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Supersymmetry Anomalies: Further Results

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ABSTRACT

We consider anomalies in local supersymmetry. An expression for the gauge(Lorentz) noninvariant part of this anomaly for $2n$ dimensional spacetime is given. The result agrees with explicit calculations in two dimensions. Complete expressions for rigid supersymmetry anomalies in two, four and six dimensions are also given.

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In a previous paper [1], we discussed supersymmetry anomalies for $N = 1$ supersymmetric theories. The analysis was done in the component language or the Wess-Zumino gauge to facilitate the discussion of theories where no superfield formalism is known to exist (such as would occur in dimensions higher than four). It was shown how gauge and gravitational anomalies generically give rise to supersymmetry anomalies. The expression for the supersymmetry anomaly separates naturally into a gauge noninvariant part and a gauge invariant part. We gave an expression for the gauge noninvariant part for rigid supersymmetry in arbitrary even dimensional spacetime and the invariant part in four dimensions. In this paper, we report on further calculations along the same lines *. First of all, we discuss the case of local supersymmetry which was only briefly touched upon in reference 1. As in the case of rigid supersymmetry, a general expression for the noninvariant part is derived. The invariant part depends specifically on the spacetime dimensions. For the case of two dimensions, the invariant part is zero and we reproduce known results on supersymmetry anomalies [3]. We also consider supersymmetric Yang-Mills theory in two and six dimensions. The invariant part is computed for these dimensions; this, alongwith the results in reference 1, completes the expression for rigid supersymmtry anomaly in two, four and six dimensions. We conclude the paper with a few remarks on ten dimensional theories.

Local Supersymmetry

We start our discussion by writing down the consistency conditions on the anomalies for local supersymmetry. The consistency conditions are the infinitesimal version of the group composition law of the transformations, which are, in this case, supersymmetry, Lorentz and general coordinate transformations. Consider the commutation rule for two supersymmetry transformations with parameters ϵ and ϵ' [4]:

$$[\delta_S(\epsilon), \delta_S(\epsilon')] = \delta_G(\xi) + \delta_L(\xi\omega) + \delta_S(-\xi\psi). \tag{1}$$

$\delta_G(\xi)$ denotes a general coordinate transformation with the parameter $\xi_\mu = 2(\overline{\epsilon'}\gamma_\mu\epsilon - \bar{\epsilon}\gamma_\mu\epsilon')$, $\delta_L(\xi\omega)$ is a Lorentz transformation with field dependent parameter $\xi_\mu\omega^\mu_{ab}$ and $\delta_S(-\xi\psi)$

* There are a number of papers dealing with anomalies in supersymmetric theories using superfields [2]. Comparison to our approach is made in [1].

is again a supersymmetry transformation with parameter $\xi_\mu \psi^\mu$. The ω_μ^{ab} is the spin connection and ψ_μ denotes the gravitino field. Denoting the anomalies by $S(\epsilon), G_L(\lambda)$ and $G_G(\xi)$, the consistency condition implied by the commutation rule (1) is

$$\delta_S(\epsilon)S(\epsilon') - \delta_S(\epsilon')S(\epsilon) = G_G(\xi) + G_L(\xi\omega) + S(-\xi\psi). \tag{2}$$

It is possible to write down local counterterms in terms of the frame field e_μ^a such that there is no anomaly in general coordinate transformations [5]. The effect of the gravitational anomaly is completely captured in the anomaly of local Lorentz transformations. (Other choices are possible but this is the most convenient for our purposes.). The consistency condition (2) thus becomes simply:

$$\delta_S(\epsilon)S(\epsilon') - \delta_S(\epsilon')S(\epsilon) = G_L(\xi\omega) + S(-\xi\psi). \tag{3}$$

Compared to the case of rigid supersymmetry, this condition has an extra term viz. $S(-\xi\psi)$.

