

**Meson Spectra from a Dynamical Three-Field Model of  
AdS/QCD**

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# Dedication

To my parents, Larry and Colleen, my first teachers. To my sister, Haley, my first student. To my wife, Alicia, my greatest support.

## Abstract

Gauge/gravity dualities are a tool that allow the analytic analysis of strongly-coupled gauge theories. The Anti-de Sitter Space/Conformal Field Theory conjecture posits a duality between ten-dimensional string theory and a super Yang-Mills theory. A phenomenologically-motivated modification of this correspondence is known as AdS/QCD, a duality between strongly-coupled QCD-like theories and weakly-coupled gravitational theories in an additional dimension. QCD is not scale-invariant, so the dual theory must be modified in the conformal dimension to reflect this.

This thesis examines “soft-wall” models of AdS/QCD, wherein the conformal symmetry is broken by a field known as a dilaton. The dynamics of the dilaton and other background fields are examined, and a potential for these fields is determined. The background fields are numerically derived from this potential and used in the calculation of meson spectra, which match well to experiment.

The work presented here is based upon previously-published work by the author:

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# Chapter 1

## Introduction

Each of the four fundamental forces of the universe is described by a particular theory. Gravity is understood through the techniques of general relativity, while electromagnetism and the weak nuclear force are united in the electroweak theory of the standard model. Quantum chromodynamics (QCD) describes the behavior of the strong nuclear force, which is the subject of this thesis.

The strong nuclear force binds together the constituents of nuclear matter. The fundamental constituents of the theory are quarks and the force-carrying particles known as gluons. The strength of an interaction is characterized by a parameter known as the coupling constant. Traditional perturbation techniques involve a series expansion in this constant, which works well when the coupling constant is small. When the coupling constant is large, each successive term in this expansion is larger than the preceding one, so the perturbation expansion is not useful. The value of the coupling constant of QCD varies with energy scale. At high energies, the coupling constant is small, but the strength of the interaction grows at lower energies. At these energies, quarks and gluons are confined within particles known as hadrons, and investigating their dynamics requires new techniques.

Beginning in the late 1990's, new techniques were proposed to tackle these non-perturbative problems. In the framework of string theory it was conjectured that a duality exists between strongly-coupled gauge theories and weakly-coupled gravitational theories. Calculations that are analytically intractable in the field theory can be related to more easily calculated results from the gravity theory. These models are known as

gauge/gravity dualities. One proposed gauge/gravity model is the AdS/CFT (Anti-de Sitter space/Conformal Field Theory) correspondence, which relates certain strongly-coupled field theories to weakly-coupled gravitational theories with an extra dimension. The potential application of the AdS/CFT correspondence to non-perturbative aspects of QCD was quickly recognized, although there are some assumptions in the correspondence that do not apply to QCD. While a gravitational dual to QCD has not been discovered, there has been much success over the past fifteen years in developing five-dimensional models that capture key features of hadron phenomenology. These effective models are known as AdS/QCD.

This thesis is organized as follows

- In Chapter 3, we describe previous work done on soft-wall AdS/QCD models. We cover the application of the model to the calculation of meson spectra and the modeling of chiral symmetry breaking.
- In Chapter 3, we detail the construction of a three-field dynamical AdS/QCD model. The need for a third background field is motivated by the analysis of chiral symmetry breaking. We derive the equations of motion for the background fields, and construct a potential with the correct behavior. We describe the numerical techniques needed to derive the form of the background fields from this potential.
- In Chapter 4, we present the meson spectra that result from the dynamical AdS/QCD model. We describe the method for calculating the meson spectra and compare the model to experimental results.

## Chapter 2

# Background

In this chapter, we cover the background of gauge/gravity dualities and the AdS/CFT correspondence. Our focus is on the phenomenological applications of these models, so many of the technical details of the correspondence will not be presented here. A variety of review articles cover the topic in more mathematical detail [?]. Once we have motivated the AdS/CFT correspondence, we will review the applications of the duality to hadronic physics. We discuss the two major approaches to alter the conformal field theory to more closely resemble quantum chromodynamics. We also introduce a scheme to put these models on a more consistent theoretical basis.

### 2.1 What is AdS/CFT?

The general principle of dualities is to describe the same physical systems with two different but equivalent theories. The goal of this approach is to find that one of the theories is more analytically tractable than the other, and therefore more useful in describing physical phenomena. The original work on the AdS/CFT correspondence began with a duality between ten-dimensional string theory on  $AdS_5 \times S^5$  and  $\mathcal{N} = 4$  Super Yang-Mills Theory. The details of the string theory are not crucial to the understanding of the derivation of the correspondence. The constituents of the theory are extended objects known as  $p$ -branes, where  $p$  refers to the spatial dimension of the object. One-dimensional  $p$ -branes are more commonly known as strings, which can be either closed or open. Closed strings propagate freely through space, while

open strings have their endpoints on  $p$ -dimensional Dirichlet-branes (D-branes). These strings are characterized by their length  $l_s$  and the coupling constant  $g_s$  that describes the strength of their interactions. These parameters are related to the  $D$ -dimensional Newton's constant

$$G_D \sim g_s^2 l_s^{D-2}. \quad (2.1)$$

The string length can also be related to the Regge slope parameter

$$\alpha' = l_s^2. \quad (2.2)$$

We motivate the AdS/CFT correspondence by examining a stack of  $N$  D3 branes as a background for a type IIB string theory. We consider the dynamics of this system by taking the limits of low energy and strong coupling. The order in which these limits are taken determines the appearance of the result, yielding the two sides of the duality.

Taking the low-energy limit first gives an  $SU(N)$  gauge theory,  $\mathcal{N} = 4$  Super Yang-Mills (SYM) Theory.<sup>1</sup> The gauge coupling in this theory is related to the string coupling by

$$g_{YM}^2 = 4\pi g_s. \quad (2.3)$$

Assuming the string length is small, any interactions involving the open strings in the bulk (the region away from the branes) can be ignored. We are now free to make the assumption of strong coupling in the SYM theory. There are now two decoupled components of the low-energy system: the string theory governing the open strings in the bulk, and the strongly-coupled SYM theory on the branes.

Now, let us take the strong coupling limit first. In this instance, we have a large number of coincident D-branes resulting in a large energy density, and we must use general relativity to take into account its effect on the curvature of spacetime. The classical metric that solves the supergravity equations is

$$ds^2 = \frac{1}{\sqrt{1 + \frac{L^4}{r^4}}} (-dt^2 + d\vec{x}^2) + \sqrt{1 + \frac{L^4}{r^4}} (dr^2 + r^2 d\Omega_5), \quad (2.4)$$

where the curvature radius  $L$  sets the scale and is given by[?]

$$L^4 = 4\pi g_s N \alpha'^2. \quad (2.5)$$

---

<sup>1</sup> N.B.  $N$  refers to the number of D3 branes and the number of colors in the gauge theory, while  $\mathcal{N}$  indicates the number of supercharges in the SYM theory.

The vector  $\vec{x}$  runs over three spatial dimensions, and  $d\Omega_5$  is the five-dimensional angular element.

The stack of  $N$  D3 branes is located at  $r = 0$ , and to an observer located at  $r = \infty$ , any excitations near the branes appears to be low energy due to gravitational red-shifting. As a consequence, taking the low-energy limit is equivalent to taking the limit  $r \ll L$ . The metric becomes

$$ds^2 = \frac{r^2}{L^2}(-dt^2 + d\vec{x}^2) + \frac{L^2}{r^2}dr^2 + L^2 d\Omega_5, \quad (2.6)$$

where  $r$  is strictly positive. This metric is a product of five dimensional anti-de Sitter space with a five-dimensional sphere ( $AdS_5 \times S_5$ ). We can re-write this metric using the coordinate transformation

$$z = \frac{L^2}{r}. \quad (2.7)$$

This transformation gives a metric with a single warp factor,

$$ds^2 = \frac{L^2}{z^2}(-dt^2 + d\vec{x}^2 + dz^2), \quad (2.8)$$

where  $z > 0$ , with a UV cut-off at an infinitesimal  $z$ -value,  $z = z_0$ . The closed string excitations in the bulk are decoupled from the excitations near the branes.

In each of these descriptions, we began with a stack of  $N$  D3 branes and ended with a low-energy, strongly-coupled system. Taking the low-energy limit first results in an  $\mathcal{N} = 4$  SYM gauge theory, while taking the strong coupling limit first yields a system of excitations on an  $AdS_5 \times S_5$  metric. Maldacena's conjecture is that these two systems should describe the same physics. Each system has closed strings in the bulk that are decoupled from the rest of the system, so the conjecture results in a proposed duality between the  $\mathcal{N} = 4$  SYM theory and type IIB string theory in anti-de Sitter space.

This duality is particularly useful if we can ignore all stringy effects and treat the type IIB string theory as a classical supergravity theory on the metric (2.6). This approximation is valid if the string length is much less than the radius of curvature,

$$l_s \ll L. \quad (2.9)$$

We can relate these quantities to the Yang-Mills coupling using (2.2), (2.3), and (2.5), yielding

$$\frac{L^4}{l_s^4} = g_{YM}^2 N. \quad (2.10)$$

The quantity  $g_{YM}N$  is known as the 't Hooft coupling,  $\lambda$ . Thus, the requirement that the string length be small is equivalent to requiring the 't Hooft coupling to be large,  $\lambda \gg 1$ . In other words, the classical approximation is valid when the dual gauge theory is strongly coupled. The 't Hooft coupling acts as the effective gauge coupling, which is often expressed as a scalar field called the dilaton,  $\Phi$ , where

$$\Phi = \log \lambda. \quad (2.11)$$

The behavior of the dilaton field will take a central role in this thesis.

In summary, Maldacena conjectured that a ten-dimensional string theory is dual to  $\mathcal{N} = 4$  super Yang-Mills theory. We showed that the regime in which stringy effects can be neglected is equivalent to the strong-coupling limit of the gauge theory. The usefulness of this duality is evident, because the strong-coupling regime of a gauge theory is the limit in which it is difficult to perform calculations. Conveniently, this regime is dual to the string theory regime where calculations are easy because they can be done classically.

For our purposes, we reduce the 10-dimensional metric (2.6) to a five-dimensional space by ignoring the  $S_5$  manifold and keeping only the  $AdS_5$  metric (2.8). It turns out that discarding  $S_5$  removes the supersymmetry from the SYM theory, resulting in a non-supersymmetric conformal field theory. In the end, we have a duality between a classical gravity theory in five dimensional anti-de Sitter space and a four-dimensional conformal field theory, our AdS/CFT correspondence. It is this relationship between a field theory in four dimensions and a gravitational theory in five dimensions that lends these models the evocative name *holography*.

## 2.2 Applying AdS/CFT to Quantum Chromodynamics

Phenomenologists would like to apply gauge/theory dualities to physical theories like quantum chromodynamics, rather than supersymmetric Yang-Mills theories that do not directly relate to real-world phenomena. These theories are closely-related enough to encourage a variety of attempts to bridge this gap. To do so, we must examine the differences between SYM theories and QCD.

We have shown how to remove the supersymmetry aspect of the gauge theory in the correspondence in the preceding section, which is necessary because QCD is not



supersymmetric. However, we are still left with a conformal field theory, with no intrinsic energy scale. QCD has an energy scale called  $\Lambda_{QCD}$ , which is related to the phenomenon of confinement. Confinement is caused by the running of the QCD coupling with respect to energy scale. Conformal field theories have coupling constants which do not run, and therefore are not confining. Finally, the Yang-Mills theory has  $N \gg 1$ , while  $N_c$ , the number of colors in QCD is 3. Despite these differences, there are a variety of approaches to adapt the AdS/CFT correspondence to apply to QCD.

## 2.3 Top-Down Approach

One approach, known as the top-down approach, consists of modifying the string theory in some way in an attempt to produce a gauge theory that more accurately models QCD. Models that break conformal symmetry can study confinement [?]. The pure gauge theory with  $N$  D3 branes consists of only gluons, leading to the insertion of D7 “flavor branes” in the bulk to include quarks in the theory. [?] Strings that begin and end on the D7 branes represent quark/anti-quark pairs. These strings lack color indices and are therefore color singlets because they do not begin and end on the D3 branes. The excitations of this system corresponds to hadronic states. A prominent example that adds flavor and chiral symmetry breaking to the confining models is the Sakai-Sugimoto model [?].

## 2.4 Bottom-Up Approach

Another approach, known as bottom-up or AdS/QCD models, begin from a phenomenological viewpoint, modifying the existing AdS background to capture some essential features of QCD. Because of this motivation, it is not known whether such models could eventually be derived from string theory. However, it is useful to investigate AdS/QCD as an effective phenomenological model, as well as a means to gain insights that may help in constructing an eventual string theory dual.

The modifications to the AdS background are constructions imposed by hand to introduce features such as confinement and chiral dynamics. Fields that are dual to the operators of the field theory are introduced by hand into the bulk of the AdS

background. A Lagrangian is constructed from these fields, and the Euler-Lagrange equations that result are the equations of motion of the gauge field excitations. The normalizable solutions to these equations of motion are the hadronic states of the theory. The masses of the excited states of the mesons are calculated immediately from these eigenvalue problems, and other factors including decay constants and form factors can be calculated as well.

The backgrounds of these bottom-up models are imposed by hand, and not dynamically generated from any equations of motion. They are not derived from string theory, nor is it likely that they could be somehow embedded within any such theory. Despite the *ad hoc* nature of these models, the phenomenological results are often accurate to 10% or better. Importantly, they also give insight into how to capture the major features of QCD, such as confinement and chiral symmetry breaking. There are two major approaches to AdS/QCD, distinguished by the means used to break the conformal symmetry.

### The Hard Wall

The simplest way to break the conformal symmetry of the gauge theory is to impose a hard cutoff in the conformal ( $z$ ) dimension. The model was proposed by [?] and further explored by [?]. These so-called hard-wall AdS/QCD models insert both a UV brane located at  $z \rightarrow 0$  and an IR brane at

$$z_1 = \frac{1}{\Lambda_{QCD}}. \quad (2.12)$$

The fields can propagate only between these two branes. The confinement scale is introduced by the IR brane.

These hard-wall models capture a variety of features of QCD, including correlation functions and form factors. The major failing of the hard-wall model is in describing the spectrum of radially excited mesons. It is well established experimentally that these excited states scale as  $m_n^2 \sim n$  as  $n$  becomes sufficiently large, a phenomenon known as linear Regge trajectories or linear confinement. Hard-wall models produce a spectrum with a  $m_n \sim n$  scaling, in conflict with experiment.

## The Soft Wall

Instead of cutting off the metric, soft-wall models insert a  $z$ -dependent scalar dilaton field that breaks the conformal symmetry by acting as an effective cutoff. The dilaton multiplies the Lagrangian

$$\mathcal{S} = \int d^5x e^{-2\Phi} \sqrt{-g} \mathcal{L}, \quad (2.13)$$

modifying the equations of motion of the fields contained in  $\mathcal{L}$ . A common and simple choice for the dilaton's behavior is a power-law,

$$\Phi = (\mu z)^\nu, \quad (2.14)$$

where  $\mu$  introduces an energy scale into the model on the order of  $\Lambda_{QCD}$ . It was shown by Karch, *et al* [?] that a quadratic dilaton field

$$\Phi \sim z^2 \quad (2.15)$$

yields meson spectra with the desired  $m_n^2 \sim n$  behavior. Because these linear Regge trajectories apply only when  $n$  becomes large, the power-law behavior is necessary strictly when  $z$  becomes large. Improved predictions for the lower meson states can be obtained by modifying the UV behavior of the dilaton field. Such modifications to this model were explored in [?]. The details of past work on soft-wall models will be explored further in Chapter 3.

## 2.5 Dynamical AdS/QCD

The early AdS/QCD models rely on parameterizations for the background fields such as the dilaton. A power law is the simplest choice for the dilaton's behavior, and was suggested in [?]. Later models including [?] improved phenomenology by modifying the dilaton in the UV limit. However, the models rely on *ad hoc* choices for the parameterization of the background fields and do not examine the dynamics of these fields.

