Remarks on the M5-Brane

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Abstract

The fivebrane of M theory – the M5-brane – is an especially interesting object. It plays a central role in a geometric understanding of the Seiberg–Witten solution of N=2 D=4 gauge theories as well as in certain new 6d quantum theories. The low energy effective action is an interacting theory of a (2,0) tensor multiplet. The fact that this multiplet contains a two-form gauge field with a self-dual field strength poses special challenges. Recent progress in addressing those challenges is reviewed.

1 Super P-Branes

Type II superstring theories in 10d and M theory in 11d have 32 supercharges. They admit various BPS $p$-brane configurations that break half of the supersymmetry. The effective world-volume theory of any of these $p$-branes contains 16 Goldstone fermion fields, corresponding to the broken supersymmetries, which describe 8 propagating fermionic modes in the world volume. Also, associated to the broken translation symmetries, there are Goldstone boson fields describing $D - p - 1$ propagating bosonic modes. Since the world-volume theory has unbroken supersymmetries, it must contain an equal number of bosonic and fermionic degrees of freedom.

The discrepancy in the numbers given above is made up by the addition of $9 + p - D$ bosonic modes that correspond to world-volume gauge fields. There are three basic cases of

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interest. The first case is when no additional degrees of freedom are required, in other words when $D = p + 9$. This is the case for Type IIA or Type IIB strings in 10d or the M2-brane (also called the supermembrane) of M theory in 11d. In each case there are eight dimensions transverse to the brane. In the Type IIB case there is actually an infinite family of strings, labelled by a pair of integers (corresponding to two-form charges). Stability requires that these integers are relatively prime \[1, 2\]. These \((m, n)\) strings transform into one another under the nonperturbative $SL(2, \mathbb{Z})$ duality group of the IIB theory. Building on earlier work that studied the relation between the (1,0) fundamental string and the (0,1) D string \[3, 4\], Townsend has recently shown \[5\] how to construct world-volume theories for this family of strings that makes their $SL(2, \mathbb{Z})$ duality properties manifest. The key step is to introduce an $SL(2, \mathbb{Z})$ doublet of $U(1)$ gauge fields. This is possible, because a gauge field in 2d is nonpropagating and does not affect the counting of physical degrees of freedom.

The second case is when the additional $9 + p - D$ bosonic modes are provided by a $U(1)$ gauge field. Since such a gauge field introduces $p - 1$ propagating degrees of freedom, the counting works precisely when $D = 10$, for all values of $p$. These are the celebrated type II D-branes \[6\], which carry charges of Ramond–Ramond gauge fields. The dimension $p$ is even in the IIA case and odd in the IIB case. Explicit supersymmetric world-volume actions with local kappa symmetry were constructed last autumn for these D-branes by a number of groups \[4, 8, 9\].

The third case, which is the one I will focus in the remainder of this talk, is the M5-brane. In this case the $9 + 5 - 11 = 3$ extra bosonic degrees of freedom are provided by a two-form potential $B_{\mu\nu}$, whose field strength is self dual in the linearized approximation. To understand the counting, note that massless particles in 6d are classified by the $Spin(4) = SU(2) \times SU(2)$ little group. The states in question belong to the $(3, 1)$ representation of this group. Since parity interchanges the two $SU(2)$’s, this is a chiral boson.

The M5-brane has a simple relation to two of the $p$-branes of Type IIA superstring theory. Recall that at strong coupling type IIA string theory on $R^{10}$ should be reinterpreted as M theory on $R^{10} \times S^1$ \[10, 11\], where the radius of the circular 11th dimension is proportional to the 2/3 power of the type IIA string coupling constant (in the string frame). There are two possibilities for the fate of the M5-brane. Either one of its dimensions is wrapped on the circle or none of them are. In the former case, one obtains the D4-brane and in the latter the NS solitonic fivebrane. To understand how the D4-brane arises, one first carries out a
“double dimensional reduction” in which one of the world-volume coordinates is identified with the circular spatial dimension and then the zero modes in this direction are extracted. Next a duality transformation in the 5d world volume is required to replace the two-form potential by the $U(1)$ gauge field that is characteristic of the D-brane. These relationships were crucial in Witten’s recent analysis of a IIA brane configuration that corresponds to an $N = 2 \, D = 4$ gauge theory [12]. The 10d picture involved a number of interconnected D4-branes and NS5-branes, but the configuration could be reinterpreted as a single M5-brane in 11d. This led to a simple geometrical interpretation of the Seiberg–Witten Riemann surface as two of the dimensions of the M5-brane! (See [13, 14] for reviews of SW theory and related string theoretic approaches to understanding them geometrically.) One implication of this result is that the M5-brane action, which we will describe, encodes a great deal of quantum information even though all of our manipulations appear to be classical.

