Abstract: We consider super-membranes ending on M5-branes, with the aim of deriving the appropriate matrix theories describing different situations. Special attention is given to the case of non-vanishing (selfdual) $C$-field. We identify the relevant deformation of the six-dimensional super-Yang–Mills theory whose dimensional reduction is the matrix theory for membranes in the presence of M5-branes. Possible applications and limitations of the models are discussed.

1. Introduction

Supermembrane [1] theory [2] is a very promising candidate for a microscopic description of M-theory. Although it is not background invariant, it gives a completely new picture of the nature of space and time at small scales, together with a description of quantum-mechanical states that goes beyond local quantum field theory. These features are most clear in the matrix [3] truncation [2] of the membrane. It is widely appreciated that first-quantised supermembrane theory through its continuous spectrum [4] is capable of describing an entire (“multi-particle”) Fock space. For reviews on the subject of membranes and matrices, see ref. [5]. Due to the immense technical difficulties associated with actual calculations in the theory, which is non-linear and inherently non-perturbative, few quantitative features are known in addition to the general picture, which is supported by many qualitative arguments. The maybe most important one is the proof that $su(N)$ matrix theory has a unique supersymmetric ground state [6,7], which gives the relation to the massless degrees of freedom of $D = 11$ supergravity.

Many situations in M-theory backgrounds involve membranes that are not closed. Supermembranes may end on “defects”, i.e., 5-branes and 9-branes [8,9,10,11,12,13,14]. It is urgent to have some mathematical formulation of these situations in order to understand the microscopic properties of physics in such backgrounds. One old enigma is the nature of the theory on multiple 5-branes, which we address in the present talk. There are several
issues to be resolved. The membrane may be stretched between multiple 5-branes, and
the truncation has to be consistent with this situation. In addition, the C-field, the 3-form
potential of M-theory, may take some non-vanishing selfdual value on the 5-brane. The
new results contained in this talk refer to the latter situation. There are several reasons
to consider this specific situation. It should be connected to the theory on multiple M5-
branes, which is some kind of non-abelian theory of selfdual tensors [15]. It should be
possible to verify the decoupling limit of OM-theory [16] from microscopic considerations.
There might also be information about the open membrane metric [17,18] and maybe even
some clue concerning the proper generalisation of the string endpoint non-commutativity
to membranes [19,18].

We start out by reviewing the consistent truncation of membranes to matrices via
non-commutativity in section 2. Section 3 describes how this construction is generalised
to situations where the membrane has a boundary [20,21,14,22]. Here we review the
alternative constructions present in the literature, and discuss their relative applicability.
In section 4, we generalise the picture to include non-vanishing C-field, both light-like [23]
and general. We identify the deformation of the 6-dimensional super-Yang–Mills theory
whose dimensional reduction is the matrix theory associated to turning on the C-field. In
section 5, we discuss the possible applications and limitations of the model.

2. From membranes to matrices

We start from the action for the supermembrane coupled to an on-shell background of
$D = 11$ supergravity,

$$S = -T \int d^3 \xi \sqrt{-g} + T \int C. \tag{2.1}$$

Here, the metric and C-field are pullbacks from superspace to the bosonic world-volume.
In what follows, we will consider flat backgrounds, but allow for non-zero constant $C$.

Let us first remind of the consistent truncation to matrix theory of a closed membrane
(we just display the bosonic degrees of freedom; fermions are straightforwardly included).
Here, the $C$-field is irrelevant. In light-cone gauge, where reparametrisation invariance is
used up except for area-preserving diffeomorphisms of the membrane “space-sheet”. The
light-cone hamiltonian $p^-$ is given by

$$\frac{p^+ p^-}{A} = \int d^2 \xi \left( P_I P_I + \frac{T^2}{4} \{X^I, X^J\} \{X^I, X^J\} \right), \tag{2.2}$$

where $A$ is the parametric area of the space-sheet, and $\{A, B\} = \varepsilon^{ij} \partial_i A \partial_j B$ is the “Poisson
bracket” on the space-sheet. The remaining gauge invariance is generated by the Poisson
bracket as $\delta_f A = \{ f, A \}$ [24]. Even though it is known that the algebra of area-preserving diffeomorphisms in a certain sense is $su(\infty)$ [2,25], $su(N)$ is not contained as a subalgebra, and there is no way of getting to $su(N)$ matrix theory as a consistent truncation.

