Pure Spinors and $D=6$ Super-Yang–Mills

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Abstract: Pure spinor cohomology has been used to describe maximally supersymmetric theories, like $D = 10$ super-Yang–Mills and $D = 11$ supergravity. Supersymmetry closes on-shell in such theories, and the fields in the cohomology automatically satisfy the equations of motion. In this paper, we investigate the corresponding structure in a model with off-shell supersymmetry, $N = 1$ super-Yang–Mills theory in $D = 6$. Here, fields and antifields are obtained as cohomologies in different complexes with respect to the BRST operator $Q$. It turns out to be natural to enlarge the pure spinor space with additional bosonic variables, subject to some constraints generalising the pure spinor condition, in order to accommodate the different supermultiplets in the same generalised pure spinor wave-function. We construct another BRST operator, $s$, acting in the cohomology of $Q$, whose cohomology implies the equations of motion. We comment on the possible use of similar approaches in other models.
1. Introduction

Pure spinors and pure spinor cohomology (or spinorial cohomology) seems to be a deep structure underlying supersymmetric gauge theories, including supergravity. In a more pragmatic sense, pure spinors act as a book-keeping device for superspace forms with spinorial indices. The pure spinor constraint, which is generically of the form $(\lambda \gamma^a \lambda) = 0$, essentially projects out torsion from the anti-commutator of two fermionic derivatives, and ensures the nilpotency of the pure spinor BRST operator $Q = \lambda^a D_a$.

Finding superspace formulations of the component dynamics of supersymmetric Yang–Mills theory [1] has to a large extent been a question of trial and error. It was recognised early that pure spinors has a rôle to play [2,3], but it took much longer to realise how to make systematic use of them. Pure spinor techniques arise naturally from the formulation of gauge theory on superspace, and have been used successfully in $D = 10$ [4,5,6,7], both for the undeformed theory and its supersymmetric deformations. The theory may be supplemented with extra pure spinor variables that enable the construction of a Chern–Simons like off-shell action [8] constructed in terms of $Q$. These techniques have also been very powerful in the covariant quantisation of the superstring (see refs. [9,10] and references therein). Some results have been obtained for $D = 11$ supergravity [11], and for maximally supersymmetric models, where explicit higher-derivative corrections were derived in refs. [4,12,11,13].

Standard component formulations of maximally supersymmetric theories are on-shell supersymmetric—the equations of motion are implied by $Q = 0$—whereas the pure spinor framework provides an off-shell formulation. In this sense, the property of the supermultiplets of having supersymmetry transformations closing only on shell is turned into a virtue. The dynamics can be given the simple form of vanishing curvature (the equations of motion of a Chern–Simons-like action). The corresponding situation for less supersymmetric theories has not been as clear. Attempts have been made to use pure spinors in lower dimensions [14], but mainly in the context of superstrings. Neither has it been as thoroughly investigated, since the existence of off-shell supermultiplets provides means of treating the models without sacrificing supersymmetry. Nevertheless, pure spinor cohomology works as a generic method of defining the supermultiplets, including auxiliary fields, and if we think there is some deeper significance to such a statement, we should also seriously consider the mechanisms used to go from the off-shell multiplets to the on-shell theory. This question is addressed and solved in the present paper for a specific case of a theory with 8 supercharges, $N = 1$ super-Yang–Mills theory in $D = 6$. (The superspace formulation of this theory, as well of the one in $D = 10$ is of course well known [15,16,17,18].)

In the concluding section we also comment on how similar techniques may clarify the situation for supermultiplets that are partially off-shell, like $D = 10$, $N = 1$ supergravity and heterotic supergravity, and how they may be used to address questions in $D = 11$ supergravity and M-theory.
2. Review of pure spinors in $D = 10$ and super-Yang–Mills

Consider a scalar field, or first-quantised wave function, that in addition to the superspace coordinates $x^a, \theta^a$ (we take superspace to be flat, for simplicity) depend on a pure spinor $\lambda^a$. The pure spinor constraint is $T^a \equiv (\lambda \gamma^a \lambda) \approx 0$. The fields is seen as a formal series expansion in $\lambda$, and gauge invariance (with the pure spinor constraint as generator) is implemented by factoring out the ideal generated by $T$.

