Consistent and Covariant Anomalies
in Six-dimensional Supergravity

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Abstract

In this note we clarify some issues in six-dimensional (1, 0) supergravity coupled
to vector and tensor multiplets. In particular, we show that, while the low-energy
equations embody tensor-vector couplings that contribute only to gauge anomalies,
the divergence of the energy-momentum tensor is properly non-vanishing. In addi-
tion, we show how to revert to a supersymmetric formulation in terms of covariant
non-integrable field equations that embody corresponding covariant anomalies.

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1 Introduction

Six-dimensional (1,0) supergravity has attracted a large interest in recent years for a number of reasons [1]. On the one hand, vector multiplets coupled to variable numbers of tensor multiplets arise naturally in perturbative type-I vacua [2], and therefore, via duality, play a ubiquitous role in non-perturbative string phenomena. On the other hand, the field equations have revealed the explicit realization of a peculiar aspect of the physics of branes. Namely, branes wrapped on vanishing cycles in the internal manifold may result in the exotic phenomenon of transitions related to tensionless strings [3], and indeed some peculiar singularities in the gauge couplings of (1,0) models in moduli space [4,5] can be ascribed to phase transitions [6] whereby a string becomes tensionless [7]. On a more technical side, these equations present the novel feature of a Green-Schwarz mechanism implemented by terms present in the low-energy effective action, at least for the gauge part of the residual anomaly. One is thus facing a case of unprecedented complexity in supergravity constructions, whereby the model is determined by Wess-Zumino conditions [8], rather than by the usual requirement of local supersymmetry. Moreover, as pointed out in [9], the algebra contains a two-cocycle and the resulting equations are not unique. By and large, this is a remarkable laboratory for current algebra, where one can play explicitly with anomalous symmetries and their consequences.

The present note is devoted to some aspects of current algebra related to the energy-momentum tensor that, although rather simple, are somewhat surprising and were not noticed in [9]. The corresponding analysis is carried out in the next Section. An additional, related problem has to do with the formulation of the resulting equations, that were originally derived in [4] to lowest order in the fermion couplings by requiring local supersymmetry. The subsequent work of [10] and [9] has developed the consistent formulation, but one can actually revert to a covariant formulation, at the price of having non-integrable field equations. The relation between the two sets of equations is one more instance of the link between covariant and consistent anomalies in field theory [11]. Once more, here the anomalies are induced by local couplings of the two-forms, and everything is totally explicit. The resulting covariant equations turn into one another under local
supersymmetry and complete to all orders in the fermi fields the results of [4].

2 The energy-momentum tensor of six-dimensional \((1, 0)\) supergravity

In six-dimensional \((1, 0)\) supergravity coupled to an arbitrary number of tensor and vector multiplets, the \(n\) scalars in the tensor multiplets parametrize the coset space \(SO(1, n)/SO(n)\), and are described by the \(SO(1, n)\) matrix [12]

\[ V = \begin{pmatrix} v_r \\ x^m_r \end{pmatrix} \quad \text{(2.1)} \]

All spinors are symplectic Majorana-Weyl. In particular, the tensorinos \(\chi^m\) are right-handed, while the gravitino \(\Psi_\mu\) and the gauginos \(\lambda\) are left-handed. The tensor fields \(B^r_{\mu\nu}\) are valued in the fundamental representation of \(SO(1, n)\) and satisfy the (anti)self-duality conditions

\[ G_{rs} H^{s\mu\nu} = \frac{1}{6\epsilon} e^{\mu\rho\alpha\beta\gamma} H_{r\alpha\beta\gamma} \quad \text{(2.2)} \]

where \(G_{rs} = v_r v_s + x^m_r x^m_s\). Moreover, their field strengths include Chern-Simons 3-forms for the vector fields according to

\[ H^r = dB^r - c^r z \omega_z \quad \text{(2.3)} \]

where the \(c^r_z\) are constants that determine the gauge part of the residual anomaly polynomial

\[ A = -\sum_{rs} \eta_{rs} c^r x^s \ tr_z(F \wedge F) \ tr_y(F \wedge F) \quad \text{(2.4)} \]

and \(z\) runs over the various factors of the gauge group [4]. Gauge invariance of \(H^r\) then requires that

\[ \delta B^r = c^r z tr_z(AdA) \quad \text{(2.5)} \]

