Received: *February 7, 2000,* Accepted: *February 15, 2000*

Automorphisms, non-linear realizations and branes

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Abstract: The theory of non-linear realizations is used to derive the dynamics of the branes of M-theory. A crucial step in this procedure is to use the enlarged automorphism group of the supersymmetry algebra recently introduced. The field strengths of the worldvolume gauge fields arise as some of the Goldstone fields associated with this automorphism group. The relationship to the superembedding approach is given.

Keywords: M-Theory, p-branes.

JHEP02(2000)024

Contents

1. Introduction

Many superbranes were constructed using a generalisation of the κ-symmetry [1] first found in the point particle and later in the string. These included the construction of the action for the twobrane of M-theory [2] and the D-branes in ten dimension [3]. However, the construction of the dynamics of the most sophisticated brane, the M-theory fivebrane, was achieved [4, 5, 6] using the superembedding formalism [7]. In fact, some the first brane actions to be constructed were for superbranes in six and two dimensions [8] using the method of non-linear realizations [9, 10]. This method has some advantages, namely the construction is entirely group theoretical and the fields have a very natural interpretation in terms of Goldstone particles for broken symmetries. There now exists a substantial literature [11, 12, 13, 20] on the construction of superbrane actions using this method. However, these works are largely confined to branes in dimensions lower than ten dimensions and the branes in eleven and ten dimensions have not been constructed using this method.

While the relationship between the superembedding approach and that using κ-symmetry is well understood in the sense that the κ-symmetry arises as part of the worldvolume supersymmetry which is automatically present in the superembedding approach, the relationship between these approaches and the approach of non-linear realizations is not clear although some discussion is given in reference [14].

It is the purpose of this paper to give a systematic method of deriving the dynamics of some of the more commonly used superbranes using the theory of non-linear realizations. A crucial step in this construction is to include in the group not only the usual Spin group, but its generalisation to a much larger automorphism that exists for supersymmetry algebras whose central charges form an arbitrary matrix. This group was introduced in reference [15] and applied within the context of the symmetries of the covariant equations of motion of the M-theory fivebrane. In particular, it was shown that the currents for supersymmetry translations and the second rank central charge arose by expressing the equations of motion for the fermion, scalars and second rank tensor gauge field respectively as total derivatives. This suggests that even this brane possessed a construction as a non-linear realization. It was also shown [15] that the fivebrane possessed a further symmetry with a rank three generator that was part of the GL(32) automorphism group of the eleven dimensional supersymmetry algebra with all its central charges. The extra generator was associated with the two form gauge field of the fivebrane. These large automorphism groups are only present if one includes sufficient central charges in the supersymmetry algebra. Indeed, if one includes all the possible central charges in the superalgebra, then the automorphism group is $\mathrm{GL}(c_d)$ where c_d are the number of Majorana supercharges. This is the natural generalisation to branes of the Spin and internal symmetry groups that occur if one considers only point particles and so has only a momentum generator in the anti-commutator of the supercharges. In the past, many constructions of branes have not taken into account in the coset construction any automorphisms of the supersymmetry algebra. Some notable exceptions that did include the well known Lorentz and internal symmetry groups were given in references [11, 20, 24] which treated some superbranes in six and four dimensions and the point particle in lower dimensions respectively. As we shall see, to construct the branes of M-theory using the method of non-linear realisations one must use the much larger automorphism groups. This is consistent with the old result that the topological charges of branes occur as the central charges in the supersymmetry algebra [23]. In the approach of non-linear realizations given in this paper, the worldvolume fields of the superbrane are certain of the Goldstone fields associated with this automorphism group.

We will also find the relationship between the superembedding approach and that of the non-linear realizations. In the latter one first derives the Cartan forms from the group elements and then constructs from them a group invariant dynamics. It was noticed long ago that for certain choices of groups one could set certain of the Cartan forms to vanish in a group invariant manner. This was the so-called inverse Higgs effect [16]. We will show that for the supergroups we consider, we can choose such a constraint on the Cartan forms which is precisely equivalent to the condition assumed in the superembedding formalism.

2. The theory of non-linear realizations

We first briefly review the theory of non-linear realizations set out by Coleman, Wess and Zumino [9] and extended by Volkov [10]. Rather than consider a general group we restrict ourselves to the (super) group $\underline{G}$ whose generators can be divided into the two

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sets $\underline{K}$ and $\underline{H}$. The sets $\underline{K}$ and $\underline{H}$ are both supgroups of $\underline{G}$ and $\underline{H}$ is the automorphism group, of $\underline{K}$. We denote the generic generators of $\underline{K}$ also by the same symbol, $\underline{K}$ and label them by the index $\underline{A}$ i.e. by $K_{\underline{A}}$, for an appropriate index range. We label the generic generators of $\underline{H}$ by the symbol $\underline{R}$ with appropriate indices. The generators in each of the above two classes are further divided as $\underline{K} = \{K, K'\}$ and $\underline{R} = \{R, R'\}$. The generators K and R both form subalgebras of the Lie (super) algebra G whose associated groups we denote by K and H respectively. The generators of H are an automorphisms of K. The division of the generators of $\underline{G}$ into these four classes corresponds to:

- K: unbroken generators associated with positions in (super) space-time,
- K': spontaneously broken generators associated with positions in (super) space-time,
- R: unbroken automorphism generators,
- R': spontaneously broken auotmorphism generators.

The automorphism generators include those of the Spin group or its covering group the Lorentz, in some cases internal group generators, but also other generators which are not usually considered.

From now on we will drop the prefix (super), but it is to be understood to be present. As will be clear, the objects associated with the group $\underline{G}$ carry an underline while those associated with the subgroup G carry no underline. The decomposition of the former into the latter and the remainder is achieved using unprimed and primed indices respectively.

We wish to consider the coset $\underline{G}/H$. For simplicity we will use an exponential description of group elements and we may then choose the coset representatives to be chosen as

$$g = e^{X \cdot K} e^{X' \cdot K'} e^{\phi' \cdot R'} . \tag{2.1}$$

In this equation $\cdot$ denotes the relavent summation over the indices. Under any rigid group transformation g_0 we find that

$$g \to \; g_0 g = \hat{g} = e^{\hat{X} \cdot K} e^{\hat{X}' \cdot K'} e^{\hat{\phi}' \cdot R'} h_o \, , \tag{2.2}$$

where h_0 is an element of H. The Cartan forms are given by

$$g^{-1} dg = F \cdot K + F' \cdot K' + \omega' \cdot R' + \omega \cdot R \, . \tag{2.3}$$

Under equation (2.2) the Cartan forms transform as

$$\hat{g} d\hat{g} = h_0 (g^{-1} dg) h_0^{-1} + h_0 d h_0^{-1} \, . \tag{2.4}$$

It can happen that the Cartan forms carry a reducible representation of H, in which case, certain of the forms can be set to zero. This is the so called inverse Higgs effect [16]. It has the effect of eliminating some of the Goldstone fields in terms of the others. The action or equations of motion are to be constructed from the Cartan forms in such a way that they are invariant under the transformations of equation (2.2).

