Spinors and Supersymmetry

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Abstract. In this paper, we survey the nature of spinors and supersymmetry (SUSY) in various types of spaces. We treat two distinct types of spaces: flat spaces and spaces of constant (non-zero) curvature. The flat spaces we consider are either three or four dimensional of signatures 3 + 1, 4 + 0, 2 + 2 and 3 + 0. In each of these cases, SUSY generators anti-commute to yield the generators of translations in the non-compact flat spaces. The spaces of constant curvature we consider are two-dimensional: the surface of the sphere $S_2$ and the Anti-deSitter space $AdS_2$. $S_2$ is embedded in a 3 + 0 Euclidean space while $AdS_2$ is embedded in 2 + 1 Minkowski space. The SUSY generators in these cases anti-commute to yield the generators of the isometry groups ($SO(3)$ or $SO(2,1)$) of the space involved.

We also report on some recent developments in looking for superspace realizations of these SUSY algebras. We can report good progress in the 3 + 0 Euclidean and in the $AdS_2$ case, somewhat less in the $S_2$ case. In each of the compact cases, we can construct field multiplet models carrying invariance under the full SUSY algebra.

1 Flat Space SUSY Analysis

1.1 3 + 1 Dimensions

The analysis of spinors and supersymmetry (SUSY) in three dimensional Minkowski space is quite standard. (See for example ref. [1].) In a representation in which Dirac matrices are given by

\[ \gamma^\mu = \left[ \begin{array}{cc} 0 & 1 \\ 1 & 0 \end{array} \right], \left[ \begin{array}{cc} 0 & \vec{\sigma} \\ -\vec{\sigma} & 0 \end{array} \right] \]

and a charge conjugation matrix $C$ is defined by

\[ C^{-1} \gamma^\mu C = -\gamma^{\mu T}, \]

we have spinors

\[ \Psi = \left( \begin{array}{c} \psi_\alpha \\ \chi_\dot{\alpha} \end{array} \right), \quad \Psi^\dagger \gamma^0 = \left( \chi_\alpha, \bar{\psi}_\dot{\alpha} \right) \]
and
\[ \Psi_C \equiv C \bar{\Psi}^T = \left( \begin{array}{c} \chi^\alpha \\ \psi_\alpha \end{array} \right) \quad \bar{\Psi}_C = (\psi^\alpha, \chi_\alpha) \] (3b)

forming representations of the Lorentz group. The spinorial generator \( Q \) of the \( N = 1 \) extension of the Poincaré group is Majorana (ie, \( Q = Q_C = \left( \begin{array}{c} Q^\alpha \\ \bar{Q}^{\dot{\alpha}} \end{array} \right) \)) and satisfies the algebra
\[ \{Q_\alpha, Q_\beta\} = 0 \] (4a)
\[ \{Q_\alpha, \bar{Q}_{\dot{\beta}}\} = 2 \sigma^\mu_{\alpha\dot{\beta}} P_\mu. \] (4b)

The two spinorial generators of \( N = 2 \) SUSY extension of the Poincaré group are both Majorana (ie, \( Q_i = Q_C \) for \( i = 1, 2 \)) and satisfy the algebra
\[ \{Q^i_\alpha, \bar{Q}_{\dot{\beta}}\} = 2 \delta^i_j \sigma_{\alpha\dot{\beta}} P_\mu \] (5a)
\[ \{Q^i_\alpha, Q^j_\beta\} = \epsilon_{\alpha\beta} \epsilon^{ij} Z \] (5b)

incorporating a central charge \( Z \) which commutes with all the other generators of this algebra. A representation of this algebra can be found using Fermionic creation and annihilation operators
\[ a_\alpha = \frac{1}{\sqrt{2}} \left( Q^1_\alpha + \epsilon_{\alpha\beta} Q^{2\dot{\beta}}_\beta \right), \quad b_\alpha = \frac{1}{\sqrt{2}} \left( Q^1_\alpha - \epsilon_{\alpha\beta} Q^{2\dot{\beta}}_\beta \right). \] (6)

