Holographic QCD predictions for production and decay of pseudoscalar glueballs

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The top-down holographic Witten-Sakai-Sugimoto model for low-energy QCD, augmented by finite quark masses, has recently been found to be able to reproduce the decay pattern of the scalar glueball candidate $f_0(1710)$ on a quantitative level. In this Letter we show that this model predicts a narrow pseudoscalar glueball heavier than the scalar glueball and with a very restricted decay pattern involving η or η' mesons. Production should be either in pairs or in association with $\eta(')$ mesons. We discuss the prospect of discovery in high-energy hadron collider experiments through central exclusive production by comparing with η' pair production.

I. INTRODUCTION

Quantum chromodynamics, the established theory of the strong interactions, predicts [1] the existence of flavor singlet mesons beyond those required by the quark model, because in the absence of quarks gluons by themselves can form bound states. However, the status of such "glueball" states in the observed meson spectrum is still unclear and controversal [2–5].

In 1980, an isoscalar pseudoscalar with a mass of around 1.44 GeV which is copiously produced in the gluon-rich radiative decays of J/ψ was proposed as the first glueball candidate [6]. Once named $\iota(1440)$ [7], this is now listed by the Particle Data Group [8] as the two states $\eta(1405)$ and $\eta(1475)$. Together with $\eta(1295)$, this indeed would give rise to a supernumerary state beyond the first radial excitations of the η and η' mesons, with $\eta(1405)$ singled out as glueball candidate [9].

The situation thus appears to be analogous to the case of the scalar glueball, which is widely considered to be responsible for a supernumerary state in the set of isoscalar scalar resonances $f_0(1370)$, $f_0(1500)$, and $f_0(1710)$, where only two are expected from the quark model (corresponding to $\bar{u}u+\bar{d}d$ and $\bar{s}s$). Here the discussion is divided on the question which of the two heavier resonances has the larger glueball contribution [10–14].

However, only the case of the scalar glueball candidates is supported by existing lattice QCD calculations [15, 16] which consistently find that the lowest-lying glueball state has a mass of around 1.7 GeV and quantum numbers $J^{PC} = 0^{++}$. The lowest-lying pseudoscalar glueball state is instead found to have a mass of around 2.6 GeV, somewhat higher than the 2⁺⁺ tensor glueball with a mass of around 2.4 GeV. Most lattice results have been obtained in the quenched approximation¹, i.e. without dynamical quarks, but recent unquenched lattice calculations [18–20] have found no evidence for significant unquenching effects, which however should be expected if the pseudoscalar glueball were to mix strongly with radially excited $\eta(')$ mesons. Moreover, Ref. [20] recently reported that correlation functions of pseudoscalar gluonic operators built from Wilson loops did not show any trace of the flavor singlet pseudoscalar meson states which can be found in the topological charge density correlator. In fact, on the experimental side it is still a controversial issue whether as many as three states $\eta(1295)$, $\eta(1405)$, and $\eta(1475)$ (and thus indication of the involvement of the pseudoscalar glueball in this mass region) really exist.²

We therefore assume that (contrary to the models used in Ref. [22–24]) the pseudoscalar glueball does not make its appearance in the known η mesons in the 1400 MeV region, but that it still has to be discovered and that it should be searched for in the mass range 2–3 GeV. Unfortunately, lattice QCD does not (yet) give information on the production and decay patterns of a pseudoscalar glueball, whereas phenomenological models are weakly constrained with regard to the particular form of pseudoscalar glueball interactions.³

In this work we show that rather specific predictions can be obtained from the Witten-Sakai-Sugimoto (WSS) model for low-energy QCD, which is a top-down stringtheoretic construction in the large color number (N_c) limit with only one free dimensionless parameter. Extrapolated to $N_c = 3$, it reproduces several experimental results in hadron physics to within 10-30% [27, 28]. In Ref. [28] we have applied this model to calculate decay rates of scalar and tensor glueballs in the chiral limit, and in [29, 30] with quark masses included. In the latter case we found a strong "nonchiral enhancement" of the decay of a predominantly dilatonic glueball into kaons

¹ It has been argued that the pseudoscalar sector may be particularly sensitive to unquenching in Ref. [17], but the estimated effects on the mass were of the order of 15%, whereas almost 50% would be needed to bring the lattice result down to the mass of $\eta(1405)$.

² E.g., the existence of $\eta(1295)$ is questioned in [3], while Ref. [21] came to the conclusion that there is "no evidence for two separate $\eta(1405)$ and $\eta(1475)$ from the present data" and only one $\eta(1440)$ is actually required.

