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Geometrical aspects of superbrane dynamics

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Abstract

The geometrical (superembedding) approach is used as a tool for deriving from the worldvolume dynamics of superbranes field theoretical models exhibiting partial supersymmetry breaking. In this way we obtain nonlinear actions for Goldstone superfields associated with physical degrees of freedom of the superbranes.

1 Introduction

In this contribution we discuss recent results in the construction of superfield actions for effective field theories on superbrane worldvolumes. These effective supersymmetric field theories describe (in a static or physical gauge) superbrane fluctuations in a supergravity background and the worldvolume dynamics of pure brane fields, such as vector gauge fields of the Dirichlet branes.

The present interest in the field theories associated with branes is at least threefold:

- these are models which exhibit (partial) breaking of supersymmetry, i.e. a property which is important for phenomenological applications of supersymmetry;
- in the case of Anti–de–Sitter backgrounds these effective theories are associated with superconformal field theories on the AdS boundary, and thus are used to test the AdS/CFT correspondence conjecture;
- the knowledge of an explicit form of brane effective actions should be useful for the development of various brane world sceanrii, which are under intensive study these days.

In this paper we will concentrate on the consideration of the first of the items above, namely, on the relationship of superbrane worldvolume actions with models of partial breaking of global supersymmetries. The mechanism of partial supersymmetry breaking caused by superbrane solutions has been under study over a long period of time starting from References [1, 2].

Recently, superfield actions for Goldstone supermultiplets which cause 1/2 supersymmetry breaking in $D=4$ and $D=3$ theories have been constructed in [3, 4, 5].

For instance, it has been shown [3] that spontaneous breaking of global $N=2$ supersymmetry down to $N=1$ in $D=4$ caused by an $N=1$ vector supermultiplet is described by the action for a Goldstone–vector supermultiplet which produces the Dirac–Born–Infeld action in the bosonic sector of the theory [6]. This implies that the supersymmetric model thus obtained should have something to do with a space–filling Dirichlet 3–brane. However, the relation between the fermionic sector of this model and the fermionic fluctuations of the D3–brane (described by a Green–Schwarz–like action [7]) has not been established, since the methods used for the construction of the Goldstone superfield actions differ from the methods used for the construction of superbrane actions. The former are based on methods of non–linear realizations of spontaneously broken symmetries.

The same situation is in a simpler case of breaking $N=2$ supersymmetry down to $N=1$ in $D=3$ by a Goldstone scalar supermultiplet [5] which is believed to be associated with supermembrane fluctuations in an $N=1$, $D=4$ superspace.

In what follows we propose a generic procedure of how one can arrive at the Goldstone superfield actions starting directly from corresponding covariant actions for superbranes. To reach this goal we shall use the superembedding approach which is a generic geometrical method for describing super–p–branes, Dp-branes and M–branes (see [8] for a review and references).

To be concrete we shall describe the basic properties of the superembedding approach and the derivation of Goldstone superfield actions exhibiting partial supersymmetry breaking with the example of a supermembrane propagating in an $N=1$, $D=4$ target superspace, though we should note that the same reasoning is also applicable to more physically interesting, but more complicated, cases of other branes, such as a D3–brane.

2 The $N=1$, $D=4$ supermembrane

Our strategy will be

i) to define the conditions of superembedding describing supermembrane dynamics;

ii) to construct a worldvolume superfield action for the $N=1$, $D=4$ supermembrane;

iii) to gauge fix all local worldvolume symmetries of the action by imposing a physical gauge and to solve for the superembedding condition in terms of an independent (Goldstone) superfield;

iv) to reduce the supermembrane action to a Goldstone superfield action with $N=2$, $d=3$ supersymmetry broken down to $N=1$, $d=3$.

2.1 The superembedding condition

For superembedding to be relevant to the description of the dynamics of superbranes one should find an appropriate condition of how a worldvolume supersurface is embedded into target superspace. The basic superembedding condition has a very simple geometrical meaning and is generic for all types of superbranes.

