Geometrical formulation for the Siegel superparticle

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Abstract

In the superspace $z^M = (x^\mu, \theta_R, \theta_L)$ the global symmetries for $d = 10$ superparticle model with kinetic terms both for Bose and Fermi variables are shown to form a superalgebra, which includes the Poincaré superalgebra as a subalgebra. The subalgebra is realized in the space of variables of the theory by a nonstandard way. The local version of this model with off-shell closed Lagrangian algebra of gauge symmetries and off-shell global supersymmetry is presented. It is shown that the resulting model is dynamically equivalent to the Siegel superparticle.

1 Introduction

By now a problem of manifestly Poincaré covariant quantization of the Green–Schwarz superstring\(^1\) and of massless relativistic superparticles\(^2-5\) has no fully satisfactory solution. In the Lagrangian framework we are faced with the Siegel $k$-symmetry\(^5\) of the Lagrangians, which is written for spinor variables as $\delta \theta_\alpha = \Pi_\nu (\Gamma^\nu k)_\alpha$. On the constraints surface of the theories, only half of the $k$-parameters contributes to $\delta \theta_\alpha$ by virtue of the condition $\Pi^2 = 0$. Since spinor $k_\alpha$ is chosen in the minimal spinor representation of $SO(1, 9)$ (we discuss $d = 10$ case only), this leads in particular to well-known difficulties in imposing a covariant gauge. Despite the importance of $k$-symmetry, there is no its clear geometrical formulation which would be suitable for quantization (in comparing with gauge or general coordinate transformations).

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In the Hamiltonian framework, there are eight first class constraints corresponding to $k$-transformations, which are combined with eight second class constraints in the following covariant equations $\bar{p}_\mu + i\bar{\Gamma}^\mu p_\mu = 0$. To apply standard methods of the Hamiltonian quantization, it is necessary to divide the constraints in a fully covariant manner. This task was solved by using covariant projectors for the cases of Green–Schwarz superstring and of massive superparticle with linearly dependent sets of first and second class constraints in the results. Thus, a real problem in these cases is quantization of a theory with infinite rank of reducibility.

One of possibilities to avoid the problems is to build modified and more convenient for the quantization formulations (in particular, a necessity of some modifications follows from absence of terms corresponding to the covariant propagators for Fermi variables in these theories). Siegel superparticle (SSP) is the least radical modification of such a kind, where in comparing with the Brink–Schwarz superparticle only first class constraints are retained. Let us enumerate the relevant facts for $AB$-formulation (see Ref. 7 and references therein). The defining system of constraints for SSP has the form

$$p_e \approx 0, \quad \pi^2 \approx 0, \quad \bar{p}_\theta \approx 0, \quad \pi_\mu (\bar{\theta}_R \Gamma^\mu)_\alpha \approx 0;$$

$$p_{L\alpha} \approx 0, \quad \bar{p}_{R\alpha} + i\pi_\mu (\theta_R \Gamma^\mu)_\alpha - i\theta_L \approx 0;$$

where $(\pi_\mu, \bar{\theta}_R, \bar{\theta}_L, p_e, \bar{p}_\theta)$ are canonically conjugate momenta for the configuration space variables $(x^\mu, \theta_{R\alpha}, \theta_{L\beta}, e, \psi_{L\alpha})$ and $\Gamma^{11}\theta_{R,L} = \pm\theta_{R,L}$. We use the Majorana representation for $\Gamma^\mu$-matrices in $d = 10$ and the notations from Ref. 15. The first- and second-class constraints (Eqs. (1) and (2) respectively) are separated in manifestly covariant manner and there are only 8 linearly-independent among 16 constraints $\pi_\mu (\bar{\theta}_R \Gamma^\mu)_\alpha \approx 0$, as a consequence of $\pi^2 \approx 0$. Thus, there are 16 dynamical variables in the Fermi-sector of the model. Lagrangian action, reintroducing (1) and (2), can be written in the form

$$S = \int d\tau \left( \frac{1}{2e} \Pi^\mu \dot{\Pi}^\mu + \bar{\theta}_R \dot{\theta}_L \right),$$

