



Taking a vector supermultiplet apart: Alternative Fayet–Iliopoulos-type terms



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ABSTRACT

Starting from an Abelian $\mathcal{N} = 1$ vector supermultiplet V coupled to conformal supergravity, we construct from it a nilpotent real scalar Goldstino superfield \mathfrak{V} of the type proposed in arXiv:1702.02423. It contains only two independent component fields, the Goldstino and the auxiliary D -field. The important properties of this Goldstino superfield are: (i) it is gauge invariant; and (ii) it is super-Weyl invariant. As a result, the gauge prepotential can be represented as $V = \mathcal{V} + \mathfrak{V}$, where \mathcal{V} contains only one independent component field, modulo gauge degrees of freedom, which is the gauge one-form. Making use of \mathfrak{V} allows us to introduce new Fayet–Iliopoulos-type terms, which differ from the one proposed in arXiv:1712.08601 and share with the latter the property that gauged R -symmetry is not required.

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1. Introduction

In quantum field theory with a symmetry group G spontaneously broken to its subgroup H , the multiplet of matter fields transforming according to a linear representation of G can be split into two subsets: (i) the massless Goldstone fields; and (ii) the other fields that are massive in general. Each subset transforms nonlinearly with respect to G and linearly under H . Each subset may be realised in terms of constrained fields transforming linearly under G [1,2]. In the case of spontaneously broken supersymmetry [3], every superfield U containing the Goldstino may be split into two supermultiplets, one of which is an irreducible Goldstino superfield¹ and the other contains the remaining component fields [4], in accordance with the general relation between linear and nonlinear realisations of $\mathcal{N} = 1$ supersymmetry [5]. It is worth recalling the example worked out in [4]. Consider the irreducible chiral Goldstino superfield \mathcal{X} , $\bar{D}_{\dot{\alpha}}\mathcal{X} = 0$, introduced in [5,6]. It is defined to obey the constraints [6]

$$\mathcal{X}^2 = 0, \quad f\mathcal{X} = -\frac{1}{4}\mathcal{X}\bar{D}^2\bar{\mathcal{X}}, \quad (1.1)$$

where f is a real parameter characterising the scale of supersymmetry breaking. As U we choose the reducible chiral Goldstino superfield X , $\bar{D}_{\dot{\alpha}}X = 0$, proposed in [7,8]. It is subject only to the constraint

$$X^2 = 0. \quad (1.2)$$

It was shown in [4] that X can be represented in the form

$$X = \mathcal{X} + \mathcal{Y}, \quad f\mathcal{X} := -\frac{1}{4}\bar{D}^2(\bar{\Sigma}\Sigma), \quad \Sigma := -4f\frac{\bar{X}}{\bar{D}^2\bar{X}}, \quad (1.3)$$

where the auxiliary field F of X is the only independent component of the chiral scalar \mathcal{Y} . Originally, the irreducible Goldstino superfield Σ was introduced in [9] to be a modified complex linear superfield, $-\frac{1}{4}\bar{D}^2\Sigma = f$, which is nilpotent and obeys a holomorphic nonlinear constraint,

$$\Sigma^2 = 0, \quad fD_{\alpha}\Sigma = -\frac{1}{4}\Sigma\bar{D}^2D_{\alpha}\Sigma. \quad (1.4)$$

These properties follow from (1.3).

The approach advocated in [4] may be pursued one step further with the goal to split any unconstrained superfield U into two supermultiplets, one of which is a reducible Goldstino supermultiplet. This has been implemented in [10] for the reducible chiral Goldstino superfield X . There exist two other reducible Goldstino superfields: (i) the three-form variant of X [11,12]; and (ii) the nilpotent real scalar superfield introduced in [13]. In the present

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¹ The notion of irreducible and reducible Goldstino superfields was introduced in [4]. For every irreducible Goldstino superfield, the Goldstino is its only independent component. Reducible Goldstino superfields also contain auxiliary field(s) in addition to the Goldstino.

paper we make use of (ii) in order to split a $U(1)$ vector supermultiplet into two constrained superfields. Our construction makes it possible to introduce new Fayet–Iliopoulos-type terms, which differ from the one recently proposed in [14] and share with the latter the property that gauged R -symmetry is not required.