So-called dynamical AdS/QCD models attempt a more rigorous examination of the vacuum dynamics of the dual model, as an attempt to remedy some of the shortcomings of the bottom-up approach discussed above. These models examine terms in the Lagrangian involving the dilaton field as well as a tachyonic field that may be related to

the chiral symmetry breaking of the model. Previous work examines the construction of the potential terms in the Lagrangian that will give rise to the desired behavior of the background fields. There has also been some work on examining meson spectra in dynamical AdS/QCD models, with some success.

## 2.6 Summary

In this chapter, we have presented the basics of AdS/CFT, including a conceptual motivation for the correspondence. We have also discussed the application of the correspondence to QCD, including the main research programs attempting to make this connection: the top-down and bottom-up approaches. Finally, we have introduced two areas of research in bottom-up AdS/QCD: the soft-wall model and dynamical AdS/QCD. These topics will be elaborated upon in the following chapters.

## Chapter 3

# Soft-Wall Model of AdS/QCD

In this chapter, we will provide detail on previous work done on the soft-wall AdS/QCD model. We will present the set-up of the background and the fields included in the model. The method for modeling the chiral symmetry breaking of QCD are discussed. The Euler-Lagrange equations for the gauge fields are derived, and the eigenvalue equations for the masses of the meson excited states are presented.

### 3.1 Metric and Field Content

The four-dimensional fields of QCD live in the five-dimensional anti-de Sitter space, with a metric given by

$$ds^2 g_{MN} dx^M dx^N = a^2(z)(\eta_{\mu\nu} dx^\mu dx^\nu + dz^2), \quad (3.1)$$

where  $a(z) = L/z$  is the warp factor and  $L$  is the curvature radius of the anti-de Sitter space. We will often work in units where the curvature radius is unity. The Minkowski metric is given by

$$\eta_{\mu\nu} = \begin{pmatrix} -1 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \\ 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix}. \quad (3.2)$$

The coordinate  $z$  has a range  $0 \leq z < \infty$ .

The bulk coordinate  $z$  is associated with inverse energy scales, with the ultraviolet limit of QCD represented by fields at  $z \rightarrow 0$  [?]. The AdS/CFT dictionary [?, ?] states that each operator  $\mathcal{O}(x)$  in the 4D conformal field theory is associated with a bulk field  $\psi(x, z)$ . The values of the bulk fields at the UV boundary act as sources for the corresponding 4D currents. Global symmetries of the 4D field theory become gauged symmetries for the bulk fields.

The generating functional of gauge-invariant operators of the gauge theory is dual to the minimum of the supergravity action. On the boundary, the supergravity fields must coincide with the sources of the gauge theory [?, ?]. That is,

$$\langle e^{\int_{\partial AdS} d^4x \phi^0(\vec{x}) \mathcal{O}(\vec{x})} \rangle_{\text{CFT}} = e^{iS_{\text{SUGRA}}(\phi)}|_{\phi=\phi_0}, \quad (3.3)$$

where  $\phi^0$  is a source term on the boundary of the five-dimensional supergravity (SUGRA) model. The gauge theory exists on the boundary,  $\partial AdS$ , of the dual gravitational theory. For each operator of conformal dimension  $\Delta$  in the gauge theory, there is a coupling,  $\phi^0 \mathcal{O}$ . For each of these terms, a scalar field with mass  $m$  is inserted into the  $AdS_5$  gravity dual.

We can write down the Lagrangian of such a free massive scalar field,

$$\mathcal{L} = g^{MN} \partial_M \phi \partial_N \phi - m^2 \phi^2. \quad (3.4)$$

The Euler-Lagrange equation is found by varying the action with respect to  $\phi$ ,

$$\begin{aligned} \delta \mathcal{S} &= \delta (\sqrt{-g} g^{MN} \partial_M \phi \partial_N \phi - \sqrt{-g} m^2 \phi^2) \\ &= \sqrt{-g} g^{\mu\nu} \partial_\mu \delta \phi \partial_\nu \phi + \sqrt{-g} g^{zz} \partial_z \phi \delta \partial_z \phi - \sqrt{-g} m^2 \phi \delta \phi \\ &= -a(z)^3 \eta^{\mu\nu} \partial_\mu \partial_\nu \phi \delta \phi - \eta^{zz} \partial_z (a(z)^3 \partial_z \phi \delta \phi) - a(z)^5 m^2 \phi \delta \phi \\ &= (z^5 \partial_z (z^{-3} \partial_z \phi) + z^2 \partial_z^2 \phi - m^2 \phi) \delta \phi, \end{aligned} \quad (3.5)$$

resulting in the equation of motion

$$z^5 \partial_z (z^{-3} \partial_z \phi) + (m^2 + z^2 \partial_z^2) \phi = 0. \quad (3.6)$$

Using the substitution  $\partial_z^2 \phi = -k^2 \phi$ , we can write this equation in momentum space

$$z^5 \partial_z (z^{-3} \partial_z \phi) + (m^2 - k^2 z^2) \phi = 0. \quad (3.7)$$

Near the  $z = 0$  boundary, the  $k^2$  term is neglected, and the asymptotic field behavior becomes [?, ?]

$$\phi \approx \phi^0 z^{\Delta_-} + \langle \mathcal{O} \rangle z^{\Delta_+}, \quad (3.8)$$

where the two solutions are

$$\Delta_{\pm} = \frac{d}{2} \pm \sqrt{\frac{d^2}{4} + m^2 R^2}. \quad (3.9)$$

The mass is determined by the spatial dimension  $d$  and the operator dimension  $\Delta$

$$m^2 = \Delta(\Delta - d), \quad (3.10)$$

giving the two possible solutions of (3.9),

$$\Delta_+ = d, \quad (3.11)$$

$$\Delta_- = d - \Delta. \quad (3.12)$$

An operator  $\mathcal{O}$  on the boundary may also be coupled to a  $p$ -form  $\mathcal{C}$  in AdS space, via the coupling

$$\int_{M_d} \mathcal{C} \wedge \mathcal{O}. \quad (3.13)$$

In this case, the mass of the scalar field in the gravity dual is [?]

$$m^2 = (\Delta - p)(\Delta - p - d). \quad (3.14)$$

It seems that we should require  $\Delta > p + d$ , to avoid instabilities from tachyons with  $m^2 < 0$ . However, it has been shown [?] that tachyons with mass above the bound

$$m^2 > -\frac{d^2}{4} \quad (3.15)$$

are allowed.

The symmetries of the gauge field are also included in the gravitational dual through a prescription provided by the correspondence dictionary. A global symmetry of the gauge theory is represented by invariance under a transformation

$$U = e^{i\eta}, \quad (3.16)$$

with  $\eta$  constant. This means that  $\eta$  should remain constant on the boundary of the gravity dual, but in the bulk, nothing prevents  $\eta$  from becoming a function of spacetime

coordinates,  $x$  and  $z$ . The symmetry in the gravity dual is then represented by the transformation

$$U = e^{i\eta(x,z)}. \quad (3.17)$$

In this way, the global symmetry of the gauge theory is represented by a local symmetry of the gravitational theory. A gauge field is inserted into the gravity dual for each relevant global symmetry.

For our AdS/QCD model, the field content of the five-dimensional gravitational dual theory is dictated by the operators relevant to the chiral dynamics of QCD. The gauge fields  $L_\mu$ ,  $R_\mu$  correspond to the left- and right-handed currents of the  $SU(N_f)_L \times SU(N_f)_R$  chiral symmetry, where  $N_f$  is the number of quark flavors in the model. The scalar field  $X$  is associated with the chiral operator  $\bar{q}q$  [?]. The masses of the bulk fields are set by the AdS/CFT relation [?]

$$m_5^2 L^2 = (\Delta - p)(\Delta + p - 4), \quad (3.18)$$

where  $\Delta$  is the dimension of the  $p$ -form QCD operator. Table 3.1 illustrates the fields and operators of our model, showing that the scalar field is the only field in this model that is not massless.

4D Operator	5D Field	$p$	$\Delta$	$m_5^2 L^2$
$\bar{q}_L \gamma^\mu t^a q_L$	$L_\mu^a$	1	3	0
$\bar{q}_R \gamma^\mu t^a q_R$	$R_\mu^a$	1	3	0
$\bar{q}_R^a q_L^b$	$\frac{2}{z} X^{ab}$	0	3	-3

Table 3.1: Operators and fields of the model. The matrices  $t^a$  are the generators of the  $SU(N_f)$  symmetry.

The simplest soft-wall action involving the fields from Table 3.1 is given in [?] as

$$S_5 = \int d^5x \sqrt{-g} e^{-\Phi(z)} \text{Tr} \left[ |DX|^2 + m_X^2 |X|^2 + \frac{1}{4g_5^2} (F_L^2 + F_R^2) \right]. \quad (3.19)$$

The 5D gauge coupling constant  $g_5$  is fixed by calculating the vector current two-point function using this model and then comparing this to the leading order result from QCD, leading to the identification  $g_5^2 = 12\pi^2/N_c$  [?].



The field  $X$  includes both the scalar and pseudoscalar fields, as well as a non-trivial vacuum expectation value (VEV)

$$X_e^{ab} = \left( \frac{\chi(z)}{2} + S^a(x, z)t^b \right) I e^{2i\pi_e(x, z)^a t^b} \quad (3.20)$$

$$X_l^{ab} = \left( \frac{\chi(z)}{2} + S^a(x, z)t^b \right) I + i\pi_l(x, z)^a t^b, \quad (3.21)$$

with  $I$  the  $N_f \times N_f$  identity matrix and  $t^a$  the  $SU(N_f)$  generators, which are normalized as

$$\text{Tr}[t^a t^b] = \delta^{ab}/2. \quad (3.22)$$

The indices  $a, b$  of the field  $X$  will be suppressed unless needed. The details of the representation choice will be discussed in Section ???. The field strength tensors are defined as

$$F_L^{MN} = \partial^M L^N - \partial^N L^M - i[L^M, L^N] \quad (3.23)$$

$$F_R^{MN} = \partial^M R^N - \partial^N R^M - i[R^M, R^N], \quad (3.24)$$

where we have used the shorthand notation  $L^M = L^M t^a$ . The covariant derivative becomes

$$D^M X = \partial^M X - iL^M X + iXR^M. \quad (3.25)$$

The physical vector ( $V$ ) and axial-vector ( $A$ ) fields are defined in terms of the  $L$  and  $R$  gauge fields

$$L^M = V^M + A^M, \quad (3.26)$$

$$R^M = V^M - A^M. \quad (3.27)$$

Substituting equations (3.26) and (3.27) into the last field strength tensors,

$$\begin{aligned} F_L^2 + F_R^2 &= 2 \left( \partial^M L^N \partial_M L_N - \partial^M L^N \partial_N L_M + \frac{1}{2} [L^M, L^N] [L_M, L_N] \right) \\ &+ 2 \left( \partial^M R^N \partial_M R_N - \partial^M R^N \partial_N R_M + \frac{1}{2} [R^M, R^N] [R_M, R_N] \right), \\ &= 4 \left( \partial^M V^N \partial_M V_N - \partial^M V^N \partial_N V_M + \frac{1}{2} [V^M, V^N] [V_M, V_N] \right) \\ &+ 4 \left( \partial^M A^N \partial_M V_N - \partial^M A^N \partial_N A_M + \frac{1}{2} [A^M, A^N] [A_M, A_N] \right), \\ &= 2 (F_V^2 + F_A^2), \end{aligned} \quad (3.28)$$

where the vector and axial field-strength tensors have the form

$$F_V^{MN} = \partial^M V^N - \partial^N V^M - \frac{i}{\sqrt{2}}[V^M, V^N], \quad (3.29)$$

$$F_A^{MN} = \partial^M A^N - \partial^N A^M - \frac{i}{\sqrt{2}}[A^M, A^N]. \quad (3.30)$$

The covariant derivative can also be written in terms of the vector and axial fields

$$D_M X = \partial_M X - i\{A_M^a, X\} + i[V_M^a, X]. \quad (3.31)$$

This relation allows us to re-write the action (3.19),

$$S_5 = \int d^5x \sqrt{-g} e^{-\Phi(z)} \text{Tr} \left[ |DX|^2 + m_X^2 |X|^2 + \frac{1}{2g_5^2} (F_A^2 + F_V^2) \right]. \quad (3.32)$$

The scalar field  $X$  takes on a  $z$ -dependent vacuum expectation value (VEV), breaking the chiral symmetry. In a flavor-symmetric model, the VEV has the form

$$\langle X \rangle = \frac{\chi(z)}{2} I, \quad (3.33)$$

where  $I$  is the  $N_f \times N_f$  identity matrix.

### 3.1.1 Equations of Motion

By varying the action (3.19), one obtains the equations of motion for the scalar, pseudoscalar, vector, and axial-vector mesons, as well as for the scalar vacuum expectation value.

Substituting the chiral field into the Lagrangian using (3.20), we vary the Lagrangian with respect to  $\chi(z)$ ,

$$\begin{aligned} \delta \mathcal{L}_{VEV} &= -\delta \left( e^{-\Phi} \sqrt{-g} \text{Tr} \left[ \frac{g^{zz}}{4} \partial_z \chi \partial_z \chi + m_X^2 \frac{\chi^2}{4} \right] \right) \\ &= -\delta \left( \frac{1}{2} e^{-\Phi} a(z)^5 (a(z)^{-2} \partial_z \chi \partial_z \chi + m_X^2 \chi^2) \right) \\ &= -e^{-\Phi} a(z)^3 \partial_z \chi \delta \partial_z \chi - e^{-\Phi} a(z)^5 m_X^2 \chi \delta \chi \\ &= (\partial_z (e^{-\Phi} a^3 \partial_z \chi) - e^{-\Phi} a^5 m_X^2 \chi) \delta \chi, \end{aligned} \quad (3.34)$$

The equation of motion becomes

$$\partial_z^2 \chi - \partial_z \Phi \partial_z \chi + \frac{\partial_z a(z)}{a(z)} \partial_z \chi - a(z)^2 m_X^2 \chi = 0. \quad (3.35)$$

In this model, the warp factor  $a(z) = 1/z$  and the mass of the scalar field  $m_X^2 L^2 = -3$ . With these substitutions, the equation of motion becomes

$$\chi'' - \Phi' \chi' - \frac{3}{z} \chi' + \frac{3}{z^2} \chi = 0. \quad (3.36)$$

In the hard wall model, there is no dilaton, so we can set  $\Phi' = 0$ , finding the exact solution to (3.36) to be

$$\chi_{\text{hw}}(z) = c_1 z + c_2 z^3, \quad (3.37)$$

where  $c_1$  and  $c_2$  are integration constants. Comparing this solution to the UV behavior of the bulk fields in the AdS/CFT dictionary (3.8),  $c_1$  and  $c_2$  correspond to the source term and the operator expectation value, respectively,

$$c_1 \sim m_q, \quad (3.38)$$

$$c_2 \sim \langle q\bar{q} \rangle \equiv \sigma. \quad (3.39)$$

Here  $m_q$  is the quark mass and  $\sigma = \langle \bar{q}q \rangle$  is the chiral condensate, the variation of the vacuum energy with respect to  $m_q$ .

In the simplest soft-wall model, it has been shown [?] that the solution to (3.36) is given by

$$\chi_{sw}(z) = \frac{m_q}{L} z \Gamma\left(\frac{3}{2}\right) U\left(\frac{1}{2}, 0, \lambda z^2\right), \quad (3.40)$$

where  $U(a, b, y)$  is the Tricomi confluent hypergeometric function. (There is also another solution that is disregarded because it leads to an action that is not finite in the IR.) In the small- $z$  limit, (3.40) expands to [?]

$$\chi_{sw}(z) \rightarrow \frac{m_q}{L} z - \frac{\lambda m_q}{2L} z^3 \left(1 - 2\gamma_E - 2\log(\sqrt{\lambda}z) - \psi\left(\frac{3}{2}\right)\right), \quad (3.41)$$

where  $\psi$  is the Euler function. However, comparing this solution to the UV behavior of the chiral condensate in the AdS/CFT dictionary (3.8), we see that  $\sigma \sim m_q$ . This is undesirable in a theory with spontaneous chiral symmetry breaking. In addition, when  $z$  becomes large, we see that the chiral condensate becomes a constant,  $\chi \sim m_q$ , linking the IR and UV behavior. As we will see, the IR behavior of the chiral field governs the behavior of the axial meson spectrum, another limitation of the solution (3.40). More generally, we can see that this issue is caused because the equation (3.36) is linear in  $\chi$ . With only one normalizable solution, only one constant survives the application of

boundary conditions. As a result,  $m_q$  and  $\sigma$  cannot be independent. To remedy this, it was suggested [?] to examine higher-order terms in the scalar potential. A particular example of this will be examined in Section 3.2, and a more general approach will be discussed in Chapter ??.

