Bulg. J. Phys. 42 (2015) 288–295

Scalar D-brane Fluctuations and Holographic Mesons in Pilch-Warner Background*

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Received *21 June 2015*

Abstract. In this paper we study the D7 probe brane scalar fluctuations in global Pilch-Warner supergravity background solution. Our choice of constant classical embeddings for the probe D-brane considerably simplifies the analysis of its fluctuations. The corresponding holographic meson spectra, obtained by the fluctuations along the transverse directions, admit equidistant structure for the higher modes and a ground state given by the conformal dimension of the operator dual to the fluctuations.

PACS codes: 11.25.Tq, 11.25.Uv

1 Introduction

The AdS/CFT correspondence is a duality conjecture relating higher dimensional string theory in the weak coupling regime to a 4 dimensional $SU(N)$ gauge field theory with strong coupling constant, and vice-versa. The lower dimensional gauge field theory lives on the asymptotic boundary of the space-time, where the strings move. This correspondence gives us the opportunity to study non-perturbative phenomena in Yang-Mills theory with the classical tools available in superstring theory and supergravity.

On the gauge theory side of the original Maldacena setup [1] one has $\mathcal{N} = 4$ superconformal Yang-Mills theory, while on the string side of the correspondence there is a stack of N $D3$-branes with open strings attached to the same stack of $D3$-branes. This configuration of fundamental strings and D-branes generates states, which transform only in the adjoint representation of the gauge group. Adding flavours in the fundamental representation can be achieved by introducing an additional stack of N_f $D7$ flavour probe branes [2]. Here, the $SU(N_f)$ is a global flavor symmetry.

*This work is in commemoration of 75th anniversary of Prof. Matey Mateev.

If we require large number N_c of the color $D3$-branes, we can no longer neglect their gravitational influence on the background. Therefore, the stack of the N_c $D3$-branes sources a supergravity background geometry, which in the original version of the AdS/CFT correspondence, is the $AdS_5 \times S^5$ space-time. In fact, in the limit $N_f \ll N_c$, we can effectively replace the stack of N_c $D3$-branes by their background geometry. What is left is a strongly coupled dual gauge theory, living on the asymptotic boundary of the 10-dimensional space-time, and a number of flavour $D7$-branes, which we can think of as probe branes. In this set up the low-energy fluctuations of the flavour D-branes correspond to meson excitations in the dual gauge field theory.

Variety of supersymmetric meson spectra were found in [3, 4] (for a review see [5]), To produce more realistic QCD like string theories, deformations of the initial $AdS_5 \times S^5$ geometry [6, 7], or introducing external magnetic or electric fields [8–11], have to be considered. Such configurations will break the supersymmetry and theories with less supersymmetry will emerge.

An example of such deformed supergravity background is the Pilch-Warner geometry [12, 13]. It is a solution of five-dimensional $\mathcal{N} = 8$ gauged supergravity lifted to ten dimensions, which, in its infrared critical point, preserves 1/4 of the original supersymmetry. On the SYM side of the duality the IR fixed point corresponds to large N limit of the superconformal $\mathcal{N} = 1$ theory of Leigh-Strassler [14]. In this study we will restrict ourselves to the IR critical point.

2 Flavours in Global Pilch-Warner Geometry

The ten-dimensional Pilch-Warner metric in global coordinates is given by [15]

$$\mathrm{d}s^2_{1,4}(IR) = L^2\,\Omega^2\,\left(-\cosh^2\!\rho\,\mathrm{d}\tau^2 + \mathrm{d}\rho^2 + \sinh^2\!\rho\,\mathrm{d}\Omega_3^2\right)\,, \qquad \text{(1a)}$$

$$\begin{aligned}\mathrm{d}s^2_5(IR) = \frac{2}{3}L^2\Omega^2\Big[\mathrm{d}\theta^2 &+ \frac{4\cos^2\theta}{3-\cos 2\theta}(\sigma_1^2+\sigma_2^2) + \frac{4\sin^2 2\theta}{(3-\cos 2\theta)^2}(\sigma_3+\mathrm{d}\phi)^2\\ &+\frac{2}{3}\Big(\frac{1-3\cos 2\theta}{\cos 2\theta-3}\Big)^2\Big(\mathrm{d}\phi-\frac{4\cos^2\theta}{1-3\cos 2\theta}\sigma_3\Big)^2\Big]\,,\end{aligned} \qquad \text{(1b)}$$