The conventional interpretation of the above equations is to regard the X as the coordinates of (super) space-time and to take the fields X' and ϕ to depend on them. This leads to a field theory on the coset space G/H. This approach has been almost universally adopted. However, when considering branes it is more instructive to consider a more general possibility. The brane is moving through the coset space $\underline{G}/\underline{H}$ with tangent group $\underline{H}$ and sweeps out a submanifold that has the dimensions of the coset G/H and a tangent space group H. We therefore consider the representatives of the coset $\underline{G}/H$ of equation (2.2) to depend on the variable ξ which parametrises the embedded submanifold. Since the Cartan forms involve the exterior derivative d they are independent of the coordinate system used. The Cartan forms associated with K are given by $F \cdot K = d\xi \cdot F \cdot K$ and similarly for the other Cartan forms. We can think of the F in this equation as the vielbein on the embedded submanifold. The covariant derivatives of the Goldstone fields associated with K' are defined by

$$F' \cdot K' \equiv F \cdot \Delta X' \cdot K' = d\xi \cdot F^{-1} \cdot F' \cdot K' \,. \tag{2.5}$$

The $\Delta X'$ are independent of reparamenterisations. A similar construction can be made for the covariant derivatives of ϕ'. When identifying the fields that can be set to zero, i.e. use the inverse Higgs mechanism, we must not only maintain the $\underline{G}$ invariance of equation (2.2), but also reparameterisation invariance. In effect, this means setting only those covariant derivatives of the Goldstone fields that transform in a covariant manner under h_0 to zero. The equations of motion, or action, are to be constructed from the vielbein and the covariant derivatives of the Goldstone fields that remain. In this way one can find a formulation of brane dynamics that is reparameterisation invariant and also invariant under the rigid $\underline{G}$ transformations of equation (2.2). From this approach we can recover the more conventional approach by using the reparameterisation invariance to choose static gauge, i.e. $X = \xi$ for those coordinates that lie in the brane directions.

3. Bosonic branes

In this section we will show that the dynamics of bosonic p-branes in a flat background in D dimensional space-time arises as a non-linear realization in the sense of the previous section. We take $\underline{G} = \mathrm{ISO}(D-1,1)$ with $\underline{K} = \{P_{\underline{n}}\}$ and $\underline{H} = \{J_{\underline{nm}}\}$

where $\underline{n}, \underline{m} = 0, 1, \ldots, D-1$ and $G = \mathrm{ISO}(p,1)$ with $K = \{P_n\}$ and $H = \{J_{nm}\}$ where $n, m = 0, 1, \ldots, p+1$. This is to be expected as the presence of the p-brane clearly breaks the background space-time group $\mathrm{ISO}(D-1,1)$ to $\mathrm{ISO}(p,1)$. The Lie algebra of $\mathrm{ISO}(D-1,1)$ is given by

$$[J_{\underline{nm}}, J_{\underline{pq}}] = -\eta_{\underline{np}} J_{\underline{mq}} - \eta_{\underline{mq}} J_{\underline{np}} + \eta_{\underline{nq}} J_{\underline{mp}} + \eta_{\underline{mp}} J_{\underline{nq}} \tag{3.1}$$

$$[P_{\underline{n}}, J_{\underline{pq}}] = +\eta_{\underline{np}} P_{\underline{q}} - \eta_{\underline{nq}} P_{\underline{p}} \,. \tag{3.2}$$

We can write the coset representatives in the form

$$\begin{aligned} g(X,\phi) &= \exp(X^n P_n + X^{n'} P_{n'}) \exp(\phi^{nm'} J_{nm'} + \phi^{n'm'} J_{n'm'}) \\ &\equiv \exp(X^{\underline{n}} P_{\underline{n}}) \exp(\phi \cdot J) \,. \end{aligned} \tag{3.3}$$

We distinguish X from X' by the indices they carry, in other words the prime on the X is understood to be present and we just write $X^{n'}$. The Cartan forms are given by

$$g^{-1} dg \equiv e^n P_n + f^{n'} P_{n'} + \Omega^{nm'} J_{nm'} + \omega^{n'm'} J_{n'm'} + w^{nm} J_{nm} \,, \tag{3.4}$$

which we may express as

$$g^{-1} dg = \exp(-\phi \cdot J)(dX^{\underline{n}} P_{\underline{n}}) \exp(\phi \cdot J) + \exp(-\phi \cdot J) d \exp(\phi \cdot J) \,. \tag{3.5}$$

A straightforward calculation shows that

$$\begin{aligned} e^n &= dX^{\underline{m}} \Phi_{\underline{m}}{}^n = -2\phi^{nm'} d\phi_{m'} + dX^n + \cdots \\ f^{n'} &= dX^{\underline{m}} \Phi_{\underline{m}}{}^{n'} = 2\phi^{m'n'} d\phi_{m'} - 2\phi^{n'm} dX_m + dX^{n'} \\ f^{nm'} &= d\phi^{nm'} + 2\left(-\phi^{p'm'} d\phi_{p'}{}^{n} + \phi^{p'n} d\phi_{p}{}^{m'}\right) \\ f^{n'm'} &= d\phi^{n'm'} + \left(\left(\phi^{p'n'} d\phi_{p'}{}^{m'} + \phi^{pn'} d\phi_{p}{}^{m'}\right) - (n' \leftrightarrow m')\right) \\ w^{nm} &= \left(\phi^{p'n} d\phi_{p'}{}^{m} - (n \leftrightarrow m)\right) , \end{aligned} \tag{3.6}$$

where $\Phi_{\underline{n}}{}^{\underline{m}}$ is defined by $\exp(-\phi \cdot J) P_{\underline{n}} \exp(\phi \cdot J) = \Phi_{\underline{n}}{}^{\underline{m}} P_{\underline{m}}$ and $+\cdots$ means higher order terms in $\phi^{\underline{nm}}$.

Under a group transformation, $g_0 g(X,\phi) = \hat{g}(\hat{x}\hat{\phi}) h_0$, the Cartan forms transform by the J_{nm} transformations with associated parameter r^{nm} in accord with equation (2.4). The effect of the $P_{\underline{n}}$ transformations in g_0 is to simply transform the $X^{\underline{n}}$. Writing $h_0 = 1 + r^{nm} J_{nm}$. we find that

$$\hat{e}^n = e^n - e^p r_p{}^n \,, \qquad \hat{f}^{n'} = f^{n'} \,, \qquad \hat{\Omega}^{nm'} = \Omega^{nm'} + 2r^{pn} \Omega_p{}^{m'} - 2r^{pm'} \Omega_p{}^{n} \,, \tag{3.7}$$

$$\hat{w}^{nm} = w^{nm} - 2(r^{np} w_p{}^m - (n \leftrightarrow m)) + dr^{nm} \tag{3.8}$$

to lowest order in r^{nm}.