### Scalar Sector

We now examine the fluctuations of the scalar field  $X$ , rather than its VEV, to determine the spectrum of the  $f_0$  scalar mesons. We use the expression in (3.20), though it does not matter which representation for the pseudoscalar component of the field is used,

$$X(x, z) = \left( S(x, z) + \frac{\chi(z)}{2} \right) e^{2i\pi(x, z)}. \quad (3.42)$$

To obtain the equations of motion, we vary the action with respect to the scalar field  $S(x, z)$ ,

$$\begin{aligned} \delta \mathcal{L}_S &= \delta \left( e^{-\Phi} \sqrt{-g} \text{Tr} [g^{MN} \partial_M S(x, z) \partial_N S(x, z) + m_X^2 S(x, z)^2] \right) \\ &= \delta \left( e^{-\Phi} \sqrt{-g} \text{Tr} (t^a t^b) (g^{\mu\nu} \partial_\mu S \partial_\nu S + g^{zz} \partial_z S \partial_z S + m_X^2 S^2) \right) \\ &= e^{-\Phi} a(z)^3 \eta^{\mu\nu} \partial_\mu S \delta \partial_\nu S + e^{-\Phi} a(z)^3 \partial_z S \delta \partial_z S + e^{-\Phi} a(z)^5 m_X^2 S \delta S \\ &= (-e^{-\Phi} a(z)^3 \partial_\mu \partial^\mu S - \partial_z (e^{-\Phi} a(z)^3 \partial_z S) + e^{-\Phi} a(z)^5 m_X^2 S) \delta S, \end{aligned} \quad (3.43)$$

leaving the equation of motion,

$$e^\Phi a^{-3} \partial_z (e^{-\Phi} a^3 \partial_z S) + \partial_\mu \partial^\mu S - a^2 m_X^2 S = 0. \quad (3.44)$$

We separate the  $z$ -dependent part of the field by using Kaluza-Klein decomposition,

$$S(x, z) = \sum_{n=0}^{\infty} S_n(z) \mathcal{S}_n(x). \quad (3.45)$$

We use the AdS warp factor  $a(z) = 1/z$  and use Proca's equation

$$\partial_i \partial^i S_n = m_n^2 S_n, \quad (3.46)$$

to obtain the equation of motion

$$-\partial_z^2 S_n + \omega'_s \partial_z S_n = m_{S_n}^2 S_n, \quad (3.47)$$

where we have defined

$$\omega_s \equiv \Phi(z) + 3 \log z, \quad (3.48)$$

and  $(')$  represents differentiation with respect to  $z$ . We can eliminate the first derivative of the field  $S_n$  and put the equation of motion in Schrödinger-like form with the substitution

$$S_n = e^{\omega_s/2} s_n. \quad (3.49)$$

The final equation of motion is now

$$-s_n'' + \left( \frac{1}{4} \omega_s'^2 - \frac{1}{2} \omega_s'' + \frac{m_X^2}{z^2} \right) s_n = m_{S_n}^2 s_n. \quad (3.50)$$

### Pseudoscalar Sector

The pseudoscalar eigenstates correspond to the pions, the pseudo-Goldstone bosons of chiral symmetry. This sector is the most difficult to analyze because the pseudoscalar field couples to the longitudinal component of the axial-vector field,

$$A_\mu = A_{\mu\perp} + \partial_\mu \varphi. \quad (3.51)$$

This results in two coupled differential equations, one that comes from varying the action with respect to the pseudoscalar field  $\pi$ , and one that comes from varying with respect to  $\varphi$ . There are also subtleties in the choice of representation, as illustrated above in (3.20-3.21). The various issues involved in analyzing the pseudoscalar sector are discussed in Section 3.2.2.

### Vector Sector

We can derive the mass spectrum of the vector  $\rho$  mesons by varying the vector field and using the axial gauge condition  $V_z = 0$ . Varying the action, we find that

$$\begin{aligned} \delta \mathcal{S}_V &= -\delta \left( e^{-\Phi} \sqrt{-g} g^{\mu\rho} g^{\nu\sigma} (\partial_\mu V_\nu \partial_\rho V_\sigma - \partial_\mu V_\nu \partial_\sigma V_\rho) + e^{-\Phi} \sqrt{-g} g^{zz} g^{\mu\nu} \partial_z V_\mu \partial_z V_\nu \right) \\ &= -e^{-\Phi} \sqrt{-g} g^{\mu\rho} g^{\nu\sigma} (\partial_\mu V_\nu \delta \partial_\rho V_\sigma - \partial_\mu V_\nu \delta \partial_\sigma V_\rho) - e^{-\Phi} \sqrt{-g} g^{zz} g^{\mu\nu} \partial_z V_\mu \delta \partial_z V_\nu \\ &= e^{-\Phi} a(z) (\partial^2 V_\mu \delta V^\mu - \partial_\mu \partial^\nu V_\nu \delta V^\mu) + \partial_z (e^{-\Phi} a(z) \partial_z V_\mu) \delta V^\mu \\ &= (e^{-\Phi} a(z) \partial_\nu \partial^\nu V_\mu + \partial_z (e^{-\Phi} a(z) \partial_z V_\mu)) \delta V^\mu, \end{aligned} \quad (3.52)$$

We can separate out the  $z$ -dependence of the field using Kaluza-Klein (KK) decomposition

$$V_\mu^n = \sum_{n=0}^{\infty} \mathcal{V}_\mu^n(x) V_n(z), \quad (3.53)$$

where  $V_n(z)$  are the Kaluza-Klein modes. The equation of motion is now one dimensional

$$-\partial_z^2 V_n + \omega' \partial_z V_n = m_{V_n}^2 V_n, \quad (3.54)$$

where we have defined

$$\omega \equiv \Phi(z) + \log z. \quad (3.55)$$

We can eliminate the first derivative of the field, bringing the equation of motion into Schrödinger-like form, using the substitution

$$V_n(z) = e^{\omega/2} v_n(z). \quad (3.56)$$

The equation of motion is now

$$-v_n'' + \left( \frac{1}{4} \omega'^2 - \frac{1}{2} \omega'' \right) v_n = m_{V_n}^2 v_n. \quad (3.57)$$

This is the general method for finding the equations of motion for the various fields. Using the asymptotic form of the dilaton  $\Phi = \lambda z^2$ , the mass eigenvalues can be found exactly for this model. The equation of motion is now

$$-v_n'' + \left( \lambda z^2 + \frac{3}{4z^2} \right) v_n = m_n^2 v_n. \quad (3.58)$$

The eigenvalues for this equation are  $m_n^2 = \lambda(4n+4)$  for  $n = 0, 1, 2, 3, \dots$ . The parameter  $\lambda$  is set by matching this trajectory to experimental data.

We can also determine the value of the 5D coupling constant  $g_5$  used in the action (3.19) by analyzing the vector sector. This is done by matching the vector two-point function  $\Pi_V(q^2)$  calculated from this model to the calculation from the operator product expansion of QCD [?], As shown in [?, ?], we calculate the two-point function near the UV boundary. We begin by re-writing (3.52) in the hard-wall case  $\Phi' = 0$ ,

$$\partial_z \left( \frac{1}{z} \partial_z V_\mu(q, z) \right) + \frac{q^2}{z} V_\mu(q, z) = 0, \quad (3.59)$$

where  $V_\mu(q, z)$  is the 4D Fourier transform of  $V_\mu(x, z)$ , and we have used the Fourier-transformed version of the Proca equation,

$$\partial_i \partial^i V_\mu(q, z) = -q^2 V_\mu(q, z) \quad (3.60)$$

Evaluating the vector part of the action (3.19) leaves the boundary action

$$\mathcal{S}_b = -\frac{1}{2g_5^2} \int d^4x \left( \frac{1}{z} V_\mu \partial_z V^\mu \right)_{z=\epsilon}, \quad (3.61)$$

where  $\epsilon$  is a UV boundary close to  $z = 0$ . We define  $V_0^\mu$  to be the Fourier-transformed source of the vector current  $J^\mu = \bar{q} \gamma_\mu t^a q$  at the UV boundary,

$$V_0^\mu \equiv \int d^4x e^{iqx} J^\mu. \quad (3.62)$$

Re-writing the vector field in a separable form

$$V^\mu(q, z) = V(q, z) V_0^\mu(q), \quad (3.63)$$

we see that we should choose the UV boundary condition  $V^\mu(q, \epsilon) = 1$ . Using this separable form, (3.61) becomes

$$\mathcal{S}_b = -\frac{1}{2g_5^2} \int d^4x V_{0\mu} \left( \frac{\partial_z V(q, z)}{z} \right)_{z=\epsilon} V_0^\mu. \quad (3.64)$$

This expression shows why  $V(q, z)$  is often known as the *bulk-to-boundary* propagator. Twice differentiating the boundary action with respect to  $V_0$ , we obtain the vector current two-point function,

$$\int d^4x e^{iqx} \langle J_\mu^a(x) J_\nu^b(0) \rangle = \delta^{ab} (q_\mu q_\nu - q^2 g_{\mu\nu}) \Pi_V(q^2), \quad (3.65)$$

where

$$\Pi_V(Q^2) = -\frac{1}{g_5^2 Q^2} \left. \frac{\partial_z V(q, z)}{z} \right|_{z=\epsilon}, \quad (3.66)$$

where  $Q^2 = -q^2$ . In the limit of large  $Q^2$ , we can expand  $V(q, z)$  near the UV boundary,

$$V(Q, z) = 1 + \frac{Q^2 z^2}{4} \log Q^2 z^2 + \dots, \quad (3.67)$$

and to first order, the correlation function becomes

$$\Pi_V(Q^2) = -\frac{1}{2g_5^2} \log Q^2. \quad (3.68)$$

Matching to the large- $N_c$  QCD perturbative result,

$$\Pi_V(k^2) = -\frac{N_c}{24\pi^2} \log k^2, \quad (3.69)$$

where  $N_c$  is the number of colors, we find that

$$g_5^2 = \frac{12\pi^2}{N_c}. \quad (3.70)$$

In this work, we take  $N_c = 3$ , and assume that this value is large enough for the large  $N_c$  results to hold.

### Axial-Vector Sector

The equation of motion for the axial sector is derived using the same method. Varying the action with respect to  $A_\mu$  and using the axial gauge  $A_z = 0$ , we get an equation of similar form to (3.52), but with an additional chiral symmetry-breaking term  $\chi^2 A_\mu$ ,

$$\delta\mathcal{S}_A = (e^{-\Phi} a(z) \partial^2 A_\mu + \partial_z (e^{-\Phi} a(z) \partial_z A_\mu) + e^{-\Phi} g_5^2 a(z)^3 v^2 A_\mu) \delta A^\mu. \quad (3.71)$$

After KK decomposition,

$$A_\mu^n = \sum_{n=0}^{\infty} \mathcal{A}_\mu^n(x) A_n(z), \quad (3.72)$$

the equation of motion becomes

$$-\partial_z^2 A_n + \omega' \partial_z A_n + \frac{g_5^2 \chi(z)^2}{z^2} A_n = m_{A_n}^2 A_n, \quad (3.73)$$

with  $\omega$  defined as in (3.55). To put the equation of motion in Schrödinger form, we make the substitution

$$A_n = e^{\omega/2} a_n, \quad (3.74)$$

yielding

$$-a_n'' + \left( \frac{1}{4} \omega'^2 - \frac{1}{2} \omega'' + g_5^2 \frac{L^2}{z^2} \chi^2(z) \right) a_n = m_{V_n}^2 a_n. \quad (3.75)$$

The only difference from (3.57) is the presence of the  $z$ -dependent mass term involving the chiral condensate. The IR asymptotic behavior of the chiral condensate  $\chi(z)$  controls the splitting between the vector and axial-vector mesons at large values of  $n$ . Experimentally, we see a constant splitting between the squared masses of the  $a_1$  and  $\rho$



mesons, which suggests that chiral symmetry is not restored for the higher excitations. In this model, we can see that the difference between the equations of motion (3.57) and (3.75) is given by

$$\Delta m^2 \equiv (m_{A_n}^2 - m_{V_n}^2)_{n \rightarrow \infty} = g_5^2 \frac{L^2 \chi^2}{z^2} (z \rightarrow \infty). \quad (3.76)$$

To obtain a constant mass splitting, this quantity must become a constant as  $z$  approaches infinity. This tells us that  $\chi \sim z$  in the IR limit.

The significant drawbacks for this simple soft-wall model are the relatively poor modeling of the ground state and lower resonances and the lack of independent spontaneous and explicit chiral symmetry breaking terms.

## 3.2 Modified Soft-Wall Model

An improvement on the soft-wall model, suggested in [?], is adding higher-order terms of  $X$  to the scalar potential, separating the spontaneous and explicit chiral symmetry breaking. The model established in [?] adds a quartic scalar term to the action:

$$\mathcal{S} = \int d^5x \sqrt{-g} e^{-\Phi(z)} \text{Tr} \left[ |DX|^2 + m_X^2 |X|^2 - \kappa |X|^4 + \frac{1}{2g_5^2} (F_A^2 + F_V^2) \right], \quad (3.77)$$

where  $\kappa$  is a dimensionless parameter to be fit to the data. To obtain the required linear Regge trajectories for the meson spectra, the dilaton field must be quadratic in  $z$  in the IR region,

$$\Phi(z \rightarrow \infty) = \lambda z^2, \quad (3.78)$$

where  $\lambda$  sets an energy scale for the model that is related to the slope of the Regge trajectories.

The chiral symmetry breaking of the model is examined by deriving the equation of motion for the vacuum expectation value of the scalar field  $X$ . Varying (3.77) with respect to  $X$ , we find

$$\begin{aligned} \delta \mathcal{S}_X = & -2e^{-\Phi} \sqrt{-g} \text{Tr} \left( g^{MN} \partial_M X \partial_N \delta X + \{A, X\} \{A, \delta X\} \right. \\ & \left. + [V, X] [V, \delta X] + m_X^2 X \delta X - 2\kappa X^\dagger X |X| \delta X \right). \end{aligned} \quad (3.79)$$

Taking the trace, integrating by parts, and using  $|X| = \chi(z)/2 + S(x, z)$ , we find the variation of the action with respect to the chiral condensate,

$$\delta\mathcal{S}_\chi = \frac{1}{2} \left( \partial_z \left( e^{-\Phi} \sqrt{-g} g^{zz} \partial_z \chi \right) + e^{-\Phi} \sqrt{-g} m_X^2 \chi + e^{-\Phi} \sqrt{-g} \frac{\kappa}{2} \chi^3 \right) \delta\chi. \quad (3.80)$$

Using the AdS metric, we find that the chiral condensate  $\chi(z)$  now has a nonlinear equation of motion

$$\partial_z \left( a^3 e^{-\Phi} \partial_z \chi(z) \right) - a^5 e^{-\Phi} \left( m_X^2 \chi(z) - \frac{\kappa}{2} \chi^3(z) \right) = 0, \quad (3.81)$$

which simplifies to

$$\chi'' - \left( \Phi' + \frac{3}{z} \right) \chi' - m_X^2 \chi + \frac{\kappa}{2} \chi^3 = 0, \quad (3.82)$$

where  $(')$  denotes a derivative with respect to  $z$ .