In order to obtain matrix theory as a consistent truncation, one introduces a non-commutativity on the membrane space-sheet (for simplicity, we consider a toroidal membrane), $[\xi^1,\xi^2] = \theta$, encoded in the Weyl-ordered star product $f \star g = f \exp(\frac{i}{2} \theta \varepsilon^{ij} \partial_i \partial_j) g$. Commutators between Fourier modes become $[e^{ik_1 \cdot \xi}, e^{ik_2 \cdot \xi}] = -2i \sin(\frac{\theta}{2} \varepsilon^{ij} k_i k_j') e^{i(k+k') \cdot \xi}$. The Poisson bracket is recovered as $\{ \cdot, \cdot \} = -i \lim_{\theta \to 0} \theta^{-1}[\cdot, \cdot]$. Choosing $\theta = \frac{2\pi}{N}$ implies that the functions $e^{iN\xi^1}, e^{iN\xi^2}$ are central. Their action on any function can consistently be modded out according to the equivalence relation $e^{iN\xi^1} \star f \approx f$. The remaining “square of functions” with mode numbers ranging from 0 to $N - 1$ generate $u(N)$ [2,25].

The model thus obtained as a consistent truncation of the supermembrane is an $su(N)$ supersymmetric matrix model, identical to the dimensional reduction to $D = 1$ of $D = 10$ super-Yang–Mills theory. This example sets the procedure we want to apply to other cases: deform by non-commutativity, replace Poisson brackets by commutators and perform a consistent truncation of the deformed theory. We would like to stress the importance of making a consistent truncation, in contrast to an approximation; the fact that the commutator obeys the same algebraic identities as the Poisson bracket means that one has control over the symmetries of the model, e.g. supersymmetry. The only symmetries that are lost in the matrix truncation are the super-Poincaré generators that are non-linearly realised in the light-cone gauge.

3. Matrix theory for membranes with boundary

Let us now turn to the first modification of the previous situation, namely when the membrane has boundaries (we think of these as lying on M5-branes, but much if what is said applies to any possible boundary). It is expected that the “no-topology” theorem that applies for closed membranes persists for membranes with boundary, so that it is irrelevant e.g. whether a membrane ending on only one 5-brane is modeled as a half sphere, a half torus or some more complicated manifold. This is an assumption we make; a proof would be desirable.

We can distinguish between two classes of approaches to this kind of configuration:

A. This approach was first physically motivated by double dimensional reduction to a D4-brane. The theory obtained after reduction is $D = 5$ super-Yang–Mills, and opening up the sixth direction should correspond to a strong coupling limit. In this limit, path
integrals are dominated by saddle points at the moduli space of “instanton” solitons. The moduli space of $N$ instantons in $U(k)$ SYM has dimension $4kN$. The matrix theory should have this space as Higgs branch. It is the dimensional reduction of a $D = 6$ $U(N)$ SYM with one adjoint and $k$ fundamental hypermultiplets [26]. We will motivate this from the point of view of the supermembrane.

B. For a fixed membrane topology (a half torus, say), the boundary conditions may be solved, at least when $C = 0$ (see [14]). For the 5 directions transverse to the (flat) $M_5$-brane one gets Dirichlet boundary conditions, which for torus topology means sine functions, and for the 4 transverse (2 have been eliminated when going to light-cône gauge) one gets Neumann boundary conditions, leading to cosine functions. The sine functions generate $SO(N)$ [13], and the cosine function transform as the symmetric representation. The matrix model obtained is the dimensional reduction of a $D = 6$ $SO(N)$ SYM with a hypermultiplet in the symmetric representation.