Clearly, we can set $f^{n'} = 0$ and preserve $\mathrm{SO}(1, D-1)$ and reparameterisation symmetry. At the linearized level, examining equation (3.6) we find it implies that $dX^{n'} = 2\phi^{n'}{}_m dX^m$ or

$$2\phi^{n'}{}_m = \frac{\partial \xi^p}{\partial X^m}\frac{\partial X^{n'}}{\partial \xi^p}\,. \tag{3.9}$$

If we choose static gauge this equation becomes,

$$2\phi^{n'}{}_m = \frac{\partial X^{n'}}{\partial X^m}\,, \tag{3.10}$$

While solving for $f^n = 0$ to all orders may be complicated it is clear that its content is to solve for $\phi^{n'}{}_m$ in terms of $\partial_m \phi^{n'}$.

What is really of interest to us is the non-linear form of f^n once we have solved this constraint. Examining equation (3.6) we find that the vector $f^{\underline{n}} \equiv (f^n, f^{n'})$ is related by a Lorentz transformation to the vector $(dX^n, dX^{n'}) = dX^{\underline{n}}$. As such,

$$e_p{}^n \eta_{nm} e_q{}^m = \partial_p X^{\underline{n}} \partial_q X^{\underline{m}} \eta_{\underline{nm}} \tag{3.11}$$

since $f^{n'} = 0$. The above expression is invariant under g_0 transformations. A world-volume reparameterisation and group invariant action is therefore given by

$$\int d^p\xi \det e = \int d^p\xi \sqrt{-\det(\partial_p X^{\underline{n}} \partial_q X^{\underline{m}} \eta_{\underline{nm}})} \tag{3.12}$$

In other words, the well known generalisation of the Nambu action for the string to a p-brane follows in a straightforward consequence of taking the non-linear realization $\mathrm{ISO}(D-1, 1)$ with subgroup $\mathrm{SO}(p, 1)$.

Clearly, had we not included the Lorentz group in our coset and introduced the corresponding Goldstone bosons the veilbein on the brane would have been trivial and we would not have found the above action.

4. Superbranes

For superbranes we consider the supergroup $\underline{G}$ considered in [15] and applied in the context of M-theory. In accord with section 2 it has the two subgroups $\underline{K}$ and $\underline{H}$. The generators of $\underline{K}$ obey the relations

$$\{Q_{\underline{\alpha}}, Q_{\underline{\beta}}\} = Z_{\underline{\alpha\beta}}\,, \qquad [Q_{\underline{\gamma}}, Z_{\underline{\alpha\beta}}] = 0\,, \qquad [Z_{\underline{\alpha\beta}}, Z_{\underline{\gamma\delta}}] = 0\,, \tag{4.1}$$

while those of the automorphism group $\underline{H}$ satisfy

$$[Q_{\underline{\alpha}}, R_{\underline{\gamma}}{}^{\underline{\delta}}] = \delta_{\underline{\alpha}}{}^{\underline{\delta}} Q_{\underline{\gamma}}\,, \qquad [Z_{\underline{\alpha\beta}}, R_{\underline{\gamma}}{}^{\underline{\delta}}] = \delta_{\underline{\alpha}}{}^{\underline{\delta}} Z_{\underline{\gamma\beta}} + \delta_{\underline{\beta}}{}^{\underline{\delta}} Z_{\underline{\alpha\gamma}}\,. \tag{4.2}$$

The generators $Z_{\underline{\alpha\beta}}$ are Grassmann even generators labeled by spinor indices. They are obviously symmetric in these indices. It is easy to verify that such an algebra obeys the generalized super Jacobi identities. Let us assume that the generators $Z_{\underline{\alpha\beta}}$ form the most general symmetric matrix. Expanding this matrix out in terms of the enveloping algebra of the generators of the relevant Clifford algebra we find that $Z_{\underline{\alpha\beta}}$ contains a set of totally anti-symmetric tensorial generators which constitute the central charges:

$$Z_{\underline{\gamma}}{}^{\underline{\delta}} = \sum_p \sum_{\underline{n}_1\cdots\underline{n}_p} (\gamma^{\underline{n}_1\cdots\underline{n}_p} C^{-1})_{\underline{\gamma}}{}^{\underline{\delta}} Z_{\underline{n}_1\cdots\underline{n}_p} . \tag{4.3}$$

The central charges include the momenta $P_{\underline{n}} \equiv Z_{\underline{n}}$. Both the eleven dimensional and the six dimensional superalgebras associated with the fivebrane equations of motion have a supersymmetry algebra in which the central charges form the most general matrix [18, 15]. In the latter case one must take an appropriate index range for the indices $\underline{\alpha}, \underline{\beta}, \ldots$ The corresponding superalgebras in the IIA and IIB theories for which the $Z_{\underline{\alpha\beta}}$ form the most general matrix in ten dimensions are well known.

The automorphism group is $\mathrm{GL}(c_d)$ where c_d is the number of supercharges. The precise properties of the matrices under complex conjugation being given by implementing the Majorana or other properties of the spinorial supercharge on the commutator relation of the supercharge with the automorphisms. For the case of the eleven dimensional superalgebra it is GL(32)

To gain a more familiar set of generators we may expand $R_{\underline{\gamma}}{}^{\underline{\delta}}$ out in terms of the elements of the enveloping Clifford algebra

$$R_{\underline{\gamma}}{}^{\underline{\delta}} = \sum_p \sum_{\underline{n}_1\cdots\underline{n}_p} (\gamma^{\underline{n}_1\cdots\underline{n}_p})_{\underline{\gamma}}{}^{\underline{\delta}} R_{\underline{n}_1\cdots\underline{n}_p} . \tag{4.4}$$

If the only central charge is the momentum then the right-hand side of the anti-commutator of two supercharges is $\gamma^{\underline{n}} P_{\underline{n}}$. In this case, the most general automorphism group is by definition the spin group. Hence, the automorphism group GL(32) is the natural generalisation to branes of the spin group relavent to point particles. In the above expansion the generators $R_{\underline{nm}}$ are those of the Spin Lie algebra.

The automorphism group has been found to be relavent to the dynamics of the fivebrane, In particular, its covariant equations of motion [4] have been found [15] to possess additional conserved currents beyond those associated with Spin transformations, central charges and supersymmetry. These new currents lead to conserved charges that generate, at the classical linearized level, a subalgebra of the automorphism group.

