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D2 to M2 procedure for D2-brane DBI effective action

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Abstract

We apply the procedure that was suggested in [B. Ezhuthachan, S. Mukhi, C. Papageorgakis, arXiv: 0806.1639] to the case of abelian D2-brane Dirac–Born–Infeld effective action and discuss its limitation. Then we suggest an alternative form of this procedure that is based on an existence of interpolating action proposed in [T. Ortin, hep-th/9707113, Y. Lozano, hep-th/9707011].

Keywords: D-branes; M-branes

1. Introduction

It was proposed by Bagger and Lambert in collection of very nice papers [1–3] and independently by Gustavsson in [4][1] following earlier works [5,6] that a certain class of $N = 8$ super-conformal theories in three dimensions are potential candidates for the world-volume description of multiple M2-branes in M-theory. These constructions are based on introducing of an algebraic structure known as Lie 3-algebra that is needed for closure of supersymmetry algebra. The metric versions of the above theories fall into two classes that depend on whether the invariant bilinear form in 3-algebra space is positive definite or indefinite. The original theories proposed by Bagger–Lambert are Euclidean theories with positive definite bilinear form while more recent proposals [24,25] contain bilinear form that is indefinite and these Lie 3-algebras are known as Lorentzian 3-algebras.

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[1] For related works, see [7–42,46–51,56,57,59,61,67,68,70,75–83,87,89–91,100,101,103]. For study of supergravity duals of these theories, see [43–45,52–55,58,60,62–66,69,71–74,84–86,88,102].

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It was claimed that the Lorentzian 3-algebra theories capture the low-energy world-volume dynamics of multiple parallel M2-branes. This model has the required classical symmetries, but has several unresolved problems. In particular, the classical theory has ghosts, $X_{\pm}$. Moreover, the ghost-free formulation seems directly equivalent to the non-conformal D2-brane theory as was argued in [39,42,46]. Explicitly, it was shown very clearly in [46] how it is possible-starting from $N=8$ SYM—systematically and uniquely recovery the theory [39,42].

Since the analysis presented in [46] is very nice and interesting it certainly deserves further study. In fact, since $(2+1)$-dimensional $N=8$ SYM theory describes low-energy dynamics of N D2-branes one can ask the question whether it is possible to extend this analysis [46][2] when we take non-linear corrections into account. As the first step in this direction we try to apply EMP procedure to the case of single Dirac–Born–Infeld (DBI) action for D2-brane.

We start our analysis with the remarkable form of $(2+1)$-dimensional action that was proposed long ago in [92,93]. This action is an interpolating action that—after appropriate integration of some world-volume fields—either describes D2-brane DBI effective action in massive Type IIA supergravity or the directly dimensional reduced gauged M2-brane action. We show that in linearized level this action is equivalent to the abelian form of the action given in [46] and hence can be considered as the starting point for non-linear generalization of EMP procedure. On the other hand, we argue that naive application of EMP procedure in this action leads to a puzzle. Explicitly, we argue that there is a unique ground state of this new action with infinite coupling constant. This is different from we would expect since M2 to D2-reduction is based on a presumption that the vacuum expectation value of $\langle X_+\rangle$ can take arbitrary constant value. Equivalently, we would expect an infinite number of ground states that differ by vacuum expectation values of X_+.

In order to resolve this problem we suggest that the natural object for the definition of non-linear EMP procedure is gauged M2-brane action. More precisely, it is well known that in the case of IIA supergravity it is possible to introduce non-zero cosmological constant proportional to m^2 with m a mass parameter [94]. Such backgrounds are essential for the existence of D8-branes whose charge is proportional to m [95]. The action for massive 11-dimensional theory has the same contain as the massless one[3]

$$\hat{g}_{MN}, \quad \hat{C}_{MNK}, \quad M,N,K=0,\dots,11. \tag{1.1}$$

The action for these fields is manifestly 11-dimensional Lorentz covariant but it does not correspond to a proper 11-dimensional theory because, in order to write down the action, we need to introduce an auxiliary non-dynamic vector field $\hat{k}^M$ such that the Lie derivatives of the metric and 3-form potential with respect to it are zero:

$$\mathcal{L}_{\hat{k}}\hat{g}_{MN}=0, \qquad \mathcal{L}_{\hat{k}}\hat{C}_{MNL}=0. \tag{1.2}$$

An existence of this Killing vector is crucial for definition of massive M2-brane. In fact, the world-volume theory of massive branes[4] was extensively studied in the past, for example [96–99]. These actions have as a common characteristic that they are gauged sigma models. The gauged isometry is the same as an isometry that is needed in order to define the massive

[2] In what follows we call this analysis as EMP procedure.

[3] We use the following notation for the hats. Hats on target space fields indicate that they are 11-dimensional.

[4] These branes-that propagate in the background with non-zero cosmological constant-are called as “massive branes” as opposed to branes that propagate in a background with zero mass parameter. It is clear that all these branes are massive in the sense that their physical mass is non-zero.

11-dimensional supergravity theory. For example, the original un-gauged M2-brane action is the same object as in the massless theory, i.e. the corresponding massless M2-brane.

Let us again return to generalized EMP procedure. We argue that it can be naturally applied for gauged M2-brane action. As opposite to the original EMP procedure, where the Yang–Mills coupling constant vector g^I_{YM} is replaced with a dynamical field X^I_+ we replace the constant Killing vector $\hat{k}^M$ with dynamical field $\hat{X}^M_+$. Then we can easily find manifestly covariant form of the generalized action with infinite number of ground states that differ by vacuum expectation values of $\hat{X}^M_+$.

It is remarkable that the gauged isometry that appears in massive M2-brane action is related to the gauge symmetry introduced in [46]. We hope that this observation will allow to find new geometrical interpretations of gauge symmetries that were introduced in [39,42,46].

The organization of this paper is as follows. In the Section 2 we introduce the interpolating D2-brane action and we argue that after appropriate redefinition of world-volume fields it agrees with the abelian version of D2-brane action introduced in [46]. In Section 3 we apply EMP prescription for gauged M2-brane action and we find covariant and non-linear version of M2-brane action that has all desired properties. In Section 4 we outline our results and suggest possible extension of our work. Finally, in Appendix A we explicitly show that the dimensional reduction in gauged M2-brane action leads to the interpolating action introduced in Section 2.

