How to detect the lightest glueball

B. P. Kosyakov\textsuperscript{a,b}, E. Yu. Popov\textsuperscript{a}, and M. A. Vronskii\textsuperscript{a,c}

\textsuperscript{a}Russian Federal Nuclear Center–VNIIEF, Sarov, 607188 Nizhni Novgorod Region;  
\textsuperscript{b}Moscow Institute of Physics & Technology, Dolgoprudni, 141700 Moscow Region;  
\textsuperscript{c}Sarov Institute of Physics & Technology, Sarov, 607190 Nizhni Novgorod Region

Abstract

We suggest a procedure for detecting the lightest glueball in a head-on collision of photons whose center-of-mass energy is in the range $1.3 - 2$ GeV. We use recent evidence for scattering of light by light in Large Hadron Collider experiments as a phenomenological basis for our suggestion. With this evidence, the cross section of the lightest glueball creation in $\gamma\gamma$ collisions is estimated to be $\sim 60$ nb. The predominant mode of the lightest glueball decay, predicted from the gauge/gravity duality, proves to be the decay into two neutral vector mesons $\rho^0\rho^0$. Because the $\rho^0$ decays into $\pi^+\pi^-$, a drastic increase of the $\pi^+\pi^-\pi^+\pi^-$ yield is expected as the center-of-mass energy approaches the value of the lightest glueball mass. This fact will be the unique signature of the lightest glueball detection.

Key words: glueball, photon collider, predominant mode of the glueball decay, gauge/gravity duality
1 Introduction

The prediction of hadrons containing no quarks, composed of gluons alone, was made [1]–[4] at the dawn of the age of quantum chromodynamics (QCD), the modern theory of the strong interactions. In the past half-century, prodigious experimental efforts went into searching such particles, now known as glueballs. However, glueballs have not yet been observed with certainty [5, 6]. Why so? Are glueballs absent from Nature altogether? Another belief prevails currently: the glueball field mixes with the quark-antiquark fields, \((u\bar{u} + d\bar{d})/\sqrt{2}\) and \(s\bar{s}\), to form the experimentally observed meson resonances [7]–[11].

Nevertheless, there are a few dissenters from the view that the mixing of glueballs and quarks is unavoidable. For example, it was stated in [12]–[14] that the depth of quark-gluon plasmas, formed in heavy-ion collisions, may have a beneficial effect on the creation of pure ground state glueballs. To verify this idea, we should be capable of distinguishing the possible yield of glueballs among many thousands of foreign tracks related to the explosion products of a quark-gluon plasma lump. This is a challenge. The multiplicities of glueballs in a single central heavy-ion collision is estimated by fitting the hadron ratios observed in \(\text{Pb} - \text{Pb}\) collisions at various energies in the Large Hadron Collider (LHC) to be \(1.5 - 4\) glueballs [15].

In the present paper, we propose an alternative procedure for detecting pure ground state glueballs. Our concern here is with the lightest glueball, a color singlet of two gluons, specified by zero total angular momentum and positive parity and charge parity, \(J^{PC} = 0^{++}\). We denote this particle by \(G\). It was anticipated in the pioneer studies [1]–[4] that a scalar glueball can decay into two photons with opposite polarizations. The reverse process provides us with a way for creating pure ground state glueballs: \(\gamma\gamma \rightarrow G\). We explicate this idea, endeavor to justify it phenomenologically and show its feasibility in Sect. 2. The central theoretical problem though is to predict the decay products of \(G\). Section 3 is devoted to this problem.

Unlike quarks, gluons are uncharged, and have vanishing weak hypercharge and isospin couplings, so that \(G\) is immune from the electromagnetic and weak interactions. We thus expect that the strong interactions are responsible for the decay of \(G\). This is tantamount to stating that the lightest glueball is subject to a deconfinement through its splitting into two gluons. Now we come into a low-energy domain in which the QCD running coupling constant \(\alpha_s(\mu)\) is large, and the perturbation technique is inoperative. Three most-used ways for the description of phenomena in this domain are:

1. semiclassical approach,
2. lattice QCD simulations,
3. gauge/gravity duality.

In the settled semiclassical picture, quarks are represented by point particles linked together by thin tubes which enclose the total color flux of the gluon field. Nucleons are colorless objects composed of three valent quarks. They are assumed to be assembled in nuclei by a residual color interaction similar to the multipole van der Waals force between neutral molecules. Meanwhile the capability to give a precise meaning to the notion of the residual color interaction turns out to be rather problematic [16]. That is why the Yukawa mechanism, thought of as a meson exchange mechanism, refined by
several technical innovations, such as spontaneously broken chiral symmetry, effective Lagrangians, and derivative expansions, still forms the basis for modern nuclear physics \[17\]–\[19\]. And yet the issue of understanding nuclei in terms of quarks is high on the agenda of the QCD developments. We do not dwell on the efforts to address this issue in the framework of different semiclassical approaches (bags, potential models, etc.), and refer the interested reader to \[20\] and references therein.

