Generating higher-derivative couplings in $\mathcal{N} = 2$ supergravity

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Using a recently developed off-shell formulation for general 4D $\mathcal{N} = 2$ supergravity-matter systems, we propose a construction to generate higher derivative couplings. We address here mainly the interactions of tensor and vector multiplets, but the construction is quite general. For a certain subclass of terms, the action is naturally written as an integral over 3/4 of the Grassmann coordinates of superspace.

1 Introduction

One of the advantages of superspace is that it makes supersymmetry manifest. Any action which is invariant under arbitrary diffeomorphisms of curved superspace automatically yields a locally supersymmetric component action. In addition, it offers the possibility to construct general couplings from Lagrangians of (essentially) arbitrary functional form. For these reasons, higher derivative supersymmetric actions, especially those coupled to supergravity, can be quite efficiently constructed in superspace. Such actions have been of interest recently [1, 2].

We review below a new class of higher derivative actions that was constructed directly in superspace in our recent paper [3]. The supergeometry which enables this construction is an off-shell formulation for general $\mathcal{N} = 2$ supergravity-matter couplings in four dimensions [4], which allows a curved-space extension of $\mathcal{N} = 2$ projective superspace [5]. For our purposes, the relevant detail is that the required superspace is $\mathcal{M}^4|8 \times \mathbb{C}P^1$. The curved supermanifold $\mathcal{M}^4|8$ is parametrized by local coordinates $z^M = (x^m, \theta^i, \bar{\theta}_{\dot{i}})$, $i = 1, 2$. Its geometry is described by covariant derivatives

$$D_A = E_A + \frac{1}{2} \Omega_A^{bc} M_{bc} + \Phi_A^{jk} J_{jk}$$

where $E_A = E_A^M \partial_M$ is the supervielbein; $M_{bc}$ and $\Omega_A^{bc}$ are the Lorentz generators and superconnections respectively; and $J_{jk}$ and $\Phi_A^{jk}$ are respectively the SU(2) generators and superconnections. The auxiliary manifold $\mathbb{C}P^1$ is parametrized by an isotwistor $v^i \in \mathbb{C}^2 \setminus 0$ defined modulo $v^i \sim cv^i$ for $c \in \mathbb{C} \setminus \{0\}$. All superfields and operators are required to have fixed homogeneity in $v^i$. For a more extensive discussion of the supergeometry, we refer the reader to the original references [4].

Our higher derivative construction is based on a duality between two basic off-shell representations of $\mathcal{N} = 2$ supersymmetry – the tensor multiplet and the vector multiplet. In curved superspace, the tensor multiplet is described by its field strength $G_{ij}$, which is a real isotriplet superfield, $\epsilon_{ij} G_{kl} = 0$. These conditions are solved in terms of an unconstrained chiral prepotential $\Psi$, $\bar{D}_{\dot{i}} \Psi = 0$, as

$$G_{ij} = \frac{1}{4} (D_{ij} + 4S_{ij}) \Psi + \frac{1}{4} (\bar{D}_{\dot{i}j} + 4\bar{S}_{\dot{i}j}) \bar{\Psi}.$$  

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The superfield $S^{ij}$ and its conjugate $\bar{S}_{ij}$ are components of the superspace torsion tensor. The superspace torsion and the algebra of covariant derivatives are given in [4].

The vector multiplet, on the other hand, is described by its field strength $W$, which is a chiral superfield, $\bar{D}_i W = 0$, obeying the Bianchi identity

$$\Sigma^{ij} := \frac{1}{4} (D^{ij} + 4 S^{ij}) \bar{W} = \frac{1}{4} (\bar{D}^{ij} + 4 \bar{S}^{ij}) W.$$ (3)

Such a superfield is called reduced chiral. Comparing this equation to (2), it is obvious that the superfield $\Sigma^{ij}$ is a composite tensor multiplet. Moreover, it is clear that the tensor prepotential $\Psi$ is defined only up to shifts $\delta \Psi = i \Lambda$ where $\Lambda$ is a reduced chiral superfield. This is just the superfield version of the component gauge transformation $\delta B_{mn} = 2 \partial_{[m} \Lambda_{n]}$ which leaves the three-form field strength $H_{mnp} = 3 \partial_{[m} B_{np]}$ invariant.