A couple of comments can be made. The first one concerns the global symmetries of the matrix theory and of the M-theory configuration it describes. A $D = 6$ SYM theory with hypermultiplets has a lagrangian

$$
\mathcal{L} = -\frac{1}{4} F_{\mu\nu}^A F^{A\mu\nu} + \frac{1}{2} \text{Re}(\lambda^A \gamma^\mu D_\mu \lambda^A) - \frac{1}{2} D^\mu \phi_I D_\mu \phi^*_I - \frac{1}{2} \text{Re}(\psi_I \gamma^\mu D_\mu \psi_I)
$$

$$
- (q^A)_{IJ} \text{Re}(\lambda^I \phi^*_J) + \frac{1}{8} (q^A)_{IJ} (q^A)_{KL} \phi^*_I \phi^*_J \phi^*_K \phi^*_L .
$$

We use the isomorphism $\text{Spin}(1, 5) \approx \text{SL}(2; \mathbb{H})$ and two-component quaternionic spinors. The scalars $\phi$ are quaternionic, $\phi = \phi_i e_i$, where $i = 1, \ldots, 4$. Indices $I, J, \ldots$ label the representation of the hypermultiplet, and $q$ is the representation matrix. For real hypermultiplets (as in case B above) there is an $SU(2)_L \times SU(2)_R$ R-symmetry realised as multiplication by unit quaternions as

$$
A \rightarrow A , \quad \lambda \rightarrow \lambda h_L ,
$$

$$
\phi \rightarrow h_L^* \phi h_R , \quad \psi \rightarrow \psi h_R .
$$

(3.2)

When the hypermultiplet is complex, the right action is occupied by the gauge group. If there is an even number of hypermultiplets in the same representation, there will however be a flavour symmetry $SU(2k)$ The representations (specified by dimensions) of the fields under the Lorentz rotations and R-symmetry are thus

$$
A : (6, 1, 1) , \quad \lambda : (4, 2, 1) ,
$$

$$
\phi : (1, 2, 2 \text{ or } 1) , \quad \psi : (4, 1, 2 \text{ or } 1)
$$

(3.3)

(the last alternative for $su(2)_R$ representations of the hypermultiplet is for the minimal content of a complex representation). The R-symmetry of the super-Yang–Mills theory is
the rotation symmetry of the membrane/matrix theory in light-cone gauge, and the SO(5) rotation symmetry remaining on the super-Yang-Mills side after dimensional reduction is the R-symmetry of the membrane/matrices.

The supersymmetry transformation rules are

\[ \delta \epsilon A^\mu = \text{Re}(\epsilon^* \gamma^\mu \lambda^A), \quad \delta \epsilon \lambda^A = \frac{1}{2} F^A_{\mu\nu} \gamma^{\mu\nu} \epsilon + \epsilon W^A, \]
\[ \delta \epsilon \phi^I = -\epsilon^* \psi^I, \quad \delta \epsilon \psi^I = \gamma^\mu \epsilon D_\mu \phi^I, \] (3.4)

where \( \epsilon \) is a spinor in the same representation as \( \lambda \). Since we later want to identify the presence of a C-field with certain deformations of SYM that leave supersymmetry unbroken, we have written the transformation of the adjoint spinor using \( W^A \), an imaginary quaternion (i.e., transforming in (1, 3, 1)) in the adjoint of the gauge group. In the undeformed case, \( W^A = W^A_0 = \frac{1}{2} (g^A)_{IJ} \phi^I \phi^*_J \). Note that the hypermultiplet potential is the square of \( W \), \( V(\phi) = \frac{1}{2} W^A W^*_A \). The deformations will be encoded in the form of \( W \).

