

SCALAR GLUEBALL IN A TOP-DOWN HOLOGRAPHIC APPROACH TO QCD*

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Identification of glueballs — bound states of gauge bosons in Quantum Chromodynamics (QCD) — is a very important open question in dynamics of the strong interaction. The search for the glueball ground state, carrying scalar quantum numbers, poses a particular challenge due to the existence of *(i)* several candidates for its realisation in the physical spectrum and *(ii)* inevitable mixing of the pure glueball state with those comprised of quarks. In this article, I discuss implications of an approach in holographic QCD where, among others, the mass and the two-pion decay of the pure scalar glueball can be studied.

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1. Introduction

The non-Abelian nature of Quantum Chromodynamics (QCD) — the theory of strong interaction — gives rise to the expectation that its gauge bosons, the gluons, form composite objects denoted as glueballs [1]. These states would have access to various quantum numbers J^{PC} , where J denotes the total spin, P the parity and C the charge conjugation; the corresponding spectrum in the QCD Yang–Mills sector has been determined in numerical simulations [2, 3] but the identification of glueballs in experimental data has proven to be a challenge, particularly in the $J^{PC} = 0^{++}$ (scalar) channel.

There are at least two reasons to study glueball states. Firstly, their masses are generated solely by the strong interaction; no influence of the Higgs mechanism — providing, *e.g.*, quarks with a current mass — is present

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in the glueball mass generation. Secondly, due to their structure, glueballs must possess integer spin assigning them to mesons; our understanding of the physical meson spectrum would not be complete without glueballs.

Glueball ground state is scalar [4]; listings of the Particle Data Group (PDG) [5] contain five isospin-zero scalar states in the energy region below 1 GeV: $f_0(500)$ or σ , $f_0(980)$, $f_0(1370)$, $f_0(1500)$ and $f_0(1710)$. While all of them possess the quantum numbers of the ground state, there are strong reasons to focus on resonances above 1 GeV as candidates for the scalar glueball: *(i)* simulations in lattice QCD determine the ground-state mass at between approximately 1.65 GeV and 1.8 GeV [2, 3]; *(ii)* various effective approaches to low-energy QCD arrive at an analogous result for the mass, while describing the overall dynamics correctly [6]. Complication is that f_0 resonances will most certainly have contributions not only from the (pure) glueball states but also from those containing quarks ($\bar{q}q$ [7], $\bar{q}\bar{q}qq$ [8] and others). This leads to various issues in both theory and experiment [9, 10] and represents the main reason why the scalar glueball has still not been clearly identified.

In this article, the question of the scalar glueball is tackled by a holographic approach to non-perturbative QCD. Such approaches are based on the idea of a duality between strongly coupled quantum gauge field theories and weakly coupled supergravity/superstring theories in one dimension higher, pursuing Maldacena's groundbreaking conjecture of a complete equivalence between the supergravity limit of type-IIB string theory on an $\text{AdS}_5 \times S^5$ space and the large- N limit of an $\mathcal{N} = 4$ supersymmetric and conformal $U(N)$ gauge theory on its boundary (AdS/CFT correspondence [11]). In Ref. [12], Witten proposed an analogous duality in type-IIA string theory, with supersymmetry and conformality not present in line with their absence in QCD. The supersymmetry is broken by compactification on a circle (S^1); for a vanishing circle radius, the 5-dimensional pure Yang–Mills theory is reduced to a 4-dimensional one. However, the supergravity approximation requires finite circle radius (whose inverse is defined as the Kaluza–Klein mass M_{KK}), and also a large coupling. Thus constructed holographic approaches are referred to as *top–down* models [13, 14]; there are also more phenomenological *bottom–up* constructions — see Ref. [15] and references therein. This article describes the implications of the top–down Witten–Sakai–Sugimoto model [14] for the glueball spectroscopy; more details can be found in Ref. [16].

2. The model and its implications

While Witten's model contained only gauge fields, the novel feature of Witten–Sakai–Sugimoto model is the inclusion of chiral quarks intro-

duced by N_f (number of flavours) probe D8 and anti-D8 branes (inducing $U(N_f) \times U(N_f)$ chiral symmetry) that extend along all dimensions of the 10-dimensional space with the exception of a (Kaluza–Klein) circle. The branes are usually antipodally separated with regard to this S^1 . The space geometry, however, is such that the branes and antibranes merge at a certain point — interpreted as realisation of the chiral-symmetry breaking.