and possesses a global supersymmetry in the standard realization $\delta \theta_R = \epsilon_R, \delta x^\alpha = \bar{\theta}_R \Gamma^\alpha \theta_R$ and local Siegel symmetry, the latter being associated with the fermionic first-class constraints from Eq. (1)

$$\delta_k \theta_R = \frac{1}{e} \Pi_\mu \Gamma^\mu k_L, \quad \delta_k \theta_L = \frac{2}{e^2} k_L \Pi^2;$$

$$\delta_k \dot{x}^\mu = \bar{\theta}_L \Gamma^\mu k_L + \bar{\theta}_R \Gamma^\mu (\delta_k \theta_R),$$

$$\delta_k e^\alpha = -\frac{4i}{e} \bar{\theta}_L \Pi_\mu \Gamma^\mu \psi_L, \quad \delta_k \psi_L = \dot{k}_L.$$
On this ground, the question about studying an algebraic structure of global transformations (the $\kappa$-algebra below) corresponding to Siegel symmetry arises quite naturally. As will be shown, the latter is the local and nonlinear version of the $\kappa$-algebra.

The paper is organized as follows. In Sec. 2 we analyze global symmetries of the simplest $d = 10$ action with kinetic terms both for Bose and Fermi variables. We show that in the space of variables of the theory it is realized a superalgebra which includes $\kappa$-transformations and the Poincaré superalgebra as subalgebras. The last one is realized in a non-standard way. Note that, in accordance with the Haag – Lopuszanski – Sohnius theorems a superalgebra which is more wide than the super Poincaré one, can be realized in nontrivial quantum theory by a special way only. Namely, the generators of transformations added to the super Poincaré ones have to vanish on physical states. This is exactly the case for our model. In Sec. 3 we derive the local version for the $\kappa$-transformations of the initial action which can be done without lose of the global off-shell supersymmetry. The resulting model is dynamically equivalent to SSP, but the structure of local symmetries turns out to be more simple in comparing with (4) (in particular, the Lagrangian algebra is off-shell closed). In Conclusion we discuss possibilities of additional modifications with the aim to get an equivalent to the Brink – Schwarz superparticle formulation.

2 $\mathbb{R}^{10,32}$ superspace

Consider $d = 10$ superspace with the coordinates $z^M = (x^\mu, \theta^R, \theta^L)$. Here the odd sector is parametrized by a pair of Majorana – Weyl spinors with opposite chirality: $\Gamma^{11}\theta_{R,L} = \pm\theta_{R,L}$. Let us analyze global symmetries which are present in the simplest Poincaré and reparametrization invariant action in this superspace

$$S = \int d\tau \left( \frac{1}{2e} \dot{x}^\mu \dot{x}^\mu + \bar{\theta} R \theta_L \right).$$

(5)

Except the trivial translations $\delta_\epsilon \theta_L = i(\gamma_L \bar{Q}_R)\theta_L \equiv \epsilon_L$ (where $\epsilon_L$ and $\bar{Q}_R$ are the global parameter and the generator respectively), the model (5) is invariant under the following transformations in the space of functions $z^M(\tau)$: $\delta_\kappa z^M = i(\bar{\kappa}_L S_R)z^M$ where:

$$\delta_\kappa \theta_L = 0, \quad \delta_\kappa \theta_R = \frac{1}{e} \dot{\theta}_\mu \Gamma^\mu \kappa_L,$$

$$\delta_\kappa x^\mu = \bar{\theta} \Gamma^\mu \kappa_L.$$  

(6)

A commutator of two $\kappa$-transformations yields a new one, which can be written in the form

$$[\delta_{\kappa_1}, \delta_{\kappa_2}] z^M = \delta_{b_2} z^M = -i(b_\mu \bar{B}^\mu)z^M, \quad b^\mu = -\bar{\kappa}_{\mu_1} \Gamma^\mu \kappa_2.$$  

(7)
The explicit realization of the superalgebra (10) – (12) for the model (5) is

\[ \delta_\theta R = \frac{1}{e} b_\mu \Gamma^\mu \bar{\theta}_L, \quad \delta_\theta L = \delta_\phi x^\mu = 0. \]  

The mixed commutators of the \( \kappa \)- and \( \epsilon \)-transformations give a Poincaré translation \( \delta_\kappa x^\mu \) with the parameter \( a^\mu = 2 \xi_\mu \epsilon L \). It is straightforward to check that the commutators for \( \delta_\kappa, \delta_\epsilon, \delta_\phi, \delta_\delta \) form a closed algebra with the Jacobi identities fulfilled.