The complete field equations were determined in [3] from the commutator of two supersymmetry transformations on the fermi fields, in the spirit of [13, 12]. The resulting model, however, has gauge and supersymmetry anomalies (to be canceled by fermion
loops) related by Wess-Zumino consistency conditions and is not unique. Aside from subtleties related to the (anti)self-dual antisymmetric tensors [14], all field equations may be derived from

\[ e^{-1} \mathcal{L} = -\frac{1}{4} R + \frac{1}{12} G_{\mu\nu} H^{\mu\nu\rho\sigma} H_{\mu\nu\rho\sigma} - \frac{1}{4} \partial_{\mu} v^{\nu} \partial^{\mu} v_{\nu} - \frac{1}{2} v_{r} c^{r} c^{z} tr_{z} (F_{\mu\nu} F^{\mu\nu}) \\
- \frac{1}{8} \epsilon^{\mu\nu\alpha\beta} c^{r} \tilde{D}_{\nu} tr_{z} (F_{\alpha\beta} F_{\gamma\delta}) \\
- \frac{i}{2} \tilde{\Psi}_{\mu} \gamma^{\mu\nu} D_{\nu} \left[ \frac{1}{2} (\omega + \hat{\omega}) \right] \Psi_{\nu} - \frac{i}{8} v_{r} [H + \hat{H}] \gamma^{\mu\nu} (\tilde{\Psi}_{\mu} \gamma_{\nu} \Psi_{\rho}) \\
+ \frac{i}{48} v_{r} [H + \hat{H}]^{\mu\nu\rho\sigma} (\tilde{\Psi}_{\mu} \gamma_{\nu} \gamma_{\rho} \gamma_{\sigma} \Psi_{\nu}) + \frac{i}{2} \chi^{m} \gamma_{\mu} D_{\mu} (\hat{\omega}) \chi^{m} \\
- \frac{i}{24} v_{r} [H + \hat{H}] \gamma^{\mu\nu} (\tilde{\Psi}_{\mu} \gamma_{\nu}) + \frac{1}{4} \epsilon^{m} [\partial_{\nu} \epsilon^{r} + \partial^{r} \epsilon^{\nu}] (\tilde{\Psi}_{\mu} \gamma^{\nu} \gamma^{m} \chi^{m}) \\
- \frac{1}{8} v_{r} [H + \hat{H}] \gamma^{\mu\nu} (\tilde{\Psi}_{\mu} \gamma_{\nu}) + \frac{1}{8} \chi^{m} \gamma_{\mu} \gamma_{\nu} \gamma_{\rho} \gamma_{\sigma} (\tilde{\Psi}_{\mu} \gamma_{\nu} \gamma_{\rho} \gamma_{\sigma} \Psi_{\nu}) \\
+ \frac{i}{2 \sqrt{2}} v_{r} c^{r} c^{z} tr_{z} [(F + \hat{F}) \gamma^{\mu\nu} \gamma^{\mu} \lambda] + \frac{1}{\sqrt{2}} v_{r} c^{r} c^{z} tr_{z} [(\tilde{\chi}^{m} \gamma^{\mu\nu} \lambda) \hat{F}_{\mu\nu}] \\
+ iv_{r} c^{r} c^{z} tr_{z} (\tilde{\chi}^{m} \gamma^{\mu\nu} \lambda) + \frac{i}{12} x_{r}^{m} x_{z}^{m} \hat{H}_{\mu\nu} c^{r} c^{z} tr_{z} (\tilde{\chi}^{m} \gamma^{\mu\nu} \lambda) \\
+ \frac{1}{16} v_{r} c^{r} c^{z} tr_{z} (\tilde{\chi}^{m} \gamma^{\mu\nu} \lambda) (\tilde{\chi}^{m} \gamma^{\mu\nu} \lambda) \\
- \frac{i}{8} \chi^{m} \gamma_{\mu} \Psi_{\rho} x^{m} c^{r} c^{z} tr_{z} (\tilde{\chi}^{m} \gamma^{\mu\rho} \lambda) - \frac{i}{2} x^{m} c^{r} c^{z} tr_{z} [(\tilde{\chi}^{m} \gamma^{\mu\rho} \lambda) (\tilde{\Psi}_{\mu} \gamma_{\rho} \lambda)] \\
- \frac{1}{8} v_{r} c^{r} c^{z} tr_{z} [(\tilde{\chi}^{m} \lambda) (\tilde{\chi}^{m} \lambda)] - \frac{3}{16} v_{r} c^{r} c^{z} tr_{z} [(\tilde{\chi}^{m} \gamma^{\mu\nu} \lambda) (\tilde{\chi}^{m} \gamma^{\mu\nu} \lambda)] \\
- \frac{1}{4} \frac{3}{4} v_{r} c^{r} c^{z} tr_{z} [(\tilde{\chi}^{m} \lambda) (\tilde{\chi}^{m} \lambda)] + \frac{1}{8} v_{r} c^{r} c^{z} tr_{z} [(\tilde{\chi}^{m} \gamma^{\mu\nu} \lambda) (\tilde{\chi}^{m} \gamma^{\mu\nu} \lambda)] \\
+ \frac{1}{4} (\tilde{\Psi}_{\mu} \gamma_{\nu} \Psi_{\rho}) (\tilde{\chi}^{m} \gamma^{\mu\rho} \lambda) - \frac{1}{2} v_{r} v_{s} c^{r} c^{z} tr_{z} [(\tilde{\chi}^{m} \gamma_{\mu} \lambda) (\tilde{\chi}^{m} \gamma_{\mu} \lambda)] \\
+ \frac{2}{2} c^{r} c^{z} tr_{z} [(\tilde{\chi}^{m} \gamma_{\mu} \lambda) (\tilde{\chi}^{m} \gamma_{\mu} \lambda)] , \quad (2.6)
while the variation of the two-cocycle with respect to vector gauge transformations gives the consistent

