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Light meson masses using AdS/QCD modified soft wall model

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We analyze light vector and scalar meson mass spectra using a novel approach where a modified soft wall model with a UV cutoff is considered. Including this cutoff introduces an extra energy scale. For this model, we found that the masses for the scalar and vector spectra are well fitted within a very small root mean square (RMS) error for 14 of these states, with nonlinear trajectories given by two common parameters, the UV locus z_0, and the quadratic dilaton profile slope κ. We concluded that in this model the $f_0(500)$ scalar resonance cannot be fitted holographically as a $q\bar{q}$ state since we could not find a trajectory that included this pole. This result is in agreement with the most recent phenomenological and theoretical methods.

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I. INTRODUCTION

The idea of using the AdS/CFT correspondence [1,2] to describe nonperturbative QCD-like phenomena has provided insight into exploring the strong interactions at strong coupling, unreachable by regular quantum field theory (QFT) methods. One possibility is considering gravity models in a given space that holographically generate low-energy QCD theories living on the conformal flat boundary. This proposal is called top-down. The second scenario considers the opposite: starting from well-known properties derived from a four-dimensional QCD, one tries to look out for a five-dimensional theory living on the AdS space, which is a holographic dual model; this is the so-called bottom-up approach. Both cases provide valuable effective models since they permit us to create a bigger landscape for a fundamental nonperturbative theory that is unknown at present.

One example of those nonperturbative phenomena is related to the dynamics of the lightest pseudoscalar mesons. A very useful effective field theory approach that describes it is given by the momentum-expansion formalism of chiral perturbation theory (ChPT), where a $SU(N_f)_L \otimes SU(N_f)_R \to SU(N_f)_V$-symmetric nonlinear sigma model (where $N_f = 2, 3$) written in terms of a meson multiplet is expanded up to a certain perturbative order; this procedure introduces a diagrammatic way to study scattering events between these particles in a particular range of energy [3–6]. A phenomenological description is attained after fitting the parameters of the model [masses, decay constants, and low-energy constants (LECs)] to an adequate set of experimental data, e.g., phase shifts, scattering lengths, or imaginary parts of the associated amplitudes.

The energy range mentioned above can be extended after unitarizing the partial waves of the scattering channels involved, thus including the respective resonances as poles in the complex plane [7]. This method checks elastic unitarity in an approximate way (order by order in the expansion), although other approaches in the momentum expansion allow us to check exactly this feature, as happens with the inverse amplitude method (IAM) [8,9], in which pole positions are quite well described, especially those respecting pion-pion scalar and vector channels.

Some of these resonances are properly analyzed as vector and tensor mesons since their structure is easily fitted as a Breit-Wigner distribution due to its $q\bar{q}$-like compositeness. However, quite the opposite happens with light scalar resonances ($I = 0$, $J^{PC} = 0^{++}$) produced below an energy close to 2 GeV since they are not easily characterized as $q\bar{q}$ mesons due to the large decay widths of some of these particles, as happens with the f_0 multiplet [10]. This lies in the model-dependent descriptions of the nature of these particles, along with the inappropriate values for their masses and widths, as happens with the $f_0(500)$ when considering it as a $q\bar{q}$ meson in a $N_f = 2$ linear sigma model [11]; nevertheless, recent approaches provide strong insight into the most likely nontrivial quark composition of this particle [12].

Regarding pole positions, a proper model-independent description of the $f_0(500)$ and $f_0(980)$ resonance parameters is achieved by using an adequate set of dispersion relations with minimal uncertainty [13,14], thus theoretically minimizing the errors both in their masses and decay widths. Pole positions for the $f_0(500)$, $f_0(980)$, and $f_0(1400/1370)$ can also be obtained through scattering matrix approaches, with the results depending on the way the couplings between scattering channels are taken into account [15].