2. D2-brane action

We start with the action that was proposed in [92,93]

$$\begin{aligned}
&S\big[X^m, X, V_\mu, B_\mu\big] \\
&\quad = -\tau_{M2}\int d^3\xi\, e^{-\Phi}\sqrt{-\det\big[g_{\mu\nu}+e^{2\Phi}F_\mu F_\nu\big]} \\
&\qquad + \frac{\tau_{M2}}{3!}\int d^3\xi\,\epsilon^{\mu\nu\rho}\big[C^{(3)}_{\mu\nu\rho}+6\pi\alpha' D_\mu X\mathcal{F}_{\nu\rho}+6m(\pi\alpha')^2 V_\mu\partial_\nu V_\rho\big],
\end{aligned} \tag{2.1}$$

where

$$\begin{aligned}
&F_\mu = D_\mu X + C^{(1)}_\mu, \\
&D_\mu X = \partial_\mu X + B_\mu, \\
&\mathcal{F}_{\mu\nu} = \partial_\mu V_\nu - \partial_\nu V_\mu - \frac{1}{2\pi\alpha'}b_{\mu\nu},
\end{aligned} \tag{2.2}$$

and where

$$\begin{aligned}
&g_{\mu\nu} = \partial_\mu X^m\partial_\nu X^n g_{mn}, \qquad b_{\mu\nu} = b_{mn}\partial_\mu X^m\partial_\nu X^n, \\
&C^{(3)}_{\mu\nu\rho} = C_{mnk}\partial_\mu X^m\partial_\nu X^n\partial_\rho X^k, \qquad C^{(1)}_\mu = C^{(1)}_m\partial_\mu X^m,,
\end{aligned} \tag{2.3}$$

where g_{mn}, b_{mn} are space–time metric and NS two form field respectively and where $C^{(3)}_{mnk}$, $C^{(1)}_m$ are Ramond–Ramond three and one forms respectively. Further, X^m, $m,n=0,\dots,9$, are world-volume modes that describe embedding of D2-brane in the target space–time. Finally, τ_{M2} is D2-brane tension defined as $\tau_{\mathrm{M2}} = 1/l_s^3$.

The action (2.1) contains extra fields as opposite to usual DBI action for D2-brane. Firstly, if we integrate out B_μ we obtain

$$-\frac{e^{\Phi}\sqrt{-g}g^{\mu\nu}F_\nu}{\sqrt{1+e^{2\Phi}g^{\mu\nu}F_\mu F_\nu}}+\pi\alpha'\epsilon^{\mu\nu\rho}\mathcal{F}_{\nu\rho}=0, \tag{2.4}$$

where we used the fact that

$$\sqrt{-\det\left[g_{\mu\nu}+e^{2\Phi}F_\mu F_\nu\right]}=\sqrt{-\det g}\sqrt{1+e^{2\Phi}g^{\mu\nu}F_\mu F_\nu}. \tag{2.5}$$

Then if we insert (2.4) into (2.1) we obtain an action in the form

$$\begin{aligned}S={}&-\tau_{\mathrm{M2}}\int d^3\xi\, e^{-\Phi}\sqrt{-\det[g_{\mu\nu}+2\pi\alpha'\mathcal{F}_{\mu\nu}]}\\&+\frac{\tau_{\mathrm{M2}}}{3!}\int d^3\xi\,\epsilon^{\mu\nu\rho}\left(C^{(3)}_{\mu\nu\rho}-6\pi\alpha' C^{(1)}_\mu\mathcal{F}_{\nu\rho}+6m(\pi\alpha')^2V_\mu\partial_\nu V_\rho\right),\end{aligned} \tag{2.6}$$

that is standard form of D2-brane in massive Type IIA background and that reduces to the massless Type IIA background when $m=0$.

In order to see that the action (2.6) is related to abelian reduction of the action given in [46] we take following background:

$$g_{mn}=\eta_{mn},\qquad \Phi=\Phi_0=\text{const},\qquad C^{(1)}_m=C^{(3)}_{mnk}=0. \tag{2.7}$$

Further, let us impose static gauge

$$X^\mu=\xi^\mu,\quad \mu=0,1,2, \tag{2.8}$$

so that

$$g_{\mu\nu}=\eta_{\mu\nu}+\delta_{ij}\partial_\mu X^i\partial_\nu X^j,\quad i,j=3,\ldots,9. \tag{2.9}$$

Then in the quadratic approximation the action (2.6) takes the form

$$\begin{aligned}S\left[X^m,X,V_\mu,B_\mu\right]={}&-\tau_{\mathrm{M2}}\int d^3\xi\, e^{-\Phi_0}\sqrt{-\eta}\\&-\tau_{\mathrm{M2}}\int d^3\xi\,\sqrt{-\eta}\left[\frac{1}{2}e^{-\Phi_0}\eta^{\mu\nu}\delta_{ij}\partial_\mu X^i\partial_\nu X^j+\frac{e^{\Phi_0}}{2}\eta^{\mu\nu}F_\mu F_\nu\right]\\&+\int d^3\xi\,\epsilon^{\mu\nu\rho}(\pi\alpha'\tau_{\mathrm{M2}})D_\mu X\mathcal{F}_{\nu\rho}.\end{aligned} \tag{2.10}$$

As the next step we introduce the gauge theory coupling constant through the standard relations

$$e^{-\Phi_0}l_s^4\tau_{\mathrm{M2}}=\frac{1}{g^2_{\mathrm{YM}}}\quad\left(\tau_{\mathrm{M2}}=\frac{1}{l_s^3},\ 2\pi\alpha'=l_s\right), \tag{2.11}$$

so that after rescaling

$$\begin{aligned}&\sqrt{\tau_{\mathrm{M2}}}e^{\Phi_0/2}X=\tilde{X},\qquad \sqrt{\tau_{\mathrm{M2}}}e^{-\Phi_0/2}X^i=\tilde{X}^i,\\&\frac{1}{l_s^{5/2}}B_\mu=\tilde{B}_\mu,\qquad l_s^{3/2}\mathcal{F}_{\mu\nu}=\tilde{\mathcal{F}}_{\mu\nu}\end{aligned} \tag{2.12}$$

the action (2.10) takes the form

$$\begin{aligned} & S\big[\tilde{X}^m, \tilde{X}, \tilde{V}_\mu, \tilde{B}\big] \\ &\quad = -\frac{1}{g_{\mathrm{YM}}^2 l_s^2} \int d^3\xi \sqrt{-\eta} \\ &\qquad - \int d^3\xi \sqrt{-\eta} \left[\frac{1}{2}\eta^{\mu\nu}\delta_{ij}\partial_\mu \tilde{X}^i \partial_\nu \tilde{X}^j + \frac{1}{2}\eta^{\mu\nu}(\partial_\mu \tilde{X} + g_{\mathrm{YM}}\tilde{B}_\mu)(\partial_\nu \tilde{X} + g_{\mathrm{YM}}\tilde{B}_\nu)\right] \\ &\qquad + \int d^3\xi\, \epsilon^{\mu\nu\rho}\left[\frac{1}{2}\tilde{B}_\mu \mathcal{F}_{\nu\rho} + \frac{1}{2 l_s^{3/2} g_{\mathrm{YM}}}\partial_\mu \tilde{X}\tilde{\mathcal{F}}_{\nu\rho}\right] \end{aligned} \tag{2.13}$$

