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ANOMALIES IN SUPERSYMMETRIC THEORIES

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ABSTRACT

The anomaly situation in supersymmetric theories is reviewed. The trace, chiral and supercurrent anomalies are discussed both in globally supersymmetric models and in models with local supersymmetry invariance.

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It was pointed out some time ago by Ferrara and Zumino¹⁾ that in a supersymmetric theory the axial current j_μ^5 , the improved energy-momentum tensor $\theta_{\mu\nu}$ and the improved supersymmetry current S_μ can be identified with the components of a superfield V_μ , so that an intimate relation exists between them. At the classical level this superfield satisfies certain relations reflecting the conservation and trace relations (modulo mass terms)

$$\partial_\mu S^\mu = \partial_\mu \theta^{\mu\nu} = 0 \quad (1)$$

$$\partial_\mu j^{\mu 5} = 0 \quad (2a)$$

$$\theta_\mu^\mu = \partial_\mu (x_\nu \theta^{\mu\nu}) = 0 \quad (2b)$$

$$\gamma^\mu S_\mu = \partial_\mu (\gamma \cdot x S^\mu) = 0 \quad (2c)$$

We note that the improved supersymmetry current is obtained by adding terms to the Noether current so as to satisfy Eq. (2c) while maintaining Eq. (1)¹⁾.

At the quantum level the matrix elements of j_μ^5 , $\theta_{\mu\nu}$, S_μ no longer satisfy Eqs. (2). Instead, they become equal to expressions which can be identified with matrix elements of operators constructed out of the basic fields. These are the anomalies (chiral, trace or supercurrent). We summarize in this lecture what is known about these anomalies and what their existence might imply, both in models of rigid (global) supersymmetry and in models containing supergravity fields.

For the scalar or vector multiplet of rigid supersymmetry, Eqs. (2) are associated with invariance of the models under conformal supersymmetry²⁾ with parameter $\alpha = \alpha(x) = \alpha_0 + \gamma \cdot x \alpha_1$ (α_0, α_1 are constant Majorana parameters). For instance, the massless scalar multiplet model (with weight $n = \frac{1}{2}$ in the notation of Ref. 2)) is invariant under

$$\delta A = i \bar{\alpha} \chi, \quad \delta B = i \bar{\alpha} \gamma_5 \chi$$

$$\delta \chi = \partial_\mu (A - \gamma_5 B) \gamma^\mu \alpha + (A - \gamma_5 B) \gamma^\mu \partial_\mu \alpha + (F + \gamma_5 G) \alpha, \quad (3)$$

$$\delta F = i \bar{\alpha} \gamma^\mu \partial_\mu \chi, \quad \delta G = i \bar{\alpha} \gamma_5 \gamma^\mu \partial_\mu \chi$$

with α given above. Furthermore, the commutator of two such supersymmetry transformations acting on the fields gives

$$\begin{aligned} [\delta_2, \delta_1] A &= \xi^\mu \partial_\mu A + \frac{1}{4} (\partial_\mu \xi^\mu) A + \frac{1}{2} \gamma B \\ [\delta_2, \delta_1] B &= \xi^\mu \partial_\mu B + \frac{1}{4} (\partial_\mu \xi^\mu) B - \frac{1}{2} \gamma A \\ [\delta_2, \delta_1] \chi &= \xi^\mu \partial_\mu \chi + \frac{3}{8} (\partial_\mu \xi^\mu) \chi + \frac{1}{2} \partial_\mu \xi_\nu \sigma^{\mu\nu} \chi - \frac{1}{4} \gamma \xi \chi \end{aligned} \quad (4)$$

with

$$\xi_\mu = 2i\bar{\alpha}_1 \gamma_\mu \alpha_2, \quad \eta = i[\partial_\mu \bar{\alpha}_1 \gamma_5 \gamma_\mu \alpha_2 - 1 \leftrightarrow 2] \quad (5)$$

Thus the model is invariant under a set of transformations containing rigid supersymmetry, conformal supersymmetry, conformal transformations and chiral transformations ($A \rightarrow B$, $B \rightarrow -A$, $\chi \rightarrow -\frac{1}{2}\gamma_5\chi$). The generators $D_\mu = x^\nu \theta_{\mu\nu}$, $I_\mu = \gamma \cdot x S_\mu$, j_μ^5 go into each other under the action of the algebra, in particular under rigid supersymmetry transformations.