The higher radially excited states of mesons have parallel Regge trajectories. However, the eigenstates of vector  $\rho$  mesons and the axial-vector  $a_1$  mesons do not become degenerate at large  $n$ , despite having the same spin. This indicates that chiral symmetry is not restored for the higher excitations. As noted in Section 3.1.1, the mass-splitting between the highly-excited states of the  $a_1$  and  $\rho$  mesons is governed by the IR behavior of the chiral condensate field  $\chi$ ,

$$\Delta m^2 \equiv (m_{A_n}^2 - m_{V_n}^2)_{n \rightarrow \infty} = g_5^2 \frac{L^2 \chi^2}{z^2} (z \rightarrow \infty). \quad (3.83)$$

Because their trajectories are parallel but not equal, we know that this  $\Delta m^2$  must be a constant. Examining the right-hand side of (3.83), we see that this requirement implies that  $\chi(z) \sim z$  for large  $z$ . To match the AdS/CFT dictionary discussed in Chapter 1, the chiral condensate must retain the same UV asymptotic form (3.37). A suitable parameterization that matches the expected UV and IR asymptotic behavior was found and justified in [?]

$$\chi(z) = \alpha z + \beta z \tanh(\gamma z^2), \quad (3.84)$$

with the parameters defined as follows

$$\alpha = \frac{\sqrt{3} m_q}{g_5 L}, \quad \beta = \sqrt{\frac{4\lambda}{\kappa L^2}} - \alpha, \quad \gamma = \frac{g_5 \sigma}{\sqrt{3} \beta}, \quad (3.85)$$

where  $m_q$  is the quark mass,  $\sigma$  is the chiral condensate,  $\lambda$  is set by the experimental Regge trajectories, and  $g_5$  is determined by (3.70) as derived Section 3.1.1. For convenience, we will work in units where the AdS curvature radius  $L$  is unity. In the small- $z$  limit, the chiral field becomes

$$\begin{aligned}\chi(z \rightarrow 0) &= \alpha z + \beta \gamma z^3 \\ &= \frac{\sqrt{3}}{g_5} m_q + \frac{g_5}{\sqrt{3}} \sigma,\end{aligned}\tag{3.86}$$

where  $\sqrt{3}/g_5$  is a normalization factor discussed in [?]. It is clear that, up to normalization, (3.86) matches the UV form required by the AdS/CFT dictionary. In the large- $z$  limit, the chiral field becomes

$$\begin{aligned}\chi(z \rightarrow \infty) &= (\alpha + \beta)z \\ &= \sqrt{\frac{4\lambda}{\kappa}} z.\end{aligned}\tag{3.87}$$

The chiral field is linear in the IR, as required, and the dimensionless parameter  $\kappa$  introduced in (3.77) becomes the parameter that controls the axial-vector mass splitting.

The quark mass and chiral condensate can each be taken to zero independently, and the non-restoration of chiral symmetry does not depend on either of these parameters. Thus, the spontaneous and chiral symmetry breaking parameters are independent, as desired. Using (3.84) in (3.81) we can solve for the derivative of the dilaton field

$$\Phi' = \frac{1}{a(z)^3 \chi'} \left( \partial_z (a(z)^3 \chi') - a^5 \left( m_X^2 \chi - \frac{\kappa}{2} \chi^3 \right) \right).\tag{3.88}$$

Using (??) in (3.88), the asymptotic behavior of the dilaton is found to be

$$\Phi(z \rightarrow 0) = \frac{\kappa}{4} \alpha^2 z^2 + \mathcal{O}(z^6),\tag{3.89}$$

$$\Phi(z \rightarrow \infty) = \frac{\kappa}{4} (\alpha + \beta)^2 z^2 = \lambda z^2,\tag{3.90}$$

where the boundary condition  $\Phi(0) = 0$  is chosen to ensure a pure AdS metric in the UV limit. We can see that the desired IR dilaton behavior is recovered by this parameterization.

### 3.2.1 Meson Spectra

Using these parameterizations for the dilaton and chiral condensate fields, we can now calculate the mass eigenvalues of the scalar, vector, and axial-vector mesons. Equations of motion are then derived using the method of Section 3. Due to the more complicated forms for  $\chi$  and  $\Phi$ , the eigenvalues are not analytically solvable, so a numerical shooting method is used to calculate the mass spectra for the scalar, vector, and axial sectors. For details on the shooting method, see Appendix A.

#### Scalar Mesons

The mass spectrum of the scalar  $f_0$  mesons is determined by deriving the equations of motion for the scalar field  $S(x, z)$  and calculating its eigenvalues. The procedure is similar to that delineated in Section 3.1.1, but the quartic term in the scalar potential in (3.77) causes the fluctuations of the scalar field to couple to its own vacuum expectation value. Using a Kaluza-Klein decomposition

$$S(x, z) = \sum_{n=0}^{\infty} \mathcal{S}_n(x) S_n(z), \quad (3.91)$$

we vary the action (3.77) with respect to  $S$ , yielding

$$\partial_z(a^3(z)e^{-\Phi}\partial_z S_n(z)) - a^5(z)e^{-\Phi}\left(m_X^2 - \frac{3}{2}\kappa\chi^2(z)\right)S_n(z) = -a^3(z)e^{-\Phi}m_{S_n}^2 S_n(z), \quad (3.92)$$

The scalar equation of motion (3.92) can be brought into a Schrödinger-like form with the substitution

$$S_n(z) = e^{\omega_s/2} s_n(z), \quad (3.93)$$

with  $\omega_s$  as defined in (3.48). The eigenvalue equation becomes

$$-\partial_z^2 s_n(z) + \left(\frac{1}{4}\omega_s'^2 - \frac{1}{2}\omega_s'' - \frac{3}{2}\frac{\kappa\chi^2(z)}{z^2} - \frac{3}{z^2}\right)s_n(z) = m_{S_n}^2 s_n(z), \quad (3.94)$$

with the boundary conditions

$$\lim_{z_0 \rightarrow 0} s_n(z_0) = 0, \quad \partial_z s_n(z \rightarrow \infty) = 0. \quad (3.95)$$

## Vector Mesons

and the equations of motion for the vector and axial fields have the same form as (3.57) and (3.75), respectively.

The results for the vector mesons are shown in Figure ??, and the scalar and axial sectors show similarly good results [?]. These results reproduce the linear trajectories found in the simple soft-wall model [?] for large  $n$ , while improving the agreement of the lower excitations with data.

### 3.2.2 Pions in Modified Soft-Wall Model of AdS/QCD

The mass spectrum for the pseudoscalar mesons was not found in the original paper [?] because the equations of motion are coupled, second order differential equations, and because of some subtleties that arise when considering the representation of the pseudoscalar field. The paper [?] attempted to circumvent these problems by reducing the equations of motion to a single second-order equation, solvable by the shooting method, but their results seemed to miss certain essential features of the pion spectrum. The authors later address this apparent discrepancy in [?]. This modified soft-wall model was later completed in a paper to which the author contributed, [?], which clarified the discrepancies between two common representations of the pseudoscalar field, calculated the pion mass spectrum to good accuracy, and derived the Gell-Mann–Oakes–Renner relation from the model.

As mentioned above, the field  $X$  contains both the field representing the scalar mesons,  $S(x, z)$ , and the field representing the pseudoscalars,  $\pi(x, z)$ , as well as a non-trivial  $z$ -dependent vacuum expectation value,  $\chi(z)$ . There are two common ways to represent this field:

$$X_e = \left( \frac{\chi(z)}{2} + S(x, z) \right) I e^{2i\pi_e(x, z)^a t^a} \quad (3.96)$$

$$X_l = \left( \frac{\chi(z)}{2} + S(x, z) \right) I + i\pi_l(x, z)^a t^a \quad (3.97)$$

with  $I$  the  $N_f \times N_f$  identity matrix and  $t^a$  the  $SU(N_f)$  generators. We will refer to  $X_e$  as the exponential representation and  $X_l$  as the linear representation. Apparent differences between the representations arise as we note that  $\pi_e$  and  $\pi_l$  are of different dimension. In addition, the linear representation has a quartic interaction term in the

Lagrangian, in contrast to the exponential representation. Despite these differences, we will show that the equations of motion derived from each representation are equivalent.

### Exponential Representation

Let us take (3.97) and substitute it into (3.77), keeping terms that include the field  $\pi(x, z)$ , as well as terms that will mix with  $\pi$ ,

$$\begin{aligned}\mathcal{L}_e = & -\sqrt{-g}e^{-\Phi(z)}\frac{1}{2}\delta^{ab}\left(g^{MN}(\chi^2\partial_M\pi\partial_N\pi + \chi^2A_MA_N - 2\chi^2\partial_MA_N)\right. \\ & \left. + \frac{g^{MP}g^{NR}}{g_5^2}(\partial_MA_N\partial_PA_R - \partial_MA_N\partial_RA_P)\right) + \dots\end{aligned}\quad (3.98)$$

We work in the axial gauge,  $A_z = 0$ , and separate the action (3.98) into four-dimensional components and  $z$ -dependent terms,

$$\begin{aligned}\mathcal{L}_e = & -\frac{1}{2}e^{-\Phi(z)}\left[\sqrt{-g}g^{\mu\nu}(\chi^2\partial_\mu\pi\partial_\nu\pi + \chi^2A_\mu A_\nu - 2\chi^2\partial_\mu\pi A_\nu)\right. \\ & + \sqrt{-g}g^{zz}\chi^2\partial_z\pi\partial_z\pi + \frac{\sqrt{-g}g^{\mu\nu}g^{\rho\sigma}}{g_5^2}(\partial_\mu A_\rho\partial_\nu A_\sigma - \partial_\mu A_\rho\partial_\sigma A_\nu) \\ & \left. + \frac{\sqrt{-g}g^{zz}g^{\mu\nu}}{g_5^2}(\partial_z A_\mu\partial_z A_\nu)\right].\end{aligned}\quad (3.99)$$

We separate  $A_\mu$  into its transverse and longitudinal components:  $A_\mu = A_{\mu\perp} + \partial_\mu\varphi$ , where  $\partial_\mu A_\perp^\mu = 0$ . Expressing the action in terms of the longitudinal component,  $\varphi$  gives

$$\begin{aligned}\mathcal{L}_e = & -\frac{1}{2}e^{-\Phi(z)}\left[\sqrt{-g}g^{\mu\nu}(\chi^2\partial_\mu\pi\partial_\nu\pi + \chi^2\partial_\mu\varphi\partial_\nu\varphi - 2\chi^2\partial_\mu\pi\partial_\nu\varphi)\right. \\ & \left. + \sqrt{-g}g^{zz}\chi^2\partial_z\pi\partial_z\pi + \frac{\sqrt{-g}g^{zz}g^{\mu\nu}}{g_5^2}(\partial_z\partial_\mu\varphi\partial_z\partial_\nu\varphi)\right].\end{aligned}\quad (3.100)$$

Varying (3.100) with respect to  $\pi$  gives

$$\delta\mathcal{L}_e = \partial_z e^{-\Phi}\sqrt{-g}g^{zz}\chi^2\partial_z\pi\delta\pi + e^{-\Phi}\sqrt{-g}\chi^2g^{\mu\nu}\partial_\nu\partial_\mu(\pi - \varphi)\delta\pi. \quad (3.101)$$

Using a Kaluza-Klein decomposition,

$$\pi(x, z) = \sum_n \Pi_n(x)\pi_n(z), \quad (3.102)$$

$$\varphi(x, z) = \sum_n \Phi_n(x)\varphi_n(z), \quad (3.103)$$

and Proca's equation

$$\partial^2 \Pi_n(x) = m_n^2 \Pi_n(x), \quad \partial^2 \Phi_n(x) = m_n^2 \Phi_n(x), \quad (3.104)$$

we can express the system of equations in terms of its  $z$ -dependent parts

$$e^\Phi \partial_z \left( \frac{e^{-\Phi} \chi^2}{z^3} \partial_z \pi_n \right) + \frac{\chi^2 m_n^2}{z^3} (\pi_n - \varphi_n) = 0. \quad (3.105)$$

Varying (3.100) with respect to  $\varphi$  and breaking it down into KK modes gives the second equation of motion

$$e^\Phi \partial_z \left( \frac{e^{-\Phi}}{z} \partial_z \varphi_n \right) + \frac{g_5^2 L^2 \chi^2}{z^3} (\pi_n - \varphi_n) = 0. \quad (3.106)$$

As usual, we express (3.105) and (3.106) in a Schrödinger-like form,

$$\pi \rightarrow e^{f(z)} \pi \quad f(z) = \Phi(z) + \log \frac{z^3}{\chi(z)^2} \quad (3.107)$$

$$\varphi \rightarrow e^{\omega(z)} \varphi \quad \omega(z) = \Phi(z) + \log z. \quad (3.108)$$

After simplifying, the equations of motion become

$$\begin{aligned} -\pi_n'' + \left( \frac{\Phi'^2}{4} - \frac{\Phi''}{2} - \frac{\Phi' \chi'}{\chi} + \frac{3\Phi'}{2z} + \frac{15}{4z^2} - \frac{3\chi'}{\chi z} + \frac{\chi''}{\chi} - m_n^2 \right) \pi_n &= -m_n^2 \frac{\chi^2 L^2}{z^2} \pi_n \\ -\varphi_n'' + \left( \frac{\Phi'^2}{4} - \frac{\Phi''}{2} + \frac{\Phi'}{2z} + \frac{3}{4z^2} + \frac{g_5^2 \chi^2 L^2}{z^2} \right) \varphi_n &= g_5^2 \pi_n \end{aligned} \quad (3.110)$$

### Linear Representation

When considering the linear representation of the pseudoscalar field (3.97), there are terms quadratic and quartic in  $\pi$  that were not present in the exponential representation.

After making the appropriate substitutions in the Lagrangian, it becomes

$$\begin{aligned} \mathcal{L}_l &= -\frac{1}{2} e^{-\Phi} \sqrt{-g} \left( g^{\mu\nu} \partial_\mu \pi \partial_\nu \pi + g^{zz} \partial_z \pi \partial_z \pi - 2\chi g^{\mu\nu} \partial_\mu \pi \partial_\nu \varphi + m_X^2 \pi^2 - \frac{\kappa}{2} \chi^2 \pi^2 \right. \\ &\quad \left. + g^{\mu\nu} \chi^2 \partial_\mu \varphi \partial_\nu \varphi + \frac{g^{\mu\nu} g^{zz}}{g_5^2} \partial_z \partial_\mu \varphi \partial_z \partial_\nu \varphi \right). \end{aligned} \quad (3.111)$$

Following the same procedure as above, we derive two coupled equations. Varying with respect to  $\varphi$  produces

$$e^\Phi \partial_z \left( \frac{e^{-\Phi}}{z} \partial_z \varphi_n \right) + \frac{g_5^2 L^2 \chi}{z^3} (\pi_n - \chi \varphi_n) = 0. \quad (3.112)$$

Varying with respect to  $\pi$  gives the second equation of the linear representation

$$z^3 e^\Phi \partial_z \left( \frac{e^{-\Phi}}{z^3} \partial_z \pi_n \right) - \left( \frac{m_X^2}{z^2} - \frac{\kappa L^2 \chi^2}{2z^2} \right) \pi_n + m_n^2 \pi_n = m_n^2 \chi \varphi_n. \quad (3.113)$$

We can express (3.112) and (3.113) in a Schrödinger-like form as above with the substitutions

$$\pi_n \rightarrow e^{\omega_s/2} \pi_n \quad (3.114)$$

$$\varphi_n \rightarrow e^{\omega/2} \varphi_n, \quad (3.115)$$

where  $\omega_s$  and  $\omega$  are defined as in (3.48) and (3.55), respectively. Simplifying the equations, we find

$$-\varphi_n'' + \left( \frac{\Phi'^2}{4} - \frac{\Phi''}{2} + \frac{3}{4z^2} + \frac{\Phi'}{2z} + \frac{g_5^2 L^2 \chi^2}{z^2} \right) \varphi_n = \frac{g_5^2 L \chi}{z} \pi_n \quad (3.116)$$

$$-\pi_n'' + \left( \frac{\Phi'^2}{4} - \frac{\Phi''}{2} + \frac{3}{4z^2} + \frac{3\Phi'}{2z} - \frac{\kappa L^2 \chi^2}{2z^2} - m_n^2 \right) \pi_n = -m_n^2 \frac{\chi L}{z} \varphi_n \quad (3.117)$$

### Representation Equivalence

The choice of representation for the pseudoscalar field should not affect any physical results obtained from the model. It is therefore desirable to show that the equations of motion derived from the two representations are equivalent.