that has the same form as the abelian form of the action given in [46].[5] Motivated by this result we perform the rescaling (2.12) in the action (2.6) and we obtain

$$\begin{aligned} & S\big[\tilde{X}^m, \tilde{X}, \tilde{V}_\mu, \tilde{B}_\mu\big] = -\int d^3\xi \sqrt{-\det \mathbf{A}_{\mu\nu}} + \int d^3\xi\, \epsilon^{\mu\nu\rho}\frac{1}{2}\tilde{B}_\mu \tilde{\mathcal{F}}_{\nu\rho}, \\ & \mathbf{A}_{\mu\nu} = \frac{1}{l_s^{8/3} g_{\mathrm{YM}}^{4/3}}\eta_{\mu\nu} + g_{\mathrm{YM}}^{2/3} l_s^{4/3}\partial_\mu \tilde{X}^i \partial_\nu \tilde{X}^j \delta_{ij} \\ &\qquad + l_s^{4/3} g_{\mathrm{YM}}^{2/3}(\partial_\mu \tilde{X} + g_{\mathrm{YM}}\tilde{B}_\mu)(\partial_\nu \tilde{X} + g_{\mathrm{YM}}\tilde{B}_\nu), \end{aligned} \tag{2.14}$$

where we ignored term that contributes to the action as total derivative. Now we are ready to apply EMP procedure for (2.14). We introduce 8-dimensional vector g_{YM}^I as $g_{\mathrm{YM}}^I = (\overbrace{0, \ldots, 0}^{7}, g_{\mathrm{YM}})$, $I = 1, \ldots, 8$, and "covariant derivative $\tilde{D}$"

$$\tilde{D}_\mu \tilde{X}^i = \partial_\mu \tilde{X}^i + g_{\mathrm{YM}}^i \tilde{B}_\mu, \qquad \tilde{D}_\mu \tilde{X} = \partial_\mu \tilde{X} + g_{\mathrm{YM}}\tilde{B}_\mu. \tag{2.15}$$

Further, we rewrite g_{YM}^2 in manifest $SO(8)$ covariant manner as $g_{\mathrm{YM}}^2 = g_{\mathrm{YM}}^I g_{\mathrm{YM}}^J \delta_{IJ} = |g_{\mathrm{YM}}|^2$ and then we replace vector g_{YM}^I with dynamical field X_+^I so that the action (2.14) takes the form

$$\begin{aligned} & S\big[\tilde{X}^I, \tilde{X}_+^I, \tilde{V}_\mu, \tilde{B}_\mu, C_I^\mu\big] = -\int d^3\xi \left(\sqrt{-\det \mathbf{A}_{\mu\nu}} + \epsilon^{\mu\nu\rho}\frac{1}{2}\tilde{B}_\mu \tilde{\mathcal{F}}_{\nu\rho} + C_I^\mu \partial_\mu \tilde{X}_+^I\right), \\ & \mathbf{A}_{\mu\nu} = \frac{1}{l_s^{8/3}\big(X_+^I X_+^J \delta_{IJ}\big)^{2/3}}\eta_{\mu\nu} + \big(X_+^I X_+^J \delta_{IJ}\big)^{1/3} l_s^{4/3} \tilde{D}_\mu \tilde{X}^I \tilde{D}_\nu \tilde{X}^J \delta_{IJ}, \end{aligned} \tag{2.16}$$

where we introduced auxiliary field C_I^μ that renders $\tilde{X}_+^I$ non-dynamical.

Let us now analyze some properties of the action (2.16). We are mainly interested in the study of the ground state of this theory that has to solve the equations of motion that follow from the action (2.16). We presume that the ground state is characterized by following configuration of the world-volume fields

$$\tilde{X}_+^8 = v = \text{const}, \qquad \tilde{D}_\mu \tilde{X}^I = 0, \qquad \tilde{B}_\mu = 0, \qquad \tilde{\mathcal{F}}_{\mu\nu} = 0. \tag{2.17}$$

Firstly, the equation of motion for C_I^μ takes the form

$$\partial_\mu \tilde{X}_+^I = 0 \tag{2.18}$$

[5] This is true up to total derivative term since $\int d^3\xi\, \epsilon^{\mu\nu\rho}\partial_\mu \tilde{X}\tilde{\mathcal{F}}_{\nu\rho} = \int d^3\xi\, \partial_\mu[\epsilon^{\mu\nu\rho}\tilde{X}\tilde{\mathcal{F}}_{\nu\rho}] - \int d^3\xi\, \tilde{X}\partial_\mu(\epsilon^{\mu\nu\rho}\tilde{\mathcal{F}}_{\nu\rho}) = \int d^3\xi\, \partial(\cdots)$.

that is clearly obeyed by the ansatz (2.17). On the other hand the equation of motion for X^I_+ takes the form

$$-\frac{1}{2}\frac{\delta \mathbf{A}_{\mu\nu}}{\delta X^I_+}\left(\mathbf{A}^{-1}\right)^{\nu\mu}\sqrt{-\det\mathbf{A}}-\epsilon^{\mu\nu\rho}\frac{1}{2l_s^{1/2}(X^I_+X^J_+\delta_{IJ})^{3/2}}\partial_\mu\tilde{X}\tilde{\mathcal{F}}_{\nu\rho}=0. \tag{2.19}$$

Since for the ansatz (2.17) the matrix $\mathbf{A}_{\mu\nu}$ is equal to

$$\mathbf{A}_{\mu\nu}=\frac{1}{l_s^{8/3}v^{4/3}}\eta_{\mu\nu}, \tag{2.20}$$

Eq. (2.19) implies

$$\frac{1}{v^3}=0. \tag{2.21}$$

In other words, the ground state corresponds to the point $v\to\infty$ that implies that there is unique ground state of the theory. As we argued in introduction this is not the same what we want since we would like to have a theory with infinite number of ground states that differ by vacuum expectation values of $\tilde{X}_+$. In order to find solution of this problem we suggest an alternative procedure how to introduce X^I_+ as a new dynamical variable. In the next section we present such an alternative procedure that is based on the fact that the action (2.6) can be considered as dimensional reduction of massive M2-brane.