The commutation rule

$$[\delta_L(\lambda), \delta_S(\epsilon)] = \delta_S(\frac{1}{2}\lambda^{ab}\sigma_{ab}\epsilon) \tag{4}$$

leads to

$$\delta_L(\lambda)S(\epsilon) - \delta_S(\epsilon)G_L(\lambda) = S(\frac{1}{2}\lambda\sigma\epsilon). \tag{5}$$

$\frac{1}{2}\lambda\sigma\epsilon$ is an infinitesimal Lorentz transformation of ϵ; being only a parameter and not a dynamical field, the variation of ϵ is not included in the definition of $\delta_L(\lambda)$. Let us define a total variation by Δ, i.e. $\Delta = \delta + \delta'$ where δ' denotes variation of the parameter. Equation (5) can then be written as

$$\Delta_L(\lambda)S(\epsilon) - \delta_S(\epsilon)G_L(\lambda) = 0. \tag{6}$$

Taking care of the other commutators, the consistency conditions can be gathered together as

$$\Delta_G(\xi)S(\epsilon) = 0 \tag{7a}$$

$$\Delta_L(\lambda)S(\epsilon) - \delta_S(\epsilon)G_L(\lambda) = 0 \tag{7b}$$

$$\delta_S(\epsilon)S(\epsilon') - \delta_S(\epsilon')S(\epsilon) = G_L(\xi\omega) + S(-\xi\psi). \tag{7c}$$

The first condition viz.(7a) simply states that $S(\epsilon)$ is coordinate invariant. To solve (7b) for $2n$ dimensional spacetime, we start off with the expression for the index of the Dirac operator in $2n+2$ dimensional spacetime. The index density $\hat{A}_{2n+2}$ is a homogeneous polynomial of traces of the curvature R and is given by the appropriate term in the expansion of the $\hat{A}$-genus [6], that is, $c_{\nu_1,\ldots,\nu_i,\ldots}$ below are determined by the $\hat{A}$ -genus.

$$\hat{A}_{2n+2} = \sum_{\nu_i \geq 0, \sum \nu = \frac{n+1}{2}} c_{\nu_1,\ldots,\nu_i,\ldots}(trR^{2\nu_1})\ldots(trR^{2\nu_i})\ldots. \tag{8}$$

Although we can write down $\hat{A}$-genus in $2n+2$ dimensions, it vanishes except in $4k$ dimensions with k as an integer. We define Ω_{2n+1} by

$$d\Omega_{2n+1} = \hat{A}_{2n+2}. \tag{9}$$

A Lorentz variation of Ω_{2n+1} gives

$$\delta_L(\lambda)\Omega_{2n+1} = dg_L(\lambda), \tag{10}$$

where $g_L(\lambda)$ is a $2n$-form. Integration of $g_L(\lambda)$ over the $2n$-dimensional spacetime gives the Lorentz anomaly $G_L(\lambda)$ [5,7]. (Appropriate fall off conditions on ϵ, λ and the fields are assumed so that spacetime is, effectively, S^{2n}).

A supersymmetry variation of Ω_{2n+1} can be written as

$$\begin{aligned}\delta_S(\epsilon)\Omega_{2n+1} &= (dl_\epsilon + l_\epsilon d)\Omega_{2n+1} \\ &= ds(\epsilon) + l_\epsilon d\Omega_{2n+1} \\ &= ds(\epsilon) + l_\epsilon \hat{A}_{2n+2},\end{aligned} \tag{11}$$

where $l_\epsilon R = \delta_S(\epsilon)\omega$, $l_\epsilon\omega = 0$ [5] and $s(\epsilon) = l_\epsilon\Omega_{2n+1}$. Making a Lorentz transformation on equation (11), we obtain:

$$\Delta_L(\lambda)\delta_S(\epsilon)\Omega_{2n+1} = d(\Delta_L(\lambda)s(\epsilon)) + \Delta_L(\lambda)l_\epsilon\hat{A}_{2n+2}. \tag{12}$$

$l_\epsilon\hat{A}_{2n+2}$ involves traces of products of $\delta_S(\epsilon)\omega$ and R and, since these are Lorentz covariant, $\Delta_L(\lambda)l_\epsilon\hat{A}_{2n+2} = 0$. Also we have

$$\Delta_L(\lambda)\delta_S(\epsilon)\Omega_{2n+1} = \delta_S(\epsilon)\delta_L(\lambda)\Omega_{2n+1}$$

$$= \delta_S(\epsilon) d g_L(\lambda) = d(\delta_S(\epsilon) g_L(\lambda)). \tag{13}$$

