

A Potential for a Three-Field AdS/QCD Model

Sean P. Bartz^{*} and Joseph I. Kapusta[†]

February 7, 2014

Abstract

The Anti-de Sitter Space/Conformal Field Theory (AdS/CFT) correspondence may offer new and useful insights into the non-perturbative regime of strongly coupled gauge theories such as Quantum Chromodynamics (QCD). We present an AdS/CFT-inspired model that describes the spectra of light mesons. The conformal symmetry is broken by a background dilaton field, and chiral symmetry breaking and linear confinement are described by a chiral condensate field. These background fields, along with a background glueball condensate field, are derived from a potential. We describe the construction of the potential, and the calculation of the meson spectra, which match experimental data well. We also argue that the presence of the third background field is necessary to properly describe the meson spectra. The outlook for application of this model to finite temperature systems is also discussed.

^{*}E-mail: bartz@physics.umn.edu

[†]E-mail: kapusta@physics.umn.edu

1 Introduction

The Anti-de Sitter Space/Conformal Field Theory (AdS/CFT) correspondence is a useful mathematical tool for the analysis of strongly-coupled gauge theories. This correspondence establishes a connection between an n -dimensional Super-Yang Mills Theory and a weakly-coupled gravitational theory in $n + 1$ dimensions [1–3]. Calculations that are analytically intractable in the field theory can be related to results from the gravity theory using an effective dictionary developed over the past decade. Quantum chromodynamics (QCD) is a strongly-coupled gauge theory at hadronic scales, making it a candidate for the application of the gauge/gravity correspondence. It is not known whether a gravitational dual to QCD exists, but there has been much work on models that capture its key features. The bottom-up approach assumes the existence of such a dual, modeling features of QCD by an effective five-dimensional gravity theory. Linear confinement in QCD sets a scale that is encoded in a cut-off of the fifth dimension in the AdS/QCD model [4, 5]. So-called soft-wall models use a dilaton field as an effective cut-off to limit the penetration of the meson fields into the bulk [6]. The simplest soft-wall models use a quadratic dilaton to recover the linear Regge trajectories, while models that modify the UV behavior of the dilaton more accurately model the ground state masses [7–10].

The soft-wall models typically include at least two background fields: the aforementioned dilaton, and a chiral condensate field that corresponds to the chiral symmetry breaking in the gauge theory. These models use parametrizations for the background dilaton and chiral fields that are not derived as the solution to any equations of motion. A well-defined action would provide a set of background equations from which these fields can be derived, and may suggest how the model can be derived from a top-down approach. In addition, this action provides access to the thermal properties of the model through perturbation of the geometry [11–13].

In this paper, we expand upon previous work to find a suitable potential for the background fields of a soft-wall AdS/QCD model [13–20]. After demonstrating the limitations of models including a dilaton and chiral field alone, we suggest the inclusion of a background glueball field. We then construct a potential that satisfies the necessary UV and IR limits, and use this potential to generate numerically the background fields and calculate the resulting meson spectra.

2 Review and Motivation

We assume that four-dimensional QCD can be modeled by the following five-dimensional action, written in the string frame:

$$\mathcal{S} = \frac{1}{16\pi G_5} \int d^5x \sqrt{-g} [e^{-2\Phi} \mathcal{L}_{grav} + e^{-\Phi} \mathcal{L}_{meson}] , \quad (1)$$

where Φ is the dilaton and the metric is pure AdS, $g_{MN} = z^{-2} \eta_{MN}$, with the AdS curvature defined to be unity. The gravitational Lagrangian includes the dilaton and

chiral condensate background fields

$$\mathcal{L}_{grav} = R + 4\partial_M\Phi\partial^M\Phi - \frac{1}{2}\partial_M\chi\partial^M\chi - V(\Phi, \chi). \quad (2)$$

The scalar fields Φ, χ are dimensionless.

It is easier to search for the background fields in the Einstein frame, where the gravitational action takes the canonical form

$$\mathcal{S}_E = \frac{1}{16\pi G_5} \int d^5x \sqrt{-\tilde{g}} \left(\tilde{R} - \frac{1}{2}\partial_M\phi\partial^M\phi - \frac{1}{2}\partial_M\chi\partial^M\chi - \tilde{V}(\phi, \chi) \right). \quad (3)$$

The tilde distinguishes the two frames, with $\tilde{V} = e^{4\Phi/3}V$, and the dilaton is rescaled for a canonical action $\phi = \sqrt{8/3}\Phi$. The string and Einstein frame metrics are related by the conformal transformation

$$g_{MN} = e^{2\phi/\sqrt{6}} \tilde{g}_{MN}. \quad (4)$$

Previous work showed how to construct a potential for such a gravity-dilaton-chiral system, assuming that the fields have power-law behavior, which is accurate in both the UV and IR limits, [15]. One of the equations of motion is independent of the choice of potential,

$$\dot{\chi}^2 = \frac{\sqrt{6}}{z^2} \frac{d}{dz} (z^2 \dot{\phi}). \quad (5)$$

To obtain linear confinement, the dilaton should have quadratic behavior in the IR limit, $\phi(z) = \lambda z^2$. The chiral field should have linear behavior, $\chi(z) = Az$, where A sets the large- n mass splitting between the axial-vector and vector mesons in the model. This constant mass-splitting at large n occurs because of the non-restoration of chiral symmetry [21]. Inserting this into (5), we find that the chiral field behaves as

$$\chi(z) = 6^{3/4} \sqrt{\lambda} z, \quad (6)$$

which removes one of the independent parameters of the model in [22]. Using the phenomenological value of λ , which determines the slope of the Regge trajectories, we find a mass splitting that is much too large.