3. Gauged theory for M2-brane

Let us again consider the action (2.6) and determine the equations of motion for V_μ

$$\pi m\alpha'\epsilon^{\mu\nu\rho}(\partial_\nu V_\rho-\partial_\rho V_\mu)+\epsilon^{\mu\nu\rho}(\partial_\nu B_\rho-\partial_\rho B_\nu)=0. \tag{3.1}$$

Inserting (3.1) back to (2.6) we obtain the action in the form

$$\begin{aligned}S={}&-\tau_{\mathrm{M2}}\int d^3\xi\sqrt{-\det\left(e^{-2/3\Phi}g_{\mu\nu}+e^{4/3\Phi}F_\mu F_\nu\right)}\\&+\frac{\tau_{\mathrm{M2}}}{6}\int d^3\xi\,\epsilon^{\mu\nu\rho}\left(C^{(3)}_{\mu\nu\rho}-3D_\mu XB^{(1)}_{\nu\rho}+\frac{6}{m}B_\mu\partial_\nu B_\rho\right).\end{aligned} \tag{3.2}$$

As was shown in [93] (and reviewed in Appendix A) this action is very close to the action that one gets by direct dimensional reduction of the massive M2-brane that is also known as gauged M2-brane action. This action can be defined in the background with Killing vector isometry $\hat{k}^M(\hat{X})$. Then the gauged M2-bane action takes the form [93]

$$\begin{aligned}S={}&-\tau_{\mathrm{M2}}\int d^3\xi\sqrt{-\det D_\mu\hat{X}^MD_\nu\hat{X}^N\hat{g}_{MN}}\\&+\tau_{\mathrm{M2}}\int d^3\xi\,\epsilon^{\mu\nu\rho}\left[D_\mu\hat{X}^MD_\nu\hat{X}^ND_\rho\hat{X}^K\hat{C}_{MNK}-\frac{6}{m}B_\mu\partial_\nu B_\rho\right],\end{aligned} \tag{3.3}$$

where the covariant derivative D_μ is defined as

$$D_\mu\hat{X}^M=\partial_\mu\hat{X}^M+B_\mu\hat{k}^M(\hat{X}), \tag{3.4}$$

where B_μ is world-volume gauge field related to the Killing gauge isometry. To clarify meaning of this gauged form of the action let us consider following transformation:

$$\delta_\eta\hat{X}^M(\xi)=\hat{X}'^M(\xi)-\hat{X}^M(\xi)=\eta(\xi)\hat{k}^M(\hat{X}), \tag{3.5}$$

where $\eta(\xi)$ is a parameter of gauge transformations. This transformation immediately implies following transformation rules of background fields:

$$\delta_\eta \hat{g}_{MN} = \eta \hat{k}^K \partial_K \hat{g}_{MN}, \qquad \delta_\eta \hat{C}_{KMN} = \eta \hat{k}^L \partial_L \hat{C}_{KMN}, \qquad \delta_\eta \hat{k}^K = \eta \hat{k}^L \partial_L \hat{k}^K, \tag{3.6}$$

and transformation of covariant derivative

$$\delta_\eta D_\mu \hat{X}^M = \eta D_\mu \hat{X}^L \partial_L \hat{k}^M, \tag{3.7}$$

where we postulate following transformation rule for gauge field B_μ:

$$\delta B_\mu = -\partial_\mu \eta. \tag{3.8}$$

Then

$$\delta\left(D_\mu \hat{X}^M D_\nu \hat{X}^N \hat{g}_{MN}\right) = \eta D_\mu \hat{X}^M \left(\partial_M \hat{k}^L \hat{g}_{LN} + \hat{g}_{ML} \partial_N \hat{k}^L + \partial_L \hat{g}_{MN}\right) D_\nu \hat{X}^N = 0 \tag{3.9}$$

since

$$\mathcal{L}_{\hat{k}} \hat{g}_{MN} = 0. \tag{3.10}$$

In the same way we obtain that

$$\delta_\eta\left(D_\mu \hat{X}^M D_\nu \hat{X}^N D_\rho \hat{X}^K \hat{C}_{MNK}\right) = 0, \qquad \mathcal{L}_{\hat{k}} \hat{C}_{MNK} = 0, \tag{3.11}$$

hence we see that the action is invariant under transformations (3.5) and (3.6).

Having clarified the fact that the D2-brane action (2.6) is related to the gauged M2-brane action we now introduce modified EMP procedure to the action (3.3). As the first step in our construction we will presume an existence of adapted system of coordinates where $\hat{k}^M = \text{const}$. This is always possible to achieve in flat background $\hat{g}_{MN} = \eta_{MN}$, $\hat{C}_{MNK} = 0$. Further, in analogy with EMP prescription, we replace constant $\hat{k}^M$ with dynamical field $\hat{X}^M_+$ and add to the action term $\frac{1}{2} C^\mu_M \partial_\mu \hat{X}^M_+$ to render this field non-dynamical. Further, we rewrite the Wess–Zumino term in (3.3) as

$$\pi\alpha' \epsilon^{\mu\nu\rho} B_\mu \mathcal{F}_{\nu\rho} + (\pi\alpha')^2 m \epsilon^{\mu\nu\rho} V_\mu \partial_\nu V_\rho. \tag{3.12}$$

In fact it is easy to see that now the equation of motion for V_μ that follow from (3.12) implies

$$\frac{1}{m\pi\alpha'}(\partial_\mu B_\nu - \partial_\nu B_\mu) = -(\partial_\mu V_\nu - \partial_\nu V_\mu), \tag{3.13}$$

and hence when we insert it back to (3.12) we obtain the last term in (3.3). Note also that this expression is invariant under η transformations (up to total derivative) since

$$\delta\left(\epsilon^{\mu\nu\rho} B_\mu \mathcal{F}_{\nu\rho}\right) = \epsilon^{\mu\nu\rho} \partial_\mu \eta \mathcal{F}_{\nu\rho} = -\eta \epsilon^{\mu\nu\rho} \partial_\mu \partial_\nu A_\rho + \eta \epsilon^{\mu\nu\rho} \partial_\mu \partial_\rho A_\nu = 0. \tag{3.14}$$

In summary we derive the action in the form

$$\begin{aligned} S = {} & -\tau_{\mathrm{M2}} \int d^3\xi \left[\sqrt{-\det \mathcal{G}_{\mu\nu}} + \frac{1}{2}\sqrt{-\det \mathcal{G}_{\mu\nu}}\, C^N_\nu \eta_{NM} \left(\mathcal{G}^{-1}\right)^{\nu\mu} \partial_\mu \hat{X}^M_+\right] \\ & + \frac{\tau_{\mathrm{M2}}}{3!} \int d^3\xi \epsilon^{\mu\nu\rho}\left[6\pi\alpha' B_\mu \mathcal{F}_{\nu\rho} + 6m(\pi\alpha')^2 V_\mu \partial_\nu V_\rho\right], \end{aligned} \tag{3.15}$$

where we added term $\frac{1}{2} C^M_\mu (\mathcal{G}^{-1})^{\mu\nu} \eta_{MN} \partial_\nu \hat{X}^N_+$ that renders $\hat{X}^M_+$ constant on-shell and where we also introduced "generalized metric" $\mathcal{G}_{\mu\nu}$

$$\mathcal{G}_{\mu\nu} = \left(\partial_\mu \hat{X}^M + B_\mu \hat{X}^M_+\right) \eta_{MN} \left(\partial_\nu \hat{X}^N + B_\nu \hat{X}^N_+\right). \tag{3.16}$$