We note for future reference that in the case of the vector multiplet (with $n = 1$) the commutator of two supersymmetry transformations has the following effect on the spinor field²⁾

$$\begin{aligned} [\delta_2, \delta_1] \lambda &= \xi^\mu \partial_\mu \lambda + \frac{3}{8} (\partial_\mu \xi^\mu) \lambda \\ &+ \frac{1}{2} \partial_\mu \xi_\nu \sigma^{\mu\nu} \lambda + \frac{3}{4} \eta \gamma_5 \lambda \end{aligned} \quad (6)$$

Thus the chiral charge of the spinor is -3 times that of the spinor in the scalar multiplet. (These chiral charge assignments can also be read off from the coupling of these fields to the axial vector auxiliary field A_μ of supergravity³⁾).

The anomalies in these theories are associated with a breaking of the (super-) conformal part of the algebra, with ordinary rigid supersymmetry remaining unbroken. A logical way to study the anomalies is to examine quantum modifications of the Ward identities for the superfield current V_μ associated with the superconformal transformations and identify the anomaly superfield. This would have the advantage of including all the anomalies at once, and would show that they are related to each other as the components of a superfield. This procedure has been carried out by Piguet and Schweda, and Clark, Piguet and Siebel⁴⁾ for the case of the self-

interacting scalar multiplet, in the framework of BPHZ supersymmetric renormalization. It has not been carried out for any other system (vector multiplet, scalar or vector multiplet in interaction with supergravity fields) where a pedestrian approach which simply calculates the (one-loop) anomalies and compares them has been used. We shall now give the results of such calculations for three examples.

Example 1 : Scalar multiplet in a background vector multiplet⁵⁾

$$\begin{aligned} \mathcal{L} = & -\frac{1}{2} (D_\mu A)^2 - \frac{1}{2} (D_\mu B)^2 - \frac{i}{2} \bar{\chi}^\alpha \gamma^\mu (D_\mu \chi)^\alpha \\ & - \frac{1}{2} g^2 [A^2 B^2 - (A^\alpha B^\alpha)^2] \\ & - g f^{abc} \bar{\lambda}^a (A^b + \gamma_5 B^b) \chi^c \\ & - \left[\frac{1}{4} v_{\mu\nu}^2 + \frac{i}{2} \bar{\lambda}^\alpha \gamma^\mu (D_\mu \lambda)^\alpha \right] \end{aligned} \quad (7)$$

where (λ^a, v_μ^a) form a Yang-Mills supersymmetric multiplet and D_μ is the Yang-Mills covariant derivative. The last two terms may be omitted if one wishes to think of (λ^a, v_μ^a) as being arbitrary background fields transforming into each other as vector multiplet fields.

The quantities of interest are

$$\begin{aligned} S_\mu = & \gamma^\lambda \gamma_\mu [D_\lambda (A + \gamma_5 B)]^\alpha \chi^\alpha \\ & + \frac{4}{3} \sigma_{\mu\lambda} \partial^\lambda [(A^\alpha + \gamma_5 B^\alpha) \chi^\alpha] \end{aligned} \quad (8a)$$

$$\begin{aligned} \Theta_{\mu\nu} = & D_\mu A D_\nu A + D_\mu B D_\nu B + \frac{i}{4} \bar{\chi} (\gamma_\mu D_\nu + \gamma_\nu D_\mu) \chi \\ & - \eta_{\mu\nu} \mathcal{L} - \frac{1}{6} (\partial_\mu \partial_\nu - \eta_{\mu\nu} \square) (A^2 + B^2) \end{aligned} \quad (8b)$$

$$d_\mu^5 = A \overleftrightarrow{D}_\mu B - \frac{i}{4} \bar{\chi} \gamma_5 \gamma_\mu \chi \quad (8c)$$

The form of the chiral current follows from the chiral properties implicit in Eqs. (4).