Quark composition of resonances like $f_0(980)$, $f_0(1370)$, $f_0(1500)$, and $f_0(1710)$ can be studied via analysis of decay widths of B mesons; in order to achieve this, these resonances have to be parametrized as superpositions of u, d, s quarkonium states and a scalar glueball so that a perturbative QCD-effective Hamiltonian is to be

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built up, using their masses as the input parameters of the model [in this case, $f_0(980)$, $f_0(1370)$, and $f_0(1710)$ are predominantly quarkonia] [16]. The $f_0(1500)$ is usually described as a glueball state since it does not decay into two photons [unlike what happens with the $f_0(1370)$] [17] and its mass coincides with lattice simulations [18].

Chiral effective models with a scalar glueball state are also used to study both compositeness and masses of scalar resonances where $m > 1.2$ GeV by considering experimental inputs such as the masses and composition of the scalar and pseudoscalar meson multiplets [19,20], with the results for quarkonia-glueball compositeness mixing for $f_0(1500)$ and $f_0(1710)$ depending on whether a glueball decay is or is not considered. In both cases, the theoretical masses of the resonances are quite close to the experimental values. Similar results for this mixing are obtained when considering lattice masses for quarkonia and glueball as input parameters [21].

Mass generation for scalar resonances can be analyzed using a linear sigma model with two quark flavors, including axial-vector mesons, a glueball degree of freedom, and two parameters that explicitly break chiral and dilation symmetries, associated respectively to the $f_0(1370)$ and the $f_0(1500)$ (referred to as a scalar field that is related with the trace anomaly) [22]. The results obtained after taking experimental inputs for quarkonia and glueball masses and widths come along with the compositeness mixing for these particles, hence giving that the $f_0(1500)$ is mostly a glueball state. For this case, the $f_0(1370)$ has a theoretical mass less than the lower experimental bound. A better result is obtained if the $f_0(1710)$ is considered as a glueball; however, this is discarded since the predicted value for its four-pion decay width is large (something that has not been observed). If three quark flavors are to be taken [23], then the $f_0(1500)$ is considered as a heavy strange quarkonium, whereas the $f_0(1710)$ is largely composed of a glueball state. In this case, the glueball is coupled to meson states and mixed with two quarkonia states. After taking proper experimental inputs, the masses and widths of these three scalar particles are predicted within less than 10% and 5% of uncertainty, respectively. These results, along with the quarkonia-glueball mixing, are independent of the fit considered.

In the case of bottom-up approaches, the most successful ones describing nonperturbative phenomena are the so-called AdS/QCD models, such as the hard wall (HW) [24] or the soft wall (SW) models [25] that are able to describe mass spectra, electromagnetic form factors, some decay constants, and other mesonic properties. The main idea behind these models is to break the conformal invariance in AdS by placing a cutoff, thus introducing an energy scale. When the cutoff is a D-brane, the model is called a hard wall, and when a quadratic dilaton is used instead, a soft wall model is obtained.

Results in the soft wall model show that masses grow linearly with the excitation number, which gives a Regge trajectory. This mass spectrum appears due to the confining potential created by the quadratic dilaton profile [26]. When dealing with the hard wall model, the masses are given by the zeroes of Bessel functions generated by the Dirichlet boundary conditions imposed at the wall/brane, yielding nonlinear trajectories [27]. Light vector meson masses are described in [25] and scalar light mesons were described in [28] in the soft wall model framework. These descriptions are not so good since they do not fit the particle mass spectra well, although mesons are organized in Regge trajectories [25,28].

Other soft wall approaches that consider scalar fields with variable masses (along with chiral symmetry effects) reproduce remarkable theoretical predictions for the light scalar sector when parameters such as quark masses and chiral condensates are introduced; however, in order to obtain these results, nonphysical values have to be taken into account for these sets of parameters [29]. This issue is properly solved when a scalar potential is introduced [30]. Both of these results, besides reproducing quite well the light vector sector, consider the f_0 multiplet belonging to a Regge trajectory. A previous development [31] takes into account a constant scalar mass in the action, although in order to reproduce correctly the scalar sector, one of the first two resonances [either the $f_0(500)$ or the $f_0(980)$] has to be removed.