Following [46] we introduce field $\hat{X}^M_-$ and add to the action an expression

$$\frac{1}{2}\sqrt{-\det\mathcal{G}}\,\partial_\mu\tilde{X}^M_-\eta_{MN}\big(\mathcal{G}^{-1}\big)^{\mu\nu}\partial_\nu\tilde{X}^N_+$$

in order the action will be invariant under additional shift symmetry

$$\delta C^M_\mu=\partial_\mu\lambda^M,\qquad \delta\hat{X}^M_-=\lambda^M. \tag{3.17}$$

Then the final form of the action takes the form

$$\begin{aligned} S={}&-\tau_{\mathrm{M2}}\int d^3\xi\,\sqrt{-\det\mathcal{G}_{\mu\nu}}\bigg(1+\frac{1}{2}\big(C^M_\mu-\partial_\mu\hat{X}^M_-\big)\eta_{MN}\big(\mathcal{G}^{-1}\big)^{\mu\nu}\partial_\nu\hat{X}^M_+\bigg)\\ &+\frac{\tau_{\mathrm{M2}}}{3!}\int d^3\xi\,\epsilon^{\mu\nu\rho}\big[6\pi\alpha' B_\mu\mathcal{F}_{\nu\rho}+6m(\pi\alpha')^2V_\mu\partial_\nu V_\rho\big]. \end{aligned} \tag{3.18}$$

Note also that in order to achieve that $\hat{X}^M_+$ is constant on-shell and that the action possesses additional shift symmetry we can consider more general form of the action

$$\begin{aligned} S={}&-\tau_{\mathrm{M2}}\int d^3\xi\,\sqrt{-\det\mathcal{G}_{\mu\nu}}\sqrt{1+\big(C^M_\mu-\partial_\mu\hat{X}^M_-\big)\eta_{MN}\big(\mathcal{G}^{-1}\big)^{\mu\nu}\partial_\nu\hat{X}^N_+}\\ &-\frac{\tau_{\mathrm{M2}}}{3!}\int d^3\xi\,\epsilon^{\mu\nu\rho}\big[6\pi\alpha' B_\mu\mathcal{F}_{\nu\rho}+6m(\pi\alpha')^2V_\mu\partial_\nu V_\rho\big] \end{aligned} \tag{3.19}$$

that can be finally written in a suggestive form as

$$\begin{aligned} &S=-\tau_{\mathrm{M2}}\int d^3\xi\,\sqrt{-\det\mathbf{A}_{\mu\nu}}-\frac{\tau_{\mathrm{M2}}}{3!}\int d^3\xi\,\epsilon^{\mu\nu\rho}\big[6\pi\alpha' B_\mu\mathcal{F}_{\nu\rho}+6m(\pi\alpha')^2V_\mu\partial_\nu V_\rho\big],\\ &\mathbf{A}_{\mu\nu}=\mathcal{G}_{\mu\nu}+\big(C^M_\mu-\partial_\mu\hat{X}^M_-\big)\eta_{MN}\partial_\nu\hat{X}^N_+. \end{aligned} \tag{3.20}$$

Let us now study properties of the action (3.20). Clearly it is invariant under shift symmetry (3.17). Further, the variation of this action with respect to C^M_μ implies

$$\eta_{NM}\partial_\nu\hat{X}^M_+\big(\mathbf{A}^{-1}\big)^{\nu\mu}\sqrt{-\det\mathbf{A}}=0 \tag{3.21}$$

that implies $\partial_\nu\hat{X}^M_+=0$. Let us again presume the ground state of the theory in the form

$$B_\mu=\hat{X}^M_-=C^M_\mu=V_\mu=0,\qquad \hat{X}^M_+=v^M. \tag{3.22}$$

It is easy to see that the equations of motion for $\hat{X}^M$, B_μ, C^M_μ and $\hat{X}^M_-$ are obeyed for this ansatz. Finally, the problematic equation of motion for $\hat{X}^M_+$ takes the form

$$\begin{aligned} &B_\mu\eta_{MN}\big(\partial_\nu\hat{X}^N+B_\nu\hat{X}^N_+\big)\big(\mathbf{A}^{-1}\big)^{\nu\mu}\sqrt{-\det\mathbf{A}}\\ &\quad+\frac{1}{2}\partial_\mu\big[\eta_{MN}\partial_\nu\hat{X}^N_-\big(\mathbf{A}^{-1}\big)^{\nu\mu}\sqrt{-\det\mathbf{A}}\,\big]=0 \end{aligned} \tag{3.23}$$

that is clearly solved by (3.22) for any v^M. Finally, let us impose the static gauge in the following form:

$$\begin{aligned} &\hat{X}^\mu=\xi^\mu,\qquad \mu,\nu=0,1,2,\\ &C^\nu_\mu=0,\qquad \hat{X}^\mu_+=\hat{X}^\mu_-=0, \end{aligned} \tag{3.24}$$

so that the matrix $\mathbf{A}_{\mu\nu}$ takes the form

$$\mathbf{A}_{\mu\nu} = \eta_{\mu\nu} + \left(\partial_\mu \hat{X}^I + B_\mu \hat{X}_+^I\right)\left(\partial_\nu \hat{X}^J + B_\nu \hat{X}_+^J\right)\delta_{IJ} + \left(C_\mu^I - \partial_\mu \hat{X}_-^I\right)\delta_{IJ}\partial_\nu \hat{X}_+^J, \quad I, J = 1, \dots, 8. \tag{3.25}$$

Then the action up to quadratic approximation can be written as

$$S = -\tau_{\mathrm{M2}} \int d^3\xi \sqrt{-\eta} - \tau_{\mathrm{M2}} \int d^3\xi \sqrt{-\det\eta}\, \mathcal{L}, \tag{3.26}$$

where the Lagrangian density takes the form

$$\mathcal{L} = \frac{1}{2}\eta^{\mu\nu}\left(\partial_\mu \hat{X}^I + B_\mu \hat{X}_+^I\right)\delta_{IJ}\left(\partial_\nu \hat{X}^J + B_\nu \hat{X}_+^J\right)\delta_{IJ} + \frac{1}{2}\eta^{\mu\nu}\left(C_\mu^I - \partial_\mu \hat{X}_-^I\right)\partial_\nu \hat{X}_+^J \delta_{IJ} - \epsilon^{\mu\nu\rho}\left(\pi\alpha' B_\mu \mathcal{F}_{\nu\rho} + m(\pi\alpha')^2 V_\mu \partial_\nu V_\rho\right), \tag{3.27}$$

that is again very close to the abelian form of the action given in [46] and provides further support of our construction.

