BPS preons and higher spin theory in $D = 4, 6, 10$\(^1\)

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Abstract
We briefly review here the notion of BPS preons, the hypothetical constituents of M-theory, emphasizing its generalization to arbitrary dimensions $D$ and its relation to higher spin theories in $D = 4, 6$ and 10.

1 BPS preons in M-theory and supergravity ($D=11$)

In $D=11$, BPS preons\(^1\) are M-theory BPS states preserving all but one of the 32 supersymmetries, $|BPS\; preon\rangle = [31/32\; BPS\rangle$. This implies that there exist 31 bosonic spinors $\epsilon_\alpha^I$, such that

$$\epsilon_\alpha^I Q_\alpha |BPS\; preon\rangle = 0, \quad \alpha = 1, \ldots, 32, \quad I = 1, \ldots, 31,$$

where the $Q_\alpha$'s are the 32 supersymmetry generators of M-theory. The $\epsilon_\alpha^I$ characterize the 31 preserved supersymmetries given by $\epsilon = \kappa^I \epsilon_\alpha^I$, where $\kappa^I$ are fermionic parameters, and correspond to the Killing spinors in the supergravity description.

Equivalently, a BPS preon may be characterized by one bosonic spinor $\lambda_\alpha$ such that

$$Q_\alpha |BPS\; preon\rangle \propto \lambda_\alpha \Rightarrow |BPS\; preon\rangle = [31/32\; BPS\rangle = |\lambda\rangle .$$

The preonic spinor $\lambda_\alpha$ is clearly orthogonal to the 31 Killing spinors in $|\lambda\rangle$, $\epsilon^I_\alpha \lambda_\alpha = 0, \quad \alpha = 1, \ldots, 32, \quad I = 1, \ldots, 31.$

The preonic nature of the 31/32 states comes from the fact\(^1\) that a $k/32$-supersymmetric BPS state can be considered as a composite of $\tilde{n} = 32 - k$ different BPS preons, schematically

$$|k/32\; BPS\rangle = |\lambda^{(1)}\rangle \otimes \ldots \otimes |\lambda^{(32-k)}\rangle \equiv \bigoplus_{l=1}^{\tilde{n}=32-k} |\lambda^{(l)}\rangle \cong |\text{preon}^{(1)}\rangle \otimes |\text{preon}^{(2)}\rangle \otimes \ldots \otimes |\text{preon}^{(32-k)}\rangle$$

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and it is characterized by \( \tilde{n} \) preonic spinors \( \lambda^{(l)} \) orthogonal to the \( k \) Killing spinors. The completely supersymmetric, \( 32/32 \) BPS states, which are usually identified with supersymmetric vacua, do not contain any preons. Adding a preon to a \( k/32 \) state one obtains a state breaking one more supersymmetry, \( (32 - k) + 1 \) in all. Thus preons can be thought of as the fundamental constituents of M-theory [1] (see [2, 3, 4, 5] for further discussion). Building a more composite BPS state from a given one corresponds to breaking of one or more of the originally preserved supersymmetries. In this picture, the fully non-supersymmetric states appear as the most complicated ones: they are composites of the maximal number, 32, of independent preons.

The \( k \)-supersymmetric M-theory BPS states, \( |k/32 \text{ BPS} > \), are usually associated with supersymmetric solutions of its low energy limit, which is identified with the \( D = 11 \) or \( D = 10 \) type II supergravities. The most important solutions were considered to be the 1/2 ones, corresponding to \( |16/32 \text{ BPS} > \) states, which contain the \( D = 10 \) Dirichlet \( p \)-branes and the \( D = 11 \) M-branes (see [6]). The less than 1/2 supersymmetric states were identified with the intersecting branes (see [7]). Thus, before the discovery (in 2002) of solutions preserving more than 1/2 supersymmetries \( (k/32 > 1/2; \text{see refs. in [8, 4, 9]} \) ), the main classification of the M-theory BPS states was based on their (intersecting) brane contents, and so it included the 1/2 BPS states and the intersecting branes when \( k/32 < 1/2 \). The 2001 preonic conjecture [1] provided an alternative algebraic classification of all possible BPS states in terms of their preons contents. It included all the then known BPS states and allowed for the existence of any \( k/32 \) states (see also [10] and [11]). In this sense it predicted the appearance of the supergravity solutions preserving more than 1/2 of the supersymmetries.