gives the supersymmetry anomaly

while the variation of $\mathcal{L}$ with respect to vector gauge transformations gives the consistent
gauge anomaly

The commutator of two supersymmetry transformations on the fermi fields closes on
the equations of (2.6), generating all local symmetry transformations, as well as the extra
two-cocycle

\[
\delta_{\text{extra}(\alpha)} \lambda = [\delta_1, \delta_2]_{\text{extra}(\alpha)} \lambda = \frac{c^x c^x}{v_s c^x} tr_z [\Lambda \delta_\mu A_\nu tr_z'(F_{\alpha\beta} F_{\gamma\delta})] .
\]

\[
\delta_{\text{extra}(\alpha)} \lambda \equiv [\delta_1, \delta_2]_{\text{extra}(\alpha)} \lambda = \frac{c^x c^x}{v_s c^x} tr_z [\Lambda \delta_\mu A_\nu tr_z'(F_{\alpha\beta} F_{\gamma\delta})] .
\]
for the gauginos. The presence of the arbitrary parameter $\alpha$ reflects the freedom of adding to the anomaly the variation of a local functional, consistently with all Wess-Zumino conditions. In six dimensions these close only on the field equations of the gaugini, and the two-cocyle grants the consistency of the construction for all values of $\alpha$ \cite{9}. It would be interesting to perform a cohomological analysis in superspace of this system along the lines of \cite{15}.