Within the projective superspace formulation of conformal supergravity [4], both of these multiplets find a natural realization. The tensor multiplet is described in terms of a real $O(2)$ multiplet $G^{ij}(v) := G^{ij} v_j v_i$ obeying $D^{ij}_v G^{ij} = \bar{D}^{ij}_\alpha G^{ij} = 0$ where $D_v^{ij} := v_i D^j \delta_\alpha$ and $\bar{D}^{ij}_\alpha := v_i \bar{D}^j \delta_\alpha$. The vector multiplet can be described by a tropical prepotential $V$, which is a real projective multiplet, $D^{ij}_\alpha V = \bar{D}^{ij}_\alpha V = 0$, of weight zero, $V(cv^i) = c V(v^i)$. The reduced chiral superfield $W$ is related to $V$ by

$$W = \frac{1}{8 \pi} \oint_C v^b dv_k \frac{u_i u_j}{(v^i u)} (D^{ij} + 4 \bar{S}^{ij}) V.$$ (4)

The fixed isotwistor $u_i$ introduced here is subject only to the condition that $v^b u_k \neq 0$ along the contour $C$ in $\mathbb{C} P^1$; the construction proves to be independent of the choice of $u_i$.

The complementary nature of these properties allows us to generate tensor multiplets from vector multiplets and vice-versa. Consider a system of $n_V$ Abelian vector multiplets with field strengths $W_I$, $I = 1, \ldots, n_V$. Given a holomorphic homogeneous function $F(W_I)$ of degree one (to guarantee the correct super-Weyl transformation of $G^{ij}$), we can define a composite tensor multiplet

$$G^{ij} := \frac{1}{4} (D^{ij} + 4 S^{ij}) F(W_I) + \frac{1}{4} (\bar{D}^{ij} + 4 \bar{S}^{ij}) \bar{F}(\bar{W}_I).$$ (5)

This is a standard construction and a trivial application of (4) using the chiral function $F$ as a composite tensor prepotential.

Conversely, we may take a system of $n_T$ tensor multiplets described by their field strengths $G^{ij}_A$, with $A = 1, \ldots, n_T$, with $G^{ij}_A := v_i v_j G^{ij}_A$ the corresponding $O(2)$ multiplets. We may similarly construct a composite vector prepotential from any function $V(G^{ij}_A)$ that is real and homogeneous of degree zero. Applying (4) leads immediately to a composite field strength $\mathbb{W}$,

$$\mathbb{W} := \frac{1}{8 \pi} \oint_C v^b dv_k \frac{u_i u_j}{(v^i u)} (D^{ij} + 4 \bar{S}^{ij}) V.$$ (6)

These two observations together enable us to take any action involving tensor and/or vector multiplets and convert it to a higher derivative action by identifying one or both multiplets as composite. Iterating the procedure leads to increasingly complex higher derivative interactions. We first demonstrate how this procedure allows the construction of general two-derivative interactions of tensor multiplets; then we extend to two-derivative interactions of real $O(2n)$ multiplets; and then finally we extend the argument to higher derivative theories.