Up to a Chern–Simons term, the corresponding action for D8-branes reads

$$S_{\text{D8}} = -T_{\text{D8}} \text{Tr} \int d^9 x e^{-\Phi} \sqrt{-\det(\tilde{g}_{MN} + 2\pi\alpha' F_{MN})}, \quad (1)$$

where $T_{\text{D8}} = (2\pi)^{-8} l_s^{-9}$ (and $l_s^2 = \alpha'$, with α' the string coupling), g_{MN} is the metric of the D-brane world volume, Φ is the dilaton field and F_{MN} a field strength tensor whose components are, upon dimensional reduction, identified as meson fields of interest. No backreaction of the Witten-model background to D8-branes is considered; consequently, N_f is fixed and, as can be argued [14], significantly smaller than the number of colours (large- N_c limit).

The above action is expanded up to the second order in fields as

$$S_{\text{D8}}^{(2)} = -\kappa \text{Tr} \int d^4 x \int_{-\infty}^{\infty} dZ \left[\frac{1}{2} K^{-\frac{1}{3}} \eta^{\mu\rho} \eta^{\nu\sigma} F_{\mu\nu} F_{\rho\sigma} + M_{\text{KK}}^2 \eta^{\mu\nu} F_{\mu Z} F_{\nu Z} \right], \quad (2)$$

where $\kappa = \lambda N_c / (216\pi^3)$ [16], $\lambda = g_{\text{YM}}^2 N_c$ is the 't Hooft coupling (and g_{YM} the 4-dimensional coupling), Z is the holographic radial coordinate (and $K = 1 + Z^2$) and $\eta^{\mu\nu}$ is the flat metric $\text{diag}(-, +, +, +)$. The Kaluza–Klein mass M_{KK} sets the model scale; beside the scale, the model contains only one unknown quantity: the coupling λ . They are usually determined such that the mass of the rho meson and the pion decay constant correspond to their physical values. This yields $M_{\text{KK}} = 949$ MeV and $\lambda = 16.63$; alternative methods for their determination do not alter model conclusions [16].

Once M_{KK} and λ are known, the model describes various experimental quantities surprisingly well [14]. Glueballs are of interest in this article; building on the work of Ref. [17], the following masses in the scalar channel are obtained

$$\begin{aligned} M_{G_{\text{E}}} &= 855 \text{ MeV}, & M_{G_{\text{E}}^*} &= 2168 \text{ MeV}, \\ M_{D_{\text{E}}} &= 1487 \text{ MeV}, & M_{D_{\text{E}}^*} &= 2358 \text{ MeV}. \end{aligned} \quad (3)$$

Two types of glueball states are present: one, denoted with G_{D} , is predominantly dilaton, while the other, denoted with G_{E} , involves a graviton polarisation in a fourth spatial dimension, unlike the dilaton, and has therefore

been termed *exotic* [18]. Excited states are denoted with an asterisk. We observe that the exotic ground state is approximately 50% lighter than the expectation from lattice QCD; the dilaton mode mass is only approximately 15% smaller than the lattice result. Numerical simulations also indicate the mass of the first excited scalar state to be $\simeq 2600$ MeV with errors amounting to $\simeq 300$ MeV [2] (although mass corrections in the unquenched case may be substantial [3]); the excited dilaton state is, within errors, consistent with the lattice result, while the excited exotic mode is approximately 15% too light. Hence, already from the mass results the indication is that the exotic mode could be discarded, while the dilaton mode appears compatible with simulations of the Yang–Mills sector of QCD.

This is corroborated by the decay ratios of the modes. The corresponding Lagrangians are obtained by inserting 10-dimensional metric fluctuations (whose explicit forms together with further details are presented in Ref. [16]) into the action for D8 branes and integrating over the bulk coordinates (see also Ref. [19]); ratios of decay widths Γ and the respective masses M read

$$\begin{aligned} \Gamma_{G_E \rightarrow \pi\pi}/M_{G_E} &= 0.092, & \Gamma_{G_E^* \rightarrow \pi\pi}/M_{G_E^*} &= 0.149, \\ \Gamma_{G_D \rightarrow \pi\pi}/M_{G_D} &= 0.009, & \Gamma_{G_D^* \rightarrow \pi\pi}/M_{G_D^*} &= 0.011. \end{aligned} \quad (4)$$

The two main candidates for the scalar glueball are the $f_0(1500)$ and $f_0(1710)$ resonances (see Ref. [10] and references therein). Experimental data imply $\Gamma_{f_0(1500) \rightarrow \pi\pi}/M_{f_0(1500)} = 0.025 \pm 0.003$ [5] and $\Gamma_{f_0(1710) \rightarrow \pi\pi}/M_{f_0(1710)} \simeq 0.009\text{--}0.017$ [20]. The exotic mode is thus too broad, again indicating that its interpretation as a physical state is uncertain; contrarily, the ratio Γ/M for the dilaton mode is within the experimental interval for $f_0(1710)$.