Recently, a new approach was developed in [32], where the usual soft wall model is upgraded by including an extra UV cutoff given by a D-brane. This extra brane will work as the boundary where the particles live and also will fix, together with the dilatonic energy scale, the mass and decay constant spectra of the particles. The application of this idea gives good results describing the first four vector states of charmonium and also the first four of bottomonium with a total error close to 30% for fitting eight quarkonium states with three parameters [32]. The extension to finite temperature of this model gives a complete holographic view of the melting processes of these heavy quarkonium states, with results in agreement with the observed phenomenology [33].

This paper is organized as follows. We introduce the holographic bottom-up model in Sec. II to describe the light scalar and vector meson resonances as poles of a two-point function. We show the main results of the model in Sec. III, regarding scalar and vector mass spectra, along with their respective error percentages when compared with experimental data. Finally, we present our conclusions in Sec. IV.

II. HOLOGRAPHIC MODEL FOR LIGHT MESONS

In order to describe light mesons, we will consider the usual SW model action [25,28]

$$I = -\frac{1}{2g_S^2}\int d^5x\sqrt{-g}\exp[-\Phi(z)][\partial_n S\partial^n S + m_5^2 S^2] - \frac{1}{4g_V^2}\int d^5x\sqrt{-g}\exp[-\Phi(z)]F_{mn}F^{mn}, \tag{1}$$

where $S(z, x^\mu)$ is a massive scalar field dual to the scalar mesons and $F_{mn} = \partial_m A_n - \partial_n A_m$ is given in terms of the massless Abelian gauge field $A_m(z, x^\mu)$.

The bulk mass fixes the conformal dimension Δ of the p-form QCD operator $\mathcal{O}_s$ dual to the S field as $m_5^2 R^2 = (\Delta - p)(\Delta + p - 4)$. In the simplest case, the scalar operator has the form $\mathcal{O}_s = \bar{q}(x)q(x)$ with dimension 3, where q is any light quark. Thus, we can fix $\Delta = 3$ and $p = 0$ such that $m_5^2 R^2 = -3$ [28].

The geometric background is given by the sliced AdS Poincaré patch [32,33]

$$dS^2 = \Theta(z - z_0)\frac{R^2}{z^2}[dz^2 + \eta_{\mu\nu}dx^\mu dx^\nu], \tag{2}$$

with $\Theta(z)$ the Heaviside step function that gives the UV D-brane (D-wall) locus. The Minkowski metric has the signature $(-, +, +, +)$.

This particular choice of boundary for AdS breaks explicitly the conformal invariance by introducing an energy scale z_0, which can be associated to the nature of the strong interaction inside the meson [32]. Such behavior is expected since when we recover the conformal boundary by setting $z_0 \to 0$, the mass spectrum is given by a usual Regge trajectory defined by the form of the dilaton profile, i.e., $M_n^2 = c(n + s + 1)$ [34]. In this case, such a profile corresponds to $\Phi(z) = \kappa^2 z^2$, which is static as in the regular soft wall model.

The constants g_S and g_V fix the units of the action in terms of the number of colors N_c as usual. Since we are not interested in the calculation of the decay constants or any form factor, these constants do not interest us. The proper value for these couplings is read from the large four-momentum expansion of the two-point function in the QCD side compared to the same kind of expansion in the gravity side [25,28].

Following the ideas expressed in [32], we will define the mass spectrum of light scalar and vector mesons as functions of two energy scales, namely, the D-wall locus z_0 and the dilaton constant κ.