The Coleman theorem \[21\] dramatically hampers semiclassical treatments of glueballs. By this theorem, no localized (kink-like) finite-energy solutions of pure Yang–Mills theory is available. Strictly speaking, there is no rendition of a glueball as a localized object; we have not the foggiest notion of what the size and structure are peculiar to the lightest glueball. For lack of color fields that issue out of fixed points and cancel each other, the notion of residual color interactions becomes quite farfetched. We thus have to state that the lightest glueball does not interact with its environment until it splits into two gluons.

According to lattice and sum rule calculations, the lightest glueball has mass in the range of about \(1.3–2\ \text{GeV}\) \[7\]–\[11\]. Furthermore, lattice QCD predicts the glueball mass spectrum. However, the glueball coupling with ordinary hadrons hitherto eluded reliable analyzing.

In this paper, the decay of \(G\) is examined in the context of gauge/gravity duality, aka the correspondence between a quantum theory of gravity in anti-de Sitter space and conformal field theories in Minkowski space, AdS/CFT, and the holographic principle \[22\]–\[24\]; for a full coverage of ideas and methods of gauge/gravity duality see \[25, 26\]. Loosely speaking, gauge/gravity duality is a doctrine whereby a good part of subnuclear physics in a four-dimensional realm is modelled on physics of black holes and similar objects (black rings, black branes, etc.) in five-dimensional anti-de Sitter space, AdS\(_5\), whose boundary is just this four-dimensional realm. However, this understanding of gauge/gravity duality apparently indulges in wishful thinking. In their popular science paper \[27\] Klebanov and Maldacena remind the reader of the physics joke about the spherical cow as an idealization of a real one, and admit that “in the AdS/CFT correspondence, theorists have really found a hyperbolic cow”. To remedy the situation, a major portion of the standard holographic mapping is to be amputated.

The main stream in the high-energy-physics research today develops the idea that a black hole in AdS\(_5\) is mapped holographically onto a quark-gluon plasma lump in a four-dimensional realm \[28\]. This realization of holography in the framework of the Bekenstein–Hawking thermal treatment of gravitational phenomena in AdS\(_5\) was offered in \[29\]. The line of reasoning is as follows. Let us compare the values that the gravitational action

\[
I = -\frac{1}{2\kappa^2} \int d^5x \sqrt{g} \left( R + \frac{12}{L^2} \right)
\]

(1)

takes when two solutions are substituted in it, one describing thermal AdS\(_5\) and the other describing thermal AdS\(_5\) with a Schwarzschild black hole \[\] It then transpires that the former is less than the latter at \(T < T_c\), where \(T_c\) is the Hawking-Page phase transition

\(^1\)To eliminate divergences of the action \[\] arising from these substitutions, the integral must be regularized by means of a cut-off, or using a nontrivial dilaton field expectation value \[29\].
temperature, and conversely, the latter is less than the former at $T > T_c$. It follows that the thermal AdS$_5$ becomes unstable at $T > T_c$ to yield the black hole formation. The holographic image of this process on the four-dimensional boundary of AdS$_5$ is a QCD phase transition implying the quark-gluon plasma formation at a critical temperature associated with $T_c$.

In an alternative approach [30]–[33], black Dp-branes are mapped holographically onto a subnuclear realm in the confinement phase. The set of significant entities suitable for the holographic mapping was defined in [34] to be that containing only extremal black holes and other extremal black objects, which correspond to their holographic counterparts, stable microscopic systems.

These two seemingly incompatible concepts of gauge/gravity duality are in fact quite consistent. The Bekenstein–Hawking thermal analysis of gravitational phenomena forms the basis of the former, whereas thermal treatments are unrelated to the latter because extremal black objects live in a cold realm, $T = 0$; they do not experience Hawking evaporation. Their associated holographic counterparts from different QCD phases may belong to distinct sectors of the holographic mapping.

In Sect. 3, following the general ideology of [34], which narrows the holographic context down to extremal black holes and their stable subnuclear counterparts, we yet seek for extracting useful information from the frontier zone separating the domain in which the holographic principle holds and that in which this principle fails.

2 How to create the lightest glueball

It was already noted that kinematics and symmetries allow the decay of $G$ into two gamma-quanta with opposite spiralities. The lower orders in $\alpha$ and $\alpha_s$ of this process is depicted in Figure 2 as diagram (c). The inverted diagram corresponds to the reverse of this process: a head-on $\gamma\gamma$ collision at a center-of-mass energy $\sqrt{s}$ in the range $1.3 − 2$ GeV, with the helicity of the $\gamma\gamma$ system being zero, may result in the lightest glueball creation $^2$.