We may also consider a subgroup of the full automorphism group. One such is given by the set of symmetric matrices, these form Lie algebra of $\mathrm{Sp}(c_d)$. The generators of this subalgebra are $S_{\underline{\gamma\delta}} = R_{\underline{\gamma\delta}} + R_{\underline{\delta\gamma}}$ where $R_{\underline{\gamma\delta}} = R_{\underline{\gamma}}{}^{\underline{\beta}} (C^{-1})_{\underline{\beta\delta}}$. Clearly, in the decomposition of equation (4.4) this means keeping only those terms for which the symmetric matrices enter. In eleven dimensions these are the generators $R_{\underline{n}_1\cdots\underline{n}_p}$ of ranks one, two and five. The commutator of the generators $S_{\underline{\gamma\delta}}$ with the supercharges is given by

$$[Q_{\underline{\alpha}}, S_{\underline{\gamma\delta}}] = (C^{-1})_{\underline{\alpha\delta}} Q_{\underline{\gamma}} + (C^{-1})_{\underline{\alpha\gamma}} Q_{\underline{\delta}} \,. \tag{4.5}$$

We divide the generators in $\underline{K}$ into $Q_{\underline{\alpha}} = (Q_\alpha, Q_{\alpha'})$ and $Z_{\underline{\alpha\beta}} = (Z_{\alpha\beta}, Z_{\alpha\beta'}, Z_{\alpha'\beta'})$ and the generators of $\underline{H}$ as $R_{\underline{\gamma}}{}^{\underline{\beta}} = (R_\gamma{}^\beta, R_{\gamma'}{}^\beta, R_\gamma{}^{\beta'}, R_{\gamma'}{}^{\beta'})$. We take the Lie superalgebra of the group G to contain the generators Q_α and $Z_{\alpha\beta}$ as well as some of the generators $R_\gamma{}^\delta$

Initially, we will take the non-linear realization approach to the superbrane in which the coset representatives to be parameterised by the coordinates on the brane and the transverse fields. Proceeding in this way, we will find the expressions for the superbrane in static gauge. As such, the coset representatives of $\underline{G}/H$ are written as

$$g = e^{X^{\alpha\beta} Z_{\alpha\beta} + \theta^\alpha Q_\alpha} e^{X^{\alpha\beta'} Z_{\alpha\beta'} + X^{\alpha'\beta'} Z_{\alpha'\beta'} + \Theta^{\alpha'} Q_{\alpha'}} e^{\phi\cdot R} \,, \tag{4.6}$$

where $\phi' \cdot R'$ is a sum that includes all the generators in $\underline{H}$, except for those in H. The Cartan forms are given by

$$\begin{aligned} g^{-1}dg = {} & e^{-\phi'\cdot R'} de^{\phi'\cdot R'} + e^{-\phi'\cdot R'} \Big(Y^{\alpha\beta} Z_{\alpha\beta} + Y^{\alpha\beta'} Z_{\alpha\beta'} + Y^{\alpha'\beta'} Z_{\alpha'\beta'} + \\ & \qquad + d\theta^\alpha Q_\alpha + d\Theta^{\alpha'} Q_{\alpha'} \Big) e^{\phi'\cdot R'} \\ \equiv {} & F^{\alpha\beta} Z_{\alpha\beta} + F^{\alpha\beta'} Z_{\alpha\beta'} + F^{\alpha'\beta'} Z_{\alpha'\beta'} + F^\alpha Q_\alpha + F^{\alpha'} Q_{\alpha'} + e^{-\phi'\cdot R'} de^{\phi'\cdot R'} \,, \end{aligned} \tag{4.7}$$

where $Y^{\alpha\beta} = dX^{\alpha\beta} + \frac{1}{4}(\theta^\alpha d\theta^\beta + \theta^\beta d\theta^\alpha)$, $Y^{\alpha\beta'} = dX^{\alpha\beta'} - d\theta^\alpha d\Theta^{\beta'}$ and $Y^{\alpha'\beta'} = dX^{\alpha'\beta'} + \frac{1}{4}(\Theta^{\alpha'} d\Theta^{\beta'} + \Theta^{\beta'} d\Theta^{\alpha'})$

In order to gain some understand the physical content of the formalism we will begin by giving a linearized analysis. To lowest order in $\phi_{\underline{\delta}}{}^{\underline{\gamma}}$ the Cartan forms are given by

$$\begin{aligned} F^\alpha &= d\theta^\alpha + d\Theta^{\delta'} \phi_{\delta'}{}^\alpha + d\theta^\delta \phi_\delta{}^\alpha \,, \qquad F^{\alpha\beta} = Y^{\alpha\beta} + Y^{\alpha\delta'} \phi_{\delta'}{}^\beta + Y^{\alpha\delta} \phi_\delta{}^\beta + Y^{\beta\delta} \phi_\delta{}^\alpha \,, \\ F^{\alpha'} &= d\Theta^{\alpha'} + d\Theta^{\delta'} \phi_{\delta'}{}^{\alpha'} + d\theta^\delta \phi_\delta{}^{\alpha'} \,, \\ F^{\alpha'\beta'} &= Y^{\alpha'\beta'} + Y^{\alpha'\delta'} \phi_{\delta'}{}^{\beta'} + Y^{\beta'\delta'} \phi_{\delta'}{}^{\alpha'} + Y^{\alpha'\delta} \phi_\delta{}^{\beta'} + Y^{\beta'\delta} \phi_\delta{}^{\alpha'} \,, \\ F^{\alpha\beta'} &= Y^{\alpha\beta'} + Y^{\alpha\delta'} \phi_{\delta'}{}^{\beta'} + Y^{\beta'} \phi_{\delta'}{}^\alpha + Y^{\alpha\delta} \phi_\delta{}^{\beta'} + Y^{\delta\beta'} \phi_\delta{}^\alpha \,. \end{aligned} \tag{4.8}$$

We wish to place constraints on the Cartan forms that respect the group $\underline{G}$ and lead to the equations of motion of the superbrane. Taking the Goldstone fields to be of order one and the coordinates $X^{\alpha\beta}$ and θ^α to be of order zero, we will first carry out a linearized analysis. As explained in section 2, the equation of motion are given by setting certain of the covariant derivatives in the Goldstone fields to vanish. As such, we will only require the supervielbein to order zero. It is straightforward to read it off from equation (4.8) and we find that

$$F = \begin{pmatrix} \frac{1}{2}(\delta_\gamma^\alpha \delta_\delta^\beta + \delta_\gamma^\beta \delta_\delta^\alpha) & 0 \\ -\frac{1}{4}(\delta_\gamma^\alpha \theta^\beta + \delta_\gamma^\beta \theta^\alpha) & \delta_\gamma^\alpha \end{pmatrix} . \tag{4.9}$$

