Abstract We discuss a holographic soft-wall model developed for the description of mesons and baryons with adjustable quantum numbers $n, J, L, S$. This approach is based on an action which describes hadrons with broken conformal invariance and which incorporates confinement through the presence of a background dilaton field.

Keywords Soft-wall holographic model, dilaton, hadrons, mass spectrum

1 Introduction

Based on the correspondence of string theory in anti-de Sitter (AdS) space and conformal field theory (CFT) in physical space-time [1], a class of AdS/QCD approaches was recently successfully developed for describing the phenomenology of hadronic properties. In order to break conformal invariance and incorporate confinement in the infrared (IR) region two alternative AdS/QCD backgrounds have been suggested in the literature: the “hard-wall” approach [2], based on the introduction of an IR brane cutoff in the fifth dimension, and the “soft-wall” approach [3], based on using a soft cutoff. This last procedure can be introduced in the following ways: i) as a background field (dilaton) in the overall exponential of the action, ii) in the warping factor of the AdS metric, iii) in the effective potential of the action. These methods are in principle equivalent to each other due to a redefinition of the bulk field involving the dilaton field or by a redefinition of the effective potential. In the literature there exist detailed discussions of the sign of the dilaton profile in the dilaton exponential \( \exp(\pm \phi) \) [3, 4, 5, 6, 7] for the soft-wall model (for a discussion of the sign of the dilaton in the warping factor of the AdS metric see Refs. [8]). The negative sign was suggested in Ref. [3] and recently discussed in Ref. [7]. It leads to a Regge-like behavior of the meson spectrum, including a straightforward extension to fields of higher spin $J$. Also, in Ref. [7] it was shown that this choice of the dilaton sign guarantees the absence of a spurious massless scalar mode in the vector channel of the soft-wall model. We stress that alternative versions of this model with a positive sign are also possible. One should just redefine the bulk field $V$ as $V = \exp(\varphi) \tilde{V}$, where the transformed field corresponds to the dilaton with an opposite profile. It is clear that the underlying action changes, and extra potential terms are generated depending on the dilaton field (see detailed discussion in [9]).
2 Bosonic case

First we demonstrate the equivalence of both versions of the soft-wall model with a positive and negative dilaton profile in the case of the bosonic field, and afterwards we consider the fermionic field and extension to higher values of total angular momentum $J$. We consider the propagation of a scalar field $S(x,z)$ in $d+1$ dimensional AdS space. The AdS metric is specified by $ds^2 = g_{M N} dx^M dx^N = \eta_{a b} e^{2A(z)} dx^a dx^b = \epsilon^{2A(z)} (\eta_{\mu \nu} dx^\mu dx^\nu - dz^2)$, where $M$ and $N = 0, 1, \cdots, d$ are the space-time (base manifold) indices, $a = (\mu, z)$ and $b = (\nu, z)$ are the local Lorentz (tangent) indices, $g_{M N}$ and $\eta_{a b}$ are curved and flat metric tensors, which are related by the vielbein $e^a_M (z) = \epsilon^{a(2)} \delta^a_M$ as $g_{M N} = e^a_M e^b_N \eta_{a b}$. Here $z$ is the holographic coordinate, $R$ is the AdS radius, and $g = |\text{det} g_{M N}| = e^{2A(z)(d+1)}$. In the following we restrict ourselves to a conformal-invariant metric with $A(z) = \log(R/z)$.

The actions for the scalar field ($J = 0$) with a positive and negative dilaton are \[ S_0 = \frac{1}{2} \int d^d x dz \sqrt{g} e^{\Phi(z)} \left[ g^{M N} \partial_M S^+(x,z) \partial_N S^+(x,z) - \mu_0^2 S^+(x,z) S^+(x,z) \right] \] \[ S_0 = \frac{1}{2} \int d^d x dz \sqrt{g} e^{-\Phi(z)} \left[ g^{M N} \partial_M S^-(x,z) \partial_N S^-(x,z) - \left( \mu_0^2 + \Delta V_0(z) \right) S^-(x,z) S^-(x,z) \right] . \]