³ In Ref. [25] a unique form of the interaction Lagrangian for extended linear sigma models has been posited, where only the coupling strength is left undetermined, but in a subsequent extension [26] more possibilities were introduced.

and η mesons which quantitatively agrees remarkably well with the data for the glueball candidate $f_0(1710)$ as far as presently known (provided the not-yet-measured decay rate into $\eta\eta'$ pairs is sufficiently small [30]). This suggests that $f_0(1710)$ could be a nearly pure glueball, in agreement with recent phenomenological models that favor $f_0(1710)$ as the scalar glueball [11, 12] with comparatively small admixture of light quarkonia.

While in Ref. [28] our WSS model prediction for the width of the tensor glueball of mass $\gtrsim 2$ GeV was very large, perhaps too large to be clearly observable, here we arrive at the prediction of a narrow pseudoscalar glueball state with a very restricted decay pattern, which will be a conspicuous feature as long as mixing with quarkonia is small. The specific interactions also suggest that the pseudoscalar glueball may be difficult to produce in radiative charmonium decay, but could be a very interesting object for glueball searches in central exclusive production (CEP) experiments at sufficiently high energies.

II. EFFECTIVE LAGRANGIAN FOR PSEUDOSCALAR GLUEBALL INTERACTIONS

The WSS model [31–33] is a gauge/gravity-dual model for nonsupersymmetric low-energy QCD based on D4 branes in type-IIA supergravity compactified on a circle and subjected to a consistent truncation of Kaluza-Klein states, with $N_f \ll N_c$ chiral quarks added through probe D8 branes. It possesses an interesting spectrum of glueball states with $J^{PC} = 0^{++}, 2^{++}, 0^{-+}, 1^{+-}, 1^{--}$ [34] whose mass scale is set by the Kaluza-Klein mass $M_{\rm KK}$. The resulting effective theory involves Goldstone pseudoscalars for nonabelian chiral symmetry breaking and a tower of vector and axial vector mesons.

Fixing $M_{\rm KK}$ through the experimental value of the ρ meson mass and varying the 't Hooft coupling $\lambda = 16.63...12.55$ such that either the pion decay constant or the string tension in large- N_c lattice simulations [35] is matched leads to quantitative predictions which are in the right ballpark when extrapolated to $N_c = 3$ QCD, including a value for the gluon condensate

$$C^{4} \equiv \left\langle \frac{\alpha_{s}}{\pi} G^{a}_{\mu\nu} G^{a\mu\nu} \right\rangle = \frac{1}{2\pi^{2}} \left\langle \operatorname{Tr} F^{2} \right\rangle = \frac{4N_{c}}{3^{7}\pi^{4}} \lambda^{2} M^{4}_{\mathrm{KK}}$$
(1)

that is close to that obtained by SVZ sum rules [27]. Moreover, it reproduces remarkably well the observed hadronic decay rates of the ρ and the ω mesons, which motivates the use of the WSS model also as a model for glueball decay [28, 36]. In Ref. [28] we argued, however, that the lightest scalar glueball mode considered in Ref. [36] which comes from an "exotic polarization" of the dual graviton along the compactified direction (denoted by G_E in the following) should be discarded and that instead the predominantly dilatonic mode (G_D) be identified with the glueball ground state.

The WSS model correctly incorporates the nonabelian chiral anomaly of QCD and the resulting Wess-ZuminoWitten term as well as the $U(1)_A$ anomaly and the Witten-Veneziano mechanism for giving mass to the flavor singlet pseudoscalar η_0 with [32, 37, 38]

$$m_0^2 = \frac{N_f}{27\pi^2 N_c} \lambda^2 M_{\rm KK}^2,$$
 (2)

leading to $m_0 = 730 - 967$ MeV for λ between 12.55 and 16.63. Introducing explicit quark mass terms in the effective Lagrangian such that physical pion and kaon masses are matched leads to η and η' masses that agree with real QCD to within $\leq 10\%$ [29, 30]. As mentioned above, the flavor-asymmetric decay pattern observed for the scalar glueball candidate $f_0(1710)$ can be reproduced quantitatively with G_D , if the (as yet undetermined) parameter for scalar glueball couplings to explicit quark mass terms is chosen such that the rate of decay into mixed $\eta\eta'$ pairs remains small.

The interaction Lagrangian of the pseudoscalar glueballs is the same for both, the chiral and the massive version of the WSS model. The pseudoscalar glueball modes are provided by a Ramond-Ramond (RR) 1-form field C_1 which plays the central role in producing the Witten-Veneziano mass m_0 . Following the notation of Ref. [32], the action for C_1 is given by

$$S_{C_1} = -\frac{1}{4\pi (2\pi l_s)^6} \int d^{10}x \sqrt{-g} |\tilde{F}_2|^2.$$
(3)

As reviewed in the Appendix, anomaly cancellation requires that \tilde{F}_2 is a gauge invariant combination of $F_2 = dC_1$ and the field

$$\eta_0(x) = \frac{f_\pi}{\sqrt{2N_f}} \int dz \,\mathrm{Tr}A_z(z, x) \tag{4}$$

with z parametrizing the radial extent of the joined D8 and anti-D8 branes on which the flavor gauge field A lives.