2 Self-couplings of tensor multiplets

Consider the simplest interaction between vector and tensor multiplets: the supersymmetric $BF$ coupling. This action can be written either as a chiral superspace integral involving the tensor prepotential $\Psi$ and
vector field strength \( \mathcal{W} \),
\[
S = \int d^4x \, d^4\theta \, \mathcal{E} \, \Psi \mathcal{W} + \text{c.c.} = -\frac{1}{2} \int d^4x \, \varepsilon^{mnpq} \, B_{mn} F_{pq} + \cdots ,
\]
(7)
or as a projective superspace integral involving the tensor field strength \( \mathcal{G}^{(2)} \) and the vector prepotential \( \mathcal{V} \),
\[
S = \frac{1}{2\pi} \int_C v^i dv_i \int d^4x \, d^4\theta \, d^4\tilde{\theta} \, \frac{E}{\mathcal{S}^{(2)}(\mathcal{S}^{(2)})} \mathcal{G}^{(2)}(\mathcal{V}) = \frac{1}{3} \int d^4x \, \varepsilon^{mnpq} \, H_{mnp} A_q + \cdots .
\]
(8)
This action is topological and involves no propagating degrees of freedom; however, by identifying one or both multiplets as composite, we may construct nontrivial actions.

For example, the unique superconformal action for a single tensor multiplet, known as the improved tensor multiplet action [6], can be written in either form. In projective superspace, it is given by (8) where \( \mathcal{V} \) is replaced by the composite vector prepotential \( \mathcal{V} = \ln(\mathcal{G}^{(2)}/i\bar{\mathcal{Y}}^{(1)} \mathcal{Y}^{(1)}) \). The arctic multiplet \( \mathcal{Y}^{(1)} \) appearing here can be shown to be a pure gauge degree of freedom. Equivalently, the action is given by (7) with the corresponding replacement \( \mathcal{W} \rightarrow \mathcal{W} \). This form is simplest to use in principle since the techniques to evaluate the superspace integral (7) are well known: we simply require \( \mathcal{W} \). To construct it, we must evaluate the contour integral in (6). The way this is usually done is to make a choice for \( u_i \) and then to evaluate the contour explicitly, e.g. by representing \( v^i = v^i(1, \zeta) \). This breaks manifest \( SU(2) \) covariance.

It is possible, however, to keep \( SU(2) \) covariance along the way [4]. One does this by first evaluating the spinor derivatives on their argument and exploiting the contour integral to rewrite the integrand as
\[
\mathcal{W} = \frac{1}{8\pi} \int_C v^i dv_i \, \left( \frac{\tilde{M}}{3} \mathcal{G}^{(2)} - \frac{4}{9} \frac{\chi_a^{(1)} \chi^{(1)}}{(\mathcal{G}^{(2)})^2} \right)
\]
(9)
where \( \chi^{(1)}_a := v_i \bar{D}_a g^{ki} \) and \( \tilde{M} := (\bar{D}_{jk} + 12 \bar{S}_{jk}) \mathcal{G}^{jk} \). This has the advantage that all dependence on the auxiliary isotwistor \( u_i \) vanishes. The remaining contour integral can be done in an \( SU(2) \) covariant way [3], with the result recast as
\[
\mathcal{W} = -\frac{G}{8}(\bar{D}_{ij} + 4\bar{S}_{ij}) \left( \frac{\mathcal{G}^{ij}}{G^2} \right) , \quad \mathcal{G} := \sqrt{\frac{1}{2} \mathcal{G}^{ij} \mathcal{G}_{ij}} .
\]
(10)
It is remarkable that this expression is chiral and obeys (5). (An equivalent but less compact expression for \( \mathcal{W} \), obtained by brute force, was given in [7].)