We begin by expanding  $X_e$  to first order in the fields

$$\begin{aligned} X_e &= \left( \frac{\chi}{2} + S \right) (1 + 2i\pi_e + \dots) \\ &= \frac{\chi}{2} + S + i\pi_e \chi. \end{aligned} \quad (3.118)$$

Comparing (3.118) to (3.97), we infer that

$$\pi_e \chi(z) \rightarrow \pi_l \quad (3.119)$$

is the relationship between the two representations. Let us substitute  $\pi_e \rightarrow \pi_l/\chi(z)$  into the equations of motion of the exponential representation and attempt to obtain the equations of motion of the linear representation. The substitution into (3.106) immediately yields

$$e^\Phi \partial_z \left( \frac{e^{-\Phi}}{z} \partial_z \varphi \right) + \frac{g_5^2 \chi}{z^3} (\pi_l - \chi \varphi) = 0, \quad (3.120)$$



which is equivalent to (3.112) as expected.

Demonstrating the equivalence of the other two equations requires a bit more analysis. First we substitute for  $\pi_e$  in (3.105) and simplify the expression,

$$\frac{z^3 e^\Phi}{\chi} \partial_z \left( \frac{e^{-\Phi} \chi^2}{z^3} \left( \frac{\pi'_l}{\chi} - \frac{\pi_l \chi}{\chi^2} \right) \right) + m_n^2 (\pi_l - \chi \varphi) = 0, \quad (3.121)$$

which becomes

$$\pi''_l - \left( \Phi' + \frac{3}{z} \right) \pi'_l - \frac{\pi_l}{\chi} \left( \chi'' - \Phi' \chi' - \frac{3}{z} \chi' \right) + m_n^2 (\pi_l - \chi \varphi) = 0. \quad (3.122)$$

Recalling the equation of motion for  $\chi(z)$  (3.81), which does not depend on the pseudoscalar representation:

$$\chi'' - \left( \Phi' + \frac{3}{z} \right) \chi' + \left( \frac{3}{z^2} + \frac{\kappa L^2 \chi^2}{2z^2} \right) \chi = 0. \quad (3.123)$$

Using (3.123) in (3.122), we find

$$\pi''_l - \left( \frac{3}{z} + \Phi' \right) \pi'_l + \left( \frac{3}{z^2} + \frac{\kappa L^2 \chi^2}{2z^2} \right) \pi_l + m_n^2 (\pi_l - \chi \varphi) = 0, \quad (3.124)$$

which is equivalent to the other equation of motion of the linear representation (3.113). The equations of motion are equivalent, confirming that physical results do not depend on the representation.

### Pseudoscalar Mass Eigenvalues

The mass eigenvalues for the pions can be calculated in either the exponential or linear representation, using the numerical matrix method detailed in Appendix A. However, it turns out that the boundary conditions

$$\pi(z_0 \rightarrow 0) = 0 \quad \partial_z \pi(z \rightarrow \infty) = 0 \quad (3.125)$$

$$\varphi(z_0 \rightarrow 0) = 0 \quad \partial_z \varphi(z \rightarrow \infty) = 0 \quad (3.126)$$

are easier to enforce using the linear representation (3.97). This may be due to the relation (??) and the fact that the chiral field  $\chi$  also goes to zero in the UV, making it difficult to enforce the boundary condition on  $\pi_e$  simultaneously. The numerical results for the  $\pi_l$  eigenvalues are shown in Table 3.2.

n	$\pi$ Data (MeV)	$\pi_l$ (MeV)	Large- $n$ $\pi_l$
1	140	143	-
2	$1300 \pm 100$	1557	-
3	$1816 \pm 14$	1887	-
4	2070*	2095	-
5	2360*	2298	2245
6	-	-	2403
7	-	-	2551

Table 3.2: The observed masses [?] and calculated masses using the linear representations. The large- $n$  limit solutions are valid from  $n \approx 4$ . From that point onward, the numerical method used is increasingly inaccurate and fails to find the linear Regge trajectories expected. \*=Appears only in the further states of [?].

For large- $n$  excitations the numerical technique develops problems with the boundary conditions. As the number of oscillations in the eigenfunctions increase for higher  $n$  modes, the routine finds eigenvalues that are skewed to larger values. To uncover the correct asymptotic behavior for large  $n$ , we take the large- $z$  limit of (3.116) and (3.117). As  $n$  increases, the eigenfunction is largely determined by the behavior of the effective potential at large  $z$ . At large  $z$ , the VEV and dilaton behave as

$$\chi(z) = (\alpha + \beta)z \equiv \Gamma \frac{z}{L} \quad (3.127)$$

$$\phi(z) = \lambda z^2. \quad (3.128)$$

To take the large- $z$  limit of the linear representation, we introduce a new dimensionless parameter,  $\xi = \sqrt{\lambda}z$ , and expand in  $\xi$ . In the linear representation, we find that (3.116) and (3.117) at large  $\xi$  become

$$-\pi_k'' + \xi^2 \pi_k = \left( \frac{\kappa \Gamma^2}{2\lambda} - 2 + \frac{m_k^2}{\lambda} \right) \pi_k - \frac{m_k^2 \Gamma}{\lambda} \varphi_k \quad (3.129)$$

$$-\varphi_k'' + \xi^2 \varphi_k = \frac{g_5^2 \Gamma}{\lambda} (\pi_k - \Gamma \varphi_k) \quad (3.130)$$

where (') here indicates differentiation with respect to  $\xi$ . This set of equations has the form of coupled harmonic oscillators, the equations of motion of which are

$$-\varphi_k'' + \xi^2 \varphi_k = (2k + 1) \varphi_k \quad (3.131)$$

$$-\pi_k'' + \xi^2 \pi_k = (2k + 1) \pi_k \quad k = 0, 1, \dots \quad (3.132)$$

We make the reasonable assumption that  $\varphi_k = c_k \pi_k$ , which ensures that (3.129), and (3.130) have solutions. Using the form of (3.131) and (3.132) to solve for  $m_k^2$ , and making use of the fact that  $\Gamma^2 = \frac{4\lambda}{\kappa}$ , we find

$$m_k^2 = g_5^2 \Gamma^2 + (2k+1)\lambda. \quad (3.133)$$

Until now we have not made use of the fact that, in the AdS metric,  $z \geq 0$ . Because of this, the eigenfunctions  $\varphi_n$  and  $\pi_n$  describe *half* harmonic oscillators with half as many modes; therefore, we must take  $k \rightarrow 2k$ . The mass eigenvalues for large  $n$ , where  $n = k + 1$ , in both representations then become

$$m_n^2 = (4n-3)\lambda + g_5^2 \Gamma^2 \quad n = 4, 5, \dots \quad (3.134)$$

which are also listed in Table 3.2 and plotted in Figure 3.1. Combining (3.134) and the numerical technique, we obtain all the pseudoscalar eigenvalues. On inspection, we find that this method should be trusted over the numerical routine for  $n \geq 4$ .

### 3.2.3 Gell-Mann–Oakes–Renner Relation

In this section we explore the Gell-Mann–Oakes–Renner relation numerically. Inserting the established equivalence between the exponential and linear representations,  $\pi_e = \pi_l/\chi(z)$ , into (3.122), we obtain

$$\frac{g_5^2 L^2 \chi^2}{z^2} \partial_z \left( \frac{\pi_l}{\chi} \right) = m_\pi^2 \partial_z \varphi. \quad (3.135)$$

Following the method of [?], we construct a perturbative solution in  $m_\pi$  where  $\varphi(z) = A(0, z) - 1$  and use the established relation

$$f_\pi^2 = -L \frac{\partial_z A(0, z)}{g_5^2 z} \Big|_{z \rightarrow 0}. \quad (3.136)$$

Integrating (3.135) yields

$$\frac{\pi(z)}{\chi(z)} = m_\pi^2 \int_0^z du \frac{u^3}{\chi^2(u)} \frac{\partial_z A(0, u)}{g_5^2 u}. \quad (3.137)$$

The function  $u^3/\chi^2(u)$  has significant support only at small values of  $u \sim \sqrt{m_q/\sigma}$ , where we may use (3.136) to relate the derivative on  $A(0, u)$  to the pion decay constant, so that

$$\frac{\pi_l}{\chi} = -\frac{m_\pi^2 f_\pi^2}{2m_q \sigma}. \quad (3.138)$$

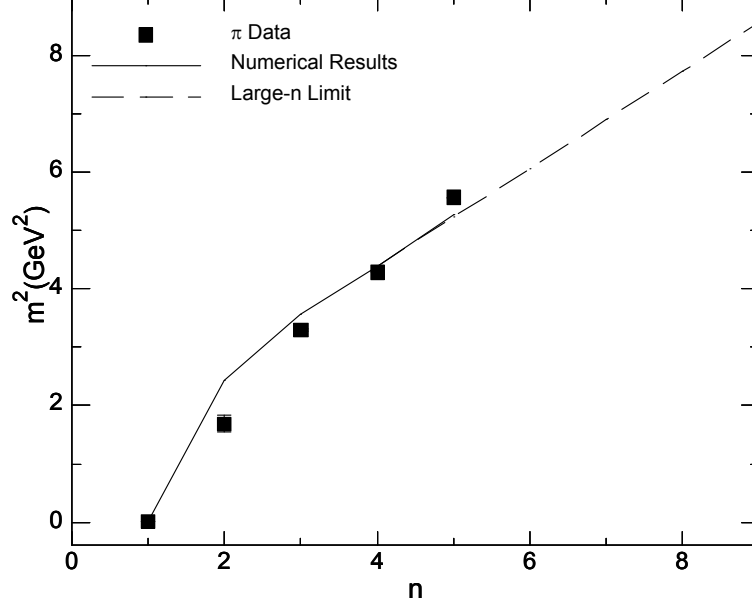


Figure 3.1: The pion mass spectrum calculated in the modified AdS/QCD model is plotted along with the experimental data [?]. The eigenvalues display two important characteristics of the experimental pion spectrum: (1) light ground state and (2) a large gap between the ground state and the first excited state. The large- $n$  mass trajectory clearly follows our calculated eigenvalues from  $n \approx 4$  when our numerical routine inadequately follows the oscillations of the higher eigenfunctions.

We find that letting  $\pi_l = -\chi(z)$  solves the axial-vector field's equation of motion

$$e^\Phi \partial_z \left( \frac{e^{-\Phi}}{z} \partial_z A_\mu(q, z) \right) - \frac{q^2}{z} A_\mu(q, z) - \frac{g_5^2 L^2 \chi^2}{z^3} A_\mu(q, z) = 0 \quad (3.139)$$

in the region of small  $z$  and as  $q \rightarrow 0$ . As a result, (3.138) becomes the expected Gell-Mann–Oakes–Renner (GOR) relation,

$$2m_q \sigma = m_\pi^2 f_\pi^2. \quad (3.140)$$

We solve for the ground-state pseudoscalar mass,  $m_\pi$ , for differing values of  $m_q$  to ensure that the numerical routine respects the GOR relation and gives a reasonable

value for  $f_\pi$ . The results are plotted in Figure 3.2. We see linear behavior in the plot, indicating that as  $m_q \rightarrow 0$  we obtain a constant ratio of  $m_q/m_\pi^2$ . The slope of the line in Figure 3.2 suggests  $f_\pi = 90$  MeV, a result consistent with the input parameters as described in [?].

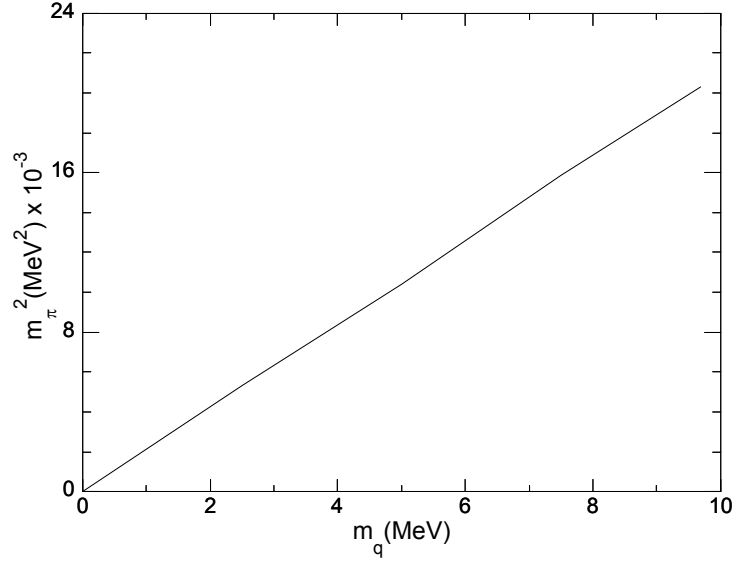


Figure 3.2: Plot of  $m_\pi^2$  vs  $m_q$  yields a straight line from which the pion decay constant  $f_\pi$  is calculated using (3.140).

### 3.3 Summary

In this chapter, we discussed the metric structure and field content of a soft-wall AdS/QCD model. We illustrated the field content of the simplest such model that describes light mesons and their spectra. We discussed the vacuum expectation value of the scalar field, which is related to the chiral symmetry breaking of the theory. The behavior of this chiral condensate was derived, and we discussed the shortcomings apparent in the simplest soft-wall models. We then derived the equations of motion for the scalar, vector, and axial-vector mesons.

We also discussed a modified soft-wall model that allows for the correct form of

chiral symmetry breaking. In this model, we discussed the derivation of the pseudoscalar equations of motion, and the equivalence between the two common representations of the pseudoscalar field. Finally, we calculated the Gell-Mann–Oakes–Renner relation in the linear representation, showing that it is equivalent to that derived in the exponential representation. We confirmed the Gell-Mann–Oakes–Renner relation numerically in this model, and calculated the value of  $f_\pi$ , in good agreement with the accepted value.

## Chapter 4

# Dynamical Three-Field Model

### 4.1 Review and Motivation

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

$$\mathcal{S} = \frac{1}{16\pi G_5} \int d^5x \sqrt{-g} e^{-2\Phi} \left( R + 4\partial_M \Phi \partial^M \Phi - \text{Tr} \left[ |DX|^2 + \partial_M \mathcal{G} \partial^M \mathcal{G} + \frac{1}{2g_5^2} (F_A^2 + F_V^2) + V_m(\Phi, X^2, \mathcal{G}) \right] \right) \quad (4.1)$$

where  $\Phi$  is the dilaton and the metric is pure AdS,  $g_{MN} = z^{-2} \eta_{MN}$ , with the AdS curvature defined to be unity. The constant  $g_5^2 = 12\pi^2/N_c$ , where  $N_c$  is the number of colors. The covariant derivative is defined as  $D_M = \partial_M + i[V_M, X] - i\{A_M, X\}$ . The scalar field,  $X$ , which is dual to the  $\bar{q}q$  operator, obtains a  $z$ -dependent vacuum expectation value (VEV):

$$\langle X \rangle = \frac{\chi(z)}{2} I, \quad (4.2)$$

where  $I$  is the  $2 \times 2$  identity matrix. The glueball field  $\mathcal{G}$  similarly obtains a  $z$ -dependent VEV,  $G(z)$ . We examine the background dynamics of the fields

$$\mathcal{S} = \frac{1}{16\pi G_5} \int d^5x \sqrt{-g} e^{-2\Phi} \left( R + 4\partial_M \Phi \partial^M \Phi - \frac{1}{2} \partial_M \chi \partial^M \chi - \frac{1}{2} \partial_M G \partial^M G - V(\Phi, \chi, G) \right), \quad (4.3)$$

where  $V = \text{Tr}[V_m]$ . The scalar fields  $\Phi, \chi, G$  are dimensionless.

It is easier to search for the background fields in the Einstein frame, where the

vacuum action takes the canonical form

$$\mathcal{S}_E = \frac{1}{16\pi G_5} \int d^5x \sqrt{-\tilde{g}} \left( \tilde{R} - \frac{1}{2} \partial_M \phi \partial^M \phi - \frac{1}{2} \partial_M \chi \partial^M \chi - \frac{1}{2} \partial_M G \partial^M G - \tilde{V}(\phi, \chi) \right). \quad (4.4)$$

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

$$g_{MN} = e^{2\phi/\sqrt{6}} \tilde{g}_{MN}. \quad (4.5)$$

Previous work showed how to construct a potential for a gravity-dilaton-chiral system without the glueball condensate. We examine the behavior assuming that the fields have power-law behavior, which is accurate in both the UV and IR limits, [?]. One of the equations of motion is independent of the choice of potential,

$$\dot{\chi}^2 = \frac{\sqrt{6}}{z^2} \frac{d}{dz} (z^2 \dot{\phi}). \quad (4.6)$$

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 [?]. Inserting this into (4.6), we find that the chiral field behaves as

$$\chi(z) = 6^{3/4} \sqrt{\lambda} z, \quad (4.7)$$

which removes one of the independent parameters of the model in [?]. Using the phenomenological value of  $\lambda$ , 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 models that do not consider the glueball condensate. Models that derive the field behavior using the superpotential method suffer from the same problem.