Inserting a mode expansion of the RR 1-form field C_1 with 4-dimensional pseudoscalar glueball fields $\tilde{G}^{(n)}(x)$, $n = 1, \ldots$, together with scalar and tensor glueball fields entering through the metric in S_{C_1} leads to the effective 4-dimensional Lagrangian

$$\mathcal{L}_{C_{1}}^{\text{eff}} = -\frac{1}{2} \partial_{\mu} \tilde{G} \, \partial^{\mu} \tilde{G} - \frac{1}{2} m_{P}^{2} \tilde{G}^{2} - \frac{1}{2} m_{0}^{2} \eta_{0}^{2} \\ + \mathcal{L}_{\eta_{0}^{2}G} + \mathcal{L}_{\tilde{G}\eta_{0}G} + \mathcal{L}_{\tilde{G}^{2}G} + O(G_{D,E,T}^{2})$$
(5)

(suppressing the summation over the mode number index (n)). Here $O(G_{D,E,T}^2)$ denotes higher-order interactions involving terms quadratic in \tilde{G} , η_0 and quadratic or higher in the glueball fields arising from metric fluctuations (the tensor glueball field $T^{\mu\nu}$ appears at most linearly, but also has interactions involving arbitrarily high powers of the scalar glueball field).

The mass of the lowest pseudoscalar glueball mode (n = 1) is [34] $M_P \approx 1.885 M_{\rm KK}$, which like in lattice QCD results is above the mass of the scalar and tensor glueballs with $M_D = M_T \approx 1.567 M_{\rm KK}$. With

coeff.	value
\bar{d}_0	$17.915 \lambda^{-1/2} N_c^{-1} M_{\rm KK}^{-1}$
$ ilde{d}_0$	$2.5833 \lambda^{1/2} N_f^{1/2} N_c^{-3/2} M_{\rm KK}$
\tilde{d}_1	$42.484 \lambda^{-1/2} N_c^{-1} M_{\rm KK}^{-1}$
\tilde{d}_2	$27.106 \lambda^{-1/2} N_c^{-1} M_{\rm KK}^{-1}$
\breve{c}_0	$15.829 \lambda^{-1/2} N_c^{-1} M_{\rm KK}^{-1}$
\bar{c}_0	$26.837 \lambda^{-1/2} N_c^{-1} M_{\rm KK}^{-1}$
\tilde{c}_0	$-4.8795 \lambda^{1/2} N_f^{1/2} N_c^{-3/2} M_{\rm KK}$
\tilde{c}_0'	$1.6306 \lambda^{1/2} N_f^{1/2} N_c^{-3/2} M_{\rm KK}$
$\tilde{c}_0^{\prime\prime}$	$2.0502 \lambda^{1/2} N_f^{1/2} N_c^{-3/2} M_{\rm KK}$

TABLE I. Coupling constants in the glueball interaction Lagrangians (6), (7), and (8).

 $M_{\rm KK} = 949$ MeV from having matched the mass of the ρ meson, $M_D \approx 1487$ MeV and $M_P \approx 1789$ MeV, but in the eventual applications we shall leave M_P a free parameter and either keep M_D at 1.5 GeV which approximately matches the mass of $f_0(1500)$ or artificially raise its mass to the mass of $f_0(1710)$.

Note that Eq. (5) contains a mass term for the flavor singlet η_0 [32], but no mixing of the pseudoscalar glueball modes $\tilde{G}^{(n)}$ with $\eta_0.^4$ As shown in the Appendix, terms proportional to $\eta_0 \tilde{G}^{(n)}$ vanish in the unperturbed background geometry, but arise in the presence of metric fluctuations dual to scalar glueballs. In the WSS model, such terms are the only ones which can mediate a decay of pseudoscalar glueballs. Explicitly they read (keeping the exotic glueball mode G_E for completeness)

$$\mathcal{L}_{\tilde{G}\eta_{0}G} = \tilde{d}_{0}\,\tilde{G}\,\eta_{0}\,G_{D} + \tilde{c}_{0}\,\tilde{G}\,\eta_{0}\,G_{E} + \frac{\tilde{c}_{0}'}{M_{E}^{2}}\,\partial_{\mu}\tilde{G}\,\eta_{0}\,\partial^{\mu}G_{E} + \tilde{c}_{0}''\,\tilde{G}\,\eta_{0}\,\frac{\Box - M_{E}^{2}}{M_{E}^{2}}G_{E} \tag{6}$$

with the numerical results for the coupling constants for the lowest pseudoscalar glueball mode listed in Table I (their integral representations will be given elsewhere).