We look at one-loop matrix elements of the above quantities between the vacuum and states of the vector multiplet (λ, v_μ) . The anomaly of j_μ^5 is the usual chiral anomaly⁶⁾, as obtained from the diagram of Fig. 1a. Only the spinor part of the current contributes.

The spinor current anomaly can be calculated from the diagrams in Figs. (1b), (1c) (1d)⁷⁾, while the trace anomaly is known also⁵⁾. One finds, with the external fields on shell:

Fig. 1 : Diagrams for one-loop chiral and supercurrent anomalies.

$$a: \quad c \, 2 \bar{\lambda} \lambda$$

$$b: \quad c \, 2 \bar{\lambda} \gamma_5 \lambda$$

$$x: \quad \gamma^\mu S_\mu = c \, 4 v_{\mu\nu} \sigma^{\mu\nu} \lambda \quad (9)$$

$$f: \quad \theta_\mu^\mu = c \, v_{\mu\nu} v^{\mu\nu}$$

$$g: \quad \partial_\mu J^{\mu 5} = -c \, \tilde{v}_{\mu\nu} v^{\mu\nu}$$

$$c = - \frac{g^2}{32\pi^2} c_v \quad (10)$$

with c_v the Casimir operator for the Yang-Mills group (for $O(3)$ $c_v = 2$). We have adjoined to the three anomalies the quantities a, b . Together they transform into each other like the components of a scalar multiplet (a, b, x, f, g) .

Example 2 : Yang-Mills vector multiplet⁸⁾

$$\mathcal{L} = -\frac{1}{4} v_{\mu\nu}^2 - \frac{i}{2} \bar{\lambda} \gamma \cdot D \lambda \quad (11)$$

$$S_\mu = -i \sigma^{\lambda\nu} \gamma_\mu \lambda^a v_{\lambda\nu}^a \quad (12a)$$

$$\theta_{\mu\nu} = v_{\mu\lambda} v_\nu^\lambda - \frac{i}{4} \bar{\lambda} (\gamma_\mu D_\nu + \gamma_\nu D_\mu) \lambda - \gamma_{\mu\nu} \mathcal{L} \quad (12b)$$

$$\partial_\mu^5 = \frac{3i}{4} \bar{\lambda}^a \gamma_5 \gamma_\mu \lambda^a \quad (12c)$$

The coefficient in j_μ^5 follow from the chiral properties implicit in Eq. (6). The trace and chiral anomalies are standard (the latter being three times the Adler anomaly because of the factor of 3 in j_μ^5), while the supercurrent anomaly has been calculated by a number of authors⁹⁾. One finds the same results as in Eq. (9) with $c \rightarrow -3c$.

Three comments are in order: As is well known, in the case of the chiral anomaly a new current can be defined whose divergence is zero:

$$\tilde{j}_\mu^5 = j_\mu^5 + \frac{3g^2 c_V}{32\pi^2} \epsilon_{\mu\nu\rho\sigma} v^{\nu\rho} v^\sigma \quad (13)$$

However, this current is not gauge invariant. Similarly one can define a new supercurrent

$$\tilde{S}_\mu = S_\mu - \frac{3g^2 c_V}{8\pi^2} v_{\rho\mu} \gamma^\rho \lambda \quad (14)$$

which satisfies $\gamma \cdot \tilde{S} = 0$, but $\partial_\mu \tilde{S}^\mu \neq 0$. In fact in an explicit calculation one will find either $\gamma \cdot S = 0$ or $\partial \cdot S = 0$ depending on what kind of regularization scheme one uses. For instance, dimensional regularization leads to $\partial \cdot S = 0$ but $\gamma \cdot S \neq 0$, while a point-splitting calculation has given $\gamma \cdot S = 0$ but $\partial \cdot S \neq 0$ ¹⁰⁾. In general it seems preferable to maintain the rigid supersymmetry relation $\partial \cdot S = 0$ and give up the $\gamma \cdot S = 0$ condition.