The superscripts $+$ and $-$ correspond to the cases of a positive or negative dilaton, respectively. The actions are equivalent to each other, which is obvious by the bulk field redefinition: $S^\pm (x,z) = e^{\pm \Phi(z)} S^\mp (x,z)$. The difference between the two actions is absorbed in the effective potential $\Delta V_0(z) = e^{-2A(z)} \Delta U_0(z)$, where $\Delta U_0(z) = \phi''(z) + (d - 1) \phi'(z) A'(z)$. The quantity $\mu_0^2 R^2 = \Delta (\Delta - d)$ is the bulk boson mass, where $\Delta$ is the dimension of the interpolating operator dual to the scalar bulk field. For the case of the bulk fields dual to the scalar fields $\Delta = 2 + L$, where $L = \max |L_j|$ is the maximal value of the $z$-component of the orbital angular momentum \[ S^\pm (x,z) = e^{\pm \Phi(z)/2} S^\mp (x,z) \] In terms of the field $S(x,z)$ the transformed action, which now is universal for both versions of the soft-wall model, reads \[ S_0 = \frac{1}{2} \int d^d x dz \sqrt{g} \left[ g^{M N} \partial_M S(x,z) \partial_N S(x,z) - (\mu_0^2 + V_0(z)) S^2 (x,z) \right] , \]

where $V_0(z) = e^{-2A(z)} U_0(z)$ and with the effective potential \[ U_0(z) = \frac{1}{2} \phi''(z) + \frac{1}{4} (\phi'(z))^2 + \frac{d - 1}{2} \phi'(z) A'(z) . \]

The last expression is identical with the light-front effective potential found in Ref. \[ S_0 = \frac{1}{2} \int d^d x dz e^{B_0(z)} \left[ \partial_\mu S(x,z) \partial^\mu S(x,z) - \partial_\mu S(x,z) \partial_\nu S(x,z) - \left( e^{2A(z)} \mu^2 + U_0(z) \right) S^2 (x,z) \right] , \]

Then we use a Kaluza-Klein (KK) expansion $S(x,z) = \sum_n S_n(x) \Phi_n(z)$, where $n$ is the radial quantum number, $S_n(x)$ is the tower of the KK modes dual to scalar mesons and $\Phi_n$ are their extra-dimensional profiles (wave-functions). We suppose a free propagation of the bulk field along the $d$ Poincaré coordinates with four-momentum $p$, and a constrained propagation along the $(d+1)$-th coordinate $z$ (due to confinement imposed by the dilaton field).
Performing the substitution $\Phi_n(z) = e^{-B_n(z)/2} \phi_n(z)$ and restricting to mass-shell $p^2 = M_{\delta 0}^2$ we derive the Schrödinger-type EOM for $\Phi_n(z)$:

$$\left[ -\frac{d^2}{dz^2} + \frac{4L^2 - 1}{4z^2} + U_0(z) \right] \phi_n(z) = M_{\delta 0}^2 \phi_n(z)$$  \hspace{1cm} (6)

where

$$M_{\delta 0}^2 = 4\kappa^2 \left( n + \frac{1}{2} \right)$$

and $M_{\delta 0}^2 = 4\kappa^2 \left( n + \frac{1}{2} \right)$ is the mass spectrum of scalar field. Here we use the generalized Laguerre polynomials $L_n^\kappa(x)$. Notice that the normalizable mode $\phi_n(z)$ has the correct behavior in both the ultraviolet (UV) and infrared (IR) limits: $\Phi_n(z) \sim z^{2+L}$ at small $z$, $\Phi_n(z) \rightarrow 0$ at large $z$. Using the KK expansion, EOM for the KK profiles $\Phi_n(z)$, the $d+1$-dimensional action for the bulk field reduces to a $d$-dimensional action for the scalar fields $S_n(x)$ dual to scalar mesons with masses $M_{\delta 0}$:

$$S_b^{(d)} = \frac{1}{2} \sum_n \int d^d x \left[ \partial_\mu S_n(x) \partial^\mu S_n(x) - M_{\delta 0}^2 S_n^2(x) \right].$$  \hspace{1cm} (8)

This last equation is a manifestation of the gauge-gravity duality. In particular, it explicitly demonstrates that effective actions for conventional hadrons in $d$ dimensions can be generated from actions for bulk fields propagating in extra $d + 1$ dimensions. The impact of the extra-dimension is encoded in the hadronic mass squared $M^2$, which is the solution of the Schrödinger equation (6) for the KK profile in extra dimension $\Phi_n(z)$.

