Available online at www.sciencedirect.com
ScienceDirect
ELSEVIER
Physics Letters B 652 (2007) 73–78
PHYSICS LETTERS B
www.elsevier.com/locate/physletb

On the light glueball spectrum in a holographic description of QCD

P. Colangelo [a,*], F. De Fazio [a], F. Jugeau [a], S. Nicotri [a,b]

[a] Istituto Nazionale di Fisica Nucleare, Sezione di Bari, Italy
[b] Università degli Studi di Bari, Italy

Received 11 April 2007; accepted 26 June 2007

Available online 12 July 2007

Editor: G.F. Giudice

Abstract

We investigate the spectra of light scalar and vector glueballs in a holografic description of QCD with a dilaton background bulk field. In particular, we study how the glueball masses depend on the conditions on the dilaton background and on the geometry of the bulk.

1. Introduction

A breakthrough in the attempt to understand strongly coupled Yang–Mills theories is represented by the AdS/CFT correspondence conjecture, stating that a connection can be established between the supergravity limit of a superstring/M-theory living on a $d+1$ anti-de Sitter (AdS) space times a compact manifold and the large N limit of a maximally $\mathcal{N}=4$ superconformal $SU(N)$ gauge theory defined in the d dimension AdS boundary [1–4]. However, the application of this conjecture to a theory such as QCD is not straightforward, being QCD neither supersymmetric nor conformal. Witten proposed a procedure to extend duality to such gauge theories [5]: the conformal invariance is broken by compactification (the compactification radius giving rise to a dimensionful parameter, namely the mass gap of QCD), while supersymmetry is broken by appropriate boundary conditions on the compactified dimensions. The AdS geometry of the dual theory is then deformed into an AdS-black-hole geometry where the horizon plays the role of an IR brane. In this (so-called top–down) approach, analyses of the glueball spectrum have been carried out, obtaining, for example, that the operator $\mathrm{Tr}\,F^2$ in four dimensions corresponds to the massless dilaton field in supergravity in ten dimensions, that the scalar glueball with $J^{PC}=0^{++}$ in QCD is related to the dilaton propagating in the black-hole geometry and its mass is computable by solving the dilaton wave equation [6–8]. The numerical results are close to the available lattice data [9].

However, one could adopt the strategy of investigating which features the dual theory should have in order to reproduce known QCD properties. In this (so-called bottom–up) approach, instead of trying to deform the high dimensional theory to obtain a theory in $4d$ with similarities with QCD, one begins with QCD and attempts to construct a five-dimensional holographic dual. A hint to follow is that, although QCD is not itself a conformal theory, it nevertheless resembles a strongly coupled conformal theory in the domain where the quark masses are neglected and the coupling is approximately constant (the possibility that the QCD β function has an infrared fixed point is discussed, e.g., in [10,11]). As pioneered by Polchinski and Strassler, it is possible to implement duality in these nearly conformal conditions defining QCD on the four-dimensional boundary, and introducing a bulk space which is a slice AdS_5, the size of which stands for an IR cutoff associated to the QCD mass gap, the so-called hard IR wall approximation [12]. This procedure was investigated in Refs. [13–17] with the calculation of the light hadron spectrum. Moreover, the glueball spectrum was studied considering various boundary conditions of the associated $5d$ field at the IR brane [17]. The static $Q\bar{Q}$ potential was also worked out [18] together with hadron wave-

* Corresponding author.
E-mail address: pietro.colangelo@ba.infn.it (P. Colangelo).

0370-2693 © 2007 Elsevier B.V. Open access under CC BY license.
doi:10.1016/j.physletb.2007.06.072

functions and form factors [19]. Besides, leaving the hard IR wall and considering a background dilaton field, it was shown that properties of QCD can be reproduced, namely the Regge behaviour of light mesons [20], at odds with what happens starting from a general string theory and attempting to deform it [21].

Even though in the bottom–up approach α_s corrections, running of the coupling constant, geometry of the compact manifold which should be considered together with AdS_5, etc., at present are left aside, there is the hope that the main features of the dual theory can be identified.