From (5) it follows that
\[ \{a_\alpha, a^{\dagger}_{\beta}\} = \delta_{\alpha\beta}(2M + Z) \] (7a)
\[ \{b_\alpha, b^{\dagger}_{\beta}\} = \delta_{\alpha\beta}(2M - Z) \] (7b)

in a frame in which \( P_\mu = (M, \vec{0}) \). By (7b) we obtain the “BPS” bound
\[ 2M \geq Z. \] (8)

### 1.2 4 + 0 Dimensions

The situation in four dimensional Euclidean space is quite different. In this case
\[ \gamma^\mu = \left[ \begin{array}{cc} 0 & 1 \\ 1 & 0 \end{array} \right], \quad \left( \begin{array}{cc} 0 & i \vec{\sigma} \\ -i \vec{\sigma} & 0 \end{array} \right) \] (9)

so that
\[ \Psi = \left( \begin{array}{c} \psi_\alpha \\ \chi^{\dot{\alpha}} \end{array} \right) \quad \bar{\Psi} = \Psi^\dagger = \left( \bar{\psi}^{\dot{\alpha}}, -\bar{\chi}_\alpha \right) \] (10a)

and
\[ \Psi_C = C \bar{\Psi}^T = \left( \begin{array}{c} -\bar{\psi}^{\dot{\alpha}} \\ \bar{\chi}_\alpha \end{array} \right) \quad \bar{\Psi}_C = (\psi_\alpha, \chi^{\dot{\alpha}}) \] (10b)
for representations of $SO(4) = SU(2) \times SU(2)$. The spinor $\psi_\alpha$ transforms under one SU(2) subgroup while $\chi^\dot{\alpha}$ transforms under the other SU(2) subgroup. It is also evident from (10) that as $(\bar{\Psi}_C)_C = -\Psi$, we cannot have Majorana spinors in $4+0$ dimensions. The simplest self conjugate SUSY algebra is now [2]

$$\{Q_a, R_b\} = i\sigma^\mu_{ab} P^\mu \quad \{Q_a, R_b\} = 0$$

(11a, b)

$$\{Q_a, \overline{Q}_b\} = \epsilon_{ab} Z \overline{Q} \overline{Q}$$

(11c, d)

or, equivalently, if we define

$$G = \begin{pmatrix} Q_a \\ R^\dot{a} \end{pmatrix} \Rightarrow S_{a1} = Q_a \\ S_{a2} = \overline{Q}_a$$

$$T_{\dot{a}1} = -R^\dot{a} \\ T_{\dot{a}2} = \overline{R}^\dot{a}$$

(12)

we obtain an equivalent algebra which displays an $SU(2)$ structure

$$\{S_{ai}, S_{bj}\} = \epsilon_{ab} \epsilon_{ij} Z \overline{Q} \overline{Q}$$

(13a)

$$\{T_{\dot{a}i}, T_{\dot{b}j}\} = \epsilon_{\dot{a}\dot{b}} \epsilon_{ij} Z \overline{R} \overline{R}$$

(13b)

$$\{S_{ai}, T_{\dot{b}j}\} = i\epsilon_{ij} \sigma^\mu_{ab} P^\mu.$$  

(13c)

This is similar in form to (5) with the roles of $Z$ and $P^\mu$ “reversed”. Thus the simplest SUSY extension of the ISO(4) group in $4+0$ dimensions is an $N = 2$ algebra. This algebra when rewritten in terms of Fermionic creation and annihilation operators becomes, in the frame where $P^\mu = (0, 0, 0, P)$,

$$\{A_a, A^\dagger_b\} = \delta_{ab} \left[ 1 + P \left( Z \overline{Q} \overline{Q} Z \overline{R} \overline{R} \right)^{-1/2} \right]$$

(14a)

$$\{B_a, B^\dagger_b\} = \delta_{ab} \left[ 1 - P \left( Z \overline{Q} \overline{Q} Z \overline{R} \overline{R} \right)^{-1/2} \right].$$

(14b)

We hence see that $P = \sqrt{P^\mu P_\mu}$ has an upper bound in $4+0$ dimensions if the Hilbert space is to be positive definite,

$$P \leq \left( Z \overline{Q} \overline{Q} Z \overline{R} \overline{R} \right)^{1/2}.$$  

(15)

As in $3+1$ dimensions, saturating the bound eliminates one half of the states.