In this paper, we make use of the simplest formulation for $\mathcal{N} = 1$ conformal supergravity in terms of the superspace geometry of [15], which underlies the Wess–Zumino approach [16] to old minimal supergravity [17,18]. This approach requires the super-Weyl transformations of [19] (defined in the appendix) to belong to the supergravity gauge group. Our notation and conventions follow [20].

2. Constructing a Goldstino superfield

Consider a massless vector supermultiplet in a conformal supergravity background. It is described by a real scalar prepotential V defined modulo gauge transformations

$$\delta_\lambda V = \lambda + \bar{\lambda}, \quad \bar{D}_{\dot{\alpha}} \lambda = 0. \quad (2.1)$$

As usual, the prepotential is chosen to be super-Weyl inert, $\delta_\sigma V = 0$. In what follows, we assume that the top component (D -field) of V is nowhere vanishing. In terms of the gauge-invariant field strength [16]

$$W_\alpha := -\frac{1}{4}(\bar{D}^2 - 4R)\mathcal{D}_\alpha V, \quad \bar{D}_{\dot{\beta}} W_\alpha = 0, \quad (2.2)$$

our assumption means that the real scalar $DW := \mathcal{D}^\alpha W_\alpha = \mathcal{D}_{\dot{\alpha}} W^{\dot{\alpha}}$ is nowhere vanishing.

It is instructive to consider a simple supersymmetric gauge theory in which the above assumption is compatible with the equations of motion. Within the new minimal formulation for $\mathcal{N} = 1$ supergravity [21,22], the dynamics of the massless vector supermultiplet with a Fayet–Iliopoulos (FI) term [23] is governed by the gauge invariant and super-Weyl invariant action (see, e.g., [24])

$$S[V] = \int d^4x d^2\theta d^2\bar{\theta} E \left\{ \frac{1}{16} V \mathcal{D}^\alpha (\bar{D}^2 - 4R) \mathcal{D}_\alpha V - 2fLV \right\}, \quad (2.3)$$

where L is the conformal compensator for new minimal supergravity [25] (and as such L is nowhere vanishing). It is a real covariantly linear scalar superfield,

$$(\bar{D}^2 - 4R)L = 0, \quad \bar{L} = L, \quad (2.4)$$

with the super-Weyl transformation $\delta_\sigma L = (\sigma + \bar{\sigma})L$. The second term in the action is the FI term, with f a real parameter. The equation of motion for V is $DW = -4fL$, and it implies that DW is indeed nowhere vanishing.

Since DW is nowhere vanishing, we can introduce (as an extension of the construction in section 5.2 of [13]) the following scalar superfield

$$\mathfrak{W} := -4 \frac{W^2 \bar{W}^2}{(DW)^3}, \quad W^2 := W^\alpha W_\alpha. \quad (2.5)$$

This superfield is gauge invariant, $\delta_\lambda \mathfrak{W} = 0$, and super-Weyl invariant,

$$\delta_\sigma \mathfrak{W} = 0, \quad (2.6)$$

as follows from the super-Weyl transformation laws of W_α and DW :

$$\delta_\sigma W_\alpha = \frac{3}{2} \sigma W_\alpha, \quad \delta_\sigma DW = (\sigma + \bar{\sigma})DW. \quad (2.7)$$

By construction, it obeys the following nilpotency conditions

$$\mathfrak{W}^2 = 0, \quad (2.8a)$$

$$\mathfrak{W} \mathcal{D}_A \mathcal{D}_B \mathfrak{W} = 0, \quad (2.8b)$$

$$\mathfrak{W} \mathcal{D}_A \mathcal{D}_B \mathcal{D}_C \mathfrak{W} = 0, \quad (2.8c)$$

which mean that \mathfrak{W} is the Goldstino superfield introduced in [13].² Associated with \mathfrak{W} is the covariantly chiral spinor \mathfrak{W}_α which is obtained from (2.2) by replacing V with \mathfrak{W} . As shown in [13], the constraints (2.8) imply that

$$\mathfrak{W} := -4 \frac{\mathfrak{W}^2 \bar{\mathfrak{W}}^2}{(D\mathfrak{W})^3}. \quad (2.9)$$

Choosing $V = \mathfrak{W}$ in (2.3) gives the Goldstino superfield action proposed in [13].