The part of the action which leads to the Witten-Veneziano mass term also gives rise to interactions with scalar glueballs which were obtained (on-shell) in [29]. To linear order in glueball fields the corresponding interaction Lagrangian reads (also including an extra off-shell contribution for the exotic mode G_E)

$$\mathcal{L}_{\eta_0^2 G} = \frac{1}{2} m_0^2 \eta_0^2 \left(3d_0 G_D - 5\breve{c}_0 G_E \right) \\ + \frac{1}{2} \bar{c}_0 m_0^2 \eta_0^2 \frac{\Box - M_E^2}{M_E^2} G_E.$$
(7)

There are also interaction terms of the form $(\partial \eta_0)^2 G_{D,E,T}$ coming from the DBI action of the D8 branes, which can be found in Ref. [28], as well as natural-parity violating terms $\eta_0 G_T^2$ from Chern-Simons action of the D8 branes, which have been obtained in Ref. [40].

Interaction terms involving pairs of pseudoscalar glueballs and a scalar or tensor glueball are given by

$$\mathcal{L}_{\tilde{G}^{2}G} = \tilde{d}_{1} \left[\frac{1}{2} \partial_{\mu} \tilde{G} \partial^{\mu} \tilde{G} - \frac{1}{8} \partial_{\mu} \tilde{G} \partial_{\nu} \tilde{G} \frac{\partial^{\mu} \partial^{\nu}}{\Box} \right] G_{D} + \frac{1}{2} \tilde{d}_{2} m_{P}^{2} \tilde{G}^{2} G_{D} + \frac{\sqrt{6}}{8} \tilde{d}_{1} \partial_{\mu} \tilde{G} \partial_{\nu} \tilde{G} T^{\mu\nu} + \mathcal{L}_{\tilde{G}^{2} G_{E}}.$$
(8)

(The more unwieldy expression $\mathcal{L}_{\tilde{G}^2G_E}$ will be given elsewhere.)

III. DECAY PATTERN OF THE PSEUDOSCALAR GLUEBALL

The only interaction terms arising within the WSS model that are relevant for the decay of pseudoscalar glueballs are contained in (6). They differ strongly from the leading interaction terms that have been assumed previously in phenomenological models.

Rosenzweig et al. [41, 42] have assumed that the chiral anomaly is not saturated by η_0 alone, but involves a further physical pseudoscalar field (\tilde{G}_2) ,⁵ which couples to the imaginary part of $\log \det \Sigma$, where Σ is the matrix of $q\bar{q}$ states (which is unitary in the nonlinear sigma model, involving only the pseudoscalars, but unrestricted in linear sigma models [43] so that it also accommodates scalar mesons). While a natural possibility [41] would be to identify \tilde{G}_2 with the radial excitation of η_0 , it was proposed to identify \tilde{G}_2 with the pseudoscalar glueball instead. Originally used in the context of the glueball candidate $\iota(1440)$, this approach was also adopted in the extended linear sigma model of Ref. [25] for pseudoscalar glueballs with a mass suggested by lattice QCD. The dominant decay mode of a pseudoscalar glueball in this approach turns out to be $K\bar{K}\pi$ (branching ratio $\mathcal{B}\approx 1/2$) followed by $\eta \pi \pi$ ($\mathcal{B} \approx 1/6$) and $\eta' \pi \pi$ ($\mathcal{B} \approx 1/10$).

Using large- N_c chiral Lagrangians, Gounaris et al. [44] argued that there should be no coupling of the pseudoscalar glueball to Im log det Σ . Instead, a coupling to Im tr $\mathcal{M}_q \Sigma$ was considered so that the pseudoscalar glueball is stable in the limit of massless quarks (\mathcal{M}_q being the quark mass matrix). This again gives a dominant decay mode $K\bar{K}\pi$, but with $\eta\pi\pi$ being more strongly suppressed (parametrically by a factor m_{π}^2/m_K^2).

⁴ This feature is due the fact the WSS model corresponds to QCD in the 't Hooft limit $N_c \gg 1$ but $N_f \sim 1$. In the bottom-up holographic model of Ref. [39], where the Veneziano limit $N_c \gg 1$ and $N_f/N_c \sim 1$ is taken, mixing of pseudoscalar glueballs and η_0 appears at leading order, but in a way that depends strongly on the choice of potentials.

⁵ In Ref. [41] \tilde{G}_1 is an auxiliary field with wrong-sign mass term that can be replaced by Im log det Σ , which is essentially η_0 , through its algebraic equations of motion.