4. Conclusion

Let us summarize our results. We studied EMP procedure for Dirac–Born–Infeld action for D2-brane and we found its limitation. Then we suggested an alternative form of this procedure that is based on a formulation of gauged M2-brane action. This fact however implies that the theory should be defined in background with non-zero mass parameter m and this observation certainly deserves better understanding and more detailed study. Further, it would be also interesting to develop BRST Hamiltonian treatment of the action (3.20) and compare it with the similar analysis that was given in [24]. Finally, it will be extremely interesting to see whether there exists an non-abelian extension of the action (3.20). We hope to return to these problems in future.

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Appendix A. Direct dimensional reduction for massive M2-brane

In this appendix we show that the gauged M2-brane action upon direct dimensional reduction in the direction X associated to the gauged isometry reduces to the action (3.2). To begin with we choose coordinates that are adapted to the isometry so that $\hat{k}^M = \delta^{Mx}$ and we split eleven coordinates $\hat{X}^M$ into the ten 10-dimensional $X^m, m = 0, \dots, 9$, and the extra scalar $\hat{X}^x \equiv X$. Using the relations between the 11-dimensional and 10-dimensional fields

$$\begin{aligned}
&\hat{g}_{xx} = e^{\frac{4}{3}\Phi}, \qquad \hat{g}_{mx} = e^{\frac{4}{3}\Phi} C_m^{(1)}, \\
&\hat{g}_{mn} = e^{-\frac{2}{3}\Phi} g_{mn} + e^{\frac{4}{3}\Phi} C_m^{(1)} C_n^{(1)}, \\
&\hat{C}_{mnk} = C_{mnk}^{(3)}, \qquad \hat{C}_{mnx} = B_{mn},
\end{aligned} \tag{A.1}$$

it is straightforward to see that

$$\mathbf{A}_{\mu\nu} = e^{-\frac{2}{3}\Phi}\partial_\mu X^m \partial_\nu X^n g_{mn} + e^{\frac{4}{3}\Phi}\left(\partial_\mu X + B_\mu + C^{(1)}_\mu\right)\left(\partial_\nu X + B_\nu + C^{(1)}_\nu\right). \tag{A.2}$$

Then if we insert (A.2) together with (A.1) into the action (3.3) we easily obtain that it reduces to the action (3.2).

References

[1] J. Bagger, N. Lambert, Modeling multiple M2's, Phys. Rev. D 75 (2007) 045020, hep-th/0611108.
[2] J. Bagger, N. Lambert, Gauge symmetry and supersymmetry of multiple M2-branes, Phys. Rev. D 77 (2008) 065008, arXiv: 0711.0955 [hep-th].
[3] J. Bagger, N. Lambert, Comments on multiple M2-branes, JHEP 0802 (2008) 105, arXiv: 0712.3738 [hep-th].
[4] A. Gustavsson, Selfdual strings and loop space Nahm equations, JHEP 0804 (2008) 083, arXiv: 0802.3456 [hep-th].
[5] A. Basu, J.A. Harvey, The M2–M5 brane system and a generalized Nahm's equation, Nucl. Phys. B 713 (2005) 136, hep-th/0412310.
[6] J.H. Schwarz, Superconformal Chern–Simons theories, JHEP 0411 (2004) 078, hep-th/0411077.
[7] S. Mukhi, C. Papageorgakis, M2 to D2, JHEP 0805 (2008) 085, arXiv: 0803.3218 [hep-th].
[8] M.A. Bandres, A.E. Lipstein, J.H. Schwarz, $N=8$ superconformal Chern–Simons theories, JHEP 0805 (2008) 025, arXiv: 0803.3242 [hep-th].
[9] D.S. Berman, L.C. Tadrowski, D.C. Thompson, Aspects of multiple membranes, Nucl. Phys. B 802 (2008) 106, arXiv: 0803.3611 [hep-th].
[10] M. Van Raamsdonk, Comments on the Bagger–Lambert theory and multiple M2-branes, JHEP 0805 (2008) 105, arXiv: 0803.3803 [hep-th].
[11] A. Morozov, On the problem of multiple M2 branes, JHEP 0805 (2008) 076, arXiv: 0804.0913 [hep-th].
[12] N. Lambert, D. Tong, Membranes on an orbifold, arXiv: 0804.1114 [hep-th].
[13] J. Distler, S. Mukhi, C. Papageorgakis, M. Van Raamsdonk, M2-branes on M-folds, JHEP 0805 (2008) 038, arXiv: 0804.1256 [hep-th].
[14] U. Gran, B.E.W. Nilsson, C. Petersson, On relating multiple M2 and D2-branes, arXiv: 0804.1784 [hep-th].
[15] P.M. Ho, R.C. Hou, Y. Matsuo, Lie 3-algebra and multiple M2-branes, JHEP 0806 (2008) 020, arXiv: 0804.2110 [hep-th].
[16] J. Gomis, A.J. Salim, F. Passerini, Matrix theory of type IIB plane wave from membranes, arXiv: 0804.2186 [hep-th].
[17] E.A. Bergshoeff, M. de Roo, O. Hohm, Multiple M2-branes and the embedding tensor, Class. Quantum Grav. 25 (2008) 142001, arXiv: 0804.2201 [hep-th].
[18] K. Hosomichi, K.M. Lee, S. Lee, Mass-deformed Bagger–Lambert theory and its BPS objects, arXiv: 0804.2519 [hep-th].
[19] G. Papadopoulos, M2-branes, 3-Lie algebras and Plucker relations, JHEP 0805 (2008) 054, arXiv: 0804.2662 [hep-th].
[20] J.P. Gauntlett, J.B. Gutowski, Constraining maximally supersymmetric membrane actions, arXiv: 0804.3078 [hep-th].
[21] G. Papadopoulos, On the structure of k-Lie algebras, Class. Quantum Grav. 25 (2008) 142002, arXiv: 0804.3567 [hep-th].
[22] P.M. Ho, Y. Matsuo, M5 from M2, JHEP 0806 (2008) 105, arXiv: 0804.3629 [hep-th].
[23] J. Gomis, G. Milanesi, J.G. Russo, Bagger–Lambert theory for general Lie algebras, JHEP 0806 (2008) 075, arXiv: 0805.1012 [hep-th].
[24] S. Benvenuti, D. Rodriguez-Gomez, E. Tonni, H. Verlinde, $N=8$ superconformal gauge theories and M2 branes, arXiv: 0805.1087 [hep-th].
[25] P.M. Ho, Y. Imamura, Y. Matsuo, M2 to D2 revisited, JHEP 0807 (2008) 003, arXiv: 0805.1202 [hep-th].
[26] A. Morozov, From simplified BLG action to the first-quantized M-theory, arXiv: 0805.1703 [hep-th].
[27] Y. Honma, S. Iso, Y. Sumitomo, S. Zhang, Janus field theories from multiple M2 branes, arXiv: 0805.1895 [hep-th].
[28] H. Fuji, S. Terashima, M. Yamazaki, A new $N=4$ membrane action via orbifold, arXiv: 0805.1997 [hep-th].
[29] P.M. Ho, Y. Imamura, Y. Matsuo, S. Shiba, M5-brane in three-form flux and multiple M2-branes, arXiv: 0805.2898 [hep-th].
[30] Y. Song, Mass deformation of the multiple M2 branes theory, arXiv: 0805.3193 [hep-th].

