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BPS preons and tensionless super-p-branes in generalized superspace

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Abstract

Tensionless super–p–branes in a generalized superspace with additional tensorial central charge coordinates may provide an extended object model for BPS preons, *i.e.* for the hypothetical constituents of M–theory preserving 31 of 32 supersymmetries [18].

1 Introduction

Recently, a new wave of interest for higher spin theories and their supersymmetric extensions can be witnessed [1, 2, 3, 4, 5, 6]. Moreover, the study of [7, 2, 4] exhibits the relation among massless high–spin theories and simple particle–like dynamical models [8, 9] living in generalized superspace $\Sigma^{(\frac{n(n+1)}{2}|n)}$ with local coordinates

$$Z^{\mathcal{M}} = (X^{\alpha\beta}\, ,\, \theta^{\alpha})\, , \quad X^{\alpha\beta} = X^{\beta\alpha}\, , \quad \alpha, \beta = 1, \ldots, n\, , \tag{1}$$

[10, 11]. This relation suggested a way to introduce a concept of causality in 'symplectic spacetime' $\Sigma^{(\frac{n(n+1)}{2}|0)}$ [4] (*i.e.* in the bosonic body of $\Sigma^{(\frac{n(n+1)}{2}|n)}$) parametrized by symmetric $GL(n)$–tensor coordinates $X^{\alpha\beta} = X^{\beta\alpha}$ [12].

For $n = 2^k$, where α can be treated also as a spinor index of a D–dimensional Lorentz group $SO(t, D-t)$ for some D and t, $X^{\alpha\beta} = X^{\beta\alpha}$ can be regarded as symmetric spin–tensor coordinates. For $k > 1$ the set of such bosonic coordinates includes, besides the usual D–dimensional spacetime coordinates $x^{\mu} = X^{\alpha\beta}\Gamma^{\mu}_{\alpha\beta}$, a set of antisymmetric tensorial coordinates $y^{\mu_1\cdots\mu_q} = X^{\alpha\beta}\Gamma^{\mu_1\cdots\mu_q}_{(\alpha\beta)}$ ($y^{\mu\nu}$, $y^{\mu_1\cdots\mu_5}$ for $D = 11$ generalized superspace $\Sigma^{(528|32)}$). Just the introduction of gamma–matrices or, equivalently, the distinction between vector and antisymmetric tensor coordinates breaks the manifest $GL(n)$ symmetry of the generalized superspace down to $Spin(t, D-t)$. (Note that $n = 32$ case allows also $SO(2, 10)$ interpretation, in which $X^{\alpha\beta}$ contains antisymmetric tensor coordinates only [13, 14, 15]). Such a breaking of high spin $GL(n)$ symmetry (actually of the $OSp(2n|1)$ symmetry, see [1, 2, 4, 8, 7] and Sec. 4 below) is expected to be spontaneous.

An important property of the models [8, 9], not yet reflected in higher spin theories, is that they describe BPS states preserving all but one spacetime supersymmetries. This property is closely related to the fact that these models produce the generalized Penrose relation

$$P_{\alpha\beta} = \lambda_{\alpha}\lambda_{\beta} \tag{2}$$

($cf.$ [16]) as a constraint for the momentum $P_{\alpha\beta}(\tau)$ canonically conjugate to the coordinate function $X^{\alpha\beta}(\tau)$,

$$P_{\alpha\beta}(\tau) - \lambda_\alpha(\tau)\lambda_\beta(\tau) = 0\,. \tag{3}$$

Here τ is a proper time parametrizing a worldline W^1 in generalized superspace,

$$W^1 \in \Sigma^{(\frac{n(n+1)}{2}|n)}: \quad X^{\alpha\beta} = X^{\alpha\beta}(\tau)\,, \quad \theta^\alpha = \theta^\alpha(\tau)\,, \tag{4}$$

$X^{\alpha\beta}(\tau)$ and $\theta^\alpha(\tau)$ are bosonic and fermionic coordinate functions.