The gauge anomaly $A_{\Lambda} = \delta_{\Lambda} \mathcal{L}$ naturally satisfies the condition

$$A_{\Lambda} = -tr(\Lambda D_{\mu}J^{\mu}) \ ,$$

(2.11)

where $J^{\mu} = 0$ is the complete field equation of the vector field. One can similarly show that the supersymmetry anomaly is related to the field equation of the gravitino, that we write succinctly $J^{\mu} = 0$, according to

$$A_{\epsilon} = -(\epsilon D_{\mu}J^{\mu}) \ .$$

(2.12)

We would like to stress that the Noether identities (2.11) and (2.12) relate the anomalies to the equations of the fields whose transformations contain derivatives. This observation has a natural application to gravitational anomalies, that we would now like to elucidate. In fact, in analogy with the previous cases one would expect that

$$A_{\xi} = \delta_{\xi} \mathcal{L} = 2\xi_{\mu}D_{\nu}T^{\mu\nu} \ ,$$

(2.13)

where the variation of the metric under general coordinate transformations is

$$\delta g_{\mu\nu} = -\xi^{\alpha}\partial_{\alpha}g_{\mu\nu} - g_{\alpha\nu}\partial_{\mu}\xi^{\alpha} - g_{\mu\alpha}\partial_{\nu}\xi^{\alpha} \ .$$

(2.14)

Thus, for models without gravitational anomalies one would expect that the divergence of the energy-momentum tensor vanish. Actually, this is no longer true if other anomalies are present, since all fields, not only the metric, have derivative variations under coordinate
transformations. For instance, in a theory with gauge and supersymmetry anomalies, the gravitational anomaly is actually

$$A_\xi = \delta_\xi \mathcal{L} = 2\xi_\mu D_\mu T^{\mu\nu} + \xi_\nu tr(A^\nu D_\mu J^\mu) + \xi_\nu (\bar{\Psi}^\nu D_\mu J^\mu) .$$

In particular, in our case we are not accounting for gravitational anomalies, that would result in higher-derivative couplings, and indeed one can verify that the divergence of the energy-momentum tensor does not vanish, but satisfies the relation

$$D_\mu T^{\mu\nu} = -\frac{1}{2}tr(A^\nu D_\mu J^\mu) - \frac{1}{2}(\bar{\Psi}^\nu D_\mu J^\mu) .$$

3 Covariant field equations and covariant anomalies

It is well known that consistent and covariant gauge anomalies are related by the divergence of a local functional [11]. In six dimensions the residual covariant gauge anomaly is [4]

$$A_\Lambda^{cov} = \frac{1}{2}e^{\mu\nu\alpha\beta\gamma\delta} c_r^z c'_z tr_z (\Lambda F^\mu) tr_z' (F'_{\alpha\beta} F'_{\gamma\delta}) ,$$

and is related to the consistent anomaly by a local counterterm,

$$A_\Lambda^{cons} + tr[\Lambda D_\mu f^\mu] = A_\Lambda^{cov} ,$$

where

$$f^\mu = c_r^z c'_z \left\{ -\frac{1}{4} e^{\mu\nu\alpha\beta\gamma\delta} D_\nu tr_z (F'_{\alpha\beta} F'_{\gamma\delta}) - \frac{1}{6} e^{\mu\nu\alpha\beta\gamma\delta} F_{\nu\alpha} \omega'_{\beta\gamma\delta} \right\} .$$

Comparing eq. (3.3) with eq. (2.8) one can see that, to lowest order in the fermi fields,

$$A_\xi = tr(\delta_\xi A_\mu f^\mu) ,$$

and this implies that the transition from consistent to covariant anomalies turns a model with a supersymmetry anomaly into another without any [4, 10]. Indeed, six-dimensional supergravity coupled to vector and tensor multiplets was originally formulated in this fashion in [4] to lowest order in the fermi fields, extending the results of Romans [12].