To test this idea, it is attractive to use a photon collider, the most extensively studied prospective device $^3$ having its origin in the conversion of laser photons into high-energy gamma-quanta through the Compton scattering on high-energy electrons [36]. The device consists of two beams of electrons moving towards each other to the interaction point $x_\ast$. The electrons collide with laser photons at a distance of about $1 \sim 5$ mm from $x_\ast$. After the scattering, the photons become gamma-quanta with energy comparable to that of the electrons and follow their direction to $x_\ast$ where they collide with similar counterpropagating gamma-quanta. The maximum energy $\omega$ of the gamma-quanta is

---

$^2$We speak about this event in tentative modality because our concern is not with a conversion $\gamma\gamma \rightarrow gg$, Figure 1 (b), which can be accounted for in the framework of the conventional perturbation theory, but with the occurrence of a bound state of two gluons. Now we are not in a position to look into the formation of gluon confinement in detail, but some insight into this process will be gained from gauge/gravity duality in the next section.

$^3$For a technical design report of the Photon Collider at TESLA see [35].
given by
\[ \omega = \frac{x}{x + 1} E, \quad x \approx \frac{4E\omega_0}{m^2}, \]  
(2)
where \( E \) and \( \omega_0 \) are, respectively, the energy of the electrons and laser photons, and \( m \) the electron mass. For example, \( E = 7.5 \) GeV is needed if we are to convert the photon energy \( \omega_0 = 1.17 \) eV (Nd: glass laser) into the gamma-quantum energy \( \omega = 0.85 \) GeV.

Using a laser with a flash energy of several joules one can obtain gamma-quantum whose spot size at \( x_* \) will be almost equal to that of the electrons at \( x_* \), and the total luminosity of \( \gamma\gamma \) collision will be comparable to the “geometric” luminosity of the electron beams. The energy spectrum of the gamma-quantum becomes most peaked if the initial electrons are longitudinally polarized and the laser photons are circularly polarized. This gives almost a factor of 4 increase of the luminosity in the high-energy peak. The present laser technology has all elements needed for the required photon colliders [35].

In order to evaluate the feasibility of the conversion \( \gamma\gamma \rightarrow gg \) in the discussed layout, we invoke recent evidence for the scattering of light by light in quasi-real photon interactions of ultra-peripheral Pb+Pb collisions, with impact parameters larger than twice the radius of the nuclei, at a nucleon-nucleon center-of-mass energy \( \sqrt{s} = 5.02 \) TeV by the ATLAS experiment at the LHC [37, 38]. The fiducial cross section of the process \( \text{Pb} + \text{Pb}(\gamma\gamma) \rightarrow \text{Pb}^{(*)} + \text{Pb}^{(*)}\gamma\gamma \), for diphoton invariant mass greater than 6 GeV, is measured to be \( 78 \pm 13 \) (stat.) \( \pm 7 \) (syst.) nb. This result is in agreement with the Standard Model [39]–[42], where the light-by-light scattering arises, in the leading order of \( \alpha \), via one-loop diagrams, Figure 1(a). An important point is that if we recalculate (according to [39]) the ATLAS experiment result to the elementary cross section of the light by light scattering in vacuum at \( \sqrt{s} = 1.5 \) GeV, we obtain \( \sigma_{\gamma\gamma \rightarrow \gamma\gamma} \sim 70 \) pb.

Figure 1: \( \gamma\gamma \) collisions resulting in: (a) a two-photon system; (b) a two-gluon system; (c) an aggregate of a glueball and a meson.

Another piece of information derives from the usual QCD calculations of quarkonium partial widths [43, 44]. The rule for changing two external photon lines by two external gluon lines [see Figure 1 respectively, plots (a) and (b)] refers to the factor
\[ \frac{9 \alpha_s^2(m_c)}{8 \alpha^2} \approx 845, \]  
(3)
where $\alpha_s(\mu)$ stands for the QCD running coupling constant, which being taken at $\mu$ equal to the mass of charmed quarks $m_c = 1.28$ GeV is $\alpha_s \approx 0.2$. Indeed, this value of $\alpha_s$ can be obtained \[45, 46\] from the ratio

$$\frac{\Gamma(J/\phi \rightarrow \text{hadrons})}{\Gamma(J/\phi \rightarrow e^+e^-)} = \frac{5(\pi^2 - 9)\alpha_s^2}{18\pi\alpha^2},$$

whose experimental value is $\approx 10$.

All things considered, the cross section of the lightest glueball creation in head-on $\gamma\gamma$ collisions at energy region about $\sqrt{s} = 1.7$ GeV is expected to be $845 \sigma_{\gamma\gamma \rightarrow \gamma\gamma} \approx 60$ nb.