An important distinction between $N = 2$ SUSY in Minkowski space and $N = 2$ SUSY in Euclidean space can now be drawn. In $3+1$ space the central charge provides a lower bound on the magnitude of the momentum, of the mass associated with the state. The lower bound can be zero; there is no inconsistency in considering zero central charge. In $4+0$ space on the other hand, the central charge provides an upper bound on the magnitude of the momentum. Such an upper bound on a positive definite quantity $\sqrt{P^\mu P_\mu}$ cannot be zero; the case of a zero central charge can only lead to all states having zero momentum yielding a trivial theory. We conclude that in $4+0$ space we must include a central charge.
A similar upper bound on momentum arises when one has extended SUSY in 4 + 0 dimensions with algebra [3]

\[
\{Q_{ai}, \overline{Q}_{bj}\} = \epsilon_{ab}Z^Q_{ij} \tag{16a}
\]

\[
\{R_{ai}, \overline{R}_{bj}\} = \epsilon_{ab}Z^R_{ij} \tag{16b}
\]

\[
\{Q_{ai}, \overline{R}_{bj}\} = i\sigma^m_{ab} \epsilon_{ij} P^\mu. \tag{16c}
\]

Just as \( N = 2 \) super Yang-Mills theory in 3 + 1 dimensions can be obtained by dimensional reduction of the \( N = 1 \) gauge theory in 5 + 1 dimensions, so also the supersymmetric gauge model of Zumino in 4 + 0 dimensions can be generated; it has the action

\[
S = \int d^4x E \left[ -\frac{1}{4} F^2_{\mu\nu}(A) + \frac{1}{2} (D_\mu A)^2 - \frac{1}{2} (D_\mu B)^2 - \frac{i}{2} (\psi^\dagger \gamma^\nu \overleftrightarrow{D} \psi) + ig\psi^\dagger (A - B \gamma_5) \psi + \frac{1}{2} g^2 (A \times B)^2 \right]. \tag{17}
\]

One simply drops dependence on one space variable and the time variable in the 5 + 1 dimensional model and has the corresponding components of the vector field identified with the scalar fields \( A \) and \( B \). Explicit calculation [4] shows that the \( \beta \)-function in this model is the same as that in \( N = 2 \) gauge theory in 3 + 1 dimensions despite the peculiar kinetic terms in (16) for the scalars \( A \) and \( B \).

A model with extended SUSY invariance in 4 + 0 dimensions can be obtained by dimensional reduction of \( N = 1 \) gauge theory in 9 + 1 dimensions. It is expected that the \( \beta \)-function in this model vanishes, just as it does for \( N = 4 \) gauge theory in 3 + 1 dimensions, thereby ensuring that conformal invariance is unbroken.

The \( SU(2) \) structure of (12) allows one to define a Harmonic superspace in conjunction with 4 + 0 dimensions [5]. This allows for off-shell realization of this symmetry in these models.

We also note that in 4 + 0 dimensions, one can define a model which is (a) Hermitian (b) gauge invariant under an axial \( U(1) \) gauge transformation (c) anomaly free. Its action is

\[
S = \int d^4x E \left( \frac{1}{4} F^\mu_{\nu}(A) F^\nu_{\mu}(A) + \Psi^\dagger_C (\not{p} + A \gamma_5) \Psi + \Psi^\dagger (\not{p} - A \gamma_5) \Psi_C \right). \tag{18}
\]

No analogue of this model can be defined in 3 + 1 dimensions.

The usual form of the actions considered in 4dE are

\[
L^{(1)} = \frac{1}{4} F^\mu_{\nu}(A) F^\nu_{\mu}(A) + \Psi^\dagger (\not{p} + A \gamma_5) \Psi \tag{19a}
\]

or

\[
L^{(2)} = \frac{1}{4} F^\mu_{\nu}(A) F^\nu_{\mu}(A) + \Psi^\dagger (\not{p} + i A \gamma_5) \Psi. \tag{19b}
\]

The former Lagrangian is non-Hermitian while the latter does not have a compact axial gauge invariance.
1.3 2 + 2 Dimensions