Similar is true for the excited states: $\Gamma_{G_E^* \rightarrow \pi\pi}$ is larger than $\Gamma_{G_D^* \rightarrow \pi\pi}$ and nicely comparable to widths of f_0 states near and above 2 GeV; however, the full decay width of G_E^* having contributions from $2K$, 2η , 4π and other channels is unphysically large (~ 1 GeV, see Ref. [16]). Contrarily, $\Gamma_{G_D^* \rightarrow \pi\pi}$ is small but the full decay width of G_D^* is of the order of 460 MeV [16] and thus significantly closer to the data [5] whose current uncertainties unfortunately do not allow for a clear identification of an excited scalar glueball.

3. Summary and outlook

In this article, a top–down holographic approach to low-energy QCD — Witten–Sakai–Sugimoto model — has been presented and its implications for phenomenology of scalar glueballs have been discussed. The model offers two sets of glueball states, a dilaton and an *exotic* mode that, unlike the dilaton, involves a graviton polarisation in a fourth spatial dimension. The exotic ground state has a mass approximately 50% smaller than the value expected in lattice QCD; its 2π decay width is substantially larger than that

of the two prime candidates for the scalar glueball, the resonances $f_0(1500)$ and $f_0(1710)$. Contrarily, the mass of the dilaton mode ($= 1487$ MeV) is quite close to masses of both mentioned resonances; its 2π decay width is within the experimental range for $f_0(1710)$, which therefore appears to be the preferred candidate for the glueball ground state. In the excited channel, the exotic state is unphysically broad while the dilaton width ~ 460 MeV is close to the (still ambiguous) data on f_0 states near/above 2 GeV.

Nonetheless, further pursuit of glueball dynamics in holography is called for, particularly in light of expectations from the planned PANDA experiments at FAIR [21].

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REFERENCES

- [1] A. Salomone, J. Schechter, T. Tudron, *Phys. Rev.* **D23**, 1143 (1981); C. Rosenzweig, A. Salomone, J. Schechter, *Nucl. Phys.* **B206**, 12 (1982); [*Erratum ibid.*, **B207**, 546 (1982)]; A.A. Migdal, M.A. Shifman, *Phys. Lett.* **B114**, 445 (1982); H. Gomm, J. Schechter, *Phys. Lett.* **B158**, 449 (1985); R. Gomm, P. Jain, R. Johnson, J. Schechter, *Phys. Rev.* **D33**, 801 (1986).
- [2] C.J. Morningstar, M.J. Peardon, *Phys. Rev.* **D60**, 034509 (1999) [arXiv:hep-lat/9901004]; Y. Chen *et al.*, *Phys. Rev.* **D73**, 014516 (2006) [arXiv:hep-lat/0510074]; M. Loan, X.Q. Luo, Z.H. Luo, *Int. J. Mod. Phys.* **A21**, 2905 (2006) [arXiv:hep-lat/0503038].
- [3] E. Gregory *et al.*, *J. High Energy Phys.* **1210**, 170 (2012) [arXiv:1208.1858 [hep-lat]].
- [4] G.B. West, *Phys. Rev. Lett.* **77**, 2622 (1996) [arXiv:hep-ph/9603316].
- [5] K.A. Olive *et al.* [Particle Data Group], *Chin. Phys.* **C38**, 090001 (2014).
- [6] F.E. Close, A. Kirk, *Eur. Phys. J.* **C21**, 531 (2001) [arXiv:hep-ph/0103173]; F. Giacosa, T. Gutsche, A. Faessler, *Phys. Rev.* **C71**, 025202 (2005) [arXiv:hep-ph/0408085]; F. Giacosa, T. Gutsche, V.E. Lyubovitskij, A. Faessler, *Phys. Lett.* **B622**, 277 (2005) [arXiv:hep-ph/0504033]; *Phys. Rev.* **D72**, 094006 (2005) [arXiv:hep-ph/0509247]; H.Y. Cheng, C.K. Chua, K.F. Liu, *Phys. Rev.* **D74**, 094005 (2006) [arXiv:hep-ph/0607206]; V. Mathieu, N. Kochelev, V. Vento, *Int. J. Mod. Phys.* **E18**, 1 (2009) [arXiv:0810.4453 [hep-ph]]; S. Janowski, D. Parganlija, F. Giacosa, D.H. Rischke, *Phys. Rev.* **D84**, 054007 (2011) [arXiv:1103.3238 [hep-ph]]; S. Janowski, F. Giacosa, D.H. Rischke, *Phys. Rev.* **D90**, 114005 (2014) [arXiv:1408.4921 [hep-ph]].