Because this problem arises in the equation that is independent of the potential, this issue cannot be resolved by the choice of potential in the two-field model. Models that derive the field behavior using the superpotential method suffer from the same problem.

To resolve this problem, we suggest to add an additional scalar field to the model, G , representing the glueball condensate. This field must be linear in the IR for linear confinement, and go as $G \sim z^4$ in the UV to match the operator dimension in the AdS/CFT dictionary.

It is noted that the model proposed by Huang and Li [18, 19] accurately models the non-restoration of chiral symmetry using a model with only two background fields, but their model differs from the work presented here in several respects. They allow the metric to be solved dynamically based upon the chosen parametrization for the background fields, while we maintain a metric that is strictly AdS. The IR behavior of

the chiral field in their model is set by examining the quark potential, and approaches a constant value, as opposed to the linear behavior described above. Finally, they do not determine an explicit form of the potential, which is the central goal of this work.

3 Setup

Consider the action in the Einstein frame for three fields: ϕ , χ and G representing the dilaton, a chiral field, and a glueball field with zero mass

$$\mathcal{S}_E = \frac{1}{16\pi G_5} \int d^5x \sqrt{-\tilde{g}} \left(\tilde{R} - \frac{1}{2} \partial_M \phi \partial^M \phi - \frac{1}{2} \partial_M \chi \partial^M \chi - \frac{1}{2} \partial_M G \partial^M G - \tilde{V}(\phi, \chi, G) \right). \quad (7)$$

To simplify the equations of motion, we use a transformed potential,

$$V = e^{-2\phi/\sqrt{6}} \tilde{V}. \quad (8)$$

It is noted that this is simply the potential in the string frame. We re-write the potential as

$$V = -12 + 4\sqrt{6}\phi + a_0\phi^2 - \frac{3}{2}\chi^2 + U. \quad (9)$$

Here U is more than quadratic in the fields. The dilaton mass is undetermined and is not connected to the dimension of the corresponding operator, as discussed in [15]. It is related to the parameter a_0 by $a_0 = \frac{1}{2} [(m_\phi L)^2 - 8]$. The potential should be an even function of χ .

The equations of motion can be written as

$$\dot{\chi}^2 + \dot{G}^2 = \frac{\sqrt{6}}{z^2} \frac{d}{dz} (z^2 \dot{\phi}) \quad (10)$$

$$U = \frac{1}{2} \sqrt{6} z^2 \ddot{\phi} - \frac{3}{2} (z \dot{\phi})^2 - 3\sqrt{6} z \dot{\phi} - 4\sqrt{6} \phi - a_0 \phi^2 + \frac{3}{2} \chi^2 \quad (11)$$

$$\frac{\partial U}{\partial \phi} = 3z \dot{\phi} - 2a_0 \phi \quad (12)$$

$$\frac{\partial U}{\partial \chi} = z^2 \ddot{\chi} - 3z \dot{\chi} \left(1 + \frac{z \dot{\phi}}{\sqrt{6}} \right) + 3\chi \quad (13)$$

$$\frac{\partial U}{\partial G} = z^2 \ddot{G} - 3z \dot{G} \left(1 + \frac{z \dot{\phi}}{\sqrt{6}} \right) \quad (14)$$

We assume that the potential has no explicit dependence on the coordinate z , so the equations 12-14 are not independent, and we can eliminate one.

3.1 IR Limit

The requirement of linear confinement requires a solution in the large z limit of the form

$$\phi = \lambda z^2 \quad (15)$$

$$\chi = Az \quad (16)$$

$$G = Bz. \quad (17)$$

Substitution into (10) gives

$$A^2 + B^2 = 6\sqrt{6}\lambda \quad (18)$$

The λ is fixed by the slope of the linear trajectory and A is fixed by the axial-vector – vector mass difference. It is useful to write these as

$$\begin{aligned} A &= 6^{3/4}\sqrt{\lambda}\cos\theta \\ B &= 6^{3/4}\sqrt{\lambda}\sin\theta, \end{aligned} \quad (19)$$

where θ now becomes the parameter controlling the axial-vector – vector mass splitting. Inserting (17) into (11-14) suggests the following terms in our ansatz for the potential

$$U = a_1\phi\chi^2 + a_2\phi G^2 + a_3\chi^4 + a_4G^4 + a_5\chi^2G^2 + a_6G^2\tanh(g\phi). \quad (20)$$