In 2 + 2 dimensions, spinors can be both Majorana and Weyl [6]. Spinors take the form

\[ \Psi = \begin{pmatrix} \phi_a \\ \chi^\alpha \end{pmatrix}, \quad \overline{\Psi} = \left(i \epsilon^{ab} \phi_b, i \epsilon_{ab} \chi^b \right) \]  

(20a, b)

\[ \Psi_C = \begin{pmatrix} \overline{\phi}_a \\ \overline{\chi}^\alpha \end{pmatrix}, \quad \overline{\Psi}_C = \left(\phi^a, \chi_{\dot{\alpha}} \right) \]

and the two simplest SUSY algebras are

\[ \{q_a, r^\dot{b}\} = 2 (\sigma_\mu)_{ab} P^\mu \]  

(no central charge)  

(21a)

and

\[ \{Q, \overline{Q}\} = 2 \gamma^\mu P_\mu + Z + Z_5 \overline{\gamma}_5 \]

(21b)

for Majorana and Dirac spinorial generators \( Q = \begin{pmatrix} q_a \\ r^\dot{a} \end{pmatrix} \) respectively. In [6] it is shown that both of these algebras can be rewritten in terms of Fermionic creation and annihilation operators that generate a Hilbert space with negative norm states; this is taken to indicate that SUSY is incompatible with a 2 + 2 dimensional space.

1.4 3 + 0 Dimensions

In three dimensional Euclidean space, the simplest SUSY algebra is [7,8]

\[ \{Q, \overline{Q}\} = \tilde{\sigma} \cdot \tilde{p} + Z \]

(22)

where \( Q \) is a two component Dirac spinorial generator, \( \tilde{\sigma} \) is a set of Pauli matrices and \( Z \) is a central charge operator. Forming a superspace with coordinates \((x^\mu, \zeta, \theta_i\) and \(\theta^i \)) allows one to make the identifications

\[ Q_i = \frac{\partial}{\partial \theta_i^\dagger} - \frac{i}{2} \left(\tilde{\sigma} \cdot \tilde{\nabla} \theta\right)_i - \frac{i}{2} \left(\theta \frac{\partial}{\partial \zeta}\right)_i \]

(23a)

\[ P_\mu = -i \frac{\partial}{\partial x^\mu}, \quad Z = -i \frac{\partial}{\partial \zeta}. \]

(23b, c)

This makes it possible to formulate supersymmetric models in 3+0 dimensions which are analogous to both the Wess-Zumino and \( N = 1 \) gauge models in 3 + 1 dimensions. A similar analysis can be applied to \( N = 2 \) supersymmetric models in 2 + 1 dimensions. Dimensional reduction can be used to establish a relationship between supersymmetric models in 3 + 1 dimensions and three dimensional supersymmetric models.

We note that just in (8) and (15), in 2 + 1 dimensions the central charge in extended SUSY models provides a lower bound for the momentum, while in 3 + 0 dimensions, it provides an upper bound.

An analysis of supersymmetry in five dimensions [7] reveals that much as in four dimensions, no time dimensions implies an upper bound on momentum; one time dimension implies a lower bound on momentum and two time dimensions implies that for all momentum, negative norm states occur.


2 Constant Curvature Space SUSY Analysis

2.1 $S_2$

The simplest SUSY algebra [9] associated with the two dimensional surface of a sphere embedded in three dimensions is

$$\{Q_i, Q_j\} = 0, \quad \{Q_i, Q^*_j\} = Z\delta_{ij} - 2\vec{\sigma}_{ij} \cdot \vec{P}$$

$$[J^a, Q] = -\frac{1}{2}\sigma^a Q,$$  \hspace{1em} $$[J^a, J^b] = i\epsilon^{abc} J^c, \quad [Z, Q] = -Q$$

(Note that $Z$ is no longer a “central charge” as it does not commute with $Q$.) To examine representations of this superalgebra, we define a state $|I\rangle$ such that

$$J^2|I\rangle = j(j+1)|I\rangle, \quad J_3|I\rangle = m|I\rangle.$$

$$Z|I\rangle = \zeta|I\rangle, \quad Q|I\rangle = 0.$$  \hspace{1em} (25)

