Diffractive vector meson photoproduction from dual string theory

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(Received 16 September 2008; published 22 April 2009)

We study diffractive vector-meson photoproduction using string theory via AdS/CFT. The large s behavior of the cross sections for the scattering of the vector meson V on a proton is dominated by the soft Pomeron, $\sigma_V \sim s^{2\epsilon - 2\alpha'_P/B}$, where from the string theory model of [arXiv:hep-th/0501039], ϵ is approximately $1/7$ below 10 GeV, and $1/11$ for higher, but still sub-Froissart, energies. This is due to the production of black holes in the dual gravity. In ϕ photoproduction the mesonic Regge poles do not contribute, so that we deal with a pure Pomeron contribution. This allows for an experimental test. At the gauge theory ''Planck scale'' of about 1–2 GeV, the ratios of the soft Pomeron contributions to the photoproduction cross sections of different vector mesons involve not only the obvious quark model factors, but also the Boltzmann factors e^{-4M_V/T_0}, with T_0 the temperature of the dual black hole. The presence of these factors is confirmed in the experimental data for ρ, ω, ϕ, J/ψ, and $\psi(2S)$ photoproduction and is compatible with the meager Υ photoproduction data. Throughout, we use vector-meson dominance, and from the data we obtain T_0 of about 1.3 GeV, i.e. the gauge theory ''Planck scale,'' as expected. The ratio of the experimental soft Pomeron onset scale $\hat{E}_R \sim 9$ GeV and of the gauge theory Planck scale, $T_0 \sim 1.3$ GeV, conforms to the theoretical prediction of $N_c^2/N_c^{1/4}$.

DOI: 10.1103/PhysRevD.79.086012 PACS numbers: 11.25.Tq, 12.40.Nn, 12.40.Vv, 13.60.Le

I. INTRODUCTION

AdS/CFT [1] started as a duality between the conformal $SU(N)$ super-Yang-Mills (SYM) field theory in four dimensions at large N, and string theory in the curved ten-dimensional space, $AdS_5 \times S_5$, in the low energy supergravity limit. The power of this approach derives from the fact that it relates a strong coupling theory—$SU(N)$ SYM at large N and large $g^2_{\rm YM}N$—to a weakly coupled string theory, where one can calculate. In [2] the duality was extended to actual string theory, away from the supergravity limit, but only for the gauge theory sector involving operators with very large R charge.

The duality was extended to various theories with less supersymmetry and/or conformal invariance. Interesting examples are given e.g., by [3,4]. But the difficulty in applying AdS/CFT to real QCD stems not only from the absence of a gravity dual to a nonsupersymmetric gauge theory with light quarks. The problem is that a calculable version with broken supersymmetry, with N_f/N_c fixed, with light quarks has not yet been found, much as one can impose each of these requirements by itself. Moreover, in real QCD $N = 3$ and $g^2_{\rm YM}N$ are both finite. In QCD, $g^2_{\rm YM}$ runs, but $g^2_{\rm YM}N$ never becomes large enough. In AdS/CFT the $1/N$ and $1/(g^2_{\rm YM}N)$ corrections are mapped to string α' and g_s corrections. String α' corrections are world sheet corrections, i.e. two-dimensional field theory quantum corrections, whereas g_s corrections are string spacetime quantum corrections.

String theory is defined only perturbatively, with nonperturbative definitions available only in special cases. Moreover, string theory calculations in the strong coupling regime are hard. Thus, in order to apply AdS/CFT to real QCD, we need to first make sure that string corrections are small in the dual theory. That is in general not true, as finite N and $g^2_{\rm YM}N$ translate into large quantum corrections in string theory.

Recently, a lot of work has been devoted to gravity dual models of QCD for phenomenological purposes, most notably the Sakai-Sugimoto model [5], but in that case we have effectively a quenched approximation, as the dual gravity background contains no backreaction of the D8-brane where the quarks live, thus effectively one works with $N_f/N_c \to 0$. Another popular way of applying gravity dual results to describe RHIC physics [6] uses $\mathcal{N} = 4$ SYM models at finite temperature, corresponding to a static black hole in AdS_5. While in these models one can perform explicit calculations, it is not completely clear why one should be allowed to use $\mathcal{N} = 4$ SYM instead of nonsupersymmetric QCD, even if both are at finite temperature. Moreover, neither the string theory quantum corrections nor the corresponding QCD $1/N$ and $1/(g^2_{\rm YM}N)$ corrections are under control.

However, in [7–10] it was noticed that there exists a particular class of problems for which quantum corrections in the dual theory are small, and one can use AdS/CFT for real QCD as well. The case in point is that of soft high energy scattering. It was shown that the total QCD cross section at large center-of-mass energy squared s can be calculated, and the behavior $\sigma_{\rm tot} \sim s^{\epsilon}$ and its later unitarization to $\ln^2(s/s_0)$ can be derived and matched against experiment. In the gravity dual this corresponds to the creation of black holes, which in the RHIC energy regime can be mapped to the fireball observed in the nucleus-

nucleus collisions [10,11]. Later models for the RHIC fireball introduce a time dependence, like for instance the model in [12], where a dual black hole moves towards the IR, or the model in [13], where the evolution of a black hole horizon is mapped to the cooling of the RHIC plasma, but they generically cannot account for the presence of the mass scale in QCD. It is also not clear whether an evolving horizon makes sense, which is a subject of recent debate.

A natural next step is to consider diffractive vector-meson photoproduction, which is governed by the same soft Pomeron physics, and for which a wealth of experimental data is available. In this paper, we take this step, and find that ϕ photoproduction which is known [14] to provide the cleanest test of soft Pomeron behavior, matches well with the experimental evidence. We also analyze the ratios of the production cross sections of various vector mesons in the soft regime, and find a formula which reproduces the data for all known mesons, and allows us to extract from experiment a ‘‘gauge theory Planck scale’’ of about 1.3 GeV.

In Sec. II we describe the gravity dual picture for the soft physics. In Sec. III we present the general picture of diffractive vector-meson photoproduction and show how to extract the soft Pomeron contribution. In Sec. IV we discuss the gravity dual picture for diffractive vector-meson photoproduction and compare it with experimental data. Finally, in Sec. V we present our conclusions.

II. BLACK HOLE PRODUCTION IN THE GRAVITY DUAL AND THE POMERON

At energies above a gravitational theory's Planck scale, i.e. at $\sqrt{s} > M_{\rm Pl}$, scattering of any two particles should produce black holes. The hard, large t scattering, however, will still be governed by string amplitudes, since a black hole radiates particles thermally, at generically low energies. On the other hand, soft small t scattering, or processes with many particles and small average emitted energy in the final state, will be dominated by black hole creation.