The starting point is inspired by a principle of the AdS/CFT correspondence, which establishes a one-to-one correspondence between a certain class of local operators (namely, the chiral primary operators and their superconformal descendants) in the $4d$ $\mathcal{N}=4$ superconformal gauge theory and supergravity fields representing the holographic correspondents in the $AdS_5 \times S^5$ bulk theory [2–4]. Analogously, in the bottom–up approach one attempts to construct a correspondence between QCD local operators and fields in the AdS_5 bulk space. Although in this way the five-dimensional dual of QCD contains an infinite number of fields, in correspondence to the infinite number of QCD operators, it was shown that, considering only few operators relevant for chiral dynamics, a few properties in the light meson sector can be obtained, namely the ρ meson spectrum, the axial-vector meson spectrum, the $\rho\pi\pi$ coupling and a few leptonic constants, with a small number of hadronic parameters.

In this Letter we consider the spectrum of low-lying scalar and vector glueballs in the approach, proposed in [20], where the hard IR cutoff in the AdS_5 space is replaced by a smooth cutoff obtained introducing a dilaton background bulk field. In particular, we study how the glueball masses depend on the conditions on the background dilaton and on the bulk geometry, so that they can be compared to, e.g., lattice QCD or QCD sum rule results [9,22].

2. Model for a $5d$ holographic dual of QCD

Following [20], we consider a five-dimensional conformally flat spacetime (the bulk) described by the metric

$$g_{MN} = e^{2A(z)}\eta_{MN},$$
$$ds^2 = e^{2A(z)}\left(\eta_{\mu\nu}\,dx^\mu\,dx^\nu + dz^2\right) \tag{1}$$

$(M, N = 0, \ldots, 4)$, where $\eta_{MN} = \mathrm{diag}(-1, 1, 1, 1, 1)$, x^μ $(\mu = 0, \ldots, 3)$ represent the usual spacetime (the boundary) coordinates and z is the fifth holographic coordinate running from zero to infinity. The metric function $A(z)$ satisfies the condition

$$A(z) \underset{z\to 0}{\longrightarrow} \ln\left(\frac{R}{z}\right) \tag{2}$$

to reproduce the AdS_5 metric close to the UV brane $z \to 0$; in the following we put to unity the radius R. Besides, we consider a background dilaton field ϕ which only depends on the holographic coordinate z and vanishes at the UV brane. By an appropriate choice of the ϕ dependence in the IR (large values of z) we construct a $5d$ model that can be considered similar to a cutoff AdS space: a smooth cutoff in the IR replaces the hard-wall IR cutoff that would be obtained by allowing the holographic variable z to vary to a maximum value $z_m \simeq \frac{1}{\Lambda_{\mathrm{QCD}}}$. The introduction of a background dilaton allows to avoid ambiguities in the choice of the field boundary conditions at the IR wall.

To investigate the mass spectra of the QCD scalar and vector glueballs, we consider the two lowest dimension operators with the corresponding quantum numbers and defined in the field theory living on the $4d$ boundary:

$$\begin{cases} O_S = \mathrm{Tr}\left(F^2\right), \\ O_V = \mathrm{Tr}\left(F(DF)F\right) \end{cases} \tag{3}$$

(with D the covariant derivative) having conformal dimension $\Delta = 4$ and $\Delta = 7$, respectively. The operator corresponding to the vector glueball satisfies the Landau–Pomeranchuk–Yang selection rule [23]. In the AdS/CFT correspondence the conformal dimension of a (p-form) operator on the boundary is related to the $(AdS\ mass)^2$ of its dual field in the bulk as follows [2,3]:

$$(AdS\ mass)^2 = (\Delta - p)(\Delta + p - 4). \tag{4}$$

In the following we assume that the mass m_5^2 of the bulk fields is given by this expression.

A $5d$ massless scalar field $X(x, z)$ can be constructed as the correspondent of $\mathrm{Tr}\,F^2$, described by the action in the gravitational background:

$$S = -\frac{1}{2}\int d^5x\,\sqrt{-g}\,e^{-\phi(z)}g^{MN}(\partial_M X)(\partial_N X), \tag{5}$$

with $g = \det(g_{MN})$. Scalar glueballs are identified as the normalizable modes of X satisfying the equations of motion obtained from (5), corresponding to a finite action.