Now if $|i\rangle = Q^*_1|I\rangle$ and $|F\rangle = Q^*_1 Q^*_2|I\rangle$, we find that $<1|1\rangle = (\zeta + 2m), \quad <2|2\rangle = \zeta - 2m, \quad <F|F\rangle = (\zeta - 2j)(\zeta + 2j + 2)$, showing that a positive definite Hilbert space occurs if $\zeta \geq 2j$. \hspace{1em} (26)

A model invariant under transformations generated by the SUSY algebra of (22) is

$$S = \int dA \left\{ \frac{1}{2} \Psi^\dagger (\sigma \cdot L + x) \Psi - \Phi^* \left( L^2 + x(1-x) \right) \Phi - \frac{1}{4} F^* F \right\} + \lambda_N \left( 2(1-2x)\Phi^* \Phi - (F^* \Phi + F \Phi^*) - \Psi^\dagger \Psi \right)^N. \hspace{1em} (27)$$

The off mass shell transformations are

$$\delta \Phi = \xi^\dagger \Psi, \quad \delta \Psi = 2(\sigma \cdot L + 1 - u) \Phi \xi - F \xi, \quad \delta F = -2 \xi^\dagger (\sigma \cdot L + x) \Psi \hspace{1em} (28a-c)$$

$$\delta_Z \Phi = [2(1-2x)\Phi - F], \quad \delta_Z \Psi = [1 + 2\sigma \cdot L] \Psi, \quad \delta_Z F = -4 \left[ L^2 + x(1-x) \right] \Phi + 2xF. \hspace{1em} (28d-f)$$

The symmetries of (28d-f) are in fact new symmetries.

A superspace representation of the algebra of (22) is provided by

$$Q = (\sigma \cdot \tau + \zeta) \frac{\partial}{\partial \theta^i} - \left( \frac{\partial}{\partial \zeta} - \sigma \cdot \nabla \right) \theta \hspace{1em} (29a)$$

$$Q^\dagger = \frac{\partial}{\partial \theta^i} (\sigma \cdot \tau + \zeta) + \theta^i \left( \frac{\partial}{\partial \zeta} - \sigma \cdot \nabla \right) \hspace{1em} (29b)$$

$$J^a = -i(r \times \nabla)^a + \frac{1}{2} \left( \theta^a \sigma^a \frac{\partial}{\partial \theta^i} + \frac{\partial}{\partial \theta} \sigma^a \theta \right) \hspace{1em} (29c)$$

$$Z = -\theta^i \frac{\partial}{\partial \theta^i} + \theta \frac{\partial}{\partial \theta}. \hspace{1em} (29d)$$
We note that under a supersymmetry transformation generated by (29)

\[
\delta r^a = \epsilon^\dagger \sigma^a \theta + \theta^\dagger \sigma^a \epsilon \\
\delta \theta = \vec{\sigma} \cdot \vec{\epsilon} + \zeta \epsilon.
\]

Furthermore, we see that

\[
\left[ Q, \vec{r}^2 - \zeta^2 - 2\theta^\dagger \theta \right] = 0 \\
\left[ Q, \theta^\dagger \frac{\partial}{\partial \theta} + \theta \frac{\partial}{\partial \vec{\theta}} + \vec{r} \cdot \vec{\nabla} + \zeta \frac{\partial}{\partial \zeta} \right] = 0.
\]

Currently we are attempting to formulate the model of (25) in terms of superfields using the superspace realizations (29) of the generators and (31a,b).

### 2.2 AdS$_2$

On AdS$_2$ we have the algebra

\[
[J_{ab}, J_{cd}] = \eta_{ac} J_{bd} - \eta_{bc} J_{ad} + \eta_{bd} J_{ac} - \eta_{cd} J_{ab} \\
\{Q, \tilde{Q}\} = 2\Sigma^{ab} J_{ab} \quad (\tilde{Q} = Q \gamma_2) \\
\{J_{ab}, Q\} = -\Sigma_{ab} Q
\]