We see that there must be a G^2 term in the IR limit, but this is forbidden in the weak-field limit because the glueball condensate field is massless. To circumvent this, we propose the term $G^2\tanh(g\phi)$ with $g > 0$. In the weak field limit this goes to $g\phi G^2$, which is acceptable. The \tanh is suggested by 8, and it suggests a rapid exponential transition from the weak field to the strong field limits that is supported by phenomenology. By substitution one finds the following constraints on the parameters:

$$\begin{aligned} U &\rightarrow 6 + a_0 + 6\sqrt{6}(\cos^2\theta a_1 + \sin^2\theta a_2) \\ &+ 6^3(\cos^4\theta a_3 + \sin^4\theta a_4 + \cos^2\theta \sin^2\theta a_5) = 0 \end{aligned} \quad (21)$$

$$\frac{\partial U}{\partial \chi} \rightarrow 2a_1 + 24\sqrt{6}\cos^2\theta a_3 + 12\sqrt{6}\sin^2\theta a_5 + \sqrt{6} = 0 \quad (22)$$

$$\frac{\partial U}{\partial G} \rightarrow 2a_2 + 24\sqrt{6}\sin^2\theta a_4 + 12\sqrt{6}\cos^2\theta a_5 + \sqrt{6} = 0 \quad (23)$$

$$\frac{\partial U}{\partial G} \rightarrow a_6 = -\frac{3}{2} \quad (24)$$

We have chosen to exclude (12) because it is not independent. The parameter a_6 is determined, and the others will be determined by an examination of the UV limit.

3.2 UV Limit

Next we look for a solution in the small z limit. The AdS/CFT dictionary dictates that the leading-order UV behavior of the chiral and glueball condensate fields is determined by their dimension. Note also that we are working in the chiral limit where the quark mass is zero. We start by examining only the leading-order terms

$$\chi = \Sigma_0 z^3 \quad (25)$$

$$G = G_0 z^4. \quad (26)$$

Substitution into (10) and imposing the boundary condition $\phi(0) = 0$ gives

$$\phi = \frac{\sqrt{6}}{28} \Sigma_0^2 z^6 + \frac{\sqrt{6}}{27} G_0^2 z^8 \quad (27)$$

Using only this leading-order behavior in (11-14), the system of equations is inconsistent, as there are more equations from matching powers of z than unknown parameters.

To solve this problem, try adding a term $\Sigma_n z^n$ to χ . Substituting into (10) and keeping only the lowest-order cross-term we find the additional term in ϕ

$$\Delta\phi = \frac{\sqrt{6}n\Sigma_0\Sigma_n}{(n+4)(n+3)} z^{n+3} \quad (28)$$

From (11) we find that

$$U = -\frac{3}{2}(z\dot{\phi})^2 - a_0\phi^2 + 3\frac{n^3 - 13n + 12}{(n+4)(n+3)}\Sigma_0\Sigma_n z^{n+3} \quad (29)$$

Since the ϕ^2 terms start out as z^{12} , z^{14} , z^{16} , and so do the terms in the potential, the n can only take the values 9, 11, etc. This term contributes only to the equation for $\partial U/\partial\chi$.

$$\frac{\partial U}{\partial\chi} = -9\Sigma_0 \left(\frac{3}{14}\Sigma_0^2 + \frac{8}{27}G_0^2 z^2 \right) z^9 + (n-3)(n-1)\Sigma_n z^n \quad (30)$$

By power counting both $n = 9$ and $n = 11$ can contribute.

There could also be higher order terms in G such as $G_m z^m$. This leads to the additional term in ϕ

$$\Delta\phi = \frac{8mG_0G_m}{\sqrt{6}(m+5)(m+4)} z^{m+4} \quad (31)$$

It contributes to the equation for $\partial U/\partial G$ as

$$\frac{\partial U}{\partial G} = -12G_0 \left(\frac{3}{14}\Sigma_0^2 + \frac{8}{27}G_0^2 z^2 \right) z^{10} + m(m-4)G_n z^m \quad (32)$$

The choice $m = 8$ is not possible as there is no term of the same order to balance it. Terms with $m = 10$ and $m = 12$ are possible. These new terms cannot affect the

equation for $\partial U/\partial \phi$ nor can they contribute to the equation for $\partial U/\partial \chi$. Considering higher order terms in both χ and G leads to

$$U = -\frac{3}{2}(z\dot{\phi})^2 - a_0\phi^2 + 3\frac{n^3 - 13n + 12}{(n+4)(n+3)}\Sigma_0\Sigma_n z^{n+3} + \frac{4m(m-4)}{m+4}G_0G_m z^{m+4} \quad (33)$$

The appearance of these terms can be understood by writing the following schematic expansions.

$$\begin{aligned}\chi &\sim \Sigma_0 z^3 + \Sigma_0^3 z^9 + G_0^2 \Sigma_0 z^{11} + \dots \\ G &\sim G_0 z^4 + \Sigma_0^2 G_0 z^{10} + G_0^3 z^{12} + \dots\end{aligned}$$

That is, χ is an odd function of Σ_0 and G is an odd function of G_0 . These are the symmetries in the equations of motion. They also follow the spirit of the AdS/CFT correspondence in terms of the dimensionality of the operators and the powers of z .