Dimension \( d = 10 \) is unique in the sense that to get the Eqs. like (7) it is necessary to use the Fiertz identities for Majorana – Weyl spinors (which are defined in \( d = 2 \text{Mod} \ 8 \)). Thus, introducing notations \( \bar{P}_\mu, M_{\mu\nu} \) for the Poincaré generators and \( Q_R, S_R, \bar{B}_\mu \) for generators of the above written transformations, we can conclude that the symmetry transformations of the action (5) realize some Lie superalgebra with the following commutation relations (the standard Poincaré subalgebra is omitted)

\[ \{ S_{R_\alpha}, S_{R_\beta} \} = -(\Gamma^\mu)^{\alpha\beta}_{\gamma\delta} \bar{B}_\mu, \quad \{ \bar{Q}_{R_\alpha}, \bar{Q}_{R_\beta} \} = 0, \]
\[ \{ \bar{Q}_{R_\alpha}, S_{R_\beta} \} = (\Gamma^\mu)^{\alpha\beta}_{\gamma\delta} \bar{P}_\mu, \quad \{ M_{\mu\nu}, \bar{B}_\rho \} = i(\eta_{\mu\nu} \bar{B}_\rho - \eta_{\nu\rho} \bar{B}_\mu), \]
\[ \{ M_{\mu\nu}, S_{R_\alpha} \} = -\frac{i}{4} (\Gamma_{\mu\nu} S_{R_\alpha}), \quad \{ M_{\mu\nu}, \bar{Q}_{R_\gamma} \} = -\frac{i}{4} (\Gamma_{\mu\nu} \bar{Q}_{R_\gamma}). \]

For further consideration it is useful to redefine the basis in the algebra by the rules: \( Q_R \rightarrow Q_R = Q_R - S_R, \bar{P}_\mu \rightarrow P_\mu = \bar{P}_\mu + \bar{B}_\mu/2, \bar{B}_\mu \rightarrow B_\mu = \bar{B}_\mu/2 \). In this basis every element of the algebra has the form \( \frac{i}{2} \omega^{\mu\nu} M_{\mu\nu} - a^\mu P_\mu + c_L Q_R + c_R S_R - b^\mu B_\mu \), with the following commutation relations for the generators

\[ \{ Q_{R_\alpha}, Q_{R_\beta} \} = -2(\Gamma^\mu)^{\alpha\beta}_{\gamma\delta} P_\mu, \quad [Q_\alpha, P_\mu] = [Q_\alpha, B_\mu] = 0, \]
\[ [M_{\mu\nu}, Q_{R_\alpha}] = -\frac{i}{4} (\Gamma_{\mu\nu} Q_{R_\alpha}); \]
\[ [S_{R_\alpha}, S_{R_\beta}] = -2(\Gamma^\mu)^{\alpha\beta}_{\gamma\delta} B_\mu, \quad [S_{R_\alpha}, P_\mu] = [S_{R_\alpha}, B_\mu] = 0, \]
\[ [M_{\mu\nu}, S_{R_\alpha}] = -\frac{i}{4} (\Gamma_{\mu\nu} S_{R_\alpha}); \]
\[ \{ Q_{R_\alpha}, S_{R_\beta} \} = (\Gamma^\mu)^{\alpha\beta}_{\gamma\delta} (P_\mu + B_\mu). \]