We can also read off the linearized Cartan forms form equation (4.8) and multiplying the result by the inverse supervielbein we find that the covariant derivatives of the Goldstone fields are given by

$$
\begin{aligned}
\nabla_\alpha \Theta^{\alpha'} &= D_\alpha \Theta^{\alpha'} + \phi_\alpha{}^{\alpha'} \,, & \nabla_{\gamma\delta} \Theta^{\alpha'} &= \partial_{\gamma\delta} \Theta^{\alpha'} \,, \\
\nabla_\gamma X^{\alpha\beta'} &= D_\gamma X^{\alpha\beta'} - \delta_\gamma^\alpha \Theta^{\beta'} \,, & \nabla_{\gamma\delta} X^{\alpha\beta'} &= \partial_{\gamma\delta} X^{\alpha\beta'} + \frac{1}{2}(\delta_\gamma^\alpha \phi_\delta{}^{\beta'} + \delta_\delta^\alpha \phi_\gamma{}^{\beta'}) \,,
\end{aligned}
\tag{4.10}
$$

where

$$
D_\alpha \equiv \partial_\alpha + \frac{1}{2}\theta^\delta \partial_{\gamma\delta} \,, \qquad \partial_{\gamma\delta} \equiv \frac{\partial}{\partial X^{\gamma\delta}} \,. \tag{4.11}
$$

It is a consequence of the linearized analysis that $\phi_\gamma{}^{\beta'}$ is the only part of $\phi_{\underline{\gamma}}{}^{\underline{\beta}}$ that enters any of these expressions. The covariant derivatives of the other Goldstone fields will not interest us here.

We now choose the group H to be just the Lorentz group. The Cartan forms carry a very reducible representation with respect to this subgroup and hence with respect to $\underline{G}$. Consequently, we may choose many of the covariant derivatives of the Goldstone fields to vanish. In fact, we will choose the set

$$
\nabla_\gamma \Theta^{\alpha'} = 0 \tag{4.12}
$$

$$
\nabla_\gamma X^{\alpha\beta'} (\gamma^{n'} C^{-1})_{\alpha\beta'} = 0 \tag{4.13}
$$

and

$$
\nabla_{\gamma\delta} \theta^{\alpha'} - \frac{2}{r_p} (\gamma^n C^{-1})_{\gamma\delta} (\gamma^n C^{-1})^{\alpha\beta} \nabla_{\alpha\beta} \theta^{\alpha'} = 0 \,, \tag{4.14}
$$

where r_p is the number of supercharges in G. The superspaces defined by the above coset space construction contain many central charge associated coordinates in addition to the usual Minkowski coordinates. The Grassmann even coordinates of the background superspace $\underline{G}/\underline{H}$ include the coordinates $X^{\underline{n}} \equiv X^{\underline{\alpha}\underline{\beta}'} (\gamma^{\underline{n}} C^{-1})_{\underline{\alpha}\underline{\beta}}$ of Minkowski space, but also all the other coordinates corresponding to the expansion of the supercharge of equation (4.3). Similarly, the coset space G/H on which the above superfields are defined includes the Minkoswki space coordinates $X^n \equiv X^{\alpha\beta'} (\gamma^n C^{-1})_{\alpha\beta}$ as well as a corresponding number of other Grassmann even coordinates associated with the central charges in G other than the momentum.

Equation (4.12) sets the central charge dependence of $\Theta^{\beta'}$ to be trivial except for the dependence on the coordinates X^n. Equation (4.13) then implies that $X^{n'}$ has the same dependence. As such, the above equations are equivalent to the conditions

$$
D_\alpha \Theta^{\beta'} = -\phi_\alpha{}^{\beta'} \tag{4.15}
$$

and

$$
D_\gamma X^{n'} = (\gamma^n C^{-1})_{\gamma\delta'} \Theta^{\delta'} \,, \tag{4.16}
$$

where all the superfields $\Theta^{\beta'}$ $X^{n'}$ and $\phi_\alpha{}^{\beta'}$ depend only on X^n and θ^α. It is known [17] from the superembedding approach that these equation imply the correct field content and equations of motion for most of the superbranes including the 2 and 5-branes of M-theory.

Let us study the consequences of these equations for the branes in M theory. We must decompose the eleven dimensional Clifford algebra to one that keeps manifest the Clifford algebra appropriate to the brane. This results in a corresponding decomposition of the spinor index $\underline{\alpha} = (\alpha, i)$. For the fivebrane $\alpha = 1, \ldots 4$ are the Weyl projected spinor indices of Spin$(1,5)$ and $i = 1, \ldots 4$ are the indices of the internal group USp(4) = Spin(5) while for the twobrane $\alpha = 1, 2$ are the spinor indices of Spin$(1,2)$ and $i = 1, \ldots 8$ are indices of Spin(8). The analysis of equation (4.16) depends only on the properties of this decomposition of the eleven dimensional Clifford algebra. For the fivebrane of M-theory it implies that [17]

$$D_{\alpha i}\Theta_\beta{}^j \sim -(\not\partial)_{\alpha\beta}(\gamma_{n'})_i{}^j X^{n'} + \delta_i^j(\gamma_{n_1n_2n_3})_{\alpha\beta}h^{n_1n_2n_3}\,, \tag{4.17}$$

where $h_{n_1n_2n_3}$ is the self-dual gauge field strength of the fivebrane. At the linearized level, the field $h_{n_1n_2n_3}$ obeys the Bianchi identity implying it is the curl of a second rank gauge field in addition to being self-dual. This relation follows from the constraint of equation (4.16) by applying further covariant derivatives.

For the twobrane we find that

$$D_{\alpha i}\Theta_\beta{}^j \sim -(\not\partial)_{\alpha\beta}(\gamma_{n'})_i{}^j X^{n'}\,. \tag{4.18}$$

Equation (4.15) implies that the Goldstone field $\phi_\alpha^{\beta'}$ associated with the automorphism group $\underline{H}$ is non-vanishing. Expanding

$$\phi_{\underline{\gamma}}{}^{\underline{\delta}} = \sum_p \sum_{\underline{n}_1\cdots\underline{n}_p}(\gamma^{\underline{n}_1\cdots\underline{n}_p})_{\underline{\gamma}}{}^{\underline{\delta}}\phi_{\underline{n}_1\cdots\underline{n}_p}\,. \tag{4.19}$$

we find that both for the fivebrane and the twobrane that $\phi_n{}^{m'} \sim \partial_n X^{m'}$ which is the same as the condition of equation (3.10) for the bosonic brane. It relates the Goldstone fields for the broken Lorentz generators to the derivative of Goldstone fields of broken translations. However, for the fivebrane we find in addition that the field strength for the self-dual gauge field is related to the Goldstone fields for broken generators of the automorphism group GL(32), in particular $h_{n_1n_2n_3} \sim \phi_{n_1n_2n_3}$. We note that these generators are the R_{nmp} contained in equation (4.4) and that these generators are not contained in the subgroup Sp(32). This is consistent with the existence of the additional rank three conserved current found [15] in the covariant equations of motion of the fivebrane.