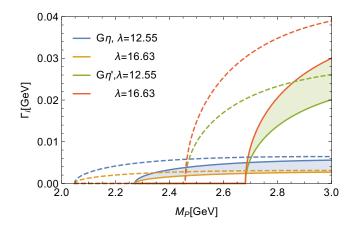


FIG. 1. Partial width of resonant decay $\tilde{G} \to G\eta(')$ (neglecting finite width of scalar glueball) for a predominantly dilatonic scalar glueball G_D with mass $m_D = 1.5$ GeV (dashed lines) and 1.723 GeV (full lines), the latter corresponding to $f_0(1710)$.

In agreement with the considerations of Ref. [44], the WSS model, which also corresponds to a large- N_c chiral Lagrangian, does not lead to a coupling of the pseudoscalar glueball to Im $\log \det \Sigma$. However, its extension to finite quark masses (either through world-sheet instantons [45] or open-string tachyon condensation [46]) does not naturally lead to a coupling to $\operatorname{Im} \operatorname{tr} \mathcal{M}_{q} \Sigma$, because Ramond-Ramond fields do not couple directly to fundamental strings. In the WSS model, the only coupling linear in G is to $\eta_0 G$. This suggests that the pseudoscalar glueball should decay dominantly in $\eta(\prime)$ and the f_0 meson which corresponds to the scalar glueball, or $\eta(\prime)$ and decay products of the latter. According to the WSS model, the decay mode $K\bar{K}\pi$ that is obtained as the dominant one in the approaches mentioned above should instead be strongly suppressed.

When the mass of the pseudoscalar glueball is larger than the mass of the scalar glueball plus the $\eta(')$ mass, the scalar glueball can be produced on-shell. The resulting decay width is displayed in Fig. 1 as a function of the pseudoscalar glueball mass for the glueball mode G_D with mass 1.5 GeV and also when raised in mass to match $f_0(1710)$, which in Ref. [29] we found to be favored by the WSS model.⁶ For the latter case, Fig. 2 shows the (not necessarily resonant) dimensionless partial decay widths Γ_i/M_P for $\tilde{G} \to G\eta(') \to PP\eta(')$ where $P = K, \pi, \eta, \eta'$ with the decay pattern for the scalar glueball $G = f_0(1710)$ obtained in Ref. [29]. With $m_P \sim 2.6$ GeV as predicted by lattice QCD, the pseudoscalar glue-

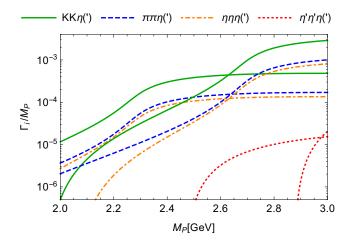


FIG. 2. Partial widths of resonant and non-resonant decays $\tilde{G} \to G\eta(') \to PP\eta(')$ where $P = K, \pi, \eta, \eta'$ assuming the decay pattern for the scalar glueball $G = f_0(1710)$ obtained in Ref. [29]. (The two cases $PP\eta$ and $PP\eta'$ are plotted in the same style but can be distinguished easily by the later onset of $PP\eta'$ which dominates at sufficiently high values of M_P .)

ball is predicted to be a rather narrow state; for $m_P \lesssim 2.3$ GeV it would be extremely narrow (in this case it is of course probable that other, subleading decay channels which are beyond the WSS model become equally if not more important).

IV. PRODUCTION OF PSEUDOSCALAR GLUEBALLS

While scalar and tensor glueballs couple directly to $q\bar{q}$ mesons, pseudoscalar glueballs do so only through the former in the WSS model, because the C_1 Ramond-Ramond field does not couple directly to the DBI and CS action of flavor D8 branes. This suggests that pseudoscalar glueballs are not as easily formed in radiative decays of J/ψ as the other glueballs, but they would have to arise from excited scalar or tensor glueballs decaying into $\eta(')\tilde{G}$ or $\tilde{G}\tilde{G}$ pairs. The thresholds for these processes are thus above the mass of the J/ψ so that excited ψ mesons or Υ would be required. Creation of $\eta(')\tilde{G}$ or $\tilde{G}\tilde{G}$ pairs via virtual scalar and tensor glueballs would also be a possibility for the planned glueball searches in proton-antiproton collisions in the PANDA experiment [47] at FAIR.⁷

⁶ As in Refs. [28–30] we discard the "exotic" scalar glueball mode G_E , assuming it has no counterpart in QCD. If we had identified the mode G_E with the lowest scalar glueball and raised its mass (which is originally only 855 MeV) to the mass of $f_0(1500)$ or $f_0(1710)$, Fig. 1 would look very similar, but the decay width would be about a factor of 10 larger.

⁷ In Ref. [48] a chirally invariant coupling of the pseudoscalar glueball to nucleons and their chiral partners in the so-called mirror assignment was considered. In the WSS model, baryons are described by Skyrmion-like solitons of the effective action of the flavor branes, which likewise excludes a direct coupling to the pseudoscalar glueball at the same order as the direct couplings to scalar and tensor glueballs.