For the case of the vector field $V_M(x,z)$ we proceed by analogy. After straightforward algebra we derive the Schrödinger-type EOM for the KK mode and find the mass spectrum of vector mesons $M_{\delta 0}^2 = 4\kappa^2 \left( n + \frac{1}{2} + \frac{1}{2} \right)$.

We further consider bulk boson fields with higher values of $J \geq 2$. This problem, in the context of soft-wall models, has been considered before in Refs. [1, 4, 10, 1, 12]. In particular, it was shown that the soft-wall model reproduces the Regge-behavior of the mesonic mass spectrum $M_{\delta 0}^2 \sim n + J$. Here, extending our results for scalar and vector fields, we show that the bound-state problem is independent on the sign of the dilaton profile.

We describe a bosonic spin-$J$ field $\Phi_{M_1 \cdots M_J}(x,z)$ by a symmetric, traceless tensor, satisfying the conditions $\nabla_M \Phi_{M_1 \cdots M_J} = 0$, $g^{M_1 M_2} \Phi_{M_1 \cdots M_J} = 0$. The actions for the bulk field $\Phi_j$ with positive and negative dilatons are [4,5,8]

$$S_J^\pm = \frac{(-1)^J}{2} \int d^d x dz \frac{e^z}{z} \left[ g^{MN} g^{M_1 N_1} \cdots g^{M_J N_J} \nabla_M \Phi_{M_1 \cdots M_J}(x,z) \nabla_N \Phi_{M_1 \cdots M_J}(x,z) \right]$$

$$- \left( \mu^2 + \Delta V_j(z) \right) g^{M_1 N_1} \cdots g^{M_J N_J} \Phi_{M_1 \cdots M_J}(x,z) \Phi_{M_1 \cdots M_J}(x,z),$$  \hspace{1cm} (9)

where $V_j^+ = 0$ and $V_j^- = V_j$. Here $\nabla_M$ is the covariant derivative with respect to AdS coordinates.

In Refs. [4,5,10,11] higher spin fields have been considered in a “weak gravity” approximation, restricting the analysis to flat metric and therefore identifying the covariant derivative with the normal derivative (i.e. neglecting the affine connection). First, we review these results and then consider the general case with covariant derivatives. In the following we call the scenario with normal derivatives scenario I and the scenario with covariant derivatives scenario II.

In scenario I the bulk mass is given by [4,5,10] $\mu_J^2 R^2 = (\Delta - J)(\Delta + J - d)$, which is fixed by the behavior of bulk fields $\Phi_j$ near the ultraviolet boundary $z = 0$. The potential $\Delta V_j(z) = e^{-2\kappa z} \Delta U_j(z)$ is given by $\Delta U_j(z) = \phi^\dagger(z) + (d - 1 + 2J) \phi(z) A(z)$.

Notice that both quantities $\mu_J^2$ and $\Delta U_j(z)$ are generalizations of the scalar ($J = 0$) and vector ($J = 1$) cases considered before. In particular, they are related to those for the scalar field as follows: $\mu_J^2 R^2 = \mu_2^2 R^2 + J(d - J)$, $\Delta U_j(z) = \Delta U_0(z) - 2J \phi(z) A(z)$. As before the two actions can be reduced to the action with a dilaton hidden in an additional potential term, using the transformation $\Phi_j^\pm(x,z) = e^{\mp \varphi_{J/2}(z)} \psi_j(x,z)$. Then the action takes the form

$$S_J = \frac{(-1)^J}{2} \int d^d x dz \sqrt{g} \left[ g^{MN} g^{M_1 N_1} \cdots g^{M_J N_J} \partial_M \Phi_{M_1 \cdots M_J}(x,z) \partial_N \Phi_{M_1 \cdots M_J}(x,z) \right]$$

$$- \left( \mu_J^2 + V_j(z) \right) g^{M_1 N_1} \cdots g^{M_J N_J} \Phi_{M_1 \cdots M_J}(x,z) \Phi_{M_1 \cdots M_J}(x,z),$$  \hspace{1cm} (10)
where \( V_J(z) = e^{-2A(z)}U_J(z) \), and with the effective potential

\[
U_J(z) = \frac{1}{2}\Phi''(z) + \frac{1}{4}(\Phi'(z))^2 + \frac{d - 1 - 2J}{2}\Phi'(z)A'(z).
\]