Recent studies seem to indicate there is, at least at the classical level, a preon ‘conspiracy’ precluding the existence of preonic solutions, both in IIB [12] and IIA [13] supergravities. According to [14], the simply-connected 31/32, preonic solutions of \( D = 11 \) supergravity are forbidden as well. This preon conspiracy in classical supergravity is nevertheless compatible with the BPS preon conjecture since preons were rather introduced as M-theoretic objects [1]. One might then establish a parallel between this preon conspiracy and quark confinement, and state that only composites of a certain number \( \tilde{n}_{\text{min}} \) of BPS preons can be ‘observed’ as supergravity solutions. An interesting point for future study is whether including quantum (‘stringy’ or ‘M-theoretic’) corrections in the supergravity equations would change the situation and allow for the existence of preonic solutions (see [13] for further discussion).

Another question concerns the level of this preon conspiracy in classical supergravity i.e., the minimal (but nonzero) allowed number \( \tilde{n} = 32 - k \) of preons which can form the BPS state described by a supergravity solution. By the above discussion, such a solution would preserve \( k = 32 - \tilde{n} \) out of 32 supersymmetries. Solutions preserving up to \( k = 28 \) supersymmetries are known in IIB and up to 26 in IIA and \( D = 11 \) supergravities (see refs. in [8, 4, 13]). Thus, the problem is whether \( \tilde{n} = 2 \) and 3 solutions, which preserve 30/32 and 29/32 supersymmetries respectively, do exist. Recent work [15] excludes the existence of solutions describing two preon states in \( D = 5 \) and in \( D = 4, N = 2 \) supergravity. The existence of \( \tilde{n} = 2 \) solutions in \( D = 10, 11 \) is still open.
2 D=4,6 and 10 BPS preons and free conformal higher spin theories

The notion of BPS preon applies [1, 3] to an arbitrary number of spacetime dimensions $D$. In $D = 4, 6$ and $10$ a preon state would be a BPS state preserving 3/4, 7/8 and 15/16 of the supersymmetries, respectively. In these ‘stringy’ dimensions a BPS preon is related to the $D = 4, 6$ and $10$ free massless conformal higher spin theory.

To exhibit this relation explicitly, we first notice that an equivalent (to (2)) definition of the BPS preon is given by the following relation [1]

$$ P_{\alpha\beta} \langle \text{BPS preon} \rangle = \lambda_\alpha \lambda_\beta \langle \text{BPS preon} \rangle, \quad \alpha, \beta = 1, \ldots, n \quad (5) $$

which implies that the preon eigenvalues matrix of the generalized momentum $P_{\alpha\beta}$ is of rank one. The $P_{\alpha\beta}$ are the abelian bosonic generators of the general supersymmetry algebra

$$ \{Q_\alpha, Q_\beta\} = 2P_{\alpha\beta}, \quad [P_{\alpha\beta}, P_{\gamma\delta}] = 0, \quad \alpha = 1, \ldots, n \quad (6) $$

which for $n = 32$ gives the M-algebra corresponding to $D=11$ (or $D=10$ type II with the proper index interpretation). The generators of this superalgebra have a natural differential operator representation

$$ P_{\alpha\beta} = i \frac{\partial}{\partial X_{\alpha\beta}} = :i\partial_{\alpha\beta}:, \quad Q_\alpha = i \frac{\partial}{\partial \theta^\alpha} + \theta^\beta \partial_{\alpha\beta}, \quad (7) $$

on an enlarged tensorial superspace parametrized by $n$ fermionic and $n(n+1)/2$ bosonic coordinates (see e.g. [3]),

$$ \Sigma^{(\frac{n(n+1)}{2})|n} : \{(X^{\alpha\beta}, \theta^\alpha)\}, \quad X^{\alpha\beta} = X^{\beta\alpha}, \quad \alpha, \beta = 1, \ldots, n \quad (8) $$

