrangian.

4. GENERALISED FI TERMS AS QUANTUM CORRECTIONS

In Section 1, it was pointed out that all healthy mechanisms to generate a cosmological constant should explain both (i) its positivity and (ii) its small value. So far we have only discussed point (i). Here a comment will be made regarding point (ii).

It is known that the standard FI term in rigid supersymmetric gauge theories does not receive any quantum corrections [36, 37]. In other words, there is no way to generate an FI term quantum mechanically. It appears that generalised FI terms may occur as pure quantum corrections. This hope is supported by the fact that the nonlinear term in the superconformal action (2.4) has a functional form typical for quantum corrections to low-energy effective actions in SYM theories, see e.g. [38, 39]. However, the important difference between the functionals (2.4) and (2.5) is that the former is even with respect to the replacement \( V \rightarrow -V \), while the latter is odd (similar to the anomalous effective action in supersymmetric gauge theories). Explicit calculations will be reported elsewhere.

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