(11)

This last expression is identical with the light-front effective potential found in Ref. [5] for \( d = 4 \) and arbitrary \( J \) [see Eq.(10)]; \( U_J(z) = \kappa^4 z^2 + 2\kappa^2 J \). Using standard algebra and restricting to the axial gauge \( \Phi_\tau \approx \Phi_\tau \approx 0 \), writing down the action in terms of fields with Lorentz indices and rescaling the fields by the boost (total angular momentum) factor \( e^{iA(z)} \Phi_\mu \rightarrow e^{iA(z)} \Phi_\mu \) we write down the action in the form

\[
S_J = \left( \frac{(-)^J}{2} \right) \int d^4z e^{B_0(z)} \left[ \partial_\mu \Phi_{\mu_1 \cdots \mu_J}(x) \partial^\mu \Phi_{\mu_1 \cdots \mu_J}(x) - \partial_\mu \Phi_{\mu_1 \cdots \mu_J}(x) \partial^\mu \Phi_{\mu_1 \cdots \mu_J}(x) \right] - \left( \frac{e^{2A(z)}}{\mu^2} + U_J(z) \right) \Phi_{\mu_1 \cdots \mu_J}(x) \Phi_{\mu_1 \cdots \mu_J}(x).
\]

(12)

Doing the KK expansion \( \Phi_{\mu_1 \cdots \mu_J}(x) = \sum_n \Phi_{\mu_1 \cdots \mu_J}(n) \Phi_n(z) \) and the substitution \( \Phi_n(z) = e^{-B_0(z)/2} \phi_n(z) \) we derive the Schrödinger-type EOM for \( \phi_n(z) \):

\[
\left[ - \frac{d^2}{dz^2} + \frac{4L^2 - 1}{4z^2} + U_J(z) \right] \phi_n(z) = M_{\mu_1 \cdots \mu_J}^2 \phi_n(z),
\]

(13)

where \( M_{\mu_1 \cdots \mu_J}^2 = 4\kappa^2 (n + \frac{J}{2} + \frac{3}{4}) \) is the mass spectrum of higher \( J \) fields, which generalizes our results for scalar and vector fields. At large values of \( J \) or \( L \) we reproduce the Regge behavior of the meson mass spectrum: \( M_{\mu_1 \cdots \mu_J}^2 \sim n + J \). Finally, using the KK expansion and the EOMs for the wave functions, we derive the \( d \)-dimensional action for mesons with total angular momentum \( J \geq 2 \) and masses \( M_{\mu_1 \cdots \mu_J}^2 \):

\[
S_J^{(d)} = \left( \frac{(-)^J}{2} \right) \sum_n \int d^d x \left[ \partial_\mu \Phi_{\mu_1 \cdots \mu_J}(x) \partial^\mu \Phi_{\mu_1 \cdots \mu_J}(x) - M_{\mu_1 \cdots \mu_J}^2 \Phi_{\mu_1 \cdots \mu_J}(x) \Phi_{\mu_1 \cdots \mu_J}(x) \right].
\]

(14)

Now we consider scenario II, i.e. without truncation of covariant derivatives. The gauge-invariant actions for the totally symmetric higher spin boson fields have been considered e.g. in Refs. [13]. In this case the bulk mass is fixed by gauge invariance, and given by \( \mu^2 R^2 = J^2 + J(d - 5) + 4 - 2d \). This mass leads to the following results for the scaling of the KK profiles: \( \Phi_n(z) \sim z^{2+J} \) at \( z \to 0 \), and their masses \( M_{\mu_1 \cdots \mu_J}^2 = 4\kappa^2 (n + J) \), which are acceptable only for the limiting cases \( J = L \) and \( J \to \infty \). Notice that the soft-wall actions [2] are obtained from gauge-invariant actions for totally symmetric higher spin boson fields [13] via the introduction of the dilaton field, which breaks conformal and gauge invariance. Therefore, it is not necessary to use the bulk mass constrained by gauge invariance. In particular, in order to get correct scaling the the KK profile \( \Phi_n(z) \sim z^{2+L} \) and their masses we should use \( \mu_J^2 R^2 = (\Delta - J)(\Delta + J - d) - J = \mu_0^2 R^2 + J(d - 1 - J) \). In this case scenario II is fully equivalent to scenario I.