Consider a worldvolume supersurface $\mathcal{M}$ parametrized by $d=3$ bosonic coordinates ξ^m and 2 fermionic coordinates η^μ, which we will collectively call

$$z^M = (\xi^m,\ \eta^\mu), \quad m=0,1,2, \quad \mu=1,2. \tag{1}$$

The geometry of $\mathcal{M}$ is described, in a superdiffeomorphism invariant way, by a set of supervielbein one–forms

$$e^A(z) = dz^M e_M^{\ A} = (e^a(\xi,\eta),\ e^\alpha(\xi,\eta))\,, \tag{2}$$

which form a local basis in the cotangent space of $\mathcal{M}$. The indices a and α are, respectively, the indices of the vector and a spinor representation of the group $SO(1,2)$ of local rotations in the cotangent space.

We would like to embed this supersurface into a curved target superspace $\underline{M}$ parametrized by $D=4$ bosonic coordinates $X^{\underline{m}}$ and 4 fermionic coordinates $\Theta^{\underline{\mu}}$, which we will collectively call

$$Z^{\underline{M}} = (X^{\underline{m}},\ \Theta^{\underline{\mu}}), \quad \underline{m}=0,\ldots,3, \quad \underline{\mu}=1,\ldots,4. \tag{3}$$

Note that for embedding we have chosen a supersurface with the number of Grassmann–odd directions being half the number of target–superspace Grassmann–odd directions. This is for being able to identify **2** local worldvolume supersymmetries with **2** independent fermionic κ–symmetries of the standard (Green–Schwarz) formulation of supermembrane dynamics by Bergshoeff, Sezgin and Townsend [9].

The geometry of the target superspace is described in a superdiffeomorphism invariant way by a set of supervielbein one–forms

$$E^{\underline{A}}(Z) = dZ^{\underline{M}} E_{\underline{M}}^{\underline{A}} = \left(E^{\underline{a}}(X,\Theta), E^{\underline{\alpha}}(X,\Theta)\right), \tag{4}$$

which form a local frame in the cotangent space of the target superspace. The indices $\underline{a}$ and $\underline{\alpha}$ are, respectively, the indices of the vector and a spinor representation of the group $SO(1,D-1)$ of local rotations in the $\underline{M}$ cotangent space.

Superembedding is a map of $\mathcal{M}$ into $\underline{M}$ which is locally described by $X^{\underline{m}}$ and $\Theta^{\underline{\mu}}$ as functions of the supersurface coordinates

$$z^M \quad \to \quad Z^{\underline{M}}(z) = \left(X^{\underline{m}}(\xi,\eta),\ \Theta^{\underline{\mu}}(\xi,\eta)\right). \tag{5}$$

The map induces the pullback onto the supersurface of the target superspace one–form (4). In particular, the vector supervielbein $E^{\underline{a}}$ pullback is a one–superform on the supersurface. It has the following decomposition in the local basis (2) on $\mathcal{M}$

$$E^{\underline{a}}(z) = e^a(z) E_a{}^{\underline{a}}(Z(z)) + e^\alpha(z) E_\alpha{}^{\underline{a}}(Z(z)). \tag{6}$$

The superembedding condition we are interested in is the vanishing of the worldvolume spinor components of $E^{\underline{a}}(z)$

$$E_\alpha{}^{\underline{a}}(Z(z)) = 0. \tag{7}$$

In other words eq. (7) is a superfield constraint on (5) which singles out the superembeddings such that the pullback of the supervielbein $E^{\underline{a}}$ has non–zero components only along vector directions of the supersurface.

In addition we assume that the supersurface and target–superspace geometry satisfy torsion constraints of corresponding $d=3$ and $D=4$ supergravity

$$T^a = -ie^\alpha e^\beta \gamma^a_{\alpha\beta} + \cdots, \qquad T^{\underline{a}} = -iE^{\underline{\alpha}} E^{\underline{\beta}} \Gamma^{\underline{a}}_{\underline{\alpha\beta}} + \cdots, \tag{8}$$

which are consistent with the superembedding condition (γ^a and $\Gamma^{\underline{a}}$ are, respectively, $d=3$ and $D=4$ Dirac matrices). So, we deal with two supergravity theories embedded one into another. What is the role of the superembedding condition (7) then?