The flat superspace torsion is $T_{\alpha \beta} = 2 \gamma_{\alpha \beta}$. The pure spinor constraint implies that the fermionic operator $Q = \lambda^a D_\alpha$ is nilpotent, and can be used as a BRST operator. The content of a scalar field, i.e., the structure of the complex is

$$r_0 \xrightarrow{Q} r_1 \xrightarrow{Q} r_2 \xrightarrow{Q} \ldots \xrightarrow{Q} r_n \xrightarrow{Q} \ldots \quad (2.1)$$

where the representations $r_n$ at each $n$ denotes a superfield in the representation $r_n$ of the Lorentz group. This representation is the conjugate representation to the symmetric product of $n$ pure spinors. When we describe a gauge theory, $r_n$ consists of totally symmetric and $\gamma$-traceless tensors in $n$ spinor indices. The fermionic exterior derivative is a projection on the representations $r_n$ of a (symmetrised) spinorial covariant derivative. The general interpretation is that $r_0$ contains gauge transformations (or ghosts), $r_1$ contains fields, $r_2$ deformations (or antifields) and $r_3$ antighosts of the theory. For a theory of a rank-$(p+1)$ tensor potential, such as the three-form in $D = 11$ supergravity, $r_p$ contains gauge transformations (ghosts), $r_{<p}$ ghosts for ghosts, $r_{p+1}$ fields, $r_{p+2}$ deformations (antifields), $r_{p+3}$ antighosts, etc. Although the structure seems general, it really occurs only for a limited number of cases, and relies on maximal supersymmetry.

![Figure 1: Convention for Dynkin indices for so(1,9) representations](image)

For $N = 1$ and $D = 10$, with $\lambda^a$ in $(00001)$ and $D_\alpha$ in $(00010)$, we have the representations

$$(00000) \xrightarrow{Q} (00010) \xrightarrow{Q} (00020) \xrightarrow{Q} \ldots \xrightarrow{Q} (00n0) \xrightarrow{Q} \ldots \quad (2.2)$$

The cohomology at zero momentum (the “field content”) can be calculated with purely algebraic means, and is displayed in the table 1 below. Each column represents the superfield content of the representations in the complex (2.2), and has been shifted so that $Q$ acts horizontally.
Our conventions for dimensions is that a bosonic derivative has dimension 1, $[\partial_a] = 1$. Then $[D_\alpha] = \frac{1}{2}$, $[\lambda^\alpha] = -\frac{1}{2}$ and $[Q] = 0$.

The cohomology at non-zero momentum restricts the fields to obey the linearised equations of motion. An easy way to read off the possible equations of motion from the zero-momentum cohomology is to find the cohomology in the column to the right of the one containing the fields. The full non-linear super-Yang–Mills equations of motion are reproduced by the solution at order $\lambda$ of the zero curvature condition $Q\Psi + \Psi^2 = 0$. The equations may be derived from a Chern–Simons like Batalin–Vilkovisky action where the measure involves the position of the antighost (this requires additional variables in order to make the measure non-degenerate [8], and will not be dealt with in this paper).

3. Pure spinors and extended pure spinors in $D = 6$

We now turn our attention to $N = 1$ super-Yang–Mills theory in $D = 6$. The conventions for Dynkin labels of $Spin(1,5) \times SU(2)$ used are indicated in figure 2. The pure spinor $\lambda^\alpha$ is in the representation $(001)(1)$ and $D_\alpha$ in $(010)(1)$.

We think of spinors as 2-component quaternionic objects, with $\gamma^a$ (strictly speaking $\sigma$-matrices, different for the two chiralities) acting from the left and imaginary quaternions $e_i$, generating $SU(2)$, from the right. Scalar product of two spinors $\alpha$ and $\beta$ (of opposite chiralities) is defined as $(\alpha|\beta) \equiv \text{Re} (\alpha^\dagger \beta)$, with conjugation and real part being quaternionic.
If the procedure from the previous section is attempted, a qualitatively different set of fields emerges in the cohomology. The complex of representations now is

\[
(000)(0) \xrightarrow{Q} (010)(1) \xrightarrow{Q} (020)(2) \xrightarrow{Q} \cdots \xrightarrow{Q} (0n0)(n) \xrightarrow{Q} \cdots
\]

and the zero-momentum cohomology is given in table 2.

\[
\begin{align*}
\text{dim} & = 0 & (000)(0) \\
\frac{1}{2} & \bullet & \bullet \\
1 & \bullet & (100)(0) & \bullet \\
\frac{3}{2} & \bullet & (001)(1) & \bullet & \bullet \\
2 & \bullet & (000)(2) & \bullet & \bullet & \bullet \\
\frac{5}{2} & \bullet & \bullet & \bullet & \bullet & \bullet & \bullet \\
3 & \bullet & \bullet & \bullet & \bullet & \bullet & \bullet & \bullet \\
\frac{7}{2} & \bullet & \bullet & \bullet & \bullet & \bullet & \bullet & \bullet & \bullet \\
4 & \bullet & \bullet & \bullet & \bullet & \bullet & \bullet & \bullet & \bullet & \bullet \\
\frac{9}{2} & \bullet & \bullet & \bullet & \bullet & \bullet & \bullet & \bullet & \bullet & \bullet & \bullet & \bullet & \bullet
\end{align*}
\]

Table 2. The cohomology of the $N = 1$, $D = 6$ SYM complex.