In [15] the process of black hole creation was analyzed using an old idea of ’t Hooft [16], modeling the two colliding particles by Aichelburg-Sexl (AS) gravitational shockwaves [17]:

$$ds^2 = 2dx^+dx^- - \Phi(x^i)\delta(x^+)(dx^+)^2 + d\vec{x}^2. \tag{2.1}$$

According to an argument of ’t Hooft, quantum gravity corrections are negligible because massive interactions have finite range, but at small distances the gravitational shockwave gives a diverging time delay for the interaction $[\Phi(r \to 0) \to \infty]$, making the corrections irrelevant, whereas for the massless interactions, classical gravity reproduces the correct physics. It would be important to obtain a definitive proof of this statement. In [15], the cross section for black hole formation in the higher dimensional curved space scattering was calculated, by scattering two AS waves, based on earlier work [18] in flat four-dimensional space. String corrections were also calculated, by scattering string-corrected AS shockwaves. It was found that in flat four-dimensional space string corrections are exponentially small at energies beyond $E_0 \sim M_P^2/M_s$, which is of the order of the Planck scale if the Planck and string scales are not too far apart.

On the other hand, AdS/CFT [1] relates gauge theories living on systems of D-branes with gravitational theories. In particular, a four-dimensional conformal field theory will have a gravity dual of the type

$$\begin{aligned} ds^2 &= \frac{\bar{r}^2}{R^2}d\vec{x}^2 + \frac{R^2}{r^2}d\bar{r}^2 + R^2 ds_X^2 \\ &= e^{-2y/R}d\vec{x}^2 + dy^2 + R^2 ds_X^2. \end{aligned} \tag{2.2}$$

Here the first two terms give the line element of the AdS_5 space of size R, and the last term gives the line element of the S^5 of radius R.

A nonconformal theory will have a gravity dual modified in the IR, i.e. small $\bar{r}$ or large y. In the simplest model of this type, one describes the unknown modification by an ‘‘IR brane,’’ a cutoff at $\bar{r}_{\rm min} \sim R^2\Lambda_{\rm QCD}$ with $\Lambda_{\rm QCD}$ the lightest excitation of the theory [19].

This simple model is dual to a pure glue gauge theory. To model the pion, the lightest excitation of real QCD, which is a $q\bar{q}$ state, one assumes the position of the cutoff to be a dynamical brane. Its fluctuation, the radion, corresponds to a simple scalar ‘‘pion’’ [20].

High energy scattering in the gauge theory of modes with momentum p and wave function e^{ipx} was mapped in [19] to scattering in the gravity dual with local AdS momentum $\tilde{p}_\mu = (R/\bar{r})p_\mu$ and wave function $e^{ipx}\psi(r, \Omega)$. The gauge amplitude is given by the gravitational amplitude, integrated over the extra coordinates, and convoluted with the wave functions. The string tension $\alpha' = R^2/(g_sN)^{1/2}$ corresponds to the gauge theory string tension $\hat{\alpha}' = \Lambda_{\rm QCD}^{-2}/(g_{\rm YM}^2N)^{1/2}$, and

$$\sqrt{\alpha'}\tilde{p}_{\rm string} \leq \sqrt{\hat{\alpha}'}p_{\rm QCD}. \tag{2.3}$$

This formalism was applied in [7,9] to the case of soft scattering with black hole formation. It was found that most of the integral over the coordinate $\bar{r}$ is supported in the IR region of small $\bar{r}$, close to the cutoff. For self-consistency though, it still has to be supported away from $\bar{r}_{\rm min}$. The cross section for black hole creation in the gravity dual was calculated in the classical shockwave scattering picture mentioned before, and then using a simple eikonal model, an elastic $2 \to 2$ quantum amplitude was derived, and used in the Polchinski-Strassler formalism [19]. Let us emphasize again that in the soft scattering regime, the black holes are produced on the average close to the IR cutoff, i.e., to the IR brane, but still away from it.

In the gravity dual, there are three energy scales of interest: the Planck scale M_P, the scale $E_R = N^2R^{-1}$ at which the produced black holes reach the AdS size, and the scale E_F at which the black holes are big enough to reach the IR brane. This last scale clearly cannot be calculated without knowing the details of the IR modification of the gravity dual to a given gauge theory. The simple cutoff model cannot be used to calculate it. According to Eq. (2.3), in gauge theory these energy scales correspond to $\hat{M}_P = \Lambda_{\rm QCD} N^{1/4}$, $\hat{E}_R = \Lambda_{\rm QCD} N^2$, and an unknown $\hat{E}_F$. It was found [7,9,20] that between $\hat{M}_P$ and $\hat{E}_R$, when in the gravity dual one produces black holes small enough to be considered approximately in flat space, the gauge theory cross section grows like

$$\sigma_{\rm gauge,tot} \simeq K(s/s_0)^\epsilon = K(s/\hat{M}_P^2)^{1/7}. \tag{2.4}$$

where K is a constant, as are $\bar{K}$ and K' in the next two equations. Between the energies $\hat{E}_R$ and $\hat{E}_F$, in the gravity dual the black holes are large enough to feel the AdS size, but not the size of X_5. Therefore they grow in $\text{AdS}_5 \times X_5$ with X_5 large, giving the gauge theory cross section [9]

$$\sigma_{\rm gauge,tot} \simeq \bar{K}(s/\bar{s}_0)^\epsilon = \bar{K}(s/\hat{E}_R^2)^{1/11}. \tag{2.5}$$

Above $\hat{E}_F$, in the gravity dual the black holes become so large that they reach the IR brane. The gauge theory cross section then saturates the Froissart bound [7,8,20–22]:

$$\sigma_{\rm gauge,tot} \simeq K' \frac{\pi}{M_1^2} \ln^2 \frac{s}{s_1}, \tag{2.6}$$

where the mass M_1 corresponds to the lightest state in the theory. In the simple cutoff model, this is the mass gap, or the lightest glueball in the gauge theory. In QCD, where the pion is a pseudo-Nambu-Goldstone boson, M_1 is replaced by the pion mass m_π. As mentioned, this Nambu-Goldstone boson is modeled by the radion, or equivalently, by the position of the IR brane in the gravity dual. There is then also a different Froissart onset scale $E'_F \neq E_F$, and therefore $\hat{E}'_F \neq \hat{E}_F$ in the gauge theory. This corresponds to the scale at which brane bending reaches the black hole.

As mentioned in the Introduction, for this analysis to apply at finite N and finite $g^2_{\rm YM}N$, as in real QCD, we have to make sure that string corrections are small. They are not small for hard scattering, but for soft scattering at t fixed and $s \to \infty$, these string corrections are small above an energy scale $\sim M_P^2/M_s$ for flat four-dimensional space. If 't Hooft's argument is valid, this should be actually around M_P. Making use of the optical theorem $\sigma_{\rm total}(s) = \text{Im}A(s, t = 0)/s$, the same analysis holds for the total cross section.