Including now $m = 10$ and 12 , and $n = 9$ and 11 , we have the following set of equations in the small z limit:

$$\begin{aligned}U_{\text{LHS}} &= 3\Sigma_0^4 z^{12} \left[4\frac{\Sigma_9}{\Sigma_0^3} - \frac{(54 + a_0)}{2^3 \cdot 7^2} \right] \\ &+ \frac{1}{7}\Sigma_0^2 G_0^2 z^{14} \left[120\frac{G_{10}}{\Sigma_0^2 G_0} + 120\frac{\Sigma_{11}}{\Sigma_0 G_0^2} - \frac{(72 + a_0)}{9} \right] \\ &+ 2G_0^4 z^{16} \left[12\frac{G_{12}}{G_0^3} - \frac{(96 + a_0)}{3^5} \right]\end{aligned} \quad (34)$$

$$\begin{aligned}U_{\text{RHS}} &= \Sigma_0^4 z^{12} \left[\frac{\sqrt{6}}{28} a_1 + a_3 \right] \\ &+ \Sigma_0^2 G_0^2 z^{14} \left[\frac{\sqrt{6}}{27} a_1 + \frac{\sqrt{6}}{28} (a_2 + ga_6) + a_5 \right] \\ &+ G_0^4 z^{16} \left[\frac{\sqrt{6}}{27} (a_2 + ga_6) + a_4 \right]\end{aligned} \quad (35)$$

$$\left(\frac{\partial U}{\partial \chi} \right)_{\text{LHS}} = 3\Sigma_0^3 z^9 \left[-\frac{9}{14} + 16\frac{\Sigma_9}{\Sigma_0^3} \right] + 8\Sigma_0 G_0^2 z^{11} \left[-\frac{1}{3} + 10\frac{\Sigma_{11}}{\Sigma_0 G_0^2} \right] \quad (36)$$

$$\left(\frac{\partial U}{\partial \chi} \right)_{\text{RHS}} = \Sigma_0^3 z^9 \left[\frac{\sqrt{6}}{14} a_1 + 4a_3 \right] + \Sigma_0 G_0^2 z^{11} \left[\frac{2\sqrt{6}}{27} a_1 + 2a_5 \right] \quad (37)$$

$$\left(\frac{\partial U}{\partial G} \right)_{\text{LHS}} = 6\Sigma_0^2 G_0 z^{10} \left[-\frac{3}{7} + 10\frac{G_{10}}{\Sigma_0^2 G_0} \right] + 32G_0^3 z^{12} \left[-\frac{1}{9} + 3\frac{G_{12}}{G_0^3} \right] \quad (38)$$

$$\left(\frac{\partial U}{\partial G} \right)_{\text{RHS}} = \Sigma_0^2 G_0 z^{10} \left[\frac{\sqrt{6}}{14} (a_2 + ga_6) + 2a_5 \right] \quad (39)$$

$$+ G_0^3 z^{12} \left[\frac{2\sqrt{6}}{27} (a_2 + ga_6) + 4a_4 \right] \quad (40)$$

Altogether, from both the UV and IR limits, there are ten independent equations for the twelve parameters $a_0 - a_6$, Σ_9 , Σ_{11} , G_{10} , G_{12} , and g . We take g as the free parameter to use as the rate of transition from small z to large z . The parameters in the potential are found to be

$$a_0 = \frac{3}{2} \frac{1}{6 + \sin^2 \theta} \left[120 + 62 \sin^2 \theta + 63\sqrt{6}g \sin^2 \theta \right] \quad (41)$$

$$a_1 = -\frac{3\sqrt{6}}{4} \frac{1}{6 + \sin^2 \theta} \left[12 + 8 \sin^2 \theta + 9\sqrt{6}g \sin^2 \theta \right] \quad (42)$$

$$a_2 = -\frac{\sqrt{6}}{4} \frac{1}{6 + \sin^2 \theta} \left[32 + 24 \sin^2 \theta + 3\sqrt{6}g(9 \sin^2 \theta - 2) \right] \quad (43)$$

$$2a_3 \cos^2 \theta + a_5 \sin^2 \theta = \frac{1}{24} \frac{1}{6 + \sin^2 \theta} \left[24 + 22 \sin^2 \theta + 27\sqrt{6}g \sin^2 \theta \right] \quad (44)$$

$$2a_4 \sin^2 \theta + a_5 \cos^2 \theta = \frac{1}{24} \frac{1}{6 + \sin^2 \theta} \left[20 + 22 \sin^2 \theta + 3\sqrt{6}g(9 \sin^2 \theta - 2) \right] \quad (45)$$

$$a_6 = -\frac{3}{2} \quad (46)$$

The coefficients a_0 , a_1 , a_2 and a_6 are determined, while there are two equations for the three coefficients a_3 , a_4 and a_5 . That leaves a_5 as a free parameter, to be fit numerically, along with g , θ , G_0 , Σ , and λ .

4 Numerical Solution

Using the above potential, we seek a numerical solution that simultaneously satisfies the UV and IR limits. We use equations (10, 13, 14), which allows for an additional term in the potential, ΔU , such that

$$\frac{\partial}{\partial \chi} \Delta U = \frac{\partial}{\partial G} \Delta U = 0, \quad (47)$$

which will be determined from the numerical solution.

The differential equations represent a stiff system, and treatment of the problem as an initial value problem leads to numerical instabilities. We treat it instead as a boundary value problem, using Dirichlet boundary conditions at both boundaries. A relaxation method is used in combination with input approximations for the background fields, which are then iterated to find a stable solution to the system with the given boundary conditions. Because the system is nonlinear, the solution found is not guaranteed to be unique.