The condition $Q\Psi = 0$ does not imply the equations of motion; instead the $SU(2)$ triplet of auxiliary fields appears at dimension 2. There is no cohomology at $n > 1$, i.e., no antifields. There are no representations for anti-fields, and consequently no room for equations of motion, or currents.

The antifields turn out to be found in a separate complex,

\[
(000)(2) \xrightarrow{Q} (010)(3) \xrightarrow{Q} (020)(4) \xrightarrow{Q} \cdots \xrightarrow{Q} (0n0)(n + 2) \xrightarrow{Q} \cdots
\]
with the cohomology given by table 3. Here, the dimension of the triplet field is chosen to be 2 in order to match with the dimensionalities of the current multiplet.

\[
\begin{array}{cccccc}
 n = 0 & n = 1 & n = 2 & n = 3 & n = 4 \\
\text{dim} = 2 & (000)(2) & & & \\
\frac{5}{2} & (010)(1) & \bullet & & \\
3 & (100)(0) & \bullet & \bullet & \\
\frac{7}{2} & & \bullet & \bullet & \bullet & \bullet \\
4 & \bullet & (000)(0) & \bullet & \bullet & \bullet \\
\frac{9}{2} & \bullet & \bullet & \bullet & \bullet & \bullet \\
5 & \bullet & \bullet & \bullet & \bullet & \bullet \\
\frac{11}{2} & \bullet & \bullet & \bullet & \bullet & \bullet \\
6 & \bullet & \bullet & \bullet & \bullet & \bullet \\
\frac{13}{2} & \bullet & \bullet & \bullet & \bullet & \bullet \\
\end{array}
\]

Table 3. The cohomology of the $N = 1$, $D = 6$ SYM antifield complex.

A key observation here is that both the fields and the antifields can be accommodated in the same field $\Psi$, if it, in addition to the pure spinor $\lambda$ of dimension $-\frac{1}{2}$, depends on a bosonic triplet $c_i$ of dimension $-2$ with $c_i(\lambda e_i)\alpha = 0$, where $e_i$, $i = 1, 2, 3$, are imaginary quaternions acting on the two-component $\lambda$ by right multiplication. Then the $\lambda$ expansion at order $c$ gives exactly the series of representations (3.2). Multiplying this condition one more time from the right with $c_i e_i$ gives the necessary constraint $c_i c_i = 0$.

The “first quantised Hilbert space” consists of functions of $\lambda^\alpha$, $c_i$, $x^a$ and $\theta^\alpha$, with the constraints $T^a \equiv (\lambda \gamma^a \lambda) = 0$, $t^a \equiv c_i (\lambda e_i)^a = 0$ and $\tau \equiv c_i c_i = 0$. Just like the constraint $T^a = 0$ is solved by a complexified $\lambda$, the full set of constraints are solved by complex $\lambda$ and $c$. Any state is a formal series expansion in non-negative powers of $\lambda^\alpha$ and $c_i$. If we call this space $\mathcal{P}$, the Hilbert space is $\mathcal{H} = \mathcal{P} / T \mathcal{P}$, where $T \mathcal{P}$ is the ideal in $\mathcal{P}$ generated by the constraints.

$\lambda^\alpha$ is in $(001)(1)$ and $c_i$ in $(000)(2)$. The constraints imply that $\lambda^\alpha c^\nu$ only contains $(00n)(n + 2\nu)$. We let $n$ and $\nu$ denote the degree of homogeneity in $\lambda$ and $c$, respectively, throughout the paper. The pure spinor BRST operator is $Q = (\lambda D)$. The ghost and fields sit in the cohomology of $Q$ at $c^0$ and the antifields and antighost at $c^1$. The auxiliary field sits at $c^0 \lambda^1 \theta^3$ and the corresponding antifield at $c^1 \lambda^0 \theta^0$, so one will need an operator $s \sim cwD^3$ to go from one to the other. Here, $w_\alpha$ is the derivative with respect to $\lambda^\alpha$. The derivative with respect to $c_i$ is denoted $b_i$. 
There is also cohomology of $Q$ at higher powers of $c$, although only at $\lambda^0$. The cohomologies at $c^\nu$ is given in Table 4, and the general positions of non-vanishing cohomology in the power expansion in $\lambda$ and $c$ is found in figure 3. We are not sure what they signify, but note that they sit “beyond” the scalar representation of the antifield in the power expansion, that would provide a measure. Each of these off-shell multiplets contain $8(2\nu - 1)$ bosonic fields and the same number of fermions.

\[
\begin{array}{cccccc}
 n = 0 & n = 1 & n = 2 & n = 3 & n = 4 \\
\hline
 n + \nu & 0 & 1 & 2 & 3 & 4 & 5 \\
\hline
\end{array}
\]

Table 4. The cohomology of $N = 1, D = 6$ at $\nu \geq 2$.