[31] I. Jeon, J. Kim, N. Kim, S.W. Kim, J.H. Park, Classification of the BPS states in Bagger–Lambert theory, arXiv: 0805.3236 [hep-th].
[32] M. Li, T. Wang, M2-branes coupled to antisymmetric fluxes, arXiv: 0805.3427 [hep-th].
[33] K. Hosomichi, K.M. Lee, S. Lee, S. Lee, J. Park, $N = 4$ superconformal Chern–Simons theories with hyper and twisted hyper multiplets, arXiv: 0805.3662 [hep-th].
[34] S. Banerjee, A. Sen, Interpreting the M2-brane action, arXiv: 0805.3930 [hep-th].
[35] H. Lin, Kac–Moody extensions of 3-algebras and M2-branes, arXiv: 0805.4003 [hep-th].
[36] P. De Medeiros, J. Figueroa-O'Farril, E. Mendez-Escobar, Lorentzian Lie 3-algebras and their Bagger–Lambert moduli space, arXiv: 0805.4363 [hep-th].
[37] A. Gustavsson, One-loop corrections to Bagger–Lambert theory, arXiv: 0805.4443 [hep-th].
[38] J. Figueroa-O'Farrill, Lorentzian Lie n-algebras, arXiv: 0805.4760 [math.RT].
[39] M.A. Bandres, A.E. Lipstein, J.H. Schwarz, Ghost-free superconformal action for multiple M2-branes, arXiv: 0806.0054 [hep-th].
[40] J.H. Park, C. Sochichiu, Single M5 to multiple M2: Taking off the square root of Nambu–Goto action, arXiv: 0806.0335 [hep-th].
[41] F. Passerini, M2-brane superalgebra from Bagger–Lambert theory, arXiv: 0806.0363 [hep-th].
[42] J. Gomis, D. Rodriguez-Gomez, M. Van Raamsdonk, H. Verlinde, Supersymmetric Yang–Mills theory from Lorentzian three-algebras, arXiv: 0806.0738 [hep-th].
[43] O. Aharony, O. Bergman, D.L. Jafferis, J. Maldacena, $N = 6$ superconformal Chern–Simons-matter theories, M2-branes and their gravity duals, arXiv: 0806.1218 [hep-th].
[44] C. Ahn, Holographic supergravity dual to three-dimensional $N = 2$ gauge theory, arXiv: 0806.1420 [hep-th].
[45] M. Benna, I. Klebanov, T. Klose, M. Smedback, Superconformal Chern–Simons theories and AdS_4/CFT_3 correspondence, arXiv: 0806.1519 [hep-th].
[46] B. Ezhuthachan, S. Mukhi, C. Papageorgakis, D2 to D2, JHEP 0807 (2008) 041, arXiv: 0806.1639 [hep-th].
[47] S. Cecotti, A. Sen, Coulomb branch of the Lorentzian three-algebra theory, arXiv: 0806.1990 [hep-th].
[48] E.A. Bergshoeff, M. de Roo, O. Hohm, D. Roest, Multiple membranes from gauged supergravity, arXiv: 0806.2584 [hep-th].
[49] P. de Medeiros, J. Figueroa-O'Farrill, E. Mendez-Escobar, Metric Lie 3-algebras in Bagger–Lambert theory, arXiv: 0806.3242 [hep-th].
[50] M. Blau, M. O'Loughlin, Multiple M2-branes and plane waves, arXiv: 0806.3253 [hep-th].
[51] M. Blau, M. O'Loughlin, DLCQ and plane wave matrix Big Bang models, arXiv: 0806.3255 [hep-th].
[52] T. Nishioka, T. Takayanagi, On type IIA Penrose limit and $N = 6$ Chern–Simons theories, arXiv: 0806.3391 [hep-th].
[53] Y. Honma, S. Iso, Y. Sumitomo, S. Zhang, Scaling limit of $N = 6$ superconformal Chern–Simons theories and Lorentzian Bagger–Lambert theories, arXiv: 0806.3498 [hep-th].
[54] Y. Imamura, K. Kimura, Coulomb branch of generalized ABJM models, arXiv: 0806.3727 [hep-th].
[55] J.A. Minahan, K. Zarembo, The Bethe ansatz for superconformal Chern–Simons, arXiv: 0806.3951 [hep-th].
[56] K. Furuuchi, S.Y. Shih, T. Takimi, M-theory superalgebra from multiple membranes, arXiv: 0806.4044 [hep-th].
[57] A. Agarwal, Mass deformations of super-Yang–Mills theories in $D = 2 + 1$, and super-membranes: A note, arXiv: 0806.4292 [hep-th].
[58] D. Gaiotto, S. Giombi, X. Yin, Spin chains in $N = 6$ superconformal Chern–Simons-matter theory, arXiv: 0806.4589 [hep-th].
[59] I.A. Bandos, P.K. Townsend, Light-cone M5 and multiple M2-branes, arXiv: 0806.4777 [hep-th].
[60] C. Ahn, Towards holographic gravity dual of $N = 1$ superconformal Chern–Simons gauge theory, arXiv: 0806.4807 [hep-th].
[61] J. Bedford, D. Berman, A note on quantum aspects of multiple membranes, arXiv: 0806.4900 [hep-th].
[62] G. Arutyunov, S. Frolov, Superstrings on $AdS_4 \times CP^3$ as a coset sigma-model, arXiv: 0806.4940 [hep-th].
[63] B. Stefanski Jr., Green–Schwarz action for type IIA strings on $AdS_4 \times CP^3$, arXiv: 0806.4948 [hep-th].
[64] G. Grignani, T. Harmark, M. Orselli, The $SU(2) \times SU(2)$ sector in the string dual of $N = 6$ superconformal Chern–Simons theory, arXiv: 0806.4959 [hep-th].
[65] K. Hosomichi, K.M. Lee, S. Lee, S. Lee, J. Park, $N = 5, 6$ superconformal Chern–Simons theories and M2-branes on orbifolds, arXiv: 0806.4977 [hep-th].
[66] P. Fre, P.A. Grassi, Pure spinor formalism for $Osp(N - 4)$ backgrounds, arXiv: 0807.0044 [hep-th].
[67] J. Bagger, N. Lambert, Three-algebras and $N = 6$ Chern–Simons gauge theories, arXiv: 0807.0163 [hep-th].
[68] S. Terashima, On M5-branes in $N = 6$ membrane action, arXiv: 0807.0197 [hep-th].
[69] G. Grignani, T. Harmark, M. Orselli, G.W. Semenoff, Finite size giant magnons in the string dual of $N = 6$ superconformal Chern–Simons theory, arXiv: 0807.0205 [hep-th].