2 The Bosonic Part of the M5-Brane Action

2.1 Symmetries

In a covariant formulation all global symmetries, broken and unbroken, are exhibited. In addition, there are various local symmetries of the world-volume theory (discussed below). When these local symmetries are used to fix a physical gauge, the broken global symmetries become non-linearly realized with associated Goldstone particles. Thus, the covariant M5-brane action has manifest global 11d super-Poincaré symmetry realized by the superspace transformations $\delta \theta = \epsilon, \delta X^M = \bar{\epsilon} \Gamma^M \theta + a^M$. However, the unbroken global symmetry of the gauge-fixed world-volume theory is just $(2,0)$ super-Poincaré symmetry in 6d. The two chiral spinors and five scalars in the $(2,0)$ tensor multiplet are the Goldstone particles associated to the broken global symmetries.

There are three kinds of local symmetries. Two of them are essentially the same for all super $p$-branes. They are world-volume diffeomorphism symmetry and local kappa symmetry. The diffeomorphism symmetry implies that the components of $X^M$ along the $p$-brane directions are unphysical and can be gauged away, e.g. by choosing a static gauge. Local kappa symmetry has the effect of eliminating half of the fermionic degrees of freedom, which is essential to get the counting required by supersymmetry. Since this paper only describes the bosonic sector of the theory, nothing more will be said about this symmetry, except to
not that the complete kappa-symmetric action has been constructed recently [13, 16, 17]. The third class of local symmetries are those that are required to properly describe a self-dual tensor. I am not referring to the standard gauge transformation \( \delta B = d\Lambda \), which is rather trivial, but rather to the new PST gauge symmetries [18, 19], described below, which are required to decouple the anti-chiral components of \( B \).

2.2 Topological Issues

In the analysis presented below I will only analyze infinitesimal diffeomorphisms of the 6d world-volume theory. However, when the 6d manifold is topologically nontrivial one should also analyze “large” diffeomorphisms, which are not continuously connected to the identity. In other words, one should check modular invariance.

To explain the issues let us briefly consider an example. Suppose the 11d spacetime is topologically of the form \( \mathbb{R}^7 \times K3 \) and that the 6d world volume is \( \Sigma \times K3 \), where \( \Sigma \) is a Riemann surface and the world-volume \( K3 \) is wrapped on the space-time \( K3 \). In this case one can carry out a double dimensional reduction obtaining an effective string action in a 7d spacetime. In fact, it was noted by Witten that M theory compactified on \( K3 \) is dual to the heterotic string theory compactified on \( T^3 \) [11]. Furthermore, Harvey and Strominger noted that the double dimensional reduction gives precisely the heterotic string in 7d [20]. This has been analyzed in greater detail, using the explicit M5-brane action, by Cherkis and me [21].

Because \( K3 \) has 19 anti-self-dual 2-forms (\( b^-_2 = 19 \)) and three self-dual ones (\( b^+_2 = 3 \)), the zero modes of the \( B \) field give rise to 19 left-moving and 3 right-moving chiral bosons on \( \Sigma \). They are compact and their momentum lattice is the even self-dual Narain lattice \( \Gamma_{19,3} \). This describes the heterotic string in 7d, which is well-known to have modular invariance, precisely because the lattice is even and self-dual. This strongly suggests, but does not completely prove, that the M5-brane on \( \Sigma \times K3 \) also has the requisite modular invariance. For a more detailed discussion of these issues the reader is referred to a recent analysis by Witten [22].

2.3 The Noncovariant Formulation

Ref. [23] analyzed the problem of coupling a 6d self-dual tensor gauge field to a metric field so as to achieve general coordinate invariance. It presented a formulation in which one
direction is treated differently from the other five. At the time that work was done, the author
knew of no straightforward way to make the general covariance manifest. However, shortly
thereafter a paper appeared [18] that presents equivalent results using a manifestly covariant
formulation [19], which we refer to as the PST formulation. In the following both approaches
and their relationship are described. These results have been generalized to supersymmetric
actions with local kappa symmetry [15, 16, 17], but here we will only consider the bosonic
theories.