The IR boundary is chosen to be sufficiently large to capture the infrared behavior and to give accurate Regge behavior large- n radial excitations of the mesons. The UV boundary should approach zero, but it cannot reach zero because of the singularity in the equations of motion. This becomes a problem because equation (10) allows constant and divergent terms

$$\Delta \phi(z) = c_1 + c_2 z^{-1}. \quad (48)$$

Symbolically, these terms can be set to zero by enforcing the Dirichlet boundary condition $\phi(0) = 0$, but this is impossible to enforce numerically. Creative choice of UV boundary conditions can eliminate one, but not both of these unwanted terms without affecting the chiral and glueball fields. The behavior of the numerical solutions suggests that the desired UV behavior is an unstable solution to the equations, and therefore difficult or impossible to find with this iterative method.

As an alternative to direct solution, we parameterize the fields as follows:

$$\Psi(z) = \psi(z)_{UV} f(z) + \psi(z)_{IR} (1 - f(z)), \quad (49)$$

where $f(z)$ is some function that transitions smoothly from 1 at small values of z to 0 at large z , and $\psi(z)_{xx}$ represents the known UV and IR limits of the fields ϕ , χ , and G . The switching functions f need not be the same for each field. We choose

$$f_\phi(z) = e^{-(\beta_1 z)^{10}} \quad (50)$$

$$f_\chi(z) = e^{-(\beta_2 z)^4} \quad (51)$$

$$f_G(z) = e^{-(\beta_3 z)^5}. \quad (52)$$

The powers of the exponential are chosen to be greater than the known power-law behavior of the fields in the UV limit, so as to not interfere with this behavior. The β_i will be set by numerical fitting.

In all, we have nine parameters to be set numerically. The first constraint is to obtain the best global visual fit to the meson spectrum. We do not simply do a chi-squared fitting to the experimental data because the measurement error for the ground state ρ meson is so much smaller than for the others that this would effectively act as the only constraint. Secondly, we seek to minimize the error in the finite-difference approximations to equations 10, 13, and 14.

Three of the parameters are most phenomenologically relevant: λ , which controls the slope of the meson spectra in the large- n limit; θ , which controls the mass splitting between the a_1 and ρ mesons, and β_2 , which controls the location of the “bend” in the a_1 spectrum. For each set of these parameters, the other parameters are set by a routine that minimizes the error in the equations of motion. The parameters found are shown in Table 1.

The background fields that are obtained from this analysis are shown in Figures 1-3. The asymptotic power-law behavior of the fields is evident in the linear portions of the log-log scale plots shown. The “transition” behavior is most evident in the dilaton because of the large value of β_1 , which controls the value of z at which the field transitions from the UV limit to the IR limit.

We now analyze the “extra” term in the potential, ΔU . We obtain this term numerically by subtracting the right-hand side of 11 from its left-hand side. This term can be approximated numerically as a function of the dilaton field,

$$\Delta U(\phi) = \alpha \phi e^{-(\phi-b)^2/c}. \quad (53)$$

The best-fit values for these parameters are shown in Table 4.

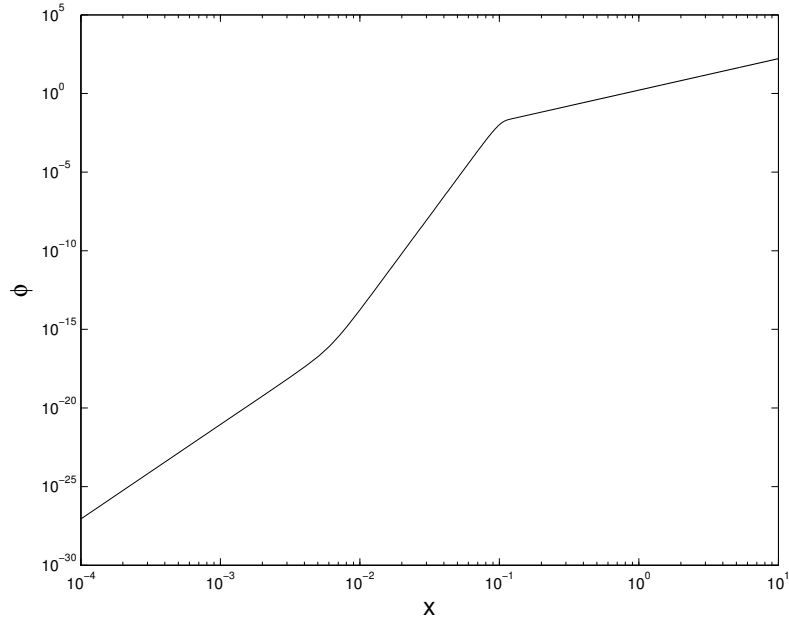


Figure 1: A plot of the dilaton field Φ generated by the parameterization (51). The UV and IR asymptotic behavior is apparent, with a transition region between them. The coordinate x is a dimensionless re-scaling of the conformal coordinate, $x = \sqrt{\lambda}z$.

5 Meson Spectra

We use the following meson Lagrangian in the string frame action (1),

$$\mathcal{L}_{meson} = |DX|^2 + m_X^2 |X|^2 + \frac{1}{2g_5^2} (F_A^2 + F_V^2). \quad (54)$$

The covariant derivative is defined as $D_M = \partial_M + i[V_M, X] - i\{A_M, X\}$. The 2×2 scalar field X contains the scalar and pseudoscalar fields (S, π) , as well as a non-trivial vacuum expectation value, dual to the $q\bar{q}$ operator, which we have identified with the field $\chi(z)$. The AdS/CFT dictionary sets the mass for the fields according to

$\lambda^{1/2}$	430 MeV	β_1	20.5 GeV
$\Sigma^{1/3}$	200 MeV	β_2	344 MeV
$G_0^{1/4}$	523 MeV	β_3	550 MeV
θ	1.47	a_5	1.18
g	1.0		

Table 1: Best fit parameters for the phenomenological model. The parameters λ , θ , and β_2 are chosen for the best visual fit to the ρ and a_1 data, with the rest set by minimizing the error in the equations of motion (10), (13-14).