Another possibility is central exclusive production (CEP) in high-energy hadron collisions through double Pomeron or Reggeon exchange (corresponding to G_T and (ρ, ω) trajectories; pion and scalar glueball exchanges are subdominant at high energies). The parametric orders of the corresponding amplitudes are shown in Fig. 3. Production of $\tilde{G}\eta_0$ occurs only via virtual scalar glueballs, whereas production of $\tilde{G}\tilde{G}$ can additionally proceed through virtual tensor glueballs. Also shown is the possibility of $G\tilde{G}$ production through the natural-parity violating coupling of η_0 to two tensor glueballs (Pomerons), which is provided by the Chern-Simons part of the action of the D8 branes and which was recently studied within the WSS model in Ref. [40].⁸

Associated production of pseudoscalar glueballs with either $\eta(')$ or other glueballs is presumably beyond the reach of the older fixed-target experiments searching for glueballs, but seem to be an exciting possibility for the new generation of CEP experiments at the LHC.

Calculation of the corresponding production cross sections within the WSS model could be attempted by employing the techniques used in Ref. [40] for η and η' production (see also [49, 50]), but will be left for future work. In this Letter we only present results for the ratio of production rates of $\hat{G}\eta'$ and $\hat{G}\hat{G}$ pairs over $\eta'\eta'$ pairs,⁹ when both are produced through a virtual G_D glueball. This ratio is fixed by the vertices obtained above together with the results obtained in Ref. [29], and the result is shown in Fig. 4 for the range of 't Hooft coupling discussed above. The amplitude $\mathcal{M}(G^* \to \tilde{G}^2) \sim \lambda^{-1/2} N_c^{-1}$ is parametrically of the same order as $\mathcal{M}(G^* \to \eta^{\prime 2})$ so that the ratio $N(\tilde{G}\tilde{G})/N(\eta'\eta')$ is particularly well determined (at least for fixed meson masses in the scenario of Ref. [29]). The results in Fig. 4 indicate that CEP of $\eta' \tilde{G}$ is only one order of magnitude below CEP of $\eta' \eta'$, while above the threshold for GG pairs, production of the latter is even up to one order of magnitude larger than CEP of $\eta' \eta'$.

Central exclusive production of η' pairs has been studied in the Durham model in Ref. [51], where its production cross section was estimated. For example, at $\sqrt{s} = 1.96$ TeV this work obtained $\sigma(\eta'\eta')/\sigma(\pi^0\pi^0) \sim$ $10^3 \dots 10^5$ assuming sufficiently high transverse momentum such that a perturbative approach becomes justified.

Since small transverse momentum is expected to provide a glueball filter [52, 53] and the production of \tilde{G} together with another \tilde{G} or $\eta(')$ according to the present model proceeds through virtual scalar glueballs, the kinematical regime of small transverse momentum (small azimuthal angle ϕ_{pp}) would be particularly interesting for the search of pseudoscalar glueballs.

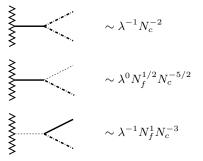


FIG. 3. Parametric orders of the production amplitudes of pseudoscalar glueballs $(\tilde{G}\tilde{G}, \eta(')\tilde{G}, \text{ and } G\tilde{G}, \text{ respectively})$ in double Pomeron or double Reggeon exchange. (Dotted, full, and dash-dotted lines represent $\eta(')$, G, and \tilde{G} , respectively. In the uppermost diagram the full line stands for G or G_T .)

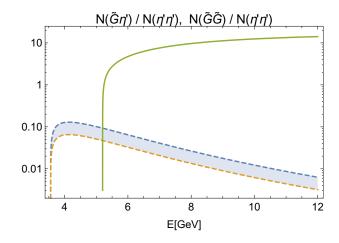


FIG. 4. Production of $\tilde{G}\tilde{G}$ and $\tilde{G}\eta'$ pairs versus $\eta'\eta'$ from a virtual scalar glueball G_D for a pseudoscalar glueball mass of 2.6 GeV as functions of the center of mass energy of the produced pair. The full line gives the ratio of the numbers of produced pairs $N(\tilde{G}\tilde{G})/N(\eta'\eta')$, which is independent of the 't Hooft coupling; upper and lower dashed lines correspond to $N(\tilde{G}\eta')/N(\eta'\eta')$ with 't Hooft coupling 12.55 and 16.63, respectively.