Superconformal actions involving several tensor multiplets can be written in an analogous fashion. In projective superspace, they are given by a general real weight-two projective Lagrangian \( \mathcal{L}^{(2)} = \mathcal{L}^{(2)}(\mathcal{G}^{(2)}_A) \) which is homogeneous of degree one. Such a Lagrangian can always be rewritten (though not uniquely) as
\[
\mathcal{L}^{(2)} = \mathcal{G}^{(2)}_A \Psi^A , \quad \Psi^A = \Psi^A(\mathcal{G}^{(2)}_B) .
\]
(11)
The corresponding projective superspace action can then be converted into a chiral action analogous to (7),
\[
S = \int d^4x \, d^4\theta \, \mathcal{E} \, \Psi^A \Psi^A + \text{c.c.}
\]
(12)
with \( \Psi^A \) given by
\[
\Psi^A = \frac{1}{3} \mathcal{F}^{A,B} \bar{M}_B + \frac{4}{9} \mathcal{F}^{A,B,C} \bar{\chi}^i_B \bar{\chi}^j_C
\]
(13)
where \( \mathcal{F}^{A,B} \) and \( \mathcal{F}^{A,B,C} \) are functions of \( \mathcal{G}^{ij}_A \) given by
\[
\mathcal{F}^{A,B} := \frac{1}{8\pi} \int_C v^i dv_i \, \frac{\partial \Psi^A}{\partial \mathcal{G}^{(2)}_B} , \quad \mathcal{F}^{A,B,C} := \frac{\partial \mathcal{F}^{A,B}}{\partial \mathcal{G}^{ij}_C} .
\]
(14)
This construction appeared originally in flat superspace [8]. Its component form coupled to conformal supergravity appeared in [9].
3 Adding $O(2n)$ multiplets

We may extend the above procedure by considering $O(2n)$ multiplets. As an example, consider a more general projective Lagrangian of the form

$$L^{(2)} = \frac{Q^{(2n)}}{(G^{(2)})^{n-1}} = G^{(2)} \cdot \frac{Q^{(2n)}}{(G^{(2)})^{n}}$$

(15)

where $Q^{(2n)}$ is a real $O(2n)$ multiplet obeying

$$Q^{(2n)} = Q^{i_{1} \cdots i_{2n} v_{i_{1}} \cdots v_{i_{2n}}} \cdot (Q^{i_{1} \cdots i_{2n}})^{*} = Q_{i_{1} \cdots i_{2n}}, \quad \mathcal{D}_{a}^{(1)} Q^{(2n)} = \mathcal{D}_{a}^{(1)} Q^{(2n)} = 0.$$  

(16)

The composite vector multiplet we construct from the above expression has the prepotential $\mathcal{V} = \frac{Q^{(2n)}}{(G^{(2)})^{n}}$. As before, it is possible to evaluate the spinor derivatives in such a way as to completely eliminate the auxiliary isotwistor $u_{i}$. The contour integral (6) can then be evaluated in an $SU(2)$ covariant way and the result recast as

$$\mathcal{W} = -\frac{(2n)!}{2^{2n+2} (n+1)! (n-1)!} \mathcal{G} (\bar{D}_{ij} + 4 \bar{S}_{ij}) \mathcal{R}^{ij},$$

(17)

where

$$\mathcal{R}^{ij} = \frac{1}{G^{2n}} \left( \delta^{ij} - \frac{1}{2G^{2n}} \mathcal{G}_{kl} \right) \mathcal{Q}^{i_{1} \cdots i_{2n-2} g_{i_{1}i_{2}} \cdots g_{i_{2n-3}i_{2n-2}}}. \quad \mathcal{H}^{ijkl} \mathcal{H}^{mn} \mathcal{G}_{mn}.$$

(18)

An interesting application of this result is for the case $Q^{(4)} = (\mathcal{H}^{(2)})^{2}$ with the Lagrangian $L^{(2)} = \frac{\mathcal{H}^{(2)}}{G^{(2)}}$. This is a curved-superspace version of that proposed in [10] to describe the classical universal hypermultiplet [11]. Applying our general formula leads to

$$\mathcal{W} = -\frac{\mathcal{G}}{16} (\bar{D}_{ij} + 4 \bar{S}_{ij}) \mathcal{R}^{ij}, \quad \mathcal{R}^{ij} = \frac{1}{G^{4}} \left( \delta^{ij} - \frac{1}{2G^{2n}} \mathcal{G}_{kl} \right) \mathcal{H}^{ijkl} \mathcal{H}^{mn} \mathcal{G}_{mn}.$$

(19)

The construction of the component Lagrangian can then be carried out by conventional means.