To resolve this problem, we consider the effects of the glueball condensate  $G$  on the background equations. 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 [?, ?] 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 place the meson fields and chiral dynamics in the open-string sector of the model. For linear confinement, this requires that the chiral field approach a constant in the IR, which necessitates a modified metric to obtain the correct chiral dynamics. Our model allows the metric to remain purely AdS in the string frame. Finally, they do not determine an explicit form of the potential, which is the central goal of this work.

## 4.2 Construction of Potential

Consider the action in the Einstein frame (4.4). To simplify the equations of motion, we use a transformed potential,

$$V = e^{-2\phi/\sqrt{6}} \tilde{V}. \quad (4.8)$$

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

$$V = -12 + 4\sqrt{6}\phi + a_0\phi^2 + \frac{m_X^2}{2}\chi^2 + U. \quad (4.9)$$

Here  $U$  is more than quadratic in the fields. The AdS/CFT dictionary sets the mass for the fields according to the dimension of the dual operator,

$$m^2 L^2 = \Delta(\Delta - 4). \quad (4.10)$$

The dimension of the  $q\bar{q}$  operator is 3, so  $m_X^2 = -3/L^2$ . The dilaton mass is undetermined and is not connected to the dimension of the corresponding operator, as discussed in [?]. 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  $\chi$ .

The equations of motion can be written as

$$\dot{\chi}^2 + \dot{G}^2 = \frac{\sqrt{6}}{z^2} \frac{d}{dz}(z^2 \dot{\phi}) \quad (4.11)$$

$$U = \frac{1}{2}\sqrt{6}z^2\ddot{\phi} - \frac{3}{2}(z\dot{\phi})^2 - 3\sqrt{6}z\dot{\phi} - 4\sqrt{6}\phi - a_0\phi^2 + \frac{3}{2}\chi^2 \quad (4.12)$$

$$\frac{\partial U}{\partial \phi} = 3z\dot{\phi} - 2a_0\phi \quad (4.13)$$

$$\frac{\partial U}{\partial \chi} = z^2 \ddot{\chi} - 3z\dot{\chi} \left( 1 + \frac{z\dot{\phi}}{\sqrt{6}} \right) + 3\chi \quad (4.14)$$

$$\frac{\partial U}{\partial G} = z^2 \ddot{G} - 3z\dot{G} \left( 1 + \frac{z\dot{\phi}}{\sqrt{6}} \right) \quad (4.15)$$

We assume that the potential has no explicit dependence on the coordinate  $z$ , so the equations 4.13-4.15 are not independent, and we can eliminate one.

#### 4.2.1 IR Limit

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

$$\phi = \lambda z^2 \quad (4.16)$$

$$\chi = Az \quad (4.17)$$

$$G = Bz. \quad (4.18)$$

Substitution into (4.11) gives

$$A^2 + B^2 = 6\sqrt{6}\lambda \quad (4.19)$$

The  $\lambda$  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

$$\begin{aligned} A &= 6^{3/4} \sqrt{\lambda} \cos \theta \\ B &= 6^{3/4} \sqrt{\lambda} \sin \theta, \end{aligned} \quad (4.20)$$

where  $\theta$  now becomes the parameter controlling the axial-vector – vector mass splitting. Inserting (4.18) into (4.12-4.15) suggests the following terms in our ansatz for the potential

$$U = a_1 \phi \chi^2 + a_2 \phi G^2 + a_3 \chi^4 + a_4 G^4 + a_5 \chi^2 G^2 + a_6 G^2 \tanh(g\phi). \quad (4.21)$$

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 4.8, 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:

$$U \rightarrow 6 + a_0 + 6\sqrt{6} (\cos^2 \theta a_1 + \sin^2 \theta a_2)$$

$$+6^3 (\cos^4 \theta a_3 + \sin^4 \theta a_4 + \cos^2 \theta \sin^2 \theta a_5) = 0 \quad (4.22)$$

$$\frac{\partial U}{\partial \chi} \rightarrow 2a_1 + 24\sqrt{6} \cos^2 \theta a_3 + 12\sqrt{6} \sin^2 \theta a_5 + \sqrt{6} = 0 \quad (4.23)$$

$$\frac{\partial U}{\partial G} \rightarrow 2a_2 + 24\sqrt{6} \sin^2 \theta a_4 + 12\sqrt{6} \cos^2 \theta a_5 + \sqrt{6} = 0 \quad (4.24)$$

$$\frac{\partial U}{\partial G} \rightarrow a_6 = -\frac{3}{2} \quad (4.25)$$

We have chosen to exclude (4.13) because it is not independent. The parameter  $a_6$  is determined, and the others will be determined by an examination of the UV limit.

#### 4.2.2 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

$$\chi = \Sigma_0 z^3 \quad (4.26)$$

$$G = G_0 z^4. \quad (4.27)$$

Substitution into (4.11) and imposing the boundary condition  $\phi(0) = 0$  gives

$$\phi = \frac{\sqrt{6}}{28} \Sigma_0^2 z^6 + \frac{\sqrt{6}}{27} G_0^2 z^8 \quad (4.28)$$

Using only this leading-order behavior in (4.12-4.15), 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  $\chi$ . Substituting into (4.11) and keeping only the lowest-order cross-term we find the additional term in  $\phi$

$$\Delta\phi = \frac{\sqrt{6}n\Sigma_0\Sigma_n}{(n+4)(n+3)} z^{n+3} \quad (4.29)$$

From (4.12) we find that

$$U = -\frac{3}{2}(z\dot{\phi})^2 - a_0\phi^2 + 3\frac{n^3 - 13n + 12}{(n+4)(n+3)}\Sigma_0\Sigma_n z^{n+3} \quad (4.30)$$

Since the  $\phi^2$  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  $\partial U/\partial\chi$ .

$$\frac{\partial U}{\partial\chi} = -9\Sigma_0 \left( \frac{3}{14}\Sigma_0^2 + \frac{8}{27}G_0^2 z^2 \right) z^9 + (n-3)(n-1)\Sigma_n z^n \quad (4.31)$$

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  $\phi$

$$\Delta\phi = \frac{8mG_0G_m}{\sqrt{6}(m+5)(m+4)} z^{m+4} \quad (4.32)$$

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

$$\frac{\partial U}{\partial G} = -12G_0 \left( \frac{3}{14}\Sigma_0^2 + \frac{8}{27}G_0^2 z^2 \right) z^{10} + m(m-4)G_n z^m \quad (4.33)$$

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  $\chi$  and  $G$  leads to

$$U = -\frac{3}{2}(z\dot{\phi})^2 - a_0\phi^2 + 3\frac{n^3 - 13n + 12}{(n+4)(n+3)}\Sigma_0\Sigma_n z^{n+3} + \frac{4m(m-4)}{m+4}G_0G_m z^{m+4} \quad (4.34)$$

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

$$\chi \sim \Sigma_0 z^3 + \Sigma_0^3 z^9 + G_0^2 \Sigma_0 z^{11} + \dots$$

$$G \sim G_0 z^4 + \Sigma_0^2 G_0 z^{10} + G_0^3 z^{12} + \dots$$

That is,  $\chi$  is an odd function of  $\Sigma_0$  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:

$$\begin{aligned}
U_{\text{LHS}} &= 3\Sigma_0^4 z^{12} \left[ 4 \frac{\Sigma_9}{\Sigma_0^3} - \frac{(54 + a_0)}{2^3 \cdot 7^2} \right] \\
&+ \frac{1}{7} \Sigma_0^2 G_0^2 z^{14} \left[ 120 \frac{G_{10}}{\Sigma_0^2 G_0} + 120 \frac{\Sigma_{11}}{\Sigma_0 G_0^2} - \frac{(72 + a_0)}{9} \right] \\
&+ 2G_0^4 z^{16} \left[ 12 \frac{G_{12}}{G_0^3} - \frac{(96 + a_0)}{3^5} \right]
\end{aligned} \tag{4.35}$$

$$\begin{aligned}
U_{\text{RHS}} &= \Sigma_0^4 z^{12} \left[ \frac{\sqrt{6}}{28} a_1 + a_3 \right] \\
&+ \Sigma_0^2 G_0^2 z^{14} \left[ \frac{\sqrt{6}}{27} a_1 + \frac{\sqrt{6}}{28} (a_2 + g a_6) + a_5 \right] \\
&+ G_0^4 z^{16} \left[ \frac{\sqrt{6}}{27} (a_2 + g a_6) + a_4 \right]
\end{aligned} \tag{4.36}$$

$$\left( \frac{\partial U}{\partial \chi} \right)_{\text{LHS}} = 3\Sigma_0^3 z^9 \left[ -\frac{9}{14} + 16 \frac{\Sigma_9}{\Sigma_0^3} \right] + 8\Sigma_0 G_0^2 z^{11} \left[ -\frac{1}{3} + 10 \frac{\Sigma_{11}}{\Sigma_0 G_0^2} \right] \tag{4.37}$$

$$\left( \frac{\partial U}{\partial \chi} \right)_{\text{RHS}} = \Sigma_0^3 z^9 \left[ \frac{\sqrt{6}}{14} a_1 + 4a_3 \right] + \Sigma_0 G_0^2 z^{11} \left[ \frac{2\sqrt{6}}{27} a_1 + 2a_5 \right] \tag{4.38}$$

$$\left( \frac{\partial U}{\partial G} \right)_{\text{LHS}} = 6\Sigma_0^2 G_0 z^{10} \left[ -\frac{3}{7} + 10 \frac{G_{10}}{\Sigma_0^2 G_0} \right] + 32G_0^3 z^{12} \left[ -\frac{1}{9} + 3 \frac{G_{12}}{G_0^3} \right] \tag{4.39}$$

$$\left( \frac{\partial U}{\partial G} \right)_{\text{RHS}} = \Sigma_0^2 G_0 z^{10} \left[ \frac{\sqrt{6}}{14} (a_2 + g a_6) + 2a_5 \right] \tag{4.40}$$

$$+ G_0^3 z^{12} \left[ \frac{2\sqrt{6}}{27} (a_2 + g a_6) + 4a_4 \right] \tag{4.41}$$

Altogether, from both the UV and IR limits, there are ten independent equations for the twelve parameters  $a_0 - a_6$ ,  $\Sigma_9$ ,  $\Sigma_{11}$ ,  $G_{10}$ ,  $G_{12}$ , 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

$$a_0 = \frac{3}{2} \frac{1}{6 + \sin^2 \theta} \left[ 120 + 62 \sin^2 \theta + 63\sqrt{6}g \sin^2 \theta \right] \tag{4.42}$$

$$a_1 = -\frac{3\sqrt{6}}{4} \frac{1}{6 + \sin^2 \theta} \left[ 12 + 8 \sin^2 \theta + 9\sqrt{6}g \sin^2 \theta \right] \quad (4.43)$$

$$a_2 = -\frac{\sqrt{6}}{4} \frac{1}{6 + \sin^2 \theta} \left[ 32 + 24 \sin^2 \theta + 3\sqrt{6}g(9 \sin^2 \theta - 2) \right] \quad (4.44)$$

$$2a_3 \cos^2 \theta + a_5 \sin^2 \theta = \frac{1}{24} \frac{1}{6 + \sin^2 \theta} \left[ 24 + 22 \sin^2 \theta + 27\sqrt{6}g \sin^2 \theta \right] \quad (4.45)$$

$$2a_4 \sin^2 \theta + a_5 \cos^2 \theta = \frac{1}{24} \frac{1}{6 + \sin^2 \theta} \left[ 20 + 22 \sin^2 \theta + 3\sqrt{6}g(9 \sin^2 \theta - 2) \right] \quad (4.46)$$

$$a_6 = -\frac{3}{2} \quad (4.47)$$

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  $g$ ,  $\theta$ ,  $G_0$ ,  $\Sigma$ , and  $\lambda$ .

### 4.3 Numerical Solution

Using the above potential, we seek a numerical solution that simultaneously satisfies the UV and IR limits. We use equations (4.11, 4.14, 4.15), which allows for an additional term in the potential,  $\Delta U$ , such that

$$\frac{\partial}{\partial \chi} \Delta U = \frac{\partial}{\partial G} \Delta U = 0, \quad (4.48)$$

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 (4.11) allows constant and divergent terms

$$\Delta\phi(z) = c_1 + c_2 z^{-1}. \quad (4.49)$$

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:

$$\Psi(z) = \psi(z)_{UV} f(z) + \psi(z)_{IR} (1 - f(z)), \quad (4.50)$$

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  $\phi$ ,  $\chi$ , and  $G$ . The switching functions  $f$  need not be the same for each field. We choose

$$f_\phi(z) = e^{-(\beta_1 z)^{10}} \quad (4.51)$$

$$f_\chi(z) = e^{-(\beta_2 z)^4} \quad (4.52)$$

$$f_G(z) = e^{-(\beta_3 z)^5}. \quad (4.53)$$

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  $\beta_i$  will be set by numerical fitting.

The chiral condensate  $\Sigma$  is set using the Gell-Mann–Oakes–Renner relation:

$$(m_u + m_d)\Sigma = f_\pi^2 m_\pi^2. \quad (4.54)$$

Using  $m_\pi = 139.6$  MeV,  $f_\pi = 92$  MeV, and  $m_u + m_d = 7.0$  MeV yields a value of  $\Sigma = (286 \text{ MeV})^3$ .

In all, we have eight 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  $\rho$  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 4.11, 4.14, and 4.15.

Three of the parameters are most phenomenologically relevant:  $\lambda$ , which controls the slope of the meson spectra in the large- $n$  limit;  $\theta$ , which controls the mass splitting between the  $a_1$  and  $\rho$  mesons, and  $\beta_2$ , 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 4.1.

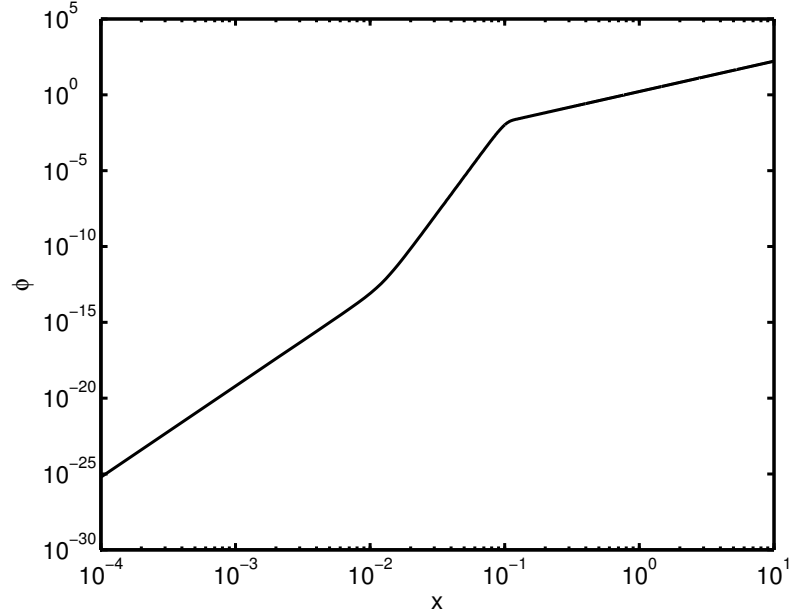


Figure 4.1: A plot of the dilaton field  $\Phi$  generated by the parameterization (4.52). The UV and IR asymptotic behavior is apparent. The coordinate  $x$  is a dimensionless re-scaling of the conformal coordinate,  $x = \sqrt{\lambda}z$ .

The background fields that are obtained from this analysis are shown in Figures 4.1-4.3. 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  $\beta_1$ , which controls the value of  $z$  at which the field transitions from the UV limit to the IR limit.