Note that string corrections to the background itself will be large, and this will translate into modifications of the various energy scales, in particular, of $\Lambda_{\rm QCD}$, the dual of the AdS size R^{-1}. *A priori*, this also entails modifications of the ratios $\hat{M}_P/\Lambda_{\rm QCD}$, $\hat{E}_R/\Lambda_{\rm QCD}$, and of $\hat{E}_F$.

By contrast, string corrections to soft scattering in a given background *are* small, meaning that we can trust the cross-section calculations. This allowed for a successful match [9] of the "soft Pomeron" exponent $\sigma_{\rm tot} \sim s^{1/11} \simeq s^{0.0909}$, expected to set in at $N_c^2 M_{1,\rm glueball} \sim 10$ GeV with the energy dependence $\sigma_{\rm tot} \sim s^{0.0933\pm0.0024}$, observed experimentally as of 9 GeV [23,24].

III. DIFFRACTIVE VECTOR-MESON PHOTOPRODUCTION (DVMP)

A. Vector-meson dominance model of vector-meson photoproduction

Before applying AdS/CFT ideas to diffractive vector-meson photoproduction processes, let us briefly recall how the "soft Pomeron" dominates them. A particularly simple and intuitive picture [25] for the photoproduction $\gamma p \to Vp$ of the vector meson V is obtained in the vector-meson dominance (VMD) approximation. One treats the photon γ as undergoing a transition to the a virtual vector meson V which then scatters elastically on the target proton. This involves the corresponding $\gamma \to V$ transition matrix element

$$\langle 0 | j_\mu(0) | V_a \rangle = f_V m_V^2 \eta_{\mu a} \tag{3.1}$$

(with $\eta_{\mu a}$ the Minkowski metric), as well as a V propagator evaluated on the photon's mass shell. This propagator then cancels the factor m_V^2, used for notational simplicity in the definition of f_V in Eq. (3.1) for the transition matrix element. Also for notational simplicity this definition already includes the electromagnetic coupling constant. This way the amplitude $A_{\gamma p \to Vp}(s, t)$, with s and t the usual Mandelstam variables, is given in terms of the Vp elastic scattering amplitude $A_{Vp}(s, t)$ as

$$A_{\gamma p \to Vp}(s, t) = f_V A_{Vp}(s, t). \tag{3.2}$$

With VMD, understanding diffractive vector-meson photoproduction reduces to understanding diffractive vector-meson-proton elastic scattering and this calls for Regge theory, and specifically for Pomeron dominance.

B. Regge theory for Vp scattering

Consider the amplitude $A_{Vp}(s, t)$ for $Vp \to Vp$ scattering. In gauge theories, it was found that one can generically describe the various Regge limits in terms of a tree level theory of effective particle ("Reggeon") interactions. In the case of QCD, the leading trajectory is the Pomeron. It is an effective particle made up of gluon interactions, and in the case of soft scattering, we have the "soft Pomeron." It was found that several effective particles can account for several regimes. For $t \leq 0$ fixed and $s \to \infty$ Regge theory gives

$$A_{Vp}(s, t) \sim \beta_V(t) \text{sig}_P(t) s^{\alpha_P(t)}, \tag{3.3}$$

where

$$\alpha_P(t) = \alpha_P(0) + \alpha'_P t + \cdots \tag{3.4}$$

is the Pomeron trajectory, and

$$\mathrm{sig}_P(t) = \frac{-1 - e^{-i\pi\alpha_P(t)}}{\sin\pi\alpha_P(t)} \tag{3.5}$$

its signature factor.[1] The Pomeron Regge residue in Vp scattering $\beta_V(t)$ is real for $t \le 0$. If $\alpha_P(0) \simeq 1$, as can be seen from total cross-section data even at energies below the onset of the Froissart regime, then $\mathrm{sig}_P(t) \simeq i$.

The total cross section for Vp scattering, $\sigma_{Vp,\mathrm{total}}(s)$ is given by the optical theorem,

$$\sigma_{Vp,\mathrm{total}}(s) = \frac{\mathrm{Im}A_{Vp}(s, t=0)}{s} \simeq \beta_V(0)s^{\alpha_P(0)-1} \equiv \beta_V(0)s^{\epsilon}. \tag{3.6}$$

Therefore $\epsilon = \alpha_P(0) - 1$, and from the gravity dual description in Sec. II, we should have $\epsilon \simeq 1/7$ below 10 GeV, and $\simeq 1/11$ between 10 GeV and $\hat{E}_F$, the onset of the Froissart saturation behavior.

On the other hand, the elastic differential cross section for $Vp \rightarrow Vp$ scattering is

$$\frac{d\sigma_{Vp}(s,t)}{dt} = \frac{1}{16\pi s^2}|A_{Vp}(s,t)|^2 \simeq \frac{1}{16\pi}|\beta_V(t)\mathrm{sig}_P(t)|^2 s^{2(\alpha_P(t)-1)} \equiv F_V(t)s^{2\epsilon+2\alpha'_P t+\cdots}. \tag{3.7}$$

For small t, such that $-1\ \mathrm{GeV}^2 < t \le 0$, the prefactor is well approximated by an exponential in t,

$$F_V(t) = \frac{1}{16\pi}|\beta_V(t)\mathrm{sig}_P(t)|^2 \simeq F_V(0)e^{Bt}. \tag{3.8}$$

At large $|t|$ it becomes a power law.

By integrating the differential cross section over t, we obtain $\sigma_{Vp,\mathrm{elastic}}(s)$, the total elastic cross section. Because of the integrand's rapid exponential falloff, one can extend the integral to $-\infty$. Thus,

$$\sigma_{Vp,\mathrm{elastic}}(s) = \int_{-\infty}^{0} dt \frac{d\sigma_{Vp}(s,t)}{dt} \simeq \int_{-\infty}^{0} dt F_V(t)s^{2\epsilon+2\alpha'_P t+\cdots} \simeq \frac{F_V(0)}{B}s^{2\epsilon}\left(1 + 2\frac{\alpha'_P}{B}\ln s\right)^{-1}. \tag{3.9}$$

On the other hand, $\alpha'_P/B < \epsilon \ll 1$, since roughly $\alpha'_P \sim 0.2\ \mathrm{GeV}^{-2}$ and $B \sim 2$–$4\ \mathrm{GeV}^{-2}$), which means that we can bring the corresponding term into the exponent, and obtain

$$\sigma_{Vp,\mathrm{elastic}}(s) \simeq \frac{F_V(0)}{B}s^{2\epsilon-2\alpha'_P/B} \equiv \frac{F_V(0)}{B}s^{2(\alpha_P(\langle t\rangle)-1)}, \tag{3.10}$$

where by definition $\langle t\rangle = -1/B$ is the average t of the integral.