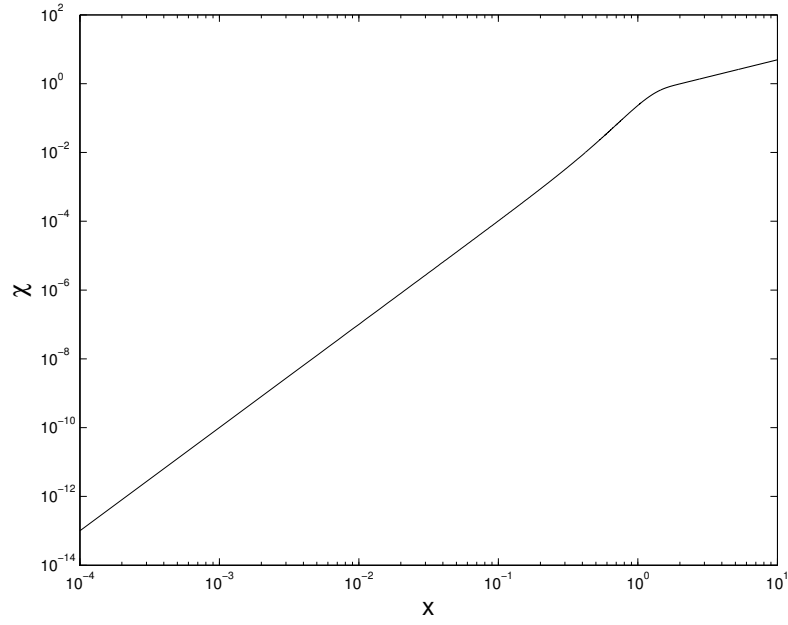


Figure 2: A plot of the chiral field χ generated by the parameterization (??). The UV and IR asymptotic behavior is apparent, with a rapid transition between them. The coordinate x is a dimensionless re-scaling of the conformal coordinate, $x = \sqrt{\lambda}z$.

α	-72.29
b	0.8326
c	0.5673

Table 2: The dimensionless parameters for the fitting to ΔU

the dimension of the dual operator,

$$m^2 L^2 = \Delta(\Delta - 4). \quad (55)$$

The dimension of the $q\bar{q}$ operator is 3, so $m_X^2 = -3/L^2$.

We find the equations of motion for the various meson fields by varying the meson action. For the vector and axial-vector fields, we assume that the Kaluza-Klein modes are separable from the 4D parts of the fields. The equation of motion in the axial gauge $\Psi_5 = 0$ is given by

$$-\partial_z^2 \Psi_n + \omega' \partial_z \Psi_n + M_\Psi^2(z) \Psi_n = m_{\Psi_n}^2 \Psi_n, \quad (56)$$

where $\omega = \Phi(z) + \log(z)$. The z -dependent mass term coefficient $M_V^2 = 0$ for the vector field, and

$$M_A^2 = \frac{g_5^2 L^2 \chi^2}{z^2} \quad (57)$$

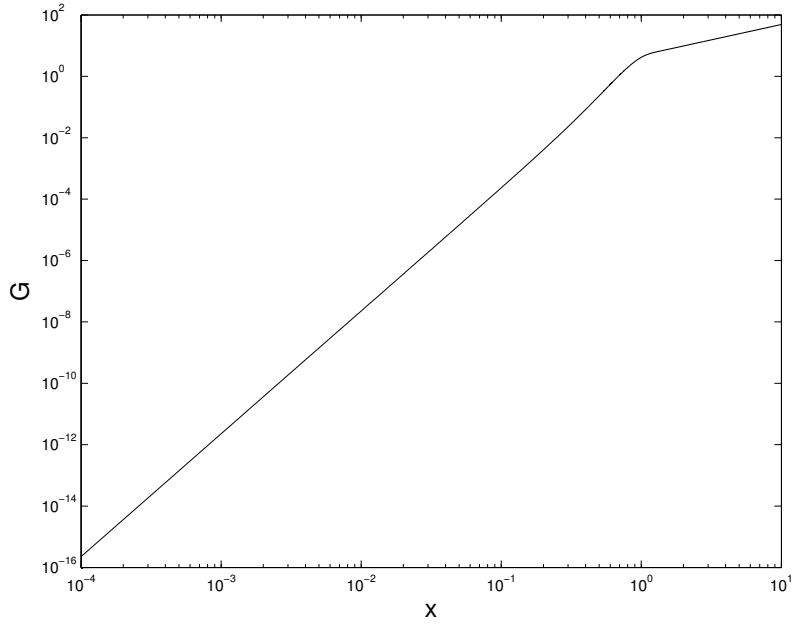


Figure 3: A plot of the glueball field G generated by the parameterization (52). The UV and IR asymptotic behavior is apparent, with a rapid transition between them. The coordinate x is a dimensionless re-scaling of the conformal coordinate, $x = \sqrt{\lambda}z$.

for the axial field. The equation can be put in the Schrödinger form with the substitution $\Psi_n = e^{\omega/2}\psi_n$, resulting in

$$-\partial_z^2 \psi_n + \left(\frac{1}{4}\omega'^2 - \frac{1}{2}\omega'' + M_\psi^2 \right) \psi_n = m_{\Psi_n}^2 \psi_n. \quad (58)$$

These equations are analytically solvable in the IR limit, but full analysis requires the use of a numerical shooting method to find the mass eigenvalues. This model finds a better phenomenological fit than the results presented in [22], particularly for the ground state ρ meson, as shown in Figure 6. The scalar mesons are expected to mix with the scalar glueball field of this model, and this analysis is not performed here.

