Variant supercurrents and linearized supergravity

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Abstract

In this paper, the variant supercurrents based on consistency and completion in off-shell \(\mathcal{N}=1\) supergravity are studied. We formulate the embedding relations for supersymmetric current and energy tensor into a supercurrent multiplet. Corresponding linearized supergravity is obtained with an appropriate choice of the Wess–Zumino gauge in each gravity supermultiplet.

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1. Introduction

According to the structure of supersymmetry algebra, the \(R\) current \(j_\mu^R\), supersymmetric current \(j_\mu\), and energy tensor \(T_{\mu\nu}\) corresponding to the \(R\) charge, supercharge and spacetime momentum respectively can be embedded into a supermultiplet. This multiplet is known as supercurrent [1]. The superfield form of a supercurrent and the constraint it satisfies are found to be quite model dependent, although some general considerations from symmetries can be taken into account [2–5].

There is a standard scheme for analyzing the structure of supercurrent and its corresponding linearized supergravity for a given physical system. The procedure is as follows.

(1) Beginning with the physics systems studied, and the conservation conditions,

\[
\partial^\mu j_\mu = 0, \quad \partial^\mu T_{\mu\nu} = 0,
\]

one has to find the embedding relations for the supersymmetric current and the energy–momentum tensor into a supercurrent. During this stage the supercurrent multiplet and the constraint it satisfies are determined at meantime.

(2) Through the constraint that the supercurrent satisfies, we obtain the constraints on the gauge transformation superfield \(L\) of a gravity supermultiplet, which tell us the analogy of the Wess–Zumino gauge in a gravity supermultiplet.
(3) Collecting the embedding relations of gravity and gravitino into a gravity supermultiplet, the action of linearized supergravity can be directly read in components.

In this paper, we study the structures of three new variant supercurrents [7] using the results obtained earlier in [8, 9]. The existence of these variant supercurrents is based on consistency and completion in $N = 1$ off-shell linearized supergravity. Other supercurrents deduced via this viewpoint include the Ferrara–Zumino (FZ) multiplet, new minimal multiplet [10] and $S$ multiplets [13] all of which have completions of quantum field theories (see also [11, 15–17]). The variant supercurrents are defined as follows:

**Case I:**

\[ \bar{D} \dot{\alpha} J^{I}_{\dot{\alpha} \alpha} = i \eta_{\alpha}, \quad \bar{D} \eta = D^{a} \eta_{a} - \bar{D}^{a} \bar{\eta}_{a} = 0 \]  
(1.1)

which is a minimal off-shell supergravity. The second case is

**Case II:**

\[ \bar{D} \dot{\alpha} J^{II}_{\dot{\alpha} \alpha} = i \eta_{\alpha} + \hat{\chi}_{\alpha}, \]  
(1.2)

with \( \bar{D}^{a} \eta_{a} = D^{a} \eta_{a} - \bar{D}^{a} \bar{\eta}_{a} = 0 \) and \( \bar{D}^{a} \hat{\chi}_{a} = D^{a} \hat{\chi}_{a} - \bar{D}^{a} \bar{\hat{\chi}}_{a} = 0 \). The last case is

**Case III:**

\[ \bar{D} \dot{\alpha} J^{III}_{\dot{\alpha} \alpha} = i \eta_{\alpha} + D_{\alpha} X, \]  
(1.3)

with \( \bar{D}^{a} \eta_{a} = D^{a} \eta_{a} - \bar{D}^{a} \bar{\eta}_{a} = 0 \) and \( \bar{D} X = 0 \).

There are some common results in variant supercurrents. Firstly, the $R$ current is not conserved, which can be easily observed from the constraint of equations (1.1)–(1.3). Secondly, there exist some special constraints for the energy tensor $T_{\mu \nu}$ as shown below. These constraints exclude some simple physical systems we are familiar with. Thus, they might serve as the necessary conditions for the existence of variant supercurrents.

The paper is organized as follows. In section 2, we discuss the minimal case I. Section 3 is devoted to study non-minimal cases II and III. The solutions to constraints (1.1)–(1.3) are obtained, with comments on conditions that the energy tensor has to satisfy. The actions of linearized supergravity are obtained after the analogy of the Wess–Zumino gauge in each case is discussed. In section 4, we conclude and discuss the difference between variant supercurrents and other supercurrents in the literature.