As we shall see, the experimentally observed value of $\langle t\rangle$ is close to $-\epsilon/(2\alpha'_P)$.

C. Regge theory of DVMP

We can now use vector-meson dominance (VMD) to relate the analysis of the previous subsection to diffractive vector-meson photoproduction (DVMP).

Combining Eqs. (3.2) and (3.7), the differential photoproduction cross section is

$$\frac{d\sigma_{\gamma p\rightarrow Vp}(s,t)}{dt} = \frac{1}{16\pi s^2}|\mathcal{A}_{\gamma p\rightarrow Vp}(s,t)|^2 = |f_V|^2\frac{d\sigma_{Vp}(s,t)}{dt} = |f_V|^2 F_V(t)s^{2\epsilon+2\alpha'_P t+\cdots}. \tag{3.11}$$

Integrating over t we get the total V meson photoproduction cross section,

$$\sigma_{\gamma p\rightarrow Vp}(s) \equiv \int_{-\infty}^{0} dt \frac{d}{dt}\sigma_{\gamma p\rightarrow Vp}(s,t) \simeq |f_V|^2\sigma_{Vp,\mathrm{elastic}}(s) = |f_V|^2\frac{F_V(0)}{B}s^{2(\alpha_P(\langle t\rangle)-1)}x. \tag{3.12}$$

On the other hand, the Compton scattering amplitude $A_{\gamma p}(s,t)$ can be obtained by a double application of VMD. We can relate the elastic process $\gamma p \rightarrow \gamma p$ with the elastic process $Vp \rightarrow Vp$ by separating a transition element on both the incoming and the outgoing γ, and summing over all vector mesons V, with the result

$$A_{\gamma p}(s,t) \simeq \sum_V f_V A_{Vp}(s,t)f_V^*, \tag{3.13}$$

where we have discarded the much smaller off-diagonal $Vp \rightarrow V'p$ terms. Making use at this point of the optical theorem gives

$$\sigma_{\gamma p,\mathrm{total}}(s) = \frac{\mathrm{Im}A_{\gamma p}(s,t=0)}{s} \simeq \sum_V |f_V|^2\frac{\mathrm{Im}A_{Vp}(s,t=0)}{s} = \sum_V |f_V|^2\sigma_{Vp,\mathrm{total}}(s), \tag{3.14}$$

so that

$$\sigma_{\gamma p,\mathrm{total}}(s) \simeq \left(\sum_V |f_V|^2\beta_V(0)\right)s^{\epsilon}. \tag{3.15}$$

In (3.12), the value of $\langle t\rangle = -1/B$ is the same for all light V mesons, but for each heavy meson both $\langle t\rangle$ and $\epsilon = \alpha_P(0) - 1$ take different values, and give a different exponent.

[1] The Pomeron has even signature. Other Regge trajectories can have odd signatures.

In order to emphasize the difference between the Pomeron parameters for heavy and for light mesons, we write $\alpha_{P,V}$ and $\langle t\rangle_V$ in (3.12), which thus becomes

$$\sigma_{\gamma p\to Vp}(s) \equiv |f_V|^2 \frac{F_V(0)}{B} s^{2(\alpha_{P,V}(\langle t\rangle_V)-1)}. \quad (3.16)$$

Allowing for such a flavor dependence in the Pomeron, heretic though it may seem from the point of view of Regge theory, has the virtue of agreeing with what is experimentally observed. Earlier work has introduced a flavor dependence of the Pomeron residues [26,27], from the point of Regge theory both acceptable and required, but what we are doing here amounts to allowing a flavor dependence of what seems to be the very position of the Pomeron singularity in the complex angular momentum plane. Among other effects, this type of ''Pomeron flavoring'' would destroy the factorization properties of the Pomeron. This Pomeron flavoring can be understood by taking note of the different kinematic regimes for light and heavy vector-meson photoproduction. Because of the large masses of the J/ψ, $\psi(2S)$, and Υ, large momentum transfers are set into play and one moves away from the diffractive soft Pomeron peak relevant for the photoproduction of the light vector mesons, into a region where the hard Pomeron takes over and where such a flavoring is not ruled out.[2]

We now normalize the s behavior to a fixed s_0, and write

$$\sigma_{\gamma p\to Vp}(s) \simeq |f_V|^2 \frac{F_V(0)}{B} s^{2(\alpha_{P,V}(<t>_V)-1)} \equiv |f_V|^2 C_V (s/s_0)^{\tilde{\epsilon}_V}, \quad (3.17)$$

where C_V depends on the choice of s_0 and the exponent of s is called $\tilde{\epsilon}_V$. For the comparison of cross sections to have predictive power, the value s_0 must be singled out by the theory. In the gravity dual calculation in the next section, s_0 will indeed be set at $s_0 \sim \hat{M}_P^2$, and we will predict ratios of C_V for this value of s_0.

IV. GRAVITY DUAL PICTURE AND MATCHING WITH EXPERIMENT

A. Theory

Let us now look at vector-meson photoproduction from the dual point of view. As we saw, for all practical purposes, we are dealing with large s Vp scattering, a soft Pomeron process. This means that it should be related to black hole production. The total QCD cross section is due to black hole production.

We are thus led to consider the process $p\bar{p} \to V\bar{V}$ around the gauge theory Planck scale. As we argued in Sec. II, above the gauge theory Planck scale this process should be governed by black hole production in the gravity dual. Specifically, a mini-black hole should be produced near the IR brane, i.e. at an average position r_{av} away from the IR brane r_{min}, but close to it. Then this black hole decays into particles.

At this scale, one might wonder whether it is possible to even speak of a black hole. Yet in string theory there are good reasons to believe that the production and decay of a black hole is a quantum process, and the apparent black hole loss of information due to the black hole temperature is just an illusion, i.e. the temperature is not indicative of information loss. How this can happen in the case of a near extremal black hole has been discussed in [29]. Furthermore in [30] for a simple scalar field theory toy model for the pion, a solution was found which mimics the properties of the gravity dual black hole, with an effective temperature and with *apparent* information loss, even though the quantum field theory from which one started is unitary. Thus, although a pure quantum process takes place, the effect of the black hole being produced is simply to give an effective temperature.