In order for our interpretation of \mathfrak{W} as a Goldstino superfield to be consistent, its D -field should be nowhere vanishing, which is equivalent to the requirement that $\mathcal{D}\mathfrak{W}$ be nowhere vanishing. As follows from (2.5), this condition implies that $\mathcal{D}^2 W^2$ is nowhere vanishing. To understand what the latter implies, let us introduce the component fields of the vector supermultiplet following [27]

$$W_\alpha| = \psi_\alpha, \quad -\frac{1}{2} \mathcal{D}^\alpha W_\alpha| = D, \quad (2.10)$$

$$\mathcal{D}_{(\alpha} W_{\beta)}| = 2i \hat{F}_{\alpha\beta} = i(\sigma^{ab})_{\alpha\beta} \hat{F}_{ab},$$

where the bar-projection, $U|$, means switching off the superspace Grassmann variables, and

$$\hat{F}_{ab} = F_{ab} - \frac{1}{2}(\Psi_a \sigma_b \bar{\psi} + \psi \sigma_b \bar{\Psi}_a) + \frac{1}{2}(\Psi_b \sigma_a \bar{\psi} + \psi \sigma_a \bar{\Psi}_b), \quad (2.11)$$

$$F_{ab} = \nabla_a V_b - \nabla_b V_a - \mathcal{T}_{ab}{}^c V_c,$$

with $V_a = e_a{}^m(x) V_m(x)$ the gauge one-form, and $\Psi_a{}^\beta$ the gravitino. The operator ∇_a denotes a spacetime covariant derivative with torsion,

$$[\nabla_a, \nabla_b] = \mathcal{T}_{ab}{}^c \nabla_c + \frac{1}{2} \mathcal{R}_{abcd} M^{cd}, \quad (2.12)$$

where \mathcal{R}_{abcd} is the curvature tensor and \mathcal{T}_{abc} is the torsion tensor. The latter is related to the gravitino by

$$\mathcal{T}_{abc} = -\frac{i}{2}(\Psi_a \sigma_c \bar{\Psi}_b - \Psi_b \sigma_c \bar{\Psi}_a). \quad (2.13)$$

For more details, see [20,27]. We deduce from the above relations that

$$-\frac{1}{4} \mathcal{D}^2 W^2| = D^2 - 2F^{\alpha\beta} F_{\alpha\beta} + \text{fermionic terms}. \quad (2.14)$$

We conclude that the electromagnetic field should be weak enough to satisfy

$$D^2 - 2F^{\alpha\beta} F_{\alpha\beta} \neq 0, \quad (2.15)$$

in addition to the condition $D \neq 0$. The D -field of \mathfrak{W} is

$$-\frac{1}{2} \mathcal{D}\mathfrak{W}| = D \left| 1 - 2 \frac{F^{\alpha\beta} F_{\alpha\beta}}{D^2} \right|^2 + \text{fermionic terms}. \quad (2.16)$$

Making use of the Goldstino superfield \mathfrak{W} leads to a new parametrisation for the gauge prepotential given by

² The Goldstino superfield constrained by (2.8) contains only two independent fields, the Goldstino and the auxiliary D -field. This can be shown by analogy with the $\mathcal{N} = 2$ analysis in three dimensions [26].

$$V = \mathcal{V} + \mathfrak{V} . \tag{2.17}$$

It is \mathcal{V} which varies under the gauge transformation (2.1), $\delta_\lambda \mathcal{V} = \lambda + \bar{\lambda}$, while \mathfrak{V} is gauge invariant by construction. Modulo purely gauge degrees of freedom, \mathcal{V} contains only one independent field, which is the gauge one-form.