In some cases the superembedding condition produces only "kinematic" constraints (e.g. the Virasoro constraints for superstrings) and does not put superbrane dynamics on the mass shell. In these cases (as the $N=1$, $D=4$ supermembrane considered here) worldvolume superfield actions can be constructed.

In other cases the superembedding condition contains all the constraints and dynamical equations of motion of corresponding superbranes (e.g. a D=11 supermembrane [10] and an M5–brane [11]).

In all the cases the superembedding condition ensures that the superworldvolume geometry is induced by the embedding [1], i.e. that there is no fully fledged propagating supergravity on the brane worldvolume.

2.2 The superembedding action

To construct the supermembrane action we should introduce one more notion, namely, a target–superspace three–form gauge superfield $A_{\underline{NML}}(X,\Theta)$, whose field strength obeys the constraint[2]

$$F^{(4)} = dA^{(3)} = \frac{i}{2} E^{\underline{a}} E^{\underline{b}} \bar{E}_{\underline{\alpha}} E^{\underline{\beta}} (\Gamma_{\underline{ab}})^{\underline{\alpha}}_{\underline{\beta}} + \frac{1}{4!} E^{\underline{a}} E^{\underline{b}} E^{\underline{c}} E^{\underline{d}} F_{\underline{abcd}}. \tag{10}$$

It is well known that the supermembrane minimally couples to this gauge superfield, which in the Green–Schwarz formulation [9] is described by the Wess–Zumino term

$$S_{WZ} = \frac{T}{2} \int d^3\xi \varepsilon^{mnp} A_{mnp}(X,\Theta), \tag{11}$$

where T is the membrane tension and $A_{mnp}(X(\xi),\Theta(\xi)) = \partial_m Z^{\underline{M}} \partial_n Z^{\underline{N}} \partial_p Z^{\underline{P}} A_{\underline{PNM}}$ is the worldvolume pullback of $A^{(3)}$. In the superembedding formulation the $A^{(3)}$ pullback is a worldvolume three superform

$$A^{(3)}|_{\mathcal{M}} = \frac{1}{3!} e^A(z) e^B(z) e^C(z)\ A_{CBA}, \tag{12}$$

and the supermembrane action is constructed as an integral over the supersurface $\mathcal{M}\ (z^M = (\xi^m, \eta^\mu))$. It can be checked that because of dimensional reasons the only component of the $A^{(3)}$ pullback which can enter the action is the one with two spinor and one vector indices $A_{\alpha\beta a}$. Thus, we assume the supermembrane action in the superembedding approach to have the following form [12] (it is a "brany" generalization of an $N=1$ superstring action proposed in [13])

$$S = \frac{1}{3!} \int d^3\xi d^2\eta\ sdet\ e\ \gamma^{a\alpha\beta} A_{\alpha\beta a} + \int d^3\xi d^2\eta P^{\alpha}_{\underline{a}} E_{\alpha}^{\ \underline{a}}, \tag{13}$$

whose second Lagrange multiplier term takes care of the superembedding condition, and $sdet\ e$ is the superdeterminant of the worldvolume supervielbein matrix $e_M^A(z)$ (2).

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[1] In particular, it can be shown that the worldvolume metric is an induced metric defined in a standard way as

$$g_{mn}(\xi) = \partial_m Z^{\underline{M}} E_{\underline{M}}^{\underline{a}} \partial_n Z^{\underline{N}} E_{\underline{N}}^{\underline{b}} \eta_{\underline{ab}}|_{\eta=0}. \tag{9}$$

[2] Note that in $D=4$, since the dual field strength ${}^*F^{(4)} = const$ on the mass shell, $A^{(3)}$ does not have physical degrees of freedom, but its vacuum energy may contribute to the value of the cosmological constant.

Integrating over the Grassmann–odd variables and eliminating auxiliary fields with the use of the superembedding condition, one can check that this action reduces to the conventional supermembrane action [9].