The most general supersymmetry algebra (called M–algebra in $D = 11$ case, *i.e.* for $n = 32$, [17])

$$\{Q_\alpha, Q_\beta\} = P_{\alpha\beta} \quad , \quad [Q_\alpha, P_{\alpha\beta}] = 0\,, \tag{5}$$

is realized in the model of [8] on the Poisson brackets (for simplicity, we ignore the i factor appearing in the Poisson brackets). After quantization, *schematically* (see [8, 7] for a precise Hamiltonian analysis and quantization), Eq. (2) could be considered as a condition on the state vector $|\lambda>$ of the quantum dynamical system,

$$P_{\alpha\beta}|\lambda> = \lambda_\alpha\lambda_\beta|\lambda> \quad . \tag{6}$$

Such a state was called *BPS preon* in [18] for reasons that will become clear below. Eq. (6) implies

$$\{Q_\alpha, Q_\beta\}|\lambda> = \lambda_\alpha\lambda_\beta|\lambda> \quad . \tag{7}$$

Then, introducing an auxiliary set of $(n-1)$ contravariant $GL(n)$ vectors w_I^α ($SO(t, D-t)$ spinors) orthogonal to the covariant $GL(n)$ vector λ_α,

$$w_I^\alpha \lambda_\alpha = 0\,, \qquad I = 1, \ldots, (n-1)\,, \tag{8}$$

one finds $w_I^\alpha\{Q_\alpha, Q_\beta\}|\lambda> = 0$. As a result, one can conclude that the BPS preon state $|\lambda>$ preserves all but one supersymmetries [18],

$$Q_I|\lambda> \equiv w_I^\alpha Q_\alpha|\lambda> \; = 0\,, \quad I = 1, \ldots, (n-1)\,. \tag{9}$$

Let us stress that the set of $(n-1)$ vectors w_I^α is pure auxiliary and has been introduced for convenience only. The preservation of $(n-1)$ of the n supersymmetries by the state $|\lambda>$ is encoded in the fact that the eigenvalue matrix $\lambda_\alpha\lambda_\beta$ of the operator $\{Q_\alpha, Q_\beta\}$, Eq. (7), has rank one.

Note that the causal structure of the symplectic spacetime $\Sigma^{(\frac{n(n-1)}{2}|0)}$ found in [4] is related to the observation that the state $|\lambda>$ obeying Eq. (6) provides the *general solution* of the conformal high–spin wave equation [2]

$$(P_{\alpha\beta}P_{\gamma\delta} - P_{\alpha\gamma}P_{\beta\delta})|\lambda> = 0\,. \tag{10}$$

The algebra similar to (5) is satisfied by the fermionic constraints $D_\alpha(\tau)$ ($\{D_\alpha, Q_\beta\} = 0$),

$$\{D_\alpha, D_\beta\} = -P_{\alpha\beta} \quad , \quad [D_\alpha, P_{\alpha\beta}] = 0\,. \tag{11}$$

Eqs. (11) and (3) imply that $(n-1)$ of n fermionic constraints, $D_I = w_I^\alpha D_\alpha$, are of the first class. These first class constraints generate $(n-1)$ local fermionic κ-symmetries through the Poisson brackets. Thus the number of κ–symmetries of the worldline actions [8] coincides with the number of supersymmetries preserved by a BPS preon state (see [18]) as one might expect (see, *e.g.* [19], and [20] for extended discussion). Thus, one can consider *the presence of* $(n-1)$ κ*–symmetries as the main characteristic property of a BPS preon model in a superspace with* n *fermionic coordinates.*

The states $|\lambda>$ were used in [18] to provide a complete algebraic classification of the BPS states in M–theory (hence the name of BPS preons). This suggests to conjecture that any BPS state $|\Psi_k>$ is a superposition of a definite number k of the BPS preons. The number k is determined by the rank of its generalized momentum matrix $p_{\alpha\beta}$,

$$P_{\alpha\beta}|\Psi_k>= p_{\alpha\beta}|\Psi_k> , \quad rank(p_{\alpha\beta}) = k \, . \tag{12}$$