- [7] D. Parganlija, F. Giacosa D.H. Rischke, *Phys. Rev.* **D82**, 054024 (2010) [arXiv:1003.4934 [hep-ph]]; D. Parganlija *et al.*, *Phys. Rev.* **D87**, 014011 (2013) [arXiv:1208.0585 [hep-ph]].
- [8] R.L. Jaffe, *Phys. Rev.* **D15**, 267 (1977); A.H. Fariborz, R. Jora, J. Schechter, *Phys. Rev.* **D72**, 034001 (2005) [arXiv:hep-ph/0506170]; F. Giacosa, *Phys. Rev.* **D75**, 054007 (2007) [arXiv:hep-ph/0611388]; A. Heinz, S. Struber, F. Giacosa, D.H. Rischke, *Phys. Rev.* **D79**, 037502 (2009) [arXiv:0805.1134 [hep-ph]]; *Acta Phys. Pol. B Proc. Suppl.* **3**, 925 (2010) [arXiv:1006.5393 [hep-ph]].
- [9] D.V. Bugg, *Phys. Rep.* **397**, 257 (2004) [arXiv:hep-ex/0412045]; E. Klempt, A. Zaitsev, *Phys. Rep.* **454**, 1 (2007) [arXiv:0708.4016 [hep-ph]].
- [10] D. Parganlija, *J. Phys. Conf. Ser.* **426**, 012019 (2013) [arXiv:1211.4804 [hep-ph]]; **503**, 012010 (2014) [arXiv:1312.2830 [hep-ph]].
- [11] J.M. Maldacena, *Int. J. Theor. Phys.* **38**, 1113 (1999) [*Adv. Theor. Math. Phys.* **2**, 231 (1998)] [arXiv:hep-th/9711200].
- [12] E. Witten, *Adv. Theor. Math. Phys.* **2**, 505 (1998) [arXiv:hep-th/9803131].
- [13] I.R. Klebanov, M.J. Strassler, *J. High Energy Phys.* **0008**, 052 (2000) [arXiv:hep-th/0007191]; J. Babington *et al.*, *Phys. Rev.* **D69**, 066007 (2004) [arXiv:hep-th/0306018]; M. Kruczenski, D. Mateos, R.C. Myers, D.J. Winters, *J. High Energy Phys.* **0405**, 041 (2004) [arXiv:hep-th/0311270].
- [14] T. Sakai, S. Sugimoto, *Prog. Theor. Phys.* **113**, 843 (2005) [arXiv:hep-th/0412141]; **114**, 1083 (2005) [arXiv:hep-th/0507073]; T. Imoto, T. Sakai, S. Sugimoto, *Prog. Theor. Phys.* **124**, 263 (2010) [arXiv:1005.0655 [hep-th]]; F. Brünner, D. Parganlija, A. Rebhan, *Acta Phys. Pol. B Proc. Suppl.* **7**, 533 (2014) [arXiv:1407.6914 [hep-ph]]; A. Rebhan, arXiv:1410.8858 [hep-th].
- [15] S.S. Afonin, *Int. J. Mod. Phys.* **A25**, 5683 (2010) [arXiv:1001.3105 [hep-ph]].
- [16] F. Brünner, D. Parganlija, A. Rebhan, arXiv:1501.07906 [hep-ph].
- [17] R.C. Brower, S.D. Mathur, C.-I. Tan, *Nucl. Phys.* **B587**, 249 (2000) [arXiv:hep-th/0003115].
- [18] N.R. Constable, R.C. Myers, *J. High Energy Phys.* **9910**, 037 (1999) [arXiv:hep-th/9908175].
- [19] K. Hashimoto, C.-I. Tan, S. Terashima, *Phys. Rev.* **D77**, 086001 (2008) [arXiv:0709.2208 [hep-th]].
- [20] D. Parganlija, arXiv:1208.0204 [hep-ph].
- [21] M.F.M. Lutz *et al.* [PANDA Collaboration], arXiv:0903.3905 [hep-ex].