Let us denote the 6d world-volume coordinates by $\sigma^\hat{\mu} = (\sigma^\mu, \sigma^5)$, where $\mu = 0, 1, 2, 3, 4$.
The $\sigma^5$ direction is singled out as the one that will be treated differently from the other five.\footnote{This is a space-like direction, but one could also choose a time-like one.}
The 6d metric $G_{\hat{\mu}\hat{\nu}}$ contains 5d pieces $G_{\mu\nu}, G_{\mu5}$, and $G_{55}$. All formulas will be written with
manifest 5d general coordinate invariance. As in refs. [24, 23], we represent the self-dual
tensor gauge field by a $5 \times 5$ antisymmetric tensor $B_{\mu\nu}$, and its 5d curl by $H_{\mu\nu\rho} = 3 \partial_{[\mu} B_{\nu\rho]}$.
A useful quantity is the dual
$$\tilde{H}^{\mu\nu} = \frac{1}{6} \epsilon^{\mu\nu\rho\lambda\sigma} H_{\rho\lambda\sigma}. \quad (1)$$
It was shown in ref. [23] that a class of generally covariant bosonic theories can be
represented in the form $L = L_1 + L_2 + L_3$, where
$$L_1 = -\frac{1}{2} \sqrt{-G} f(z_1, z_2),$$
$$L_2 = -\frac{1}{4} \tilde{H}^{\mu\nu} \partial_5 B_{\mu\nu}, \quad (2)$$
$$L_3 = \frac{1}{8} \epsilon_{\mu\nu\rho\lambda\sigma} \frac{G^{5\rho}}{G_{55}} \tilde{H}^{\mu\nu} \tilde{H}^{\lambda\sigma}.$$
The notation is as follows: $G$ is the 6d determinant ($G = \det G_{\hat{\mu}\hat{\nu}}$) and $G_5$ is the 5d determi-
nant ($G_5 = \det G_{\mu\nu}$), while $G^{55}$ and $G^{5\rho}$ are components of the inverse 6d metric $G^{\hat{\mu}\hat{\nu}}$. The
$\epsilon$ symbols are purely numerical with $\epsilon^{01234} = 1$ and $\epsilon^{\mu\nu\rho\lambda\sigma} = -\epsilon_{\mu\nu\rho\lambda\sigma}$. A useful relation is
$G_5 = GG^{55}$. The $z$ variables are defined to be
$$z_1 = \frac{\text{tr}(G\tilde{H}G\tilde{H})}{2(-G_5)},$$
$$z_2 = \frac{\text{tr}(G\tilde{H}G\tilde{H}G\tilde{H}G\tilde{H})}{4(-G_5)^2}. \quad (3)$$
The trace only involves 5d indices:
$$\text{tr}(G\tilde{H}G\tilde{H}) = G_{\mu\nu}\tilde{H}^{\nu\rho} G_{\rho\lambda} \tilde{H}^{\lambda\mu}. \quad (4)$$
The quantities $z_1$ and $z_2$ are scalars under 5d general coordinate transformations.

Infinitesimal parameters of general coordinate transformations are denoted $\xi^{\hat{\mu}} = (\xi^\mu, \xi)$. Since 5d general coordinate invariance is manifest, we focus on the $\xi$ transformations only. The metric transforms in the standard way

$$\delta_\xi G_{\hat{\mu}\hat{\nu}} = \xi \partial_5 G_{\hat{\mu}\hat{\nu}} + \partial_{\hat{\mu}} \xi G_{5\hat{\nu}} + \partial_{\hat{\nu}} \xi G_{\hat{\mu}5}. \quad (5)$$

The variation of $B_{\mu\nu}$ is given by a more complicated rule, whose origin is explained in ref. [23]:

$$\delta_\xi B_{\mu\nu} = \xi K_{\mu\nu}, \quad (6)$$

where

$$K_{\mu\nu} = 2 \frac{\partial (L_1 + L_3)}{\partial \tilde{H}_{\mu\nu}} = K_{\mu\nu}^{(1)} f_1 + K_{\mu\nu}^{(2)} f_2 + K_{\mu\nu}^{(e)} \quad (7)$$

with

$$K_{\mu\nu}^{(1)} = \sqrt{-G_{\nu}} \frac{G_{G\tilde{H}}G_{\mu\nu}}{(-G_5)} \quad (8)$$