Our second comment relates to higher order effects. The Adler-Bardeen theorem asserts that the anomaly in $\partial_\mu j_\mu^5$ receives no contributions from higher orders. It is known that this is not the case

for θ_μ^μ . (The β function is not zero in higher order for this model.) We must conclude therefore that either the multiplet pattern of Eq. (9) does not hold in higher order or that in a calculation that preserved rigid supersymmetry at each step the Adler-Bardeen or β function result (or both) would be modified. This is suggested by our next example.

Our third comment stems from the observation that for a system consisting of three scalar multiplets and one vector multiplet interacting as in the above examples the one-loop anomalies cancel¹¹⁾. Furthermore, it has been shown¹²⁾ that also the two-loop β function (trace anomaly) is zero for a model, equivalent to the above one at the one-loop level, but with additional interactions which make it $O(4)$ symmetric.

Example 3 : Self-interacting scalar multiplet¹³⁾.

$$\begin{aligned} \mathcal{L} = & -\frac{1}{2} (\partial_\mu A)^2 - \frac{1}{2} (\partial_\mu B)^2 - \frac{i}{2} \bar{\chi} \gamma \cdot \partial \chi \\ & + \frac{1}{2} F^2 + \frac{1}{2} G^2 + g [FA^2 - FB^2 + 2GAB - i\bar{\chi}(A - \gamma_5 B)\chi] \end{aligned} \quad (15)$$

This system is interesting because naïvely one would expect $\partial_\mu j_\mu^5 = 0$ while $\theta_\mu^\mu \neq 0$ at the one-loop level. This is because, according to the usual lore, external vector mesons are needed to get an anomaly, while here we are taking matrix elements to external spin zero or one-half particles. However, the actual result is regularization dependent. Using a supersymmetric regularization procedure (e.g., Pauli-Villars), one finds¹⁴⁾ at the one-loop level:

$$a : \quad d [AF - BG]$$

$$b : \quad d [AG + BF]$$

$$x : \quad \gamma \cdot S = d [(A - \gamma_5 B) \gamma \cdot \partial \chi + (F - \gamma_5 G) \chi]$$

$$f : \quad \theta_\mu^{\mu} = d [A \not{\partial} B + B \not{\partial} A - i \bar{\chi} \gamma \cdot \partial \chi + F^2 + G^2]$$

$$g : \quad \partial_\mu j^{\mu} = d [\partial_\mu (A \overleftrightarrow{\partial}_\mu B + \frac{i}{2} \bar{\chi} \gamma_5 \gamma_\mu \chi)]$$

$$d = - \frac{g^2}{\pi^2} \quad (16)$$

Unlike the Adler anomaly case, one could here redefine the current j_μ^5 by subtracting the quantity $d(A \overleftrightarrow{\partial}_\mu B + i/2 \bar{\psi} \gamma_5 \gamma_\mu \psi)$, so that the new current is anomaly free, since no gauge invariance requirements exist. However, it turns out that the resulting current $\tilde{j}_\mu^5$ would not be a finite operator (Green's functions involving $\tilde{j}_\mu^5$ would be infinite even after renormalization), so that this is not a desirable procedure. The above results, to all orders, have also been obtained in Ref. 4) using supersymmetric BPHZ renormalization techniques.

Thus for the self-interacting scalar multiplet the anomaly situation seems quite clear. It is possible, by insisting on rigid supersymmetry, to define quantum currents such that the anomalies transform as members of a scalar multiplet to all orders. For the vector multiplet, where (Yang-Mills) gauge invariance must also be imposed, the higher loop situation remains unclear. As we

mentioned above, if the anomalies are to form a multiplet either the β function results or the Adler-Bardeen theorem must be changed in a completely supersymmetric calculation. Otherwise, one should understand why the higher-order calculations do not lead to supersymmetric results.

We turn now to the anomalies of the scalar or vector multiplets in the presence of external supergravity fields. Explicit calculations have been carried out for the supercurrent¹⁵⁾ while the trace and chiral anomalies can be obtained from known results, which are summarized in the following table given by Christensen and Duff¹⁶⁾.