6 Conclusion

We discussed the construction of a potential for the background fields of a soft-wall AdS/QCD model. We showed the limitation of a model that contains only the dilaton and chiral condensate fields, and suggested a solution by adding a glueball condensate to the model. We analytically constructed a general potential $U(\phi, \chi, G)$ that recovers the necessary asymptotic behavior of the background fields. Using this as a basis, we numerically constructed a potential that solves the selected background equations to within an accuracy of 10^{-4} . There is an additional allowed term in the potential, $\Delta U(\phi)$ that does not affect the equations that were used in the numerical

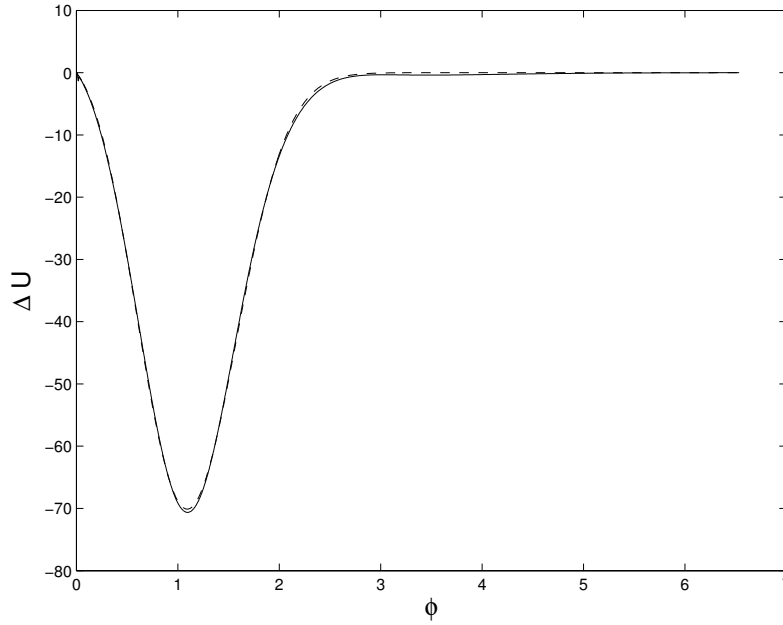


Figure 4: Plot of the “extra” term in the potential, $\Delta U(\phi)$. The solid line represents the numerical result, while the dashed line is the fitting of (53) using the parameters of Table 4.

procedure. This term was found numerically, and fit as a function of the dilaton alone. These background fields can then be used to find the spectra of the light vector and axial-vector mesons. The model shows good phenomenological agreement with the experimental data for the meson spectra.

The potential as constructed is not guaranteed to be unique. If a different set of the background equations is chosen, the extra term would be expressed as a function of fields other than the dilaton. The parameterization in (51-52) could also be chosen differently, resulting in a different potential but making little difference to the resulting meson spectra. Finally, terms can be added that do not affect the equations of motion at all, i.e. terms such that

$$\Delta U = \Delta \frac{\partial U}{\partial \phi} = \Delta \frac{\partial U}{\partial \chi} = \Delta \frac{\partial U}{\partial G} = 0. \quad (59)$$

This work demonstrates the construction of a potential for the background fields of a soft-wall AdS/QCD model that captures several key features of QCD observed through meson spectra. The radially excited states of the light mesons have linear Regge trajectories. Chiral symmetry is not restored for highly-excited mesons, as seen in the constant mass-splitting of the vector and axial-vector mesons.

Future improvements to this model will include incorporating the quark mass by adding a linear term to the UV limit of the chiral condensate field. The pseudo scalar and scalar f_0 meson spectrum should be explored, which will require analyzing how the equations of motion mix with the scalar potential. This potential also opens the

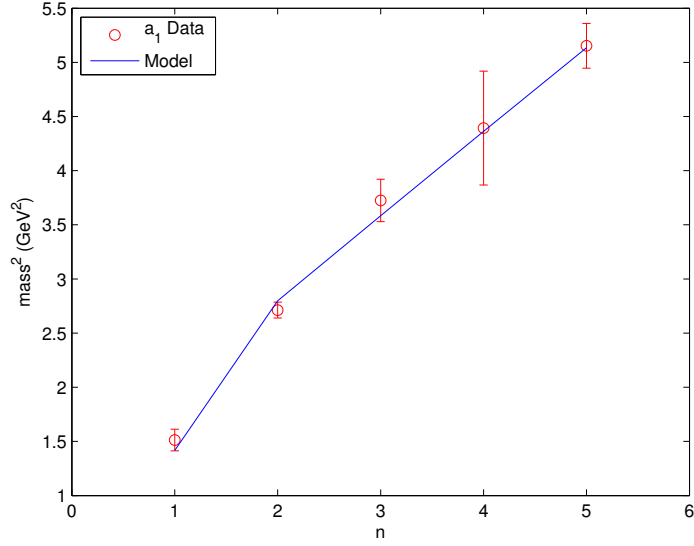


Figure 5: Comparison of the predicted mass eigenvalues for the axial-vector sector with the experimental a_1 meson spectrum [23].

possibility of exploring the thermal properties of a model that has the correct chiral symmetry breaking behavior.