Note that it is the pure ground state glueball which will be produced in the proposed experiment; the creation of $G$ is not obscured by the mixing effects with isoscalar $q\bar{q}$ mesons. To see this, we compare the probabilities of production of an unmixed scalar glueball and an aggregate of a glueball and a meson [Figure 1, respectively, plots (b) and (c)]: the latter, being represented by higher perturbative orders, is suppressed by a factor of $\alpha_s^4 \approx 1.6 \cdot 10^{-3}$ as against the former.

Finally, the fact that the Stanford Linear Collider luminosity, $\approx 3 \cdot 10^{30}\text{cm}^2/\text{s}$, is four orders of magnitude greater than the luminosity $\approx 5 \cdot 10^{26}\text{cm}^2/\text{s}$ of the light by light scattering event in Pb − Pb collisions at the LHC [37, 38] counts in favor of feasibility of photon colliders with fairly high luminosity. This is an added reason for the proposed experiment layout.

3 Decay of $G$

We assume that the lightest glueball is unstable, and hence we would like to know its decay channels. Our prime interest here is with the decay products of the predominant mode. The detection of just this yield in the photon collider experiment may evidence that it is the lightest glueball which has been created in a $\gamma\gamma$ collision.

The following criterion for discriminating between stable and unstable systems of the subnuclear world was offered in [34]: a system is stable when its gravitational dual is an extremal black object. It may appear that the converse is also true, namely if a microscopic system is unstable, then its gravitational counterpart is an ordinary black hole amenable to Hawking evaporation. But such is not the case. An obstacle for gaining a well-defined correspondence between unstable microscopic systems and evaporating black holes relates to basic tenets of quantum mechanics, the principle of reversibility and the principle of identity and indistinguishability.

Quantum mechanical processes are reversible. Suffice it to say that the probability amplitude for a decay equals that for the pertinent recombinaction. By contrast, the black hole evaporation is an irreversible process.

All microscopic systems of the same species are identical and indiscernible. Their gravitational duals must exhibit identical properties. Suppose that a Lorentz frame $O'$ moves past another Lorentz frame $O$ at constant velocity $V$, and the standard synchrony of their clocks is established. Let $O$ and $O'$ meet, and, at the instant of their meeting, both reset their clocks to 0. Suppose that $O$ carries a black hole whose properties (mass,
electric charge, and spin), at $t = 0$, are identical to those of a further black hole assigned to $O'$. The rate of evaporation, as measured on the proper time, is common for both black holes. Therefore, at $t > 0$, the black holes possessed by $O$ and $O'$ have different masses; the relativistic effect of time dilation keeps evaporating black holes from being regarded as identical objects.

Our assumption that $G$ is an unstable particle seems to forbid us from invoking the holographic principle because this principle holds for stable systems of the microscopic world and fails for unstable systems. Nevertheless, there is a frontier zone in which the holography could be handled, with extreme care, as a guiding principle, to clarify general features of the decay mechanism typical for unstable particles of this zone. What are the reasons for this anticipation?

The decay of a particle may be attributed not only to an incessant evaporation process of a black hole but also to a single act of spontaneous splitting of the black hole into two or several black holes. If this splitting gives rise to extremal black holes, then the process is apparently reversible: isolated extremal black holes remain unchanged for any length of time, and their collision can recover the initial black hole.

We next imagine that a single black hole is in five-dimensional anti-de Sitter universe, or else, a black hole is widely separated from other matter, and can be considered to be almost free from force interactions. Then the issue of quantum mechanical identity of this black hole with objects of the same kind is of little significance. A holographic counterpart of this black hole is a microscopic system which is exceptional in the sense that the vast majority or even all but one of ways for its decay are forbidden.

We will focus upon the decay of neutral spinless particles whose gravitational duals are taken to be Schwarzschild black holes.

Such black holes in $\text{AdS}_5$ have the greatest possible spatial isometry group $SO(4)$, equivalent to $SO(3) \times SO(3)$, which is holographically mapped upon the $SU(2)_L \times SU(2)_R$ chiral invariance of QCD with $N_f = 2$ flavors. But chiral invariance is spontaneously broken in the confinement phase down to the isospin $SU(2)_V$ symmetry. Therefore, it is reasonable to expect that the dual isometry group $SO(4)$ is also broken down to $SO(3)$. In other words, a Schwarzschild black hole in $\text{AdS}_5$ is amenable to spontaneous splitting into black objects whose symmetry is limited to $SO(3)$, such as spinning black holes of Myers and Perry.

By the duality argument, a neutral spinless particle dual to this Schwarzschild black hole must decay into spinning particles. We emphasize again that this rule is only valid for particles from the frontier zone. The lightest neutral spinless meson $\pi^0$, whose weak and strong interactions are suppressed, decays through the electromagnetic channel as

---

4The reader will easily find this situation to be a kind of the well-known twin paradox in which the twins are the evaporating black holes, and the responsibility for their identification rests with the quantum mechanical principle of identity and indistinguishability of particles.