For the spin 1 glueball, we introduce a 1-form A_M described by the action:

$$S = -\frac{1}{2}\int d^5x\,\sqrt{-g}\,e^{-\phi(z)} \times \left[\frac{1}{2}g^{MN}g^{ST}F_{MS}F_{NT} + m_5^2 g^{ST}A_S A_T\right], \tag{6}$$

with $F_{MS} = \partial_M A_S - \partial_S A_M$ and $m_5^2 = 24$, and study its normalizable modes. Notice that the action (6), with a different value of m_5^2, describes fields that are dual to other operators in QCD, namely those describing hybrid mesons with spin one, which is an explicit example of different QCD operators having similar bulk fields as holographic correspondents.

In Section 3 we discuss how the spectrum can be worked out. However, before such a discussion, it is interesting to comment on the pseudoscalar glueball, described in QCD by the $\Delta = 4$ operator $O_P = \mathrm{Tr}(F \wedge F)$. Identifying another 0-form in the bulk as the correspondent of O_P and describing it by the same action (5), a degenerate mass spectrum would be obtained for scalar and pseudoscalar glueballs, at odds with the results obtained, e.g., in lattice QCD where it is found that the mass of the lightest scalar glueball is smaller than the mass of the lightest pseudoscalar one.[1] A way out to such a degeneracy issue could be represented by the choice of still considering the relation (4) (keeping in mind that such a relation rigorously holds only in the AdS/CFT correspondence conditions), and attempting to describe also the field corresponding to the pseudoscalar glueball by a massive 1-form A_M. There is an indication for that, since in the top–down approach the pseudoscalar glueball is described by a massless R-R 1-form [8]. The construction followed in that approach is that, after a first compactification of an $11d$ M-theory in $AdS_7 \times S^4$, AdS/CFT establishes a duality between a type IIA string theory in $AdS_6 \times S^4$ and a low-energy effective supersymmetric theory $SU(N)$ described in terms of N coincident D4-branes. On this $5d$ D4-brane worldvolume, the field that couples to Tr $F \wedge F$ is a massless R-R 1-form A_ρ, the coupling term reading as [4]

$$\frac{1}{16\pi^2}\int d^5x\, \varepsilon^{\rho\mu\nu\alpha\beta} A_\rho F_{\mu\nu} F_{\alpha\beta}. \tag{7}$$

Then, a second compactification provides a non-supersymmetric model of QCD in terms of N coincident D3-branes dual to the type IIA string theory in an AdS-black-hole geometry. The mass spectrum of the pseudoscalar glueball is then determined by solving the equation of motion for the R-R 1-form in the $10d$ bulk.[2] We do not continue here in such an analysis, since the issue of parity of various hadronic excitations deserves a dedicated study.

3. Background fields

The metric function $A(z)$ and the background dilaton field $\phi(z)$ can be constrained. A constraint to $A(z)$ is the condition (2) which, together with the requirement $\phi(z) \xrightarrow[z\to 0]{} 0$, allows to reproduce the AdS_5 metric close to UV brane $z \simeq 0$. On the other hand, a suitable large z dependence of $\phi(z)$ can be fixed to reproduce the Regge behaviour of the low-lying mesons. The two conditions

$$\begin{cases} \phi(z) - A(z) \xrightarrow[z\to 0]{} \ln z, \\ \phi(z) - A(z) \xrightarrow[z\to \infty]{} z^2, \end{cases} \tag{8}$$

satisfy the two requirements: indeed, the first condition satisfies Eq. (2), while the second one allows to recover the Regge behaviour of ρ resonances, as shown in [20]. Moreover, the metric function $A(z)$ must not have any contribution growing as z^2 at large z, a condition coming from computing the masses of higher spin mesons [20].

The simplest choice consistent with these constraints[3]

$$\phi(z) = z^2, \qquad A(z) = -\ln z, \tag{9}$$

has been chosen to calculate the spectrum of mesons of spin S and radial quantum number n, with the result: $m_n^2 = 4(n+S)$ [20].