[70] S. Terashima, F. Yagi, Orbifolding the membrane action, arXiv: 0807.0368 [hep-th].
[71] N. Gromov, P. Vieira, The AdS_4/CFT_3 algebraic curve, arXiv: 0807.0437 [hep-th].
[72] C. Ahn, P. Bozhilov, Finite-size effects of membranes on $AdS_4 \times S_7$, arXiv: 0807.0566 [hep-th].
[73] N. Gromov, P. Vieira, The all loop AdS_4/CFT_3 Bethe ansatz, arXiv: 0807.0777 [hep-th].
[74] B. Chen, J.B. Wu, Semi-classical strings in $AdS_4 \times CP^3$, arXiv: 0807.0802 [hep-th].
[75] S. Cherkis, C. Saemann, Multiple M2-branes and generalized 3-Lie algebras, arXiv: 0807.0808 [hep-th].
[76] C.S. Chu, P.M. Ho, Y. Matsuo, S. Shiba, Truncated Nambu–Poisson bracket and entropy formula for multiple membranes, arXiv: 0807.0812 [hep-th].
[77] M.A. Bandres, A.E. Lipstein, J.H. Schwarz, Studies of the ABJM theory in a formulation with manifest $SU(4)$ R-symmetry, arXiv: 0807.0880 [hep-th].
[78] Y. Zhou, Interpreting N M2 branes, arXiv: 0807.0890 [hep-th].
[79] J. Gomis, D. Rodriguez-Gomez, M. Van Raamsdonk, H. Verlinde, A massive study of M2-brane proposals, arXiv: 0807.1074 [hep-th].
[80] M. Schnabl, Y. Tachikawa, Classification of $N = 6$ superconformal theories of ABJM type, arXiv: 0807.1102 [hep-th].
[81] T. Li, Y. Liu, D. Xie, Multiple D2-brane action from M2-branes, arXiv: 0807.1183 [hep-th].
[82] N. Kim, How to put the Bagger–Lambert theory on an orbifold: A derivation of the ABJM model, arXiv: 0807.1349 [hep-th].
[83] Y. Pang, T. Wang, From N M2's to N D2's, arXiv: 0807.1444 [hep-th].
[84] M.R. Garousi, A. Ghodsi, M. Khosravi, On thermodynamics of $N = 6$ superconformal Chern–Simons theories at strong coupling, arXiv: 0807.1478 [hep-th].
[85] A. Hashimoto, P. Ouyang, Supergravity dual of Chern–Simons Yang–Mills theory with $N = 6, 8$ superconformal IR fixed point, arXiv: 0807.1500 [hep-th].
[86] D. Astolfi, V.G.M. Puletti, G. Grignani, T. Harmark, M. Orselli, Finite-size corrections in the $SU(2) \times SU(2)$ sector of type IIA string theory on $AdS_4 \times CP^3$, arXiv: 0807.1527 [hep-th].
[87] M. Ali-Akbari, M.M. Sheikh-Jabbari, J. Simon, Relaxed three-algebras: Their matrix representations and implications for multi M2-brane theory, arXiv: 0807.1570 [hep-th].
[88] C. Ahn, R.I. Nepomechie, $N = 6$ super-Chern–Simons theory S-matrix and all-loop Bethe ansatz equations, arXiv: 0807.1924 [hep-th].
[89] K. Hanaki, H. Lin, M2–M5 systems in $N = 6$ Chern–Simons theory, arXiv: 0807.2074 [hep-th].
[90] H. Verlinde, D2 or M2? A note on membrane scattering, arXiv: 0807.2121 [hep-th].
[91] Y. Imamura, K. Kimura, $N = 4$ Chern–Simons theories with auxiliary vector multiplets, arXiv: 0807.2144 [hep-th].
[92] Y. Lozano, Eleven dimensions from the massive D2-brane, Phys. Lett. B 414 (1997) 52, hep-th/9707011.
[93] T. Ortin, A note on the D2-brane of the massive type IIA theory and gauged sigma models, Phys. Lett. B 415 (1997) 39, hep-th/9707113.
[94] L.J. Romans, Massive $N = 2a$ supergravity in ten dimensions, Phys. Lett. B 169 (1986) 374.
[95] E. Bergshoeff, M. de Roo, M.B. Green, G. Papadopoulos, P.K. Townsend, Duality of type II 7-branes and 8-branes, Nucl. Phys. B 470 (1996) 113, hep-th/9601150.
[96] E. Eyras, B. Janssen, Y. Lozano, 5-branes, KK-monopoles and T-duality, Nucl. Phys. B 531 (1998) 275, hep-th/9806169.
[97] E. Bergshoeff, J.P. van der Schaar, On M9-branes, Class. Quantum Grav. 16 (1999) 23, hep-th/9806069.
[98] E. Bergshoeff, Y. Lozano, T. Ortin, Massive branes, Nucl. Phys. B 518 (1998) 363, hep-th/9712115.
[99] E. Bergshoeff, P.M. Cowdall, P.K. Townsend, Massive IIA supergravity from the topologically massive D2-brane, Phys. Lett. B 410 (1997) 13, hep-th/9706094.
[100] Y. Honma, S. Iso, Y. Sumitomo, H. Umetsu, S. Zhang, Generalized conformal symmetry and recovery of $SO(8)$ in multiple M2 and D2 branes, arXiv: 0807.3825 [hep-th].
[101] E.A. Bergshoeff, O. Hohm, D. Roest, H. Samtleben, E. Sezgin, The superconformal gaugings in three dimensions, arXiv: 0807.2841 [hep-th].
[102] B.H. Lee, K.L. Panigrahi, C. Park, Spiky strings on $AdS_4 \times CP^3$, arXiv: 0807.2559 [hep-th].
[103] C. Krishnan, C. Maccaferri, Membranes on calibrations, arXiv: 0805.3125 [hep-th].