5This act may be thought of as a possible scenario for the completion of the history of an evaporating black hole. For the scenarios proposed in the literature see [17].

6An apparent objection to this statement is that the Schwarzschild geometry is stable against small perturbations in the classical context. However, the case in point is a quantum tunnelling of one black hole through the event horizon of the other black hole, a phenomenon which falls into the Hawking radiation pattern.
\[ \pi^0 \to \gamma \gamma. \] The lightest scalar particle immune from the electromagnetic and strong interactions, the Higgs boson \( H^0 \), decays into pairs of heavy fermions (\( b\bar{b}, \tau\bar{\tau} \)) or gauge vector bosons (\( W^+W^-, ZZ, gg, \gamma\gamma \)). The lightest glueball is likely to fall into the same category because the only conceivable way for its decay is related to a splitting into two vector particles \( G \to gg \).

Observable manifestations of this rule for the decay of \( G \) can be established in the framework of the usual perturbation theory whose lower orders in \( \alpha \) and \( \alpha_s \) are depicted in Figure 2. Photons, quarks, and gluons are shown as the sine waves, oriented lines, and spirals, respectively. The outgoing vector mesons, composed of quarks and antiquarks, are represented by couples of antiparallel rays. Diagram (a) illustrates the decay mode whose outcome is a pair of truly neutral light-quark vector mesons, \( \rho^0\rho^0 \), or, alternatively, \( \omega\omega \). The masses of these particles are, respectively, \( m_{\rho^0} = 775 \text{ MeV} \), and \( m_\omega = 783 \text{ MeV} \). Diagram (b) displays the outcome as a photon and a truly neutral vector meson, which may be given by either \( \rho^0 \), or \( \omega \), or \( \phi \) (\( m_\phi = 1019 \text{ MeV} \)). Diagram (c) sketches a two-photon decay mode. The ratio of probabilities of these modes can be roughly estimated as \( 1 : O(\alpha) : O(\alpha^2) \).

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig2.png}
\caption{Decay modes for the lightest glueball}
\end{figure}

It follows that the holographically inspired predominant decay mode of the lightest glueball is \( G \to \rho^0\rho^0 \). With the fact that \( \rho^0 \) decays into \( \pi^+\pi^- \) (\( \approx 100\% \) fraction; \( \Gamma = 149.1 \pm 0.8 \text{ MeV} \)) [6], one may expect a drastic increase in the \( \pi^+\pi^-\pi^+\pi^- \) yield as the center-of-mass energy \( \sqrt{s} \) approaches the value of the lightest glueball mass \( m_G \), with the orbital momentum quantum number \( l \) being equal to 1 for each \( \pi^+\pi^- \) pair.

4 Discussion

The use of \( \gamma\gamma \) collisions for the possible yield of glueballs has already been realized as ingredients of experimental studies on \( ee \) and \( e\bar{e} \) collisions [49–51]. However, the physics behind those experiments differs from the physics behind the hunting of the lightest glueball proposed here. Unlike the former which has to do with virtual photons, the latter bears on real photons. According to the concept of the vector meson dominance, virtual photons are capable of creating neutral vector mesons, Figure 3 (a). The most plausible
scenario for the collision of such photons relates to the conversion of the created vector mesons into a pair of scalar mesons\cite{footnote1}, represented by planar tree diagrams, Figure 3(b). If higher order diagrams, with gluon lines being taken into account, the resonances appearing in the cross section may be attributed to the interposition of a glueball. In fact, this effect is due to a glueball mixed with $q\bar{q}$ states. This comes into particular prominence from Figure 3(c): the glueball history is sandwiched between the quark and antiquark world lines.

On the other hand, the probability amplitude that a real photon spontaneously turns to a massive particle is very small, while a head-on collision of two such photons brings into existence of a massive entity $G$. If the total helicity of the colliding photons is zero, and the center-of-mass energy $\sqrt{s}$ equals the mass of the lightest glueball, this $G$ proves to be just the lightest glueball. Although the outcome of the array of reactions completed in the above experiments, $\gamma\gamma \to \phi\phi \to K\bar{K} \to 4\pi$, is similar to that of the array of the proposed reactions, $\gamma\gamma \to G \to \rho^0\rho^0 \to 4\pi$, the contents of the compared processes seem much different.

![Figure 3: The strong interaction of photons: (a) a virtual photon becomes a vector meson; (b) two virtual photons collide to give a $K\bar{K}$ system; (c) two virtual photons collide to give a $K\bar{K}$ system through the mediation of a glueball.](image)

It might be well to mention a recent event\cite{footnote2}, indirectly touching the subject of our discussion, – the experimental discovery of the odderon (a virtual colorless three-gluon state), which is likely to be of great importance in the dynamics of three-gluon vector glueball.