Then (see [18]) there exists a set of k $GL(n)$ vectors λ^a_β (spinors of the Lorentz $SO(t, D-t) \subset GL(n)$) such that

$$p_{\alpha\beta} = \sum_{a=1}^{k} \lambda^a_\alpha \lambda^a_\beta \, . \tag{13}$$

Eq. (13) allows to speculate that the BPS state $|\Psi_k>$, satisfying Eq. (12), can be considered as composite of k preon states $|\lambda^a>$, $a = 1, \ldots, k$ [18].

The existence of BPS preons and other BPS states preserving more than $1/2$ of the supersymmetry (i.e. composites of $k < n/2$ preons) is allowed from an algebraic point of view [8, 21, 22]. However, for a long time realizations of such states as solitonic solutions of the 'usual' $D \leq 11$ supergravity equations were not known and, in fact, the first search in simple models gave negative results [22]. However, such solutions (now with up to 28 of 32 supersymmetries preserved) have recently been found [23, 24, 25] as a particular case of pp–waves [26]. Thus the original expectation that BPS preons and the states composed from less than $n/2$ BPS preons cannot be realized in the 'usual' superspace (a 'BPS preon conspiracy') is broken, at least partially. The relation of such solutions with models in generalized superspace has not been clarified yet. One may assume that the (constant) 'values' of antisymmetric tensor fields, characteristic of the pp–wave background, should play there the role of some tensorial coordinates of generalized superspace (*cf.* [11]), but the details of the embedding of pp–wave spacetimes into a generalized superspace require additional study.

Here we address another problem. Only point–like models with the properties of BPS preons [8] (and composites of less than $n/2$ preons [9]) were known in the generalized superspace. On the other hand, if one takes seriously the hypothesis [18] that all the M-theory BPS–states ($M2$–brane, $M5$–brane, intersecting brane configurations, *etc.*) are composed from ($n=32$) BPS preons, one should find for the latter an extended object model (*i.e.* a model with p–dimensional worldvolume W^{p+1} rather than worldline W^1 (4)), at least in the generalized superspace. The main message of this letter is that such a model for $D=11$ BPS preons is provided by a 'twistor–like' formulation of tensionless p–branes in the generalized superspace $\Sigma^{(528|32)}$. Moreover, the model can be formulated in an arbitrary generalized superspace $\Sigma^{(\frac{n(n+1)}{2}|n)}$.

Tensionless p–branes in $D=4$ ($n=4$) generalized superspace $\Sigma^{(10|4)}$ were previously studied in [27, 28]. In [28] it was found that a twistor–like formulation of the tensionless p–brane in $\Sigma^{(10|4)}$ (which generalizes the model from [29] for the case of additional tensorial coordinates) possess 3 κ–symmetries. We will show here that for any n (or any D), including $n=32$ ($D=11$), the $\Sigma^{(\frac{n(n+1)}{2}|n)}$ generalization of the tensionless p–brane action from [29] possesses $(n-1)$ κ–symmetries. In the light of the above mentioned correspondence, this implies that the $n=32$ ($D=11$) version of this action provides a dynamical model for a BPS state which preserves 31 of 32 supersymmetries, *i.e.* an extended object model for a BPS preon.