The explicit realization of the superalgebra (10) – (12) for the model (5) is

\[ \left\{ \begin{array}{l}
\delta_\kappa \theta_L = \epsilon_L, \\
\delta_\kappa \theta_R = \frac{1}{e} b_\mu \Gamma^\mu \bar{\theta}_L, \\
\delta_\kappa x^\mu = \epsilon^\mu \\
\delta_\epsilon \theta_L = 0, \\
\delta_\epsilon \theta_R = \frac{1}{e} b_\mu \Gamma^\mu \theta_L, \\
\delta_\epsilon x^\mu = \epsilon^\mu \\
\delta_\phi \theta_L = 0, \\
\delta_\phi \theta_R = \frac{1}{e} b_\mu \Gamma^\mu \bar{\theta}_L, \\
\delta_\phi x^\mu = \epsilon^\mu \end{array} \right. \]  

\[ \left\{ \begin{array}{l}
\delta_\delta \theta_L = 0, \\
\delta_\delta \theta_R = \frac{1}{e} b_\mu \Gamma^\mu \kappa_L, \\
\delta_\delta x^\mu = \epsilon^\mu \\
\delta_\phi \theta_R = \frac{1}{e} b_\mu \Gamma^\mu \bar{\theta}_L, \\
\delta_\phi x^\mu = \epsilon^\mu = 0. \end{array} \right. \]
Let us give some comments concerning the structure of this superalgebra.

a) From an algebraic point of view, in the basis chosen the superalgebra (10)–(12) contains two subalgebras \((M, P, Q)\) and \((M, B, S)\), both satisfying the commutation relations of the Poincaré superalgebra\(^{16}\), and having nontrivial overlapping in the odd sector (12).

b) In Eqs. (13) – (17) namely the \((M, P, Q)\) subalgebra is realized as the Poincaré superalgebra, because just the \(P^\mu\)-generator corresponds to translations for \(x^\mu\)-variables. Note that only on-shell (\(\dot{\theta}_L = 0\)) the Poincaré translations (14) are realized in a standard way.

c) The transformations (17), corresponding to the \(B^\mu\)-generators, vanish on-shell and consequently \(\{S_{R\alpha}, S_{R\beta}\}|_{\text{on-shell}} = 0\). This means that \(S_{R\alpha}|a >= 0\) for any physical state \(|a >\) and, therefore, the \((S, B)\)-subalgebra is a trivial symmetry in quantum theory. However, such a construction turns out to be useful for describing of superparticle models because the transformations (16) for \(\theta_R\) in fact are the linearized Siegel transformations if we consider \(\kappa_L\) as a local parameter.

d) Since the action (5) is symmetric under the change \(\theta_{R,L} \rightarrow i\theta_{L,R}\), it follows that there exists conjugate to (10)–(12) superalgebra, which can be got from Eqs. (10)–(12) by the change \(L \leftrightarrow R\) (an explicit realization of the conjugate algebra in the action (5) is achieved by the change \(\kappa \rightarrow i\xi\) for all odd parameters and variables \(\xi\)). In the next section, the transformations (16) will have been localized in (5) in such a way that the global supersymmetry with a parameter \(\epsilon_R\) from the conjugate superalgebra will be present in the resulting action.

e) The superalgebra (10) – (12) is realized by Eqs. (13) – (17) on the space of functions \(z^M(\tau)\), because it is convenient for our goals. Let us note that it may be realized on the superspace \((x^\mu, \theta_L, \theta_R)\) as well, if one omits the multiplier \(1/e\) and \(\tau\)-derivatives in Eqs. (13) – (17).

3 The superparticle in \(\mathbb{R}^{10,32}\)-superspace

From an analysis of constraints in the Hamiltonian formalism it follows that there are 16 + 16 dynamical (phase) degrees of freedom in the Fermi sector of the model (5), instead of 16 ones in the Siegel’s model. Therefore, let us consider a local version of the transformations (16) in the action (5). The Noether procedure\(^{17}\), for example, can be used for these goals. In order to localize the \(\kappa\)-transformations it is sufficient to covariantize the derivatives \(\dot{x}^\mu\) and to modify the transformation law for \(\delta_{\epsilon, \theta_R}\) by nonlinear, in the coordinates, terms. The resulting locally-invariant action has the form

\[
S = \int d\tau \left( \frac{1}{2\epsilon} D\theta^\mu D\theta_\mu + i\theta_R^{\mu} \theta_L \right),
\] (18)
where $Dx_{\mu} \equiv x^\mu + i\bar{\psi}_L \Gamma^\mu \theta_L$ and $\psi_L$ is a Majorana–Weyl spinor playing the role of a gauge field. The local version for Eq. (16) looks like

$$\delta_{\epsilon} \theta_R = \frac{1}{e} Dx_{\mu} \Gamma^\mu \epsilon_R, \quad \delta_{\epsilon} \psi_L = \bar{\epsilon}_L,$$

and forms off-shell closed algebra together with the local $\delta_\epsilon$-transformations of Eq. (17).