A. Light vector mesons

We begin our analysis with the light vector meson action given by

$$I_V = -\frac{1}{4g_V^2}\int d^5x\sqrt{-g}\exp[-\Phi(z)]F_{mn}F^{mn}, \tag{3}$$

according to (1). After considering small variations in the A_μ field and imposing the gauge condition $A_z = 0$, we obtain the equation of motion for the space-time components as

$$\partial_z\left[\frac{\exp(-\kappa^2 z^2)}{z}\partial_z A^\mu\right] + \frac{\exp(-\kappa^2 z^2)}{z}\eta^{\rho\sigma}\partial_\rho\partial_\sigma A^\mu = 0. \tag{4}$$

Equation (4) allows us to obtain a boundary action from (3) for the vector fields that reads

$$I_{\text{V On-Shell}} = -\frac{R}{2g_V^2}\int d^5x\left\{\partial_z\left[\frac{\exp(-\kappa^2 z^2)}{z}A_n\partial_z A^n\right]\right\}. \tag{5}$$

In the latter equation, we have used again the gauge condition $A_z = 0$. According to the Minkowskian prescription, this boundary action (5) gives the two-point function, and its poles define the mass spectrum. From this same equation, we infer that the boundary term (i.e., taking $z = z_0$) is such that

$$I^{\text{Boundary}}_{\text{V On-Shell}} = -\frac{R}{2g_V^2}\int d^4x\frac{\exp(-\kappa^2 z^2)}{z}A_\mu\partial_z A^\mu\Big|_{z_0}. \tag{6}$$

Two-point functions are easily obtained after solving the equation of motion (4) by introducing Fourier transform vector fields

$$A^\mu(z, x^\mu) = \frac{1}{(2\pi)^4}\int d^4q\exp(-iq_\mu x^\mu)v_\mu(z, q), \tag{7}$$

where we write $v_\mu(z, q)$ as a function of the source term $v_\mu^0(q)$ and the bulk-to-boundary propagator $V(z)$ as follows:

$$v_\mu(z, q) = v_\mu^0(q)V(z). \tag{8}$$

Therefore, recalling that $\eta^{\rho\sigma}\partial_\sigma\partial_\rho = -\Box = q^2$, we obtain that $V(z)$ holds with the following:

$$\partial_z\left[\frac{\exp(-\kappa^2 z^2)}{z}\partial_z V(z)\right] + \frac{q^2}{z}\exp(-\kappa^2 z^2)V(z) = 0. \tag{9}$$

The regular solution of (9) reads

$$V(z) = c_1\kappa^2 z^2 {}_1F_1\left(1 - \frac{q^2}{4\kappa^2}, 2, \kappa^2 z^2\right), \tag{10}$$

where ${}_1F_1(1 - q^2/4\kappa^2, 2, \kappa^2 z^2)$ is the Kummer confluent hypergeometric function and c_1 is a normalization constant. Hence, we deduce from the on-shell boundary action

$$I^{\text{Boundary}}_{\text{V On-Shell}} = -\frac{R}{2g_V^2}\int\frac{d^4q}{(2\pi)^4}v_\mu^0(q)v^{\mu 0}(-q)\frac{\exp(-\kappa^2 z^2)}{z}V_z\partial_z V_z\Big|_{z_0} \tag{11}$$

the following vector two-point function $G^{\mu\nu}(q^2)$:

$$G^{\mu\nu}(q^2) = \eta^{\mu\nu}\Pi(q^2), \tag{12}$$

$$\Pi(q^2) = -\frac{R}{g_V^2}\left[\frac{\exp(-\kappa^2 z^2)}{z} V(z)\partial_z V(z)\right]\bigg|_{z_0}. \quad (13)$$

After normalizing (10) such that $V(z_0) = 1$, we finally obtain that $\Pi(q^2)$ reads

$$\Pi(q^2) = -\frac{R\exp(-\kappa^2 z_0^2)}{g_V^2 z_0^2} \times \left[\frac{2}{z_0} + \kappa^2 z_0\left(1 - \frac{q^2}{4\kappa^2}\right)\frac{{}_1F_1(2-\frac{q^2}{4\kappa^2}, 3, \kappa^2 z_0^2)}{{}_1F_1(1-\frac{q^2}{4\kappa^2}, 2, \kappa^2 z_0^2)}\right]. \quad (14)$$