4 Higher derivative couplings of tensor and vector multiplets

We now turn to our main goal: the construction of higher derivative actions. This can be done quite straightforwardly by iterating the above procedure. Begin with a set of tensor multiplets $\mathcal{G}^{ij}_{A}$ and Abelian vector multiplets $\mathcal{W}_{I}$. We construct a set of degree-zero functions $\mathcal{V}_{I}$ of the tensor multiplets $\mathcal{G}^{(2)}_{A}$, which lead to the composite vector multiplet field strengths

$$\mathcal{W}_{I} := \frac{1}{8\pi} \oint_{C} v^{k} dv_{k} (u^{ij}_{I}) \left( \mathcal{D}^{ij} + 4 \mathcal{S}^{ij} \right) \mathcal{V}_{I}(\mathcal{G}^{(2)}_{A}).$$

(20)

Using both sets of vector multiplets, we introduce a set of degree-one holomorphic functions $F_{A}(\mathcal{W}_{I}, \mathcal{W}_{I})$ which can be used to construct composite tensor multiplet field strengths

$$\mathcal{G}^{ij}_{A} := \frac{1}{4} \left( \mathcal{D}^{ij} + 4 \mathcal{S}^{ij} \right) F_{A}(\mathcal{W}_{I}, \mathcal{W}_{I}) + \frac{1}{4} \left( \mathcal{D}^{ij} + 4 \mathcal{S}^{ij} \right) \mathcal{F}_{A}(\mathcal{W}_{I}, \mathcal{W}_{I}).$$

(21)

Then we may introduce these composite tensor multiplets into the functions $\mathcal{V}_{I}$ in (20), leading to

$$\mathcal{W}_{I} := \frac{1}{8\pi} \oint_{C} v^{k} dv_{k} (u^{ij}_{I}) \left( \mathcal{D}^{ij} + 4 \mathcal{S}^{ij} \right) \mathcal{V}_{I}(\mathcal{G}^{(2)}_{A}, \mathcal{G}^{(2)}_{A}).$$

(22)
This new vector multiplet can be used to construct new composite tensor multiplets and so on and so forth, with each iteration adding two spinor derivatives (or one vector derivative) to the interaction.

This method of constructing higher derivative actions should be contrasted with the more traditional way of generating higher derivative structures using the chiral projection operator $\tilde{\Delta}$, which is a curved space generalization of $D^4 = \tilde{D}_i\tilde{D}^i/48$. Given any scalar superfield $U(z)$ which is inert under super-Weyl transformations, its descendant $\Delta U$ is chiral and super-Weyl weight two. Given a vector multiplet $\mathcal{W}$ that is nowhere vanishing, we can then define the chiral scalar $\mathcal{W}_-\Delta U$ which is invariant under the super-Weyl transformations. We can then construct $\mathcal{W}_-\Delta (\mathcal{W}_-\Delta U)$, and so on and so forth.

Using these composite chiral operators, one may construct higher derivative actions involving chiral superspace actions. However, it is usually possible to convert the chiral superspace action into an integral over the whole superspace by eliminating one of the chiral projection operators. Schematically, if $\mathcal{L}_c = \Phi \Delta U$ for some chiral superfield $\Phi$ and a well-defined local and gauge-invariant operator $U$, then

$$\int d^4x d^4\theta E \Phi \Delta U = \int d^4x d^4\theta d^4\bar{\theta} E \Phi U . \quad (23)$$

Thus, higher derivative actions of this type are invariably most naturally written as integrals over the entire superspace and are not intrinsically chiral. This has important ramifications for perturbative calculations, where non-renormalization theorems place strong restrictions on intrinsic chiral Lagrangians.