We use these expressions for the background dilaton and the metric function to work out the glueball spectrum. The field equations of motion obtained from the actions (5)–(6) can be reduced in the form of a one-dimensional Schrödinger equation in the variable z:

$$-\psi'' + V(z)\psi = -q^2\psi, \tag{10}$$

involving the function $\psi(z)$ obtained applying a Bogoliubov transformation $\psi(z) = e^{-B(z)/2}\tilde{Q}(q,z)$ to the Fourier transform $\tilde{Q}$ of the field Q ($Q = X, A_M$) with respect to the boundary variables x^μ. The function $B(z)$ is a combination of the dilaton and the metric function: $B(z) = \phi(z) - cA(z)$, with the parameter c given by: $c = 3$ and $c = 1$ in cases of X and A_M fields, respectively. The condition $q^2 = -m^2$ identifies the mass of the normalizable modes of the two fields.

Eq. (10) is a one-dimensional Schrödinger equation where $V(z)$ plays the role of a potential. It reads as:

$$V(z) = \frac{1}{4}\big(B'(z)\big)^2 - \frac{1}{2}B''(z) + \frac{m_5^2}{z^2} = V_0(z) + \frac{m_5^2}{z^2} \tag{11}$$

with

$$V_0(z) = z^2 + \frac{c^2+2c}{4z^2} + c - 1. \tag{12}$$

With this potential Eq. (10) can be analytically solved. Regular solutions at $z \to 0$ and $z \to \infty$ correspond to the spectrum:

$$m_n^2 = 4n + 1 + c + \sqrt{(c+1)^2 + 4m_5^2} \tag{13}$$

with n an integer (we identify it as a radial quantum number), while the corresponding eigenfunctions read as:

$$\psi_n(z) = A_n e^{-z^2/2} z^{g(c,m_5^2)+1/2} \times {}_1F_1\big(-n, g\big(c, m_5^2\big)+1, z^2\big), \tag{14}$$

with ${}_1F_1$ the Kummer confluent hypergeometric function, A_n a normalization factor, and $g(c, m_5^2) = \sqrt{\frac{(c+1)^2}{4} + m_5^2}$. From these relations we obtain the spectrum of scalar and vector glueballs:

$$m_n^2 = 4n + 8, \tag{15}$$

$$m_n^2 = 4n + 12, \tag{16}$$

respectively.

A few remarks are in order. First, both the spectra have the same dependence on the radial quantum number n as the mesons of spin S: this is a consequence of the large z behaviour chosen for the background dilaton. Second, both the lowest lying glueballs are heavier than the ρ mesons, the spectrum of which reads: $m_n^2 = 4n + 4$, as derived in [20]. Finally, the vector glueball turns out to be heavier than the scalar one.

Comparing our result to the computed ρ mass, we obtain for the lightest scalar (G_0) and vector (G_1) glueballs

$$\frac{m_{G_0}^2}{m_\rho^2} = 2, \qquad \frac{m_{G_1}^2}{m_\rho^2} = 3, \tag{17}$$

[1] A parity degeneracy has been pointed out in the light baryon spectrum in the framework of a holographic dual of QCD [13].

[2] The spectrum of the pseudoscalar glueball has also been analyzed using a massive 3-form of the $11d$ supergravity coupled to a $\Delta = 9$ operator of the $6d$ boundary theory [6].

[3] We put to one the scale parameter multiplying z^2 in the dilaton field; mass predictions will be given in units of this parameter.

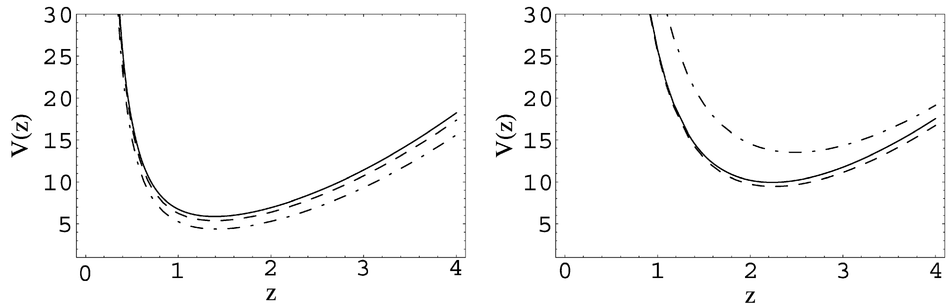

Fig. 1. The unperturbed potential (solid line), and the potential obtained perturbing the dilaton field (dashed line) and the metric (dot-dashed line) for scalar (left) and vector glueball (right) using $\lambda = -0.2$.

which implies that these glueballs are expected to be lighter than as predicted by other QCD approaches [9]. Moreover, the result $m^2_{G_1} - m^2_{G_0} = m^2_\rho$ predicts a lightest vector glueball with mass below 2 GeV.