3 Fermionic case

In the fermion case we first consider the low-lying \( J = 1/2 \) modes \( \Psi_{\pm}(x, z) \) (here the index \( \pm \) corresponds again to scenarios with positive/negative dilaton profiles, respectively.) The actions with positive and negative dilaton read [14, 15, 16]

\[
S_{1/2}^{\pm} = \int d^4x d\sqrt{g} e^{\mp \phi(z)} \left[ \frac{i}{2} \bar{\Psi}^\pm \gamma^\mu \Gamma^\mu \partial_M \Psi^\pm - \frac{i}{2} (\partial_M \Psi^\pm)^\dagger \Gamma^0 \gamma^\mu \Gamma^\mu \Psi^\pm - \bar{\Psi}^\pm \left( \mu + V_F(z) \right) \Psi^\pm \right],
\]

(15)

where \( \partial_M \) is the covariant derivative and \( \Gamma^\mu = (\gamma^\mu, -i\gamma^5) \) are the Dirac matrices.

The quantity \( \mu \) is the bulk fermion mass with \( m = \mu R = \Delta - d/2 \), where \( \Delta \) is the dimension of the baryon interpolating operator, which is related with the scaling dimension \( \tau = 3 + L \) as \( \Delta = \tau + 1/2 \). For \( J = 1/2 \) we have two baryon multiplets \( J^P = 1/2^+ \) for \( L = 0 \) and \( J^P = 1/2^- \) for \( L = 1 \). \( V_F(z) = \phi(z)/R \) is the dilaton field dependent effective potential. Its presence is necessary due to the following reason. When fermionic fields are rescaled \( \Psi^\pm(x, z) = e^{\mp \phi(z)/2} \Psi(x, z) \), the dilaton field is removed from the overall exponential. The form of the
potential $V_F(z)$ is constrained in order to get solutions of the EOMs for fermionic KK modes of left and right chirality, and the correct asymptotics of the nucleon electromagnetic form factors at large $Q^2$ [14, 15, 16].

The action in terms of a field with Lorentz indices is:

$$S_{1/2} = \int d^4x dz \sqrt{g} \Psi(x, z) \left[ i \partial \gamma^\mu \partial_\mu + \frac{d}{2} A'(z) \gamma^5 - \frac{e(z)}{R} \left( m + \varphi(z) \right) \right] \Psi(x, z),$$

(16)

where the Dirac field satisfies the following EOM [14, 15, 16]: $iz \gamma^\mu \partial_\mu - 4 \gamma^5 - m - \varphi(z) \Psi(x, z) = 0$. Based on these solutions the fermionic action should be extended by an extra term in the ultraviolet boundary (see details in Refs. [14, 15]). Here we review the derivation of the EOMs for the KK modes dual to the left- and right-chirality spinors, in the soft-wall model [14, 15, 16]. First we expand the fermion field in left- and right-chirality components: $\Psi(x, z) = \Psi_L(x, z) + \Psi_R(x, z)$. Then we perform a KK expansion for the $\Psi_L/R(x, z)$ fields: $\Psi^{n}_{L/R}(x, z) = \sum_n \Psi^{n}_{L/R}(x) F^{n}_{L/R}(z)$. The KK modes $F^{n}_{L/R}(z)$ satisfy the two coupled one-dimensional EOMs [14, 15, 16]:

$$\left[ \partial_z + \frac{e(z)}{R} \left( m + \varphi(z) \right) + \frac{d}{2} A'(z) \right] F^n_{L/R}(z) = \pm M_n F^n_{L/R}(z),$$

(17)

where $M_n$ is the mass of baryons with $J = 1/2$. After the substitution $F^{n}_{L/R}(z) = e^{-A(z)d/2} f^{n}_{L/R}(z)$ we derive decoupled Schrödinger-type EOM for $f^{n}_{L/R}(z)$

$$\left[ - \partial_z^2 + \kappa^2 z^2 + 2 \kappa \left( m + \frac{1}{2} \right) \right] f^n_{L/R}(z) = M_n^2 f^n_{L/R}(z),$$