2 The complete coupling to a single tensor multiplet, as well as to vector and hyper multiplets, was originally constructed in [16] for the special case of vanishing residual anomaly.
resulting vector equation is not integrable. Moreover, the corresponding gauge anomaly is not the gauge variation of a local functional and does not satisfy Wess-Zumino consistency conditions.

This result can be generalized naturally, if somewhat tediously, to include terms of all orders in the fermi fields. The complete supersymmetry anomaly of eq. (2.8) has the form

$$ A_\epsilon = tr(\delta_\epsilon A_\mu f^\mu) + \delta_\epsilon e_\mu^a g^\mu_a , \quad (3.5) $$

where to lowest order $f^\mu$ is defined in eq. (3.3). Modifying the vector equation so that

$$ (eq. A^\mu)_{(cov)} \equiv J^\mu_{(cov)} = \frac{\delta L}{\delta A_\mu} - f^\mu , \quad (3.6) $$

and similarly for the Einstein equation, the resulting theory is supersymmetric but no longer integrable. The covariant vector field equation is

$$ 2D_\nu(v_r F^{\mu\nu}) - 2G_r H^{\mu\nu\rho} F_{\nu\rho} - i \frac{\nu_r}{2} (\bar{\Psi}_\alpha \gamma^{\alpha\beta\mu\nu} \Psi_\beta) F_{\nu\rho} $$

$$ + i \nu_r (\bar{\chi} m \gamma^{\mu\nu} \chi m) F_{\nu\rho} - x_r^m (\bar{\Psi}_\alpha \gamma^{\mu\nu} \chi m) F_{\nu\rho} - i x_r^m x_r^m c^{sz'} tr_{s'} (\bar{\chi}' \gamma^{\mu\nu} \chi') F_{\nu\rho} $$

$$ + i \sqrt{2} D_\nu(v_r (\bar{\Psi}_\gamma \gamma^{\mu\nu} \gamma^{\rho}) \nu) + \sqrt{2} D_\nu(x_r^m (\bar{\chi} m \gamma^{\mu\nu} \chi m)) $$

$$ - i \frac{\nu}{2} F_{\nu\rho} c^{sz'} tr_{s'} (\bar{\chi}' \gamma^{\mu\nu} \chi') - i \frac{\nu}{2} c^{sz'} tr_{s'} [(\bar{\chi} \gamma^{\mu\nu} \chi') F_{\nu\rho}] - i c^{sz'} [(\bar{\chi} \gamma^{\mu\nu} \chi') F_{\nu\rho}] $$

$$ + \frac{1}{2 \sqrt{2}} c^{sz'} tr_{s'} [(\bar{\chi} \gamma^{\mu\nu} \gamma^{\rho}) (\bar{\chi}' \gamma^{s'} \nu)] $$

$$ + \frac{3i}{2 \sqrt{2}} (\bar{\lambda} \gamma^{\mu\nu} \lambda') (\bar{\lambda}' \gamma^{s'} \nu) + \frac{i}{4 \sqrt{2}} (\bar{\lambda} \gamma^{\mu\nu} \lambda') (\bar{\lambda}' \gamma^{s'} \nu) $$

$$ + \frac{i}{2 \sqrt{2}} (\bar{\lambda} \gamma^{\mu\nu} \lambda') (\bar{\lambda}' \gamma^{s'} \nu) $$

and completes the results in [4] to all orders in the fermi fields. Its divergence satisfies

$$ tr(\Delta D_\mu J^\mu_{(cov)} = - A^\mu_{(cov)} , \quad (3.8) $$
where $A_\Lambda^{\text{cov}}$ contains higher-order fermi terms:

$$A_\Lambda^{\text{cov}} = c^x c^\prime v x v x' \left\{ \frac{1}{2} \epsilon^{\mu\nu\alpha\beta\gamma\delta} (\Lambda F_{\mu\nu})(F'_{\alpha\beta} F'_{\gamma\delta}) + i e \Lambda F_{\nu\rho} (\bar{\chi} \gamma^{\rho\mu\nu} D_{\mu} \chi') + \frac{i e}{2} \Lambda D_{\mu} ((\bar{\chi} \gamma^{\rho\mu}) F'_{\nu\rho}) + i e \Lambda D_{\mu} \left[ (\bar{\chi} \gamma^{\rho\mu}) F'_{\nu\mu} \right] - \frac{e}{2\sqrt{2}} \Lambda D_{\mu} \left[ ((\bar{\chi} \gamma^{\mu\nu}) \bar{\chi} \gamma^{\rho\mu}) (\bar{\chi}' \gamma^\rho \Psi_{\rho}) \right] + e \Lambda D_{\mu} \left\{ \frac{1}{2} c^x c^\prime v x v x' \left[ - \frac{3i}{2\sqrt{2}} (\bar{\chi} \gamma^{\mu}) (\bar{\chi}') (\bar{\chi}' \gamma^m) - \frac{i}{4\sqrt{2}} (\bar{\chi} \gamma^{\mu\nu} \chi')(\bar{\chi}' \gamma_{\nu\rho} \chi^m) \right] - \frac{i}{2\sqrt{2}} (\bar{\chi} \gamma_{\nu} \chi')(\bar{\chi}' \gamma^m) \right\} \right\} .$$

Finally, one can study the divergence of the Rarita-Schwinger and Einstein equations in the covariant model. To this end, let us begin by stating that the derivation of Noether identities for a system of non-integrable equations does not present difficulties of principle, since these involve only first variations. Indeed, the only difference with respect to the standard case of integrable equations is that now $\delta \mathcal{L}$ is not an exact differential in field space. Still, all invariance principles reflect themselves in linear dependencies of the field equations. Thus, for instance, with the covariant equations obtained from the consistent ones by the redefinition of eq. (3.6) and by

$$\left( \text{eq. } e^\mu_a \right)_{(\text{cov})} = \frac{\delta \mathcal{L}}{\delta e^\mu_a} - g^\mu_a ;$$

the total $\delta \mathcal{L}$ vanishes by construction. The usual procedure then proves that the divergence of the Rarita-Schwinger equation vanishes for any value of the parameter $\alpha$. On the other hand, the divergence of the energy-momentum tensor presents some subtleties that we would now like to describe. In particular, it vanishes to lowest order in the fermi couplings, while it gives a covariant non-vanishing result if all fermion couplings are taken into account. The subtlety has to do with the transformation of the vector under general
coordinate transformations,
\[
\delta_\xi A_\mu = -\xi^\alpha \partial_\alpha A_\mu - \partial_\mu \xi^\alpha A_\alpha ,
\] (3.11)
and with the corresponding full (off-shell) form of the identity of eq. (2.15). Starting
again from the consistent equations, one finds
\[
A_\xi = \delta_\xi \mathcal{L} = 2\xi_\nu D_\mu T^{\mu\nu} + \xi_\nu tr(A_\nu D_\mu J_\mu) + \xi_\nu tr(F^{\mu\nu} J_\mu) + \xi_\nu (\bar{\Psi}^\nu D_\mu J_\mu) .
\] (3.12)
Reverting to the covariant form eliminates the divergence of the Rarita-Schwinger equa-
tion and alters the vector equation, so that the third term has to be retained. The final
result is then
\[
D_\mu T^{\mu\nu}_{(cov)} = -\frac{1}{2} tr(A_\nu D_\mu J_\mu_{(cov)}) - \frac{1}{2} tr(f_\mu F^{\mu\nu}) - \frac{1}{2} tr(A_\nu D_\mu J_\mu) - \frac{1}{2} e^{\nu a} D_\mu g^a ,
\] (3.13)
and is nicely verified by our equations. In particular, this implies that, to lowest order in
the fermi couplings, the divergence of \( T^{\mu\nu}_{(cov)} \) vanishes.

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