In addition to the $D$ vectorial, spacetime coordinates $x^a = (1/n)X^{\alpha\beta}\Gamma_a^{\alpha\beta}$ contains a number of tensorial coordinates, $y^{[D/2]} \propto X^{\alpha\beta}\Gamma_{\alpha\beta}^{a_1 \ldots a_{[D/2]}}$, the types and number of which depend on $n$ and $D$. For standard (one-time) spacetimes (see [16] and refs. therein for two-time physics), $X^{\alpha\beta}$ contains

$$ \{X^{\alpha\beta}\} = \begin{cases} 
  x^a, & \text{for } n = 2, D = 3 \\
  (x^a, y^{ab}), & \text{for } n = 4, D = 4 \\
  (x^a, y^{abc}), I = 1, 2, 3 & \text{for } n = 8, D = 6 \\
  (x^a, y^{abcdef}) & \text{for } n = 16, D = 10 \quad (9a) \\
  (x^a, y^{ab}, y^{abcdef}) & \text{for } n = 32, D = 11 \quad (9b) 
\end{cases} $$

For $n=2$, $X^{\alpha\beta} = X^{\beta\alpha}$ just provides another presentation of the $D=3$ spacetime coordinates. For $n \geq 4$ further bosonic coordinates appear, as e.g. 6 in $D = 4$ ($n=4$), 126 in $D=10$ ($n=16$) and 517=528-11 for $D=11$ ($n=32$).

In $D=4$, the superspace $\Sigma^{(10)4}$ of (9b) was proposed in [17] as a basis for description of massless higher spin theories (see [18]). A dynamical realization of these ideas was found
to be given by a generalized superparticle model, the $\Sigma^{(a_{n+1})/2}$ \textit{tensorial} superparticle \cite{10}. Its action reads

$$
S := \int d\tau \lambda_\alpha(\tau)\lambda_\beta(\tau) \Pi^{\alpha\beta} := \int d\tau \lambda_\alpha(\tau)\lambda_\beta(\tau) \left( \partial_\tau X^{\alpha\beta} - i\partial_\tau \theta^{(\alpha} \theta^{\beta)} \right)
$$

This is generalization (to $n > 2$, $\alpha, \beta = 1, \ldots, n$) of the $D=3$ version of that of the Ferber-Schirafuji superparticle action \cite{20}. Indeed, the classical mechanics counterpart of the second definition of BPS preon, Eq. \cite{5},

$$
\mathcal{P}_{\alpha\beta} - \lambda_\alpha(\tau)\lambda_\beta(\tau) \approx 0,
$$

follows from \cite{10} as a primary constraint. The $\Sigma^{(a_{n+1})/2}$ superparticle model \cite{10} possesses \((n - 1)\) $\kappa$-symmetries \cite{21} and $n$ supersymmetries. This implies \cite{10} that its ground state preserves all but one of the tangent space supersymmetries and, thus, can be identified \cite{4, 3} with a BPS preon.

Upon a quantization (which converts second class constraints into first class ones) \cite{19}, the constraint \cite{11} is imposed on the wave function in the coordinate representation. Ignoring here fermionic coordinates for simplicity, this gives the \textit{preonic equation} \cite{3, 25}

$$(\partial_{\alpha\beta} + i\lambda_\alpha(\tau)\lambda_\beta(\tau)) G(X, \lambda) = 0
$$

which admits the ‘plain wave’ solution $G(X, \lambda) = \phi(\lambda) e^{-i\lambda_\alpha X^{\alpha\beta}}$. Clearly, the bosonic spinor $\lambda_\alpha$ carries the (generalized) momentum degrees of freedom so that the (generalized) coordinate representation for the wavefunction is given by the integral of $G(X, \lambda)$ on $\lambda$ for some measure of integration. The simplest one, $d^n\lambda$, gives a scalar function $b(X) = \int d^n\lambda G(X, \lambda) = \int d^n\lambda \phi(\lambda) e^{-i\lambda_\alpha X^{\alpha\beta}}$; choosing alternatively $d^n\lambda \lambda_\alpha$ one arrives at a spinor wavefunction suitable for describing fermions, $f_\alpha(X) = \int d^n\lambda \lambda_\alpha \tilde{G}(X, \lambda) = \int d^n\lambda \lambda_\alpha \phi(\lambda) e^{-i\lambda_\alpha X^{\alpha\beta}}$. These wavefunctions obey the equations

$$
\partial_{\alpha\beta}\partial_{\gamma\delta} b(X) - \partial_{\alpha\gamma} \partial_{\beta\delta} b(X) = 0,
$$

$$
\partial_{\alpha\beta} f_\gamma(X) - \partial_{\alpha\gamma} f_\beta(X) = 0,
$$

which were proposed by Vasiliev \cite{22} to describe $D=4$ massless higher spin theories. These were generalized to (enlarged) $AdS$ superspaces ($OSp(1|n)$ supermanifolds) \cite{23}, and were shown to describe a whole tower of bosonic and fermionic free massless conformal higher spin fields also in $D=6,10$ \cite{24}.