In comparing with Eq. (5), the new terms in Eq. (18) don’t allow to achieve an invariance under the global supersymmetry transformations (13) with the parameter $\epsilon_L$ or to obtain some generalization of those formulae. However, analogous transformation for (13) in the conjugate superalgebra (with a parameter $\epsilon_R$) is a symmetry of the modified action (18). A mere realization of the symmetry has the form

$$\delta_{\epsilon} \theta_R = \epsilon_R, \quad \delta_{\epsilon} \theta_L = \frac{1}{e} Dx_{\mu} \Gamma^\mu \epsilon_R,$$

and

$$\delta_{\epsilon} x^\mu = \bar{\epsilon}_R \Gamma^\mu \theta_R, \quad \delta_{\epsilon} c = -2i(\bar{\psi}_L \epsilon_R).$$

So far as

$$[\delta_{\epsilon_1}, \delta_{\epsilon_2}] \theta_L = \frac{1}{e} \epsilon_1^R \Gamma^\mu \epsilon_2 R \Theta_R - \frac{1}{e} Dx_{\nu} \Gamma^\nu \psi_L \equiv i \epsilon_L \epsilon_R,$$

the algebra of the commutators is closed on-shell only. Following the standard ideology, it is not difficult to choose an auxiliary variables, which provide off-shell closure of thus realized superalgebra Poincaré. Let us introduce $d = 10$ vector $h^\mu(\tau)$ and modify the transformation law for $\theta_L$ in the following way: $\delta_{\epsilon} \theta_L = \frac{1}{e}(Dx_{\mu} + h_\mu) \Gamma^\mu \epsilon_R$. Requiring the off-shell closure for commutators, one can find the transformation law for $h_\mu$. It is straightforward to check that the transformations

$$\delta_{\epsilon} \theta_R = \epsilon_R, \quad \delta_{\epsilon} \theta_L = \frac{1}{e}(Dx_{\mu} + h_\mu) \Gamma^\mu \epsilon_R,$$

$$\delta_{\epsilon} x^\mu = \bar{\epsilon}_R \Gamma^\mu \theta_R, \quad \delta_{\epsilon} c = -2i(\bar{\psi}_L \epsilon_R),$$

form an off-shell realization of the superalgebra Poincaré together with the Poincaré translations $\delta_{\epsilon} x^\mu = a^\mu$, $\delta_{\epsilon} \theta_{L,R} = 0$. The variation of Eq. (18) under the (22) can be written in the form

$$\delta_{\epsilon} S = \int d\tau \left( -\frac{1}{e} \bar{\psi}_L \Gamma^\mu (\bar{\theta}_R - \frac{1}{e} Dx_{\nu} \Gamma^\nu \psi_L) h_\mu \right) \equiv \int d\tau \delta_{\epsilon} \left( \frac{i}{2e} h_\mu h^\mu \right),$$

(23)
whence it follows a final expression for globally supersymmetric under Eq. (22) and locally invariant under Eq. (19) action
\[
S = \int d\tau \left( \frac{1}{2e} D_{\mu} D^{\mu} \theta R - \frac{1}{2e} \dot{h}_\mu h^\mu \right).
\] (24)

Note that after the redefinition \( h^\mu \rightarrow e h^\mu \), the variables \( h^\mu \) are transformed as scalars under the reparametrizations. Let us check that the obtained model is equivalent to SSP. The constraints system, corresponding to the model (24) in the Hamiltonian formalism is

\[
\begin{align*}
\pi^2 &\approx 0; \\
\pi_\mu &\approx 0; \\
\theta_R &\approx 0; \\
\psi_L &\approx 0; \\
\theta_R \Gamma^\mu &\pi_\mu &\approx 0; \\
p_L &\approx 0; \\
p_R - i \theta_L &\approx 0;
\end{align*}
\] (28) (29)