$$K_{\mu\nu}^{(2)} = \sqrt{-G_{\nu}} \frac{G_{G\tilde{H}G\tilde{H}G_{\mu\nu}}}{(-G_5)^2} \quad (8)$$

$$K_{\mu\nu}^{(e)} = \epsilon_{\mu\nu\rho\lambda\sigma} \frac{G_{5\rho}}{2G_{55}} \frac{\tilde{H}_{\lambda\sigma}}{\sqrt{-G_5}} \quad (8)$$

and we have defined

$$f_i = \frac{\partial f}{\partial z_i}, \quad i = 1, 2. \quad (9)$$

Assembling the results given above, ref. [23] showed that the required general coordinate transformation symmetry is achieved if, and only if, the function $f$ satisfies the nonlinear partial differential equation [25]

$$f_1^2 + z_1 f_1 f_2 + (\frac{1}{2} z_1^2 - z_2) f_2^2 = 1. \quad (10)$$

As discussed in [24], this equation has many solutions, but the one of relevance to the M theory five-brane is

$$f = 2 \sqrt{1 + z_1 + \frac{1}{2} z_1^2 - z_2}. \quad (11)$$

For this choice $L_1$ can reexpressed in the Born–Infeld form

$$L_1 = -\sqrt{-\det \left( G_{\mu\nu} + i G_{\mu\rho} G_{\nu\lambda} \tilde{H}^{\rho\lambda} / \sqrt{-G_5} \right)}. \quad (12)$$

This expression is real, despite the factor of $i$, because it is an even function of $\tilde{H}$. 
2.4 The PST Formulation

In ref. [18] (using techniques developed in ref. [19]) equivalent results are described in a manifestly covariant way. To do this, the field $B_{\mu\nu}$ is extended to $B^\hat{\mu}\hat{\nu}$ with field strength $H^\hat{\rho}\hat{\sigma}$. In addition, an auxiliary scalar field $a$ is introduced. The PST formulation has new gauge symmetries (described below) that allow one to choose the gauge $B_{\mu\nu} = 0, a = \sigma^5$ (and hence $\partial_\mu a = \delta^5_\mu$). In this gauge, the covariant PST formulas reduce to the ones given above.

Equation (12) expresses $L_1$ in terms of the determinant of the $6 \times 6$ matrix

$$M_{\hat{\mu}\hat{\nu}} = G_{\hat{\mu}\hat{\nu}} + i \frac{G_{\hat{\mu}\hat{\rho}}G_{\hat{\nu}\hat{\lambda}}}{\sqrt{-G}} \tilde{H}^{\hat{\rho}\hat{\lambda}}. \quad (13)$$

In the PST approach this is extended to the manifestly covariant form

$$M^{\text{cov.}}_{\hat{\mu}\hat{\nu}} = G_{\hat{\mu}\hat{\nu}} + i \frac{G_{\hat{\mu}\hat{\rho}}G_{\hat{\nu}\hat{\lambda}}}{\sqrt{-G}} \tilde{H}^{\rho\lambda}_{\text{cov.}}. \quad (14)$$

The quantity

$$(\partial a)^2 = G^{\hat{\mu}\hat{\nu}} \partial_\mu a \partial_\nu a \quad (15)$$

reduces to $G^{55}$ upon setting $\partial_\mu a = \delta^5_\mu$, and

$$\tilde{H}^{\rho\lambda}_{\text{cov.}} = \frac{1}{6} \epsilon^{\rho\lambda\hat{\mu}\hat{\nu}\hat{\sigma}} H_{\hat{\mu}\hat{\nu}\hat{\sigma}} \partial_\tau a \quad (16)$$

reduces to $\tilde{H}^{\rho\lambda}$. Thus $M^{\text{cov.}}_{\hat{\mu}\hat{\nu}}$ replaces $M_{\hat{\mu}\hat{\nu}}$ in $L_1$. Furthermore, the expression

$$L' = -\frac{1}{4(\partial a)^2} \tilde{H}_{\text{cov.}}^{\hat{\mu}\hat{\nu}} H_{\hat{\mu}\hat{\nu}\hat{\rho}} G^{\hat{\rho}\hat{\lambda}} \partial_\lambda a, \quad (17)$$

which transforms under general coordinate transformations as a scalar density, reduces to $L_2 + L_3$ upon gauge fixing. It is interesting that $L_2$ and $L_3$ are unified in this formulation.