2 Tensionless p–brane action in $\Sigma^{(\frac{n(n+1)}{2}|n)}$

We consider the following action for an extended object (p–brane) moving in generalized superspace $\Sigma^{(\frac{n(n+1)}{2}|n)}$

$$S=\int d^{p+1}\xi\, L=\frac{1}{2}\int d^{p+1}\xi\, \rho^m \Pi_m^{\alpha\beta}\lambda_\alpha\lambda_\beta \tag{14}$$

(*cf.* [29] for the usual $D=4$ superspace and [30] for $p=0$). Here

$$\Pi^{\alpha\beta}\equiv d\xi^m \Pi_m^{\alpha\beta}=dX^{\alpha\beta}(\xi)-id\theta^{(\alpha}\,\theta^{\beta)}(\xi) \tag{15}$$

is the pull–back of the supersymmetric Volkov–Akulov one–form for $\Sigma^{(\frac{n(n+1)}{2}|n)}$ on the worldvolume

$$W^{p+1}\subset\Sigma^{(\frac{n(n+1)}{2}|n)}: \qquad X^{\alpha\beta}=X^{\alpha\beta}(\xi)\,, \qquad \theta^\alpha=\theta^\alpha(\xi) \tag{16}$$

parametrized by local coordinates ξ^m, $m=0,1,\ldots,p$; $\rho^m=\rho^m(\xi)$ is a Lagrange multiplier and $\lambda_\alpha=\lambda_\alpha(\xi)$ are auxiliary bosonic variables. The action (14) does not contain any dimensionful parameter, what allows us to call its associated dynamical system *tensionless super–p–brane in generalized superspace*.

The $n=4$ counterpart of the action (14), with λ treated as a Majorana representation of $D=4$ Lorentz harmonics [30], was studied in [28]. On the other hand, for $n=2^k=dim(Spin(1,D-1))$, substituting $\Gamma_\mu^{\alpha\beta}\Pi_m^\mu\equiv\Gamma_\mu^{\alpha\beta}(\partial_m x^\mu-i\partial_m\theta\Gamma^\mu\theta)$ for $\Pi_m^{\alpha\beta}$ in Eq. (14), one arrives at a D dimensional generalization of the null–super–p–brane action from [29]. Certainly, only for $D=3,4,6,10$ the momentum density $P_\mu(\xi)=\lambda\Gamma_\mu\lambda$ is light–like and the tensionless super–p–brane can be called null–super–p–brane.

The set of global symmetries of the action (14) includes $GL(n)$ transformations acting on the indices $\alpha,\beta=1,\ldots,n$. It is also invariant, by construction, under the global supersymmetry

$$\begin{aligned}\delta_\epsilon X^{\alpha\beta}(\xi)=i\epsilon^{(\alpha}\theta^{\beta)}(\xi)\,,\quad \delta_\epsilon\theta^\alpha(\xi)=\epsilon^\alpha\,,\\ \delta_\epsilon\lambda_\alpha(\xi)=0\,,\quad \delta_\epsilon\rho^m(\xi)=0\,,\end{aligned} \tag{17}$$

The generators Q_α of the supersymmetry (17) satisfy the algebra (5) involving the generator $P_{\alpha\beta}$ of the translations: $\delta_a X^{\alpha\beta}(\xi)=a^{\alpha\beta}$, $\delta_a\theta^\alpha(\xi)=0$, $\delta_a\lambda_\alpha(\xi)=0$, $\delta_a\rho^m(\xi)=0$.

A straightforward calculation of canonical momentum for $X^{\alpha\beta}(\tau)$, $P_{\alpha\beta} = \frac{\partial L}{\partial_0 X^{\alpha\beta}}$, results in the primary constraint

$$\Phi_{\alpha\beta} = P_{\alpha\beta}(\xi) - \rho^0(\xi)\lambda_\alpha(\xi)\lambda_\beta(\xi) = 0 \ , \tag{18}$$

(*cf.* Eq. (3)) which implies the propagation of the extended object in the directions characterized by $\lambda_\alpha(\xi)$. Such directions could be regarded as a $\Sigma^{(\frac{n(n+1)}{2}|n)}$ generalization of the light–like directions of the usual D–dimensional superspace.