### 2. Minimal case I

In this paper, we follow the conventions of Wess and Bagger [6]. The real vector superfield $J_{\mu}$ is defined in the bi-spinor representation as

\[ J_{\mu a} = \sigma^{\mu} J_{\mu a}, \quad \text{and} \quad J_{\mu} = \frac{i}{2} \sigma^{\mu} J_{\mu a}. \]  
(2.1)

The components are expressed as

\[ S = C^S + i \theta \chi^S - i \bar{\theta} \bar{\chi}^S + \frac{i}{2} \theta^2 (M^S + i N^S) - \frac{i}{2} \bar{\theta}^2 (M^S - i N^S) - \theta \sigma^m \bar{\theta} \upsilon^S_m \]
\[ + i \bar{\theta} \bar{\theta} \left( \bar{\chi}^S + \frac{i}{2} \sigma^m \partial_m \chi^S \right) - i \theta \bar{\theta} \left( \chi^S + \frac{i}{2} \sigma^m \partial_m \bar{\chi}^S \right) + \frac{1}{2} \theta^2 \bar{\theta}^2 \left( D^S + \frac{1}{2} \Box C^S \right). \]  
(2.2)

Note that the lowest component field $C^S$ in the supercurrent superfield $J$ is the $R$ current $j^S_5$.

We deduce a new constraint on the supercurrent from constraint (1.1):

\[ \bar{D}^{\beta} \bar{D}^{\alpha} J^{I}_{\alpha \alpha} = 0 \]  
(2.3)

The first equation in constraint (1.1) can be classified into its real and imaginary parts, respectively\(^3\). Explicit expressions for these components can be found in [14].

\(^3\) Similar methods are applied to the other two cases that we will discuss in this paper.
Solving equations (2.3) and (1.1), we obtain

\[ J_I^\mu = C_\mu + \theta \left( j_\mu + \frac{1}{3} \sigma^\nu j_\nu \right) + \bar{\theta} \left( \bar{j}_\mu + \frac{1}{3} \bar{\sigma}^\nu \bar{j}_\nu \right) \]

\[ + (\theta \sigma^\nu \bar{\theta}) \left( a T_{\nu \mu} + b T_{\eta \mu} + \frac{1}{4} \epsilon_{\mu \nu \rho \lambda} (\partial^\rho C^\lambda - \partial^\lambda C^\rho) - \frac{1}{2} \Phi_{\nu \mu} \right) \]

\[ + \theta^2 \bar{\theta} \left( \frac{i}{3} \bar{\sigma}^\nu \partial^\mu j_\nu \right) + \bar{\theta}^2 \theta \left( \frac{i}{3} \sigma^\nu \partial^\mu \bar{j}_\nu \right) + \theta^2 \bar{\theta} \left( - \frac{1}{2} \Delta C_\mu - \frac{1}{2} \partial^\mu \partial^\nu C_{\nu} \right) \] (2.4)

and

\[ \eta_{\alpha} = -i \Lambda_\alpha(y) + \left( \delta_\alpha^\beta \Delta - 2i \bar{\sigma}^\mu \sigma^\nu \Phi_{\mu \nu} \right) \theta^\beta + \theta^2 (\sigma^\mu \partial_\mu \bar{\Lambda}(y))_{\alpha} \] (2.5)

where the coefficients \( a, b \) are introduced to define \( \hat{T}_{\mu \nu} \mid s = a T_{\mu \nu} + b T_\eta \mu \mid s \). In this case, constants \( a \) and \( b \) are given by

\[ a = -4b, \quad 2b \partial_\nu T_\eta \mu = -\partial^\mu /\Phi_1 \mu \nu, \quad \Box T = 0. \] (2.6)

The lower indices \( s, a \) in \( \hat{T}_{\mu \nu} \mid s \) refer to the symmetric and anti-symmetric parts, respectively. \( \Phi^{\rho \sigma} \) and \( \Delta \) are the tensor field and \( D \)-term in the \( \eta \) superfield, respectively. The degrees of freedom of \( \hat{T}_{\mu \nu} \mid s \) can be considered as totally provided by \( \Phi^{\rho \sigma} \). Physical systems with the energy tensor \( T_{\mu \nu} \) that satisfies these special constraints are extraordinary. The non-existence of these conditions might serve as a proof that the first kind of constrained supercurrent is not physical. This question will be investigated further.