The most convenient way of checking supersymmetry is to note that \( W \) is contained in the same supermultiplet as the hypermultiplet gauge current:

\[ \delta \epsilon W^A = -\text{Im}(\epsilon^* \mu^A), \]
\[ \delta \epsilon \mu^A = J^{A\mu} \gamma^\mu \epsilon + \gamma^\mu \epsilon D_\mu W^A, \] (3.5)

where \( \mu^A = \mu^A_0 = (g^A)_{IJ} \psi^I \phi^*_J \), \( J^{A\mu} = \frac{1}{2} (g^A)_{IJ} (D_\mu \phi^I \phi^*_J - \phi^I D_\mu \phi^*_J - \psi^I \gamma^\mu \psi^*_J). \)

Before turning to the derivation of case A from the supermembrane, let us discuss the advantages and limitations of the two approaches and some aspects of their physical content. Both cases are defined as dimensional reductions of \( D = 6 \) SYM with matter. The expression for the potential is a sum of positive semidefinite terms, so the Higgs branch is determined by \( W = 0 \). In light of the correspondence with five-dimensional physics mentioned above, it is interesting to investigate the geometry of the Higgs branch. The low-energy limit of adiabatic motion on the Higgs branch is also the situation when bulk excitations (gravity) decouple. Counting the dimension of the Higgs branch as \#(scalar matter fields) - \#W - dim(gauge group), one gets in case A: \( 4N^2 + 4kN - 3N^2 - N^2 = 4kN \), and in case B: \( 4 \frac{N(N+1)}{2} - 3 \frac{N(N-1)}{2} - \frac{N(N-1)}{2} = 4N \). Closer investigation reveals that the spaces agree for \( k = 1 \), and that the Higgs branch then is \( \mathbb{R}^4 \times (\mathbb{R}^{(N-1)}/P_N) \), interpreted as the space of loci of N indistinguishable partons/D0-branes. This is a flat hyper-Kähler space with conical singularities where partons coincide, which is where the Higgs branch intersects the Coulomb branch.
There is an index theorem [27] stating that the matrix theory has a unique super-symmetric ground state. The eight fermion zero modes lie in the representation \((4,1,2)\), so the ground state is the breaking to \(\text{SO}(5)\times\text{SU}(2)\) of an \(\text{SO}(8)\) spinor \(8_s \oplus 8_c\) when the vector decomposes as \(8_v \to (4,2)\) (the “Hopf breaking”). Then \(8_s \to (1,3) \oplus (5,1)\) and \(8_c \to (4,2)\), giving the bosonic and fermionic fields of the selfdual \(D = 6\) tensor multiplet in the light-cone gauge.

Note that approach \(B\) does not seem to accommodate multiple \(M_5\)-branes in a natural way. On the other hand, approach \(A\), as we will see, is less adaptable to incorporate the stringy nature of the membrane boundary. This is connected to the way it is derived from the membrane below; no boundary conditions are solved, the nature of the boundary is rather point-like. It is also unclear how \(A\) generalises to separated \(M_5\)-branes. Concerning the incorporation of a non-vanishing \(C\)-field (following section), approach \(A\) has the advantage of being more or less directly applicable, while approach \(B\) encounters problems, due to the difficulty (impossibility?) of solving the boundary conditions in the presence of a \(C\)-field.

Let us sketch briefly how case \(A\) is derived as a consistent truncation from the supermembrane. As we already mentioned, no boundary conditions are solved before performing the matrix truncation. Instead we introduce the “boundary” through the truncated \(\delta\)-function \(\Delta \equiv \sum_{n=0}^{N-1} e^{\imath n\varrho}\) (we consider a boundary located at \(\varrho = 0\), where \(\varrho\) for simplicity is a coordinate on a torus). Due to the identities \(\Delta^2 = N\Delta\) and \(\Delta \ast f \ast \Delta = 0\), left and right star multiplication with \(\Delta/\sqrt{N}\) projects on two “boundary representations” \(N\) and \(\bar{N}\) with opposite \(U(1)\) charge under adjoint action of \(\Delta\), e.g., \([\Delta, \Delta \ast f] = \Delta \ast \Delta \ast f - \Delta \ast f \ast \Delta = N\Delta \ast f\). Introduction of the “boundary” breaks \(su(N)\) to \(su(N-1) \oplus u(1)\). Higher rank \(\delta\)-functions (sums of \(\Delta\)'s with \(\varrho\) shifted by \(\frac{2\pi}{N}\) times an integer) gives \(su(N-k) \oplus su(k) \oplus u(1)\).