There exists a different way to construct a reducible Goldstino superfield in terms of V , which is given by

$$\hat{\mathfrak{V}} := -4 \frac{W^2 \bar{W}^2}{\mathcal{D}^2 W^2 \bar{\mathcal{D}}^2 \bar{W}^2} \mathcal{D}W . \tag{2.18}$$

Unlike \mathfrak{V} defined by (2.5), this gauge-invariant superfield is not manifestly super-Weyl invariant. Nevertheless, it proves to be invariant under the super-Weyl transformations,

$$\delta_\sigma \hat{\mathfrak{V}} = 0 , \tag{2.19}$$

as follows from the observation [28] (see also [27]) that

$$\left(\mathcal{D}^2 - 4\bar{R} \right) \frac{W^2}{\Upsilon^2} \tag{2.20}$$

is super-Weyl invariant for any compensating (nowhere vanishing) real scalar Υ with the super-Weyl transformation law

$$\delta_\sigma \Upsilon = (\sigma + \bar{\sigma}) \Upsilon . \tag{2.21}$$

In the new minimal supergravity, we can identify

$$\Upsilon = L . \tag{2.22a}$$

In the case of the old minimal formulation for $\mathcal{N} = 1$ supergravity [16–18], we choose

$$\Upsilon = \bar{\Phi} \Phi , \quad \bar{\mathcal{D}}_{\dot{\alpha}} \Phi = 0 , \tag{2.22b}$$

where the chiral compensator Φ has the super-Weyl transformation law $\delta_\sigma \Phi = \sigma \Phi$.

By construction, the superfield (2.18) obeys the nilpotency conditions (2.8). It may be shown that the composites (2.5) and (2.18) coincide if V is chosen to be \mathfrak{V} or $\hat{\mathfrak{V}}$. Thus the two Goldstino superfields \mathfrak{V} or $\hat{\mathfrak{V}}$ differ only in the presence of a gauge field. It follows from the definition (2.18) that $\hat{\mathfrak{V}}$ is well defined provided the condition (2.15) holds. The same definition tells us that the D -field of $\hat{\mathfrak{V}}$ is equal to

$$-\frac{1}{2} \mathcal{D} \hat{\mathfrak{V}} | = D + \text{fermionic terms} . \tag{2.23}$$

The composite (2.18) was introduced in a recent paper [14]. The authors of [14] put forward the supersymmetric invariant

$$\hat{\mathcal{J}}_{\text{FI}} = \int d^4x d^2\theta d^2\bar{\theta} E \Upsilon \hat{\mathfrak{V}} \tag{2.24}$$

as a novel FI term that does not require gauged R -symmetry. The compensating superfield Υ was chosen in [14] to be the old minimal expression (2.22b).

We propose an alternative FI-type invariant

$$\mathcal{J}_{\text{FI}} = \int d^4x d^2\theta d^2\bar{\theta} E \Upsilon \mathfrak{V} . \tag{2.25}$$

It also does not require gauged R -symmetry. In addition, it does not require (2.15).

Actually, the above constructions can be generalised by introducing a gauge-invariant Goldstino superfield of the form

$$\mathfrak{V}_n := \mathfrak{V} \frac{(\mathcal{D}W)^{4n}}{[\mathcal{D}^2 W^2 \bar{\mathcal{D}}^2 \bar{W}^2]^n} , \tag{2.26}$$

for some integer n . The superfield (2.18) correspond to $n = 1$. It is obvious that \mathfrak{V}_n obeys the constraints (2.8). Moreover, \mathfrak{V}_n is super-Weyl invariant. The superfields \mathfrak{V} and \mathfrak{V}_n coincide if the gauge prepotential V is chosen to be \mathfrak{V} . New FI-type invariants are obtained by making use of \mathfrak{V}_n instead of \mathfrak{V} in (2.25),

$$\hat{\mathcal{J}}_{\text{FI}}^{(n)} = \int d^4x d^2\theta d^2\bar{\theta} E \Upsilon \mathfrak{V}_n . \tag{2.27}$$