The degrees of freedom in this case are described by \( (C_\mu, \chi_\mu, \hat{T}_{\mu \nu} \mid s, \hat{T}_{\mu \nu} \mid a) \), which imply that supersymmetric theories correspond to 12/12 off-shell supergravity. Gauging the supercurrent \( J^1 \) in supergravity via coupling

\[ \int d^4x \int d^4\theta J^1_{au} H^{au}. \] (2.7)

the gauge invariance of action (2.7) under the transformation \( H_\mu \rightarrow H_\mu + \Delta_\mu \), or equivalently via its bi-spinor expression

\[ H_{au} \rightarrow H_{au} + D_a L_{\alpha} - D_a L_a, \] (2.8)

leads to

\[ D_a D^2 L_a + D_a D^2 L^a = 0. \] (2.9)

Here, the superfield \( L \) is defined as

\[ \Delta_\mu = -\frac{1}{2} \delta^{\alpha \beta} (D_\alpha L_{\beta} - D_{\beta} L_{\alpha}), \] (2.10)

where \( \Delta_\mu \) is a general real superfield. Equation (2.10) suggests that the relations of the embedding graviton and gravitino into the supergravity multiplet \( H_\mu \) follow those of [12, 13]^4. \( H_\mu \mid \phi \) is divided into the symmetric part \( v^{H}_{\mu \nu} \) and the anti-symmetric part \( B_{\mu \nu} \). The gauge transformations are as follows:

\[ \delta h_{\mu \nu} = \partial_\mu \xi_\nu + \partial_\nu \xi_\mu, \quad \delta \Psi_{\mu a} = \partial_\mu \omega_a \]

Constraint (2.9) imposes some equations on components in \( L \), which implies a set of constraints to the components of \( \Delta_\mu \) via equation (2.10). These constraints determine the analog of the Wess–Zumino gauge for a supermultiplet \( H_\mu \). Defining

\[ L_a = i D_a V. \] (2.11)

Following the conventions we take, one can see that these embedding relations are independent of constraints on \( L_a \). They are valid throughout this paper.
equation (2.9) leads to the identification of $V$ as a Wess–Zumino gauged vector superfield. The constraints on components in $\Delta_\mu$ are

$$L_\alpha | = L_\alpha |_\theta = L_\alpha |_{\bar{\theta}} = 0$$

and

$$\partial^m (L_\alpha |_{\phi^m}) = -2 (L_\alpha |_{\phi^m})$$

which imply that the $B_{\mu\nu}$ field in a gravity supermultiplet cannot be set to zero.

One can see that the analogy of the Wess–Zumino gauge is as follows:

$$H_\mu | = H_\mu |_\theta = H_\mu |_{\bar{\theta}} = H_\mu |_{\theta^2} = H_\mu |_{\bar{\theta}^2} = 0.$$  

The residual degrees of freedom in a gravity supermultiplet are represented by $(h_{\mu\nu}, B_{\mu\nu}, \Psi_{1\mu\alpha}$ and $D_\mu^a)$, which describe 12/12 minimal supergravity. They match with that of the supercurrent.

Following the notation in equation (2.2), we obtain the action in components:

$$S = -\epsilon^H_{\mu\nu} \tilde{T}^{\mu\nu} | - \frac{1}{2} j^A_{\mu} D_\mu^H + \left( \frac{i}{2} \chi^{(\mu)}(H) \chi^{(H)\mu} + \text{c.c.} \right)$$

The kinetic term of graviton can be constructed via an appropriate derivative operator [7]. Starting with the constraint on a gauge transformation superfield, the results in [7] can be reproduced. Similar results can be found in non-minimal cases.