The first step follows from $\Delta_L(\lambda)\delta_S(\epsilon) - \delta_S(\epsilon)\delta_L(\lambda) = 0$ and we use (10) in the second step. Using (13), equation (12) becomes

$$d(\delta_S(\epsilon) g_L(\lambda) - \Delta_L(\lambda) s(\epsilon)) = 0. \tag{14}$$

Integrating this over a $2n+1$ dimensional disc whose boundary is spacetime S^{2n}, we get

$$\delta_S(\epsilon) G_L(\lambda) - \Delta(\lambda) \int_{S^{2n}} s(\epsilon) = 0. \tag{15}$$

This equation tells us that (7b) can be solved by:

$$S(\epsilon) = S_0(\epsilon) + S_{inv}(\epsilon)$$

$$S_0(\epsilon) = \int_{S^{2n}} s(\epsilon) = \int_{S^{2n}} l_\epsilon \Omega_{2n+1}, \tag{16}$$

and $S_{inv}(\epsilon)$ is invariant under local Lorentz transformations. $S_{inv}(\epsilon)$ has to be determined by (7c), to which we now turn. Using $\delta_S(\epsilon) = dl_\epsilon + l_\epsilon d$, we can write

$$(\delta_S(\eta)\delta_S(\epsilon) - \delta_S(\epsilon)\delta_S(\eta))\Omega_{2n+1} = d[\delta_S(\eta)s(\epsilon) - \delta_S(\epsilon)s(\eta)] + (\delta_S(\eta) l_\epsilon A_{2n+2} - \delta_S(\epsilon) l_\eta A_{2n+2}). \tag{17}$$

From the commutation rule (1) (and remembering that differential forms are invariant under coordinate tansformations) the left hand side of (17) can be written as

$$d g_L(\xi\omega) + \delta_S(-\xi\psi)\Omega_{2n+1} = d g_L(\xi\omega) + ds(-\xi\psi) + l_{-\xi\psi} d\Omega_{2n+1}$$

$$= d(g_L(\xi\omega) + s(-\xi\psi)) + l_{-\xi\psi} A_{2n+2}. \tag{18}$$

Equation (17) thus simplifies as

$$d(\delta_S(\eta)s(\epsilon) - \delta_S(\epsilon)s(\eta) - s(-\xi\psi) - g_L(\xi\omega)) = -(\delta_S(\eta) l_\epsilon - \delta_S(\epsilon) l_\eta - l_{-\xi\psi}) A_{2n+2}. \tag{19}$$

A_{2n+2} is a polynomial in traces of the products of curvatures.

$$l_\epsilon A_{2n+2} = A(\alpha_\epsilon, R), \tag{20}$$

where $A(\alpha_\epsilon, R)$ is the sum of terms obtained by replacing each R in A_{2n+2} in succession by $\alpha_\epsilon = \delta_S(\epsilon)\omega$. It is a $2n+1$ form. For example, when $n = 3$

$$A_8 = c_2 tr R^4 + 2c_{0,1}(tr R^2)^2 \tag{21a}$$

$$A_7(\alpha_\epsilon, R) = 4(c_2\alpha_\epsilon R^3 + c_{0,1} trR^2 tr\alpha_\epsilon R). \tag{21b}$$

Thus

$$\delta_S(\eta) l_\epsilon \mathcal{A}_{2n+2} = A(\delta_S(\eta)\alpha_\epsilon, R) + A(\alpha_\epsilon, D\alpha_\eta, R)$$

$$\text{and} \quad (\delta_S(\eta) l_\epsilon - \delta_S(\epsilon) l_\eta)\mathcal{A}_{2n+2} = i_\xi \mathcal{A}_{2n+2} + l_{-\xi\psi}\mathcal{A}_{2n+2} - dA(\alpha_\epsilon, \alpha_\eta, R), \tag{22}$$

where i_ξ denotes interior contraction with ξ_μ. $A(\alpha_\epsilon, \alpha_\eta, R)$ is the sum of terms obtained by replacing each R in $A(\alpha_\epsilon, R)$ in succession by α_η. We have used

$$[\delta_S(\epsilon), \delta_S(\eta)]\omega = i_\xi R + \delta_S(-\xi\psi)\omega. \tag{23}$$