To summarize, the WSS model, which is based on the 't Hooft limit of large $N_c \gg N_f$ where mixing of glueballs with $q\bar{q}$ states is suppressed, suggests a very restricted decay pattern of a rather narrow pseudoscalar glueball G, namely decay into an η' meson together with a scalar glueball, with the latter decaying mostly into pairs of pseudoscalar mesons. In particular, the $K\bar{K}\pi$ decay mode obtained in many other models is found to be suppressed, because the WSS model does not directly couple the pseudoscalar glueball mode carried by the Ramond-Ramond gauge field C_1 to fundamental strings and flavor branes. This is certainly not a universal feature of holographic models and thus will not necessarily hold in other (e.g., bottom-up) holographic approaches to QCD, but the (top-down) WSS model appears to be particularly attractive because it incorporates nonabelian

⁸ A natural-parity violating coupling of η_0 also exists with Reggeons. Fig. 3 gives the parametric order for double Pomeron exchange, which is down by a factor $1/N_c$ compared to Reggeons, but becomes dominant at sufficiently high energies.

⁹ The production rate of $\tilde{G}\eta$ has a smaller threshold and thus larger phase space but is reduced by a factor $(\tan \theta_P)^2 \sim 0.1$.

chiral symmetry breaking as well as the anomaly structure of QCD in a most natural way.

By the same token, the WSS model predicts the production of pseudoscalar glueballs to proceed through excited scalar or tensor glueballs decaying into $\eta' \tilde{G}$ or $\tilde{G} \tilde{G}$ pairs so that the threshold is above radiative J/ψ decays. This could explain why no pseudoscalar glueball candidate with mass in the range of lattice predictions has as yet been found there. Instead, searches in excited charmonium or Υ decays and CEP experiments at high-energy hadron colliders as well as proton-antiproton collisions at FAIR should have the potential for finally discovering the pseudoscalar glueball, with production cross-sections comparable to those of $\eta' \eta'$ pairs.

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Appendix A: Mode expansion of the C_1 Ramond-Ramond field sector

In this appendix we recapitulate some fundamental properties of the WSS model, in particular concerning the sector involving the C_1 Ramond-Ramond field, together with the mode expansion of the latter that is needed to study pseudoscalar glueball interactions.

The metric in the WSS model reads

$$ds^{2} = \left(\frac{u}{R_{\rm D4}}\right)^{3/2} \left[\eta_{\mu\nu}dx^{\mu}dx^{\nu} + f(u)(dx^{4})^{2}\right] \\ + \left(\frac{R_{\rm D4}}{u}\right)^{3/2} \left[\frac{du^{2}}{f(u)} + u^{2}d\Omega_{4}^{2}\right]$$
(A1)

with $f(u) = 1 - (u_{\rm KK}/u)^3$; the nonconstant dilaton is given by

$$e^{\Phi} = (u/R_{\rm D4})^{3/4}.$$
 (A2)

The parameters of the dual field theory are given by [32, 33, 54, 55]

$$g_{\rm YM}^2 = \frac{g_5^2}{2\pi R_4} = 2\pi g_s l_s M_{\rm KK}, \quad R_{\rm D4}^3 = \pi g_s N_c l_s^3.$$
 (A3)

The RR 1-form field $C_1 = C_{\tau}(u, x)d\tau$ contains pseudoscalar glueball modes. For nonvanishing θ -parameter, it also carries nonvanishing flux through the 2-plane parametrized alternatively by (u, τ) or (y, z) with y = 0 being the position of the stack of D8 branes. Anomaly cancellation requires that C_1 transforms nontrivially under U(1) flavor gauge field transformations. This can

be taken into account by replacing its field strength $F_2 = dC_1$ in the 10-dimensional action for C_1 by the gauge invariant combination

$$\tilde{F}_2 = dC_1 + \operatorname{tr}(A) \wedge \delta(y) dy \qquad (A4)$$
$$= dC_1' + \frac{c}{u^4} \left(\theta + \frac{\sqrt{2N_f}}{f_\pi} \eta_0(x) \right) du \wedge d\tau$$

where C'_1 is a reduced RR 1-form field with zero net flux through the (u, τ) -plane and

$$c = \frac{3u_{\rm KK}^3}{\delta\tau}, \quad \delta\tau \equiv 2\pi/M_{\rm KK}$$
 (A5)

such that in the absence of C'_1

$$\partial_u(\sqrt{-g}g^{uu}g^{\tau\tau}\tilde{F}_{u\tau}) = \partial_u(u^4\tilde{F}_{u\tau}) = 0 \qquad (A6)$$

and

$$\int \tilde{F}_{u\tau} du \wedge d\tau = \left(\theta + \frac{\sqrt{2N_f}}{f_\pi} \eta_0(x)\right), \qquad (A7)$$

with $f_{\pi}^2 = \lambda N_c M_{\rm KK}^2 / (54\pi^4)$. (Since we shall be setting $\theta = 0$ in the end, we are ignoring here that a finite θ leads to backreactions on the metric, which have been worked out in [56]. A priori, terms involving higher powers of θ require also contributions with higher powers of η_0 fields. We have checked, however, that for $\theta = 0$ inclusion of this backreaction does not lead to additional vertices involving η_0 and pseudoscalar glueball modes beyond those worked out below.)