The non-linear realization described above contains more Goldstone superfields than those of equations (4.15) and (4.16). One could consider imposing more constraints generalizing that of equation (4.16) to include the covariant derivatives of the

Goldstone fields associated with the central charges. These would determine these additional fields in terms of the fields that we have already. In particular, one could impose a constraint the implies that the Goldstone field associated with the central charge Z_{nm} is the second rank gauge field of the fivebrane.

It is not clear that the choice of subgroup H and the related constraints of equations (4.12) to (4.14) are the unique possible choice that leads to the correct dynamics. Indeed, it would be interesting to explore other choices. In particular, one could also consider constraints which allow the superfields to depend on the central charges of the six dimensional subalgebra, but in a rather restricted way.

To make contact with the superembedding formalism, and to carry out an all orders analysis, it will be convenient to adopt a different parameterisation of the coset representatives. As explained in section 2, we can take the Goldstone fields to depend on the coordinates of the superembedded manifold and treat the transverse and longitudinal coordinates of the superbrane on an equal footing. This allows us to work in a superreparameterisation invariant manner and not choose superstatic gauge as above. As such, we choose our coset representatives to be

$$g = e^{X^{\underline{\alpha\beta}} Z_{\underline{\alpha\beta}} + \Theta^{\underline{\alpha}} Q_{\underline{\alpha}}} e^{\phi' \cdot R'} \ , \tag{4.20}$$

where the coordinates $Z^{\underline{M}} \equiv (X^{\underline{\alpha\beta}}, \Theta^{\underline{\alpha}})$, not be confused with the central charges $Z_{\underline{\alpha\beta}}$, depend on the coordinates, $z^M = (\xi^{\alpha\beta}, \theta^\alpha)$ that parameterise the superbrane. The transformation from the fields of equation (4.6) to that in the above equation is obtained by choosing superstatic gauge $X^{\alpha\beta} = \xi^{\alpha\beta}$, $\Theta^\alpha = \theta^\alpha$ and making the shift

$$X^{\alpha\beta'} \to X^{\alpha\beta'} + \frac{1}{2}\theta^\alpha \Theta^{\beta'} \ , \tag{4.21}$$

with all the other fields being the same. Constructing the Cartan forms as before, we find that

$$\begin{aligned} g^{-1}dg &= e^{-\phi' \cdot R}(E^{\underline{A}} K_{\underline{A}}) e^{\phi' \cdot R'} + e^{-\phi' \cdot R'} d e^{\phi' \cdot R'} \\ &\equiv F^{\underline{A}} K_{\underline{A}} + e^{-\phi' \cdot R'} d e^{\phi' \cdot R'} \ . \end{aligned} \tag{4.22}$$

In this equation the index $\underline{A} = (\underline{\gamma}\underline{\delta}, \underline{\gamma})$ and the one forms $E^{\underline{A}} \equiv dZ^{\underline{\Pi}} E_{\underline{M}}{}^{\underline{A}}$ contain the supervielbeins $E_{\underline{M}}{}^{\underline{A}}$ of the background superspace and have the form

$$E_{\underline{M}}{}^{\underline{A}} = \begin{pmatrix} \frac{1}{2}(\delta_{\underline{\gamma}}^{\underline{\alpha}} \delta_{\underline{\delta}}^{\underline{\beta}} + \delta_{\underline{\gamma}}^{\underline{\beta}} \delta_{\underline{\delta}}^{\underline{\alpha}}) & 0 \\ -\frac{1}{4}(\delta_{\underline{\gamma}}^{\underline{\alpha}} \Theta^{\underline{\beta}} + \delta_{\underline{\gamma}}^{\underline{\beta}} \Theta^{\underline{\alpha}}) & \delta_{\underline{\gamma}}^{\underline{\alpha}} \end{pmatrix} . \tag{4.23}$$

Defining $e^{-\phi' \cdot R'}(K_{\underline{A}}) e^{\phi' \cdot R'} \equiv \Phi_{\underline{A}}{}^{\underline{B}} K_{\underline{B}}$ we find that

$$F^{\underline{A}} = E^{\underline{B}} \Phi_{\underline{B}}{}^{\underline{A}} . \tag{4.24}$$

The dynamics of the non-linear theory should be given by imposing the constraints of equations (4.12) to (4.14) where the derivatives are now defined for the full theory. In the next section we will argue that these constraints correctly encode the non-linear dynamics of the brane. It would be interesting to examine the implications of these constraints in detail and to see in particular what they imply for all the Goldstone fields associated with the automorphism. For those superbranes that have a straightforward action, it is possible that the supervielbein on the superbrane that arises in the non-linear realization will play an essential role in defining the action as it does for the case of the bosonic brane considered in section 2.

5. Relation to the superembedding formalism

In the superembedding approach [7] we consider one supermanifold M embedded in another $\underline{M}$. The preferred vector fields on the former are given by $E_A = E_A{}^M \partial_M$, where $E_A{}^M$ are the inverse supervielbein, while those on the latter are $E_{\underline{A}} = E_{\underline{A}}{}^{\underline{M}} \partial_{\underline{M}}$. The fields on M must point somewhere in $\underline{M}$ as is expressed by the equation

$$E_A{}^M \partial_M = E_A{}^{\underline{A}} E_{\underline{A}}{}^{\underline{M}} \partial_{\underline{M}} . \tag{5.1}$$

The assumption in the superembedding approach is that

$$E_\alpha{}^{\underline{a}} = 0 . \tag{5.2}$$

This states that the Grassmann odd preferred vector fields in M can be expressed in terms of only the Grassmann odd vector fields in $\underline{M}$. Although the elegant simplicity of this assumption is apparent, its deeper meaning is still unclear. For many branes, and certainly many of the most interesting ones, this single assumption is sufficient to put the theory on-shell and determine its complete equations of motion [17]. This is the case for the twobrane and fivebrane of M-theory.