The calculation of the other canonical momenta, $\mathcal{P}^\alpha(\xi) = \frac{\partial L}{\partial(\partial_0\lambda_\alpha)}$, $\mathcal{P}_m = \frac{\partial L}{\partial(\partial_0\rho_m)}$ and $\pi_\alpha(\xi) = \frac{\partial L}{\partial_0\theta^\alpha}$ also results in the constraints: $\mathcal{P}^\alpha(\xi) = 0$, $\mathcal{P}_m = 0$ and

$$D_\alpha = \pi_{\alpha\beta}(\xi) + iP_{\alpha\beta}\theta^\beta(\xi) = 0 \ . \tag{19}$$

The fermionic constraints (19) obey the algebra (11) on the Poisson brackets. This already indicates the presence of $(n-1)$ local fermionic κ–symmetries, which we now describe explicitly in the Lagrangian approach.

3 κ–symmetry and other gauge symmetries

It is convenient to write the general variation of the action (14) as

$$\begin{aligned} \delta S \ &= \int d^{p+1}\xi \, [\frac{1}{2}\delta\rho^m \Pi^{\alpha\beta}_m \lambda_\alpha\lambda_\beta + \rho^m \Pi^{\alpha\beta}_m \lambda_\beta \, \delta\lambda_\alpha] - \\ &- \frac{1}{2}\int d^{p+1}\xi \ \partial_m(\rho^m\lambda_\alpha\lambda_\beta) i_\delta \Pi^{\alpha\beta} - i\int d^{p+1}\xi \ \rho^m \partial_m \theta^\alpha \lambda_\alpha \delta\theta^\beta \lambda_\beta \ , \end{aligned} \tag{20}$$

where $i_\delta \Pi^{\alpha\beta} \equiv \delta X^{\alpha\beta} - i\delta\theta^{(\alpha}\,\theta^{\beta)}$, and integration by parts has been performed. Eq. (20) makes evident that the action (14) possesses $(n-1)$ κ–*symmetries*

$$\delta_\kappa \rho^m = 0 \ , \qquad \delta_\kappa \lambda_\alpha = 0 \ , \tag{21}$$

$$\delta_\kappa X^{\alpha\beta}(\xi) = i\delta_\kappa \theta^{(\alpha}\theta^{\beta)}(\xi) \ , \tag{22}$$

$$\delta_\kappa \theta^\alpha(\xi) \ = \kappa^I(\xi) w^\alpha_I(\xi) \ , \qquad I = 1, \ldots, (n-1) \ , \tag{23}$$

with parameters $\kappa^I(\xi)$. In Eq. (23) the $(n-1)$ auxiliary $GL(n)$ vector fields $w^\alpha_I(\xi)$ are defined as in Eq. (8), $w^\alpha_I(\xi)\lambda_\alpha(\xi) = 0$. In other words, the κ–symmetry transformation of the Grassmann coordinate function (23) is provided by the general solution of the equation

$$\delta_\kappa \theta^\alpha(\xi) \, \lambda_\alpha(\xi) = 0 \ . \tag{24}$$

Thus, we are not enforced to consider an extension of the phase space of our dynamical system by incorporation of auxiliary variables $w^\alpha_I(\xi)$ and their momentum: we can keep Eqs. (22), (21), (24) instead as the definition of the κ–symmetry (but we may use $w^\alpha_I(\xi)$ as a convenient tool to present the results in a transparent form).

The bosonic 'superpartner' of the fermionic κ–symmetry is provided by the b–symmetry transformations

$$\begin{gathered} \delta_b \theta^\alpha(\xi) \ = 0 \ , \qquad \delta_b \rho^m = 0 \ , \qquad \delta_b \lambda_\alpha = 0 \\ \delta_b X^{\alpha\beta}(\xi) = b^{IJ}(\xi) w^\alpha_I(\xi) w^\beta_J(\xi) \ , \end{gathered} \tag{25}$$

with $\frac{n(n-1)}{2}$ parameters $b^{IJ}(\xi) = b^{JI}(\xi)$, $I, J = 1, \ldots, (n-1)$. The only nontrivial part of the b–symmetry transformations, Eq. (25), is the general solution of the equation