The action for the 1-form RR field is given by

$$S_{C_1} = -\frac{1}{4\pi (2\pi l_s)^6} \int d^{10}x \sqrt{-g} |\tilde{F}_2|^2.$$
 (A8)

The reduced C_1' will be expanded in pseudoscalar glueball modes, $C_1' = C_\tau' d\tau$ and

$$C'_{\tau}(u,x) = \sum_{n=1}^{\infty} V^{(n)}(\bar{u})\tilde{G}^{(n)}(x)$$
 (A9)

with radial mode functions $V^{(n)}(\bar{u}) = f(\bar{u})\bar{V}^{(n)}(\bar{u})$ satisfying

$$-\bar{u}^{-1}\frac{d}{d\bar{u}}\left(\bar{u}^{4}\frac{d}{d\bar{u}}\left[f(\bar{u})\bar{V}^{(n)}(\bar{u})\right]\right) = \frac{9}{4}\frac{(M_{P}^{(n)})^{2}}{M_{\rm KK}^{2}}\bar{V}^{(n)}(\bar{u})$$
(A10)

where $\bar{u} = u/u_{\rm KK}$ and $f(\bar{u}) = 1 - \bar{u}^{-3}$. The two lowest normalizable solutions with $V^{(n)}(\bar{u} = 1) = V^{(n)}(\infty) =$ 0 but $\bar{V}^{(n)}(\bar{u} = 1) \neq 0$ have the eigenvalues $M_P^{(1)} \approx$ 1.885 $M_{\rm KK}$ and $M_P^{(2)} \approx 2.838 M_{\rm KK}$, respectively.

With this mode expansion which keeps all fields independent of the compactified coordinate τ and the coordinate τ

dinates of the S^4 we have

$$S_{C_1} = -\frac{1}{4\pi (2\pi l_s)^6} \int d^{10}x \sqrt{-g} \left\{ g^{mn} g^{\tau\tau} \partial_m C'_{\tau} \partial_n C'_{\tau} + g^{uu} g^{\tau\tau} \frac{c^2}{u^8} \left(\theta + \frac{\sqrt{2N_f}}{f_{\pi}} \eta_0(x) \right)^2 + 2g^{mu} g^{\tau\tau} \partial_m C'_{\tau} \frac{c}{u^4} \left(\theta + \frac{\sqrt{2N_f}}{f_{\pi}} \eta_0(x) \right) \right\}, (A11)$$

with indices $m, n \in \{0, 1, 2, 3, u\}$.

Inserting the background metric of the WSS model and setting $\theta = 0$ produces the kinetic terms in (5) with Witten-Veneziano mass (2). The last term in (A11) which is proportional to $\tilde{G}\eta_0$ does not give rise to a mixing of \tilde{G} and η_0 because it vanishes after radial integration. However, in the presence of metric fluctuations it no longer vanishes and gives rise to the interaction terms in $\mathcal{L}_{\tilde{G}\eta_0 G}$ listed in Eq. (6).

In order to determine the values of interaction vertices,

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we need to normalize the pseudoscalar glueball fields. Demanding that the pseudoscalar glueball fields $\tilde{G}^{(n)}(x)$ appearing in the mode expansion of C'_{τ} have canonical kinetic terms fixes the normalization of the radial mode functions through

$$\frac{\Omega_4 \delta \tau}{2\pi (2\pi l_s)^6} R_{\mathrm{D4}}^3 \int_{u_{\mathrm{KK}}}^{\infty} du \, u f^{-1}(u) [V^{(n)}(u)]^2$$
$$= \frac{\lambda^3}{4 \cdot 3^5 \pi^4} \int_1^{\infty} d\bar{u} \, \bar{u} f^{-1}(\bar{u}) [V^{(n)}(\bar{u})]^2 = 1 \quad (A12)$$

with $\Omega_4 = 8\pi^2/3$ being the volume of the unit S^4 . For the lightest and the first excited pseudoscalar glueball mode this implies

$$[\bar{V}^{(1)}(\bar{u}=1)]^{-1} = 0.002046\dots\lambda^{3/2},$$
 (A13)

$$[\bar{V}^{(2)}(\bar{u}=1)]^{-1} = 0.001157\dots\lambda^{3/2}.$$
 (A14)

Using the mode expansions of the metric fields given in Ref. [28], the effective Lagrangian for pseudoscalar glueball interactions can be obtained by numerical integrations over products of the relevant radial mode functions.

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