It is interesting to investigate how it is possible to modify the z dependence of the background dilaton field and of the metric function A, and how the spectra change, an issue discussed in the following section.

4. Perturbed background

There are other choices for the background dilaton ϕ and the metric function A which satisfy the constraints in (8) and may modify the predictions for the scalar and vector glueball masses. As a matter of fact, it is possible to add to the background fields terms of the type z^α with $0 \leqslant \alpha < 2$. Considering the simplest case: $\alpha = 1$, this can be done in two different ways. Instead of using (9), we can modify the dilaton field including a linear contribution which is subleading in the IR regime $z \to \infty$:

$$\begin{aligned}\phi(z) &= z^2 + \lambda z,\\ A(z) &= -\ln z,\end{aligned} \tag{18}$$

with λ a real parameter. Another possibility consists in modifying the metric function,

$$\begin{aligned}\phi(z) &= z^2,\\ A(z) &= -\ln z - \lambda z,\end{aligned} \tag{19}$$

which now acquires a linear term subleading in the UV regime $z \to 0$. The two choices produce different results.

Using the expressions (18), i.e., modifying the dilaton field, the potential (10) becomes:

$$V(z) = V_0(z) + \lambda V_1(z) + \frac{\lambda^2}{4} + \frac{m_5^2}{z^2} f(z,\lambda), \tag{20}$$

with $V_0(z)$ given in Eq. (11) and

$$\begin{aligned}V_1(z) &= z + \frac{c}{2z},\\ f(z,\lambda) &= 1.\end{aligned} \tag{21}$$

On the other hand, using the expressions in (19), i.e., modifying the metric in the IR, the potential term reads as:

$$V(z) = V_0(z) + \lambda \tilde{V}_1(z) + \frac{c^2\lambda^2}{4} + \frac{m_5^2}{z^2} f(z,\lambda), \tag{22}$$

where

$$\begin{aligned}\tilde{V}_1(z) &= c\left(z + \frac{c}{2z}\right),\\ f(z,\lambda) &= e^{-2\lambda z}.\end{aligned} \tag{23}$$

Considering Eqs. (18)–(23) one sees that the mass term is the main responsible of the difference between the scalar and vector cases when the geometry is perturbed, while its effect turns out to be the same when the background dilaton is modified. The obtained potentials are depicted in Fig. 1.

Eq. (10) with the new potentials (20) and (22) can be solved perturbatively, and for small values of the parameter λ the spectra are modified:

$$m_n^2 = m_{n,(0)}^2 + \lambda m_{n,(1)}^2. \tag{24}$$

For the scalar glueball, modifying the dilaton field according to (18) we obtain for the first three states:

$$\begin{aligned}m_0^2 &= 8 + \lambda\frac{3\sqrt{\pi}}{2},\\ m_1^2 &= 12 + \lambda\frac{27\sqrt{\pi}}{16},\\ m_2^2 &= 16 + \lambda\frac{237\sqrt{\pi}}{128}.\end{aligned} \tag{25}$$

On the other hand, modifying the geometry according to (19) the masses of the first three states the spectrum are given by:

$$\begin{aligned}m_0^2 &= 8 + \lambda\frac{9\sqrt{\pi}}{2},\\ m_1^2 &= 12 + \lambda\frac{81\sqrt{\pi}}{16},\\ m_2^2 &= 16 + \lambda\frac{711\sqrt{\pi}}{128}.\end{aligned} \tag{26}$$

Also for vector glueballs a different spectrum is obtained, depending on the perturbations (18) or (19). Modifying the dilaton field the values of the first three states of the spectrum are:

$$\begin{aligned}m_0^2 &= 12 + \lambda\frac{189\sqrt{\pi}}{128},\\ m_1^2 &= 16 + \lambda\frac{105\sqrt{\pi}}{64},\\ m_2^2 &= 20 + \lambda\frac{14667\sqrt{\pi}}{8192},\end{aligned} \tag{27}$$