On a $2n+1$ dimensional disc $i_\xi \mathcal{A}_{2n+2} = 0$. Use of (22) in equation (19) then gives

$$d(\delta_S(\eta)s(\epsilon) - \delta_S(\epsilon)s(\eta) - s(-\xi\psi) - g_L(\xi\omega) - A(\alpha_\epsilon, \alpha_\eta, R)) = 0. \tag{24}$$

Upon integration over a $2n+1$ dimensional disc with spacetime as boundary,

$$\delta_S(\eta)S_0(\epsilon) - \delta_S(\epsilon)S_0(\eta) - S_0(-\xi\psi) - G_L(\xi\omega) = \int_{S^{2n}} A(\alpha_\epsilon, \alpha_\eta, R). \tag{25}$$

Using this for $S_0(\epsilon)$, equation (7c) simplifies as

$$\delta_S(\eta)S_{inv}(\epsilon) - \delta_S(\epsilon)S_{inv}(\eta) - S_{inv}(-\xi\psi) = -\int_{S^{2n}} A(\alpha_\epsilon, \alpha_\eta, R). \tag{26}$$

This is the equation to be satisfied by the invariant part of the anomaly. The expression for $A(\alpha_\epsilon, \alpha_{\epsilon'}, R)$ is very much dimension dependent and there seems to be no general way of solving this equation. In two dimensions, $\int A(\alpha_\epsilon, \alpha'_\epsilon, R)$ is zero. This can be seen as follows. The index density in four dimensions is [6]

$$\mathcal{A}_4 = \frac{1}{192\pi^2} trR^2, \tag{27}$$

This gives immediately

$$A_2(\alpha_\epsilon, \alpha_{\epsilon'}, R) = \frac{1}{96\pi^2} Str(\delta_S(\epsilon)\omega \; \delta_S(\epsilon')\omega). \tag{28}$$

Using $\delta_S(\epsilon)\omega^{ab}_\mu = 2\bar{\epsilon}\gamma_\mu\psi^{ab}$, we find

$$A_2 = \frac{1}{12\pi^2}(\epsilon^{\mu\nu}\bar{\epsilon}\gamma_\mu\psi^{01}\bar{\epsilon'}\gamma_\nu\psi^{01})d^2x$$

$$= \frac{det(e)}{12\pi^2}(\bar{\epsilon}\gamma_0\psi^{01}\bar{\epsilon'}\gamma_1\psi^{01} - \bar{\epsilon}\gamma_1\psi^{01}\bar{\epsilon'}\gamma_0\psi^{01})d^2x.$$

Using $\gamma_1 = \gamma_0\gamma_5$ and the Weyl nature of the spinors, we see that A_2 is zero. Thus in two dimensions, equation (26) is homogeneous with respect to $S_{inv}(\epsilon)$. A solution is $S_{inv}(\epsilon) = 0$ and the entire anomaly is given by $S_0(\epsilon)$ which can be written out as

$$S_0(\epsilon) = \frac{1}{96\pi^2}\int Str\delta_S(\epsilon)\omega \;\; \omega = \frac{i}{96\pi^2}\int d^2x(\epsilon^{\mu\nu}\omega^{01}_\mu\bar{\epsilon}\gamma_\nu\psi^{01}). \tag{29}$$

This result agrees with explicit computation of the anomaly [3].