The constructions we are considering are interesting partly because they include higher derivative terms which cannot be written as full superspace integrals without introducing prepotentials. We take the example of a projective Lagrangian $\mathcal{L}^{(2)} = \mathcal{G}^{(2)} \mathcal{V}(\mathcal{G}^{(2)}_A)$ of several tensor multiplets $\mathcal{G}^{(2)}_A$ and one composite tensor multiplet

$$\mathcal{G}^{(2)} = \frac{1}{4}((D^{(1)})^2 + 4S^{(2)})F(\mathcal{W}_I) + \frac{1}{4}((\bar{D}^{(1)})^2 + 4\bar{S}^{(2)})\bar{F}(\bar{W}_I) . \quad (24)$$

The action can be rewritten

$$S = \frac{1}{2\pi} \oint_C v^i dv_i \int d^4x d^4\theta d^4\bar{\theta} E \mathcal{G}^{(2)} F(\mathcal{W}_I) + \text{c.c.} \quad (25)$$

which is a special case of a more general action

$$S = \frac{1}{2\pi} \oint_C v^i dv_i \int d^4x d^4\theta d^4\bar{\theta} E \mathcal{G}^{(2)}_A \Omega(\mathcal{W}_I, \mathcal{G}^{(2)}_A) + \text{c.c} . , \quad \bar{D}^{(1)}\bar{\Omega} = 0 . \quad (26)$$

The complex integrand $\Omega$ is required to be annihilated by only 1/4 of the spinor derivatives. Such an action is the locally supersymmetric version of

$$S = -\frac{1}{8\pi} \oint_C v^k dv_k \int d^4x \frac{u_i u_j}{(v^i u_i)^2} \bar{D}^{ij} D^4 \Omega(\mathcal{W}_I, \mathcal{G}^{(2)}_A) + \text{c.c} . , \quad D^4 := \frac{1}{48} \bar{D}^{ij} D_{ij} . \quad (27)$$

In other words, it is an integral over 3/4 of the Grassmann coordinates of superspace.\footnote{This is a generalization of the construction\cite{12} of rigid superconformal invariants containing $F^m$.\footnote{Special holomorphic three-derivative contributions to $\mathcal{N} = 2$ supersymmetric Yang-Mills effective actions, which are given as an integral over 3/4 of superspace, have been discussed in\cite{13}.}} Because this class of higher derivative action cannot be written as an integral over the full superspace without introducing gauge-dependent prepotentials, an $\mathcal{N} = 2$ extension of the well known $\mathcal{N} = 1$ non-renormalization theorems likely applies.

In order for the action (26) to be super-Weyl invariant, $\Omega$ must be degree 1 in $\mathcal{W}_I$ and degree zero in $\mathcal{G}^{(2)}_A$. The simplest example is $\Omega = \mathcal{W} \mathcal{V}$ where $\mathcal{V} = \ln(G^{(2)}_A / \mathcal{Y}^{(1)}(\mathcal{Y}^{(1)})$. (In this case, the arctic multiplets are...}
again pure gauge degrees of freedom.) A straightforward argument shows that the action can be rewritten
\[
\int d^4x \, d^4\theta \, \mathcal{E} \, W^{\mathbb{V} \mathbb{W}}
\]
which is the supersymmetric generalization of
\[
\int d^4x \, e F_{pq} = \frac{1}{4} \int d^4x \, e F_{pq}(\varepsilon_{kl} G_{ij} D_p G_{ik} D_q G_{jl}) / G^3 + \cdots .
\]
(28)

Thus the actions (26) describe three-derivative couplings in the bosonic sector. Needless to say, the integrand \( \Omega \) can be generalized to include composite vector and tensor multiplets, which leads to four-derivative and higher interactions.

For simplicity we have restricted our higher-derivative discussion to vector and tensor multiplets. However, as emphasized in [3], the composite vector multiplets (6) may be built out of not only tensor multiplets, but also more general \( \mathcal{O}(2n) \) multiplets and even polar multiplets. The corresponding chiral super-space action (7) can then be used to give a new and very general class of higher-derivative supergravity-matter couplings.

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