Acknowledgments

SB thanks Tony Gherghetta and Aleksey Cherman for useful discussions. This research is supported in part by the Department of Energy Office of Science Graduate Fellowship Program (DOE SCGF), made possible in part by the American Recovery and Reinvestment Act of 2009, administered by ORISE-ORAU under contract no. DE-AC05-06OR23100.

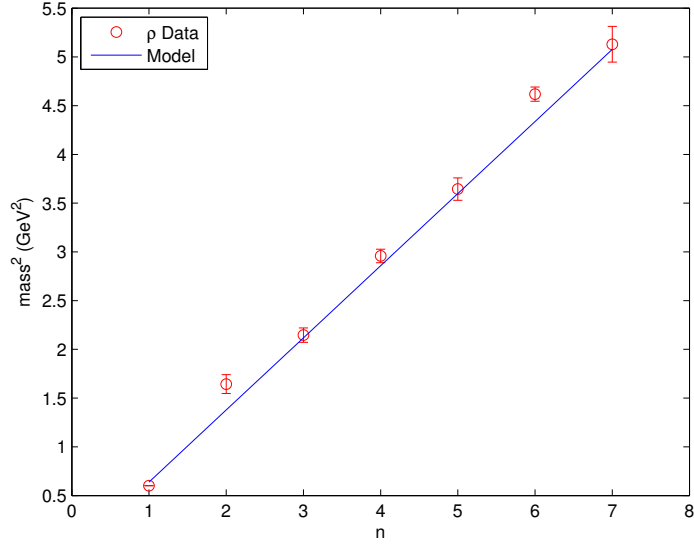


Figure 6: Comparison of the predicted mass eigenvalues for the vector sector with the experimental ρ meson spectrum [23].

References

- [1] J. M. Maldacena, *Advances in Theoretical Mathematical Physics* **2**, 20 (1997), 9711200.
- [2] S. Gubser, I. Klebanov, and A. Polyakov, *Physics Letters B* **428**, 105 (1998), 9802109.
- [3] E. Witten, *Adv.Theor.Math.Phys.* **2**, 253 (1998), 9802150.
- [4] J. Erlich, E. Katz, D. T. Son, and M. A. Stephanov, *Physical Review Letters* **95** (2005), 0501128.
- [5] L. Da Rold and A. Pomarol, *Nuclear Physics B* **721**, 79 (2005), 0501218.
- [6] A. Karch, E. Katz, D. Son, and M. Stephanov, *Physical Review D* **74**, 015005 (2006), 0602229.
- [7] T. Gherghetta, J. I. Kapusta, and T. M. Kelley, *Physical Review D* **79**, 076003+ (2009).
- [8] T. M. Kelley, S. P. Bartz, and J. I. Kapusta, *Physical Review D* **83**, 016002 (2011), 1009.3009.
- [9] P. Colangelo, F. De Fazio, F. Giannuzzi, F. Jugeau, and S. Nicotri, *Physical Review D* **78**, 17 (2008), 0807.1054.
- [10] L.-X. Cui, Z. Fang, and Y.-L. Wu, p. 15 (2013), 1310.6487.
- [11] C. Herzog, *Physical Review Letters* **98**, 091601 (2007), 0608151.

- [12] C. Bayona, H. Boschi-Filho, N. Braga, and L. Zayas, *Physical Review D* **77**, 046002 (2008), 0705.1529.
- [13] U. Gürsoy, E. Kiritsis, L. Mazzanti, and F. Nitti, *Physical Review Letters* **101**, 181601 (2008), 0804.0899.
- [14] B. Batell and T. Gherghetta, *Physical Review D* **78**, 026002+ (2008).
- [15] J. I. Kapusta and T. Springer, *Physical Review D* **81**, 086009 (2010).
- [16] U. Gürsoy, E. Kiritsis, and F. Nitti, *Journal of High Energy Physics* **2008**, 019 (2008), 0707.1349.
- [17] C. Csáki and M. Reece, *Journal of High Energy Physics* **2007**, 062 (2007), 0608266.
- [18] D. Li and M. Huang, *Journal of High Energy Physics* **2013**, 88 (2013), 1303.6929.
- [19] D. Li, M. Huang, and Q.-S. Yan, *The European Physical Journal C* **73**, 2615 (2013), 1206.2824.
- [20] S. He, S.-Y. Wu, Y. Yang, and P.-H. Yuan, p. 18 (2013), 1301.0385.
- [21] M. Shifman and A. Vainshtein, *Physical Review D* **77**, 034002 (2008).
- [22] T. Gherghetta, J. Kapusta, and T. Kelley, *Physical Review D* **79**, 076003 (2009), 0902.1998.
- [23] J. Beringer *et al.*, *Physical Review D* **86**, 010001 (2012).