### 3. Reducible cases

Now we discuss the non-minimal cases II and III. Their supercurrent multiplets both include 16 + 16 degrees of freedom (supermultiplets with 16 + 16 degrees of freedom are also discussed in [18, 19]), which are manifested by their corresponding gravity supermultiplets.

In comparison with the minimal case I, the gauge transformation superfield $L_\alpha$ is more constrained, which is the origin of more degrees of freedom in gravity supermultiplets.

#### 3.1. Reducible case II

Constraint (1.2) implies that

$$\bar{D}^\delta \tilde{D}^\mu J_{\alpha}^{I\mu} = 0.$$  

Solving equations (3.1) and (1.2) gives

$$J_{\mu}^{II} = C_\mu + \theta \left( j_\mu + \frac{1}{3} \sigma_\mu \bar{\sigma}_\nu j_\nu + \frac{1}{3} \sigma_\nu \bar{\sigma}_\mu \bar{\psi} \right) + \bar{\theta} \left( j_\mu + \frac{1}{3} \bar{\sigma}_\nu \sigma_\mu j_\nu - \frac{1}{3} \bar{\sigma}_\mu \psi \right)$$

$$+ (\delta^\nu \bar{\sigma}^\mu \bar{\sigma}_\nu \bar{\psi}) \left( a T_{\nu\mu} - \frac{b}{a + 4b} Z_{\nu\mu} + \frac{1}{4} \epsilon_{\nu\mu\rho\lambda} (\partial^\rho C^\lambda - \partial^\lambda C^\rho + \Sigma^{\rho\lambda}) - \frac{1}{2} \Phi_{\nu\mu} \right)$$

$$+ \bar{\theta}^2 \left( - \frac{2i}{3} \partial_\mu \bar{\psi} + \frac{i}{3} \sigma^\mu \bar{\sigma}_\nu j_\nu \right) + \bar{\theta}^2 \left( \frac{2i}{3} \partial_\mu \psi + \frac{i}{3} \sigma^\mu \bar{\sigma}_\nu j_\nu \right)$$

$$+ \bar{\theta}^2 \left( \frac{1}{2} \partial_\mu Z - \frac{1}{2} \square C_\mu - \frac{1}{2} \partial_\mu \partial^\nu C_\nu + \frac{3}{2} \partial^\nu \Sigma_{\mu\nu} \right)$$

and

$$\eta_\alpha = -i \Lambda_\alpha(y) + (\delta^\mu_\alpha A - 2i \bar{\sigma}^\mu \sigma_\nu \Phi_{\nu\mu}(y)) \theta_\alpha + \bar{\theta}^2 (\sigma^\mu \partial_\mu \Lambda(y)) \alpha$$

$$\tilde{\eta}_\alpha = -i \bar{\psi}_\alpha(y) + (\delta^\mu_\alpha Z - 2i \bar{\sigma}^\mu \sigma_\nu \Sigma_{\nu\mu}(y)) \theta_\alpha + \bar{\theta}^2 (\sigma^\mu \partial_\mu \bar{\psi}(y)) \alpha.$$
The coefficients $a$ and $b$ satisfy

$$\begin{align*}
(a + 4b) T &= -Z, \\
2b\partial_i T &= -\partial^\mu \Phi_{\mu
u}.
\end{align*}$$

(3.4)

As emphasized above, the existence of $a$ and $b$ is necessary for physical systems described by case II. The multiplet $J^\mu_{II}$ contain $12 + 12$ degrees of freedom, a Weyl spinor $\psi$, a closed two-form $\Sigma_{\mu\nu}$, and a real scalar $Z$. Thus, it describes a $16 + 16$ supermultiplet.