### 5 Concluding remarks

Glueballs are enigmatic objects. A major portion of their properties are associated with infrared effects which are beyond the control of the conventional perturbation theory. The derivative expansion in chiral perturbations is also irrelevant here because the lightest glueball $G$ is heavier than 1 GeV.

Semiclassical treatments are hampered by the fact that no rendition of a glueball is available due to the Coleman theorem\cite{footnote3}. We thus have not the slightest notion what

\footnote{Say into $K\bar{K}$, which is typical for the experiments analyzed in\cite{footnote4}–\cite{footnote5}.}
the size and structure are peculiar to the lightest glueball. Lattice QCD predicts the 
glueball mass spectrum. However, the glueball coupling with ordinary hadrons still does 
not have a good lattice-based method of solution.

It remains to try to address the gauge/gravity duality or, what is the same, the 
holography. As was made clear in Sect. 3, the gauge/gravity duality bears no relation to 
unstable microscopic systems. However, there is a frontier zone separating the domain in 
which the holography holds and that in which it fails. We argued that the holography can 
be used as a guiding principle to clarify qualitative properties of the decay mechanism 
typical for unstable particles from the frontier zone. The lightest glueball is just among 
the set of particles from the frontier zone. The holographic reasoning together with 
calculations of $m_G$ predict the predominant decay mode of $G$ to be $G \rightarrow \rho^0\rho^0$.

In Sect. 2 we proposed a procedure of detecting the lightest glueball in a head-on $\gamma\gamma$ 
collision, and adduced some phenomenological justifications. Our proposal holds much 
practical promise because the photon collider experiments are an extensively elaborated 
project, ready to be implemented (for the state of the art see, e. g., [35]). The cross section 
of the process $\gamma\gamma \rightarrow G$ estimated as $\sim 60$ nb is in general consistent with the results of 
experiments for exploring the process $\gamma\gamma \rightarrow \rho^0\rho^0$ in the same energy range which had 
already been made in the 80s (see, e. g., [53] and references therein).

In addition, the idea to detect $G$ in $\gamma\gamma$ collisions has truly theoretical virtues. There 
is good reason to believe that an unmixed scalar glueball will be created in the proposed 
experiment, which is confirmed by the estimated cross section of the $G$ creation and 
feasible luminosity of photon colliders. By now, five isoscalar resonances are established: $f_0(500)$, $f_0(980)$, $f_0(1370)$, $f_0(1500)$, and $f_0(1710)$ \[^8\] whose masses are in the range of 
about $0.5 - 2$ GeV \[^8\]. Among them, only $f_0(1710)$ can be suspected to be an unmixed 
scalar glueball \[^54, 55\], but this view was challenged \[^56\]. If the proposed experiment 
will exhibit a comparatively narrow state at the energy $\sqrt{s}$ other than the mass of any 
one of these resonances, this will strongly suggest that a new kind of hadron matter 
without quarks, the lightest glueball, has been discovered. A precise measurement of $m_G$ 
will provide further impetus and guide the way for improving lattice QCD. The expected 
drastic increase in the $\pi^+\pi^-\pi^+\pi^-$ yield at $\sqrt{s} = m_G$ will be a direct experimental evidence 
in support of the gauge/gravity duality prediction.

Acknowledgment

We are grateful to Masud Chaichian, Wolfgang Ochs, Valery Tel’nov, and Vicente Vento 
for useful discussions.

References

[1] H. Fritzsch and M. Gell-Mann. Current algebra: Quarks and what else? EConf 
C720906V2 2, 135-165, (1972); hep-ph/0208010. 

\[^8\]Note that these resonances decay into spinless (rather than spinning) particles $\pi\pi$, $KK$, $\eta\eta$, $\eta\eta’$ \[^6\].
[2] H. Fritzsch and P. Minkowski. $\Psi$-resonances, gluons and the Zweig rule. Nuovo Cimento A 30, 393-429 (1975).

[3] P. G. O. Freund and Y. Nambu. Dynamics of the Zweig–Iizuka rule and a new vector meson below 2 GeV/$c^2$. Phys. Rev. Lett. 34, 1645-1649 (1975).

[4] R. L. Jaffe and K. Johnson. Unconventional states of confined quarks and gluons. Phys. Lett. B 60, 201-204 (1976).

[5] V. Crede and C. A. Meyer. The experimental status of glueballs. Prog. Part. Nucl. Phys. 63, 74-116 (2009); hep-ex/0812.0600v3.

[6] P. A. Zyla et al. (Particle Data Group). Review of Particle Physics. Prog. Theor. Exp. Phys. 2020: 083C01 (2020).