where \((\pi_\mu, \pi_R, \psi_L, \psi_R, \theta_R, \theta_L, \pi^2, e)\) are momenta for the variables \((x^\mu, \theta_R, \theta_L, \psi_L, \psi_R, e)\), respectively. The constraints (25) are standard and together with Eq. (27) they mean that (in the light-cone gauge) there are the following dynamical degrees of freedom in the Bose sector of the model: \(x^- , x^i, i = 1, \ldots, 8\). The first-class constraints (26) allow us to impose the gauge \(\psi_L \approx 0\), thus \(\psi_L\) is a pure gauge field. In the system of the first-class constraints (28), there are only 8 linearly independent (as a consequence of \(\pi^2 \approx 0\)). Taking into account the second-class constraints (29), we can conclude that there are 16 physical fermionic degrees of freedom in the theory.

Thus, in Eqs. (19), (22), and (24) we obtained the model dynamically equivalent to SSP, but possessing more simple algebra of gauge symmetries in comparing with (4). To achieve these, the crucial observation is a nonstandard realization of the Poincaré superalgebra in space of variables of the theory.

One can note that there is another possibility to achieve the off-shell global supersymmetry in the action (18). The change \(D_{\mu} \rightarrow \Pi_{\mu} \equiv x^\mu - i \theta_R \Gamma^\mu \theta_R + i \psi_L \Gamma^\mu \theta_L\) leads to the SSP action (3), which is invariant under supersymmetry transformations in the standard realization \(\theta_R = \epsilon_R, \delta x^\mu = \epsilon_R \Gamma^\mu \theta_R\), and under the local transformations (4). As we have shown above, this additional modification is not necessary.

4 Conclusion

In comparing with the Brink – Schwarz superparticle, there are absent eight second class constraints in the Siegel model. As a consequence, it leads to undesirable negative norm states after quantization\(^\text{10}\). The action (3), therefore, can be considered only as a model one to study a task of a manifestly covariant quantization of a theory with infinite rank of reducibility. For this reason, in a series of papers an intriguing possibility
of modification of SSP, so as to get an equivalent to the Brink–Schwarz superparticle model, was considered (ABC and ABCD-models\textsuperscript{10–13}). An analogous modification can be done as well, if one starts from Eq. (24) instead of Eq. (3). In particular, the formulation equivalent to the ABC-model arises after adding the term $S_1 = \int d\tau \bar{\theta} R \chi \theta_L$ to the action (24) (where $\chi_{\alpha\beta} = -\chi_{\beta\alpha}$ are the Lagrange multipliers). The transformation laws for $\psi_L$ and $\chi$ can be chosen in such a way that the full action maintains some global supersymmetry. However, the additional first-class constraints arising from the $S_1$-term (the same as in the ABC-model) are quadratic in fermions and it leads to well-known difficulties both for Dirac quantization and for BFV one\textsuperscript{8,9}.

Let us note, therefore, one more possibility. It is straightforward to check that the addition of the term $S_2 = \int d\tau \bar{\xi}_L (\theta_R - x_\mu \Gamma^\mu \theta_L)$ to Eq. (24) leads to the theory which is equivalent to the Brink–Schwarz superparticle (technically, the $S_2$-term supply an addition to the constraints system (25) – (29) of 8 gauge conditions for Eq. (28) and of an additional 8 second-class constraints which are combined in one Lorentz-covariant equation $\theta_R - x_\mu \Gamma^\mu \theta_L \approx 0$). Since there are second-class constraints only, the problems with the Lorentz-covariant gauge fixing is absent in this model. However, covariance under the Poincaré translations takes place on the equations of motion for dynamical variables only. Also, the supertranslations are realized only on the $SO(8)$ component of $\theta_L$ variables.

We couldn’t find additional modifications leading to Poincaré invariant action.

Thus, in the present paper it has been shown that in constructing of $d = 10$ superparticle models with terms corresponding to the covariant propagator for Fermi variables, it is useful to consider the superalgebras which are more wide than the superalgebra Poincaré. The last one in this case is a subalgebra, and is realized in the space of variables of the theory by a nonstandard way. In the case of the Siegel’s model it allows us to obtain a formulation with more simple (in particular, off-shell closed) algebra of the Lagrangian gauge symmetries.

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