Figure 3. Position of cohomologies in the power series.
4. The new BRST operator

Derivatives are not well defined on $\mathcal{H}$, since they can map an element of the ideal to something outside the ideal. For example, $w_\alpha \cdot (\lambda \gamma^a \lambda) = 2(\gamma^a \lambda)\alpha \notin T\mathcal{P}$. One can define modified “quantum” derivatives that map the ideal to itself, so that they are well-defined on $\mathcal{H}$. A good derivative, normalised so that it acts exactly as $w$ on the representations $(0n0)(n+2\nu)$, is

$$\tilde{w}_\alpha = \frac{1}{2(1+n+2\nu)}[(2+n+2\nu)w_\alpha - j_i(we_i)\alpha] , \quad (4.1)$$

where $j_i$ is the $su(2)$ generator $j_i = (\lambda we_i) + 2\epsilon_{ijk}c_jb_k$ and $n$ and $\nu$ are the number operators for $\lambda$ and $c$, $n = (\lambda w)$, $\nu = c_i b_i$. However, a careful analysis of the most general Ansatz for an operator $s$ shows that it cannot be constructed using such gauge invariant operators only, if one demands that it must act in the cohomology of $Q$. Instead, as a “last resort”, one may allow for an operator acting in the cohomology of $Q$ to be not strictly gauge invariant, but gauge invariant modulo $Q$-exact terms.

There is essentially two terms to write down. The Ansatz is

$$s = [Q, f(n)c_i(w\gamma^aw)(D\gamma_a De_i)] + g(n)c_i(w\gamma^a Da_i)\partial_a , \quad (4.2)$$

where $f$ and $g$ are some functions of $n$ (the functions could in principle depend also on $\nu$, but everything else in $s$ commutes with $\nu$). The first term contains $(w\gamma^aw)$, which is not gauge invariant, but $[T^a, (w\gamma^bw)] = -4(j^{ab} + (n+4)\eta^{ab})$, where $j^{ab} = (\lambda \gamma^{ab} w)$, which is a gauge invariant operator (i.e., well-defined on the ideal). Since $[Q, T^a] = 0$, the first term in eq. (4.2) is gauge invariant modulo $Q$-exact terms. An analogous statement holds for invariance under $t^a$. The second term in eq. (4.2) is gauge invariant.

We will now check if it is possible to obtain $\{Q, s\} = 0$ on $\mathcal{H}$ for some functions $f$, $g$. Let us write $s = s_1 + s_2 = [Q, u] + s_2$ for the terms above. We then have

$$\{Q, s_1\} = \{Q, [Q, u]\} = \frac{1}{2}([Q, Q], u] = -[T^a, u]\partial_a$$

$$= 4c_i f(n)(j^{ab} + (n+4)\eta^{ab})(D\gamma_b De_i)\partial_a . \quad (4.3)$$

Here, we have dropped $[T^a, f(n)] = (f(n-2) - f(n))T^a$, which standing on the left gives pure gauge (an element in the ideal $T\mathcal{P}$). We also have

$$[Q, s_2] = (g(n-1) - g(n))(\lambda D)c_i(w\gamma^a De_i)\partial_a - g(n)c_i(D\gamma^a De_i)\partial_a . \quad (4.4)$$

† In this and following expressions, no “normal ordering” is assumed. Ordering is of course relevant, and all expressions are ordered as they are written.
The term from the anticommutator of the $D$’s in $Q$ and $s_2$ has been dropped; it vanishes thanks to $t^\alpha = 0$. Now, we want to do a Fierz rearrangement of the first of these terms, i.e., of $(\lambda D)c_i(w\gamma^a D_e_i)$. Doing a general Fierz in the two $D$’s, one finds that only $(D\gamma^b D_e_i)$ contributes (again, due to $t^\alpha = 0$), so that $(\lambda D)c_i(w\gamma^a D_e_i) = -\frac{1}{4}(j^{ab} - n\eta^{ab})c_i(D\gamma^b D_e_i)$. Taken together, this gives two types of terms in $[Q, s_1]$ and $[Q, s_2]$, those with $j$ and those with $\eta$. Demanding that they cancel gives the equations

\begin{align}
  j : & \quad 4f(n) + \frac{1}{4}(g(n) - g(n - 1)) = 0 , \\
  \eta : & \quad 4(n + 4)f(n) - g(n) - \frac{1}{4}(g(n) - g(n - 1)) = 0 .
\end{align}