3. U(1) duality invariant models and BI-type terms

Ref. [27] presented a general family of U(1) duality invariant models for a massless vector supermultiplet coupled to off-shell supergravity, old minimal or new minimal. Such a theory is described by a super-Weyl invariant action of the form

$$S_{\text{SDVM}}[V; \Upsilon] = \frac{1}{2} \int d^4x d^2\theta \mathcal{E} W^2 + \frac{1}{4} \int d^4x d^2\theta d^2\bar{\theta} E \frac{W^2 \bar{W}^2}{\Upsilon^2} \Lambda \left(\frac{\omega}{\Upsilon^2}, \frac{\bar{\omega}}{\Upsilon^2} \right) . \tag{3.1}$$

Here \mathcal{E} is the chiral density, $\omega := \frac{1}{8} \mathcal{D}^2 W^2$, and $\Lambda(\omega, \bar{\omega})$ is areal analytic function satisfying the differential equation [29,30]

$$\text{Im} \left\{ \Gamma - \bar{\omega} \Gamma^2 \right\} = 0 , \quad \Gamma := \frac{\partial(\omega \Lambda)}{\partial \omega} . \tag{3.2}$$

These self-dual dynamical systems are curved-superspace extensions of the globally supersymmetric systems introduced in [29, 30]. The curved superspace extension [28], $S_{\text{SBI}}[V]$, of the supersymmetric Born-Infeld action [28,31] corresponds to the choice

$$\Lambda(\omega, \bar{\omega}) = \frac{g^2}{1 + \frac{1}{2} A + \sqrt{1 + A + \frac{1}{4} B^2}} , \quad A = g^2(\omega + \bar{\omega}) , \tag{3.3}$$

$$B = g^2(\omega - \bar{\omega}) ,$$

with g a coupling constant.

In flat superspace (which, in particular, corresponds to $\Upsilon = 1$), the fermionic sector of (3.1) was shown [27] to possess quite remarkable properties. Specifically, only under the additional restriction

$$\Lambda_{u\bar{u}}(0, 0) = 3\Lambda^3(0, 0) , \tag{3.4}$$

the component fermionic action proves to coincide, modulo a non-linear field redefinition, with the Volkov-Akulov action [3]. This ubiquitous appearance of the Volkov-Akulov action in such models was explained in [32]. If the FI term is added to the flat-superspace counterpart of (3.1), then the auxiliary scalar D develops a non-vanishing expectation value, in general, for its algebraic equation of motion has a non-zero solution. As a result, the supersymmetry becomes spontaneously broken, and thus the photino action should be related to the Goldstino action, due to the uniqueness of the latter.

In supergravity, the situation is analogous to the rigid supersymmetric case. Let us add a standard FI term to the vector multiplet action (3.1) coupled to new minimal supergravity,

$$S = \frac{3}{\kappa^2} \int d^4x d^2\theta d^2\bar{\theta} E L \text{In}L + S_{\text{SDVM}}[V; L] - 2f \int d^4x d^2\theta d^2\bar{\theta} E L V . \tag{3.5}$$

Here the first term is the supergravity action. In general, this system describes spontaneously broken supergravity. It suffices to

consider the case of vanishing gauge field, which corresponds to $V = \mathfrak{W}$. Using the nilpotency conditions (2.8) and relation (2.9), one may show that

$$\begin{aligned} S_{\text{SDVM}}[\mathfrak{W}; L] &= \frac{1}{2} \int d^4x d^2\theta \mathcal{E} \mathfrak{W}^2 \\ &+ \frac{1}{4} \int d^4x d^2\theta d^2\bar{\theta} E \frac{\mathfrak{W}(\mathcal{D}\mathfrak{W})^3}{L^2} \Lambda\left(\frac{\zeta}{L^2}, \frac{\zeta}{L^2}\right), \end{aligned} \quad (\text{A.3.6})$$