The action (13) is invariant under the following symmetries:

i) worldvolume and target space superdiffeomorphisms

$$z^M \;\to\; f^M(\xi,\eta), \quad Z^{\underline{M}} \;\to\; f^{\underline{M}}(X,\Theta), \tag{14}$$

ii) super–Weyl transformations

$$e'^a = W^2(z)e^a, \quad e'^\alpha = W(z)e^\alpha - ie^a\gamma_a^{\alpha\beta}\mathcal{D}_\beta W, \tag{15}$$

iii) local $SO(1,2)$ rotations in the tangent space of the superworldvolume.

The action (13) describes the supermembrane in an arbitrary $N=1$, $D=4$ supergravity background[3], but since here we are interested in effects of spontaneous breaking of global supersymmetry, let us choose the superbackground to be flat. Then the supervielbeins and the gauge superfield take the form

$$E^{\underline{a}} = dX^{\underline{a}} - id\bar\Theta\Gamma^{\underline{a}}\Theta, \quad E^{\underline{\alpha}} = d\Theta^{\underline{\alpha}}, \tag{16}$$

$$A^{(3)} = i\bar\Theta\Gamma_{\underline{ab}}d\Theta(E^{\underline{a}}E^{\underline{b}} - iE^{\underline{a}}\bar\Theta\Gamma^{\underline{b}}d\Theta - \frac{1}{3}\bar\Theta\Gamma^{\underline{a}}d\Theta\bar\Theta\Gamma^{\underline{b}}d\Theta) \tag{17}$$

and the superembedding condition is

$$E_\alpha{}^{\underline{a}} = \mathcal{D}_\alpha X^{\underline{a}} - i\mathcal{D}_\alpha\bar\Theta\Gamma^{\underline{a}}\Theta = 0, \quad \mathcal{D}_A = (\mathcal{D}_\alpha, \mathcal{D}_a) = e_A{}^M(z)\partial_M, \tag{18}$$

where $e_A^M(z)$ is the inverse of the supervielbein matrix (2).

The integrability of (18) requires

$$\gamma^a_{\alpha\beta}E_a{}^{\underline{a}} = \gamma^a_{\alpha\beta}(\mathcal{D}_a X^{\underline{a}} - i\mathcal{D}_a\bar\Theta\Gamma^{\underline{a}}\Theta) = \mathcal{D}_\alpha\bar\Theta\Gamma^{\underline{a}}\mathcal{D}_\beta\Theta. \tag{19}$$

Using eqs. (18), (19) and the symmetriesed gamma–matrix identities in $D=4$

$$(\Gamma_{\underline{ab}}\Gamma^{\underline{b}})_{\{\underline{\alpha\beta\gamma\delta}\}} = 0, \quad (\Gamma_{\underline{a}}\Gamma^{\underline{a}})_{\{\underline{\alpha\beta\gamma}\}\delta} = 0, \tag{20}$$

one can reduce the supermembrane action in the flat target superspace to the following simple form

$$S = -\frac{i}{3!}\int d^3\xi d^2\eta\; det^{-1}(e_a{}^m)\;(\bar\Theta\Theta)\; E_a{}^{\underline{a}}E^{\;\underline{b}a}\eta_{\underline{ab}} + \int d^3\xi d^2\eta P^\alpha_{\underline{a}}E_\alpha{}^{\underline{a}}, \tag{21}$$

which resembles the Howe–Tucker–Polyakov term of the action for the p–branes, though it contains both the Nambu–Goto and the Wess–Zumino part of the conventional action.

[3]Actually, the action (13) also describes supermembranes in $D=5,7$ and 11 superspaces [12].

2.3 The physical gauge

We now gauge fix all the local symmetries of the supermembrane action.