Gauging the supercurrent $J^\mu_{II}$ in supergravity via coupling

$$\int d^4x \int d^4\theta J^\mu_{II} H^{\mu\alpha},$$

(3.5)

the gauge invariance of the action under transformation (2.8) leads to

$$D^\mu D^2 L_{\alpha} = D^\mu D^2 L_{\alpha} = 0.$$  

(3.6)

The embedding relations of the graviton and gravitino into the $H_{\mu}$ superfield is the same as in case I. Note that the equation of motion of a field strength chiral superfield without the FI term is exactly the same with equation (3.6). The analogy of the Wess–Zumino gauge is given by

$$H_{\mu} |\sigma = H_{\mu} |\sigma = H_{\mu} |\sigma = H_{\mu} |\sigma = 0.$$  

(3.7)

The residual degrees of freedom in a gravity supermultiplet are represented by $(h_{\mu\nu}, B_{\mu\nu}, \Psi_{\mu\alpha})$ and $(\bar{D}_{\mu})$, which describe $16/16$ linearized supergravity. They match with that of the supercurrent. The corresponding action is in components with notation in equation (2.2):

$$S = -\upsilon_{\mu} H_{\mu} |\sigma - B_{\mu\nu} \hat{T}_{\mu\nu} |\sigma + \frac{1}{2} j^{5}_{\mu} D_{\mu} H + \left(\frac{i}{2} \chi^{(4)}(H)_{\mu} + \text{c.c.}\right).$$

(3.8)

### 3.2. Reducible case III

Finally, we address the third possible constraint satisfied by the supercurrent. Solving equation (1.3), we obtain $J^\mu_{III}$.

$$J^\mu_{III} = C_{\mu} + \theta^\alpha \left( j_{\mu} + \frac{1}{3} \sigma_{\sigma} \sigma^{\sigma} j_\nu \right) + \bar{\theta} \left( \bar{j}_{\mu} + \frac{1}{3} \bar{\sigma}_{\bar{\sigma}} \sigma^{\bar{\sigma}} \bar{j}_\nu \right) - i\theta^2 \partial_{\mu} \phi + i\bar{\theta}^2 \partial_{\mu} \bar{\phi}^*$$

$$+ (\theta \sigma^{\bar{\sigma}} \bar{\theta}) \left( aT_{\nu} - 2\text{Re}(F)\eta_{\nu\mu} + \frac{1}{2} \bar{\epsilon}_{\nu\mu\alpha\beta} \partial^\alpha C_{\lambda} - \frac{1}{2} \Phi_{\nu\mu} \right)$$

$$+ \bar{\theta}^2 \bar{\theta} \left( \frac{1}{3} \bar{\sigma}_{\bar{\sigma}} \partial_{\bar{\nu}} j_{\rho} - \sqrt{2} \bar{\partial}_{\mu} \bar{\psi} \right) + \theta^2 \theta \left( \frac{1}{3} \sigma_{\sigma} \partial_{\sigma} \bar{j}_{\rho} - \sqrt{2} \partial_{\mu} \psi \right)$$

$$+ \bar{\theta}^2 \theta^2 \left( -2\partial_{\mu} (\text{Im}(F)) + \frac{1}{2} \Box C_{\mu} - \frac{3}{2} \partial_{\mu} \partial^\rho C_{\rho} \right)$$

(3.9)

and

$$X = \phi(y) + \sqrt{2} \theta \psi(y) + \theta^2 F$$

$$\eta_{\mu} = -i\Lambda_{\alpha}(y) + (\delta^\alpha_{\bar{\alpha}} \Delta - 2i \sigma^\alpha \sigma_{\nu} \Phi_{\nu\mu}(y)) \theta_{\bar{\alpha}} + \theta^2 (\sigma^\alpha \partial_{\mu} \Lambda(y))_{\bar{\alpha}}.$$  

(3.10)

The components fields in $\eta_{\mu}$ satisfy extra constraints

$$\Delta = -\partial^\mu C_{\mu} - 2\text{Im}(F),$$

$$\Lambda_{\alpha} = \frac{1}{3} (\sigma^\alpha j_{\mu})_{\bar{\alpha}} - \sqrt{2} \psi_{\bar{\alpha}}.$$  

(3.11)