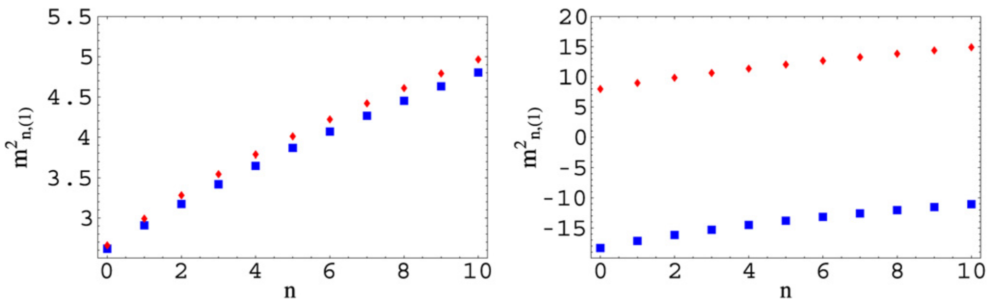

Fig. 2. Mass shifts in Eq. (24) for scalar (red diamonds) and vector glueballs (blue boxes) obtained by modifying the dilaton field (left) or the metric function (right). (For interpretation of the references to colour in this figure legend, the reader is referred to web version of this Letter.)

while modifying the geometry we obtain:

$$m_0^2 = 12 - \lambda \frac{1323\sqrt{\pi}}{128},$$
$$m_1^2 = 16 - \lambda \frac{1239\sqrt{\pi}}{128},$$
$$m_2^2 = 20 - \lambda \frac{74685\sqrt{\pi}}{8192}. \quad (28)$$

Therefore, when the dilaton field is modified, the mass shifts have the same sign in case of scalar and vector glueballs, while the sign is opposite when the geometry is changed. This result is depicted in Fig. 2, where the mass shifts $m^2_{n,(1)}$ defined in (24) are plotted for the first 11 states in case of modified dilaton or geometry.

Different predictions at $\mathcal{O}(\lambda)$ for the vector and scalar glueball mass difference are obtained modifying either the dilaton or the geometry. Modifying the dilaton, we get

$$m^2_{G_1} - m^2_{G_0} = 4 - \frac{3\sqrt{\pi}}{128}\lambda, \quad (29)$$

while, modifying the metric function, we obtain:

$$m^2_{G_1} - m^2_{G_0} = 4 - \frac{1899\sqrt{\pi}}{128}\lambda. \quad (30)$$

Therefore, the mass splitting between vector and scalar glueballs increases if λ is negative, and the maximum effect is produced for the same value of λ when the metric function is perturbed. This can be considered as an indication on the type of constraints the background fields in the bulk must satisfy.

5. Conclusions

We have discussed how the QCD holographic model proposed in [20], with the hard IR wall replaced by a background dilaton field, allows to predict the light glueball spectrum. Scalar and pseudoscalar glueballs turn out to be degenerate if the fields representing the holographic correspondent of the respective QCD operators are both massless zero forms. Vector glueballs turn out to be heavier than the scalar ones, and the dependence of their masses on the radial quantum number is the same as obtained for ρ and higher spin mesons. Combining the calculations of the glueball and ρ masses in the same holographic model, the glueballs turn out to be lighter than predicted in other approaches.

We have investigated how the masses change as a consequence of perturbing the dilaton in the UV or the bulk geometry in the IR, finding that constraints in the bottom–up approach can be found if information on the spectra from other approaches is considered. Such constraints should be taken into account in the attempt to construct the QCD gravitational dual.

Acknowledgements

We are grateful to D.T. Son and L. Yaffe for discussions. We thank F. Canfora for information. This work was supported in part by the EU Contract No. MRTN-CT-2006-035482, “FLAVIAnet”.