[7] V. V. Anisovich. Exotic mesons: the search for glueballs. Phys. Usp. 38, 1179-1201 (1995) [Transl. from Usp. Phys. Nauk 165, 1225-1247 (1995)]; The lightest scalar glueball. Ibid. 41, 419-439, 1998 [Transl. from Usp. Phys. Nauk 168, 481-502 (1998)]; Systematics of quark-antiquark states and scalar exotic mesons. Ibid. 47, 45-67 (2004) [Transl. from Usp. Phys. Nauk 174, 49-72 (2004)].

[8] F. E. Close and N. A. Törnqvist. Scalar mesons above and below 1 GeV. J. Phys. G 28, R249-R267 (2002); hep-ph/0204053v3.

[9] C. Amsler and N. A. Törnqvist. Mesons beyond the naive quark model. Phys. Rep. 389, 61-117 (2004).

[10] V. Mathieu, N. Kochelev, and V. Vento. The physics of glueballs. Int. J. Mod. Phys. E 18, 1-49 (2009); hep-ph/0810.4453.

[11] W. Ochs. The status of glueballs. J. Phys. G 40: 043001 (2013); hep-ph/1301.5183.

[12] S. Kabana and P. Minkowski. Glueball production in hadron and nucleus collisions. Phys. Lett. B 472, 155-160 (2000); hep-ph/9907570.

[13] N. Kochelev and D. P. Min. Role of glueballs in non-perturbative quark-gluon plasma. Phys. Lett. B 650, 239-243 (2007); hep-ph/0611250.

[14] V. Vento. Glueball enhancement by color de-confinement. Phys. Rev. D 75:055012, 380-385 (2007); hep-ph/0609219.

[15] I. N. Mishustin, L. M. Satarov, and W. Greiner. Possible glueball production in relativistic heavy-ion collisions. J. Phys. G. 32, L59-L63 (2006); hep-ph/0606251.

[16] B. P. Kosyakov, E. Yu. Popov, and M. A. Vronska. The bag and the string: Are they opposed? Phys. Lett. B 744, 281P33 (2015).

[17] S. Weinberg. Nuclear forces from chiral lagrangians. Phys.Lett. B 251, 288-292 (1990)
[18] E. Epelbaum, H.-W. Hammer, and U.-G. Meißner. Modern theory of nuclear forces. Rev. Mod. Phys. 81, 1773-1825 (2009); nucl-th/0811.1338.

[19] R. Machleidt and D. R. Entem. Chiral effective field theory and nuclear forces Phys. Rept. 503, 1-75 (2011); nucl-th/1105.2919.

[20] B. P. Kosyakov, E. Yu. Popov, and M. A. Vronskaï. Could the static properties of nuclei be deduced from the dynamics of a single quark? Eur. Phys. J. A 53: 82 (2017); hep-ph/1604.06613.

[21] S. Coleman. There are no classical glueballs. Commun. Math. Phys. 55, 113-116 (1977).

[22] J. Maldacena. The large $N$ limit of superconformal field theories and supergravity. Adv. Theor. Math. Phys. 2, 231-252 (1998); hep-th/9711200.

[23] E. Witten. Anti-de Sitter space and holography. Adv. Theor. Math. Phys. 2, 253-291 (1998); hep-th/9802150.

[24] S. S. Gubser, I. R. Klebanov, and A. M. Polyakov. Gauge theory correlators from noncritical string theory. Phys. Lett. B 428, 105-114 (1998); hep-th/9802109.

[25] M. Ammon and J. Erdmenger. Gauge/Gravity Duality (CUP, Cambridge, 2015).

[26] H. Nâstase. Introduction to the AdS/CFT Correspondence (CUP, Cambridge, 2015).

[27] I. R. Klebanov and J. M. Maldacena. Solving quantum field theories via curved spacetime. Phys. Today 62 (1), 28-33 (2009).

[28] I. Ya. Aref’eva. Holographic approach to quark-gluon plasma in heavy ion collisions Phys.-Usp. 57 527-555 (2014) [Transl. from Usp. Fiz. Nauk 184, 569-598 (2014)].

[29] C. P. Herzog. A holographic prediction of the deconfinement temperature. Phys. Rev. Lett. 98: 091601 (2007); hep-ph/0608151.

[30] T. Sakai and S. Sugimoto. Low energy hadron physics in holographic QCD. Prog. Theor. Phys. 113, 843-882, (2005); hep-ph/0412141.

[31] T. Sakai and S. Sugimoto. More on a holographic dual of QCD. Prog. Theor. Phys. 114, 1083-1118 (2006); hep-ph/0507073.

[32] K. Hashimoto, Ch.-I Tan, and S. Terashima. Glueball decay in holographic QCD. Phys. Rev. D 77, 086001 (2008); hep-ph/0709.2208.

[33] F. Bruenner, D. Parganlija, and A. Rebhan. Glueball decay rates in the Witten–Sakai–Sugimoto model. Phys. Rev. D 91, 106002 (2015); hep-ph/1501.07906.