Therefore since the energy is small enough for the decay to produce on the average just two particles—the least number needed for momentum conservation—the simplest decay mode of the Planck-sized black hole is into a particle-antiparticle pair. The decay of the dual black hole is mapped in QCD to a gluonic interaction followed by decay into vector mesons. Thus, the dual black hole's only observable effect in the QCD $p\bar{p}$ scattering is the existence of well-defined relative probabilities of decay for various vector-meson pairs, so that

$$\frac{\sigma_{V\bar{V}}(t)}{\sigma_{V'\bar{V}'}(t)} \sim \frac{P_V}{P_{V'}} = e^{-(2(M_V - M_{V'})/T_0)}, \quad (4.1)$$

where t is the Mandelstam variable for the $p\bar{p} \to V\bar{V}$ ''t channel,'' M_V and $M_{V'}$ are the vector-meson masses, T_0 is the temperature of the Planck-sized (M_P) black hole, and the factor 2 in the exponent corresponds to the meson pair.[3]

[2]Note that hard scattering physics was the first to be understood in AdS/CFT using the toy model for QCD used in [19], of AdS space with an IR cutoff or $r_{min} = R^2\Lambda_{QCD}$. However, having both hard and soft scattering in the same AdS/CFT description is also very difficult. Such a description was attempted, for instance, in [28], where hard scattering was connected to a Balitsky-Fadin-Kuraev-Lipatov (BFKL) evolution equation. We will therefore not give an explanation for the flavor dependence of the Pomeron within our model, but we will leave it as an experimental fact.

[3]Of course, any kind of particles can be produced in the $p\bar{p}$ collision, but we will be focusing on the case of production of a meson pair, as it can be related by crossing to $Vp \to Vp$. Within AdS/CFT, it is clear that the black hole creation and decay is actually unitary, so the $V\bar{V}$ creation should be described by a well-defined quantum amplitude $A_{p\bar{p}\to V\bar{V}}$. Alternatively, one can think of the production of quark-antiquark pairs in the t channel. This way one can account for the observed equality of the πN and say ρN diffractive amplitudes, because they involve the same quarks. In terms of the black hole Boltzmann factors, this would have the sole effect of renormalizing the temperature by a factor 1/2, given that the constituent quarks weigh in at essentially half the corresponding vector-meson masses. Otherwise all arguments would remain unchanged.

Since the black hole is minimal (of Planck size), it cannot evolve or cool off, like in the alternative models for the RHIC fireball in [12,13], where the dual black hole is large, but rather it has to quickly decay into hadrons.

In QCD this involves a highly nonperturbative process, so that the effect of the dual black hole would in principle be reproduced by very complicated QCD interactions.

QCD being a quantum theory, the process is unitary, as we said, and therefore the effect on the quantum amplitude for $p\bar{p} \to V\bar{V}$ is that it factorizes into a universal piece $a(s,t)$ independent of V, and a V-dependent Boltzmann factor,

$$A_{p\bar{p}\to V\bar{V}}(s,t) \sim e^{-(2M_V/T_0)}a(s,t). \quad (4.2)$$

Let us now look at this process in the crossed $Vp \to Vp$ s channel. If this is in the soft regime, due to gluon interactions, or equivalently "Pomeron exchange," then the above picture should apply, and we have

$$A_{Vp}(s,t) \equiv A_{Vp\to Vp}(s,t) \sim e^{-(2M_V/T_0)}a(s,t), \quad (4.3)$$

where $a(s,t)$ is a universal $2 \to 2$ amplitude in the Regge regime. We can now use this relation in (3.3) and (3.8) to obtain $\beta_V(t) = e^{-(2M_V/T_0)}\beta(t)$ and $F_V(t) = e^{-(4M_V/T_0)}F(t)$.

Then, since the process we are interested in is actually $\gamma p \to Vp$, we can use VMD as in Secs. III A and III B, obtaining the amplitude

$$\mathcal{A}_{\gamma p\to Vp}(s,t) \sim f_V e^{-(2M_V/T_0)}a(s,t) \quad (4.4)$$

and the V meson photoproduction cross section

$$\sigma_{\gamma p\to Vp}(s) \simeq |f_V|^2 e^{-(4M_V/T_0)}C(s/s_0)^{2(\alpha_{P,V}(\langle t\rangle_V)-1)}$$
$$= |f_V|^2 e^{-(4M_V/T_0)}C(s/s_0)^{2\epsilon_V - 2(\alpha'_P/B_V)} \quad (4.5)$$

or $C_V = Ce^{-(4M_V/T_0)}$ in (3.17). This is our main result, which we will confront with experiment.

First, however, let us go back and reexamine the assumption of the black hole whose temperature T_0 is of Planck size, and also understand the scale s_0 at which the comparison of ratios is to be made.

Why are Planck-sized black holes with $T_0 \sim M_P$ relevant? After all, we argued that for the total QCD cross section in the gravity dual one has black holes of growing size, thus decreasing temperature, that stops at $T_{\rm min} = a4M_1/\pi$ (with a a suitable numerical factor) corresponding in QCD to $T = a4m_\pi/\pi$ [10]. The growing size of the produced black holes is responsible for the s^ϵ behavior of $\sigma_{\rm tot}(s)$.

But for the temperature factor we have gone from the $p\bar{p} \to V\bar{V}$ amplitude to the s channel, exchanging the original t with s. The original t had to be $\geq M_P$ to create black holes, but for $pV \to pV$ we need $|t| \leq M_P$, thus $t \simeq M_P$ for $p\bar{p} \to V\bar{V}$ and so indeed $T_0 \sim M_P$. A similar argument can be made exchanging s with t, yielding a Planck scale value $s_0 \sim \hat{M}_P^2$ of the Mandelstam variable s at which the soft Pomeron behavior sets in for the $pV \to pV$ s channel.

Thus, for V photoproduction at $\sqrt{s} = \sqrt{s_0} \sim \hat{M}_P$, the soft Pomeron contribution, dual to black hole production, should be distributed according to temperature $T_0 \sim \hat{M}_P$. Beyond that energy, one should have the $s^{\tilde{\epsilon}_V}$ behavior. For light vector mesons, with mass $M_V < \hat{M}_P$, $\tilde{\epsilon}_V$ should be the soft Pomeron exponent $2\epsilon - 2\alpha'_P/B$. As we saw, this is numerically close to ϵ. On the other hand, for heavy vector mesons, with $M_V > \hat{M}_P$, $\tilde{\epsilon}_V$ should be an exponent defined by hard scattering. The reason for that is purely kinematic, as it is well known: the photon has $Q^2 = 0$, whereas the vector meson has $P^2 = M_V^2$, as if M_V^2 had been effectively added to it. Thus, if $M_V > \hat{M}_P$, the physics is no longer dominated by the lightest glueball.

Also, for kinematic reasons, production of a heavy meson V with $M_V > \hat{M}_P$ will start above $\sqrt{s} = M_V V$ instead of at $\sqrt{s_0} \sim \hat{M}_P$, and as we said will be dominated by hard scattering. We expect that if we extrapolate the hard scattering formula $\sigma_{V,\rm hard} \sim s^{\tilde{\epsilon}_{\rm hard}}$ down to $\sqrt{s_0} = \hat{M}_P$, the soft formula should apply. In other words, the formula has $\sqrt{s_0} \sim \hat{M}_P$ for heavy mesons also, even though it only applies for $s > M_V^2$.