The equations have a solution which is unique up to an overall normalisation:

\begin{align}
  f(n) &= -\frac{k}{(n+1)(n+2)(n+3)(n+4)} , \\
  g(n) &= \frac{8k}{(n+3)(n+4)} .
\end{align}

An explicit evaluation of the commutator $[Q, u]$ gives

\begin{align}
  s = \frac{3k}{(n+1)(n+2)(n+3)(n+4)}Qc_i(w\gamma^a w)(D\gamma^a D_e_i) \\
  - \frac{2k}{(n+2)(n+3)(n+4)}c_i(wD^3_i) + \frac{8k}{(n+2)(n+3)}c_i(w\gamma^a D_e_i)\partial_a ,
\end{align}

where $(D^3)_{i}^\alpha = (\gamma^a D)^\alpha(D\gamma^a D_e_i)^\alpha\partial_a$ is the totally antisymmetric product $\wedge^3 D$ in $(001)(3)$. The first term vanishes when one chooses a gauge $(w\gamma^a w)\Psi = 0$ for the wave function, and vanishes in the cohomology of $Q$. The remaining terms are the essential ones. It is obvious from the structure of the cohomology of $Q$ (figure 3) that $s^2 = 0$ on any element of the cohomology.

The calculations performed in this section make use of the properties of antisymmetric tensor products of spinors, that dictate the content of superfields. The upper half of a scalar superfield is shown in figure 4 (the remaining part is of course given by the conjugate representations).

\texttt{Figure 4: The antisymmetric products of chiral spinors in $D = 6$}
5. Conclusions and outlook

We have enlarged the pure spinor space in \( D = 6 \) to include also a bosonic \( SU(2) \) triplet \( c_i \). The generalised pure spinor constraints are \( T^a \equiv (\lambda \gamma^a \lambda) = 0 \), \( t^\alpha \equiv c_i (\lambda e_i)^\alpha = 0 \) and \( \tau \equiv c_i c_i = 0 \). The cohomology of \( Q = \lambda^\alpha D_\alpha \) in a scalar field then includes not only ghosts and fields but also antifields and antighosts. We have found a new BRST operator \( s \) whose cohomology imply the equations of motion. However, in order to define \( s \), we had to restrict to the cohomology of \( Q \), outside of which \( s \) is not well defined. This means that it is not possible to form a modified BRST operator \( Q + s \) on the scalar field. If one wants to write a BV action for \( D = 6, N = 1 \) super-Yang–Mills (after introducing some non-minimal set of pure spinor variables [8]), it would need to be formulated on a field satisfying \( Q \Psi + \Psi^2 = 0 \). This is clearly a weakness of the formalism, but it is not at all clear that it is a generic one.

It will be interesting to continue the investigation for other non-maximal supersymmetric theories. Supergravity with \( N = 1 \) supersymmetry in \( D = 10 \) is a natural candidate. Here, the off-shell/on-shell discussion is more complicated, since the superspace formulation contains some, but not all, equations of motion. Preliminary results show that the structure is clearly reflected in the corresponding pure spinor cohomology. There is also indication that the situation may be better than for \( D = 6, N = 1 \) super-Yang–Mills theory, in the sense that it may be possible to find a single modified BRST operator on the unconstrained field depending on a generalised pure spinor [19]. It is also conceivable that a generalised pure spinor space provides a natural setting for heterotic supergravity. Other interesting cases are half-maximal euclidean supergravities in \( D = 6 \) and \( D = 7 \) in which backgrounds for topological string theory and M-theory may be embedded.

A superspace treatment of deformation of a non-maximally symmetric theory, such as \( D = 6 \) super-Yang–Mills, would rely on deformation of the condition \( s \Psi = 0 \). This would apply for Born–Infeld dynamics, or its derivative expansion.

Another issue that may be addressed using similar methods is the relation between the two superspace versions of (linearised) 11-dimensional supergravity. The same on-shell multiplet is found using either a vector field (the “vielbein complex”) or a scalar field (the “\( C \)-field complex”). Only the \( C \)-field complex is appropriate for writing a linearised action. If one can enlarge the pure spinor space to encompass the two complexes in the same field, it may provide a starting point for the search for a superspace action. Similar thoughts have been touched on in refs. [20,21].

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