The field strength of the spacetime higher spin fields can be extracted, e.g., by decomposing the $b(X) = b(x, y)$ and $f_\alpha(x, y)$ in a power series on the tensorial coordinates, $y^{(D/2)}$ ($y^{[2]} = y^{mn}$, $y^{[3]} = y^{mnk}$ and $y^{[5]} = y^{mnklp}$ for $D = 4, 6, 10$, Eqs. \cite{24}–\cite{21}). Schematically (see \cite{24, 25} for the precise expressions),

$$
b(x, y) = \phi(x) + y^{[D/2]} F_{[D/2]}(x) + y^{[D/2]} R_{[D/2]}(x) + \sum_{s=3}^{\infty} y^{[D/2]} \cdots y^{[D/2]} R_{[D/2]} \cdots [D/2]_{s+1}(x),
$$

$$
f_\alpha(x, y) = \psi_\alpha(x) + y^{[D/2]} \Psi_\alpha [D/2](x) + y^{[D/2]} y^{[D/2]} \Psi_\alpha [D/2](x) + \cdots.
$$
As is well known, in D=4 all the free massless equations are conformally invariant. Consequently, all the massless field strengths are included in the \( n = 4 \) version of the decomposition (15) on \( y^{[2]} = y^{mn} \). In particular, \( F_{[2]} := F_{mn}(x) = - F_{nm}(x) \) is the field strength of Maxwell field, \( R_{m_1n_1m_2n_2}(x) = R_{m_2n_2m_1n_1}(x) \) is the linearized Riemann tensor, \( \Phi_{[2]}(x) := \Psi_{[pmn]}(x) \) is the Rarita-Schwinger field strength (spin 3/2), etc. Eqs. (13) and (14) fix both the algebraic properties of these and other higher spin field strengths (such as the Bianchi identities \( R_{[m_1n_1m_2n_2]} = 0 \) for the linearized Riemann tensor) and also define the linear differential equations for these field strengths (such as the Bianchi identities for the Maxwell field strength and the equations of motion) [22, 23, 24].

In contrast, in \( D=6 \) and 10 not all the massless fields are conformal and, consequently, not all massless fields but only the conformal ones enter the decomposition (15) for \( n=8 \) and 16. In \( D=10 \) these are, in addition to the usual scalar and spinor fields \( \phi(x) \) and \( \psi_{\alpha}(x) \), the basic self-dual five form field strength \( F_{[m_1m_2...m_5]} = \frac{1}{5!} \epsilon_{m_1m_2...m_5n_1n_2...n_5} F^{n_1n_2...n_5} \) (characteristic of type IIB supergravity) and the tensors with several symmetrized groups of ‘five’s’ i.e., with symmetrized sets of five antisymmetric self-dual indices, \( R_{[n_1][n_2]} = R_{[n_2][n_1]} \) etc., as well as their fermionic counterparts [24].

The bosonic \( b(X) \) and fermionic \( f_{\alpha}(X) \) fields are the two lowest components of a superfield on the \( \Sigma^{\frac{n(n+1)}{2}} \) superspace, \( \Phi(X, \theta) = b(X) + \theta^\alpha f_{\alpha}(X) + \mathcal{O}(\theta \theta) \). Then the free conformal higher spin equations (13) and (14) follow from the simple linear differential equation [25]

\[
D_{[\alpha} D_{\beta]} \Phi(X, \theta) = 0 , \quad D_\alpha := \frac{\partial}{\partial \theta^\alpha} + i \theta^\beta \partial_{\alpha\beta} .
\]  

A calculation shows that Eq. (16) also implies the vanishing of all higher components of \( \Phi(X, \theta) \). The group-theoretical meaning of this equation was discussed in [26], while its curved space (generalized AdS) generalization and supergravity in tensorial superspace was the subject of [25] to which we also refer for a discussion on the problems and perspectives for an interacting higher spin theory in this framework.