The poles of the two-point function (14) can be read from the roots of the hypergeometric confluent function in the denominator

$${}_1F_1(1-\chi_n, 2, \kappa^2 z_0^2) = 0, \quad (15)$$

with $\chi_n = q_n^2/4\kappa^2$ the root spectrum and $q_n^2 = M_n^2$ the physical masses. Thus, the mass spectrum for the light vector mesons is given by

$$M_{n,\mathrm{V}}^2 = 4\kappa^2\chi_n(z_0, \kappa). \quad (16)$$

The result above assures us that the mass spectrum (16) is given by a nonlinear Regge trajectory defined by the parameters z_0 and κ. In general, the roots of the hypergeometric confluent function increase with n [[35] Sec. 13.9], so the masses increase with the excitation number, as we expected. The results for the light vector masses are shown in Table I.

B. Light scalar mesons

We see that the scalar case follows a procedure similar to that of the vector fields shown in Sec. II A. Thus, we define from (1) the scalar action as

$$I_S = -\frac{1}{2g_S^2}\int d^5x\sqrt{-g}\exp[-\Phi(z)][\partial_n S\partial^n S + m_5^2 S^2], \quad (17)$$

TABLE I. Mass spectrum for ρ vector mesons with $\kappa = 0.45$ GeV and $z_0 = 5\ \mathrm{GeV}^{-1}$. Experimental values are obtained from [10].

ρ	M_{th} (GeV)	M_{exp} (GeV)	$\%M$
$\rho(775)$	0.975	0.775	20.53
$\rho(1450)$	1.455	1.465	0.66
$\rho(1570)$	1.652	1.570	4.96
$\rho(1700)$	1.829	1.720	5.97
$\rho(1900)$	1.992	1.909	4.15
$\rho(2150)$	2.142	2.153	0.50

whose associated equation of motion, after taking small variations in S, taking the gauge condition $A_z = 0$, and replacing the definition of the conformal dimension in terms of m_S, is given by

$$\partial_z\left[\frac{\exp(-\kappa^2 z^2)}{z}\partial_z S\right] - \frac{\exp(-\kappa^2 z_2)}{z^3}\Box S + \frac{3\exp(-\kappa^2 z^2)}{z^5} S = 0, \quad (18)$$

where $\Box = -\eta^{\mu\nu}\partial_\mu\partial_\nu$. We obtain the solution of (18) by considering the Fourier transform of the scalar field as

$$S(x_\mu, z) = \frac{1}{(2\pi)^4}\int \exp(-ix_\mu q^\mu)S(z, q), \quad (19)$$

$$S(z, q) = S^0(q)\bar{v}(z). \quad (20)$$

In this case, the, bulk-to-Boundary propagator is labeled as $\bar{v}(z)$, while the scalar source term is given by $S^0(q)$. Hence, (18) changes into

$$\partial_z\left[\frac{\exp(-\kappa^2 z^2)}{z^3}\partial_z\bar{v}(z)\right] + \frac{\exp(-\kappa^2 z^2)}{z^3}q^2\bar{v}(z) + \frac{3\exp(-\kappa^2 z^2)}{z^5}\bar{v}(z) = 0, \quad (21)$$

whose regular solution is given in terms of the Kummer confluent hypergeometric function as follows:

$$\bar{v}(z) = \bar{c}_1\kappa^3 z^3\,{}_1F_1\left(\frac{3}{2} - \frac{q^2}{4\kappa^2}, 2, \kappa^2 z^2\right). \quad (22)$$