where $\mathrm{d}\Omega_3^2 = \mathrm{d}\phi_1^2 + \sin^2\phi_1\,(\mathrm{d}\phi_2^2 + \sin^2\phi_2\,\mathrm{d}\phi_3^2)$ is the metric on the 3-sphere, and

$$\Omega^2 = \frac{2^{1/3}}{\sqrt{3}}\sqrt{3-\cos(2\,\theta)} \qquad (2)$$

is the warp factor at the IR point. Our left-invariant one-forms satisfy $\mathrm{d}\sigma_i = \varepsilon_{ijk}\,\sigma_j \wedge \sigma_k$ in such a way that $\mathrm{d}\tilde{\Omega}_3^2 = \sigma_i\,\sigma_i$ gives the metric on the unit 3-sphere. In global Pilch-Warner coordinates one has:

$$\sigma_1 = \frac{1}{2}(\sin\beta\mathrm{d}\alpha - \cos\beta\sin\alpha\,\mathrm{d}\gamma)\,, \qquad (3)$$

$$\sigma_2 = -\frac{1}{2}\left(\cos\beta\,\mathrm{d}\alpha + \sin\beta\,\sin\alpha\,\mathrm{d}\gamma\right), \tag{4}$$

$$\sigma_3 = \frac{1}{2}\left(\mathrm{d}\beta + \cos\alpha\,\mathrm{d}\gamma\right). \tag{5}$$

The Pilch-Warner background includes non-trivial Ramond-Ramond (R-R) and Neveu-Schwarz (NS-NS) form fields entering the $D7$ probe brane action. The full action is given by two terms: a Dirac-Born-Infeld (DBI) term and a Wess-Zumino (WZ) term [16, 20]:

$$\begin{aligned} S_{D_7} = S_{DBI} + S_{WZ} &= -T_7 \int d^8\xi\, e^{-\Phi}\sqrt{-\det(P[G]+\mathcal{F})} \\ &- T_7 \int \left(P[C_8] - P[C_6]\wedge\mathcal{F} + \frac{1}{2}P[C_4]\wedge\mathcal{F}\wedge\mathcal{F} + \cdots\right), \end{aligned} \tag{6}$$

where $\mathcal{F} = P[B_2] + 2\,\pi\,\alpha'\,F_2$ is the invariant gauge two-form, F_2 is the world volume gauge field, B_2 is the Kalb-Ramond 2-form, T_7 is the $D7$-brane tension, Φ is the dilaton, and P denotes the pullback of the bulk space-time tensor to the world-volume of the brane:

$$P[G]_{ab} = G_{AB}\,\frac{\partial X^A}{\partial\xi^a}\,\frac{\partial X^B}{\partial\xi^b}\,. \tag{7}$$

The indices $a,\, b = 0,\ldots,7$ span the world volume of the $D7$-brane, while $A,\, B = 0,\ldots,9$ span the whole space-time. In static gauge the pullback is given by:

$$P[G]_{ab} = g_{ab} + G_{mn}\,\frac{\partial X^m}{\partial\xi^a}\,\frac{\partial X^n}{\partial\xi^b}\,, \tag{8}$$

where g_{ab} is the induced metric on $D7$, and G_{mn} $(m, n = \theta,\, \phi)$ are the metric components in front of the transverse coordinates governing the $D7$-brane fluctuations.

In the IIB string theory the R-R potentials are C_0, C_2, C_4, C_6, C_8 and corresponding field strengths, which satisfy certain Bianchi identities and equations of motion [16]:

$$\Phi = C_0 = 0,\; F_1 = dC_0 = 0\,, \tag{9a}$$

$$C_2 = \Re e(A_2),\; B_2 = \Im m(A_2)\,, \tag{9b}$$

$$H_3 = dB_2,\; F_3 = dC_2 - C_0\wedge H_3 = dC_2\,, \tag{9c}$$

$$\mathrm{d}F_3 = \mathrm{d}H_3 = 0,\; \mathrm{d}F_5 = H_3\wedge F_3\,, \tag{9d}$$

$$\mathrm{d}(\star F_3) = -H_3\wedge F_5,\; \mathrm{d}(\star H_3) = F_3\wedge F_5,\; F_5 = \star F_5\,, \tag{9e}$$

$$dC_4 + d\tilde{C}_4 = F_5 + C_2\wedge H_3\,, \tag{9f}$$

$$F_7 = \star F_3 = dC_6 - C_4\wedge H_3\,, \tag{9g}$$

$$F_9 = \star F_1 = 0 = dC_8 - C_6\wedge H_3 = C_6\wedge H_3,\; \chi = C_8 = 0\,. \tag{9h}$$