In the same way as the scalar bosonic and the spinor fermionic wavefunctions, \( b(X) \) and \( f_{\alpha}(X) \) in Eqs. (13), (14), are constructed from the solution of the preonic equation (12) [3], one can express the solution of the superfield equation (16) as an integral \( \Phi(X, \theta) = \int d^\nu \lambda \mathcal{G}_0(X, \theta, \lambda) \) of the \( \lambda \)-dependent (‘phase space’) superfield \( \mathcal{G}_0(X, \theta, \lambda) \) which obeys the following superfield generalization [25] of the preonic equation (12)

\[
(D_{\alpha} D_{\beta} - 2 \lambda_{\alpha} \lambda_{\beta}) \mathcal{G}_0(X, \theta, \lambda) = 0 .
\]  

The antisymmetric part of this equation gives rise to (16) while the symmetric part has the form of the preonic equation (12) \( \{ D_{\alpha} , D_{\beta} \} = 2 i \partial_{\alpha\beta} \). The phase space superfield \( \mathcal{G}_0 \), in its turn, appears as the leading component of the Clifford superfield \( \mathcal{G}(X, \theta, \lambda, \chi) = \mathcal{G}_0(X, \theta, \lambda) + \chi \mathcal{G}_1(X, \theta, \lambda) \), \( \chi \chi = 1 \) [27], which obeys the first order Clifford superspace equation [19, 25]

\[
(D_{\alpha} + i \chi \lambda_{\alpha}) \mathcal{G}(X, \theta, \lambda, \chi) = 0 , \quad \chi \chi = 1 .
\]
Eq. (18) implies \( D_\alpha G_0 + i \lambda_\alpha G_1 = 0 \) and \( D_\alpha G_1 + i \lambda_\alpha G_0 = 0 \). These mean, besides that both components \( G_0, G_1 \) of the Clifford superfield \( G \) obey the preonic equation (12) (representing in torsional spacetime the second definition (5) of a BPS preon), that \( Q_\alpha G_0 \propto \lambda_\alpha \) and \( Q_\alpha G_1 \propto \lambda_\alpha \) are valid. Any of these two equations provide a torsional superspace representation (see Eq.(7)) of the first definition (2) of a BPS preon.

3 Concluding remarks: fermionic preons?

The above discussion suggests the possibility of considering BPS preons with not only bosonic, but also with fermionic (and, perhaps, even with exotic) statistics; this requires further study. Here we only notice that the two equivalent definitions of a BPS preon, Eqs.(2) and (5), imply the existence of an ultrashort preonic supermultiplet containing one bosonic and one fermionic state, \( |\lambda, b\rangle \) and \( |\lambda, f\rangle \), characterized by the same bosonic spinor \( \lambda_\alpha \), such that

\[
Q_\alpha |\lambda, b\rangle = \lambda_\alpha |\lambda, f\rangle, \quad Q_\alpha |\lambda, f\rangle = \lambda_\alpha |\lambda, b\rangle.
\] (19)

For their associated fields \( \phi(\mathbb{X}) = < X|\lambda, b > \) and \( \psi(\mathbb{X}) = < X|\lambda, f > \) (where \( \mathbb{X} \) may be \( X^{\alpha\beta}, \lambda_\alpha \) or different), the supersymmetry transformations in (19) read \( \delta \phi(\mathbb{X}) = \varepsilon^\alpha \lambda_\alpha \psi(\mathbb{X}), \delta \psi(\mathbb{X}) = \varepsilon^\alpha \lambda_\alpha \phi(\mathbb{X}) \). As a ground state is taken to be bosonic, \( \psi(\mathbb{X}) = 0 \), such a state is clearly invariant under the 31 supersymmetries associated with the 31 Killing spinors \( \epsilon^\alpha \) of Eq. (3). This bosonic ground state configuration is identified with a BPS preon. However, one sees that the same 31 supersymmetries are preserved by the purely fermionic \( (\phi(\mathbb{X}) = 0) \) state characterized by the Grassmann odd function \( \psi(\mathbb{X}) \), the fermionic counterpart of the bosonic BPS preon. It would be interesting to understand whether such a simple algebraic construction of a fermionic BPS preon also has a dynamical realization.

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