[34] B. P. Kosyakov, E. Yu. Popov, and M. A. Vronskii. Correspondence between the physics of extremal black holes and that of stable heavy atomic nuclei. Class. Quantum Grav. 36: 135001 (2019); hep-th/1802.03545.
[35] B. Badelek et al. The photon collider at TESLA. Int. J. Mod. Phys. A 19, 5097-5186 (2004); hep-ex/0108012.

[36] Ginzburg, I. F., G. L. Kotkin, V. G. Serbo, and V. I. Tel’nov. Production of high-energy colliding $\gamma\gamma$ and $\gamma e$ beams with a high luminosity at VLLEP accelerators. JETP Lett. 34, 491-495 (1981) [Pis’ma Zh. Éksp. Teor. Fiz. 34, 508-510, 1981].

[37] ATLAS Collaboration. Evidence for light-by-light scattering in heavy-ion collisions with the ATLAS detector at the LHC. Nature Physics. 13, 852-858 (2017); hep-ex/1702.01625v2.

[38] G. Aad et al. Observation of light-by-light scattering in ultraperipheral Pb + Pb collisions with the ATLAS detector. Phys. Rev. Lett. 123: 052001 (2019); hep-ex/1904.03536.

[39] D. d’Enterria and G. G. Silveira. Observing light-by-light scattering at the Large Hadron Collider. Phys. Rev. Lett. 111: 080405 (2013) [Erratum: Phys. Rev. Lett. 116: 129901 (2016)]; hep-ph/1305.7142.

[40] G. Jikia and A. Tkabladze. Photon-photon scattering at the photon linear collider. Phys. Lett. B 323, 453-458 (1994); hep-ph/9312228.

[41] G. J. Gounaris, P. I. Porfyriadis, and F.M. Renard. Light by light scattering at high energy: a tool to reveal new particles. Phys. Lett. B 452, 76-82 (1999); hep-ph/9812378.

[42] Z. Bern et al. QCD and QED corrections to light-by-light scattering. JHEP 0111: 031 (2001); hep-ph/0109079.

[43] T. Appelquist and H. D. Politzer. Heavy quarks and $e^+e^-$ annihilation. Phys. Rev. Lett. 34, 43-45 (1975).

[44] R. Barbieri, R. Gatto, and R. Kögerler. Calculation of the annihilation rate of P wave quark-antiquark bound states. Phys. Lett. B 60, 183-188 (1976).

[45] T. Appelquist, A. De Rujula, S. L. Glashow, and H. D. Politzer. Spectroscopy of the new mesons. Phys. Rev. Lett. 34, 365-368 (1975).

[46] A. I. Vainshtein et al. Charmonium and quantum chromodynamics. Sov. Phys. Usp. 20, 796-818 (1977) [Transl. from Usp. Fiz. Nauk 123, 217-255 (1977)].

[47] P. Chen, Y. C. Ong, and D.-h. Yeom. Black hole remnants and the information loss paradox. Phys. Rep. 603, 1-45 (2015); gr-qc/1412.8366.

[48] R. C. Myers and M. J. Perry. Black holes in higher dimensional space-times. Ann. Phys. 172, 304-347 (1986).

[49] L3 Collaboration. B. Acciarri et all. $K^0_S K^0_S$ final state in two-photon collisions and implications for glueballs. Phys. Lett. B 501, 173-182, (2001); hep-ex/0011037.
[50] Belle Collaboration. B. Abe et al. Measurement of $K^+K^-$ production in two-photon collisions in the resonant-mass region. Eur. Phys. J. C 32, 323-336, (2003); hep-ex/0309077.

[51] S. Uehara et al. High-statistics study of $K^0_S$ pair production in two-photon collisions. Prog. Theor. Exp. Phys. 2013, 123C01 (2013); hep-ex/1307.7457.

[52] T. Csörgő et al. Evidence of Odderon-exchange from scaling properties of elastic scattering at TeV energies. Eur. Phys. J. C 81: 180 (2021); hep-ph/1912.11968.

[53] H. Kolanoski. Hadronic final states in soft photon-photon scattering. 5th International Workshop on Photon Photon Collisions. Lect. Notes Phys. 191, 175-205, (1983).

[54] S. Janowski, F. Giacosa, and D. H. Rischke. Is $f_0(1710)$ a glueball? Phys. Rev. D 90, 114005 (2015); hep-ph/1408.4921.

[55] M. Albaladejo and J. A. Oller. Identification of a scalar glueball. Phys. Rev. Lett. 101, 252002 (2008); hep-ph/0801.4929.

[56] L. S. Geng and E. Oset. Vector meson – vector meson interaction in a hidden gauge unitary approach. Phys. Rev. D 79, 074009 (2009); hep-ph/0812.1199.