The super–Weyl (15) and local $SO(1,2)$ transformations allow one to choose the spinor–spinor part of the worldvolume supervielbein to be the unit matrix

$$e_\alpha{}^\mu = \delta_\alpha{}^\mu.$$

The worldvolume diffeomorphisms are fixed by imposing a physical gauge. To this end we split the target space coordinates into the ones along and transverse to the membrane

$$X^{\underline{a}} = (X^a, X^3(\xi,\eta)), \quad \Theta^{\underline{\alpha}} = (\theta^\alpha, \Psi_\alpha(\xi,\eta)) \tag{22}$$

and identify X^a and θ^α with the superworldvolume coordinates

$$X^a = \xi^a, \quad \theta^\alpha = \eta^\alpha. \tag{23}$$

Note that in this way we also identify 1/2 of space–time supersymmetry (which shifts θ^α) with global worldvolume supersymmetry

$$\delta\theta^\alpha = \delta\eta^\alpha = \epsilon_1^\alpha. \tag{24}$$

It is this supersymmetry that remains unbroken.

The spontaneously broken part of space–time supersymmetry is the one which shifts $\Psi_\alpha(\xi,\eta)$. Thus $\Psi_\alpha(\xi,\eta)$ is the Goldstone fermion superfield associated with this symmetry. The form of its transformation under spontaneously broken supersymmetry is dictated by the requirement that the physical gauge conditions remain invariant under this symmetry. This transformation is easily found to be

$$\delta\Psi_\alpha = \epsilon_\alpha^2 + i(\bar\epsilon^2\gamma^m\Psi)\partial_m\Psi_\alpha. \tag{25}$$

We observe that broken supersymmetry is nonlinearly realized in the transformations of the Goldstone fermion.

The superfield $X^3(\xi,\eta)$ is the Goldstone scalar associated with spontaneously broken translations transverse to the membrane. $\Psi_\alpha(\xi,\eta)$ and $X^3(\xi,\eta)$ are not independent. The transverse ($\underline{a}=3$) component of the superembedding condition (18) relates them as follows

$$E_\alpha^3 = 0 \quad \Rightarrow \quad \mathcal{D}_\alpha\Phi(z) = 2i\Psi_\alpha(z), \tag{26}$$

where

$$\Phi = X^3 + i\eta^\alpha\Psi_\alpha, \quad \mathcal{D}_\alpha = \partial_\alpha + e_\alpha{}^m(z)\partial_m. \tag{27}$$

And the part of the superembedding condition parallel to the brane expresses the remaining independent supervielbein components $e_\alpha{}^m(z)$ in terms of Ψ_α. $e_\alpha{}^m(z)$ thus become induced by the superembedding, as we discussed in the Introduction

$$E_\alpha^m = 0 \quad \Rightarrow \quad e_\alpha{}^m = i\gamma^m_{\alpha\beta}\eta^\beta + iD_\alpha\bar\Psi\gamma^m\Psi - D_\alpha\bar\Psi\gamma^b\Psi\partial_b\bar\Psi\gamma^m\Psi$$

$$= i\gamma^m_{\alpha\beta}\eta^\beta + iD_\alpha\bar{\Psi}\gamma^b\Psi(\delta^m_b + i\partial_b\bar{\Psi}\gamma^m\Psi), \tag{28}$$

where

$$D_\alpha = \frac{\partial}{\partial\eta^\alpha} + i\eta^\beta\gamma^a_{\beta\alpha}\frac{\partial}{\partial\xi^a}, \quad \{D_\alpha, D_\beta\} = 2i\gamma^a_{\alpha\beta}\frac{\partial}{\partial\xi^a} \tag{29}$$

are covariant derivatives in a flat $N=1$, $d=3$ superspace.

We have thus shown that in the physical gauge the fluctuations of the $N=1$, $D=4$ supermembrane are described by a single unconstrained scalar superfield $\Phi(\xi,\eta)$, the Goldstone boson associated with broken translations in the direction transverse to the membrane. The Goldstone spinor $\Psi_\alpha(\xi,\eta)$ is expressed through $\Phi(\xi,\eta)$ and its derivatives in a highly nonlinear way. This is an example of a so called inverse Higgs effect [14]. The effect is that under certain covariant conditions the number of Goldstone fields gets reduced by making some of them dependent on the others. We should note that the inverse Higgs effect is only part of the superembedding condition, which is more general, and in particular also ensures the superworldvolume geometry to be induced by the embedding, as we have seen above.