Table

Spin	$180(4\pi)^2 \theta_\mu^\mu =$	$180(4\pi)^2 \partial_\mu j^{\mu 5} =$
0	$R_{\mu\nu\alpha\beta} R^{\mu\nu\alpha\beta}$	$0 \times \tilde{R}_{\mu\nu\alpha\beta} R^{\mu\nu\alpha\beta}$
$\frac{1}{2}$	$\frac{7}{4} \times \quad "$	$\frac{-15}{2} \times \quad "$
1	$-13 \quad "$	$0 \quad "$
$\frac{3}{2}$	$\frac{-233}{4} \quad "$	$\frac{315}{2} \quad "$
2	$212 \quad "$	$0 \quad "$

(here j_μ^5 is defined with unit chiral charge). For instance, the trace anomaly for the scalar multiplet is given by twice the spin zero contribution plus the spin $\frac{1}{2}$ contribution leading to a coefficient of $15/4$, which matches that of the chiral anomaly. Similarly for the vector multiplet the sum of spinor and vector contributions gives $-45/4$, to be compared to $(-3) \times (-15/2)$ for the chiral

anomaly (remembering that the spinor has chiral charge -3). We find, for the scalar multiplet, with external supergravity fields on shell

$$\begin{aligned}
 a: & \quad c \bar{\Psi}_{\nu, \alpha} (\Psi_{\nu, \alpha} - \Psi_{\alpha, \nu}) + \dots \\
 b: & \quad c \bar{\Psi}_{\nu, \alpha} \gamma_5 (\Psi_{\nu, \alpha} - \Psi_{\alpha, \nu}) + \dots \\
 x: & \quad \gamma \cdot S = c R^{\alpha \nu \beta \lambda} \sigma_{\beta \lambda} \Psi_{\nu, \alpha} + \dots \\
 f: & \quad \theta_{\lambda}^{\lambda} = c R^{\alpha \nu \beta \lambda} R_{\alpha \nu \beta \lambda} + \dots \\
 g: & \quad \partial_{\mu} d^{SM} = c R^{\alpha \nu \beta \lambda} \tilde{R}_{\alpha \nu \beta \lambda} + \dots
 \end{aligned} \tag{17}$$

with

$$c = \frac{\chi^2}{768\pi^2}$$

For the vector multiplet one obtains the same results times a factor of -3.

Finally we turn to the anomalies of the supergravity multiplet³⁾ itself. The trace and chiral anomalies can be read from the above table and give factors of 615/4 (the sum of spin 3/2 and spin 2 trace anomalies) and 945/4, respectively. (We note that the spin 3/2 field carries a chiral charge -3, just like the spinor of the vector multiplet. This can be seen most easily from its coupling to the axial vector auxiliary field. In fact the fermion

of the scalar multiplet is exceptional in carrying unit chiral charge, apparently because it belongs to a chiral superfield.) The anomaly of the spinor current has not yet been calculated but the numbers above are hard to reconcile with the supposition that the anomalies should belong to a multiplet. A priori there is no reason why they should; ordinary Poincaré supergravity is not (super-) conformally invariant and the currents do not satisfy conservation equations. A study of the quantum Ward identities is hampered by the off-shell non-renormalizability of the theory. The situation may be different in conformal supergravity³⁾ but the anomaly situation has not been investigated there.

Finally, we discuss some implications of the anomalies. First of all, one can look at their topological significance. For the chiral anomalies this is well known, and instanton effects appear to be present in supersymmetric theories¹⁹⁾. The trace anomalies have similar well-known significance. In this context it is interesting to look at models with vanishing trace anomalies. We have already mentioned the $O(4)$ supersymmetric Yang-Mills theory which has a vanishing one- and two-loop β function. The same model clearly has vanishing (one-loop) anomalies when coupled to external supergravity fields. Finally, it is easy to check from the table above that in $O(3)$ supergravity (one spin 2, 3 spin $3/2$, 3 spin 1 and one spin $\frac{1}{2}$) the trace anomaly is zero¹⁶⁾¹⁸⁾.