\((Q\text{ is Majorana, } \eta_{ab} = \text{diag}(+,-,+), \gamma_a \gamma_b = -\eta_{ab} - i\epsilon^{abc} \gamma_c, \Sigma_{ab} = \frac{1}{4} \left[ \gamma_a, \gamma_b \right].)\) This algebra can be realized by

\[
J_{ab} = \frac{\partial}{\partial \theta} \Sigma_{ab} \theta - (x_a \partial_b - x_b \partial_a) \\
Q = \gamma^a \partial_a \theta + \gamma^a x_a \frac{\partial}{\partial \theta} \\
\tilde{Q} = -\tilde{\theta} \gamma^a \partial_a + \frac{\partial}{\partial \theta} \gamma^a x_a
\]

where \(\theta\) is a two component Grassmann Majorana spinor. We also define

\[
D = -\gamma^a \partial_a \theta + \gamma^a x_a \frac{\partial}{\partial \theta} \\
\tilde{D} = \tilde{\theta} \gamma^a \partial_a + \frac{\partial}{\partial \theta} \gamma^a x_a.
\]

We note that

\[
\left[ Q_i, x^a \partial_a + \theta_j \frac{\partial}{\partial \theta_j} \right] = 0 \\
\left[ Q_i, x^a x_a - \tilde{\theta} \right] = 0.
\]

Applying the condition

\[
\Delta \Phi = \omega \Phi
\]
so that if

$$\Phi = \phi + \tilde{\lambda}\theta + F\tilde{\theta}\theta$$

then we have $(x \cdot \partial - \omega)\phi = (x \cdot \partial + 1 - \omega)\lambda = (x \cdot \partial + 2 - \omega)F = 0$.

Some suitable supersymmetric actions are

$$S_1 = \int d^3x d^2\theta \delta \left(x^2 - \tilde{\theta}\theta - a^2\right) \Phi(\tilde{D}D + \rho)\Phi$$

$$S_2 = \int d^3x d^2\theta \delta \left(x^2 - \tilde{\theta}\theta - a^2\right) \left(\tilde{D}\Phi D\Phi + \rho\Phi^2\right)$$

$$S_3 = \int d^3x d^2\theta \delta \left(x^2 - \tilde{\theta}\theta - a^2\right) \left(\Phi\tilde{Q}Q\Phi + \rho\Phi^2\right)$$

$$S_4 = \int d^3x d^2\theta \delta \left(x^2 - \tilde{\theta}\theta - a^2\right) \left(\tilde{Q}\Phi Q\Phi + \rho\Phi^2\right).$$

In component form, for example, (38a) reduces to

$$S_1 = \int d^3x \left\{ \delta \left(x^2 - a^2\right) \left[F(-2x^2F + 2(\rho - 1)\phi) + 1 \right. \right.$$

$$+ \frac{1}{2x^2}\phi \left(L^{ab} L_{ab} + 2\omega(1 + \omega)\right)\phi - \tilde{\lambda} \left(\Sigma^{ab} L_{ab} - \frac{3 - \rho}{2}\right)\lambda \right.$$

$$\left.+ \delta' \left(x^2 - a^2\right) \left[2\phi \left(x^2F + \left(\frac{\rho}{2} - \omega\right)\phi\right)\right]\right\}. \quad (39)$$

Other supersymmetric actions on $AdS_2$ can be devised in component form; for example

$$S = \int d^2x \left\{ \tilde{\Psi} \left(\Sigma^{ab} L_{ab} + \chi\right)\Phi + \Phi \left(\frac{1}{2}L^{ab} L_{ab} + \chi(1 + \chi)\right)\Phi \right.$$

$$- FF] + \lambda_N \left[(1 + 2x)\Phi\Phi + \tilde{\Psi}\Psi + 2\Phi F\right]^N \right\} \quad (40)$$

possesses the invariance

$$\delta\Psi = \left[\left(\Sigma^{ab} L_{ab} - (1 + x)\right)\Phi - F\right] \xi \quad (41a)$$

$$\delta\Phi = \tilde{\xi}\Psi, \quad \delta F = - \tilde{\xi} \left(\Sigma^{ab} L_{ab} + \chi\right)\Psi. \quad (41b,c)$$

The relation between the models of (38) and (40) is not apparent.

The role of $\zeta$ in (29) is not at all clear. However, it is necessary to introduce $\zeta$ in order for $Q$ to be the “square root” of the non-Abelian operator $J^a$.

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