As expected, our solution depends on a normalization constant $\bar{c}_1$. Before showing the normalized solution of the bulk-to-Boundary propagator, we deduce from (17) that the on-Shell Boundary action reads

$$I_{\mathrm{SOn\text{-}Shell}}^{\mathrm{Boundary}} = \frac{R^3}{g_S^2}\int d^4q\frac{\exp(-\kappa^2 z^2)}{z^3}S^0(q)S^0(-q)\bar{v}(z)\partial_z\bar{v}(z)\bigg|_{z_0}. \quad (23)$$

Hence, the scalar two-point function $\Pi_S(q^2)$ is such that

$$\Pi_S(q^2) = -\frac{R^3}{g_S^2}\frac{\exp(-\kappa^2 z^2)}{z^3}\bar{v}(z)\partial_z\bar{v}(z)\bigg|_{z_0}. \quad (24)$$

Our solution for (24), written in terms of a normalized $\bar{v}(z)$ function, is given by

TABLE II. Mass spectrum for f_0 scalar resonances with $\kappa = 0.45$ GeV and $z_0 = 5.0\ \mathrm{GeV}^{-1}$. Experimental values for the masses are read from [10].

f_0	M_{th} (GeV)	M_{exp} (GeV)	$\%M$
$f_0(980)$	1.070	0.99	7.46
$f_0(1370)$	1.284	1.370	5.11
$f_0(1500)$	1.487	1.504	1.13
$f_0(1710)$	1.674	1.723	2.93
$f_0(2020)$	1.846	1.992	7.94
$f_0(2100)$	2.153	2.101	2.39
$f_0(2200)$	2.292	2.189	4.49
$f_0(2330)$	2.424	2.314	4.52

$$\Pi_S(q^2) = -\frac{R^3}{g_S^2}\frac{\exp(-\kappa^2 z_0^2)}{z_0^3} \times \left[\frac{3}{z_0} + \kappa^2 z_0\left(\frac{3}{2} - \frac{q^2}{4\kappa^2}\right)\frac{{}_1F_1(\frac{5}{2} - \frac{q^2}{4\kappa^2}, 3, \kappa^2 z_0^2)}{{}_1F_1(\frac{3}{2} - \frac{q^2}{4\kappa^2}, 2, \kappa^2 z_0^2)}\right]. \tag{25}$$

As in the vector case, we obtain the pole expansion from the roots of the denominator in (25),

$${}_1F_1\left(\frac{3}{2} - \bar{\chi}_n, 2, \kappa^2 z_0^2\right) = 0, \tag{26}$$

with $\bar{\chi}_n = q_n^2/4\kappa$. Therefore, the mass spectrum is given by

$$M_{n,S}^2 = 4\kappa^2 \bar{\chi}_n(z_0, \kappa). \tag{27}$$

Notice that (27) is also nonlinear and defined by the increasing $\bar{\chi}_n$ and the parameters κ and z_0. The results for these mesons are shown in Table II.

III. RESULTS

The respective spectra for vector and scalar resonances is generated after finding the associated poles of the two-point functions (14) and (25). In order to obtain them, we only need to fix two parameters: the boundary radius z_0 and the dilaton slope κ. Following [32], we will fix κ as flavor independent, so we will use the same κ for scalar and vector mesons since they are made of up and down quarks, which in the chiral limit have the same mass. The z_0 parameter is defined as a quantity related to the nature of the strong interactions inside the mesons. Thus, we could use the same value reported in [32], but due to the color screening it is expected that the z_0 parameter would be different for light and heavy quarks.

In this case, we have that the best values that fit the experimental masses [10] correspond to

$$z_0 = 5\ \mathrm{GeV}^{-1}, \tag{28}$$

$$\kappa = 0.45\ \mathrm{GeV}. \tag{29}$$

In Table I, we present the theoretical values calculated with the model proposed in [32], along with the experimental masses and the corresponding uncertainties for the ρ vector meson trajectory. It is interesting to notice that the spectrum is not linear, as in the case of the regular soft wall model [25].