Let us also show how approach \(A\) generalises to a situation where the membrane is stretched between two separated parallel \(M_5\)-branes (separation \(L/2\)) or where the membrane is wound on a non-contractible circle (length \(L\)) [14]. The mode expansion of a coordinate of a cylindrical membrane in the separation direction then contains a linear term in addition to the oscillators: \(Y(\sigma, \varrho) = \frac{L}{2\pi} \varrho + \frac{1}{\pi} \sum_{n=-\infty}^{\infty} \sum_{m=1}^{\infty} y_{nm} e^{\imath m\sigma} \sin m\varrho\). The star-adjoint action of \(\frac{\varrho}{2\pi}\) is identical to \(-\frac{i}{N} \frac{\partial}{\partial \sigma} \cdot \frac{\varrho}{2\pi}\) is an outer derivation on the algebra of functions, and its presence means that it is not consistent to truncate in the \(\sigma\)-direction. Truncating in the \(\varrho\)-direction only leads to an affine \(SO(N)\) algebra. The matrix theory is an “affine matrix theory”, or a matrix string theory, which is the dimensional reduction to \(D = 2\) of a \(D = 6\) SYM theory with a hypermultiplet in the symmetric representation.
Note that the coupling constant is relevant, since there is a dimensionless quotient between the eleven-dimensional Planck length and the brane separation.

4. Non-vanishing $C$-field

We now turn to the situation where there is a non-vanishing $C$-field on the $M_5$-brane (the gauge invariant statement is in terms of the selfdual 3-form field strength $F_{(3)} = dB_{(2)} - C_{(3)}$ on the brane). In the process of choosing a light-cone gauge for the membrane, the same is done for the $C$-field. We choose $C_{-ij} = 0$. Then a selfdual $C$ falls in either of the two classes, modulo choice of frame:

1. “Light-like”: $C_{ijk} = 0$, $C_{+-i} = 0$, $C_{+ij}$ selfdual in four dimensions. The transverse rotation are broken as $so(4) \approx su(2) \oplus su(2) \rightarrow su(2) \oplus u(1)$.

2. “Space-like”: $C_{+ij} = 0$, $C_i = \varepsilon_{ijkl} C_{-l}$. Transverse rotations broken as $so(4) \approx su(2) \oplus su(2) \rightarrow (su(2))^\text{diag}$.

We now have to include the Wess–Zumino term of eq. (2.1) in the canonical analysis. The light-cone membrane hamiltonian becomes

$$H \equiv \frac{p^+p^-}{A} = \int d^2\xi \left[ \frac{1}{2} \Pi_I \Pi_I + \frac{T^2}{4} \{X^I, X^J\}\{X^I, X^J\} - \frac{T^2}{2} C_{+JK} C_{+L} \{X^I, X^J\}\{X^K, X^L\} \right.\left. - \frac{E^+T_A}{T} C_{+IJ} \{X^I, X^J\} \right] ,$$

where $\Pi_I (= \dot{X}_I) = P_I - \frac{T}{2} C_{IJK} \{X^J, X^K\} - T C_{+IJ} \{X_I, X^J\}$. In order to identify the connection with SYM, it is useful to form the lagrangian

$$L = \int d^2\xi \left[ \frac{1}{2} \dot{X}^I \dot{X}_I + \frac{T}{2} C_{IJK} \dot{X}^I \{X^J, X^K\} + T C_{+IJ} \dot{X}_I \{X^J, X^K\} - V(X) \right]$$