References

[1] J.M. Maldacena, Adv. Theor. Math. Phys. 2 (1998) 231;
J.M. Maldacena, Int. J. Theor. Phys. 38 (1999) 1113.
[2] E. Witten, Adv. Theor. Math. Phys. 2 (1998) 253.
[3] S.S. Gubser, I.R. Klebanov, A.M. Polyakov, Phys. Lett. B 428 (1998) 105.
[4] For reviews see: O. Aharony, S.S. Gubser, J.M. Maldacena, H. Ooguri, Y. Oz, Phys. Rep. 323 (2000) 183;
J.L. Petersen, Int. J. Mod. Phys. A 14 (1999) 3597.
[5] E. Witten, Adv. Theor. Math. Phys. 2 (1998) 505.
[6] C. Csaki, H. Ooguri, Y. Oz, J. Terning, JHEP 9901 (1999) 017.
[7] R. de Mello Koch, A. Jevicki, M. Mihailescu, J.P. Nunes, Phys. Rev. D 58 (1998) 105009;
M. Zyskin, Phys. Lett. B 439 (1998) 373;
J.A. Minahan, JHEP 9901 (1999) 020;
C. Csaki, Y. Oz, J. Russo, J. Terning, Phys. Rev. D 59 (1999) 065012;
R.C. Brower, S.D. Mathur, C.I. Tan, Nucl. Phys. B 587 (2000) 249;
R. Apreda, D.E. Crooks, N.J. Evans, M. Petrini, JHEP 0405 (2004) 065.
[8] N. Evans, J.P. Shock, T. Waterson, Phys. Lett. B 622 (2005) 165.
[9] C.J. Morningstar, M.J. Peardon, Phys. Rev. D 60 (1999) 034509;
H.B. Meyer, hep-lat/0508002.
[10] L. von Smekal, R. Alkofer, A. Hauck, Phys. Rev. Lett. 79 (1997) 3591;
D.M. Howe, C.J. Maxwell, Phys. Lett. B 541 (2002) 129;
D.M. Howe, C.J. Maxwell, Phys. Rev. D 70 (2004) 014002;
D. Zwanziger, Phys. Rev. D 69 (2004) 016002.
[11] A.C. Mattingly, P.M. Stevenson, Phys. Rev. D 49 (1994) 437;
S.J. Brodsky, S. Menke, C. Merino, J. Rathsman, Phys. Rev. D 67 (2003) 055008.
[12] J. Polchinski, M.J. Strassler, Phys. Rev. Lett. 88 (2002) 031601.
[13] G.F. de Teramond, S.J. Brodsky, Phys. Rev. Lett. 94 (2005) 201601.
[14] J. Erlich, E. Katz, D.T. Son, M.A. Stephanov, Phys. Rev. Lett. 95 (2005) 261602.
[15] E. Katz, A. Lewandowski, M.D. Schwartz, Phys. Rev. D 74 (2006) 086004.
[16] L. Da Rold, A. Pomarol, Nucl. Phys. B 721 (2005) 79;
L. Da Rold, A. Pomarol, JHEP 0601 (2006) 157;

S. Hong, S. Yoon, M.J. Strassler, JHEP 0604 (2006) 003;
K. Ghoroku, N. Maru, M. Tachibana, M. Yahiro, Phys. Lett. B 633 (2006) 602;
J. Hirn, N. Rius, V. Sanz, Phys. Rev. D 73 (2006) 085005.

[17] H. Boschi-Filho, N.R.F. Braga, JHEP 0305 (2003) 009;
H. Boschi-Filho, N.R.F. Braga, H.L. Carrion, Phys. Rev. D 73 (2006) 047901.

[18] O. Andreev, V.I. Zakharov, Phys. Rev. D 74 (2006) 025023.

[19] A.V. Radyushkin, Phys. Lett. B 642 (2006) 459;
H.M. Choi, C.R. Ji, Phys. Rev. D 74 (2006) 093010;
S.J. Brodsky, G.F. de Teramond, Phys. Rev. Lett. 96 (2006) 201601;
H.R. Grigoryan, A.V. Radyushkin, hep-ph/0703069.

[20] A. Karch, E. Katz, D.T. Son, M.A. Stephanov, Phys. Rev. D 74 (2006) 015005.

[21] E. Schreiber, hep-th/0403226;
M.A. Shifman, hep-ph/0507246.

[22] C.A. Dominguez, N. Paver, Z. Phys. C 31 (1986) 591;
S. Narison, Nucl. Phys. B 509 (1998) 312.

[23] V.A. Novikov, M.A. Shifman, A.I. Vainshtein, V.I. Zakharov, Nucl. Phys. B 191 (1981) 301.