We show in Table II the results for the f_0 trajectory. Again, the spectrum is nonlinear. Notice that the $n = 1$ state is not associated to the $f_0(500)$ state. In this model, it is not possible to fit this resonance into the trajectory (27) with any parameter choice. Thus, since we have related κ and z_0 with the color structure inside mesons, we can conclude that, holographically, the $f_0(500)$ resonance is not a $q\bar{q}$ state. This is in agreement with theoretical phenomenology [12].

Following [32], we can test the predictability of the model developed here with the RMS error for estimating N parameters using N_p parameters as

$$\delta_{\mathrm{RMS}} = \sqrt{\frac{1}{N - N_p}\sum_i^N \left(\frac{\delta\mathcal{O}_i}{\mathcal{O}_i}\right)^2}, \tag{30}$$

where $\mathcal{O}_i$ is the experimental mean value of a given observable and $\delta\mathcal{O}_i$ is the absolute uncertainty given by the model. In our case, we fit up to 14 resonant states with two parameters, thus obtaining an RMS error δ_{RMS}

$$\delta_{\mathrm{RMS}} = 7.64\%. \tag{31}$$

As it can be seen from Tables I and II, the resonances we obtain are not degenerate, as expected from the usual Regge theory. We attain this after carefully choosing the pole positions of the two-point functions (14) and (25) according to their q^2 dependence.

We also want to point out that the approach considered here minimizes the number of parameters to be taken into account since the model (both in the scalar and vector sector) does not deal directly with a certain meson internal structure, as shown in [29,30] [all this information is summarized in the choosing of the (κ, z_0) parameter space]. Pions and axial states are not reproduced since we do not take into account chiral symmetry breaking effects.

IV. CONCLUSIONS

The model we considered does not deal directly with the composition of the scalar mesons, as in the case of the $f_0(1500)$ and $f_0(1710)$, which are glueball candidates. This was not necessary since the poles only depend on the model parameters κ and z_0. Also, the errors we obtained are within the phenomenological bounds given in [23]. We also obtained a remarkable result for the $f_0(980)$ mass, a possible non-$q\bar{q}$ state. However, it was not possible to fit the $f_0(500)$ since our model only considered ordinary light $q\bar{q}$ mesons (both scalar and vector). This means that the model needs to be extended somehow to describe these sorts of scalar particles; a proper description of exotic states such as scalar glueballs can be found in [36]. On the other hand, light vector mesons were well fitted, with the $\rho(770)$ state having the biggest error. As a matter of fact, unlike what happened with the scalar multiplet, the ground state could be determined up to the higher error bound allowed by these sorts of nonconformal models. We also note that all of our results did not need to consider either experimental or lattice input parameters.

We showed here that these AdS/QCD approaches could reproduce light meson spectra after minimizing the amount of holographic and physical parameters; we attained this by analyzing the respective poles of the scalar and vector propagators in such a way that only the dilaton profile κ and the D-wall locus z_0 are needed, thus avoiding the introduction of nonphysical quark masses and condensates, as in [29]. Furthermore, internal properties of mesons were also avoided here since quark masses and condensate-dependent confining potentials [30] were not directly treated. These parameters are, by some unknown form, related with the constituent quark mass and to the naturalness of the strong interaction.

Despite having different values for κ and z_0 for heavy [32] and light mesons, a universality class can be established for these sort of models. In fact, there is a huge phenomenological difference between heavy and light quarks due to the heavy quark symmetry: heavy quark systems are considered nonrelativistic, e.g., Schrödinger-like heavy quarkonium potentials. Also, color screening effects in both systems are different since they strongly depend on the quark masses [37].

In a future work, we want to study finite-temperature chiral symmetry restoration effects in these sorts of models [33] after properly introducing pseudoscalar and axial particles. Our objective is to check if these holographic approaches properly describe phase transitions, as happens with large-N nonlinear sigma models [38,39].

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