At the present time it is unclear what topological significance, if any, the supercurrent anomaly might have either in ordinary space-time, or as part of the anomaly superfield in superspace.

In the context of conventional field theory, one has to worry about the effect of anomalies on renormalization. In rigid supersymmetry the currents with anomalous properties do not appear in any couplings so the anomalies have no effect there. In supergravity the same is the case (with one exception that we discuss below). While one cannot talk about renormalizability in the

conventional sense, an anomaly in the $\partial_\mu S^\mu = 0$ equation would destroy rigid supersymmetry and affect finiteness arguments beyond one-loop. If the anomaly is present only in the $\gamma \cdot S = 0$ equation it presumably has no effect.

The anomalies seem to play a curious role in two models of extended supergravity with $SU(4)$ invariance¹⁹⁾. One of the models is $SO(4)$ supergravity²⁰⁾ in which the action is $SO(4)$ invariant and only the equations of motion exhibit $SU(4)$ invariance. In the other one²¹⁾, recently constructed, the action itself is $SU(4)$ invariant.

At the classical level the two theories are equivalent in the sense that their equations of motion can be transformed into each other by a combination of point field and duality transformations. At the quantum level the equivalence is not clear. In fact, in the $SU(4)$ theory decay of the pseudoscalar meson of the model into gravitons appears possible because it couples to the divergence of chiral currents of the spin $\frac{1}{2}$ and spin $\frac{3}{2}$ fields, and also to a particular combination of spin 1 fields. In the $SO(4)$ model the couplings do not permit such decays. Remarkably enough, the strength of the couplings and the numerical values of the chiral anomalies are such that the total decay amplitude is zero in $SU(4)$ theory. Thus, at least for this process "anomalies" are necessary in order to maintain the equivalence between $SO(4)$ and $SU(4)$ supergravity¹⁹⁾.

In $SU(4)$ supergravity there are also couplings of scalars and spin $\frac{1}{2}$ fermions to objects related to θ_μ^μ and $\gamma \cdot S$. The precise role these couplings play is under investigation.

We conclude with the following observation. For the case of the scalar or vector multiplets, the fact that the anomalies themselves form multiplets could have been used to obtain one from the other, if the numbers were not already known. To a certain extent also, rigid supersymmetry "explains" the relative size of the

numbers 1, $7/4$, -13 and $-15/2$ in the table. ($2 \times 1 + 7/4$ must match $1/2 \times 15/2$, and $7/4 - 13$ must match $3/2 \times 15/2$.)

Another relation, between the Adler anomaly and the spin $\frac{1}{2}$ axial anomaly in the presence of gravitons is revealed by examining the situation for the coupling of the vector multiplet to off-shell external supergravity, including the auxiliary fields³⁾. In this case it is best to think of the external supergravity fields as being fields of conformal supergravity. The couplings are the same but the system is also invariant under chiral transformations of the spin $\frac{1}{2}$ field, together with gauge transformations on the external axial vector field A_μ (and chiral transformations of ψ_μ). The anomaly $\partial_\mu j^{\mu 5}$ must be proportional to the G component of a scalar multiplet and it is known³⁾ that the gravitational and axial fields A_μ must appear in the combination

$$c \left[\tilde{R}_{\mu\nu\alpha\beta} R^{\mu\nu\alpha\beta} - \frac{1}{6} \tilde{F}_{\mu\nu}(A) F^{\mu\nu}(A) \right]$$

But c and $c/6$ is what one would compute by conventional Feynman diagrams, with two external gravitons, or spin one mesons, respectively (the fact that A_μ couples with $\gamma_5 \gamma_\mu$ rather than γ_μ simply reverses the sign compared to the usual Adler anomaly). Taking into account the unusual normalization of the field A_μ , and the strength of its coupling to the spinor field, one obtains a result for the ratio of Adler to gravitational anomaly which agrees with the conventional values.

Thus, if one were to rewrite history, one could argue that a knowledge of the Adler anomaly, together with rigid and local supersymmetry is all that is required in order to compute almost all the one-loop anomalies of matter fields (spin 0, $\frac{1}{2}$, 1).

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