In conclusion, the theoretical expectation for the Pomeron contribution to σ_V is as sketched in Fig. 1. For light mesons $M_V \leq \hat{M}_P$, we have the soft Pomeron exponent $\tilde{\epsilon} \equiv 2\epsilon - 2\alpha'_P/B$, which as we saw is numerically close to ϵ. For heavy mesons $M_V \geq \hat{M}_P$, we have the hard Pomeron exponent, but both the soft and hard Pomeron lines extrapolated down to about $\hat{M}_P$ should be distributed according to (4.5).

B. Comparison with experiment

Let us now compare the main formula (4.5) to the experimental data, and extract the experimental value of T_0. For all light mesons, the fits show the same soft Pomeron exponent, so that we can compare the ratios of cross sections at any s, in particular large values of s, where neglecting Regge exchanges beyond the Pomeron is a good approximation. By contrast, different heavy mesons yield different Pomeron "trajectories," so that the powers of s differ and these different exponents are not predicted by theory. We will therefore extrapolate the data to the energy $\sqrt{s_0} \sim \hat{M}_P \sim 1$–$2$ GeV, even though, as we already said, for heavy mesons the formula itself starts applying only at a higher energy. We will then work out the relevant cross section ratios from the so extrapolated data.

Our procedure will amount to a simple graphical analysis taking advantage of best fit lines to the high energy data. A more extensive data analysis is not warranted at this point, since there are large theoretical uncertainties. For instance, what is the precise value of s_0 at which we make the comparison, and what kind of corrections do we expect to the black hole exponential factor e^{-4M_V/T_0}? Given these

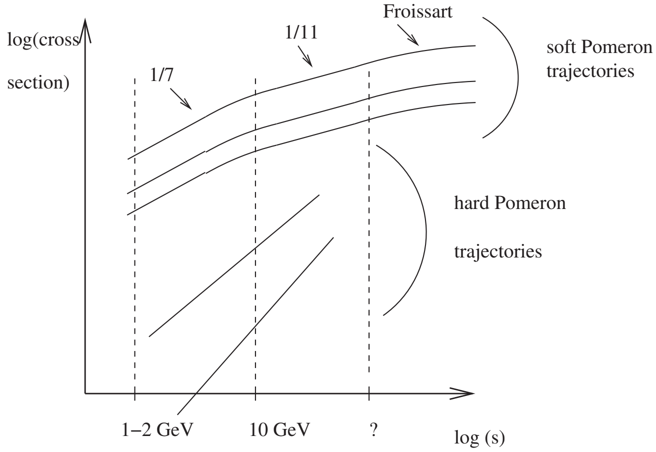

FIG. 1. Predicted soft Pomeron contribution to the scattering cross section (due to production of dual black holes). The first energy scale is the gauge theory Planck scale, $\hat{M}_P = N_c^{1/4} M_{1,\text{glueball}} \sim 1$–$2$ GeV. The second is $\hat{E}_R = N_c^2 M_{1,\text{glueball}} \sim 10$ GeV, and the third is the *a priori* unknown Froissart scale, $\hat{E}_F$, above which $\sigma \sim \log^2(s)$, thus $\log(\sigma) \sim 2\log(\log(s))$. For vector-meson photoproduction, the exponent is $\tilde{\epsilon} = 2\epsilon - 2\alpha'_P/B$, which as was already mentioned is seen to be close to the numerical value of ϵ (the equality is true numerically, though not clear if true theoretically). For vector mesons with masses above $M_{1,\text{glueball}}$, we have hard scattering, with a larger exponent. One could have *a priori* also a scale at which the exponent changes, as for the soft Pomeron exponent. At $\hat{M}_P$, the various mesons should be distributed according to the corresponding Boltzmann factors.

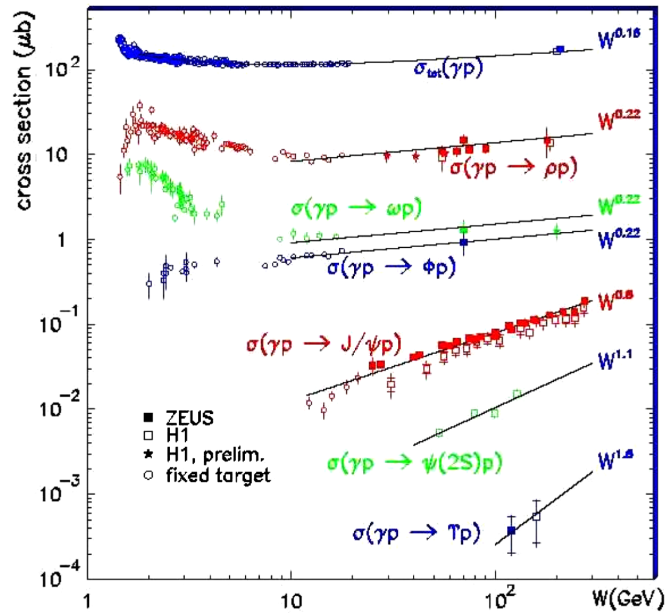

FIG. 2 (color online). Energy (W) dependence of the exclusive photoproduction of light and heavy vector mesons, Fig. 2 of [31].

uncertainties, it makes little sense to try a more sophisticated statistical analysis of the data at this point.

The graph that we will be using for our analysis is Fig. 2 of Ref. [31] (see also for instance, the review [32]), reproduced here for convenience in Fig. 2. The values of the exponents for the best fit lines at high energies are $\tilde{\epsilon}_\rho = \tilde{\epsilon}_\omega = \tilde{\epsilon}_\phi = 0.11$, $\tilde{\epsilon}_{J/\Psi} = 0.41$, $\tilde{\epsilon}_{\psi(2S)} = 0.55$, and $\tilde{\epsilon}_\Upsilon = 0.9$.

Among the light mesons, ϕ photoproduction is the cleanest test of the soft Pomeron exponent, because in this case all mesonic Regge poles decouple [14]. We therefore identify the cross-section curve for ϕ photoproduction with the soft Pomeron contribution. For the other light vector mesons, we look at data at higher energies, $\sqrt{s} \gg 10$ GeV, where the Pomeron contribution dominates. We then extrapolate down to $\hat{M}_P$, parallel with the ϕ line. The next cleanest line, also well split from ϕ, is the J/ψ line, where also all mesonic Regge poles decouple. We therefore use the J/ψ:ϕ split to determine T_0 and then with the so-obtained value of T_0 check the formula for the vector mesons for which the photoproduction cross-section measurements carry larger errors. We also extrapolate the heavy meson lines down to $s_0 \sim \hat{M}_P^2 \sim 1$–$2$ GeV^2. For concreteness, we at first choose $s_0 = 2$ GeV^2.