(18)

where

$$f^n_{L}(z) = c_L \kappa^{L+3} \epsilon^{L+5/2} e^{-\kappa^2 z^2/2} L_{n}^{L+2}(\kappa z^2), \quad f^n_{R}(z) = c_R \kappa^{L+2} \epsilon^{L+3/2} e^{-\kappa^2 z^2/2} L_{n}^{L+1}(\kappa z^2)$$

(19)

and $M_n^2 = 4 \kappa^2 (n + L + 2)$ with $L = m - 3/2 = 0, 1$ for $J = 1/2$ fermions. Here $c_L = \sqrt{2} \Gamma(n + 1)/\Gamma(n + L + 3)$ and $c_R = \sqrt{2} \Gamma(n + 1)/\Gamma(n + L + 2)$. One can see that the functions $F^{n}_{L/R}(z)$ have the correct scaling behavior for small $z$: $F^{n}_{L}(z) \sim \epsilon^{3/2+L}$, $F^{n}_{R}(z) \sim \epsilon^{3/2+L}$ and vanish at large $z$ (confinement). As in the bosonic case, integration over the holographic coordinate $z$ gives a $d$-dimensional action for the fermion field $\Psi^n(x) = \Psi^n_L(x) + \Psi^n_R(x)$:

$$S_{1/2}^{(d)} = \sum_n \int d^4x \bar{\Psi}^n(x) [i \partial - M_n] \Psi^n(x).$$

(20)

Extension of our formalism for higher spin states $J$ is straightforward. In particular, the actions for higher spin fermions with positive and negative dilaton are written as

$$S^\pm_J = \int d^4x dz \sqrt{g} e^{\pm \varphi(z)} g^{K_{J-1/2}N_{J-1/2}} \left[ \frac{i}{2} \bar{\Psi}^\pm_{K_{J-1/2}N_{J-1/2}} \sigma^a M^a \sigma^b N^b \Psi^\pm_{K_{J-1/2}N_{J-1/2}} - \frac{i}{2} (\Theta_M \Psi^\pm_{K_{J-1/2}N_{J-1/2}})^T \Gamma^0 e^M \Gamma^a \Psi^\pm_{K_{J-1/2}N_{J-1/2}} - \Psi^\pm_{K_{J-1/2}N_{J-1/2}} \left( \mu + V_F(z) \right) \bar{\Psi}^\pm_{K_{J-1/2}N_{J-1/2}} \right].$$

(21)

As before, we remove the dilaton field from the exponential prefactor, perform the boost of the spin-tensor field and restrict ourselves to the axial gauge. Next, after a straightforward algebra (including KK expansion), we derive the same equation of motion for the KK profile and mass formula as for fermions with lower spins. The action for physical baryons with higher spins is written as

$$S^{(d)}_J = \sum_n \int d^4x \bar{\Psi}^{\mu_1...\mu_{J-1/2}}(x) \left[ i \partial - M_n \right] \Psi^{\mu_1...\mu_{J-1/2}}(x).$$

(22)

Therefore, the main difference between the bosonic and fermionic actions is that in the case of bosons the mass formula depends on the combination $(J + L)/2$, while in the baryon case it depends only on $L$. Also in the fermion case the dilaton prefactor and possible warping of conformal-invariant AdS metric can be easily absorbed in the field, without the generation of extra potential terms.
4 Conclusions

We performed a systematic analysis of extra-dimensional actions for bosons and fermions, which give rise to actions for observable hadrons. Masses are calculated analytically from Schrödinger type equations of motion with a potential which provides confinement of the Kaluza-Klein (KK) modes in extra $(d + 1)$ dimension. The tower of KK modes with radial quantum number $n$ and total angular momentum $J$ has direct correspondence to realistic mesons and baryons living in $d$ dimensions. For such correspondence the sign of the dilaton profile is irrelevant, because the exponential prefactor containing the dilaton is finally absorbed in the bulk fields. On the other hand, the sign of the dilaton profile becomes important for the definition/calculation of the bulk-to-boundary propagator — i.e. the Green function describing the evolution of bulk field from inside of AdS space to its ultraviolet boundary. The corresponding sign should be negative in order to fulfill certain constraints discussed recently in Refs. [7].

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