Here the field strengths are defined in terms of the corresponding potentials as

$$H_3 \equiv \mathrm{d}B_2, \quad F_\mathrm{p} \equiv \mathrm{d}C_{\mathrm{p}-1} - C_{\mathrm{p}-3} \wedge H_3. \tag{10}$$

In this setup the axion/dilaton system of scalars (9a) and (9h) is trivial along the flow. We also have an ansatz for the self-dual five form:

$$F_5 = -\frac{2^{5/3}}{3}\, L^4\, \cosh\rho \sinh^3\rho\,(1+\star)\,\mathrm{d}\tau \wedge \mathrm{d}\rho \wedge \epsilon(S^3_\phi)\,, \tag{11}$$

where $\epsilon(S^3_\phi) = \sin^2\phi_1\, \sin\phi_2\, \mathrm{d}\phi_1 \wedge \mathrm{d}\phi_2 \wedge \mathrm{d}\phi_3$ is the volume element of the unit 3-sphere S^3_ϕ, and $\star$ is the Hodge star operator. The ansatz for the 2-form potential A_2 at the IR point is also known:

$$\begin{aligned} A_2(IR) = C_2 + i\,B_2 = \\ -\frac{i}{2} e^{-2i\phi} L_0^2 \cos\theta \Big(\mathrm{d}\theta - \frac{2i\sin 2\theta}{3-\cos 2\theta}(\sigma_3 + \mathrm{d}\phi)\Big) \wedge (\sigma_1 + i\sigma_2)\,. \end{aligned} \tag{12}$$

Two additional constraints are necessary for (9a) to be consistent, namely

$$F_3 \wedge \star H_3 = 0, \qquad F_3 \wedge \star F_3 = H_3 \wedge \star H_3\,. \tag{13}$$

The explicit form of the R-R and NS potentials in global Pilch-Warner coordinates has been calculated in [17].

3 Scalar Fluctuations of the $D7$-Brane and the Meson Spectra

We begin by choosing a classical embedding of D7-probe brane in global Pilch-Warner background geometry. The most general D7 embedding is of the form

$$\xi^a = (\tau, \rho, \phi_1, \phi_2, \phi_3, \alpha, \beta, \gamma), \quad \theta = \theta(\xi^a), \quad \phi = \phi(\xi^a)\,, \tag{14}$$

where ξ^a are the D7-brane world volume coordinates, and (θ,ϕ) are normal to the brane. In [17] the authors study the following classical kappa-symmetric embedding of the D7 probe brane:

$$\theta = 0, \quad \phi = -\beta + c\,. \tag{15}$$

In this study we are going to consider a slightly different embedding of the form:

$$\theta = \text{const}\,, \quad \phi = -b\,\beta + \text{const}\,. \tag{16}$$

where b is some constant. This choice for θ and ϕ solves the classical D-brane equations of motion, but do not, in general, satisfy the kappa-symmetry or supersymmetry preserving conditions [17–19].

3.1 Fluctuation equations

In order to calculate the holographic meson spectrum we will use the classical D7-brane embedding (16), where we can take $\theta = 0$ and redefine the ϕ coordinate as $\phi \to \phi + b\,\beta = const$. Making the proper shift in the 10-dimensional Pilch-Warner metric we can study the D-brane fluctuations around the constant embedding $\phi = 0$. For simplicity we set the world volume gauge field $F = 0$. One can choose the following ansatz for the fluctuations along θ and ϕ:

$$\theta = 0 + \eta\,\Theta(\xi^a), \quad \phi = 0 + \eta\,\Phi(\xi^a)\,. \tag{17}$$