Due to the difficulties with solving the non-linear boundary conditions in the presence of a $C$-field, we choose to work in the approach A. The light-like case is much simpler, and already well known (although not, to our knowledge, derived from the membrane). There, the last line in eq. (4.1) represents the only deformation. We note that in the membrane hamiltonian a term $\int d\sigma d\varrho \{A, B\} = \int_{\varrho=0}^{\varrho=\pi} d\varrho \partial_\varrho B$ is a cocycle that is not well defined in the matrix truncation (since it is defined using the derivation $\varrho$). Any boundary term should be represented by a cocycle, defined by a derivation $\partial$ as $\text{tr}(A[\partial, B])$. Since a finite-dimensional Lie algebra only has inner derivations, one may be lead to conclude that it is necessary to use the affine matrix theory mentioned earlier.
This is however not true. The relevant derivation is the truncated $\delta$-function $[\Delta, \cdot]$, which is inner, so that the cocycle $\text{tr}(A, [\Delta, B])$ is exact in the space of functions. We get two equivalent pictures, one with a deformed algebra $[A, B]_k = [A, B] + k \text{tr}(A, [\Delta, B])$, and one with an undeformed algebra (obtained from the deformed generators by a redefinition containing $\Delta$) and a modified trace involving $\text{tr}\Delta \neq 0$. This gives a coupling containing the boundary representations $N$ and $\bar{N}$. It amounts to the introduction in the SYM theory of a Fayet–Iliopoulos term [23] by $W^A = W_A^0 + \zeta^A$, $\zeta$ being a fixed vector in the $U(1)$ direction defined by $\Delta$. It breaks the rotational $so(4)$ symmetry to $su(2) \oplus u(1)$ and leaves supersymmetry unbroken. Its effect is to resolve the singularities of the Higgs branch.

Turning to space-like $C$-field, we use the selfduality condition on $C$ to rewrite the terms in the lagrangian (4.2) linear in time derivatives as $\frac{1}{2} C_{-4J} (2 \dot{X}^I \{X^I, X^J\} + \varepsilon_{KLMJ} \dot{X}^K \{X^L, X^M\})$. Choosing a basis where $C_{-4} = \gamma = C_{123}$ and splitting quaternions in real and imaginary parts with $X^4 = \text{Re} X$, this can be rewritten as proportional to $f^{\mathcal{A} \mathcal{B} \mathcal{C}} \text{Re}(X^{\mathcal{A}} X^{\mathcal{B}} X^{\mathcal{C}} \star \phi)$, where indices $\mathcal{A}, \mathcal{B}, \ldots$ enumerate the truncated basis of functions. The only contribution from this term which is not a total derivative comes from the cocycle mentioned earlier, and the relevant part is then proportional to $\text{tr}(\dot{\phi} W_0)$, which leads to the conclusion that space-like $C$-field corresponds to a deformation of the SYM theory given by

$$W = W_0 + \gamma \text{Im} \phi, \quad \mu = \mu_0 + \gamma \psi,$$

where $W$ is a proper deformation, $\mu$ is a gauge-field deformation, and the deformations take values in $u(1)$. Of course, also the potential terms have to be matched against the SYM theory. It is straightforward to show, using the supersymmetry transformations of eq. (3.4), that this deformation preserves supersymmetry. The details of this are left for a future publication [28], where a fuller account will be given.

5. Conclusions

We have reviewed and constructed matrix theories describing situations where supermembranes end on M-theory 5-branes. Special emphasis has been put on non-vanishing $C$-field, which is also where the new results are found.

There are some potential applications of the results, that will be investigated in a future publication. One is to obtain the decoupling from gravity in the limit of maximal $C$-field, the OM limit. For any value of the $C$-field, we should be able to use our formulation

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† We use a linear self-duality condition, although the self-duality on an M5-brane should really be non-linear. It is not obvious to us why a linear relation seems to produce the right result.
to derive the open membrane metric, which should arise naturally after certain rescalings in the process of matching the truncated membrane hamiltonian to the SYM one.

One of our motivations for initiating this work was the prospect of treating membrane boundary conditions in the presence of non-vanishing $C$. In order for this to work, and to get information of the generalisation of the string end-point non-commutativity to membrane end-strings, one would need to find a generalisation of the approach B described above, so that the string nature of the boundary is preserved. We have not been able to do this. An intriguing observation is that there are two inequivalent cocycles extending an $su(N)$ loop algebra to an affine algebra—the untwisted and the twisted one. The zero-modes of the twisted affine algebra form an $so(N)$ algebra, which certainly indicates a connection to approach B. Further investigations along this line of thought might provide interesting results.

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