We may re-express equation (5.1) as

$$dz^M E_M{}^A E_A{}^{\underline{B}} = dz^M \partial_M Z^{\underline{\Lambda}} E_{\underline{\Lambda}}{}^{\underline{B}} . \tag{5.3}$$

Multiplying by $\Phi_{\underline{B}}{}^{\underline{C}}$, identifying the supervielbien on the embedded super manifold with the Cartan form in the K_A directions, i.e. $E_M{}^A = F_M{}^A$, and recognizing $F^{\underline{A}}$ from equation (4.22) we conclude that

$$F_A{}^{\underline{B}} \equiv E_A{}^M F_M{}^{\underline{B}} = \nabla_A Z^{\underline{B}} = E_A{}^{\underline{C}} \Phi_{\underline{C}}{}^{\underline{B}} . \tag{5.4}$$

The Φ defined above equation (4.24) are constructed from the Goldstone fields associated with the automorphism group. They form an invertible matrix and also must have indices such that they are always Grassmann even; implying for example that $\Phi_{\underline{\beta}}{}^{\underline{b}} = 0$. It follows from equation (5.4) that the inverse Higgs condition

$$F_\alpha^{\underline{a}} = 0 , \tag{5.5}$$

is the necessary and sufficient condition for the embedding condition of equation (5.2). However, the superembedding formalism uses superfields that do not depend on the central charges other than the usual coordinates of Minkowski space. provided we take the central charge dependence of the superfields to be trivial. Hence, provided we adopt the inverse Higgs conditions of equations (4.12), (4.13) and (4.15) we should find that they describe the dynamics of the full non-linear theory. this follows because equation (4.13) is just the condition of equation (4..5), equation (4.14) eliminates the dependence of the superfields on the coordinates associated with the central charges and equation (4.12) identifies the automorphism Goldstone fields in terms of the fields of the brane. We note that the first two of the inverse Higgs conditions can be written as $F_\alpha{}^{\underline{A}} = \delta_{\underline{\alpha}}^{\underline{A}}$.

The relationship between the inverse Higgs mechanism and the superembedding condition has been previously discussed in the context of superbranes in lower dimensions and their construction as non-linear realizations [14] [20]. However, in these cases the dynamics of the branes are not determined by the embedding condition of equation (5.5). It would be interesting to apply the methods developed in these papers to branes of this type.

Given that we have identify the supervielbien F_A^M in the non-linear realization approach with the supervielbein of the superembedding approach implies that $F_A{}^B = \delta_A{}^B$ and as a consequence $E_A{}^{\underline{B}}\Phi_{\underline{B}}{}^B = \delta_A{}^B$. Taking fermonic indices and including the above condition of equation (5.5) we conclude that $E_\alpha{}^{\underline{\beta}}\Phi_{\underline{\beta}}{}^\gamma = \delta_\alpha{}^\gamma$. We can however choose a set of frames normal to the superbrane. By expressing these vector fields in terms of the preferred frames of the background superspace, we can define an extension of the $E_A{}^{\underline{B}}$ to $E_{\underline{A}}{}^{\underline{B}}$. Making this extension in an appropriate way we conclude that $E_{\underline{A}}{}^{\underline{B}}$ is the inverse of the $\Phi_{\underline{A}}{}^{\underline{B}}$. Thus the Goldstone fields associated with the automorphism group specify the relation between preferred frames associated with the branes and those in the background superspace.

Although the superfields do not in the end depend on the central charges, without the central charges in the original algebra one would not have the large automorphism group which is essential for the theory to contain the fields of the brane. This is apparent as a consequence of the discussion of the above paragraph, however, another way to observe this fact is to note that the supervierbeins on the embedded manifold would be essentially trivial if we had not included the Goldstone fields for the automorphism group.

6. Conclusion

We have shown that the branes of M-theory and, by reduction, those in ten dimensions can be described by the theory of non-linear realization provided we use a superalgebra that has the large automorphism group considered in [15]. Indeed, the field strengths of the worldvolume fields of the brane arise as the Goldstone bosons

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for the automorphisms. It is a widely held belief that branes correspond to the low energy effects of a defect in space-time the breaks certain symmetries. We have shown in this paper that their dynamics is essentially determined by a knowledge of which symmetries are broken.

In reference [19] it was explained how the fields of the branes arise as zero modes of the the corresponding background supergravity theory. It would be interesting understand the connections between that work and that of this paper. It would also be interesting to find if there are connections between the non-linear realization proposed in this paper and the coset consider in reference [14] which was based on the OSP groups.

The necessity of including at least parts of the GL(32) automorphism group in the construction of the non-linear realization of the branes of M-theory supports the idea proposed in [15] that this symmetry may play an important role in M-theory.

Acknowledgments

The author would like to thank Jon Bagger for a number of useful discussions at an early stage of this work and the hospitality of the Ecole Normale Superieure, where some of this work was carried out. This work was supported in part by the EU network on Integrability, Non-perturbative effects, and Symmetry in Quantum Field theory (FMRX-CT96-0012).

References

[1] W.Siegel, *Hidden local supersymmetry in the supersymmetric particle action*, *Phys. Lett.* **B 128** (1983) 397.

[2] E. Bergshoeff, E. Sezgin and P.K. Townsend, *Supermembranes and eleven-dimensional supergravity*, *Phys. Lett.* **B 189** (1987) 75;
Properties of the eleven-dimensional super membrane theory. Ann. Phys. (NY) **185** (1988) 330.

[3] M. Cederwall, A. von Gussich, B.E.W. Nilsson and A. Westerberg, *The dirichlet super-three-brane in ten-dimensional type-IIB supergravity*, *Nucl. Phys.* **B 490** (1997) 163 [hep-th/9610148];
M. Cederwall, A. von Gussich, B.E.W. Nilsson, P. Sundell and A. Westerberg, *The dirichlet super-p-branes in ten-dimensional type-IIA and -IIB supergravity*, *Nucl. Phys.* **B 490** (1997) 179 [hep-th/9611159];
M. Aganagic, C. Popescu and J.H. Schwarz, *D-brane actions with local κ-symmetry*, *Phys. Lett.* **B 393** (1997) 311 [hep-th/9610249];
E. Bergshoeff and P.K. Townsend, *Super D-branes*, *Nucl. Phys.* **B 490** (1997) 145 [hep-th/9611173];
M. Aganagic, J. Park, C. Popescu and J.H. Schwarz, *World-volume action of the M-theory five-brane*, *Nucl. Phys.* **B 496** (1997) 191 [hep-th/9701166].

[4] P.S. Howe, E. Sezgin and P.C. West, *Covariant field equations of the M-theory five-brane*, *Phys. Lett.* **B 399** (1997) 49 [hep-th/9702008].

[5] I. Bandos, K. Lechner, A. Nurmagambetov, P. Pasti, D. Sorokin and M. Tonin, *Covariant action for the super-five-brane of M-theory*, *Phys. Rev. Lett.* **78** (1997) 4332 [hep-th/9701149]; *On the equivalence of different formulations of the M-theory five-brane*, *Phys. Lett.* **B 408** (1997) 135.

[6] M. Perry and J.H. Schwarz, *Interacting chiral gauge fields in six-dimensions and Born-Infeld theory*, *Nucl. Phys.* **B 489** (1997) 47 [hep-th/9611065];
J.H. Schwarz, *Coupling a selfdual tensor to gravity in six-dimensions*, *Phys. Lett.* **B 395** (1997) 191 [hep-th/9701008];
M. Aganagic, J. Park, C. Popescu and J.H. Schwarz, *World-volume action of the M-theory five-brane*, *Nucl. Phys.* **B 496** (1997) 191 [hep-th/9701166].