Here, the coefficient $a$ is given by $aT = 6 \text{Re}(F)$, with $F = \text{Re}(F) + i\text{Im}(F)$. The multiplet $J^\mu_{III}$ contains $12 + 12$ degrees of freedom, a Weyl spinor $\psi$, a complex scalar $\phi$, and a complex scalar $F$ (or equivalently $\text{Re}(F)$ and $\Delta$), which imply that it actually is a $16 + 16$ supermultiplet.
Compared with the $S$-multiplet that is introduced to solve problem of FI term in supergravity [13], the scalar $\text{Re}(F)$ is now replaced by $F$. The embedding relations are also very different. Gauging the supercurrent $J^{III}$ in supergravity via coupling
\[ \int d^4x \int d^4\theta J^{III}_{\alpha\dot{\alpha}} H^\alpha{}^\dot{\alpha}, \] (3.12)
the gauge invariance of the action under transformation (2.8) leads to
\[ \bar{D}^2 D^\mu L_\alpha = 0, \quad \bar{D}_{\dot{\alpha}} D^2 \bar{L}^\dot{\alpha} = D_\mu \bar{D}^2 L^\mu. \] (3.13)
As more constraints are imposed, less component fields in a gravity supermultiplet can be set to zero. Constraint (3.13) suggests that the analog of the Wess–Zumino gauge is
\[ H_\mu |\mu| \theta = H_\mu |\dot{\theta}| \bar{\theta} = 0. \] (3.14)
The action can be read in components with notation in equation (2.2):
\[ S = -\psi^H T^{\mu\nu} |\nu| - B_{\mu\nu} \bar{T}^{\mu\nu} |\nu| + \frac{1}{2} j^H_{\mu\nu} D^H_\mu \]
\[ + \left[ \frac{i}{2} \eta_{\mu\nu} (H/\mu) + \frac{1}{4} (M^I + iN^I)(M^H - iN^H) + c.c. \right]. \] (3.15)

4. Conclusions

In this paper, we have studied a set of variant supercurrents that arise from consistency and completion in $N = 1$ off-shell supergravity. We have used the component languages of the superfield to obtain the embedding relations of the supersymmetric current and the energy–momentum tensor into a formalism of linear supergravity. The analogy of the Wess–Zumino gauge in each case has been analyzed in detail.

Instead of the superfield formalisms used to describe variant supercurrents, we find more physical results that are uncovered in the component expressions. First, the consistent conditions for the energy–momentum tensor of supersymmetric theories that can be described by variant supercurrent multiplets are determined explicitly. Second, the component results help identifying corresponding linearized supergravity.

Although supercurrents that include the $S$-multiplet [13], FZ multiplet and minimal multiplet have rich constructions of quantum field theories and important applications, the consistent conditions (2.6) and (3.4) for variant supercurrents studied in this paper imply $\square T = 0$. It can be verified that the variant supercurrents are not viable for simple supersymmetric field theories including pure supersymmetric Yang–Mills theories and SQCD-like theories as a result of $\square T \neq 0$. The main reason for this difference between the variant supercurrent and the $S$-multiplet is that the $i$ factor in front of the linear superfield in equations (1.1)–(1.3) leads to the real and imaginary parts of the $\theta^2\bar{\theta}$ component in these equations exchanged. In the case of the $S$-multiplet, the energy–momentum tensor depends on the $D$-term of the $\eta_\mu$ superfield, and an anti-symmetric tensor field $\Phi_{\mu\nu}$ is related to the anti-symmetric part of $\bar{T}_{\mu\nu}$. In the case of variant supercurrents, however, exchanging the real and imaginary parts of the component $\theta^2\bar{\theta}$ in equations (1.1)–(1.3) leads to the fact that the dependence on an anti-symmetric tensor field $\Phi_{\mu\nu}$ is transferred to the derivative of the energy–momentum tensor trace $\partial_\mu T$, which is the origin of the severe constraint $\square T = 0$ for variant supercurrents.

5 Recently, it was found in [13] that the $S$-multiplet is useful to embed the Fayet–Iliopoulos term for an Abelian gauge supermultiplet into supergravity.

6 As shown in [13], $\square T$ is proportional to $\square \text{Re} F_X$ and $\square \Delta$. Although the component expressions of $F$ and $D$ terms are quite involved, it is expected that both $\square \text{Re} F_X \neq 0$ and $\square \Delta \neq 0$ in terms of equations of motion of relevant fields. In other words, $\square T = 0$ is forbidden in the case of the $S$ multiplet.
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