$$\delta_b X^{\alpha\beta}(\xi)\, \lambda_\beta(\xi) = 0 \ . \tag{26}$$

Note also an evident scaling gauge symmetry of the action (14),

$$\begin{aligned} \delta_s \theta^\alpha(\xi) \ &= 0 \ , \qquad \delta_s X^{\alpha\beta}(\xi) = 0 \ , \\ \delta_b \rho^m = -2s(\xi)\rho^m \ , &\qquad \delta_s \lambda_\alpha = s(\xi)\lambda_\alpha(\xi) \ , \end{aligned} \tag{27}$$

as well as the symmetry under worldvolume general coordinate transformations (in their variational version $\tilde{\delta}_{g.c.}$ characterized by $\tilde{\delta}_{g.c.}\xi^m = 0$, see [20] and refs. therein)

$$\tilde{\delta}_{gc} X^{\alpha\beta}(\xi) = t^m(\xi)\partial_m X^{\alpha\beta} \ , \quad \tilde{\delta}_{gc}\theta^\alpha(\xi) = t^m(\xi)\partial_m \theta^\alpha \ , \quad \tilde{\delta}_{gc}\lambda_\alpha(\xi) = t^m(\xi)\partial_m \lambda_\alpha \ , \tag{28}$$

$$\tilde{\delta}_{gc}\rho^m(\xi) = \partial_n(\rho^m t^n) - \rho^n \partial_n t^m \ . \tag{29}$$

4 Supertwistor representation and $OSp(64|1)$ symmetry of the BPS preon model

Let us use the Leibniz rule ($\lambda\partial_m X \equiv \partial_m(\lambda X) - (\partial_m \lambda)X$, *etc.*, no integration by parts and no gauge fixing) to present the action (14) in the equivalent form

$$\begin{aligned} S = \ & \tfrac{1}{2}\int d^{p+1}\xi \, (\lambda_\alpha \rho^m \partial_m \mu^\alpha - \rho^m \partial_m \lambda_\alpha \, \mu^\alpha) - \\ & -\tfrac{i}{2}\int d^{p+1}\xi \, \rho^m \partial_m \eta \, \eta) \ , \end{aligned} \tag{30}$$

where

$$\mu^\alpha = X^{\alpha\beta}\lambda_\beta - \tfrac{i}{2}\theta^\alpha\theta^\beta\lambda_\beta \ , \quad \eta = \theta^\beta \lambda_\beta \ . \tag{31}$$

λ_α, μ^α and η can be regarded as components of an $OSp(2n|1)$ supertwistor $\mathcal{Y}^\Sigma$ [8],

$$\mathcal{Y}^\Sigma = (\, \lambda_\alpha \, , \quad \mu^\alpha \, , \quad \eta \,) \tag{32}$$

In terms of $\mathcal{Y}^\Sigma$ the action (30) reads

$$S = -\frac{1}{2}\int d^{p+1}\xi \, \rho^m \partial_m \mathcal{Y}^\Sigma \, C_{\Sigma\Lambda} \, \mathcal{Y}^\Lambda \ , \tag{33}$$

where

$$C_{\Sigma\Lambda} = \begin{pmatrix} 0 & \delta^\alpha{}_\beta & 0 \\ -\,\delta_\alpha{}^\beta & 0 & 0 \\ 0 & 0 & i \end{pmatrix} = -(-)^{\Sigma\Lambda} C_{\Lambda\Sigma} \tag{34}$$

is the orthosymplectic ($OSp(2n|1)$ invariant) 'metric' tensor.