[7] For two reviews see
P.S. Howe, E. Sezgin and P.C. West, *Aspects of superembeddings* in *D.V. Volkov Memorial volume*, Lecture notes in physics, vol. 509, Springer-Verlag, Berlin 1998:
D. Sorokin, *Superbranes and superembeddings*, hep-th/9906142.

[8] J. Hughes and J. Polchinski, *Partially broken global supersymmetry and the superstring*, *Nucl. Phys.* **B 278** (1986) 147;
J. Hughes, J. Liu and J. Polchinski, *Supermembranes*, *Phys. Lett.* **B 180** (1986) 370.

[9] S. Coleman, J. Wess and B. Zumino, *Structure of phenomenological lagrangeans, 1*, *Phys. Rev.* **177** (1969) 2239;
K. Callan, S. Coleman, J. Wess and B. Zumino, *Structure of phenomenological lagrangeans, 2*, *Phys. Rev.* **177** (1969) 2247.

[10] D.V. Volkov, *Sov. J. Nucl. Phys.* **4** (1973) 3;
D.V. Volkov and V.P. Akulov, *Possible universal neutrino interaction*, *Sov. Phys. JETP Lett.* **16** (1972) 438;
Is the neutrino a Goldstone particle?, *Phys. Lett.* **B 46** (1973) 109.

[11] J. Bagger and A. Galperin, *Matter couplings in partially broken extended supersymmetry*, *Phys. Lett.* **B 336** (1994) 25 [hep-th/9406217]; *A new goldstone multiplet for partially broken supersymmetry*, *Phys. Rev.* **D 55** (1997) 1091 [hep-th/9608177]; *The tensor goldstone multiplet for partially broken supersymmetry*, *Phys. Lett.* **B 412** (1997) 296 [hep-th/9707061].

[12] F. Gonzalez-Rey, I.Y. Park and M. Roček, *On dual 3-brane actions with partially broken $N = 2$ supersymmetry*, *Nucl. Phys.* **B 544** (1999) 243 [hep-th/9811130];
M. Roček and A.A. Tseytlin, *Partial breaking of global $D = 4$ supersymmetry, constrained superfields and 3-brane actions*, *Phys. Rev.* **D 59** (1999) 106001 [hep-th/9811232];
S. Bellucci, E. Ivanov and S. Krivonos, *Partial breaking of $N = 1$ $D = 10$ supersymmetry*, *Phys. Lett.* **B 460** (1999) 348 [hep-th/9811244];
E. Ivanov and S. Krivonos, *$N = 1$ $D = 4$ supermembrane in the coset approach*, *Phys. Lett.* **B 453** (1999) 237 [hep-th/9901003].

[13] S. Bellucci, E. Ivanov and S. Krivonos, *Partial breaking $N = 4$ to $N = 2$: hypermultiplet as a Goldstone superfield*, *Fortsch. Phys.* **48** (2000) 19, [hep-th/9809190];

S. Belucci, E. Ivanov and S. Krivonos, *Partial breaking of $N=1$, $D=10$ supersymmetry*, *Phys. Lett.* **B 460** (1999) 348 [hep-th/9811244;
F. Gonzalez-Rey, I.Y. Park and M. Rocek, *On dual 3-brane actions with partially broken $N=2$ supersymmetry*, *Nucl. Phys.* **B 544** (1999) 243 [hep-th/9811130];
M. Roček and A.A. Tseytlin, *Partial breaking of global $D=4$ supersymmetry, constrained superfields and 3-brane actions*, *Phys. Rev.* **D 59** (1999) 106001 [hep-th/9811232];
E. Ivanov and S. Krivonos, *Phys. Lett.* **B 453** (1999) 237;
S.V. Ketov, *Mod. Phys. Lett.* **A 14** (1999) 501; *Born-Infeld-Goldstone superfield actions for gauge-fixed D-5 and D-3 branes in 6D*, *Nucl. Phys.* **B 553** (1999) 250 [hep-th/9812051].

[14] T. Adawi, M. Cederwall, U. Gran, M. Holm and B.E.W. Nilsson, *Superembeddings, nonlinear supersymmetry and five-branes*, *Int. J. Mod. Phys.* **A 13** (1998) 4691 [hep-th/9711203].

[15] O. Barwald and P.C. West, *Brane rotating symmetries and the fivebrane equations of motion*, hep-th/9912226.

[16] E.A. Ivanov and V.I. Ogievetsky, *The inverse Higgs phenomenon in nonlinear realizations*, Teor. Mat. Fiz. **25** (1975) 164.

[17] P.S. Howe and E. Sezgin, *Superbranes*, *Phys. Lett.* **B 390** (1997) 133 [hep-th/9607227]; $D=11$, $p=5$, *Phys. Lett.* **B 394** (1997) 62 [hep-th/9611008].

[18] D. Sorokin and P.K. Townsend, *M-theory superalgebra from the M5-brane*, *Phys. Lett.* **B 412** (1997) 265 [hep-th/9708003].

[19] T. Adawi, M. Cederwall, U. Gran, B.E.W. Nilsson and B. Razaznejad, *Goldstone tensor modes*, *J. High Energy Phys.* **02** (1999) 001.

[20] E. Ivanov and S. Krivonos, *$N=1$ $D=4$ supermembrane in the coset approach*, *Phys. Lett.* **B 453** (1999) 237 [hep-th/9901003].

[21] F. Delduc, E. Ivanov and S. Krivonos, 1/4 *PBGS and superparticle actions*, hep-th/9912292; 1/4 *partial breaking of global supersymmetry and new superparticle actions*, hep-th/9912222.

[22] I. Bandos, J. Lukierski and D. Sorokin, *The* OSp(1|4) *superparticle and exotic BPS states*, hep-th/9912264;
P. Pasti, D. Sorokin and M. Tonin, *Branes in super-ADS backgrounds and superconformal theories*, hep-th/9912076;
I. Bandos, J. Lukierski and D. Sorokin, *Generalized superconformal symmetries and supertwistor dynamics*, hep-th/9912051.

[23] J.A. de Azcarraga, J.P. Gauntlett, J.M. Izquierdo and P.K. Townsend Topological extensions of the superalgebra for extended objects *Phys. Rev. Lett.* **63** (1989) 2443.

[24] J.P. Gauntlett, K. Itoh and P.K. Townsend, *Superparticle with extrinsic curvature*, *Phys. Lett.* **B 238** (1990) 65;
J.P. Gauntlett, J. Gomis and P.K. Townsend, *Particle actions as Wess-Zumino terms for space-time (super)symmetry groups*, *Phys. Lett.* **B 249** (1990) 255.

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