After expanding the D7 Lagrangian (6) up to quadratic order in the fluctuations (keeping only terms quadratic in $\eta = 2\,\pi\,\alpha'$), one finds the following set of equations for the fluctuation fields Φ and Θ:

$$-\partial_\tau^2\Phi + \cosh^2\!\rho\,\tilde{\Delta}_\rho\Phi + \coth^2\!\rho\,\Delta_{\phi_i}\Phi\, + 3\cosh^2\!\rho\,\tilde{\Delta}_{\alpha_i}\Phi = 0\,, \tag{18}$$

$$-\partial_\tau^2\Theta + \cosh^2\!\rho\,\tilde{\tilde{\Delta}}_\rho\Theta + \coth^2\!\rho\,\Delta_{\phi_i}\Theta + 3\cosh^2\!\rho\,\tilde{\Delta}_{\alpha_i}\Theta + \frac{6}{1+b}\cosh^2\!\rho\,\Theta = 0\,, \tag{19}$$

where $\phi_i = (\phi_1,\, \phi_2,\, \phi_3)$, $\alpha_i = (\alpha,\, \beta,\, \gamma)$, and

$$\tilde{\Delta}_\rho\Phi = \partial_\rho^2\Phi + (3\,\coth\rho + \tanh\rho)\,\partial_\rho\Phi\,, \tag{20}$$

$$\tilde{\tilde{\Delta}}_\rho\Theta = \partial_\rho^2\Theta + \frac{\big(7 + 4\,b + 3\,(1+b)\,\mathrm{csch}^2\rho\big)\,\tanh\rho}{1+b}\,\partial_\rho\Theta\,, \tag{21}$$

$$\Delta_{\phi_i}\mathcal{P} = \frac{\partial_{\phi_1}\left(\sin^2\!\phi_1\,\partial_{\phi_1}\mathcal{P}\right)}{\sin^2\!\phi_1} + \frac{\partial_{\phi_2}\left(\sin\phi_2\,\partial_{\phi_2}\mathcal{P}\right)}{\sin^2\!\phi_1\,\sin\phi_2} + \frac{\partial_{\phi_3}^2\mathcal{P}}{\sin^2\!\phi_1\,\sin^2\!\phi_2}\,, \tag{22}$$

$$\tilde{\Delta}_{\alpha_i}\mathcal{P} = \frac{\partial_\alpha\Big(\sin\alpha\,\partial_\alpha\mathcal{P}\Big)}{\sin\alpha} + \frac{1}{\sin^2\!\alpha}\Big(\frac{\cos 2\alpha + 7}{8(1+b)^2}\partial_\beta^2\mathcal{P} + \partial_\gamma^2\mathcal{P} - \frac{2\cos\alpha}{1+b}\partial_{\beta\gamma}^2\mathcal{P}\Big). \tag{23}$$

Here $\mathcal{P}$ stands for Θ or Φ.

3.2 Meson spectrum along ϕ

Next we proceed with finding the meson spectrum along the ϕ transverse direction. One can separate the variables in eq. (18) as:

$$\Phi(\xi^a) = e^{i\,\omega\,\tau}\,R(\rho)\,\mathcal{Y}^\ell(S^3_{\phi_i})\,Z(\tilde{S}^3_{\alpha_i})\,, \tag{24}$$

which leads to the following set of spectral equations:

$$\ddot{T}(\tau) = -\omega^2\,T(\tau)\,, \tag{25}$$

$$\tilde{\Delta}_{\alpha_i}Z(\alpha_i) = -\nu\,Z(\alpha_i)\,, \tag{26}$$

$$\Delta_{\phi_i}\mathcal{Y}^\ell(\phi_i) = -\ell\,(\ell+2)\,\mathcal{Y}^\ell(\phi_i)\,, \tag{27}$$

$$R''(\rho) + (3\,\coth\rho + \tanh\rho)\,R'(\rho) + \left(\frac{\omega^2}{\cosh^2\!\rho} - \frac{\ell\,(\ell+2)}{\sinh^2\!\rho} - 3\,\nu\right)\,R(\rho) = 0\,, \tag{28}$$

where ω is the energy of the fluctuations and $\mathcal{Y}^\ell(\phi_i)$ are the hyperspherical harmonics, $\ell \in \mathbb{N}_0$. In order to facilitate the calculation of the spectrum we change the radial variable $\sinh\rho = r$ in eq. (28):

$$R''(r) + \frac{3+5\,r^2}{r\,\left(r^2+1\right)}\,R'(r) + \left(\frac{\omega^2}{\left(r^2+1\right)^2} - \frac{\ell\,(\ell+2)}{r^2\,\left(r^2+1\right)} - \frac{3\,\nu}{r^2+1}\right)\,R(r) = 0\,. \tag{29}$$

The only regular solution (up to a normalization constant) at the origin $r = 0$ is given by

$$R(r) = c\,r^\ell\,\left(r^2+1\right)^{-\frac{\omega}{2}}\,{}_2F_1\left(a+\ell+1,\,b+\ell+1;\,2+\ell;\,-r^2\right)\,, \tag{30}$$

where

$$a = \left(-\ell-\omega-\sqrt{3\,\nu+4}\right)/2, \quad b = \left(-\ell-\omega+\sqrt{3\,\nu+4}\right)/2\,.$$