The Lagrange multiplier ρ^m does not carry physical degrees of freedom. Indeed, using the general coordinate transformations $\tilde{\delta}_{gc}$, Eq. (29), and the scaling symmetry, Eq. (27), one can fix, *e.g.*, the gauge $\rho^m(\xi) = \delta^m_0$. The generalized Penrose correspondence (31) clearly does not restrict μ^α (as the first term in *r.h.s* contains the $\frac{n(n+1)}{2}$ parametric

$X^{\alpha\beta}$). Hence the tensionless super–p–brane model allows a description in terms of $2n$ bosonic and 1 fermionic components of the unconstrained *orthosymplectic supertwistor (32)* which *describes all the physical degrees of freedom of the system* and makes the global $OSp(2n|1)$ symmetry manifest. In particular, this implies that for $n = 32$ (i.e. $D = 11$) the extended BPS preon model (14) possesses an $OSp(64|1)$ generalized superconformal symmetry, which is characteristic both for high–spin theories (see [1, 2, 4]) and for the two–time physics approach to M-theory [31, 32].

5 Conclusion and outlook

We have shown that the dynamical system described by the action (14) possesses $(n-1)$ local fermionic κ–symmetries. Hence, in $n = 32$ ($D = 11$) such a dynamical system can be considered as an extended object model for BPS preons, the hypothetical constituents of M-theory [18]. We have seen as well that the BPS preon model possesses $OSp(64|1)$ symmetry, which was suggested to be a generalized conformal symmetry of M-theory (see [31, 32, 18] and refs. therein); this becomes transparent after passing to the equivalent supertwistor representation, Eq. (30) or (33), of the action (33). This simple transformation also exhibits the physical degrees of freedom of the dynamical system.

We call the object described by the action (14) a *tensionless super–p–brane in generalized superspace* $\Sigma^{(\frac{n(n+1)}{2}|n)} = \{(X^{\alpha\beta}, \theta^\alpha)\}$. The reasons are that the action (14) does not contain dimensionful parameters, and that the constraints (18) imply propagation in the generalized light–like directions of $\Sigma^{(\frac{n(n+1)}{2}|n)}$ (*cf.* Ref. [2]). Moreover, for $n = 2, 4, 8, 16$, one converts Eq. (14) into the action of null–super–p–brane in the usual $D = 3, 4, 6, 10$ superspaces (see [29] for D=4) by substituting $\Gamma_\mu^{\alpha\beta}\Pi_m^\mu \equiv \Gamma_\mu^{\alpha\beta}(\partial_m x^\mu - i\partial_m\theta\Gamma^\mu\theta)$ for $\Pi_m^{\alpha\beta}$.

Tensionless strings and p–branes in usual spacetime and usual superspace were discussed many times in the context of superstring/M–theory [33]–[36], [29], [37]–[43] (see [29, 41] for more references). In particular, they appear as singularities in K3 compactification of superstring theory down to six dimensions which connect all known supersymmetric six dimensional vacua [39]. An interesting perturbative approach to search for solutions of nonlinear superstring equations in the curved spacetime background was developed in [36]. It is based on a power series decomposition in the p–brane tension T_p and is close in spirit to earlier propositions [44] to obtain the quantum propagator of a p–brane by starting from the propagator of null–p–brane and summing up the perturbative series in T_p. The leading order of such expansion, null–string for $p = 1$, should dominate string amplitudes describing short distance string physics [45]. The tension generation mechanism, which allows one to obtain a tensionful superbrane action from a null–super–p–brane action was studied in [46, 47]. In this frame the (super)brane tension T_p appears as an integration constant in the solution of the superstring equations of motion. This allows for its different values in regions of a universe separated by a domain wall and unifies null–p–branes ($T_p = 0$) with tensionful p–branes ($T_p \neq 0$) . Furthermore, it was shown in [47] that the tension T_p can appear also as a result of a dimensional reduction. A development of this approach for the case of generalized superspaces might be useful for establishing mechanisms of tension generation and of the formation of the fundamental M–branes from our extended BPS preons.