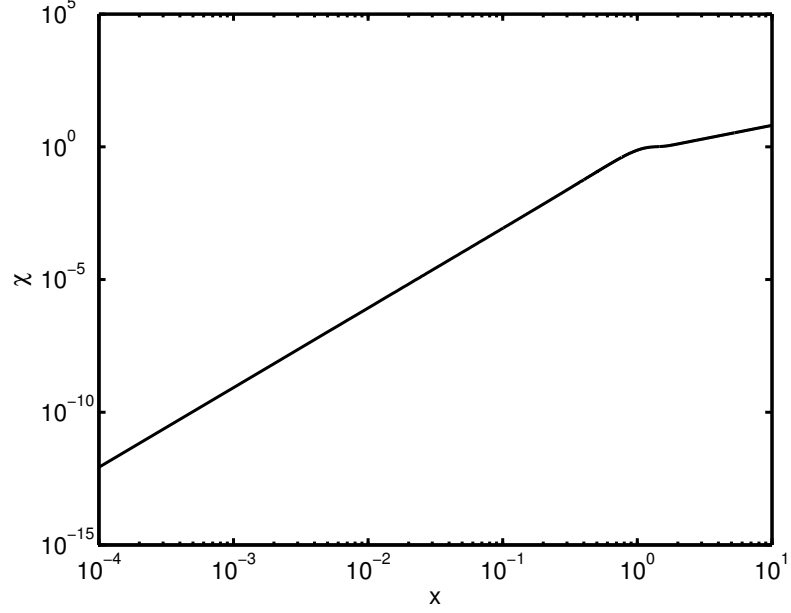


Figure 4.2: A plot of the chiral field  $\chi$  generated by the parameterization (4.53). The UV and IR asymptotic behavior is apparent, with a rapid transition between them. The coordinate  $x$  is a dimensionless re-scaling of the conformal coordinate,  $x = \sqrt{\lambda}z$ .

We now analyze the “extra” term in the potential,  $\Delta U$ . We obtain this term numerically by subtracting the right-hand side of 4.12 from its left-hand side. This term can be approximated numerically as a function of the dilaton field,

$$\Delta U(\Phi) = \alpha_1 \Phi^2 e^{-(\Phi-\gamma_1)^2/\delta_1} + \alpha_2 \Phi^2 e^{-(\Phi-\gamma_2)^2/\delta_2}. \quad (4.55)$$

The best-fit values for these parameters are shown in Table 4.3.

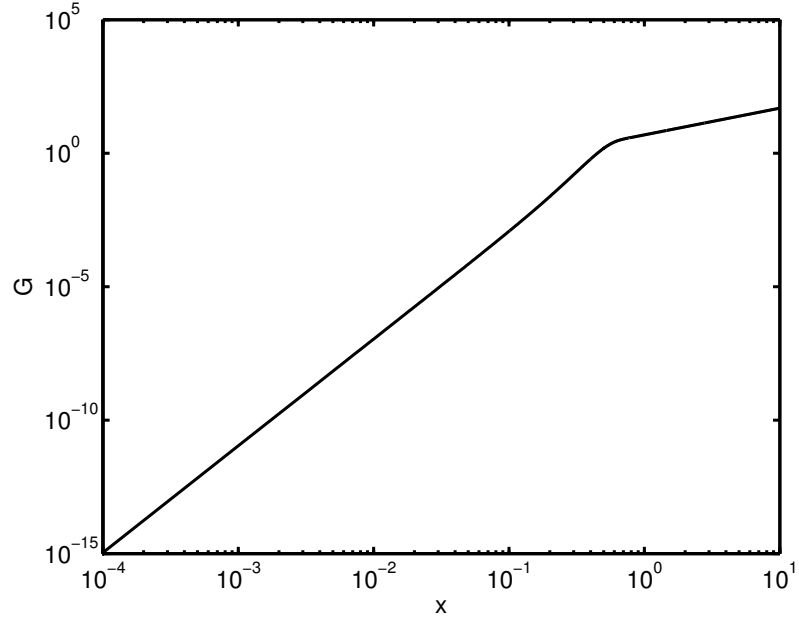


Figure 4.3: A plot of the glueball field  $G$  generated by the parameterization (4.53). The UV and IR asymptotic behavior is apparent, with a rapid transition between them. The coordinate  $x$  is a dimensionless re-scaling of the conformal coordinate,  $x = \sqrt{\lambda}z$ .

$\lambda^{1/2}$	304 MeV	$\beta_1$	3.04 GeV
$G_0^{1/4}$	552 MeV	$\beta_2$	274 MeV
$\theta$	1.44	$\beta_3$	558 MeV
$g$	3.20	$a_5$	1.63

Table 4.1: Best fit parameters for the phenomenological model. The parameters  $\lambda, \theta$ , and  $\beta_2$  are chosen for the best visual fit to the  $\rho$  and  $a_1$  data, with the rest set by minimizing the error in the equations of motion (4.11), (4.14-4.15).

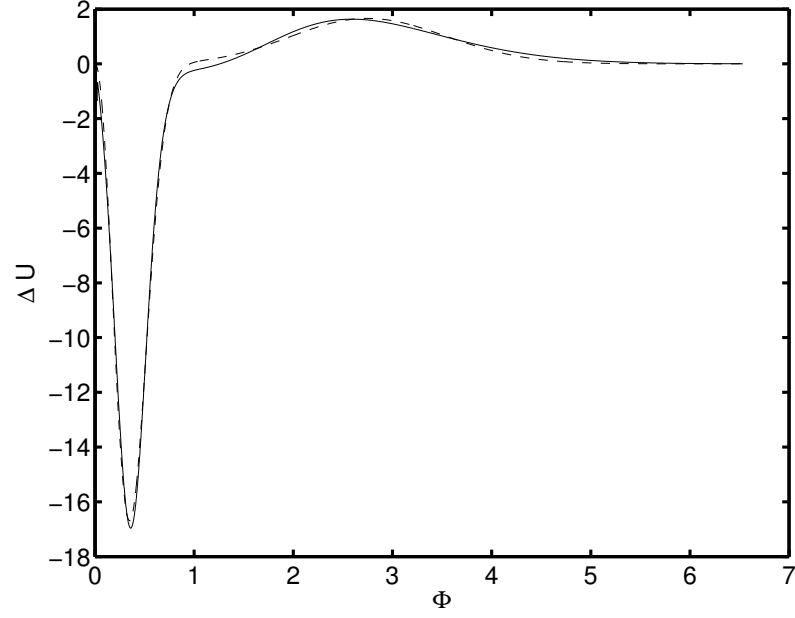


Figure 4.4: Plot of the “extra” term in the potential,  $\Delta U(\phi)$ . The solid line represents the numerical result, while the dashed line is the fitting of (4.55) using the parameters of Table 4.3.

$\alpha_1$	$-304.3 \pm 5 \times 10^{-1}$
$\gamma_1$	$0.07086 \pm 4 \times 10^{-3}$
$\delta_1$	$0.09699 \pm 1.3 \times 10^{-3}$
$\alpha_2$	$0.2671 \pm 6.4 \times 10^{-3}$
$\gamma_2$	$2.213 \pm 4 \times 10^{-2}$
$\delta_2$	$1.471 \pm 9 \times 10^{-2}$

Table 4.2: The dimensionless parameters for the fitting to  $\Delta U$

## Chapter 5

# Meson Spectra

To calculate the spectra of the radial excitations of the mesons, we examine the relevant terms from the string frame action (4.1),

$$\mathcal{S}_{meson} = -\frac{1}{16\pi G_5} \int d^5x \sqrt{-g} e^{-2\Phi} \text{Tr} \left[ |DX|^2 + V_m(\Phi, X^2, \mathcal{G}) + \frac{1}{2g_5^2} (F_A^2 + F_V^2) \right]. \quad (5.1)$$

The  $2 \times 2$  scalar field  $X$  contains the scalar and pseudoscalar fields  $(S, \pi)$ , as well as the VEV.

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

$$-\partial_z^2 \Psi_n + \omega' \partial_z \Psi_n + M_\Psi^2(z) \Psi_n = m_{\Psi_n}^2 \Psi_n, \quad (5.2)$$

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

$$M_A^2 = \frac{g_5^2 L^2 \chi^2}{z^2} \quad (5.3)$$

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

$$-\partial_z^2 \psi_n + \left( \frac{1}{4} \omega'^2 - \frac{1}{2} \omega'' + M_\psi^2 \right) \psi_n = m_{\Psi_n}^2 \psi_n. \quad (5.4)$$

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 [?], particularly for the ground state  $\rho$  meson, as shown in Figure 5.2. The scalar mesons are expected to mix with the scalar glueball field of this model, and this analysis is not performed here.

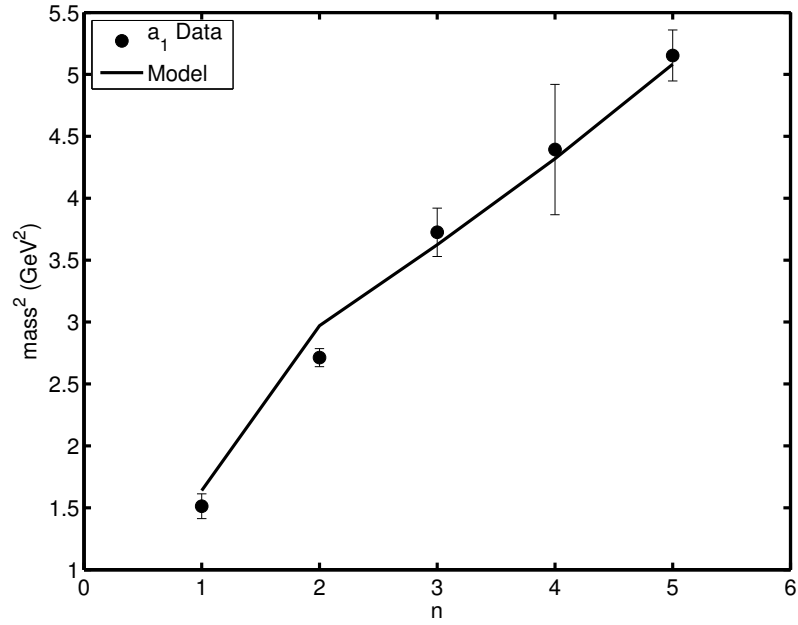


Figure 5.1: Comparison of the predicted mass eigenvalues for the axial-vector sector with the experimental  $a_1$  meson spectrum [?].

n	$a_1$ experimental (MeV)	$a_1$ model
1	$1230 \pm 40$	1280
2	$1647 \pm 22$	1723
3	$1930^{+30}_{-70}$	1904
4	$2096 \pm 122$	2078
5	$2270^{+55}_{-40}$	2254

Table 5.1: The experimental and predicted values for the masses of the axial-vector mesons.

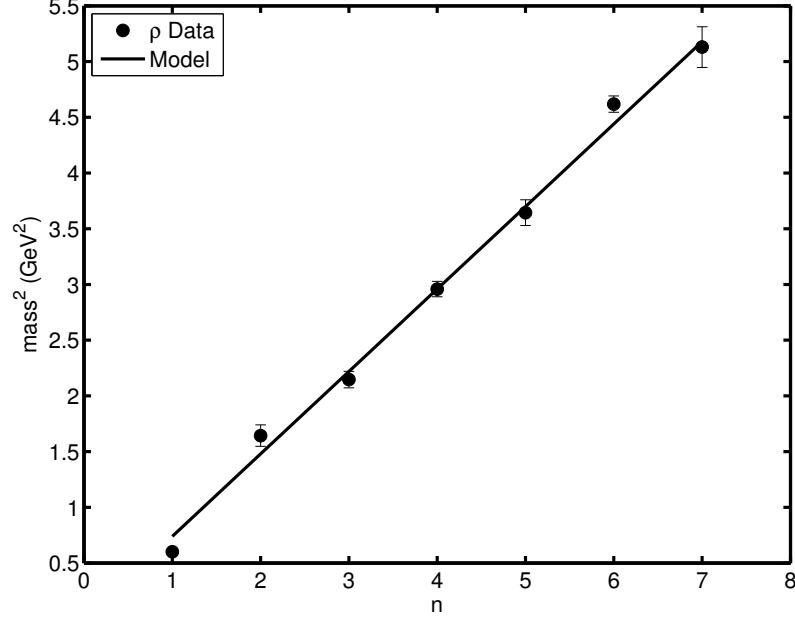


Figure 5.2: Comparison of the predicted mass eigenvalues for the vector sector with the experimental  $\rho$  meson spectrum [?].

## 5.1 Pseudoscalar Sector

When using the exponential representation for the scalar field, the terms from the potential do not contribute to the equations of motion for the pion field. This can be easily seen by noting that  $|X_e|^n$  does not contain any terms involving the  $\pi_e$  field when  $n$  is even. We have required the potential to be an even function of  $X$ , so there are no such terms. This would seem to suggest that we use the exponential representation to calculate the pion mass spectrum. However, as noted in [?],  $\pi_e$  is extremely sensitive to boundary conditions, and the numerical results are not reliable. For this reason, we seek to work with an equation of motion written in the linear representation.

For convenience, we begin by deriving the equations of motion in the exponential representation. Working in the axial gauge  $A_z = 0$ , we rewrite the axial meson field in terms of its perpendicular and longitudinal components:  $A_\mu = A_{\mu\perp} + \partial_\mu\varphi$ . Only the longitudinal component of the axial field,  $\varphi$ , contributes to the pion equations of

n	$\rho$ experimental (MeV)	$\rho$ model
1	$775.5 \pm 1$	860
2	$1282 \pm 37$	1216
3	$1465 \pm 25$	1489
4	$1720 \pm 20$	1720
5	$1909 \pm 30$	1923
6	$2149 \pm 17$	2107
7	$2265 \pm 40$	2276

Table 5.2: The experimental and predicted values for the masses of the vector mesons.

motion. We use (5.1), keeping only the relevant terms

$$\mathcal{L} = e^{-2\Phi} \sqrt{-g} [\chi^2 (\partial_\mu \pi_e \partial^\mu \pi_e + \partial_\mu \varphi \partial^\mu \varphi - 2 \partial_\mu \pi_e \partial^\mu \varphi + \partial_z \pi_e \partial^z \pi_e) + \frac{1}{g_5^2} \partial_z \partial_\mu \varphi \partial^z \partial^\mu \varphi]. \quad (5.5)$$

Varying with respect to  $\varphi$  yields

$$e^{2\Phi} \partial_z \left( \frac{e^{-2\Phi}}{z} \partial_z \varphi \right) + \frac{g_5^2 \chi^2}{z^3} (\pi_e - \varphi) = 0, \quad (5.6)$$

while varying  $\pi_e$  gives

$$\frac{e^{2\Phi} z^3}{\chi^2} \partial_z \left( \frac{e^{-2\Phi} \chi^2}{z^3} \partial_z \pi_e \right) + m_n^2 (\pi_e - \varphi) = 0. \quad (5.7)$$

It was shown in [?] that the equations of motion are equivalent under the substitution  $\pi_e \rightarrow \pi_l / \chi(z)$ , so we make the appropriate substitution and expand the equations:

$$-\varphi'' + \left( 2\Phi' + \frac{1}{z} \right) \varphi' = \frac{g_5^2 \chi}{z^2} (\chi \varphi - \pi_l) \quad (5.8)$$

$$-\pi_l'' + \left( 2\Phi' + \frac{3}{z} \right) \pi_l' + \left( \chi'' - 2\chi' \Phi' - \frac{3\chi'}{z} \right) \frac{\pi_l}{\chi} = m_n^2 (\pi_l - \chi \varphi). \quad (5.9)$$

We can put these equations into Schödinger-like form with the following substitutions

$$\varphi = e^{\omega/2} \varphi_n \quad (5.10)$$

$$\pi_l = e^{\omega_s/2} \pi_n, \quad (5.11)$$

with  $\omega = 2\Phi + \log(z)$  and  $\omega_s = 2\Phi + 3\log(z)$ . This yields

$$-\varphi_n'' + \left( \frac{1}{4} \omega'^2 - \frac{1}{2} \omega'' + \frac{g_5^2 \chi^2}{z^2} \right) \varphi_n = \frac{g_5^2 \chi}{z} \pi_n \quad (5.12)$$

$$-\pi_n'' + \left( \frac{1}{4}\omega_s'^2 - \frac{1}{2}\omega_s'' + \frac{\chi''}{\chi} - \frac{2\chi'\Phi'}{\chi} - \frac{3\chi'}{z\chi} - m_n^2 \right) \pi_n = -m_n^2 \frac{\chi}{z} \varphi_n. \quad (5.13)$$