Gauge theory in two and six dimensions

We now consider (rigid) supersymmetric gauge theories. In this case, the supersymmetry anomaly is given by [1]

$$S(\epsilon) = S_0(\epsilon) + S_{inv}(\epsilon), \tag{30a}$$

where, upto a overall normalization factor,

$$S_0(\epsilon) = n(n+1)\int_{S^{2n}}\int_0^1 dttStr(\delta_S(\epsilon)A, A, F_t^{n-1}), \tag{30b}$$

and $S_{inv}(\epsilon)$ satisfies the equation

$$\delta_S(\eta)S_{inv}(\epsilon) - \delta_S(\epsilon)S_{inv}(\eta) = -n(n+1)\int_{S^{2n}} Str(\delta_S(\epsilon)A, \delta_S(\eta)A, F^{n-1}). \tag{30c}$$

In the reference 1, equation (30c) was solved for four dimensions to give*

$$S(\epsilon) = \int_{S^4} Str\delta_S(\epsilon)A(AdA + dAA + \frac{3}{2}A^3 + \frac{1}{4}\bar{\psi}\gamma^{(3)}\psi). \tag{31}$$

We shall now do the same for two and six dimensions.

The two dimensional gauge multiplet has the fields $(A^a_\mu, \lambda^a, P^a)$ where P^a is a pseudoscalar field [8], λ^a is a Majorana spinor. The transformation rules are

$$\delta_S(\epsilon)A_\mu = i\bar{\epsilon}\gamma_\mu\lambda$$

$$\delta_S(\epsilon)\lambda = \frac{1}{2}\gamma_{\mu\nu}F_{\mu\nu}\epsilon - i\gamma_5\not{\partial}P\epsilon$$

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* In reference 1, the coefficient of the fermionic trilinear terms of the anomaly in four dimensions was given as one half. This is corrected here to one quater.

$$\delta_S(\epsilon)P = \bar{\epsilon}\gamma_5\lambda. \tag{32}$$

This is in fact a $N = 2$ model and so the equation for $S_{inv}(\epsilon)$ is changed. The modification comes from the modified commutation rule

$$[\delta_S(\eta), \delta_S(\epsilon)] = \delta_G(\xi) + \delta_{\tilde{\Lambda}}, \tag{33}$$

$$\text{where} \quad \tilde{\Lambda} = 2i\bar{\eta}\gamma_\mu\epsilon A_\mu - 2\bar{\epsilon}\gamma_5\eta P.$$

The term in $\tilde{\Lambda}$ depending on P changes the equation for $S_{inv}(\epsilon)$ to

$$\delta_S(\eta)S_{inv}(\epsilon) - \delta_S(\epsilon)S_{inv}(\eta) = -4\bar{\epsilon}\gamma_5\eta \int TrPF. \tag{34}$$

A similar equation holds for other extended supersymmetric theories.

The solution to this equation is found to be

$$S_{inv}(\epsilon) = \int d^2x tr(\bar{\epsilon}\lambda P) \tag{35}$$

giving

$$S(\epsilon) = \int Str[\delta_S(\epsilon)AA + \bar{\epsilon}\lambda P]. \tag{36}$$

The $N = 1$ case is obtained by truncation of this theory: λ is now taken to be a Majorana-Weyl spinnor and P is set to zero.

We now turn to six dimensions. In this case, since we do not have auxiliary fields, $S_{inv}(\epsilon)$ satisfies equation (30c) for $n = 3$ only up to terms which vanish when the equations of motion are satisfied. The complete solution is given by (up to overall normalization).

$$S(\epsilon) = S_0(\epsilon) - \frac{2}{3}\int_{S^6} Str(\delta_S(\epsilon)A\overline{\psi}\gamma^{(3)}\psi F)$$

$$= \int_{S^6} Str\delta_S(\epsilon)A[3A(dA)^2 + \frac{12}{5}A(A^2dA + dAA^2) + 2A^5 - \frac{2}{3}\overline{\psi}\gamma^{(3)}\psi F]. \tag{37}$$

A few words about the algebra of verifying equation (30c) for this case are in order. By taking variations of the invariant part of expression (37) and using standard Fierz identities, we get

$$\delta_S(\eta)S_{inv}(\epsilon) - \delta_S(\epsilon)S_{inv}(\eta)$$

$$= -12\int_{S^6} Str(\delta_S(\epsilon)A, \delta_S(\eta)A, F^2) - \frac{2}{3}\int_{S^{2n}} Str(\delta_S(\epsilon)A\delta_S(\eta)AD\overline{\psi}\gamma^{(3)}\psi). \tag{38}$$