Suggestions about a possible relation between tensionless strings and higher spin field theories can be found already in [48]. Recently this possible relation was discussed in the context of AdS/CFT correspondence [42]. The key observation is that, on one hand, both string field theory and the interacting higher spin theory contain infinite number of fields of higher spins, but, on the other hand, the latter has much more powerful gauge symmetry. This allowed Vasiliev to discuss in [48] the possibility that higher spin theories are more fundamental than string theory and that string theory can be viewed as a spontaneously broken phase of the higher spin theories. Then a possibility of an identification of the null strings (or null–p–branes) with higher spin theories was suggested by the enhancement of the symmetry in the tensionless limit of (super)string model (see [42] for further reasons).

The BPS preon conjecture [18] as a whole and, particularly, the tensionless superstring and super–p–brane models (14) *in generalized superspace*, can be considered also as a development of the above ideas. The models (14), being formulated in the generalized superspace which allows for a formulation of higher spin theories [12, 2, 4], respect, by construction, at least the $GL(n)$ part of the higher spin symmetry (see [27, 28] for other models in $D = 4$). The physically relevant M–branes and D–branes in usual $D = 11$ and $D = 10$ superspaces are expected to appear in a spontaneously broken phase of the BPS preon models, which should imply the breaking of $GL(n)$ symmetry down to some $Spin(1, D - 1) \subset GL(n)$ (e.g. $Spin(1, 10) \subset GL(32)$ for M-branes, $Spin(1, 9) \subset GL(32)$ for D–branes). Moreover, for $n = 2, 4, 8, 16$ the models (14) are directly related the $D = 3, 4, 6, 10$ massless higher spin theories: for $p > 0$ they describe an extended object generalization of the classical mechanics description of the free higher spin theories. Indeed, as it was shown in [7], the quantum state spectra of the $n = 2, 4, 8, 16$ generalized superparticle models [8], which are identical to the point–like ($p = 0$) models (14) (ρ^0 can be removed by rescaling of λ, Eq. (27)), consist of towers of massless fields of all possible 'spins' in $D = 3, 4, 6$ *and* 10. The special property of the $n = 2, 4, 8, 16$ ($n = 2(D-2)$, $D = 3, 4, 6, 10$) point–like models (14) are related with the existence of the Hopf fibrations $S^{2D-5}/S^{D-2} = S^{D-3}$ (see [7]). This provides a mechanism for 'momentum space compactification' of the additional (with respect to usual spacetime) degrees of freedom. The situation with $n = 32$ ($D = 11$) model is still unclear: classically it describes a particle with a dynamically generated mass [8, 9] and there is a problem in interpretation of the quantum state spectrum because no counterpart of the Hopf fibration is known for this case (see [9] for some discussion). However, in the framework of the BPS preon conjecture [18], which refers to superbrane rather than to field theories, the problem is rather a search for possibilities to construct a 'physical' BPS objects defined in the standard superspace, like M–branes and D–branes, from BPS preons in generalized superspace. In principle, one could explore the composite nature of the M–branes in terms of the point–like BPS preon model, but in an indirect way similar to the Matrix model description of supermembrane [49]. The tensionless p–brane models (14) provide a new basis to search for a possible composed nature of the M–branes: this search might be carried out by the quasi-classical methods for the extended object action, *e.g.* by studying solutions of the equations of motion (see [27] for some results in $D = 4$) and specific interactions with background fields in generalized superspace.

Recently an explicit relation between superparticle wavefunctions in generalized superspace and Vasiliev's 'unfolded equations' for higher spin field was established in [50] for $n = 4$ ($D = 4$). Moreover, the quantization of a counterpart of the $p = 0$, $n = 4$ model (14) defined in the generalized AdS_4 superspace has been also considered in [50].

This is of particular importance as the nontrivial interactions of higher spin fields can be constructed in a selfconsistent way only in a spacetime with nonvanishing cosmological constant (see [48] and refs. therein). It is interesting that the proper generalized AdS_4 superspace was proved to be just the supergroup manifold $OSp(1|4)$ [50, 51]. These results provide a reason to study also the AdS generalizations of the $n = 4$ ($D = 4$) versions of the model (14): the tensionless superstring and supermembrane on $OSp(1|4)$ supergroup manifold.

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