Numerically, as mentioned already, the energy dependence of total cross section is close to that of the light meson photoproduction cross sections. This way for light mesons $\tilde{\epsilon} \simeq \epsilon \simeq 0.11$, even though all one had a right to expect was $\tilde{\epsilon} = 2\epsilon - 2\alpha'_P/B$. In some sense it is as if we were testing the behavior of $\sigma_V \sim s^\epsilon$. The fact that in the photoproduction cross sections the switch from one power law to the other occurs around 9 GeV is particularly significant in that it does not depend on whether $\tilde{\epsilon} \simeq \epsilon$ is true or not.

The ϕ photoproduction cross-section data have the general features that we expect. Most of the data are around 10 GeV, which from previous work [9] corresponds to the scale $\hat{E}_R$ at which in the dual gravity picture the black hole size reaches the AdS size. This is the transition region between the $\epsilon = \frac{1}{7}$ and $\epsilon = \frac{1}{11}$. Consequently, one measures mostly the energy dependence at 10 GeV, which should be of the form s^ϵ with ϵ between $\frac{1}{7}$ and $\frac{1}{11}$, and indeed one finds $\epsilon = 0.11 \simeq (\frac{1}{7} + \frac{1}{11})/2$. Moreover, from the graph one can clearly see that ϵ is larger below 10 GeV. In fact one nicely fits the data with $\epsilon = 0.15$, which is close to $\frac{1}{7} \simeq 0.143$.

Encouraged by this fit, let us turn to the formula at $s = s_0$. The masses of the relevant mesons are

$$\begin{aligned} M_{\rho^0} &= 775.8 \pm 0.5 \text{ MeV}; \\ M_\omega &= 782.59 \pm 0.11 \text{ MeV}; \\ M_\phi &= 1019.456 \pm 0.020 \text{ MeV}; \\ M_{J/\Psi} &= 3096.916 \pm 0.011 \text{ MeV}; \\ M_{\psi(2S)} &= 3686.093 \pm 0.034 \text{ MeV}; \\ M_\Upsilon &= 9460.30 \pm 0.26 \text{ MeV}. \end{aligned} \tag{4.6}$$

For mesons that are close by in mass and wave functions, the relevant VMD factor f_V^2 should be dominated by the group theory factor $H_V \equiv (TrQ\mathcal{V})^2$, where in (u, d, s, c, b) space, $Q = \mathrm{diag}(2/3, -1/3, -1/3, 2/3, -1/3)$ is the electric charge matrix of the quarks, and $\mathcal{V}$ is the quark content matrix of the vector-meson. Since we have approximately

$$\rho^0 \simeq \frac{u\bar{u} - d\bar{d}}{\sqrt{2}}, \qquad \omega \simeq \frac{u\bar{u} + d\bar{d}}{\sqrt{2}}, \qquad \phi \simeq s\bar{s} J/\psi \simeq c\bar{c}, \qquad \psi(2S) \simeq c\bar{c}; \qquad \Upsilon \simeq b\bar{b}, \tag{4.7}$$

it follows that

$$\mathcal{V}_{\rho^0} \simeq \frac{1}{\sqrt{2}}\tau_{3,(u,d)}, \qquad \mathcal{V}_\omega = \frac{1}{\sqrt{2}}1_{(u,d)}, \qquad \mathcal{V}_\phi = 1_s, \qquad \mathcal{V}_{J/\psi} = \mathcal{V}_{\psi(2S)} = 1_c, \qquad \mathcal{V}_\Upsilon = 1_b \tag{4.8}$$

and so

$$H_{\rho^0}:H_\omega:H_\phi:H_{J/\psi}:H_{\psi(2S)}:H_\Upsilon = 9:1:2:8:8:2. \tag{4.9}$$

The ρ and ω mesons being almost mass degenerate, the ratio of their photoproduction cross sections at high s should be 9:1, as has been known for a long time, see e.g. [25], and can also be read off the experimental data.

The ϕ:J/Ψ split.—If we extrapolate the J/ψ photoproduction cross section down to about 2 GeV, we find $\frac{\sigma_\phi}{\sigma_{J/\psi}} \sim 100$, which when compared with Eqs. (4.5) and (4.9) sets the temperature T_0 at 1.3 GeV.

The ω:ϕ split.—With the temperature determined this way, we predict

$$\frac{\sigma_\omega}{\sigma_\phi} = \frac{1}{2}e^{(4(M_\phi - M_\omega)/T_0)} \sim 1.03 \tag{4.10}$$

in agreement with experiment.

Here it is worth mentioning that the meson-dominated Pomeron picture [26,27], calls for a suppression of heavier vector-meson photoproduction cross sections inverse proportional to the fourth power of the vector-meson mass. This replaces our Boltzmann factors. For the ϕ versus ρ photoproduction, this means that the Boltzmann factor $e^{(4(M_\phi - M_\rho)/T_0)} \sim 2$ is to be replaced by $\frac{M_\phi^4}{M_\rho^4} \sim 3$. Including the group theory factors, this yields $\frac{\sigma_\rho}{\sigma_\phi} \sim 13.5$, as opposed to our result of 9. At high energies, where the soft Pomeron dominates, both pictures are compatible with experiment, within the stated errors.

The J/ψ:$\psi(2S)$ split.—Here J/ψ and $\psi(2S)$ are $n = 1$, and $n = 2$ $c\bar{c}$ S states, respectively. Yet the corresponding $f_V M_V^2$ are not the same, as the J/ψ and $\psi(2S)$ wave functions differ. The $n = 2$ state will localize the quarks further from the origin, for instance. In order to extract the matrix element ratio, we use VMD, and start from the decay rates of J/ψ:$\psi(2S)$ into a lepton pair. In fact,

$$\frac{\Gamma_{J/\psi}(e^+e^-)}{\Gamma_{\psi(2S)}(e^+e^-)} = \frac{|f_{J/\psi}|^2}{|f_{\psi(2S)}|^2}\frac{M_{J/\psi}}{M_{\psi(2S)}}, \tag{4.11}$$

where the ratio of masses on the right-hand side comes from the phase space factors of these S-wave decays. Experimentally this ratio of widths is $5.37\ \mathrm{keV}/2.10\ \mathrm{keV} = 2.55$. This way

$$\frac{\sigma_{J/\psi}}{\sigma_{\psi(2S)}}\Bigg|_{\sqrt{s}=\sqrt{s_0}\sim\hat{M}_P\sim 2\ \mathrm{GeV}} = \frac{|f_{J/\psi}|^2}{|f_{\psi(2S)}|^2}e^{(4(M_{\psi(2S)} - M_{J/\psi})/T_0)} = 10^{1.27}. \tag{4.12}$$

The experimental graph gives a split of about $10^{1.25}$.