The dependence of these equations of motion on the scalar potential can be made explicit by using the background equation for the chiral field, written here in the string frame

$$z^2\chi'' - 3z\chi' \left( 1 + \frac{z\Phi'}{\sqrt{6}} \right) = m_X^2\chi + \frac{\partial U}{\partial \chi}. \quad (5.14)$$

Substituting, we can re-write (5.13) as

$$-\pi_n'' + \left( \frac{1}{4}\omega_s'^2 - \frac{1}{2}\omega_s'' + \frac{m_X^2}{z^2} + \frac{1}{z^2} \frac{\partial U}{\partial \chi} - m_n^2 \right) \pi_n = -m_n^2 \frac{\chi}{z} \varphi_n \quad (5.15)$$

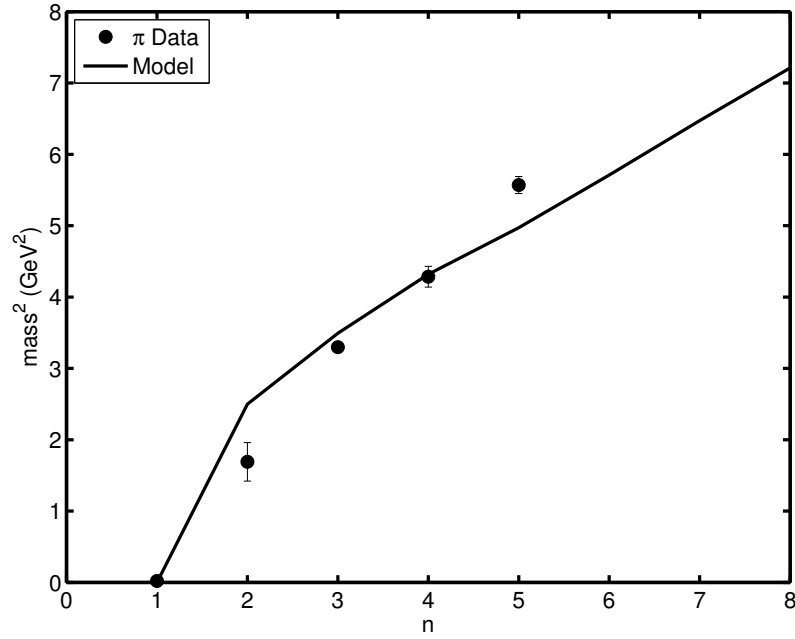


Figure 5.3: Comparison of the predicted mass eigenvalues for the scalar sector with the experimental  $\pi$  meson spectrum.



n	$\pi$ experimental (MeV)	$\pi$ model
1	140	0
2	$1300 \pm 100$	1580
3	$1816 \pm 14$	1868
4	$2070 \pm 35^*$	2078
5	$2360 \pm 25^*$	2230
6	—	2389
7	—	2544
8	—	2686

Table 5.3: The experimental and predicted values for the masses of the pseudoscalar mesons. \*Appears only in the further states of the PDG.

## Chapter 6

# Scalar and Glueball Sectors

The terms of the action relevant for the scalar meson and glueball fields have the form

$$\mathcal{S} = -\frac{1}{16\pi G_5} \int d^5x \sqrt{-g} e^{-2\Phi} \text{Tr} [|DX|^2 + \partial_M \mathcal{G} \partial^M \mathcal{G} + V_m(\Phi, X^2, \mathcal{G})]. \quad (6.1)$$

Here, we are using the exponential representation for the scalar field,

$$X = \left( S(x, z) + \frac{\chi(z)}{2} \right) e^{2i\pi(x, z)}, \quad (6.2)$$

where  $S(x, z)$  is the scalar field fluctuations and  $\chi(z)$  is the chiral condensate. The glueball field  $\mathcal{G}$  also has a condensate, written as

$$\mathcal{G} = \left( H(x, z) + \frac{G(z)}{2} \right), \quad (6.3)$$

where  $H(x, z)$  are the fluctuations of the glueball field and  $G(z)$  is the glueball condensate.

The potential  $V_m$  is defined in such a way that  $\langle \text{Tr} V_m \rangle = V$ , where  $V$  is the potential for the background fields in the string frame,

$$V = -12 + 4\sqrt{6}\phi + a_0\phi^2 + \frac{m_X^2}{2}\chi^2 + U + \Delta U, \quad (6.4)$$

where  $\phi$  is the re-scaled version of the dilaton,  $\phi = \sqrt{8/3}\Phi$ . Here,  $U$  is more than quadratic in the fields,

$$U = a_1\phi\chi^2 + a_2\phi G^2 + a_3\chi^4 + a_4G^4 + a_5\chi^2 G^2 + a_6 G^2 \tanh(\phi), \quad (6.5)$$

and  $\Delta U$  is assumed to be a function of  $\phi$  only.

The terms of  $\text{Tr}V_m$  that are relevant for the analysis of the scalar sector are those terms involving  $X, \mathcal{G}$ ,

$$\text{Tr}V_m \sim m_X^2 X^2 + 2a_1 \phi X^2 + 2a_2 \phi \mathcal{G}^2 + 8a_3 X^4 + 8a_4 \mathcal{G}^4 + 8a_5 X^2 \mathcal{G}^2 + 2a_6 \mathcal{G}^2 \tanh g\phi. \quad (6.6)$$

Expanding and keeping only terms that are quadratic in the fields  $S, H$ , we obtain

$$\text{Tr}V_m \sim (m_X^2 + 2a_1 \phi + 12a_3 \chi^2 + 2a_5 \mathcal{G}^2) S^2 + (2a_2 \phi + 12a_4 \mathcal{G}^2 + 2a_5 \chi^2 + 2 \tanh(g\phi)) H^2 + 4a_5 G \chi H S. \quad (6.7)$$

We can simplify the expression by writing

$$\text{Tr}V_m \sim \left( m_X^2 + \frac{\partial^2 U}{\partial \chi^2} \right) S^2 + \frac{\partial^2 U}{\partial G^2} H^2 + 4a_5 G \chi H S, \quad (6.8)$$

where

$$\frac{\partial^2 U}{\partial \chi^2} = 2a_1 \phi \chi + 4a_3 \chi^3 + 2a_5 \chi \mathcal{G}^2 \quad (6.9)$$

$$\frac{\partial^2 U}{\partial G^2} = 2a_2 \phi G + 4a_4 G^3 + 2a_5 \chi^2 G + 2a_6 G \tanh(g\phi) \quad (6.10)$$

The equations of motion are as follows. Varying with respect to  $S$  gives the equation of motion

$$\partial_z(z^{-3} e^{-2\Phi} S') - z^{-5} e^{-2\Phi} \left( m_X^2 + \frac{\partial^2 U}{\partial \chi^2} \right) S - 4a_5 z^5 e^{-2\Phi} G \chi H = -z^3 e^{-3} m_S^2 S. \quad (6.11)$$

Varying with respect to  $H$  yields the equation of motion

$$\partial_z(z^{-3} e^{-2\Phi} H') - z^{-5} e^{-2\Phi} \frac{\partial^2 U}{\partial G^2} H - 4a_5 z^5 e^{-2\Phi} G \chi S = -z^3 e^{-3} m_H^2 H. \quad (6.12)$$

To put the equations in Schrödinger form, we make the substitutions

$$S \rightarrow e^{\omega_s/2} S \quad (6.13)$$

$$H \rightarrow e^{\omega_s/2} H, \quad (6.14)$$

where  $\omega_s = 2\Phi + 3 \log z$ . This reduces the equations of motion to

$$-S'' + \left( \frac{1}{4} \omega_s'^2 - \frac{1}{2} \omega_s'' + m_X^2 - \frac{\partial^2 U}{\partial \chi^2} \right) S + 4a_5 G \chi H = m_S^2 S \quad (6.15)$$

$$-H'' + \left( \frac{1}{4} \omega_s'^2 - \frac{1}{2} \omega_s'' - \frac{\partial^2 U}{\partial G^2} \right) H + 4a_5 G \chi S = m_H^2 H \quad (6.16)$$

These equations need to be solved simultaneously for the eigenvalues  $m_S^2, m_H^2$ .

## Chapter 7

# Conclusion and Discussion

## Appendix A

# Numerical Methods for solving Ordinary Differential Equations

Ordinary differential equations can always be reduced to a system of coupled first-order differential equations, which is advantageous for numerical solutions. For example, the second-order equation

$$y''(x) + f(x)y'(z) = g(x) \tag{A.1}$$

can be rewritten as

$$y'(x) = z(x) \tag{A.2}$$

$$z'(x) = g(z) - f(x)z(x) \tag{A.3}$$

Written generically, an ordinary differential equations system is reduced to a set of  $N$  first-order differential equations

$$y'_i(x) = f_i(x, y_1, \dots, y_N) \quad i = 1, \dots, N \tag{A.4}$$

where the functions  $f_i$  are some known set of equations involving the coordinate  $x$  and the dependent variables  $y_i$ . For a unique solution, one must specify a set of  $N$  boundary conditions. If these conditions are all specified at a single point,  $x_0$ , this is known as an initial value problem, and one can simply choose an integration method and calculate the values for  $f_i$  over the desired domain. When some of the conditions are specified at point  $x_0$  and the rest are specified at the point  $x_f$ , this is known as a boundary value problem, and one must use one of the methods detailed below.

## A.1 Shooting Method

The shooting method turns boundary value problems into initial value problems. The classic example of a shooting method can be visualized as the firing of a cannon, with boundary conditions given by the location of the cannon and the location of the target. This is changed to an initial value problem by selecting an arbitrary value for the angle of the cannon and firing. The angle of the cannon is incremented until the cannonball hits the target, which matches the final boundary condition.

More concretely, the boundary value problem is written as in (A.4), with boundary conditions

$$B_{1j}(x_0, y_1, \dots, y_N) = 0 \quad j = 1, \dots, n_1 \quad (\text{A.5})$$

$$B_{2k}(x_f, y_1, \dots, y_N) = 0 \quad k = 1, \dots, n_2 \quad (\text{A.6})$$

with  $n_1$  boundary conditions at point  $x_0$  and  $n_2 = N - n_1$  at point  $x_f$ . These boundary conditions generically can be any algebraic combination of the variables.

For concreteness, consider a second-order eigenvalue problem of the type considered in this thesis,

$$\psi_n''(z) + V(z)\psi_n(z) = m_n^2\psi_n(z). \quad (\text{A.7})$$

This can be reduced to two first-order differential equations by making the substitutions  $y_1 = \psi$ ,  $y_2 = \psi'$ , resulting in the system

$$y_1' = y_2 \quad (\text{A.8})$$

$$y_2' = (m_n^2 - V(z)) y_1. \quad (\text{A.9})$$

The boundary conditions are  $y_1(z_0) = y_1(z_f) = 0$ .<sup>1</sup> In this instance, the eigenvalues are of greater importance than the behavior of the solutions  $y_1, y_2$ , so the overall normalization of the solutions is arbitrary.

This boundary value problem is changed to an initial value problem by setting

$$y_1(z_0) = 0 \quad (\text{A.10})$$

$$y_2(z_0) = c, \quad (\text{A.11})$$

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<sup>1</sup> In the soft-wall model, we should take the limit  $z_f \rightarrow \infty$ . This is impossible to do numerically, but it is easy to make  $z_f$  sufficiently large that the eigenvalues are not effected.

where  $c$  is an arbitrary constant. A small test value for  $m_n$  is chosen, and the initial value problem is solved.

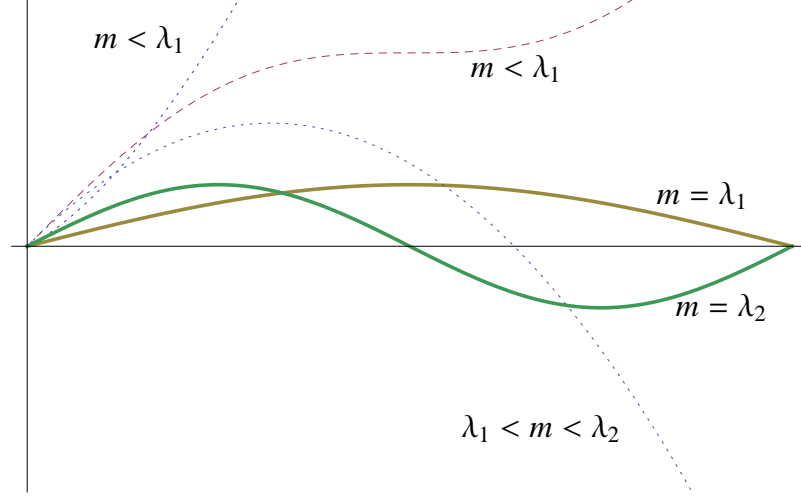


Figure A.1: An illustration of the shooting method.

As illustrated in Figure A.1, when the test value for  $m_n$  is less than the first eigenvalue  $\lambda_1$ , the value of  $y_1(x_f)$  is positive. As  $m_n$  is incremented, eventually  $y_1(x_f)$  becomes negative. A root-finding routine is used to approximate the values  $\lambda_n$  such that  $y_1(x_f) = 0$ . The number of antinodes in the wavefunction indicates to which excitation mode the eigenvalue corresponds.

## A.2 Matrix Method

The coupled second-order The equations of motion can be reduced to a set of second-order differential equations,

$$-\varphi'' + V_1(z)\varphi + f(z)\pi = 0, \quad (\text{A.12})$$

$$-\pi'' + V_2(z)\pi + g(z)\varphi = 0, \quad (\text{A.13})$$

where the eigenvalues are contained within the coefficient functions. These equations can be expressed as a system of first-order differential equations

$$\Phi' + W(z)\Phi = 0, \quad (\text{A.14})$$

where  $W$  is the matrix

$$W = \begin{pmatrix} 0 & 1 & 0 & 0 \\ V_1(z) & 0 & f(z) & 0 \\ 0 & 0 & 0 & 1 \\ g(z) & 0 & V_2(z) & 0 \end{pmatrix} \quad (\text{A.15})$$

and  $\Phi$  is the vector

$$\Phi_{\alpha i} = \begin{pmatrix} \varphi_i \\ -\varphi'_i \\ \pi_i \\ -\pi'_i \end{pmatrix} \quad (\text{A.16})$$

that forms an orthonormal basis of solutions. We can propagate the solution  $\Phi$  between two boundary points

$$\Phi(z_f) = U(z, z_f, z_0, m_n^2) \Phi(z_0), \quad (\text{A.17})$$

where we solve (A.14) with the appropriate boundary condition at  $z_0$ . The eigenvectors and eigenvalues of  $U$  are then calculated. Two of the solutions are divergent, but the solutions that correspond to the two smaller eigenvalues do not diverge. We label the non-divergent solutions as the eigenvectors  $u_3$  and  $u_4$ . The non-divergent solutions for  $\Phi_i$  can be written as a linear combination of  $u_3$  and  $u_4$

$$\Phi_i = \alpha u_3 + \beta u_4. \quad (\text{A.18})$$

Non-trivial values of  $\alpha$  and  $\beta$  satisfy

$$\begin{pmatrix} u_3^1 & u_4^1 \\ u_3^3 & u_4^3 \end{pmatrix} \begin{pmatrix} \alpha \\ \beta \end{pmatrix} = 0 \quad (\text{A.19})$$

for Dirichlet or

$$\begin{pmatrix} u_3^2 & u_4^2 \\ u_3^4 & u_4^4 \end{pmatrix} \begin{pmatrix} \alpha \\ \beta \end{pmatrix} = 0 \quad (\text{A.20})$$

for Neumann boundary conditions. The values of  $m_n^2$  that minimize the determinant of the matrix (A.19) or (A.20) are the eigenvalues of the system (A.12-A.13). An abrupt change in its behavior of the plot of the determinant vs. the value of  $m_n$  signals the location of an eigenvalue.