We can now substitute the expressions, which we have found in the physical gauge, into the covariant supermembrane action and upon some calculations get the nonlinear Goldstone superfield action in the following form

$$S = iT\int d^3\xi d^2\eta\, \frac{\Psi^2}{1-\frac{1}{4}D^2\Psi^2} + T\int d^3\xi\cdot 1\,, \tag{30}$$

where the second term in (30) is the membrane ground state ($\Psi_\alpha = 0$) energy, $D^2 = D^\alpha D_\alpha$, $\Psi^2 = \Psi^\alpha\Psi_\alpha$ and the Goldstone fermion Ψ_α depends on the Goldstone scalar Φ (26), i.e. $\Psi_\alpha = D_\alpha\Phi + (\textit{nonlinear terms})$. Though the explicit expression of Ψ in terms of Φ is rather involved, it is nevertheless possible to get from (30) the action for the independent superfield $\Phi(\xi,\eta)$ [12]

$$S = -\frac{iT}{2}\int d^3\xi d^2\eta\, \frac{D^\alpha\Phi D_\alpha\Phi}{1-\frac{1}{8}(D^2\Phi)^2 + \sqrt{1+\partial_a\Phi\partial^a\Phi(1-\frac{1}{16}(D^2\Phi)^2)}} + T\int d^3\xi\cdot 1\,. \tag{31}$$

Eq. (31) is the superfield form of the gauge fixed component action for the $N=1$, $D=4$ supermembrane obtained in [2].

On the other hand, the action (30) can be reduced to the Goldstone superfield action constructed in [5] with the use of the method of a ‘linear’ realization of spontaneously broken supersymmetry. The action of [5] describes the dynamics of a Goldstone scalar superfield $\rho(\xi,\eta)$ different from $\Phi(\xi,\eta)$ of (31). The explicit relation between $\rho(\xi,\eta)$ and $\Phi(\xi,\eta)$ has not been presented in the literature. However, there exists [5] the expression of $\Psi_\alpha(\xi,\eta)$ in terms of $\rho(\xi,\eta)$

$$\Psi_\alpha = \frac{\zeta_\alpha}{1+D^2\mathcal{F}}, \quad \zeta_\alpha = D_\alpha\rho(\xi,\eta), \tag{32}$$

where

$$\mathcal{F} = \frac{1}{2}\frac{\zeta^2}{1+\sqrt{1+D^2\zeta^2}}. \tag{33}$$

Substituting (32) into (30) we get the action for $\rho(\xi,\eta)$ found in [5]

$$S = 2Ti\int d^3\xi d^2\eta \frac{(D\rho)^2}{1+\sqrt{1+D^2(D\rho)^2}} + T\int d^3\xi \cdot 1\,.$$

3 Conclusion

Starting from a covariant superembedding formulation of supermembrane dynamics in $N=1$, $D=4$ superspace and having gauge fixed the superworldvolume local symmetries we have got an effective nonlinear superfield theory on the brane superworldvolume which exhibits partial supersymmetry breaking. We have also demonstrated how the supermembrane is related to a model of partial breaking of $N=2$, $d=3$ supersymmetry discussed in [5]. The superembedding approach has provided us with a systematic way of doing this.

As a generalization of these results, it should be possible to get a $D=4$ Dirac–Born–Infeld action with partially broken $N=2$ supersymmetry from a superembedding formulation of a space–filling D3–brane in $N=2$, $D=4$ superspace, and thus to establish the direct relationship between the D3–brane and the Goldstone superfield actions of [3, 4, 6].

One may also address a problem of whether the conditions and symmetries associated with superembeddings may allow one to overcome ambiguities in the construction of the non–Abelian generalization of the Dirac–Born–Infeld theory, and hopefully to make a progress in the target space covariant description of the system of N coincident D–branes and its embedding into target *super*space. In a conventional "kappa–symmetric" approach a study of this problem has been undertaken in [15].

The methods of superembeddings can also be used to study the possibility of obtaining a simpler form of superconformal field theory actions on the Anti–de–Sitter boundary, as well as studying effects of local supersymmetry breaking in supergravity theories.

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