Thus we have to show that the second term on the right hand side vanishes when the fermionic equations of motion $\not{D}\psi = 0$ are imposed. Fierz transformations again reduce this term as follows:

$$Str(\delta_S(\epsilon)A\delta_S(\eta)AD(\overline{\psi}\gamma^{(3)}\psi)) \approx 3!Str(\delta_S(\epsilon)A_\mu\delta_S(\eta)A_\nu(D_\alpha\overline{\psi}\gamma_\beta\psi - \overline{\psi}\gamma_\beta D_\alpha\psi)\delta^{\mu\nu}_{\alpha\beta}), \quad (39)$$

where $\approx$ means equality up to terms which vanish when $\not{D}\psi = 0$. The following Fierz identities in six dimensions (which can be proved by judicious use of the standard one) then show that the above expression is a total derivative, upto terms which are zero when $\not{D}\psi = 0$.

$$(\overline{\epsilon}\gamma_\alpha\psi\overline{\eta}\gamma_\beta\psi - \overline{\epsilon}\gamma_\beta\psi\overline{\eta}\gamma_\alpha\psi)\overline{\psi}\gamma_\beta D_\alpha\psi \approx \text{total derivative}$$

$$(\overline{\psi}\gamma_\alpha\epsilon\overline{\psi}\gamma_\beta\eta - \overline{\psi}\gamma_\beta\epsilon\overline{\psi}\gamma_\alpha\eta)\overline{\psi}\gamma_\beta D_\alpha\psi \approx \text{total derivative}$$

$$(\overline{\psi}\gamma_\alpha\epsilon\overline{\eta}\gamma_\beta\psi - \frac{1}{2}\overline{\psi}\gamma_\beta\epsilon\overline{\eta}\gamma_\alpha\psi)\overline{\psi}\gamma_\beta D_\alpha\psi \approx \text{total derivative}. \quad (40)$$

Thus (37) is indeed a solution to the consistency conditions. One can also show that (37) is unique upto terms which can be written as the supervariation of the local terms.

In ten dimesions, one needs the analogues of the identities (40) to obtain a solution for $S_{inv}(\epsilon)$. This is under investigation. We are considering supersymmetry anomalies induced by gauge and gravitational anomalies. Thus theories which are free of gauge and gravitational anomalies will be free of these anomalies too. More specifically, consider the field theory limit of the anomaly free superstring theories[9]. The vacuum persistence amplitude for these theories is given by

$$< 0|0 >= \int d\mu(\Phi)\exp(-\int(\mathcal{L}_0 + \mathcal{L}_1 + \mathcal{L}_2 + \mathcal{L}_3 + \mathcal{L}')) \quad , \quad (41)$$

where $\mathcal{L}_0$ is the supersymmetrized tree level action, $\mathcal{L}'$ gives gauge fixing and ghost terms and $\mathcal{L}_1, \mathcal{L}_2, \mathcal{L}_3$ are the local counterterms constructed by Green and Schwarz [9]. The measure $d\mu(\Phi)$ contains, in particular the gaugino and gravitino terms; these are to be defined by standard eigenmode expansions. Equation (41) is to be interpreted as an effective action endowed with a cutoff for the zero mass sector of the string theory. The interpretation of the cancellation of anomalies would then be that the natural measure of integration emerging from string calculations is not $d\mu(\Phi)$ but $d\mu(\Phi)\exp(-\int(\mathcal{L}_1 + \mathcal{L}_2 + \mathcal{L}_3))$. This measure is invariant under gauge and local Lorentz transformations (for gauge groups $SO(32)$ and

$E_8 \times E_8$). It is then invariant under supersymmetry transformations also, no additional counterterms are necessary.(see also [10]). The argument of course does not constitute a proof of the absence of supersymmetry anomaly since other candidates for $S(\epsilon)$ not linked to gauge and Lorentz transformations have also to be ruled out. These candidates must be gauge and Lorentz invariant and must satisfy a homogeneous consistency condition (eq.(26) or eq.(30c) with the right hand side put to zero). In two and four dimensions, one can show that there are no such solutions [11]. In ten dimensions, although it is very likely that this is the case, there is no conclusive proof yet.

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