where $\zeta = -\frac{1}{8}(\mathcal{D}\mathfrak{W})^2$. The auxiliary field may be eliminated by requiring the functional

$$S_{\text{SDVM}}[\mathfrak{W}; L] - 2f \int d^4x d^2\theta d^2\bar{\theta} E L \mathfrak{W} \quad (\text{A.3.7})$$

to be stationary under local rescalings $\delta\mathfrak{W} = \rho\mathfrak{W}$, with ρ an arbitrary real superfield (compare with [13]). The resulting algebraic equation proves to coincide with the one derived in [32].

An important property of the standard FI term, which was pointed out in [33], is that it remains invariant under the second nonlinearly realised supersymmetry of the rigid supersymmetric Born–Infeld action [31]. This property implies the supersymmetric Born–Infeld action deformed by a FI term still describes partial $\mathcal{N} = 2 \rightarrow \mathcal{N} = 1$ supersymmetry breaking [32,34]. As for the novel FI-type terms (2.24) and (2.25), they do not appear to share this fundamental property.

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Appendix A. Super-Weyl transformations

It was first realised by Howe and Tucker [19] that the Grimm–Wess–Zumino algebra of covariant derivatives [15] is invariant under super-Weyl transformations of the form

$$\delta_\sigma \mathcal{D}_\alpha = (\bar{\sigma} - \frac{1}{2}\sigma)\mathcal{D}_\alpha + \mathcal{D}^\beta \sigma M_{\alpha\beta}, \quad (\text{A.1a})$$

$$\delta_\sigma \bar{\mathcal{D}}_{\dot{\alpha}} = (\sigma - \frac{1}{2}\bar{\sigma})\bar{\mathcal{D}}_{\dot{\alpha}} + (\bar{\mathcal{D}}^{\dot{\beta}}\bar{\sigma})\bar{M}_{\dot{\alpha}\dot{\beta}}, \quad (\text{A.1b})$$

$$\begin{aligned} \delta_\sigma \mathcal{D}_{\alpha\dot{\alpha}} &= \frac{1}{2}(\sigma + \bar{\sigma})\mathcal{D}_{\alpha\dot{\alpha}} + \frac{i}{2}\bar{\mathcal{D}}_{\dot{\alpha}}\bar{\sigma}\mathcal{D}_\alpha + \frac{i}{2}\mathcal{D}_\alpha\sigma\bar{\mathcal{D}}_{\dot{\alpha}} + \mathcal{D}^\beta_\alpha\sigma M_{\alpha\beta} \\ &+ \mathcal{D}_\alpha^{\dot{\beta}}\bar{\sigma}\bar{M}_{\dot{\alpha}\dot{\beta}}, \end{aligned} \quad (\text{A.1c})$$

accompanied by the following transformations of the torsion superfields

$$\delta_\sigma R = 2\sigma R + \frac{1}{4}(\bar{\mathcal{D}}^2 - 4R)\bar{\sigma}, \quad (\text{A.2a})$$

$$\delta_\sigma G_{\alpha\dot{\alpha}} = \frac{1}{2}(\sigma + \bar{\sigma})G_{\alpha\dot{\alpha}} + i\mathcal{D}_{\alpha\dot{\alpha}}(\sigma - \bar{\sigma}), \quad (\text{A.2b})$$

$$\delta_\sigma W_{\alpha\beta\gamma} = \frac{3}{2}\sigma W_{\alpha\beta\gamma}. \quad (\text{A.2c})$$

Here the super-Weyl parameter σ is a covariantly chiral scalar superfield, $\bar{\mathcal{D}}_{\dot{\alpha}}\sigma = 0$.

A tensor superfield \mathcal{T} (with its indices suppressed) is said to be super-Weyl primary of weight (p, q) if its super-Weyl transformation law is

$$\delta_\sigma \mathcal{T} = (p\sigma + q\bar{\sigma})\mathcal{T}, \quad (\text{A.3})$$

for some parameters p and q .

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