To assure normalizability at infinity one has to terminate the series of the hypergeometric function at some finite non-negative integer power n. The hypergeometric function becomes a polynomial of degree n, if one of its first two arguments is set to a negative integer $-n$, $n \geq 0$. Therefore, setting $b - c + 1 = -n$, gives the quantization condition and the form of the scalar meson spectrum[1]:

$$\omega = \Delta + \ell + 2\,n\,. \tag{31}$$

Here $\Delta = 2 + \sqrt{3\nu+4}$ is the conformal dimensions of the operators dual to the fluctuations. We conclude that the energy of the ground state $(n, \ell = 0)$ is given by the conformal dimension of the operator dual to the fluctuations. For the higher modes the spectrum is equidistant.

3.3 Fluctuations along θ

The analysis of the D7 fluctuations along θ is similar to the one done in the previous subsection. Separation of variables of the form (24) leads to the following radial equation:

$$R''(r) + \left(\frac{3}{r} + \frac{(5+2\,b)\;r}{(1+b)\;(1+r^2)}\right)\,R'(r) + \\ + \left(\frac{\omega^2}{\left(1+r^2\right)^2} - \frac{\ell\;(\ell+2)}{r^2\;(1+r^2)} - \frac{3}{1+r^2}\;\left(\nu - \frac{2}{1+b}\right)\right)\,R(r) = 0\,. \tag{32}$$

[1] The hypergeometric function is symmetric in its first two arguments, therefore, without any loss of generality, one can chose to work with one argument or the other.

Its regular solution is given by

$$R(r) = C\, r^{\ell} \left(1 + r^2\right)^{\frac{-3-\sqrt{9+4\,(1+b)^2\,\omega^2}}{4(1+b)}} F\left(a,\, b,\, 2+l,\, -r^2\right)\,, \tag{33}$$

where

$$a = \frac{1}{4\,(1+b)}\,\left(4 + 2\,\ell + b\,(4 + 2\,\ell)\right.$$
$$\left. + \sqrt{25 + 12\,\nu + b^2\,(16 + 12\,\nu) + b\,(32 + 24\,\nu)} - \sqrt{9 + 4\,(1+b)^2\,\omega^2}\right)\,,$$

and

$$b = \frac{1}{4\,(1+b)}\,\left(4 + 2\,\ell + b\,(4 + 2\,\ell)\right.$$
$$\left. - \sqrt{25 + 12\,\nu + b^2\,(16 + 12\,\nu) + b\,(32 + 24\,\nu)} - \sqrt{9 + 4\,(1+b)^2\,\omega^2}\right)\,.$$

Imposing normalizability one finds the following form of the meson spectrum along θ:

$$\omega^2 = (\Delta + \ell + 2\,n)^2 - \frac{9}{4\,(1+b)^2}\,, \tag{34}$$

where the conformal dimension Δ is given by

$$\Delta = -\frac{1}{2\,(1+b)\,(-2+\nu+b\,\nu)}\,\left(8\,(1+b) - 4\,(1+b)^2\,\nu\right.$$
$$\left. + \sqrt{2-(1+b)\,\nu}\,\sqrt{(2-(1+b)\,\nu)\,(25 + 12\,\nu + 4\,b\,(2+b)\,(4+3\,\nu))}\right)\,,$$

If we set $b = 1$, the result for the spectrum coincides with the kappa-symmetric one obtained in [17]. However there is no value for b, which can cancel the mysterious additional constant shift appearing in the ground state $(n, \ell = 0)$ of the meson spectrum (34).

4 Conclusions

In this study we considered the D7-brane embedding compatible with the brane equations of motion, but, in general, not compatible with the kappa-symmetry or supersymmetry preserving conditions [17–19]. Working in the limit $N_f \ll N_C$ allowed us to study the flavour D7-branes as probe branes, thus neglecting their gravitational back-reaction on the background. All meson spectra are obtained analytically and admit equidistant structure in their higher modes. The corresponding ground states are proportional to the conformal dimension of the operators dual to the D7-brane fluctuations.

Acknowledgments

The authors would like to thank V. Filev and H. Dimov for insightful discussions during the preparation of the paper. This work was partially supported by the Bulgarian NSF grant DFNI T02/6.

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