Let us now describe the new gauge symmetries of ref. [18]. Since degrees of freedom $a$ and $B_{\mu\nu}$ have been added, corresponding gauge symmetries are required. One of them is

$$\delta B_{\mu\nu} = 2\phi_{[\mu} \partial_{\nu]} a, \quad (18)$$

where $\phi_{\bar{\mu}}$ are infinitesimal parameters, and the other fields do not vary. In terms of differential forms, this implies $\delta H = d\phi \wedge da$. $\tilde{H}^{\hat{\rho}\hat{\lambda}}_{\text{cov.}}$ is invariant under this transformation, since it corresponds to the dual of $H \wedge da$, but $da \wedge da = 0$. Thus the covariant version of $L_1$ is invariant under this transformation. The variation of $L'$, on the other hand, is a total derivative.
The second local symmetry involves an infinitesimal scalar parameter $\varphi$. The transformation rules are $\delta G_{\hat{\mu}\hat{\nu}} = 0$, $\delta a = \varphi$, and

$$
\delta B_{\hat{\mu}\hat{\nu}} = \frac{1}{(\partial a)^2} \varphi H_{\hat{\rho}\hat{\mu}\hat{\nu}} G^{\hat{\rho}\hat{\lambda}} \partial_{\hat{\lambda}} a + \varphi V_{\hat{\mu}\hat{\nu}},
$$

where the quantity $V_{\hat{\mu}\hat{\nu}}$ is to be determined. Rather than derive it from scratch, let’s see what is required to agree with the previous formulas after gauge fixing. In other words, we fix the gauge $\partial_{\hat{\mu}} a = \delta_{\hat{\mu}}$ and $B_{\mu\hat{\nu}} = 0$, and figure out what the resulting $\xi$ transformations are. We need

$$
\delta a = \varphi + \xi \partial_{\hat{\mu}} a = \varphi + \xi = 0,
$$

which tells us that $\varphi = -\xi$. Then

$$
\begin{align*}
\delta_\xi B_{\mu\nu} & = \frac{1}{(\partial a)^2} \varphi H_{\mu\nu\rho} G^{\rho\lambda} \partial_{\lambda} a + \varphi V_{\mu\nu} + \xi H_{5\mu\nu} \\
& = -\xi \left( \frac{G^{\rho\delta}}{G^{55}} H_{\mu\nu\rho} + V_{\mu\nu} \right) = \xi (K_{\mu\nu}^{(e)} - V_{\mu\nu}).
\end{align*}
$$

Thus, comparing with eqs. (19) and (20), we need the covariant definition

$$
V_{\hat{\mu}\hat{\nu}} = -2 \frac{\partial L_I}{\partial H_{\hat{\mu}\hat{\nu}}},
$$

(22)

to achieve agreement with our previous results.

The presence of the factor $(\partial a)^2$ in various denominators is potentially problematical unless the six manifold on which the action is defined satisfies an appropriate topological condition. The precise determination of what that condition is, and how it relates to the condition of modular invariance discussed earlier, deserve further study.

3 Discussion

The explicit formula for the M5-brane action should be useful for a number of purposes. For one thing, Witten’s derivation [12] of the Seiberg–Witten quantum effective action only required knowledge of the quadratic (free) approximation. The complete nonlinear formula should allow one to extend the analysis to derive some of the higher dimension terms in the 4d quantum effective action. Exactly which ones could be derived is not entirely clear, but I suppose that the requisite mathematical control would be best for those that are encoded in holomorphic functions.
Evidence has been mounting that there are several classes of interacting non-gravitational 6d quantum theories. Recently, Seiberg has argued that starting from a set of $k$ parallel M5-branes one can define a class of these theories by means of a suitable limit in which they decouple from the bulk degrees of freedom while still remaining interacting \cite{26}. Such a theory is certainly not a conventional quantum field theory, and it is not yet entirely clear how it should best be formulated. In any case, if $k > 1$ there are string-like excitations due to supermembranes attaching between a pair of the fivebranes. If $k = 1$, and if one of the transverse dimensions is compact, so that there is IIA interpretation, then there are also strings due to a supermembrane that starts and ends on the fivebrane and wraps around the circular transverse dimension. These considerations have led to suggestions that these might be viewed as “tensionless” or “noncritical” string theories. In any case, whatever the proper fundamental formulation will turn out to be, the formulas we have obtained should be viewed as the low-energy effective action for the $k = 1$ case, just as 11d supergravity is the low-energy effective description of M theory. In fact, if these 6d theories can be formulated as matrix models, the analogy would be a good one.

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