The J/ψ:Υ split.—In this case the group theory and Boltzmann factors give

$$\frac{\sigma_{J/\psi}}{\sigma_\Upsilon}\Bigg|_{\sqrt{s}=\sqrt{s_0}\sim 2\ \mathrm{GeV}} = \frac{8}{2}e^{(4(M_\Upsilon - M_{J/\psi})/T_0)} \simeq 10^9, \tag{4.13}$$

i.e. nine decades difference. The experimental graph, blindly extrapolated to 2 GeV, yields a difference of about 4.5 decades. But the Υ "line" consists of just two close-by points with huge error bars, so that its slope could easily reach double the value suggested by the best fit. The meson-dominated *soft* Pomeron picture is hardly applicable here, given the large momentum transfers kinematically required on account of the large Υ mass. This calls for the hard Pomeron in Υ photoproduction and eliminates the fourth-power mass dependence characteristic of the meson-dominated *soft* Pomeron picture, which though valid for ρ photoproduction, fails for Υ photoproduction.

A better determination of the Υ line through more accurate Υ photoproduction data will allow for an important additional test of these ideas.

On the whole, so far, our main formula with its characteristic Boltzmann factors compares quite well with experimental results.

We were led to a Planck temperature of about 1.3 GeV by the data. This should be thought of as the QCD Planck scale $\hat{M}_P = N^{1/4}\Lambda_{\mathrm{QCD}}$.

Also experimentally, the $\sigma_{\mathrm{tot}} \sim s^{1/11}$ behavior sets in at about 9 GeV [23,24], which we can thus identify with $\hat{E}_R = N^2\Lambda_{\mathrm{QCD}}$. Again, from the data we found $\hat{E}_R/\hat{M}_P = 9\ \mathrm{GeV}/1.3\ \mathrm{GeV}$, which agrees embarrassingly well with the theoretical prediction $N^2/N^{1/4} = 9/1.316$ for $N = 3$ colors. This is all the more mysterious, since, as mentioned in Sec. II, the values of energy scales can get possibly large string corrections. Could it be that, in the ratios of energy scales, the string corrections cancel?

We could have made a more rigorous analysis of the data, but as we already pointed out, there are several theoretical uncertainties, that could slightly modify the results, so it is not clear that it is useful getting better fits from the current experimental data, not until there is a better control on the theoretical uncertainties, even just in evaluating their size.

On the theoretical side we have already mentioned the uncertainties in the exact value of s_0 and the unknown potential corrections to the black hole Boltzmann factor. Also the hard Pomeron exponents $\tilde{\epsilon}_V$ for heavy vector mesons could change at a lower scale along our extrapolation. This could parallel what happens for the soft Pomeron in the gravity dual model. This would slightly affect the extrapolation to $\hat{M}_P$.

Then, the behavior of the extrapolation of the hard Pomeron from M_V down to $\hat{M}_P$ could also be slightly changed, as one needs at least $\sqrt{s} = M_V$ to even produce the mesons. For J/ψ and $\psi(2S)$ that difference between the $\hat{M}_P$ and M_V extrapolation is relatively small, but for Υ it could be larger, and it could contribute along with the large experimental errors to the observed discrepancy. Finally, we have assumed that the hard Pomeron extrapolated down to $\hat{M}_P$ gives the soft Pomeron contribution at $\hat{M}_P$, but the soft Pomeron contribution could be slightly smaller.

V. CONCLUSIONS

In this paper we have analyzed diffractive vector-meson photoproduction using AdS/CFT. The soft Pomeron behavior of QCD was argued to be due to production of dual black holes. The photoproduction of the ϕ meson is the cleanest test of the soft Pomeron. Indeed the ϕ mass is sufficiently small for the applicability of the soft Pomeron picture, and moreover mesonic Regge poles do not contribute to ϕ photoproduction [14].

For vector-meson photoproduction, the power $\tilde{\epsilon} = 2\epsilon - 2\alpha'_P/B$ differs in principle from the power ϵ determined from the energy behavior of total cross sections, but they are numerically surprisingly close. Assuming outright equality, $\tilde{\epsilon} = \epsilon$, the ϕ data match the s^ϵ-like theoretical prediction: $\sigma_V \sim s^{1/7}$ below 9 GeV and $\sigma_V \sim s^{1/11}$ above 9 GeV. The measured exponent of 0.11 is mostly due to data in the transition region around 9 GeV, and can be fit with an exponent of $(1/7 + 1/11)/2 = 0.117$, though of course a more detailed analysis of the AdS/CFT scattering is needed to determine if the matching is relevant. We emphasize though the absence of any theoretical understanding whatsoever for such an equality $\tilde{\epsilon} = \epsilon$ in the case of light vector mesons.

Then we have tested the general formula at $s = s_0$. The predicted ratios of soft Pomeron contribution to σ_V at about 1–2 GeV were successfully compared with experiment. This allowed us to extract the temperature $T_0 \sim 1.3$ GeV. This T_0 has to be close to $\hat{M}_P$, and we therefore identified their values. This led us to the value 9 GeV/1.3 GeV for the ratio of the soft Pomeron onset scale, $\hat{E}_R$, to $\hat{M}_P$. As was pointed out, this agrees well beyond all expectation with the prediction $\hat{E}_R/\hat{M}_P = N_c^2/N_c^{1/4} = 9/1.316$, though it is unclear why we have such good agreement, when *a priori* there could be dual string corrections.

We should emphasize that, while the Fermi-Landau statistical model also gives an effective temperature for the soft high energy scattering, the AdS/CFT description provides an understanding of the temperature T_0 in a quantum theory, as being due to a *unitary* process of dual black hole creation and decay. AdS/CFT also gives a prediction for $T_0 \sim \hat{M}_P$ that we tested experimentally. Of course, QCD is a good description for the scattering, so it should also be possible to modify the Fermi-Landau model [33,34] to obtain the same result.

An important further test of the ideas of this paper could be obtained from a more precise determination of the s dependence of the Υ photoproduction cross section, allowing for a better determination of the extrapolated J/Ψ:Υ split at $s_0 \sim 1$–$2\ \text{GeV}^2$.

Finally, we have used AdS/CFT, together with Regge theory, to describe vector-meson photoproduction, and both were needed to derive our results. But it should be in principle possible to find the same formulas from QCD, at least on the lattice, if not analytically.

ACKNOWLEDGMENTS

We would like to thank R. Enberg, whose graph we have used for a preliminary analysis, and for pointing us to the correct sources for the data. One of us (H. N.) would like to acknowledge valuable discussions with Chung-I Tan and the late Kyungsik Kang and thank the University of Chicago for the opportunity to visit, which resulted in the beginning of this work. The research of H. N. has been supported in part by MEXT's program "Promotion of Environmental Improvement for Independence of Young Researchers" under the Special Coordination Funds for Promoting Science and Technology, and also with partial support from MEXT KAKENHI Grant No. 20740128.

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